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Electroweak Interactions at High Energies

Proceedings of the 1982 DESY Workshop

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PREFACE

The 1982 DESY Workshop on Electroweak Interactions at High Energies was held at DESY, Hamburg, from September 28 to 30, 1982.

The objective of the meeting was twofold: Firstly, to investigate the present status of the Glashow-Weinberg-Salam theory (including radiative corrections and best fits to the various parameters) together with the room available for extensions (in particular Grand Unification schemes and neutrino masses); secondly, to discuss novel interpretations of electroweak interactions within the framework of supersymmetry as well as connection with composite models of leptons and quarks.

The main outcome of the meeting was, that although the Standard Model is in a very good shape, there do exist promising alternative approaches, which can provide us with new structures in the TeV energy region.

Many of the topics discussed during the meeting lay in the center of Prof. J.J. Sakurai's interest. Being a prominent member of the Organizing Committee he not only contributed much to formulating the program, but also inspired the discussions during the workshop by his profound remarks and his challenging questions. His tragic and completely unexpected death a short time after the workshop has taken from us a very active and creative contributer in this subject matter. We have dedicated these Proceedings to his memory.

We are grateful to many people for their help in organizing the meeting and producing these Proceedings. Our special thanks go to the DESY secretaries for their organizational work in connection with the workshop and to Elisabeth Bähr and Bärbel Jahns for their helpful and competent assistance in preparing the manuscripts for the Proceedings.

R. Kögerler
D. Schildknecht

TO THE MEMORY OF

J. J. SAKURAI

J.J. Sakurai 1933 - 1982

The world of high energy physics was shocked by the sudden and untimely death of J.J. Sakurai in Ferney Voltaire during a three-month visit to CERN at the end of October 1982.

J.J. Sakurai was born in Tokyo, Japan, on January 31, 1933. He obtained his university education in the U.S. Having received a B.A. degree from Harvard University in 1955 and a Ph.D. from Cornell University in 1958, he joined the faculty of physics of the University of Chicago and was appointed Full Professor in 1964. He left the University of Chicago for UCLA in 1969 as a Visiting Professor, and was Professor at UCLA since 1970. He spent the academic years 1966/67 and 1975/76 at CERN, and the academic year 1981/82 at the Max-Planck-Institut für Physik und Astrophysik in Munich. Apart from these long term stays, he was a frequent visitor to Europe as well as the Far East. He contributed to numerous international conferences and summer schools all over the world and had intensive scientific contacts with theoretical and experimental colleagues in all major high energy physics laboratories.

J.J. Sakurai left a list of publications with 124 entries including three books. Right after having obtained his Ph.D., he became one of the founders of the V - A Theory of weak interactions by writing a widely quoted paper on the derivation of the V - A form from "mass reversal invariance" (1958). Subsequently he turned to strong interactions. In his 1960 Annals of Physics paper he made the "attempt to construct a theory of strong interactions in analogy with electromagnetism, in which the notions of conserved currents and universality play a fundamental role" (quoted from J.J. Sakurai, "Vector Mesons 1960 - 1968", in "Lectures in Theoretical Physics", Gordon and Breach, Inc., 1969, ed. K.T. Mahanthappa, W.E. Brittin and A.O. Barut). This paper contains the first application of a nonabelian gauge theory to strong interaction physics. The predicted vector mesons were subsequently discovered in the early 1960's, and are known as ρ^0, ω, ϕ . In the 1960's, in particular after high energy electron accelerators had been put into operation, it soon became recognized that ρ^0, ω, ϕ play a crucial role in photon and electron interactions at high energies.

"Vector Meson Dominance" and later "Generalized Vector Dominance" and "q^2-duality" were strongly advocated by J.J. Sakurai and became most important tools for interpreting the experimental data. After the discovery of weak neutral currents in 1973, J.J. Sakurai devoted his full energy to an analysis of neutral current phenomena "without gauge theory prejudices". Our thorough knowledge of neutral current parameters owes much to his persistence and endurance in interpreting the data. With $\gamma W^{(o)}$ mixing as an alternative explanation of neutral current structure (1978), which implies boson mass relations less stringent than the canonical ones of the spontaneously broken $SU(2) \times U(1)$ gauge theory, he was eagerly awaiting the discovery of charged and neutral weak bosons at the CERN $p\bar{p}$ collider, which he expected to happen during his stay in Geneva.

J.J. Sakurai was a very successful teacher and published the books "Invariance Principles and Elementary Particles" (1964), "Advanced Quantum Mechanics" (1967) and "Currents and Mesons" (1969). He left an almost complete manuscript on "Nonrelativistic Quantum Mechanics" on his desk. It is hoped that it will be possible to complete the book and publish it. Many of J.J. Sakurai's students have meanwhile developed a reputation of their own in the field of high energy physics.

I first met J.J. Sakurai during his first visit to DESY after the 1969 Vienna conference, and since then I saw him on numerous occasions in Europe as well as in the U.S. We saw each other last in Hamburg at a delightful dinner given for the lecturers right after the end of the workshop.

It has been the wish of his widow, Noriko, and of his two sons, Ken and George, that J.J. Sakurai's unbounded interest in and love for physics be perpetuated by establishing a fund in his memory, the J.J. Sakurai Memorial Fund, to be administered by UCLA.

Dieter Schildknecht

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I. THE STANDARD MODEL AND ITS EXTENSIONS

STANDARD MODEL, LEFT-RIGHT AND OTHER EXTENSIONS

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ABSTRACT

The standard model is reviewed with emphasis on the neutrino masses and the baryon and lepton number violation , arising from peculiarities of the Higgs structure. A review of neutral current interactions is also given in the context of left-right symmetric and other gauge models. Theoretical indications for possible non-standard behavior at high energies are also discussed.

1. Introduction

According to the present philosophy the dynamics in Nature should be such that the corresponding Lagrangian remains invariant under certain transformations, even when these are performed locally (say only in Hamburg) and change with time. This invariance requirement under local transformations naturally leads to Yang-Mills interactions[1] . Their generators T_i, as far as the matter-fields are concerned, are given by $T_i = \int J_o^i(x)d^3x$, in terms of the currents $J_\mu^i(x)$. They couple to the corresponding massless gauge vector fields like $W_\mu^i J^{i\mu}$.

For the gauge Lagrangian describing the electroweak interactions the charge operator Q should be a linear combination of T_i, and the photon field A_μ should be a linear combination of the fields $W_{i\mu}$. The Lagrangian symmetry defined by T_i, is not also a symmetry of the vacuum. This is usually achieved through the Higgs mechanism[2], which gives

$$T_i|o> \neq 0$$

for <u>some</u> T_i. The states $T_i|o>$ are massless spin-zero states which are eaten by the corresponding gauge bosons $W_{i\mu}$, thereby giving them mass.

2. The Standard Model (general)[3-5]

The gauge group is $SU(2)_L \times U(1)_Y$. The corresponding currents are $\vec{J}_{L\mu}$, $J_{Y\mu}$, and couple respectively to the gauge boson fields $\vec{W}_{L\mu}$, B_μ. The normalization is chosen such that

$$J_\mu^{em} = J_\mu^{(3)} + J_{Y\mu} .$$

The group assignment of any left-handed fermion f_L is

$$(1/2 \; , \; \frac{B-L}{2}) \; , \tag{1}$$

where B , L are the baryon and lepton numbers respectively, while the representation of a right-handed fermion f_R is

$$(0 \; , \; Q_f) \; . \tag{2}$$

Using these we get

$$\begin{aligned} \vec{J}_{L\mu} &= \sum_{f_L} \bar{f}_L \gamma_\mu \frac{\vec{\tau}}{2} f_L \; , \\ J_{Y\mu} &= \sum_{f_L} \left(\frac{B-L}{2}\right)_f \bar{f}_L \gamma_\mu f_L + \sum_{f_R} Q_f \bar{f}_R \gamma_\mu f_R \; , \end{aligned} \tag{3}$$

and the fermion part of the interaction

$$\begin{aligned} \mathcal{L}_f = &-eJ_\mu^{em} A^\mu - \frac{e}{\sin\theta_w \cos\theta_w} Z^\mu (J_{L\mu}^{(3)} - \sin^2\theta_w J_\mu^{em}) - \\ &- \frac{e}{\sqrt{2}\,\sin\theta_w} (W_\mu J^{\mu +} + \text{h.c.}), \end{aligned} \tag{4}$$

where Z_μ , W_μ are the neutral and charged weak boson fields respectively. The local current-current interaction, valid for $q^2 \ll M_W^2$, M_Z^2 is

$$\mathcal{H}_{eff} = 2\sqrt{2} G_F \{ J_{L\mu}^+ J_L^{-\mu} + \rho [J_{L\mu}^{(3)} - \sin^2\theta_w J_\mu^{em}]^2 \} \; , \tag{5}$$

where

$$4\sqrt{2} G_F = \frac{e^2}{M_W^2 \sin^2\theta_w} \; , \tag{6}$$

and

$$\rho \equiv \frac{M_W^2}{M_Z^2 \cos^2\theta_W} \simeq 1 \tag{7}$$

tolowest order in electroweak interactions,provided that the Higgs fields transform like the (1/2 , 1/2)- representation of $SU(2)_L \times U(1)_Y$. The theory is broken either spontaneously or dynamically ; i.e.

$$\begin{aligned} \vec{T}_L|0\rangle &\neq 0 \ , \\ T_Y|0\rangle &\neq 0 \ . \end{aligned} \tag{8}$$

But $U(1)_{em}$ gauge invariance is preserved since the electric charge operator satisfies

$$Q\ |0\rangle \equiv (T_L^{(3)} + T_Y)|\ 0\rangle = 0 \tag{9}$$

Eqs. (5) is in perfect agreement with low energy interactions,and I refer to various recent reviews for details[6,7] . Here I only note that (5) is consistent with the suppression of the flavor-changing neutral currents observed in $K_2 \to \mu^+\mu^-$, $K_1 - K_2$ mass difference, etc. Comparing lowest order theoretical calculations with neutral-current data on νN and $\bar{\nu} N$ inclusive interactions, $\nu N \to \nu\pi X$, $\nu p \to \nu p$, $\nu_\mu e \to \nu_\mu e$, $\bar{\nu}_\mu e \to \bar{\nu}_\mu e$, $e^+e^- \to \mu^+\mu^-$ and hadrons, $eD \to eX$ - asymmetries, and atomic physics experiments we get[6-9]

$$\sin^2\theta_W = 0.227 \pm 0.015 \tag{10}$$

and

$$\rho^2 = 1.00 \pm 0.08 \tag{11}$$

to be compared with the same-order-theoretical value (7).

If results from the next higher order are compared with the data we get[10]

$$\sin^2\theta_W^M \equiv 1 - \frac{M_W^2}{M_Z^2} = 0.217 \pm 0.014 \ , \tag{12}$$

and[6,10]

$$\rho^2_{exp} = 0.998 \pm 0.050 \tag{13a}$$

while the corresponding theoretical one is[10,6]

$$\rho^2_{th} = 0.983 \tag{13b}$$

Low energy charged-current interactions among quarks and leptons are also consistent with (5). Assuming three generations, the quark contribution to the weak charged current is given by

$$J^+_{L\mu} = (\bar{u}_L \ , \ \bar{c}_L \ , \ \bar{t}_L)\gamma_\mu U \begin{pmatrix} d_L \\ s_L \\ b_L \end{pmatrix} \tag{14}$$

where the Cabibbo-Kobayashi-Maskawa Matrix U has been determined through the efforts of J.J. Sakurai and others[11]. A representative form of U is[6]

$$U = \begin{pmatrix} 0.97 & 0.22 & 0.068 \\ -0.22 & 0.85-0.66\text{x}10^{-3}i & -0.48+2.1\text{x}10^{-3}i \\ -0.046 & 0.48+3.2\text{x}10^{-3}i & -0.88-1.0\text{x}10^{-3}i \end{pmatrix} \tag{15}$$

The imaginary parts in (15) supply a CP-violating interaction in agreement with observations. We also note that CP-violation can also come from instantons, due to the interaction[12]

$$\Theta_{QCD}/(32\pi^2)F^i_{\mu\nu} \ \tilde{F}^{\mu\nu} \ . \tag{16}$$

$F^i_{\mu\nu}$ is the field-strength for gluons. Interaction (16) implies an electric dipole moment for neutron, and experimental upper limits on those give[13,6]

$$|\Theta_{QCD}| < 2\text{x}10^{-9}. \tag{17}$$

It has recently been claimed that certain supersymmetry breakings can naturally guarantee such a small value[14].

3. Fermion masses[3-5].

The standard model accommodates the masses and mixings of the quarks and charged leptons through interactions of the Higgs fields. They have the form[15]

$$\mathcal{L}_{f} = -\sum_{\substack{f,f'=\\=d,s,b}} \left(H^{N}_{ff'}(\bar{f}_L\Phi)f'_R + h.c. \right) - \sum_{\substack{f,f'=\\=u,c,t}} \left(H^{P}_{ff'}(\bar{f}_L\tilde{\Phi})f'_R + h.c. \right) - \sum_{\substack{l,l'=\\=e,\mu,\tau}} \left(H^{L}_{ll'}(\bar{l}_L\Phi)l'_R + h.c. \right) \tag{18}$$

where Φ is the Higgs doublet and $\tilde{\Phi} = i\tau_2\Phi^*$. Eqs. (18) supplies Dirac masses to all quarks and charged leptons. A Majorana mass is impossible for charged particles since it would violate charge conservation. On the other hand no firm predictions are available for the mass matrices[15] H^N, H^P, H^L in the simplest version of the model. Extra assumptions are needed for the horizontal structure and /or assumptions concerning the way that the theory (supersymmetrically) unifies.

We now turn to the neutrino masses. Here a Majorana mass implying L-violation is possible, since no-breaking of any gauge symmetry is ever involved. There are many ways to get ν-masses in the standard model.

i) In the simplest version one introduces ν_R in the theory transforming according to (2) under $SU(2)_L \times U(1)_Y$. We then obtain the gauge invariant interaction[16]

$$\mathcal{L}_{\nu} = H_d \nu_L^{c\tau} C\, \tilde{\Phi}^{+} \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} + h.c. + M_m \nu_L^{c\tau} C \nu_L^{c} + h.c. \tag{19}$$

which involves explicit L-breaking. Here C is the usual charge conjugation matrix ; i.e. $\nu_L^{c\tau}C \equiv \bar{\nu}_R$. The neutrino mass matrix then is[16]

$$\mathcal{L} = \ldots - (\nu_L^{\tau}C\ ,\ \nu_L^{c\tau}C) \begin{pmatrix} 0 & m_d \\ m_d & M_\nu \end{pmatrix} \begin{pmatrix} \nu_L \\ \nu_L^{c} \end{pmatrix} + h.c. \tag{20}$$

where

$$m_d = \frac{H_d v}{2} \quad , \quad \langle\Phi\rangle_0 = \frac{1}{2}\binom{0}{v}$$

is the Dirac neutrino mass expected to satisfy

$$m_d \sim O(m_\rho) \quad , \tag{21}$$

and M_m is a Majorana mass. Since the M_m-term in (19) is invariant under $SU(2)_L \times U(1)_Y$, M_m can be as large as, or even much larger than M_W. Grand unification schemes or Horizontal symmetry models can easily suggest

$$M_m \sim 10^{19}\ \text{GeV} \quad \text{or} \quad \sim 10^{13}\ \text{GeV}$$

Diagonalizing (20) we get a light neutrino[17]

$$\nu_m \simeq \nu_L - \left(\frac{m_d}{M_m}\right)\nu_L^c \tag{22a}$$

with mass

$$m_{\nu_m} = m_d^2 / M_m \ll m_\rho \tag{22b}$$

in agreement with observations, and a heavy neutrino

$$N = \nu_L^c + \left(\frac{m_d}{M_m}\right)\nu_L \tag{23a}$$

with mass

$$m_N \simeq M_m \tag{23b}$$

A non-vanishing value for m_{ν_m} makes possible neutrino oscillations, as well as the appearance of neutrino-less double $(\beta\beta)_{0\nu}$ decays[19]

$$(A\ ,\ Z) \to (A\ ,\ Z+2) + e^- + e^- \tag{24a}$$

even in the absence of right handed charged current interactions. In Fig. 1, the Feynman diagrams for the $(\beta\beta)_{0\nu}$ and the competing process $(\beta\beta)_{2\nu}$

$$(A\ ,\ Z) \rightarrow (A\ ,\ Z+2) + e^- + e^- + \nu + \nu\ , \tag{24b}$$

are shown. Other related processes which can in principle give information on m_{ν_m} include neutrino-less double positron emission, electron positron conversion

$$e^- + (A\ ,\ Z) \rightarrow (A\ ,\ Z+2) + e^+\ , \tag{25}$$

and double electron capture

$$e^- + e^- + (A\ ,\ Z) \rightarrow (A\ ,\ \underbrace{Z-2)^*}_{\hookrightarrow\ \gamma} \tag{26}$$

Of course with each of these processes there is a 2ν related process competing[18]. It is therefore very hard to draw useful conclusions for m_{ν_m} from such measurements, particularly because theoretical calculations, being sensitive to nuclear wave functions, are rather uncertain[6]. Notice also that it may be possible to enhance e.g. process (24a) with respect to (24b), by restricting to e^-e^- pairs having the highest possible energy.

ii) There may exist an (unknown) reason prohibiting any explicit breaking of the global L or B or L-B symmetries[17,19]. There have consequently been papers where these symmetries were broken spontaneously through the introduction of extra Higgs fields. In such cases the second term in (19) is replaced by[17,19,20]

$$\nu_L^{c\tau} C \nu_L^c \Delta_0 \tag{27a}$$

where Δ_0 is a new Higgs field transforming like (0,0) and having $L_{\Delta_0}=0$, or by

$$(-e_L^\tau C\ ,\ \nu_L^\tau C) \begin{pmatrix} \Delta^+ & \Delta^{++}/\sqrt{2} \\ \Delta^0/\sqrt{2} & -\Delta^+/\sqrt{2} \end{pmatrix} \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} \tag{27b}$$

where $L_\Delta = -2$. Such couplings give Majorana masses to the neutrinos, and at the same time imply the existence of some real Goldstone bosons (majorons) or very light Higgs particles with interactions of the form[21]

$$\mathcal{L} = -\frac{1}{2} m_N \bar{\nu}\nu + \frac{f}{2} \bar{\nu}(i\gamma_5 M + \chi)\nu. \tag{28}$$

Here M , χ are such very light bosons and $\nu = \nu_L + \nu_R^c$ is a Majorana neutrino field. Now, in addition to ν-oscillations and $(\beta\beta)_{0\nu}$, we may also have varying missing mass against l^+ in the decay-processes

$$K^+(\pi^+) \rightarrow l^+\nu \; M \text{ or } \chi , \tag{29}$$

as well as astrophysical implications. Experimental data for process (29) suggest[21] $f \lesssim 7\times10^{-3}$.

4. Baryon and Lepton number breaking.

Our freedom to enrich the Higgs structure of the standard model is close to infinity (!). Using it to introduce even more Higgs fields, and assigning suitable B , L eigenvalues to them, we can obtain processes involving $\Delta B = 2$ and/or $\Delta L = 2$ transitions[22]. The diagram of Fig. 2a gives an example for the possible n-$\bar{n}$ oscillations, while the one in Fig. 2b induces the transition $pp \rightarrow e^+e^+$. The later may be very interesting, implying the decay[23]

$$(A \; , \; Z) \rightarrow (A-2 \; , \; Z-2)e^+e^-$$

with a lifetime $\sim 10^{33}$ years, and a clear signal of back to back positrons each carrying an energy of $\sim$1 GeV.

These $\Delta B = 2$ and $\Delta L = 2$ transitions are <u>not</u> necessarily related to the GUT scale, since they are not generated by superheavy gauge bosons, which are hoped to play a dominant role in genuine p-decays like $p \rightarrow e^+\pi^0$. They just depend on the existence of some extra peculiar Higgs fields with conveniently assigned B and L quantum numbers. I believe that it is important to look for such processes. Their discovery will prove that there is something in the Higgs forces which is much deeper than our present level of understanding. It is also safe to state that the present great flexibility of the Higgs structure may reduce the size of the window[24] that B-violation will presumably supply to the physics at superheavy energies. Of course, if we stick to the principle of simplicity for the Higgs-potential, then the above processes, as well as the Majoron-type particles of the previous section, should not exist.

5. Simplest Higgs coupling in the standard model[3-5]

The simplest version of the Higgs potential in the standard model includes only the famous douplet Φ which[6] satisfies

$$\langle\Phi\rangle_0 = \frac{1}{\sqrt{2}} \binom{0}{v} \quad , \quad v = (\sqrt{2}\, G_F)^{-1/2} = 250 \text{ GeV} \tag{30}$$

The physical Higgs boson H has a mass

$$m_H = \sqrt{2}\, \mu = v\sqrt{\frac{\lambda}{2}} \quad , \tag{31}$$

and μ , λ are defined through

$$V(\Phi) = -\mu^2 |\Phi|^2 + \frac{\lambda}{2} |\Phi|^4 \quad , \tag{32}$$

while its interactions with fermions f and the weak bosons are given by

$$\begin{aligned} \mathcal{L}_H &= -g_{Hf}\, H\bar{f}f + \frac{g_{HZ}}{2} HZ^2 + g_{HW}\, HW^+_\mu W^{-\mu} \,, \\ g_{Hf} &\quad m_f(\sqrt{2}\, G_F)^{1/2} \\ g_{HV} &= 2M_V^2 \{\sqrt{2}\, G_F\}^{1/2} \quad , \quad V = Z \text{ or } W. \end{aligned} \tag{33}$$

According to (33), if $2m_b \ll m_H < 2m_t$ the most important H-decay modes satisfy

$$\begin{aligned} &\Gamma(H \to \bar{b}b) \div \Gamma(H \to \bar{c}c) : \Gamma(H \to \bar{\tau}\tau) : \Gamma(H \to \bar{s}s) : \Gamma(H \to \bar{\mu}\mu) \\ &\simeq 3m_b^2 \quad : \quad 3m_c^2 \quad : \quad m_\tau^2 \quad : \quad 3m_s^2 \quad : \quad m_\mu^2 \end{aligned}$$

If m_H is smaller than the ortho-$\bar{t}t$ mass, then H could be seen in the decay[25] ${}^3S_1(\bar{t}t) \to H\gamma$. We expect $B({}^3S_1 \to H\gamma) > 1\%$. If $m_H \lesssim 100$ GeV, then the best way to discover H is $e^+e^- \to ZH$ since[6]

$$\frac{\sigma(e^+e^- \to ZH)}{\sigma_{pt}} \gtrsim 0.1 \qquad (34)$$

for $\sqrt{S} \simeq 200$ GeV. This mechanism will not work if the Higgs is a composite particle like the neutral pseudo-Goldstone bosons $P^{0,3}$ in the technicolor model. In the later case there is no strong HZ^2 coupling, and the usual Adler-Bell-Jackiw anomaly contribution from a techniquark loop gives[26]

$$\frac{\sigma(e^+e^- \to Z^0 + P^{0,3})}{\sigma(e^+e^- \to \mu^+\mu^-)_{pt}} \lesssim 10^{-3} - 10^{-4} \qquad (35)$$

So if you don't see anything produced against the Z in $e\bar{e}$-annihilation, then the Higgs-particle is probably composite. A study of the photon polarization in ${}^3S_1(\bar{t}t) \to H\gamma$ could[27] also discriminate between the interactions $H\bar{t}t$ and $iH\bar{t}\gamma_5 t$.

6. The Masses and interactions of the weak bosons.

To lowest order and for $\sin^2\theta_W = 0.227 \pm 0.015$, the standard model gives[6-9]

$$M_W = 78.2 \,^{+2.7}_{-2.5} \text{ GeV} ,$$

$$M_Z = 89.0 \,^{+2.2}_{-2.0} \text{ GeV} ,$$

to be compared with the results expected if first order radiative corrections are included

$$M_W = 83.1 \,^{+3.1}_{-2.8} \text{ GeV} ,$$

$$M_Z = 93.9 \,^{+2.5}_{-2.2} \text{ GeV} .$$

Moreover gauge theories demand a specific form for the γWW and the ZWW interactions. In particular their γW^+W^- coupling demands canonical $W^\pm$ dipole magnetic moment, and no electric quadrupole moment; i.e. $\mu_W^\gamma = \frac{e}{M_W}$, $Q_W^\gamma = 0$. The ZW^+W^- vertex has a similar form with

$$g_{ZWW} = \frac{e}{tg\theta_W} ,$$

canonical weak W-magnetic moment

$$\mu_W^Z = \frac{g_{ZWW}}{M_W} ,$$

and no weak W-quadrupole moment; i.e.

$$Q_W^Z = 0 .$$

In principle we could also have C, P or T-violating interactions in the ZWW vertex, but no such terms are allowed in gauge theories[28].

Measurements in $e^+e^- \to W^+W^-$ may give[28] information on μ_W^γ, Q_W^γ, μ_W^Z, Q_W^Z as well as on the existence of C, P or T violating terms in ZWW. If some strong interaction is present at the 1 TeV range, then final state interaction effects will also look like T-violation in $e^+e^- \to W^+W^-$; (Fig. 3).

7. Are there any viable alternatives?

Can something non-standard be expected above 100 GeV? I first address myself to this question by remaining within the context of gauge theories.

i) Left-Right symmetric models [29]: The gauge group is taken to be $SU(2)_L \times SU(2)_R \times U(1)_{\frac{B-L}{2}}$ or $SU(2)_L \times U(1)_L \times U(1)_{\frac{B-L}{2}}$, under which the fermions transform like

$$f_L \sim (1/2 , 0 , \tfrac{B-L}{2}) , \tag{35}$$

$$f_R \sim (0 , 1/2 , \tfrac{B-L}{2}) .$$

The electromagnetic current is

$$J_\mu^{em} = J_{L\mu}^{(3)} + J_{R\mu}^{(3)} + \frac{1}{2} J_\mu^{B-L} \quad ,$$

and the interaction Lagrangian for the gauge bosons is

$$-\mathcal{L} = \ldots\ g_L \vec{W}_{L\mu} \cdot \vec{J}_L^\mu + g_R \vec{W}_{R\mu} \cdot \vec{J}_R^\mu + g_c B_\mu \frac{1}{2} J^\mu_{B-L} \tag{36}$$

where

$$g_L \equiv \frac{e}{\sqrt{x_L}} \quad , \quad g_R \equiv \frac{e}{\sqrt{x_R}} \quad , \quad \frac{1}{g_R^2} + \frac{1}{g_L^2} + \frac{1}{g_c^2} = \frac{1}{e^2}$$
$$0 < x_L + x_R < 1 \ . \tag{37}$$

The effective neutral current interaction at low q^2 is [30,7]

$$\mathcal{H}_{eff}^{NC} = 2\sqrt{2}\, G_F\{(\rho_1^2 + \rho_2^2)[J_{L\mu}^{(3)} - x_L J_\mu^{em}]^2 + \eta^2 [J_{R\mu}^{(3)} - x_R J_\mu^{em}]^2 +$$
$$+ 2\eta\rho_2 [J_{L\mu}^{(3)} - x_L J_\mu^{em}][J_{R\mu}^{(3)} - x_R J_\mu^{em}]\} \quad , \tag{38}$$

having as independent parameters x_L , x_R , ρ_1 , ρ_2 , η. We note that (38) becomes identical to Weinberg-Salam if $\rho_1^2 + \rho_2^2 = 1$, $\eta \to 0$ and $x_L \equiv \sin^2\theta_W$. The η^2-term is blind to ν_L-scattering[7]. The contribution of the $\eta\rho_2$-term to the e-q parity violating interaction is proportional to $x_L - x_R$, so that it vanishes if the condition $g_L = g_R$ is imposed. Barger, Ma and Whisnant[30], and also Sehgal[7] presented fits to the neutral current data, which are shown in Figs.4, 5, 6. There exist important correlations among the independent parameters ρ_1^2 , ρ_2^2 , η , x_L, x_R. The allowed range of these values for any symmetry breaking is shown in Figs. 4, 5, 6. This range is significantly reduced if the symmetry breaking is required to arise from a Higgs mechanism[30]. The data are consistent with $x_L \simeq 0.23$, $\rho_1^2 \simeq 1$, $\rho_2^2 \lesssim 0.05$, $\eta \sim 0.17 \pm 0.17$, $0 \lesssim x_R \lesssim 0.8$, and demand that the masses of the two neutral gauge bosons satisfy

$$83 \text{ GeV} < M_{Z_1} < 115 \text{ GeV} \ ,$$
$$205 \text{ GeV} < M_{Z_2}$$

For the charged gauge bosons the most useful constraint seems to come from K_S, K_L mas difference which gives[31] $M_{W_R} > 1.6$ TeV; (M_{W_L} should be very close to the Weinberg-Salam value). Summarizing, if a left - right symmetric model holds in Nature and $M_{W_R} > M_{Z_2}$, then a possible scenario of the symmetry breakings may be

$$\underset{(\text{with } g_L = g_R)}{SU(2)_L \times SU(2)_R \times U(1)_{\frac{B-L}{2}}} \longrightarrow$$

$$\xrightarrow[M_{W_R}]{} SU(2)_L \times U(1)_R \times U(1)_{\frac{B-L}{2}} \xrightarrow[(M_{W_L}=M_W)]{} U(1)_{em} .$$

ii) Minimal Extentions[32-35]: These extentions of the standard model are based on $SU(2)_L \times U(1)_Y \times G$, where G is an arbitrary group affecting only the Higgs fields. All quarks and leptons are invariant under G, and transform under $SU(2)_L \times U(1)_Y$ in the standard way. The spontaneous breaking mixes the G-gauge bosons with the $SU(2)_L \times U(1)_Y$-gauge bosons. So there may be many Z's. At low energies, the general structure of the NC interaction is

$$\mathcal{H}^{NC}_{eff} = 2\sqrt{2}\, G_F \{(J^{(3)}_{L\mu} - \sin^2\theta_w J^{em}_\mu)^2 + C(J^{em}_\mu)^2\} \tag{39}$$

where $C \geq 0$, and vanishes for the standard model. Experiment[36] suggests $C \leq 0.02$ (95% C.L.).

Examples of such models are : (M_Z is the standard Z-mass)[33,34,37]

$SU(2)_L \times U(1)_Y \times U(1)'$ — 80 GeV $< M_{Z_1} \quad M_Z \longrightarrow C \propto \cos^4\theta_w$,

$SU(2)_L \times U(1)_Y \times SU(2)'$ — 60 GeV $< M_{Z_1} \quad M_Z$ — $C \propto \sin^4\theta_w$,

$SU(2)_L \times U(1)_Y \times SU(2')$ — $M_{Z_1} < 115$ GeV — $C \propto \sin^2\theta_w \cos^2\theta_w$.

Concluding this section we remark that within the context of the gauge theories, the experimental data suggest that the lightest Z cannot be heavier than[37] 115 GeV.

On the other hand Bjorken's general treatment based on $W^{(0)}$-γ mixing leads also to (39), with $C \propto \sin^4\theta_W$ and vanishing in the one neutral weak boson case[38]. This case has been thoroughly studied by Hung and Sakurai[39]. It seems that if the lightest Z is found experimentally above 115 GeV, then gauge theories are excluded and only the mixing schemes survive[40]. Such schemes would suggest the existence of some new strong interactions in the TeV scale effecting the γ-$W^{(0)}$ mixing, and being responsible for $W^{\pm}$, Z in the sence that QCD is responsible for $\rho^{\pm}$, ρ^{0}, But the existence of such new strong interactions is not necessarily inconsistent with standard Z and W masses.

7. (Theoretical) Difficulties of the Standard Model.

Since up to now there are no experimental difficulties of the standard model, we can only use theoretical arguments to motivate the idea that something may go wrong at high energies. These arguments start from the simplest Higgs potential

$$V(\Phi) = -\mu^2|\Phi|^2 + \frac{\lambda}{2}|\Phi|^4 .$$

Higher order corrections to $V(\Phi)$ coming from one-loop diagrams give a quadratically divergent contribution to μ^2 ; i.e.

$$\mu^2 \longrightarrow \mu^2 + \delta\mu^2 = \mu^2 + g^2\Lambda^2 ,$$

where Λ is the usual ultraviolet cut-off[41]. This is the hierarchy problem. The radiative corrections to μ^2 tend to be very large, and they can easily push μ and m_H (see Eqs. 31) in the TeV range where the Higgs-forces are essentially a new kind of strong interactions[42]. There are two attitudes to take toward this. Either you like these new strong interactions or you don't like them.

i) If you like them, then you can also find other traces of their existence. This is done by a suitable interpretation of the $W^{(3)}$- γ mixing which in the standard model creates the M_W, M_Z spilitting. Following Feynman and Sakurai we start from the kinetic energy part for the fields $W_\mu^{(3)}$ and B_ν in the Weinberg-Salam Lagrangian.

$$\mathcal{L}_G = -\frac{1}{4}(\partial_\mu W_\nu^{(3)} - \partial_\nu W_\mu^{(3)})^2 - \frac{1}{4}(\partial_\mu B_\nu - \partial_\nu B_\mu)^2 =$$

$$= -\frac{1}{4}(W_{\mu\nu}^{(3)})^2[1 + (\frac{g'}{g})^2] - \frac{1}{4}(F_{\mu\nu})^2(\frac{g'}{e})^2 -$$

$$- \frac{1}{4}(\frac{g'^2}{eg})\ [W_{\mu\nu}^{(3)} F^{\mu\nu} + F^{\mu\nu}W_{\mu\nu}^{(3)}] \quad , \tag{40}$$

$$W_{\mu\nu}^{(3)} = \partial_\mu W_\nu^{(3)} - \partial_\nu W_\mu^{(3)} \quad ,$$

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \ .$$

The wavefunction renormalization constants for $W_\mu^{(3)}$, A_ν are

$$Z_{3\gamma} = (\frac{e}{g'})^2 = \cos^2\theta_W \ ,$$

$$Z_{3W^{(3)}} = \{1 + (\frac{g'}{g})^2\}^{-1} = \cos^2\theta_W \ ,$$

leading, in terms of the renormalized fields, to

$$\mathcal{L}_G = -\frac{1}{4}(W_{\mu\nu}^{(3)R})^2 - \frac{1}{4}(F_{\mu\nu}^R)^2 - \frac{1}{4}\lambda(F_{\mu\nu}^R W_{\mu\nu}^{(3)R} + W_{\mu\nu}^{(3)R}F_{\mu\nu}^R) \ , \tag{41}$$

$$\lambda = \sin\theta_W \ .$$

The γ-$W^{(3)}$ mixing implies

$$\langle 0|eJ^{em}|W^{(3)}\rangle = M_W^2 \sin\theta_W \ . \tag{42a}$$

Interpreting this mixing as due to a new kind of strong interactions, and comparing (42a) with the QCD-induced mixing

$$\langle 0|eJ_\mu^{em}|\rho\rangle = e\,\frac{m_\rho^2}{f_\rho} \ , \tag{42b}$$

we get[43]

$$\frac{\sin\theta_W M_W^2}{em_\rho^2/f_\rho} \simeq 8.6 \times 10^4 < (\frac{\Lambda_{new}}{\Lambda_{QCD}})^2 \ ,$$

which means that the scale of this new strong interactions may be as low as 1 TeV. At this point there are still two possibilities. Either $W^{\pm}$, Z are ordinary bound states created by this new strong interaction[43]; or $W^{\pm}$, Z are still elementary gauge bosons which eat some of the Goldstone bosons of the new strong interaction (technicolor idea)[44]. In both cases there may be lots of surprises at high energies. At least the W , Z - masses and the structure of $e^+e^- \to W^+W^-$ may be completely different from the standard expectations. The old days when humble theorists were listening to their experimental colleagues will come again.

ii) But many theorists don't like this. So the only way left is supersymmetry[45]. Supersymmetry can save us from the quandradic divergences. Higgs masses can be kept low; i.e. $\sim O(M_W)$. And there will be no desert up to $\sim$1 TeV, since the supersymmetric partners of the known particles should lie there.

Acknowledgments

These proceedings have been dedicated to the memory of Professor J.J. Sakurai. Personally I would like just to add here my sincere gratitude to him for the philosophy in physics that he taught me . Thanks are also due to the Greek National Research Foundation for partial support.

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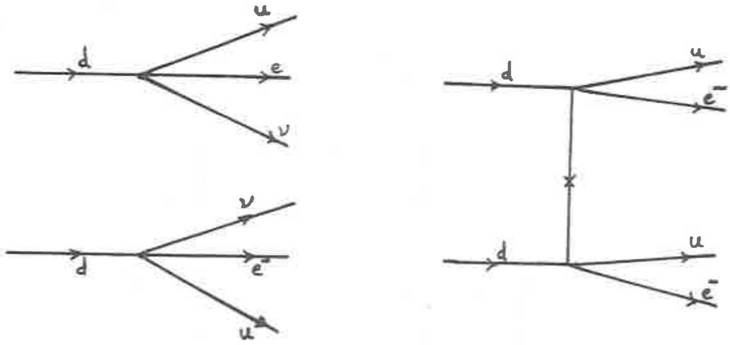


Fig. 1 Feynman diagrams for $dd \to e^-e^-\bar{\nu}\bar{\nu}uu$ and $dd \to e^-e^-uu$.

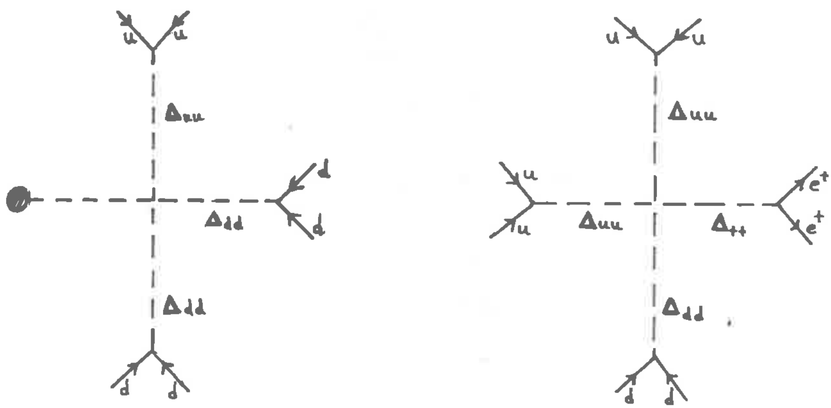
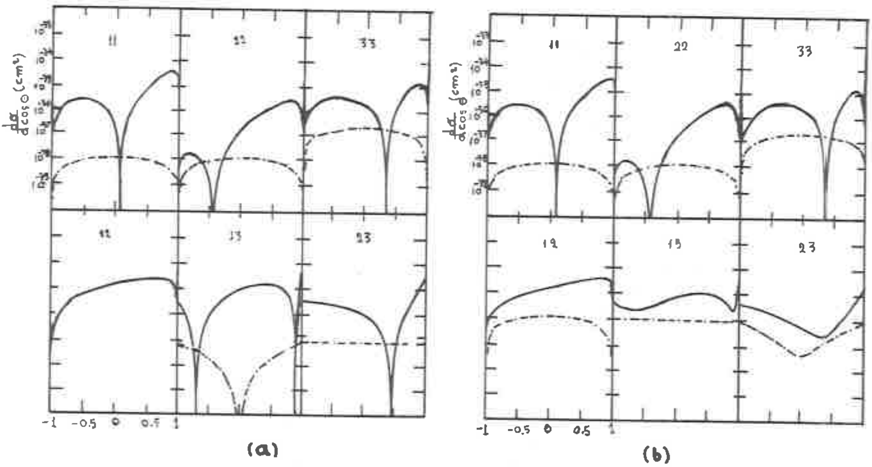


Fig. 2 Feynman diagram for n-$\bar{n}$ oscillation (a), and process $p+p \to e^+e^+$ (b).



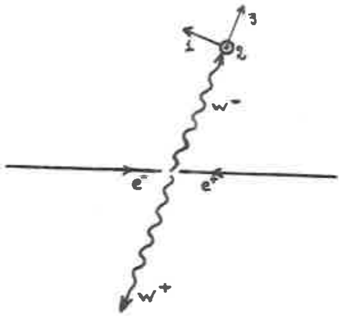


Fig. 3 Sensitivity of the $e^+e^- \to W^+W^-$ cross section to the polarization of the initial $e^\pm$ at 100 GeV per beam. The solid curves are for $e^-_L e^+_R$, the dashed curves for $e^-_R e^+_L$. The notation for different W-helicity is shown in the bottom part of the figure. (a) gives the standard model prediction. (b) includes a T-violating contribution in ZWW vertex.(see Ref. 28)

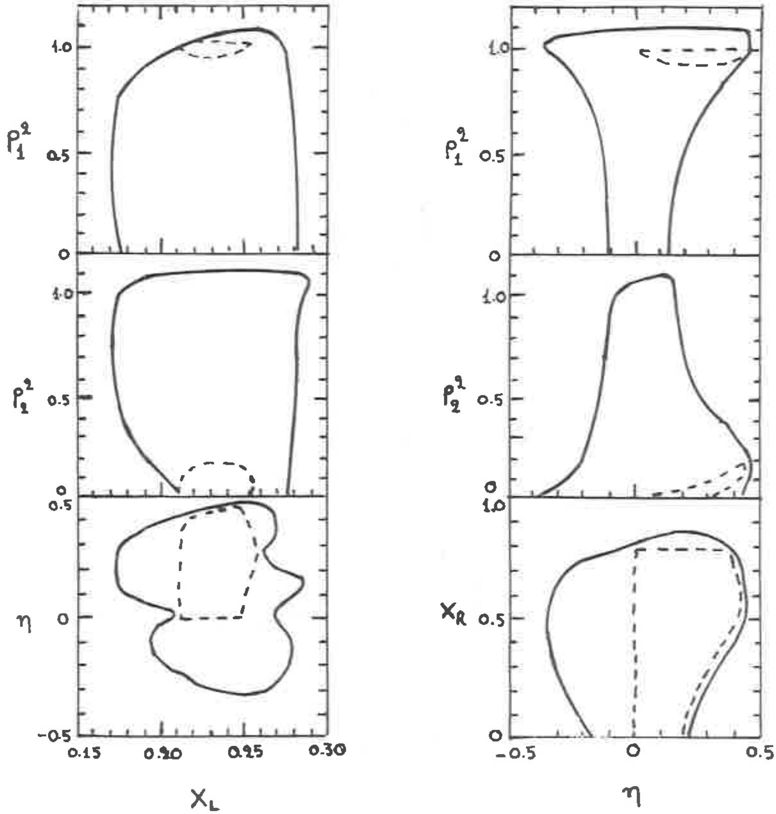


Fig. 4 Correlation of parameters in left-right models. Solid curves denote boundary of the most general model. Dashed lines denote boundary when the left-right gauge symmetry is broken by a Higgs mechanism.

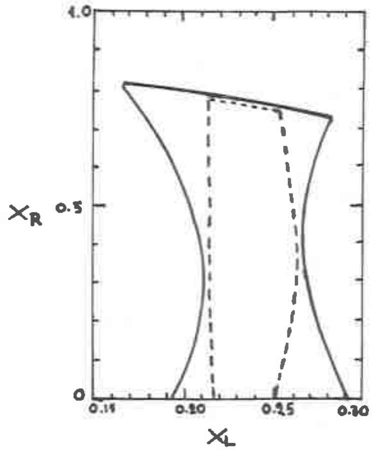


Fig. 5 Correlation of parameters in left-right models. Dashed and solid curves as in Fig. 4.

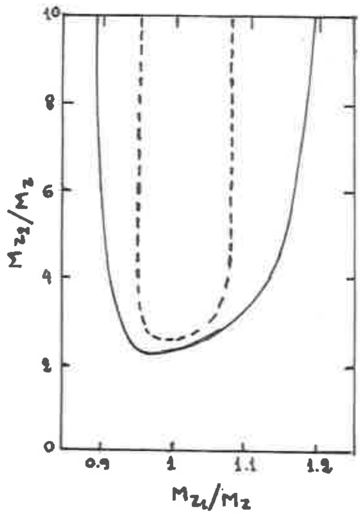


Fig. 6 Allowed region of M_{Z_1}/M_Z versus M_{Z_2}/M_Z (M_Z = 93.8 GeV). Dashed and solid curves as in Fig. 4.

STANDARD MODEL REVISITED

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ABSTRACT

The standard model (of Glashow, Salam and Weinberg) of electroweak interactions at low energies is reviewed. The extent to which "it" (i.e., the standard model) is fixed is discussed. The extent to which "it" is tested, in both charged as well as neutral currents, is described. CP violation in standard and near-standard models is discussed. Potential problems are mentioned.

1. The Model

I shall summarize the current status of the standard model with a brief discussion of "near-standard" models.

By the standard model, [1)] I mean sequential SU(2) x U(1) with one Higgs doublet. The low q^2 effective Lagrangian is

$$\mathcal{L}_{eff} = \mathcal{L}_{free}(W^{\pm}, Z^{\circ}, \gamma, q_i, \ell_i, \nu_i, H^{\circ}) + \mathcal{L}_{CC} + \mathcal{L}_{NC} + \mathcal{L}_{e.m.} + \mathcal{L}_{H}, \quad (1)$$

where $\mathcal{L}_{free}$ contains kinetic energy and mass terms for all particles and

$$\mathcal{L}_{e.m.} = e\, J_{\mu}^{e.m.} A_{\mu} \quad (1b)$$

$$\mathcal{L}_{CC} = \frac{4G_F}{\sqrt{2}} J_{L\mu}^{\dagger} J_{L\mu} \quad (1c)$$

$$\mathcal{L}_{NC} = \frac{4G_F}{\sqrt{2}} (J_{L\mu}^{3} - \sin^2\Theta_W J_{\mu}^{e.m.})^2 \quad (1d)$$

$$\mathcal{L}_H = -\frac{m_H^2}{4v^2} H^4 + \sum_i \frac{m_i}{v} \bar{f}_i f_i H + \ldots \tag{1e}$$

$$J_{L\mu} = (\bar{u}\ \bar{c}\ \bar{t})_L \gamma_\mu U_{KM} \begin{pmatrix} d \\ s \\ b \end{pmatrix}_L + (\bar{\nu}_1 \bar{\nu}_2 \bar{\nu}_3)_L \gamma_\mu U \begin{pmatrix} e \\ \mu \\ \tau \end{pmatrix}_L \tag{1f}$$

and the dots indicate trilinear Higgs coupling and Higgs couplings to W and Z. The relationship of the gauge couplings (of SU(2) x U(1)) g and g', mixing angle θ_W, Higgs vacuum expectation value v to the Fermi constant and the masses is given by

$$\frac{1}{2v^2} = \frac{g^2}{8m_W^2} = \frac{G_F}{\sqrt{2}} \quad , \quad g \sin\theta_W = e \ , \tag{2}$$

$$\tan\theta_W = \frac{g'}{g}$$

$$M_Z^2 = \frac{M_W^2}{\cos^2\theta_W} \quad . \tag{3}$$

The last is a prediction of the model, which follows from the iso-doublet nature of the symmetry breaking, to be tested.

How many free parameters are there in the model?

(g,g',v) or $(e,\sin\theta_W,G_F)$	3
$m_u,m_c,m_t,m_d,m_s,m_b;\ \theta_i,\delta$	10
$m_e,m_\mu,m_\tau,m_{\nu_1},m_{\nu_2},m_{\nu_3};\ \phi_i,\eta$	10
m_H	1

So there are 24 (25 if photon mass is counted) free parameters which have to be fixed.

2. Fixing It

How many are fixed so far? We know $(e,\sin\theta_W,G_F)$, $(m_u,m_c,m_d,m_s,m_b,U_{ud},U_{us})$ and (m_e,m_μ,m_τ); i.e., 13 out of 24! Actually, we do know a little more about some of the other parameters, in particular, about

U^{ij}_{KM}, U^{ij} and m_{ν_i}. Consider U_{KM} first. To extract information about G_F and U_{KM} it has to be assumed[2] that a) all $m_{\nu_i} = 0$ or b) $U = 1$ or c) all $m_{\nu_i} \ll Q$ in μ-decay and other semi-leptonic decays. Then G_F is fixed by μ-lifetime and U_{ud} can be measured via ft-values of allowed nuclear β-decays. Radiative corrections play an important role[3] and the final result for U_{ud} can be written as[3-5]

$$|U_{ud}|^2 = \frac{192(\ln 2)\tau_\mu(m_\mu/m_e)^5 f(m_e/m_\mu)\ [1 - \alpha/8\pi(\pi^2 - 25)]}{ft(1 + \delta_R)(1 - \delta_c)(1 + \Delta_R - \delta_\mu)} \tag{4}$$

where δ_c and δ_R are Coulumb and internal radiative nuclear corrections, Δ_R and δ_μ are electro-weak corrections to β-decay and μ-decay. The most accurate ft values are those of O^{14} and Al^{21}. Combining the recent analyses, we conclude:

$$|U_{ud}| = 0.9734 \pm 0.0025. \tag{5}$$

The largest uncertainties are in δ_c and in evaluating hadronic matrix elements of quark operators.

U_{us} can be extracted in a similar way from $K \to \pi e^- \nu_e$ and $Y \to N e^- \nu_e$, albeit with less accuracy. In case of K_{e3}, apart from U_{us} and G_F there are two form factors $f_-(q^2)$ and $f_+(q^2)$, whose ratio and q^2 dependence is measured. $|f_+(0)|^2$ is determined theoretically[6] (broken chiral SU(3) x SU(3)) to be about 0.98. Then one finds[5,7] $|U_{us}| = 0.219 \pm 0.003$. The hyperon semileptonic decays[8] are slightly problematic. One potential problem is the rate for $\Gamma(\Sigma^- \to \Lambda e^- \nu_e)$ which is lower than SU(3) fits would like and the other is the electron asymmetry in polarized Σ^--decay($\Sigma^- \to n e^- \nu_e$) which may have the sign opposite to the expected value. Now the first can be accommodated with SU(3) breaking in the axial-vector current matrix element without affecting the overall fit drastically whereas the second, if correct, would require a profound change in our understanding of charged weak currents. Ignoring this second problem, recent fits[4,9] to hyperon decays yield values for $|U_{us}|$ ranging from 0.225 ± 0.002, to 0.227 ± 0.003 or even 0.228, all of which are somewhat higher than values obtained earlier[5], viz. 0.222 ± 0.003. For the present purposes, we adopt a crude average from K_{e3} and hyperon decays of

$$|U_{us}| = 0.224 \pm 0.006. \tag{6}$$

A better way to use the hyperon data is to make a two angle fit (θ_V and θ_A) and use only θ_V to determine $|U_{us}|$, thus minimizing dependence on the uncertain axial vector matrix elements.

From the bounds (at 1σ level) on $|U_{ud}|$ and $|U_{us}|$ in Eq. (3) and (4) and the unitarity of U_{KM} we find

$$|U_{ub}| < 0.1 \tag{7}$$

in agreement with the earliest bound[10)] obtained on $|U_{ub}|$. In principle, $|U_{ub}|$ can also be bounded by the same-sign dimuon rate in $\bar{\nu}_\mu N$. This has a contribution from $\nu_\mu u \to \mu^+ b \to \mu^+ c \to \mu^+\mu^+ x$ which can be bounded by the unknown rate. However, due to the high $\sigma(\bar{\nu}_\mu \to \mu^+\mu^+)$, the threshold suppression in b-production and the unfavorable y-dependence, the bound obtained on $|U_{ub}|$ is weaker than (7).[11)] A more promising way of constraining $|U_{ub}|$ may be to use the "unexplained" part of $\sigma(\bar{\nu}_\mu \to \mu^+\mu^-\mu^+)$ i.e., the observed cross-section after subtracting the cross-section calculated from all conventional sources.

The unitarity limit on U_{cd} from U_{ud} in (5) is $|U_{cd}| < 0.24$. To get a non-zero lower bound it was suggested[12)] that one use the valence contribution to $\sigma(\nu_\mu \to \mu^- l^+)$ which comes from $\nu_\mu d \to \mu^- c \to \mu^- l^+ x$. Using the recent CDHS data for the fraction of valence in $\sigma(\nu_\mu \to \mu^-\mu^+)$ viz. 52%, one finds $|U_{cd}| > 0.2$. A more complete analysis by CDHS[13)] using the identity

$$|U_{cd}|^2 = \frac{2}{3\langle B\rangle} \; \frac{\left[\dfrac{\sigma(\nu \to \mu^-\mu^+)}{\sigma(\nu \to \mu^-)} - R\dfrac{\sigma(\bar{\nu} \to \mu^+\mu^-)}{\sigma(\bar{\nu} \to \mu^+)}\right]}{1 - R} \tag{8}$$

where $R = \sigma(\bar{\nu} \to \mu^+)/\sigma(\nu \to \mu^-)$ and using fragmentation function they measured finds $|U_{cd}| = 0.24 \pm 0.03$. So summarizing

$$0.2 < |U_{cd}| < 0.24. \tag{9}$$

In the same analysis CDHS[13)] also find that $(|U_{cs}|^2/|U_{cd}|^2) \times (2S/U + D) = 0.92 \pm 0.06$ and from their charged current data $(\bar{U} + \bar{D})/(U + D) = 0.13 \pm 0.02$. So if $2S < \bar{U} + \bar{D}$ they deduce the most conservative bound $|U_{cs}| \geq 0.59$. But from CHARM analysis[14)] of their neutral current data (assuming s and d have identical neutral current couplings) we know that $2S/(\bar{U} + \bar{D})$ is close to 0.5. From this, it is possible to use CDHS data to deduce a stronger bound[15] on $|U_{cs}|$ viz. $|U_{cs}| > 0.7$ to 0.74. A similar bound on $|U_{cs}|$ can be obtained[12)] from $\Gamma(D^+ \to K^\circ e^+\nu_e)$. Using F*- dominance model for the form factors and a calculated value[16] for $|f_+(0)|^2 \approx \frac{1}{2}$, one finds $|U_{cs}| > 0.8$. The main uncertainties are the values for $B(D^+ \to e^+\bar{K}^\circ e^+\nu_e)$ and $\tau(D^+)$. Including the unitarity limit from $|U_{cs}|$, then[4]

$$0.8 < |U_{cs}| < 0.98. \tag{10}$$

It would be nice to check that indeed $|U_{cs}| < 1$ and that unitarity is satisfied.

There are several ways of measuring the ratio $|U_{ub}/U_{cb}|$ as reviewed[17)] in Madison. CLEO has reported[18)] on a new measurement of $\langle N_K\rangle_{T''}/\langle N_K\rangle_c = 1.88 \pm 0.28$ from which a limit $|U_{ub}/U_{cb}| < 0.4$ can be deduced. A stronger result from a study of the shape and the end-point of the electron spectrum in $B \to e\nu_e x$ is reported to this conference by CUSB[19)]:

$$|U_{ub}/U_{cb}| < 0.2. \tag{11}$$

If B-decay is described well by quark decay, as seems likely, the lifetime τ_B is given approximately by

$$\tau_B \cong \tfrac{1}{3}\tau_\mu \left(\frac{m_\mu}{m_B}\right)^5 \frac{1}{(|U_{cd}|^2 + 2.5|U_{cb}|^2)} . \tag{12}$$

With the new limit[20] on τ_B from JADE $\tau_B < 1.4 \times 10^{-12}$ s, and the smallness of $|U_{ub}|$, $|U_{cb}|$ should satisfy

$$0.57 > |U_{cb}| > 0.05. \tag{13}$$

where the upper bound comes from unitarity.

We can now use unitarity to limit U_{ti}, e.g., we find $|U_{td}| < 0.13$, $|U_{ts}| < 0.56$ and $0.99 > |U_{tb}| > 0.82$. Hence $|U_{ts}\, U_{td}| < 0.073$. However, a better bound on $|U_{ts}\, U_{td}|$ is obtained from saturating the dispersive part of $K_L \to \mu^+\mu^-$ amplitude[21] by one loop graphs which are dominated by t. If the fourth generation heavy quark contribution is small compared to t, then for $m_t > 20$ GeV, it can be shown that[2] $|U_{ts}\, U_{td}| < 0.035$. $|U_{tb}|^2 + |U_{ts}|^2 + |U_{td}|^2$ which should be ≤ 1 can be measured in principle in weak decays of topponium: $(t\bar{t}) \to \bar{t} + b + l\nu$ etc. The rate goes as $G_F^2 m_t^5 \sum_i |U_{ti}|^2$ and the branching ratio can be as high as 5 to 6% for $M_{t\bar{t}} \sim 50$ GeV. A possible signature is prompt e/μ accompanied by hadrons.[22] Since $\tau_T < 10^{-17}$ s, this may be the only way to measure the absolute value of t-couplings. Relative strengths such as $|U_{ts}/U_{tb}|$ can only be measured in $t \to bl\nu$, $t \to sl\nu$ etc.

A summary of our present knowledge of U_{KM} elements is:

$$U_{KM} = \begin{pmatrix} |U_{ud}| = 0.9734 \pm 0.024 & |U_{us}| = 0.224 \pm 0.006 & 0 < |U_{ub}| < 0.1 \\ 0.2 < |U_{cd}| < 0.24 & 0.8 < |U_{cs}| < 0.98 & 0.05 \leq |U_{cb}| < 0.57 \\ 0 < |U_{td}| < 0.13 & 0 \leq |U_{ts}| < 0.56 & 0.82 \leq |U_{tb}| \leq 0.99 \end{pmatrix} \tag{14}$$

If and only if there are three families, U_{KM} can be parameterized á la Kobayashi-Maskawa[23]

$$U_{KM} = \begin{pmatrix} C_1 & S_1C_3 & S_1S_3 \\ -S_1C_2 & C_1C_2S_3 - S_2S_3e^{i\delta} & C_1C_2S_3 + C_3S_2e^{i\delta} \\ S_1S_2 & -C_1S_2C_3 - C_2S_3e^{i\delta} & -C_1S_2S_3 + C_2C_3e^{i\delta} \end{pmatrix} \tag{15}$$

where C_i and S_i are $\cos\theta_i$ and $\sin\theta_i$ with $0 \leq \theta \leq \pi/2$ and $0 \leq \delta \leq 2\pi$. The value of $|U_{ud}|$ fixes $S_1 = 0.229$. The other constraints on U_{ij}

restrict S_2 and S_3 as shown in Figs. 1 to 3. In Fig. 1 and Fig. 2, $\sin\delta$ is small and δ is near 0 and π respectively. In Fig. 3, $\sin\delta$ is large (>0.1). These constraints can be converted into other parameterizations[24] for U_{ij}.

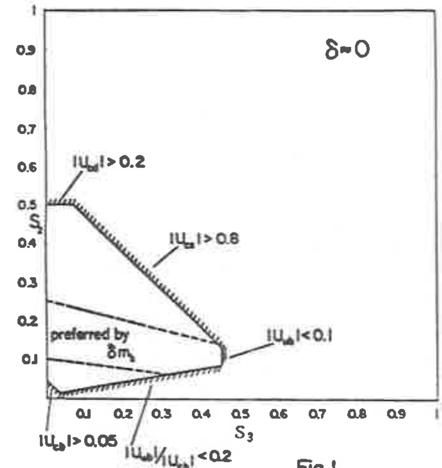

Fig.1

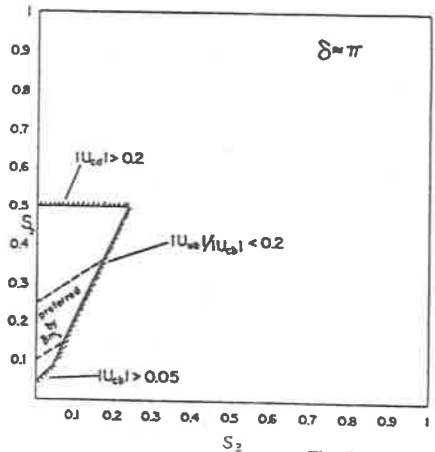

Fig.2

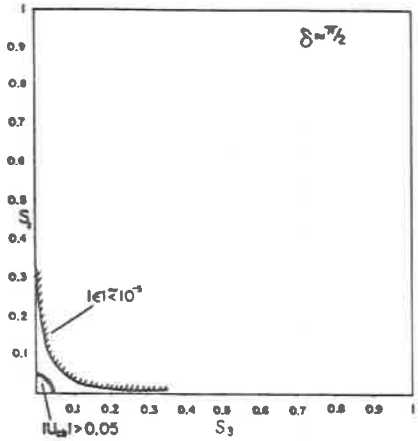

Fig.3

Bounds on θ_i and δ can also be derived from evaluating δm_{L-S} and Re ε for the K_L - K_S system.[25] The principle uncertainty is the evaluation of the hadronic matrix element. The bag-model evaluation is too sensitive to the parameters chosen and even the sign is unstable.[26] There is also the factor of 2 to 3 uncertainty due to possible long-range effects[27] (e.g., π^0, η, η', intermediate states). Two recent attempts to estimate this matrix element are very encouraging. One[28] relates the $\Delta S = 2$, $\Delta T = 1$ matrix element for $K^0 \leftrightarrow \bar{K}^0$ to the $\Delta S = 1$, $\Delta T = 3/2$ matrix element $\bar{K}^0 \to \pi^0$ in $K^+ \to 2\pi$ and obtains a suppression factor of 0.33 compared to vacuum saturation. The other[29] uses QCD sum rules to place a bound on the form factor for vacuum $\to$ KK yielding a limit of the suppression factor <0.6. Allowing for this uncertainty and allowing m_t to range between 20 and 60 GeV, the range for S_2 and S_3 preferred[25,30] by the K_L - K_S parameters is shown on the figures. Similar constraints[21,31] can be obtained from $K_L \to \mu^-\mu^+$.

In the lepton sector thing are not so good. For neutrino masses, we have the limits

$$\begin{aligned} m_{\nu_e} &< 50 \text{ eV} \\ m_{\nu_\mu} &< 500 \text{ keV} \qquad (16) \\ m_{\nu_\tau} &< 200 \text{ MeV} , \end{aligned}$$

where m_{ν_e} refers to the mass eigenstate coupled dominantly to e, and so on. From neutrino oscillation studies we have correlated bounds[32] on

δm_{ij}^2 and $|U_{i\alpha}\ U^*_{j\alpha}|$. Roughly speaking, for large (> 10 eV^2) δm^2, large off-diagonal mixings are ruled out ($4|U_i U_j|^2 < 0.1$).

3. Testing It

In charged currents there are many vulnerable tests of the V-A nature of the coupling. Unfortunately, most of them have not been repeated in a long time, and the limits are not very good, e.g., in nuclear β-decay a V+A amplitude at 10% is not ruled out.[33] From a combined analysis of β-decay and μ-decay, m_{W_R} (coupled to V+A with the $g_R = g_L$) can be constrained to be $\gtrsim 3\ m_W$. Recent checks of V-A structure in $\nu_\mu N \to \mu^- x$ are almost as good as in μ-decay; they constrain admixture of V+A to be less than 0.095 at the amplitude level.[34] Limits on S, P, T couplings are not very strong;[35] e.g., in μ-decay g_T can be as large as 0.2 for vanishing g_S and g_P. In $\pi \to e\nu_e$ and $\pi \to \mu\nu_\mu$ g_S and g_P can be quite large if they are proportional to m_ℓ. These limits can be made tighter by difficult experiments e.g., improved measurement of e^+ asymmetry in polarized μ-decay (current limit: $\xi P_\mu = 0.975 \pm 0.015$); measurement of helicity of ν_e and ν_μ in $\pi \to e\nu_e$, $\mu\nu_\mu$!

The lifetime of τ should be given by

$$\tau_\tau = \tau_\mu \left(\frac{m_\mu}{m_\tau}\right)^5 B(\tau \to e)$$
$$= (2.81 \pm 0.25)10^{-13} \quad \text{sec.} \tag{17}$$

to be compared to the current experimental value

$$\tau_\tau = (3.31 \pm 0.57 \pm 0.6\ )10^{-13} \quad \text{sec.} \tag{18}$$

This leaves room[24] for ν_τ mixing with (say) a fourth neutrino ν_L at a level of $\varepsilon \sim 0.4$ or a non-zero mass for ν_τ (e.g., M_{ν_τ} of 250 MeV lengthens τ_τ by about 15%).

From the limit on $|U_{cb}| < 0.57$ B mesons should live longer than 8.10^{-15} sec. i.e.

$$1.4 \times 10^{-12}\ \text{sec.} > \quad \tau_B \quad > 8.10^{-15}\ \text{sec.} \tag{19}$$

We know quite a bit about b from CESR data[36] now. In particular
(a) all models in which b decays only semi-leptonically are ruled out;
(b) topless models with b_L, b_R $SU(2)_L$ singlets; b_L, d_L, s_L mixing allowed need $B \to e^+e^-x$ at least at 2.5% level and are now ruled out;
(c) models in which b decays dominantly into $\bar{p}$'s or τ's are ruled out;
(d) a topless model in which b_L is an $SU(2)_L$ singlet but $\begin{pmatrix} c \\ b \end{pmatrix}_R$ is a

doublet, while not pleasing can not be ruled out.[37] The decay pattern is just like the standard model with $U_{ub} \sim 0$. This model would predict $\tau_B \sim 3.10^{-15}$ sec. and a V+A form factor in decays such as $B \to D\pi\ell\nu$.

In the neutral current sector, x_W^2 has been measured in the ν_μ couplings to u- and d-quarks. From the observed strength, $\rho=1$ is tested to about 1.5%. Hence all neutral current couplings are fixed. This means vector and axial-vector couplings of all 12 fermions to each other are fixed (counting g_V and g_A this is a total of 156 otherwise free parameters). How many others have been checked? $\nu_\mu e$, $\bar{\nu}_\mu e$ and $\bar{\nu}_e e$ are consistent with expected values. The coupling of e to u and d was first seen in the celebrated SLAC asymmetry experiment and is also consistent with x_W^2 from ν_μ data. Last year, at PETRA, couplings of e to e, e to μ and also e to u,d,s,c,b were checked to some extent. Two new tests were reported[15] this year: one was the EMC experiment testing μ to u,d couplings and the other is the check of ν_μ to c couplings in the diffractive ψ-production by ν_μ. In view of all these successes, it is sobering to be reminded that[38] $\nu_e - \nu_\mu$ universality in neutral currents is not tested accurately at all and that ξ(the ratio of ν_e to ν_μ couplings) can vary from 1.6 to -0.4.

From the above discussion of B-decays, except for the remote possibility of $\begin{pmatrix} c \\ b \end{pmatrix}_R$ we expect a t quark to accompany b. At what mass do we expect t? One surviving speculation for m_t is

$$m_c/m_t = m_\mu/m_\tau \ . \qquad (20)$$

With this mass formula, making a QCD correction with a gauge invariant quark propagator, with $m_c(m_c)$ taken from analysis[39] of E_e-spectrum in $D \to e\nu x$ to be between 1.74 GeV < m_c < 1.87 GeV, one finds the threshold for naked t to be given by[40]

$$52 \text{ GeV} > 2M_t > 46 \text{ GeV} \qquad (21)$$

If this estimate is wrong and $2M_t$ is below 46 GeV PETRA should find the t-quark in the near future. I think it will be delightful to find t in either case.

4. Aside On D,F Decays

As stressed by Maiani,[24] $D^+ \to \mu\nu$, $\tau\nu$, and $F \to \mu\nu$, $\tau\nu$ are very useful and interesting. For $f_F \sim f_D \sim 0.15$ GeV, B.R.$(D \to \mu\nu) \sim 10^{-4}$, B.R.$(D \to \tau\nu_\tau) \sim 3.10^{-4}$, B.R.$(F \to \mu\nu) \sim 0.1\%$, B.R.$(F \to \tau\nu_\tau) \sim 1.4\%$. Detections of these modes serve variety of purposes. Absolute rates or branching ratios into $\mu\nu$ modes are good f-meters and tell us about the wave functions at the origin. Ratio of rates such as $\tau\nu_\tau/\mu\nu_\mu$ is a very sensitive test[41] of m_{ν_τ} e.g. for F, τ/μ ratio goes from 14 at $m_{\nu_\tau} = 0$ to 1.1 at $m_{\nu_\tau} = 220$ MeV. In case of D, a mere observation of $\tau\nu_\tau$ mode limits $m_{\nu_\tau} < 60$ MeV. Perhaps an attempt to find these modes at ψ'' is not hopeless.

5. CP Violation in Standard Model

In the standard model, setting $\theta_{QCD} \equiv 0$, the phase δ in U_{KM} is the only source of CP violation. Using the bounds on KM parameters one can estimate CP violation in processes besides $K_L \to 2\pi$. The results are disappointing.

In neutral meson systems, analogues of K°-$\bar{K}^\circ$, one defines two parameters: the mixing parameter Δ:

$$\Delta = \frac{(\delta m/\Gamma)^2 + \frac{1}{4}(\delta\Gamma/\Gamma)^2}{2 + (\delta m/\Gamma)^2 - \frac{1}{4}(\delta\Gamma/\Gamma)^2} \tag{21}$$

and the CP parameter ε:

$$\varepsilon = \frac{i\mathrm{Im}M_{12} + \frac{1}{2}\mathrm{Im}\Gamma_{12}}{i/2\delta\Gamma - \delta m} \tag{22}$$

Then in the process $e^+e^- \to M^\circ\bar{M}^\circ \to \ell^+\ell^-, \ell^-\ell^-, \ell^+\ell^+, \ell^-\ell^+$, the observables r and a are related to Δ and ε

$$\begin{aligned} &\text{by} \quad r = (N^{++} + N^{--})/N^{+-} = 2\Delta/(1 + \Delta^2) \\ &\text{and} \quad a = (N^{++} - N^{--})/(N^{++} + N^{--}) \sim 4\mathrm{Re}\varepsilon \text{ (for small } \varepsilon). \end{aligned} \tag{23}$$

For CP violating effect, in this case the asymmetry a, to be observably large, the mixing i.e. r must be large enough so the interesting quantity is ar. To estimate a and r; δm and $\mathrm{Im}M_{12}$ are calculated from box diagram with vacuum insertion and some assumption about the matrix element i.e. for f_D, f_B etc. Similarly $\mathrm{Im}\Gamma_{12}$, $\delta\Gamma$ are estimated from the appropriate absorptive parts of box diagrams. Typical values from a recent calculation[30] are:

System	r	a
D°-$\bar{D}^\circ$	10^{-7} to 10^{-4}	10^{-2} to 10^{-4}
B°-$\bar{B}^\circ$	10^{-5} to 0.7	10^{-2} to 10^{-4}
B°_s-$\bar{B}^\circ_s$	0.5 to 0.9	10^{-4} to 10^{-6}

(24)

Generally for B-systems when r increases a decreases and it is difficult to find ar larger than 10^{-4}. So if a significant charge asymmetry is seen in any neutral meson system it will be from outside the standard model. This is equally true of T_u, T_c systems.

It has been suggested[42] that there might be significant CP violating effects to be found in differences in particle-antiparticle decay rates to CP-conjugate channels. In principle such effects exist but are likely to be small as they need a) at least two channels in the same final state (e.g., $I = \frac{1}{2}$ and 3/2) b) non-zero final state phase-shift difference between these channels and c) non-zero CP-phase difference between the two channels. e.g., $\Gamma(B^- \to \pi^-\pi^\circ) = \to \Gamma(B^+ \to \pi^+\pi^\circ)$ has to vanish while $\Gamma(B^- \to K^-\pi^\circ) - \Gamma(B^+ \to K^+\pi^\circ)$ can be non-zero. One of the more promising possibilities[43] with possible large $\delta\Gamma$ is $B_c \to \pi\bar{D}$. One expects rather small effects ($<10^{-5}$) in hyperon-anti-hyperon decay rate differences.

ε' in K_L decay in Zweig suppressed and expected to be very small. If Penguin graphs play an important role in $K_S \to 2\pi$ then they can also contribute significantly to ε'. Some estimates[44] yield

$$|\varepsilon'/\varepsilon| \sim (1 \text{ to } 4)\, 10^{-3} \tag{25}$$

to be compared to the present limit $|\varepsilon'/\varepsilon| = 0.003 \pm 0.015$. On the other hand it is possible that Penguin graphs do not play an important role in $K_S \to 2\pi$, in that case $|\varepsilon'/\varepsilon|$ may be practically zero.[45]

As shown by Shabalin[46] electric dipole moment of the neutron vanishes up to two loops. At three loops he estimates

$$d_n^e \lesssim 10^{-34} \text{ e cm.} \tag{26}$$

There are other contributions which are estimated[46] to be $\sim 10^{-33}$e cm and optimistically[47] to be $\sim 10^{-30}$ e cm. I think it is safe to say $d_n^e \lesssim 10^{-32}$ e cm for the standard model to be compared to the bound[48] $<4.10^{-25}$ e cm.

Manifestly symmetric left-right $SU(2)_L \times SU(2)_R \times U(1)$ model has interesting implications for CP violation. From β and μ-decay[49] M_{W_R} has to be larger than (2 to 3) M_{W_L}, whereas a more model dependent analysis of K_L-K_S mass difference constrains[50] $M_{W_R} \gtrsim 1$ to 1.5 TeV. The first interesting result is a connection between neutron electric dipole moment and ε' viz.[51]

$$d_n^e \gtrsim 6.10^{-25}\, |\varepsilon'/\varepsilon| \text{ e cm.} \tag{27}$$

Further, in a four flavor model it can be shown that[52] $\varepsilon' \equiv 0$ and this can be generalized by requiring U_{KM} to be real to six flavors.[53] It can be further shown that in this case to reproduce[53,54] ε, W_R cannot be too heavy viz. $M_{W_R} \lesssim 30$ TeV.

In models where CP violation occurs only in the Higgs couplings there are two cases. In class A, neutral Higgs are allowed to couple as they may. Then at least two Higgs doublets are needed. The general features of these models are: neutral Higgs couplings are important and they have to be heavy enough to respect $K_L \to \mu\mu$ rates and light enough to account for $K_L \to 2\pi$. In general $K_L \to \mu e$ will be expected at accessible ($\sim 10^{-11}$ in branching ratio) rates; ε' is nearly 0, neutron electric dipole moment $\lesssim 10^{-30}$ e cm. In general, superweak is mimicked more closely than any other scheme.

In the other Higgs model class B (Weinberg model[55]) flavor changing couplings of neutral Higgs are forbidden and at least three doublets are needed.

In this model charged Higgs play the important role and cannot be too heavy. The model is on the verge of being tested experimentally in its predictions[46,56] of $|\varepsilon'/\varepsilon| \sim 0.01$, $d_n^e \sim 10^{-25}$ e cm and relatively light charged Higgs (current limit from PETRA is $M_H \pm > 15$ Gev). It will be very interesting to watch these bounds being improved in the near future.

6. Potential Problems

There are two experimental results[57] which potentially pose problems for the standard model. One is the persistent indication from various beam dump experiments that fluxes of prompt ν_e's and ν_μ's are not identical. This cannot be explained away within the standard model. Many simple explanations have been ruled out: ν-oscillations are ruled out for the relevant values of L/E, light neutral particles decaying preferentially into $\nu_\mu\bar{\nu}_\mu$ are ruled out, charged Higgs exchange giving higher rate for $c \to s\mu\nu_\mu$ than $c \to se\nu_e$ is ruled out .... In any case a possible violation of e-μ universality in charm decays should be checked by comparing branching ratios for $D \to \mu$ and $D \to e$ directly.

The second is the observed rate for ν_μ induced $\mu^-\mu^-$ production has defied explanation by any known mechanism and remains an outstanding puzzle[57]

I, for one, look forward to many such effects to give us indications of future directions beyond the minimal standard model.

I dedicate this paper to the memory of Jun John Sakurai, friend and mentor.

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MASSES OF $W^{\pm}$ AND Z^{o} BOSONS

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ABSTRACT

In this talk a review is presented of recent work on higher order corrections to the masses of the intermediate bosons W and Z. QCD corrections to order α_s are considered and some of the problems that strong interactions may pose are pointed out. In section I. the standard model and its (many) parameters are discussed. In section 2. one loop corrections are discussed. In section 3. the order α_s corrections are presented. Finally in section 4. a discussion of the hadronic corrections is given.

1. Parameters of the standard model and lowest order results.

If we assume that the standard SU(3) ⊗ SU(2) ⊗ U(1) model (with a simple $I=\frac{1}{2}$ Higgs-sector) describes nature accurately, we are faced with the problem that this model needs very many input-parameters in order to be able to make predictions. To start with we need three coupling constants, one for each gauge group. Furthermore for three families of quarks and leptons we need twelve masses and possibly eight mixing angles. On top of this there are two extra mass parameters, the vacuum expectation value of the Higgs-field. This parameter sets the scale for the masses of the W and Z. Secondly the mass of the physical Higgs-boson is a free parameter, and is needed to define the standard model. Finally we have the QCD vacuum parameter Θ. In total we see that we need 26 parameters. Unfortunately not all of these parameters are known accurately. The Higgs-mass (M_H) for instance is unknown because we have not seen the Higgs-boson yet. On the other hand it is difficult to give a precise meaning to the quark masses, because there is confinement.

The masses of the intermediate bosons (M_W and M_Z) depend on the parameters of the theory. In lowest order they can actually be expressed in only three parameters i.e. the coupling-constants of U(1) and SU(2) and the vacuum expectation value of the Higgs-field $<\phi>$. Again in lowest order we can express these three parameters in terms of three low-energy parameters $(\alpha, G_F, \sin^2\Theta_W)$, in terms of these we have the well-known expressions

$$M_W^{(0)} = \frac{1}{\sin\theta_W}\left(\frac{\pi\alpha}{G_F\sqrt{2}}\right)^{1/2}, \quad M_Z^{(0)} = \frac{1}{\sin\theta_W\cos\theta_W}\left(\frac{\pi\alpha}{G_F\sqrt{2}}\right)^{1/2} \quad (1)$$

A mass prediction of this type can be used to test the model. In a way it plays a role analogous to that of the prediction for the (g-2) of the electron and the muon in QED. At present it is not yet possible to see whether (1) is correct, because we do not know yet what M_W and M_Z are. A more practicle problem is the fact that the experimental value of $\sin^2\Theta_W$ has not yet been determined accurately. As we can see from (1) the uncertainty in $\sin^2\Theta_W$ determines the uncertainty in M_W and M_Z.

Another problem, which is of a purely theoretical nature is posed by the fact that the expressions (1) are lowest order results, which will be modified if loop corrections are taken into account. These corrections have been considered by many people [1)-7)]. I like to stress here that it will be a very important test of the standard model to see whether the relation (1) and the higher order corrections to it will fit the experimental data (which are of course not available yet).

2. One loop corrections.

In order to get an idea of the procedure involved, we may first look again at the anomalous magnetic moment of the electron and muon. It is well known that we can express it in a series of expansions:

$$a_e = 0.5\left(\frac{\alpha}{\pi}\right) - 0.32848\left(\frac{\alpha}{\pi}\right)^2 + 1.211\left(\frac{\alpha}{\pi}\right)^3 \cdots \quad (2)$$

This expression is used to calculate a. The value for α is obtained by a careful measurement of the electronic charge. We may note however that (2) may be used the other way around. If we have a precise measurement of the anomalous magnetic moment, we can deduce from it a value for α. This value of α will be more accurate if we include more terms in the series (2). The situation in the full standard model is more complicated than the one sketched here. The most natural situation would be where we know the masses of all particles and e.g. the coupling constant of electromagnetism. From these parameters we could then compute low energy parameters such as $\sin^2\Theta_W$ and G_F. Since we do not know yet the masses of the W,Z and Higgs-boson, we are forced to work backward (i.e. calculate α from a measurement of g-2). This procedure has been followed in [1] and [2]. A large part of the correction to the Z and W mass is due to vacuum polarisation insertions in the Z and W propagators, where the polarization tensor is calculated for leptons and quarks. Following [1] and [2] we have the following expressions for ΔM_W and ΔM_Z.

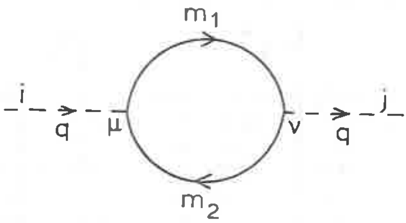

FIG. 1A

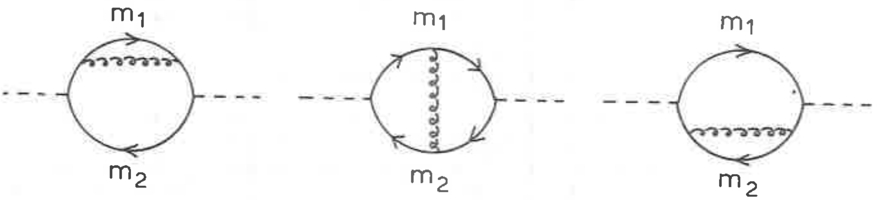

FIG. 1B

Figure 1: Contributions to the vacuum polarisation tensor from fermion loops.

$$\frac{\delta M_W^2}{M_W^2} = \left[\frac{\Pi_{1WW}(q^2)}{M_W^2} + \frac{\cos\theta_W}{\sin\theta_W} \frac{\Pi_{1\gamma Z}(q^2)}{q^2} + \frac{\Pi_{1\gamma\gamma}(q^2)}{q^2} \right]_{q^2=0} - \frac{Re\,\Pi_{1WW}(M_W^2)}{M_W^2}$$

$$\frac{\delta M_Z^2}{M_Z^2} = \left[\frac{\Pi_{1WW}(q^2)}{M_W^2} - \frac{1-2\cos^2\theta_W}{\cos\theta_W \sin\theta_W} \cdot \frac{\Pi_{1\gamma Z}(q^2)}{q^2} + \frac{\Pi_{1\gamma\gamma}(q^2)}{q^2} \right]_{q^2=0} - \frac{Re\,\Pi_{1ZZ}(M_Z^2)}{M_Z^2} \quad (3)$$

The Π_{1ij} is related to the polarisation tensor standing for the self energy graphs (figs. 1a,b). It is defined by

$$\Pi^{\mu\nu}_{ij}(q) = \Pi_{1ij}(q^2)\, g^{\mu\nu} + \Pi_{2ij}(q^2)\, q^\mu q^\nu \qquad i,j = W, Z, \gamma \quad (4)$$

It should be noted that Π_2 does not contibute to the mass shifts. These expressions can easily be evaluated for a single fermion loop and they give, when inserted into (3) an ultraviolet finite result. This result depends of course on masses of quarks and leptons. At this point we are faced with two problems i) What is the mass one should use for the up, down and strange quarks, i.e. the current quark mass or the constituent quark mass? For the c and b quarks we may neglect the difference. ii) What is the mass of the top-quark? We will look at the problem of the light quark masses first. If we take the quark masses in a doublet to be m_1 and m_2 we find in one loop the following approximate expression.

$$\frac{\delta M_W^2}{M_W^2} = \frac{g^2}{48\pi^2}\left[2 \ln \frac{M_W^2}{m_1^2} + \ln \frac{M_W^2}{m_2^2} - 5 \right] \quad (5)$$

and a similar expression for δM_Z. In (5) g is the SU(2) gauge coupling constant. From (5) we see immediately that if we take $m_1, m_2 \to 0$ we will find a diverging mass shift. If we would take current quark masses of a few MeV, (5) would predict large mass shifts [5]. As can be seen from (3) we must evaluate the polarisation tensors at $q^2=0$ as well as at $q^2=M_W^2, M_Z^2$. It is not so clear that perturbation theory will correctly give $\Pi(q^2=0)$.

Since $\Pi(q^2)$ may be given by a dispersion integral and since $\mathrm{Im}\Pi(q^2)$ is described well (in average) if we calculate it by means of perturbation theory using constituent quark masses, it seems the most plausible approach to calculate (5) using for the up and down quark masses 250 MeV [2]. The net results for the mass-shifts due to fermion loops is now given by [2]

$$\Delta M_W = 3.08 \text{ GeV}$$

$$\Delta M_Z = 3.31 \text{ GeV} \quad (6)$$

These corrections also include corrections due to Higgs-boson loops. The Higgs mass was taken to be 200 GeV

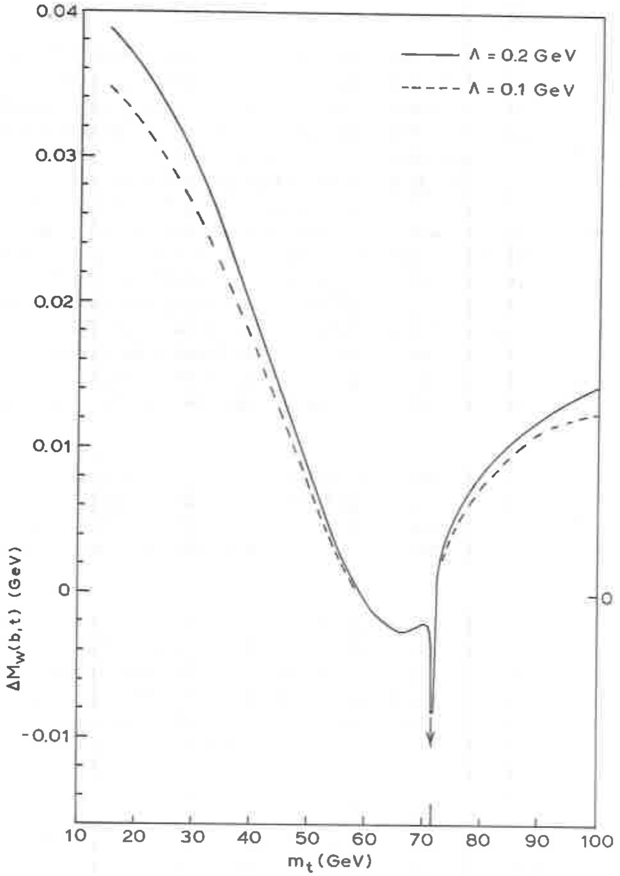

FIG. 2B

Figure 2: The mass-shift of the W to order $_s$

Note the threshold effect.

3. Order α_S corrections.

It should be noted that the effects of strong interactions can entirely be incorporated by using expression (3) and by inserting the full hadronic polarisation tensors. If we assume that these tensors obey dispersion relations, we can try to calculate hadronic effects in two ways. In the first place one could approximate the imaginary part of the hadronic tensor by saturating it with resonances and a continuum. This has been done in [5] and [9]. Another possibility is to use QCD perturbation theory and to make a systematic expansion in powers of α_S. In [10)-11)] the corrections to order α_S have been calculated. This has been done by calculating the order α_S imaginary part of these polarisation tensors. Through a dispersion relation it is then possible to calculate the real part. Since these dispersion relations have to be subtracted once, care should be taken to obey Ward-identities as with dimensional regularisation. The dispersion integrals were evaluated numerically. Details of the subtraction procedure can be found in [10)-11)].

In the results an interesting artefact of the calculation shows for particular values of the top-quark mass. If this mass is such that the mass of the W or Z is close to the $t\bar{b}$ or $t\bar{t}$ threshold respectively the real part of the vacuum polarisation amplitude becomes singular logarithmically.

In figure 2. the order α_S mass shift is given for the W, due to the top-bottom doublet only. Just above 70 GeV a sharp (negative) singularity is seen. This is due to the singularity mentioned above. A solution to this problem can be given by taking into account in some way the full strong interaction effects to all order in α_S. In the real world, the vacuum polarisation tensor does not have the singularity structure as shown in figure 2. Rather, due to the formation of resonances, singularities of this type are absent.

In practice this problem is probably not very important since it would need a fantastic coincidence in the masses of quarks and the W and Z boson masses. It shows however as with light quarks that thresholds cannot always be ignored and that non-perturbative effects can become important.

4. Conclusions.

We have seen so far that it is possible to calculate higher order corrections to the mass of the W and Z bosons. Comparison of the results of these calculations with the physical masses of the W and Z will be a good test of the standard model since the results do not only depend on the tree-level theory, but include higher order corrections as well. From a practical point of view there are uncertainties connected with the values of the quark masses. The light quarks give threshold problems close to $q^2=0$. The heavy quarks may give problems at

$q^2 = M_W^2, M_Z^2$. These problems could possibly be solved by using data in order to get a handle on the structure of the $\pi(q^2)$ tensor.

It is known that also the presence of Higgs-bosons gives corrections to the intermediate boson masses. Their effect has been calculated in [2]. For Higgs-masses smaller than 200 GeV the mass-shifts are of the order of 100 MeV.

We can conclude that the masses of the W and Z are presumably not very good instruments to get information on the Higgs-sector of the theory.

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WEAK NEUTRAL CURRENT EFFECTS IN e^+e^- ANNIHILATIONS

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ABSTRACT

Recent results from PETRA and PEP on weak neutral current effects in lepton pair production in e^+e^- annihilations are presented. Preliminary results from quark pair production are also reviewed.

1. Introduction

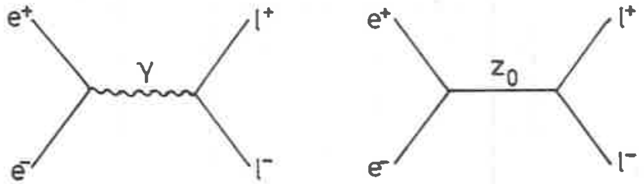
Following the discovery[1] of weak neutral current effects in neutrino scattering experiments in 1973, subsequent neutrino scattering experiments off electrons[2] and nuclei confirmed[3] the existence of the neutral current. Its expected parity violating nature was demonstrated by the polarized electron deuterium scattering experiment[4] at SLAC. These lepton-lepton and lepton-quark scattering experiments, characterized by relatively low spacelike momentum transfers between 10^{-11} and 10^2 GeV^2, indicated a neutral current with contributions from vector and axial vector currents. In brief everything agreed with expectations based on the parton model[5] and the SU(2)×U(1) simple gauge theory of Weinberg and Salam (WS)[6].

The observation of weak neutral current effects in e^+e^- annihilations became a reality with the operation of the PETRA (1978) and PEP (1980) electron positron storage rings. These machines opened up a new kinematic region with timelike momentum transfers of 10^3 GeV^2. The pair production of leptons and quarks, in the later case neglecting fragmentation, provides a very clean way of studying the properties[7] of the neutral current. Further the large timelike momentum transfer allows a search for the propagator effects expected in the WS model.

In this talk I shall present recent results from PETRA and PEP on weak neutral current effects in lepton pair production. Preliminary results from quark pair production are also reviewed. In section 2 brief details concerning the contributing experiments, in particular the integrated luminosities collected and the center of mass, $\sqrt{s}$, are given. This is followed in section 3 by a review of the expectations from the parton and WS models. The radiative corrections applied to all the data are described in section 4. In section 5 and 6 the experimental results from lepton pair production and quark pair production respectively are reviewed along with their implications. Conclusions are given in section 7.

For the data reviewed in this talk both PETRA and PEP operated with unpolarized electron and positron beams which in conjunction with the non measurement of final state fermion polarization result in differential and hence total fermion pair production cross-sections which are parity concerving quantities[15]. The classic signal of a weak interaction - parity violation is lost. With these restrictions what then can we hope to measure?

The first order diagrams for lepton ($\ell = e,\mu,\tau$) pair production via γ and Z_0 exchange are shown in Fig. 1 .



$$M = \frac{-4\pi\alpha}{s}\,\bar{\ell}\gamma_\mu \ell \bar{e}\gamma^\mu e \qquad - \frac{\sqrt{2}G_F M_Z^2}{(s-M_Z^2+iM_Z\Gamma_Z)}\,\bar{\ell}\gamma_\mu(g_V^\ell+\gamma_5 g_A^\ell)\ell\bar{e}\gamma^\mu(g_V^e+\gamma_5 g_A^e)e \qquad (1)$$

Fig.1 First order γ and Z exchange diagrams contributing to lepton pair production.

The differential cross-section would then be,

$$\frac{d\sigma_{\ell\ell}}{d\cos\Theta} = \frac{\alpha^2\pi}{2s}\left\{(1+\cos^2\Theta)\left[1+2g_V^e g_V^\ell \mathrm{Re}(\chi)+(g_V^{e^2}+g_A^{e^2})(g_V^{\ell^2}+g_A^{\ell^2})|\chi|^2\right] + 4\cos\Theta\left[g_A^e g_A^\ell \mathrm{Re}(\chi)+2g_V^e g_V^\ell g_A^e g_A^\ell|\chi|^2\right]\right\} \qquad (2)$$

where Θ is the polar angle between the e^- and ℓ^- and χ is the propagator term

$$\chi = \frac{G_F}{2\sqrt{2}\pi\alpha}\cdot\frac{s}{\left(\frac{s}{m_Z^2}-1+\frac{i\Gamma_Z}{m_Z}\right)} \xrightarrow{s \ll m_Z^2} -(2\cdot 10^{-4})s \qquad (3)$$

The $(1+\cos^2\Theta)$ term in eqn(2) is essentially the familiar QED lepton pair production cross-section plus interference and propagator contributions to the total $\sigma_{\ell\ell}$. The $\cos\Theta$ term represents a charged asymmetry. From WS at $\sqrt{s} \sim 35$ GeV,

$$\sigma_{\ell\ell} \simeq \sigma_{\mu\mu}(1 + g_A^2|\chi|^2) = \sigma_{\mu\mu}(1 + 0.003) \qquad (4)$$

2. Details of contributing experiments.

The experimental results presented here come from the four PETRA experiments CELLO[8)], JADE[9)], MARK J[10)] and TASSO[11)] and from the PEP experiments MAC[12)] and MARK II[13)].

The PETRA experiments have collected most of their data, 50 - 70 pb^{-1} per experiment, in the center of mass energy range 30 - 37 GeV ($\sqrt{s} \sim 34$ GeV). Small amounts of data, 2 - 3 pb^{-1}, have also been collected in the energy ranges, 12 - 14 and 22 - 25.GeV. (CELLO has received considerably less luminosity than the other PETRA experiments due to sharing its interaction region with the PLUTO experiment). The PEP experiments have each collected 20 - 30 pb^{-1} at 29 GeV.

All the groups have similar detectors i.e., large volume magnetic solenoids enclosing charged particle tracking chambers and surrounded by electromagnetic calorimeters, muon identifiers etc. Detailed descriptions of the individual detectors are given in the references above.

3. Model expectations

In this section the implications of the WS model for fermion pair production in e^+e^- annihilations is reviewed. The motivation for this is to introduce some of the ideas which will appear later and to introduce most of the definitions required. The WS model is choosen because of its ability to fit the low q^2 neutrino and electron scattering experiments. The WS model is an SU(2) × U(1) gauge theory whose properties ie, M_Z mass of the neutral gauge boson, Z_0, vector and axial vector coupling constants g_V and g_A, etc are fixed by specifying the mixing angle $\sin^2\Theta_W$. From low q^2 experiments $\sin^2\Theta_W = 0.22$[14)]. This value determines $M_Z \sim 90$ GeV ie well above PETRA and PEP energies (this could have been inferred from measurements of the muon pair cross-section since no deviation from QED is observed). The lepton and quark weak charges predicted by WS are shown in table 1.

Table 1. WS definitions of quark and lepton weak charges

	Q	g_A	g_V
ν_e, ν_μ, ν_τ	0	1/2	1/2
e, μ , τ	-1	-1/2	$-1/2 + 2\sin^2\Theta_W$
d, s , b	-1/3	-1/2	$-1/2 + 2/3 \sin^2\Theta_W$
u, c , t	2/3	1/2	$1/2 - 4/3 \sin^2\Theta_W$

PETRA and PEP are therefore operating in an energy region where the single photon exchange strength has a 1/s dependence and where the Z_0 exchange strength is rising linearly with s (neglecting propagator effects). In this region $\gamma - Z_0$ interference is constant.

The deviation from $\sigma_{\ell\ell}$ is probably too small to measure. The charge asymmetry, A, defined as

$$A_{\ell\ell} = \frac{\sigma_{\ell\ell}(\Theta < 90^0) - \sigma_{\ell\ell}(\Theta > 90^0)}{\sigma_{\ell\ell}} \simeq \frac{-3}{2}(2\times10^{-4})g_A^2 \frac{s}{1-s/M_Z^2} \sim -0.10 \qquad (5)$$

is large and could be measured.

Having developed the discussion for lepton lairs one can extend it to quark pairs. The total quark pair cross-section is then,

$$\sigma_{qq} = \sigma_{\mu\mu}\ 3\sum_q \left\{e_q+2g_V^\ell g_V^q \mathrm{Re}(x)+(g_V^{\ell^2}+g_A^{\ell^2})(g_V^{q^2}+g_A^{q^2})|\ \ |^2\right\} \qquad (6)$$

where e_q is the quark electric charge and 3 the flavour factor. The quark charge asymmetry is,

$$A_{qq} \simeq -A_{\mu\mu}\frac{g_A^q}{g_A^\mu e_q} \qquad (7)$$

The presence of the neutral current will lead to a change in $\sigma_{qq} \sim 1$ %, and give charge asymmetries $A_{d,s,b} \sim -0.30$ and $A_{\mu,c} \sim -0.15$.

4. Radiative corrections.

It has been pointed out that without beam polarization or measurement of final state polarization all neutral current effects in e^+e^- annihilation are intrinsically parity conserving. QED can therefore cause similar effects to those caused by Z_o exchange. For example higher order corrections contribute to $\sigma_{\mu\mu}$ and produce a forward-backward charge asymmetry by interference of graphs with different C parity. In order to extract the neutral current information it is necessary to subtract off the QED radiative correction. All groups working at PETRA and PEP handle higher order (α^3) QED processes by using the calculations of Berends and Kleiss[16)] or something similar. The QED contribution to the muon pair asymmetry evaluated in this way is ~+0.015 for $|\cos\Theta| < 0.8$, compared to the expected asymmetry from electroweak origin of ~-0.10. Although not a large correction it is useful to check the correctness of the calculations. This can be done by comparing the predicted and observed collinearity distribution of muon pairs, Fig. 2, which shows good agreement. A further check has been made by observing the $\mu^+\mu^-\gamma$ final state, this has been done by JADE, Fig. 3, again good agreement is found. In the following review the QED contribution will always be subtracted.

The full radiative corrections to the Z_o exchange graph are not available, first order corrections[17)] including hard photon emission reduce the expected asymmetry by ~+0.01. This reduction is due primarily to lowering the effective center of mass energy of the reaction

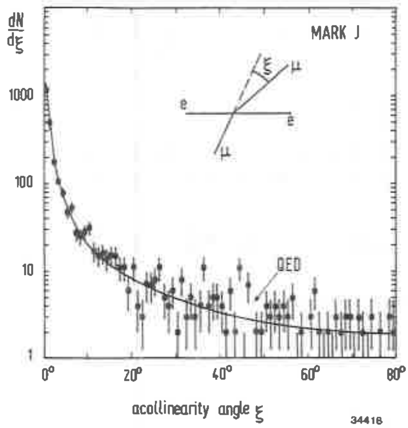


Fig.2 Acollinearity distribution for muon pairs at $\sqrt{s}$ ~ 34.6 GeV measured by MARK J. The acollinearity distribution expected from radiative corrections, QED (α^3), is also shown (full line).

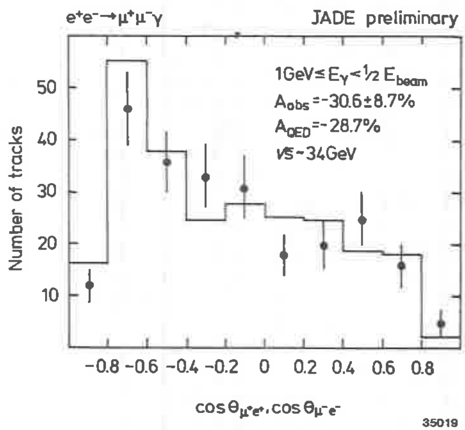


Fig.3 Angular distribution between the outgoing muon and the beam of the same charge observed by JADE in $\mu^+\mu^-\gamma$ events at $\sqrt{s}$ ≈ 34 GeV. The QED (α^3) prediction is shown (full line).

by hard photon emission and hence moving away from the Z_0 pole. This radiative correction is not applied to the results presented here.

5. Electroweak effects in lepton pair production.

5.1 Muon pairs

The criteria for selecting muon pairs are fairly standard. Events are selected containing two reconstructed charged tracks which are collinear to within 10 - 20^{o}. The main backgrounds are from Bhabha events, cosmic ray muons, tau pairs and two photon processes. Time of flight techniques are used to eliminate cosmic rays, Bhabhas and other non muon backgrounds are removed using particle identification by muon identifiers and electromagnetic calorimeters. Two photon induced muon pairs are removed by demanding both tracks to have high momentum (typically more than half the beam energy). The collinearity and momentum requirements reduce the number of tau pairs, where the taus decays leptonically into muons plus unobservable neutrinos, in the sample, but cannot completely exclude them. The residual background contaminations are very small ~2 %.

The presence of a neutral current is not expected to produce large deviations from the QED cross-section until $\sqrt{s} \sim M_Z$. The deviation predicted by WS at PETRA energies has already been shown to be ~0.3 %. For the present data samples available the statistical limitations are ~2 % and the systematic uncertainties ~3 %. The systematic uncertainty is primarily due to the luminosity measurement errors. Under these conditions a test of QED is essentially a test of neutral current contributions. The total cross-sections divided by QED are shown in Fig. 4 for the PETRA experiments. The agreement with QED and hence WS is good.

It has become traditional to measure deviations from QED by cutoff parameters ie,

$$\sigma_{\mu\mu} = \sigma_{QED} \left(1 \mp \frac{s}{s - \Lambda_{\mp}^2}\right) \qquad (8)$$

the 95 % confidence limits for Λ are around 150 GeV. Λ can be re-expressed as limits of g_V as shown in Fig. 4.

Muon pair charge asymmetries have been measured by all groups. Differences in polar acceptances are removed by fitting the measured angular distribution to the form: $1 + a\cos\Theta + \cos^2\Theta$. This is equivalent to extrapolating to the angular region $|\cos\Theta| \leq 1$. The results of each group together with the WS expectation is shown in table 2.

Table 2. Measured Muon Pair Asymmetries

Experiment	$\sqrt{s}$	#Events	$A_{\mu\mu}$ (%)	$A_{WS}^{(\%)}(\sin^2\Theta_W=0.22)$
CELLO	34.2	387	-6.4 ± 6.4	-9.1
JADE	34.2	2224	-10.8 ± 2.2	-9.2
MARK J	34.6	2435	-10.4 ± 2.1	-9.4
TASSO	34.4	2390	-10.4 ± 2.3	-9.3
combined:	34.4	7436	-10.4 ± 1.2	-9.3
MAC	29	1515	-4.4 ± 2.4	-6.3
MARK II	29	652	-9.6 ± 4.5	-6.3
combined:	29	2167	-5.6 ± 2.1	-6.3

The combined muon asymmetry seen at PETRA is -10.4 ± 1.2 % in good agreement with the value expected from WS. The differential muon pair cross-sections measured by JADE, MARK J and TASSO are shown in Fig. 5.

The combined asymmetry, eqn(5) and $|g_A^e|$ = -0.53 ± 0.035 measured[18] from neutrino electron scattering give $|g_A^\mu|$ = 0.53 ± 0.10. Further limits on g_V, g_A and $\sin^2\Theta_W$ will be described later.

5.2 Tau pairs

The selection of tau pairs is complicated by their short lifetime[19] $t_\tau = (3.4 \pm 0.7) \times 10^{-13}$s, and the wide variety of decay products including unobservable neutrinos. This makes event selection and background rejection more difficult than in the muon pair analysis and accounts for the smaller statistics.

The analysis of tau pair results is essentially the same as muon pairs. The cross-section behaviour, Fig. 6, is seen to agree with WED. The charged asymmetries measured are shown in table 3 along with the decay modes analysed.

The combined tau pair asymmetry seen at PETRA -7.9 ± 2.2 % is in good agreement with the value expected from WS. As with the muon $|g_A^\tau|$ = 0.39 ± 0.10 is in agreement with WS. In Fig. 7 tau pair angular distributions measured by the PETRA experiments are shown.

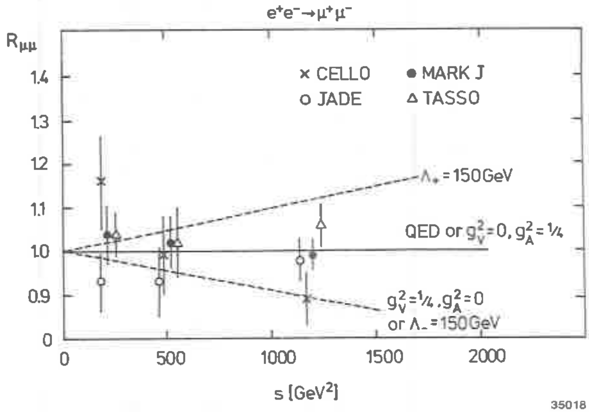


Fig.4 Ratio of $R_{\mu\mu}$ total measured μ-pair cross-section to the QED pointlike cross-section for the PETRA experiments.

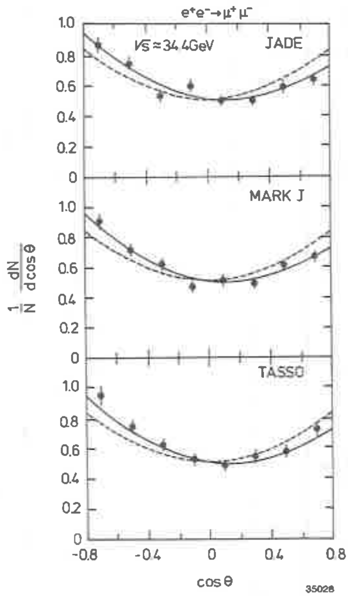


Fig.5 Differential cross-section for muon pair production measured by JADE, MARK J and TASSO.

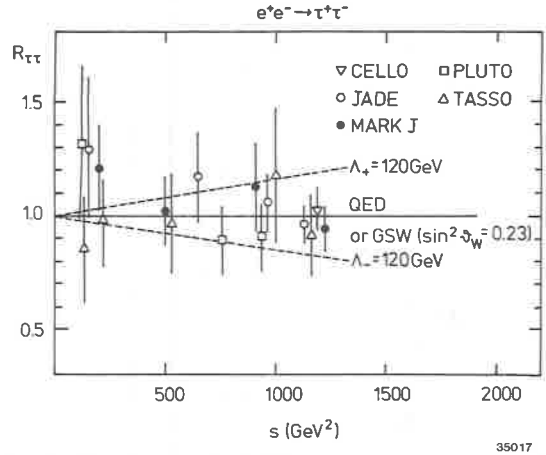


Fig.6 Ratio, $R_{\tau\tau}$, of total tau pair cross-sections measured by the PETRA experiments. PLUTO results from the first period of PETRA running are also included.

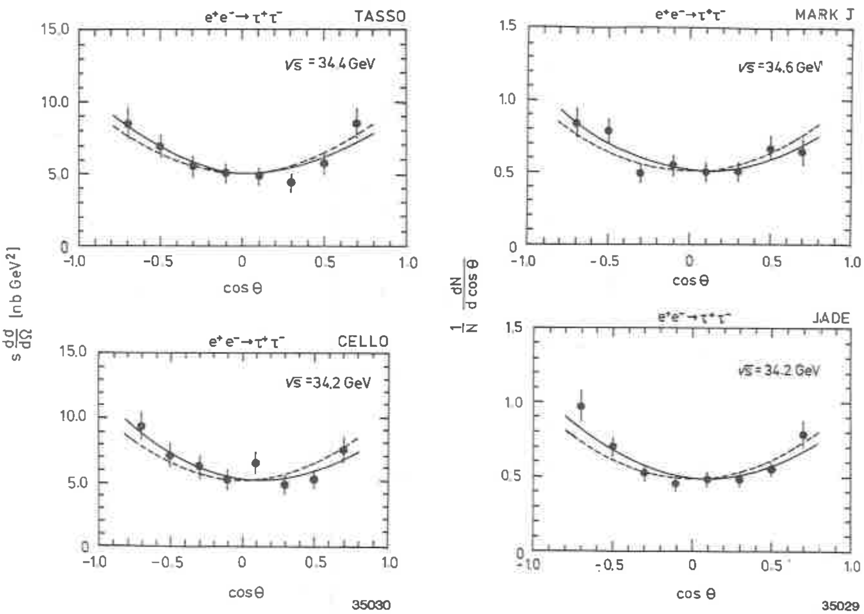


Fig.7 Differential cross-section for tau pair production measured by CELLO, JADE, MARK J and TASSO.

Table 3. Measured Tau Pair Asymmetries

Experiment	$\sqrt{s}$	#Events	$A_{\tau\tau}(\%)$	$A_{WS}(\%)$ $\sin^2\Theta_W=0.21$	Decay Modes
CELLO	34.2	434	-10.3 ± 5.2	-9.2	ALL except ee,$\mu\mu$
JADE (prelim)	34.2	853	-7.9 ± 3.2	-9.2	ALL except ee, $\pi\pi,\mu\pi$
MARK J	34.6	550	-8.4 ± 4.4	-9.4	μX
TASSO	34.4	517	-5.4 ± 4.5	-9.3	1-3
combined:	34.4	2354	-7.9 ± 2.2	-9.3	
MAC	29	-	-1.3 ± 2.9	-6.3	ALL except ee,$\mu\mu$
MARK II*	29	-	-3.2 ± 5.0	-5.0	μX

* $|\cos\Theta| < 0.7$

5.3 Electron pairs (Bhabhas)

The Bhabha behaviour is complicated by the existence of t-channel exchanges. The extra γ exchange diagram produces a large increase in the Bhabha cross-section particularly in the forward direction. This effectively swamps the neutral current effects. The QED normalized Bhabha angular distributions for MARK J and TASSO are shown in Fig. 8, along with WS prediction.

Bhabha events contain only electrons (positrons) thus $g_V^{e^2}$ and $g_A^{e^2}$ can be directly measured from the differential cross-section as opposed to the products of coupling constants measured from muon and tau pair results. The fitted values of $g_V^{e^2}$ and $g_A^{e^2}$ are shown in table 4.

Table 4. g_V^2 and g_A^2 from Bhabha data.

Experiment	g_V^2	g_A^2
MARK J	-0.12 ± 0.34	-0.16 ± 0.54
TASSO	-0.15 ± 0.14	0.01 ± 0.16

Assuming lepton universality and M_Z = 90 GeV the results from the individual leptonic channels can be added and give further constraints on g_V^2, g_A^2 and $\sin^2\Theta_W$, as shown in table 5.

The coupling constants are in excellent agreement with WS and the value of $\sin^2\Theta_W$ is found to agree with low q^2 experiments.

Table 5. Fitted g_V^2, g_A^2 and $\sin^2\Theta_W$ from leptonic channels.

Experiment	Reactions	g_V^2	g_A^2	$\sin^2\Theta_W$
CELLO	ee,μμ,ττ	-0.03 ± 0.08	0.31 ± 0.12	$0.21\,^{+0.14}_{-0.09}$
JADE	ee,μμ	0.05 ± 0.08	0.29 ± 0.06	——
MARK J	ee,μμ,ττ	-0.02 ± 0.05	0.28 ± 0.06	0.26 ± 0.09
TASSO	ee,μμ	-0.04 ± 0.06	0.26 ± 0.07	0.27 ± 0.07
MAC	ee,μμ	——	0.18 ± 0.10	0.24 ± 0.10
MARK II	ee,μμ	0.05 ± 0.10	0.24 ± 0.16	——
combined:		-0.01 ± 0.03	0.27 ± 0.03	0.25 ± 0.05
WS	$\ell^+\ell^-$	0.002	0.25	0.22

The agreement of the leptonic results with WS though impressive are no proof of the correctness of the model. The same results could also be contrived from a four fermion coupling where the couplings agree with those of WS. In order to exclude this possibility it is necessary to observe neutral current effects due to the presence of a non infinite mass exchange boson. This can be done by searching for propagator effects in the charge asymmetry A. By reexpressing eqn.(5) ie,

$$\frac{1}{M_Z^2} = \frac{1}{s} - (2.7 \times 10^{-4})\,\frac{g_A^2}{A_{\mu\mu}} \qquad (9)$$

assuming $g_A = 1/4$ and using the combined PETRA muon pair asymmetry $\frac{1}{M_Z^2} = (1.97 \pm 0.67) \times 10^{-4}$ corresponding to a 95 % limit of $M_Z > 55$ GeV.

The size of the propagator effect predicted from the WS model and that due to the four fermion interaction are shown in Fig. 9.

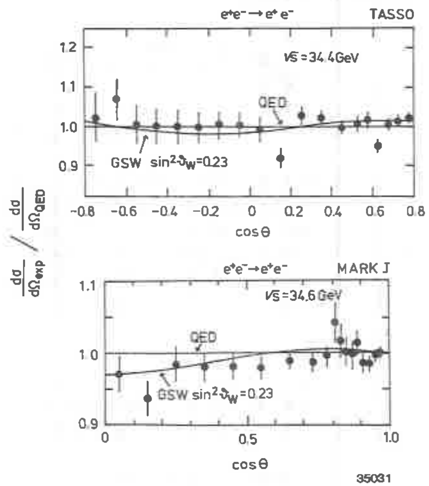


Fig.8 Differential cross-section for Bhabha production normalized by the QED cross-section measured by MARK J and TASSO.

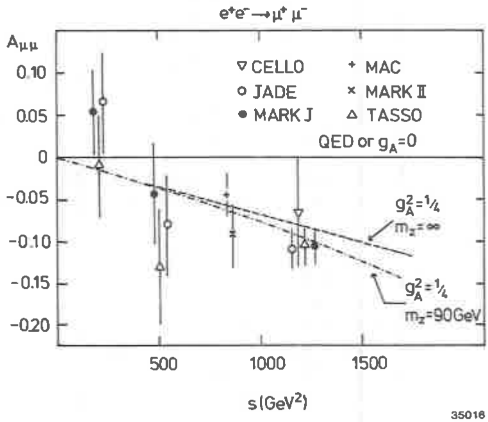


Fig.9 Muon pair charge asymmetry $A_{\mu\mu}$ measured as a function of s by PETRA and PEP experiments. The expected behavious from a four Fermi ($M_Z = \infty$) and WS exchange boson ($M_Z = 90$) contributions are also shown.

Before PETRA and PEP the SU(2)×U(1) model was often extended into richer structure multi boson models by introducing an extra electromagnetic term into the lagrangian,

$$L^{NC}_{eff} = \frac{-4G_F}{\sqrt{2}} \underbrace{\left[\left(J^3_\mu - \sin^2\Theta_W J^{em}_\mu\right)^2\right]}_{WS} + \frac{-4G_F}{\sqrt{2}} \underbrace{c\left(J^{em}_\mu\right)^2}_{extension} \qquad (10)$$

The electromagnetic nature of the extension is necessarily parity conserving and also has no couplings to the neutrino because of its charge neutrality. The addition is therefore hidden from neutrino scattering and polarized electron scattering experiments. At PETRA and PEP these extensions can now be tested most simply by looking at their modification to the vector coupling constant i.e.

$$g_V^2 \to \frac{1}{4}(1 - 4\sin^2\Theta_W) + 4c \qquad (11)$$

The magnitude of c is determined by the mass spectra of the bosons. The PETRA limits on c at 95 % confidence are given in table 6.

Table 6. Upper limits on c

Experiment	c (95% cl)
CELLO	0.031
JADE	0.089
MARK J	0.021
TASSO	0.018
combined:	0.015

The consequences of the combined result for the $SU(2)\times U(1)\times\tilde{U}(1)$[20] and $SU(2)\times U(1)\times SU(2)$[21] double boson models are shown in Fig. 10.

6. Electroweak effects in quark pair production.

With the substantial numbers of hadronic events now available at high energies, for PETRA experiments approximately 20000 events for $\sqrt{s}$ > 30 GeV, it has been possible to search for electroweak effects. In this section I shall present results on the measurement of $\sin^2\Theta_W$ from measurements of the energy dependence of the total hadronic cross-section, $\sigma_{q\bar{q}}$, and preliminary results on the measurement of charm and bottom quark asymmetries.

The selection of hadronic events typically requires the presence of 5 or more charged tracks to be found with a summed momentum greater than 30 % of the center of mass energy. Background from high multiplicity two-photon induced reactions and tau-pair decays are estimated to be 2 and 1 % respectively.

6.1 Determination of $\sin^2\Theta_W$ from the total hadronic cross-section.

The accuracy of published hadronic cross-sections is in general limited not by statistics (~1 %) but by systematic uncertainties which are typically of the order of 10 %. By carefully controlling all systematic uncertainties the JADE[22] and TASSO[23] groups have been able to reduce the overall uncertainties to ~5%. For TASSO this error has the following contributions: trigger acceptance 1.0 %, backgrounds 1.5 %, radiative correction (mostly uncertainties in the contributions from higher order graphs ($>\alpha^3$)) 2.5 %, uncertainties in selection criteria 2.5 %, and uncertainties in the measured luminosity 3.4 %. In quadrature this represents an error of 5.2 %, 4.5 % of which is an overall normalization uncertainty, independent of $\sqrt{s}$, and 2.7 % the point to point uncertainty.

We have seen in section 3 that the presence of a neutral current can give the total hadronic cross-section predicted by the QPM an energy dependence (see Fig. 11). QCD also produces an energy dependence due to contributions from the gluon bremsstrahlung graphs. This contribution has the form,

$$\sigma_{QCD} = \sigma_{\mu\mu}\ 3 \sum e_q^2 \left(1 + \frac{\alpha_s}{\pi} + C_2 \left(\frac{\alpha_s}{\pi}\right)^2 + \ldots,\right) \tag{12}$$

where α_s is the strong coupling constant and C_2 the strength of the second order gluon diagrams. The value[24] of α_s is evaluated from the rate of three jet hadronic events, C_2 from the $\overline{MS}$ scheme[25]. The contribution of the QCD correction to the total hadronic cross-section is ~6 % as shown in Fig. 11.

The energy dependence of the JADE, MARK J and TASSO cross-sections are shown in Fig. 12.

The value of $\sin^2\Theta_W$ is obtained by fixing α_s and fitting the cross-sections over the energy range with $\sin^2\Theta_W$ free. The results of this fits are shown in table 7..

Table 7. $\sin^2\Theta_W$ determined from total hadronic cross-section

Experiment	α_s(fixed)	$\sin^2\Theta_W$
JADE	0.17	0.25 ± 0.05
MARK J	0.17	$0.44\ ^{+0.11}_{-0.19}$
TASSO	0.18	$0.40 \pm 0.15 \pm 0.02$

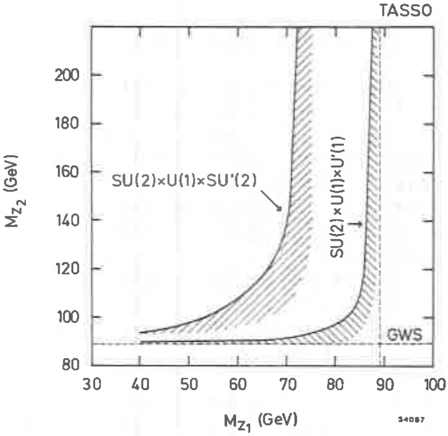


Fig.10 Excluded M_{Z1} and M_{Z2} mass range, non shaded, from TASSO for the extended gauge models SU(2)×U(1)×SU'(2) and SU(2)×U(1)×U'(1).

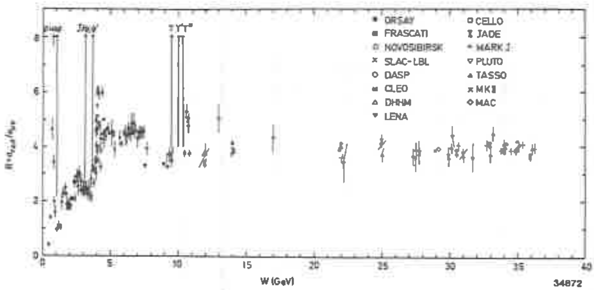


Fig.11 Measured total hadronic cross-section as a function of $\sqrt{s}$(= W) from e^+e^- annihilation. The contribution from QCD is indicated for $\sqrt{s}$ = 35 GeV.

6.2 Quark pair asymmetries.

The requirement of tagging both the charge and flavour has so far restricted the search for charge asymmetries in quark pair production to charm and bottom quarks. I shall describe the two tagging methods so far developed in searching for these asymmetries.

6.2a High P_T muons from b and c quark semileptonic decays

The semileptonic decay of b and c quarks results in a relatively hard muon transverse momentum distribution (P_T). The transverse momentum is measured with respect to the sphericity or thrust axis of the hadronic event. The Monte Carlo predictions for the muon P_T distribution from primary c and b decays are shown in Fig. 13, as well as the P_T distribution expected from hadronic punch through and decay backgrounds. By cutting at high values of P_T it would seem to be possible to tag primary heavy quark events.

According to current ideas on quark hadronization one is unlikely to produce secondary b and c quarks. It is however impossible to separate pure b and c quark samples by this method due to the large overlap in P_T distributions. The best that can be done is the enrichment of the b or c content by moving the P_T cut. The residual contaminations have consequences for the measured asymmetries because a μ^- tags a b or $\bar{c}$ and a μ^+ tags a $\bar{b}$ or c primary quark.

The MARK J and TASSO groups have presented results on heavy quark asymmetries using this method. TASSO requiring the muon to have $P_T \geq 1$ GeV with respect to the sphericity axis. MARK J requiring the angle between the muon and the thrust axis $\Theta_\mu \geq 22^o$. Their results are shown in table 8.

Table 8. Measured c and b quark asymmetries and WS model prediction.

Experiment	$A_{c\bar{c}}$ (%)	$A^{WS}_{c\bar{c}}$ (%)	$A_{b\bar{b}}$ (%)	$A^{WS}_{b\bar{b}}$ (%)
MARK J	-7.5 ± 5.0	-4 ± 1	-17.0 ± 10.0	-5 ± 1
TASSO	—	—	-17.0 ± 10.0	-8

6.2b D* Tagging

The production of charmed mesons, D^*, at high energy has been reported[26] by MARK II, DELCO and TASSO. The TASSO group have used the D* sample to estimate the c quark asymmetry. The technique used is to identify the primary D^{*+} or D^{*-} mesons using the decay,

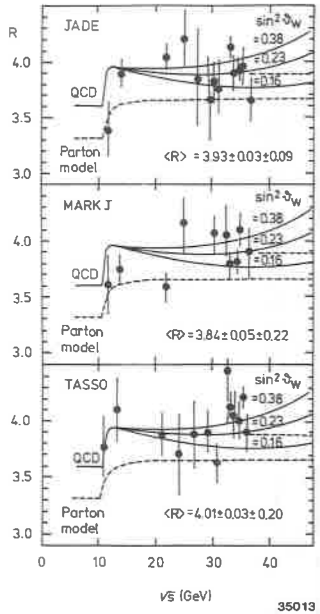


Fig.12 Ratio, R, of the total hadronic cross-section to the pointlike lepton pair cross-section measured by JADE, MARK J and TASSO. The contributions from the QPM, QCD and WS are indicated.

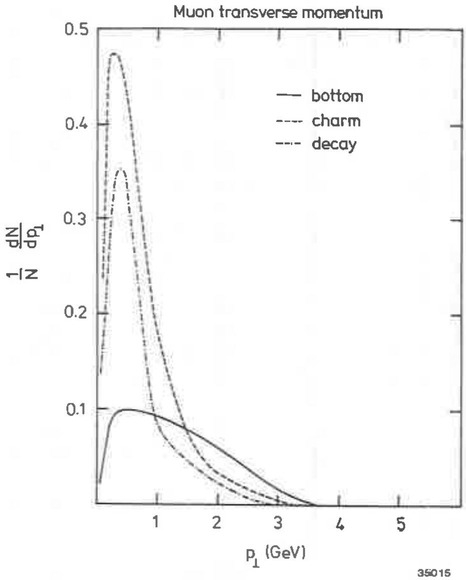


Fig.13 Transverse momentum distributions for muons from bottom and charm decay.

$$D^{*+} \xrightarrow{40\ \%} D^o\pi^+$$
$$\quad\quad \lfloor \xrightarrow{3\%} K^-\pi^+$$

An exclusive final state for the D^o decay is choosen although it only has a ~3 % branching fraction, in order to use the mass resolution from the reconstructed charged tracks to give a signal. No particle identification is required. The mass resolution is improved to $\Delta p/p^2 = 0.0105$ (p in GeV/c) by refitting the tracks to a common production vertex, the D^o lifetime can be neglected. Next the $K^-\pi^+$ mass region 1.8633 ± 0.120 GeV around the measured D^o mass is selected ie, the D^o region. To further reduce combinatorial backgrounds $x > 0.5$ is required for the $K^-\pi^+\pi^+$ system. For $D^{*+} \to D^o$ the distribution $\Delta M = M(K^-\pi^+\pi^+) - M(K^-\pi^+)$ should peak at the known mass difference of the D^* and D^o, ie at 145 MeV. Note that the π^+ produced in the D^* decay can only have a momentum of a few MeV/c, so that although the $D^*(K^-\pi^+\pi^+)$ and $D^o(K^-\pi^+)$ masses have measurement errors of ±70 MeV, the difference between their masses is well measured ±3 MeV. The ΔM distribution for the D^o region is shown in Fig. 14 together with a control region $M(K^-\pi^+) = 2.220 \pm 0.120$ GeV. A clear D^* enhancement is seen. In the range of $0.142 < M < 0.148$ MeV 47 combinations including a background of 11 is found, leaving 36 D^* events.

There is a compilation in the region $M(K^-\pi^+) < 1.8$ GeV where events of the type,

$$D^{*+} \to D^o\pi^+ \qquad \lfloor\!\to K^{*-}\pi^+ \qquad \lfloor\!\to K^-\pi^+\pi^o$$

$$D^{*+} \to D^o\pi^+ \qquad \lfloor\!\to K^-\rho^+ \qquad \lfloor\!\to \pi^+\pi^o$$

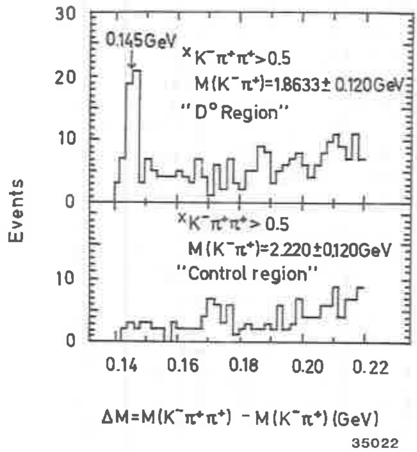
cause an enhancement in the $K^-\pi^+$ mass spectrum around 1.6 GeV. The explanation of non exclusive D^o decays was suggested by the MARK II group.

The average D^* momentum in the TASSO sample is 0.8 of the beam energy hence they are almost certainly due to primary charmed quarks since the momentum distribution from $b \to c$ cascade decays would produce a softer momentum distribution. The c quark asymmetry from this sample is found to be -35 ± 14 % compared with the WS prediction of -14 %. The angular distribution for the D^* events is shown in Fig. 15.

7. Conclusions

The results from experiments at PETRA and PEP demonstrate for the first time the presence of neutral current effects in e^+e^- annihilations in both the lepton and quark pair production channels.

The lepton pair results are found to be in good agreement with the predictions of the Weinberg-Salam model. Specifically the measured axial vector, $g_A^2 = 0.27 \pm 0.03$, and vector, $g_A^2 = -0.01 \pm 0.03$, coupling



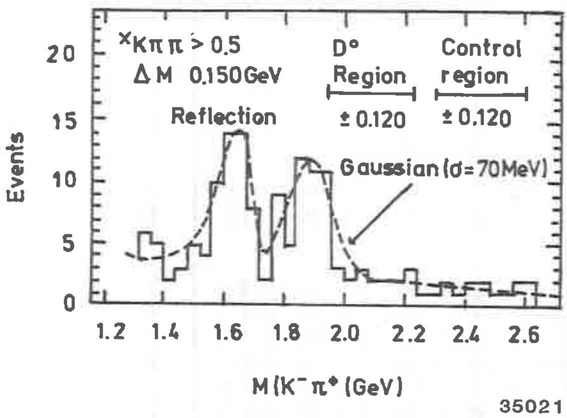


Fig.14 D^{0} region and control region ΔM distribution and $K\pi$ mass distribution in D^{0} selection.

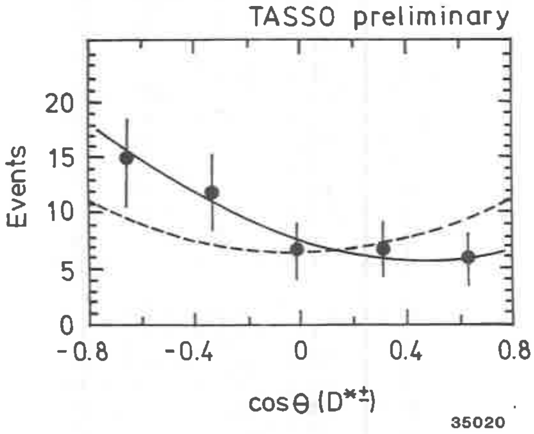


Fig.15 Differential cross-section for D* production.

constants are in excellent agreement with predictions. The lepton results are also in agreement with the results from lower q^2 electron and neutrino scattering experiments when gauged by the measured value of $\sin^2\Theta_W$ of 0.25 ± 0.05. Although the data are consistent with the Weinberg-Salam model they are not yet able to exclude the presence of a four fermion point like interaction. The 95 % confidence limit lower bound on the propagator mass measured from the muon pair asymmetry is M_Z > 55 GeV.

The quark pair results, excluding total cross-section measurements, are confined at the moment to the heavy bottom and charm quarks where charge and flavour tagging techniques have been developed. The results from these channels, although encouraging, require more data to be collected before allowing the measu.ement of quark coupling constants.

Acknowledgements.

I would like to thank all those people at PETRA and PEP who contributed to the results presented in this talk. In particular Drs. J. Branson, A. Böhm, I.C. Brock, R.J. Cashmore, P. Grosse-Wiesmann, H.-U. Martyn, B. Naroska and G. Wolf.

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EXPERIMENTAL RESULTS FROM NEUTRINO-SCATTERING

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1. Introduction

Neutrino scattering experiments provide many informations about the structure of weak currents, the main topic, on which I will report.

The weak interaction can be desribed by an effective Lagrangian of the form:

$$L_{eff} = \frac{G}{\sqrt{2}} \left[J_\lambda^+ \, J_\lambda^- + \rho \, J_\lambda^o \, J_\lambda^o \right] \qquad (1)$$

where J_λ is the "weak isospin" current, which has the following left-handed doublets

$$\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L, \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L, \begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}_L, \begin{pmatrix} u \\ d' \end{pmatrix}_L, \begin{pmatrix} c \\ s' \end{pmatrix}_L, \begin{pmatrix} t? \\ b' \end{pmatrix}_L$$

$J_\lambda^\pm$ represents the weak charged current, J_λ^o the weak neutral current and ρ determines the strength of the neutral current relative to the charged current. At low energies $(s, Q^2 << M_Z^2)$ J_λ^o is well described by the following form (standard model):

$$J_\lambda^o = J_\lambda^3 - \sin^2\Theta_W \cdot J_\lambda^{em}, \qquad (2)$$

i.e. composed of the third component of the "weak isospin" current and the electromagnetic current mixed by the Weinberg mixing angle Θ_W.

In my talk I will report on experiments testing the validity of the (V-A)-theory for charged-current reactions at high neutrino energies and on experiments determining the space-time structure of the neutral current. Many of these results are covered in some recent reviews [1-7].

2. Space-Time Structure Of The Weak Charged Current

The latest experimental evidences for the validity of the (V-A) theory at higher neutrino energies are the following:

2.2.1 μ^+-Polarization In The Reaction $\bar{\nu}_\mu + Fe \to \mu^+ + X$

Measurements of the helicity of muons produced in neutrino reactions can solve the question[8] whether the charged current (CC) is of the V,A (Vector, Axialvector) and/or S,P,T (Scalar, Pseudoscalar, Tensor)-type. Positive muons from antineutrino interactions are expected to have positive helicity if the interaction is of V- and/or A-type and negative helicity if the interaction is of S,P,T-type, because of the known helicity of the incoming $\bar{\nu}$. This was measured in the CERN wide-band beam using the CDHS-detector[9] as target (CERN - Dortmund - Heidelberg - Saclay) and the CHARM-detector[10] (CERN - Hamburg - Amsterdam - Rome - Moscow) as polarimeter for the primary produced muons. For about 17000 muons, stopping in the polarimeter, the helicity was determined by means of the time dependent forward-backward asymmetry of their decay positrons (fig 1). Details of the experimental arrangement can be found elsewhere[11]. This asymmetry can be described by:

$$R(t) = \frac{N_B(t) - N_F(t)}{N_B(t) + N_F(t)} = R_o \cos(\omega t + \phi) + \text{const} \qquad (3)$$

R_o is a product of polarization and analysing power.

The data give $R_o = 0.116 \pm 0.010$ and $\phi = -3.02 \pm 0.08$ compatible with $-\pi$ as expected for positive helicity. This can be translated into a value of the longitudinal muon polarization of $P(\mu^+) = 0.80 \pm 0.07(\text{stat.}) \pm 0.10(\text{syst.})$.

An upper limit of $\sigma_{S,P,T}/\sigma_{tot} < 0.20$ with 95% confidence level on S,P,T contributions to charged-current interactions can be given at an average momentum transfer $Q^2 = 4.0 \pm 0.04$ $(GeV/c)^2$. Fig. 2a shows the amplitude R_o as a function of the inelasticity y; there are no obvious changes with y. By analysing the asymmetry in terms of y an increased sensitivity to S- and P-contributions is obtained. Fig. 2b shows the expected polarization as a function of S- and P-contributions for various bins in y. From the ratio $R_o(y < 0.2)$ to $R_o(y < 0.5)$, which is independent of the analysing power, one gets the result

$$\sigma_{S,P}/\sigma_{tot} \leq 7\% \qquad (95\%\ CL).$$

This experiment shows, that the coupling is mainly of V- and/or A-type.

The same experiment provides a limit for the violation of time reversal invariance. Violation of time reversal invariance would manifest itself by the observation of a polarization component perpendicular to the μ^+-production plane, i.e. by the presence of a $\vec{\sigma}_\mu \cdot (\vec{p}_\nu \times \vec{p}_\mu)$ term. No such term was found. The following limit for the time-invariance violating cross-section for $Q^2 = 4.0$ $(GeV/c)^2$ can be given:

$$\sigma_{T\text{-violating}}/\sigma_{tot} < 16\% \quad (95\%\ CL)$$

2.2 The Inverse μ-Decay

The data on muon decay at low energies are consistent with a mixture of V/A-terms, but it is impossible to derive the relative amount of V- and A-contributions, since only the quantity $(|g_V|^2 + |g_A|^2)$ can be measured. Neglecting possible S,P and T-contributions the differential cross-section for the inverse muon-decay $\nu_\mu + e^- \rightarrow \mu^- + \nu_e$ can be written as[12)]

$$\frac{d\sigma}{dy} = \frac{G^2}{\pi}\,\frac{S}{8}\left[(1+P)\;(1-\lambda)\,y^2 + (1-P)\;(1+\lambda)\right] \tag{4}$$

with the polarization of the neutrino beam

$$P = \frac{N(\nu_R) - N(\nu_L)}{N(\nu_R) + N(\nu_L)} \tag{5}$$

and

$$\lambda = -\frac{2R_e(g_V^{*}g_A)}{|g_V|^2 + |g_A|^2} \tag{6}$$

A value of $\lambda = +1$ would be a proof of the (V-A)-structure of the interaction, and $P = -1$ of left-handed two-component neutrinos. The centre of mass energy of the reaction is low because of the small target mass, but the energy threshold in the laboratory system is 10.9 GeV. Fig. 3 shows the result of two CERN-experiments[13,14)] in terms of upper limits in the P-λ-plane. The results favour the (V-A)-structure of the charged currents and the two-component neutrino theory (left-handed ν).

2.3 Limit On Right-Handed Couplings In Inelastic Neutrino-Nucleon-Interactions

The CDHS-collaboration[15)] has obtained a limit on right-handed couplings using the charged-current reactions

$$\nu_\mu Fe \rightarrow \mu^- X$$

$$\bar{\nu}_\mu Fe \rightarrow \mu^+ X$$

If the current couples to left- and right-handed quarks the cross-sections of these reactions can be written in the form:

$$\frac{d^2\sigma^{\nu}}{dx\,dy} \sim q(x) + \rho^2\bar{q}(x) + (1-y)^2 \left[\bar{q}(x) + \rho^2 q(x)\right]$$

$$= q_L(x) + (1-y)^2\, q_R(x) \tag{7}$$

$$\frac{d^2\sigma^{\bar{\nu}}}{dx\,dy} \sim (1-y)^2 \left[q(x) + \rho^2\bar{q}(x)\right] + \bar{q}(x) + \rho^2 q(x)$$

$$= (1-y)^2\, q_L(x) + q_R(x)$$

where $q(x)$ and $\bar{q}(x)$ are the quark and antiquark structure functions and $\rho = |C_R/C_L|$ is the ratio of right- to left-handed coupling. The different dependence on y for ν- and $\bar{\nu}$-reactions can be used to determine the x-dependence of the structure functions for left-handed and right-handed quarks. The ratio of the cross-sections $d\sigma^{\bar{\nu}}/dy$ to $d\sigma^{\nu}/dy$ is for large y and large x sensitive to the existence of right-handed couplings. Fig. 4a shows the functions $q_L(x)$ and $q_R(x)$ as a function of $\underline{x}$ for a particular Q^2-range. Fig. 4b shows the ratio $(d\sigma^{\bar{\nu}}/dy)$ / $(d\sigma^{\nu}/dy)$ as a function of y for different x-regions. For $x > 0.5$ and $y > 0.66$ and $\langle Q^2\rangle = 33\ (\mathrm{GeV}/c)^2$ the following limit can be derived

$$\rho^2 = C_R^{\,2}/C_L^{\,2} < 0.009 \quad (90\%\ \mathrm{CL})$$

This limit can be used to set limits on the mass mixing of a second charged vector boson which is predicted in a left-right symmetric theory based on $SU(2)_L \times SU(2)_R \times U(1)$. The mixing of the masses is given by

$$C_L = \frac{\cos^2\theta}{M_1} + \frac{\sin^2\theta}{M_2} \quad ; \; C_R = \sin\theta\cos\theta\,\left(\frac{1}{M_2^{\;2}} - \frac{1}{M_1^{\;2}}\right) \tag{8}$$

In Fig. 5 the result is compared with the limits obtained from the muon decay[16)]

2.4 y-Distribution For Prompt Opposite-Sign Dimuon Events

Opposite-sign dimuon events observed in neutrino- and antineutrino interactions are produced by charm-changing charged-current reactions and can be interpreted as being due to the production and the subsequent semileptonic decay of charmed mesons. Assuming only V- and A-contributions in the Lagrangian the differential cross-section $d\sigma/dy$ can be parametrized as

$$\frac{d\sigma}{dy} = A\,\{\alpha + (1-\alpha)(1-y)^2\}\; f\;(y) \tag{9}$$

where f(y) is a threshold factor, which flattens out at higher energies. If the weak charged current is of pure (V-A) type the y-distributions for neutrino- and antineutrino-interactions should be satisfied with a value of α around unity. The CDHS-collaboration[17)] gives the following limits on $(1-\alpha)$, what is equivalent to a limit on (V+A)-contributions:

neutrinos	$(1-\alpha) < 0.10$	95% CL
antineutrinos:	$(1-\alpha) < 0.30$	95% CL

The CHARM-collaboration[18)] quoted the following value for α

antineutrinos: $\alpha = 0.85 \pm 0.10$

Fig. 6 shows some y-distributions from the CDHS-collaboration.

Taking all experimental facts together one may conclude that the weak charged current is well described by a Lagrangian dominated by (V-A)-terms.

3. Space-Time Structure Of The Weak Neutral Current

3.1 Introduction

The Sakurai tetragon[19] (fig. 7) summarizes almost all possible processes, involving neutral-current contributions:

(i) elastic scattering of neutrinos on electrons; coupling constants:
g_A, g_V;

(ii) experiments on neutrino-quark scattering; chiral coupling constants:
u_L, u_R, d_L, d_R as defined by Sehgal[20] or $\alpha, \beta, \gamma, \delta$ as defined by Hung and Sakurai[21];

(iii) reactions induced by electrons, i.e. electron-quark-scattering; coupling constants:
$\tilde{\alpha}, \tilde{\beta}, \tilde{\gamma}, \tilde{\delta}$;

(iV) pure leptonic experiments on e^+e^--storage rings ($e^+e^- \to \ell^+\ell^-, \ell^\pm \hat{=}$ lepton); coupling constants:
h_{AA}, h_{VA}, h_{VV}

(V) reactions induced by muons, i.e. the deep-inelastic scattering of polarized negative and positive muons on carbon performed by the BCDMS - (Bologna - CERN - Dubna - Munich - Saclay)-collaboration[22].

The magnitude of the interference term between parity violating weak and electromagnetic amplitudes observed in longitudinally polarized electron-deuteron scattering[23] rules out a neutral current consisting entirely of S,P and T-terms. The measured asymmetry

$$A = \frac{\sigma_{e_L d} - \sigma_{e_R d}}{\sigma_{e_L d} + \sigma_{e_L d}} = \frac{A_{em} \cdot A_{weak}}{|A_{em}|^2 + |A_{weak}|^2} \gtrsim 10^{-4} Q^2 \qquad (10)$$

is compatible with maximal interference between the weak and electromagnetic amplitudes.

Since in electromagnetic interactions the lepton helicity is conserved, the weak neutral current must consist mainly of helicity conserving parts, i.e. V,A-terms. Small admixtures of S,P and T-terms in the interaction are much harder to rule out.

For the following it is assummed, that the weak neutral current consists only of V- and A-terms.

The two sets of coupling-constants u_L, d_L, u_R, d_R and $\alpha,\beta,\gamma,\delta$ for the neutrino-quark scattering are related to each other by the following definition of the current:

$$\begin{aligned} J_\alpha^{\,q} = \bar{u}\,\gamma_\alpha & \left[u_L(1+\gamma_5) + u_R(1-\gamma_5)\right]u \; + \\ \bar{d}\,\gamma_\alpha & \left[d_L(1+\gamma_5) + d_R(1-\gamma_5)\right]d \end{aligned} \qquad (11)$$

or

$$J_\alpha^{\,q} = \alpha\, V_\alpha^3 + \beta\, A_\alpha^3 + \gamma\, V_\alpha^o + \delta\, A_\alpha^o \qquad (12)$$

where u_L, d_L, u_R, d_R are the left- and right-handed chiral coupling constants of the u- and d-quark, and α, β the vector- and axialvector couplings of the isovector and γ, δ of the isoscalar part in the current. The following relations hold:

$$u_L = \frac{1}{4}(\alpha + \beta + \gamma + \delta) \qquad u_R = \frac{1}{4}(\alpha - \beta + \gamma - \delta)$$
$$d_L = \frac{1}{4}(-\alpha - \beta + \gamma + \delta) \qquad d_R = \frac{1}{4}(-\alpha + \beta + \gamma - \delta) \tag{13}$$

and in the SU(2) x U(1) - model

$s_L = d_L$, $s_R = d_R$, $c_L = u_L$, $c_R = u_R$ (generation universality).

In any model involving a single Z-boson and assuming generation universality, one needs only seven independent parameters to describe all neutral-current phenomena corresponding to the coupling strengths of ν_L, e_L, e_R, u_L, u_R, d_L, d_R, i.e. there are six factorization relations, which connect the coupling constants for neutrino-quark and neutrino-electron scattering with the electron-quark and electron-positron scattering ones. These are[21)]

$$\tilde{\gamma}/\tilde{\alpha} = \gamma/\alpha \quad ; \quad \tilde{\delta}/\tilde{\beta} = \delta/\beta$$
$$g_V/g_A = (\alpha\cdot\tilde{\beta})/(\beta\cdot\tilde{\alpha}) \tag{14}$$
$$h_{VV} = g_V{}^2/\rho \; ; \quad h_{AA} = g_A{}^2/\rho \; ; \quad h_{VA} = g_V g_A/\rho$$

ρ is the relative strength of the neutral-current to the charged current.

Taking these relations into account Kim et al.[22)] determined two years ago the coupling constants in an unique way using all neutrino induced experiments on leptons and hadrons and the experimental result of the SLAC-electron-deuteron scattering experiment. Since then many more precise experiments, especially at the electron-positron storage rings, were carried out.

In the frame work of SU(2) x U(1) Kim et al.[2] made a more generalized analysis by allowing for left- and right-handed isospin multiplets. The coupling constants can then be expressed as:

$$\begin{aligned} \varepsilon_L(i) &= \rho \left[T_{3L}(i) - Q_i \sin^2\Theta_W\right] \\ \varepsilon_R(i) &= \rho \left[T_{3R}(i) - Q_i \sin^2\Theta_W\right] \end{aligned} \qquad (15)$$

where $T_{3L}(i)$ and $T_{3R}(i)$ represent the third component of the weak isospin for the left- and right-handed component of the fermion i with charge Q_i. For left-handed fields the canonical values $T_{3L}(u) = 1/2$, $T_{3L}(d) = T_{3L}(e) = -1/2$ are assumed. In the Weinberg-Salam-GIM model the values of $T_{3R}(u)$, $T_{3R}(d)$ and $T_{3R}(e)$ are identical to zero and ρ equal to one.

Three fits to the data where done:

(i) Within the context of generalized SU(2) x U(1) the weak neutral current is described by the five parameters:

$\rho = M_W^2/M_Z^2\cos^2\Theta_W$, $\sin^2\Theta_W$, $T_{3R}(u)$, $T_{3R}(d)$, $T_{3R}(e)$. A simultaneous fit to the data gives:

$$T_{3R}(u) = -0.010 \pm 0.040 \quad ; \quad T_{3R}(d) = -0.101 \pm 0.058$$
$$T_{3R}(e) = 0.039 \pm 0.047 \quad ; \quad \rho = 1.018 \pm 0.045 \quad ;$$
$$\sin^2\Theta_W = 0.249 \pm 0.031$$

the values of the right-handed components are consistent with zero.

(ii) Set $T_{3R}(u)$, $T_{3R}(d)$, $T_{3R}(e)$ to zero. One is left with a two-parameter fit for ρ and $\sin^2\Theta_W$. The result is $\rho = 1.002 \pm 0.015 (\pm 0.011)$ and $\sin^2\Theta_W = 0.234 \pm 0.013$ ($\pm$ 0.009). The errors in the brackets are due to model-uncertainties.

(iii) If in addition ρ is set to one (Weinberg-Salam-GIM model) one gets a value for $\sin^2\Theta_W$ of

$$\sin^2\Theta_W = 0.233 \pm 0.009\ (\pm 0.005).$$

There is good agreement of all data with the Weinberg-Salam-GIM model.

In the following the more recent experiments and results for the determination of the space-time structure of the neutral-current will be described.

3.2 Electroweak Asymmetry In Deep Inelastic Muon-Nucleon-Scattering

The BCDMS-collaboration[22] studied the interference of the neutral current with the electromagnetic current in the deep inelastic scattering of longitudinally polarized muons on carbon nuclei i.e.

$$\mu_R^- \ C \rightarrow \mu^- X$$

$$\mu_L^+ \ C \rightarrow \mu^+ X$$

They measured the cross-section asymmetry

$$B = \frac{d\sigma^{\mu^+}(-\lambda) - d\sigma^{\mu^-}(+\lambda)}{d\sigma^{\mu^+}(-\lambda) + d\sigma^{\mu^-}(+\lambda)} \tag{16}$$

for energies of 120 and 200 GeV, where λ - the longitudinal polarization of the muon beam - was 81% at 200 GeV and 66% at 120 GeV. An asymmetry arises if the helicity and charge of the incident muons are simultaneously inverted. The asymmetry is due to the parity-conserving part $A_{lepton} \cdot A_{quark}$

of the interaction and in the frame work of gauge models[24) equal to

$$B(Z,\lambda) = K(\lambda\ v_\mu - a_\mu)\ A_o \cdot Z \tag{17}$$

with $K = G/\sqrt{2} \cdot 1/2\,\pi\alpha = 1.8 \cdot 10^{-4}\,GeV^{-2}$ and $Z = (1-(1-y)^2)/(1+(1-y)^2)\ Q^2$ (y = inelasticity) $\cdot$ v_μ and a_μ are the vector and axialvector couplings of the muon to the Z^o, and A_o is a ratio of structure functions and reduces to a combination of axial-vector quark couplings[25) $A_o = 6/5 \cdot (a_d - 2\,a_u)$ and is assumed to be known from other experiments.

Fig. 8 shows the measured asymmetry at 200 and 120 GeV vs. the quantity $(1-(1-y)^2)/(1+(1-y)^2) \cdot Q^2$

In the general gauge-theory the couplings of the muon, allowing for left- and right-handed isospin multiplets, can be expressed as:

$$\begin{aligned} v_\mu &= I_3^{\,L} + I_3^{\,R} + 2\sin^2\theta_W \\ a_\mu &= I_3^{\,L} - I_3^{\,R} \end{aligned} \tag{18}$$

These couplings are related to the slope b of the asymmetry

$$\frac{b}{KA_o} = \lambda \cdot v_\mu - a_\mu = I_3^{\,L}(\lambda - 1) + I_3^{\,R}(\lambda + 1) + 2\,\lambda \sin^2\theta_W \tag{19}$$

The experimental results are:

$b = (-\ 0.147 \pm 0.037) \cdot 10^{-3}\ GeV^{-2}$ at 200 GeV

$b = (-\ 0.174 \pm 0.075) \cdot 10^{-3}\ GeV^{-2}$ at 120 GeV

to be compared with the standard model prediction of $b = 1.51 \cdot 10^{-4}\ GeV^{-2}$.

The Q^2 range was from 15 to 180 GeV^2. The mixing angle of the standard WS/GIM - model derived from the measured slope parameter b is $\sin^2\theta_W = 0.23 \pm 0.07$ (stat) ± 0.04 (syst) using $I_3^L = -1/2$ and $I_3^R = 0$. If they use $\sin^2\theta_W = 0.23$ and $I_3^L = -1/2$ they give a value for I_3^R of $I_3^R = 0.00 \pm 0.06$ (stat) ± 0.04 (syst). This result rules out a neutral heavy lepton of any mass in a right-handed weak isospin doublet with the muon. One can conclude that in agreement with the standard model the muon couples according to universality.

3.3 New Results On Purely Leptonic Processes

3.3.1 Scattering of muon-neutrinos and muon-antineutrinos off electrons

The CHARM-collaboration has performed an experiment in the horn-focussed wide-band beam of the CERN 400 GeV SPS to measure the cross-sections of the processes $\nu_\mu + e^- \to \nu_\mu + e^-$ and $\bar{\nu}_\mu + e^- \to \bar{\nu}_\mu + e^-$ [26] using the same detector. Fig. 9 shows the distributions of the $\nu_\mu e$ and $\bar{\nu}_\mu e$ candidate events versus the variable $E^2\theta^2$ (square of momentum transfer). Shown are also the shapes and magnitudes of the two main background sources; quasi-elastic $\overset{(-)}{\nu}_e$ scattering, and coherent π^0 and γ production by $\overset{(-)}{\nu}_\mu$. After background subtraction and correction for the composition of the beam of the different neutrino types, 42 ± 11 events were attributed to $\nu_\mu e$-scattering and 64 ± 16 to $\bar{\nu}_\mu e$-scattering. After normalization to the total number of charged-current $\overset{(-)}{\nu}_\mu$ interactions the following cross-sections were obtained:

$$\frac{\sigma(\nu_\mu e)}{E_\nu} = \left[2.1 \pm 0.55(\text{stat}) \pm 0.49(\text{syst})\right] \times 10^{-42}\,\text{cm}^2/\text{GeV}$$

$$\frac{\sigma(\bar{\nu}_\mu e)}{E_{\bar{\nu}}} = \left[1.6 \pm 0.35(\text{stat}) \pm 0.36(\text{syst})\right] \times 10^{-42}\text{cm}^2/\text{GeV}$$

The ratio R of the two cross-sections determines the coupling constants of the leptonic weak neutral current with an experimental uncertainty which is smaller than in a measurement of a single cross-section. The experimental value R is:

$$R = 1.37 \begin{array}{l} + 0.65 \\ - 0.44 \end{array} \text{stat.}$$

In the standard model of electroweak interaction R is a function of the weak mixing angle $\sin^2\Theta_W$. R is independent of ρ, the ratio of the over-all strength of neutral-current and charged-current couplings.

$$R = \frac{\sigma(\nu_\mu e)}{\sigma(\bar{\nu}_\mu e)} = 3\,\frac{1-4\sin^2\Theta_W + 16/3\sin^4\Theta_W}{1-4\sin^2\Theta_W + 16\sin^4\Theta_W} \qquad (20)$$

Fig. 10 shows R as a function of $\sin^2\Theta_W$.

From R it was derived:

$$\sin^2\Theta_W = 0.215 \pm 0.04(\text{stat}) \pm 0.015(\text{syst})$$

To determine the parameter ρ the simultaneous measurement of R and $\sigma(\bar{\nu}_\mu e)$ can be used with the result:

$$\rho = 1.12 \pm 0.12(\text{stat}) \pm 0.11(\text{syst})$$

In terms of the coupling constants g_A, g_V the cross-sections define elliptic domains in the g_A, g_V-plane (fig. 11). Including the experimental results [27,28] on the reaction $\bar{\nu}_e + e^- \rightarrow \bar{\nu}_e + e^-$ at the Savannah River fission reactor, two possible solutions are left. To distinguish between both on the basis of purely leptonic reactions one needs the results from the electron-positron storage rings[29,30]. At the e^+e^--storage rings the coupling constant $h_{VV} = \rho\, g_V^2$ is determined by measuring the deviation of the cross-section for the reactions $e^+e^- \rightarrow e^+e^-$, $\mu^+\mu^-$, $\tau^+\tau^-$ from the pure QED prediction and the coupling constant $h_{AA} = \rho\, g_A^2$ is derived from the forward-backward asymmetry of the outgoing lepton. Using $\rho = 1$, as determined from the scattering of muon-neutrinos and muon-antineutrinos off electrons, fig. 11 shows the possible solutions for the coupling constants g_A and g_V. Taking all purely leptonic reactions together the axialvector dominant solution is clearly favoured. The values of the coupling constants are:

$$g_A = -0.52 \pm 0.06 \qquad g_V = 0.06 \pm 0.08$$

i.e. in agreement with the standard model.

3.3.2 Neutrino-Trident Production

The CDHS-[3] and CHARM-collaborations[31] searched for neutrino trident production, i.e. the coherent production of a muon pair and a neutrino in the Coulomb field of a nucleus i.e. $\nu_\mu + Z \rightarrow \mu^- + \nu_\mu + \mu^+ + Z$. This process is allowed by charged-current and neutral-current interaction. Both groups give an upper limit on the diagonal leptonic coupling constant, i.e.

CHARM: $G_D < 1.5\ G_F$ (90% CL)

CDHS: $G_D < 1.6\ G_F$ (90% CL)

The prediction of the standard model is:

$$G_D = (g_L^2 + g_R^2)^{1/2} = 0.77\ G_F \text{ for } \sin^2\Theta_W = 0.23$$

3.4 Semi Leptonic Reactions: Neutrino-Quark-Scattering

To determine the strength and the isospin structure of the coupling, one assumes, in general, V,A-structure and the validity of the quark-parton model.

3.4.1 Inclusive Reactions On Isoscalar Targets

From the inelastic scattering of neutrinos and anti-neutrinos on isoscalar targets the coupling constants g_L^2 and g_R^2, i.e.

$$
\begin{aligned}
g_L^2 &= \rho^2 (u_L^2 + d_L^2) \\
g_R^2 &= \rho^2 (u_R^2 + d_R^2)
\end{aligned}
\tag{21}
$$

can be determined. Especially the Paschos-Wolfenstein-relation

$$R = \frac{\sigma^{\nu}(NC) - \sigma^{\bar{\nu}}(NC)}{\sigma^{\nu}(CC) - \sigma^{\bar{\nu}}(CC)} = g_L^2 - g_R^2 = \rho^2 (1/2 - \sin^2\theta_W) \tag{22}$$

can be used to analyse the data within the Weinberg-Salam model.

The CHARM-collaboration[32)] quoted the following values:

$$
\begin{aligned}
g_L^2 &= 0.305 \pm 0.013 \\
g_R^2 &= 0.036 \pm 0.013
\end{aligned}
$$

or expressing these values in terms of the parameters ρ and $\sin^2\theta_W$:

$$\rho = 1.027 \pm 0.023 \qquad \sin^2\theta_W = 0.247 \pm 0.038$$

using the value of $\rho = 1$ from the relation (22) the following value for the mixing-angle was obtained:

$$\sin^2\Theta_W = 0.230 \pm 0.023 \qquad E_H > 2\ \text{GeV}$$

Fig. 12 shows a comparison of the results of the various experiments on $R = \sigma^{\nu}(NC)/\sigma^{\nu}(CC)$ and $\bar{R} = \sigma^{\bar{\nu}}(NC)/\sigma^{\bar{\nu}}(CC)$ with the Weinberg-Salam model.

A QCD-corrected (Kim et al.[2]) value of the electroweak mixing angle from the cross-section ratios R and $\bar{R}$ of the CHARM-collaboration is $\sin^2\Theta_W = 0.220 \pm 0.014$.

The CHARM-collaboration[33] measured in addition the inelasticity-distributions $d\sigma/dy$ (fig. 13). From those they got compatible results for the coupling constants g_L^2 and g_R^2.

From the deep inelastic neutrino scattering on iso-scalar targets alone only the coupling constants g_L^2 and g_R^2 can be determined. To separate u_L^2, d_L^2, u_R^2, d_R^2 one needs in addition experiments on deep inelastic scattering of neutrinos on neutrons and protons or on semi-inclusive pion-production.

B. Inclusive Reactions On Protons And Neutrons

There are new results from the Bari-Birmingham-Brussels-London (UC)-E.C.Palaiseau-Rutherford-Saclay collaboration[7] on the process $\nu + p \to \nu + X$. The experiment was carried out with BEBC at CERN. They obtain:

$$R_p = \frac{\sigma(\nu p \to \nu X)}{\sigma(\nu p \to \mu^- x)} \sim 2\, u_L^2 + d_L^2 + \text{corrections}$$

$$= 0.49 \pm 0.05$$

The Amsterdam-Bergen-Bologna-Padua-Pisa-Saclay-Turin-collaboration[7)] carried out an experiment with BEBC filled with deuterium. They quote the following values:

$$R^p = 0.48 \pm 0.05$$

$$R^n = 0.26 \pm 0.03$$

$$R^d = 0.33 \pm 0.03$$

From these results they derive the following values for the coupling constants

$$u_L^2 = 0.14 \pm 0.04; \quad d_L^2 = 0.18 \pm 0.05;$$

$$\sin^2\theta_W = 0.210 \pm 0.030$$

Together with results from semi-inclusive pion-production $\nu N \rightarrow \mu\pi X$ also the squares of the right-handed coupling constants can be determined:[34)]

$$u_R^2 = 0.03 \pm 0.01 \qquad d_R^2 = 0.00 \pm 0.01$$

C. Exclusive Neutrino Reactions And Coupling Constants

The experiments so far allow a unique determination of the squares of the coupling constants u_L^2, d_L^2, u_R^2, d_R^2. The relative sign of the chiral coupling constants is chosen by $u_L > 0$. Since $|d_R| \approx 0$, there are four different solutions left corresponding to the two essential sign-ambiguities of the products $u_L \cdot d_L$ and $u_L \cdot u_R$, determining the isospin and V/A-properties of the neutral current.

The following exclusive neutrino-reactions distinguish between the different possible solutions.

(i) Observation of a strong Δ-enhancement in the reaction of single π^o-production, i.e. $\nu p \to \nu p \pi^o$

$$u_L \cdot d_L < o \curvearrowright \text{NC isovector dominant} \quad .$$

(ii) Neutrino-disintegration of the deuteron[35], i.e. $\bar{\nu}_e\, d \to \bar{\nu}_e\, n\, p$. Near threshold it is a pure Gamow-Teller-(isovector-axialvector) transition

$$u_L \cdot u_R < o \curvearrowright \text{NC axialvector dominant} \quad .$$

(iii) Elastic neutrino- and antineutrino-proton scattering. $\overset{(-)}{\nu} p \to \overset{(-)}{\nu} p$.

The results are consistent with the result of (ii), i.e. $u_L \cdot u_R < o$.

(iV) Coherent production of π^o, i.e. $\nu N \to \nu \pi^o N$. The Aachen-Padua-Collaboration[36] reported for the very first time the experimental evidence of this reaction. They give the following cross-section

$$\sigma_\nu = (15.7 \pm 5.1) \times 10^{-40} cm^2/\text{Al nucleus}$$

The reaction constitutes a test of the axial vector part of weak neutral currents, since in the extreme forward direction, vector couplings do not contribute. In the weak coherent interaction of neutral current is:

$$J_\mu^{coherent} = \beta A_\mu^{\;3} + \delta A_\mu^{\;o} . \tag{23}$$

The isovector coupling β describes coherent π^o and the iso-scalar coupling δ describes coherent η^o production. For β they quote a value of 1.3 ± 0.67, which is in agreement with the standard model $(\beta = 1)$. Since they saw no signal for η-production the coupling constant δ is about zero. They quote a value of $\delta < 0.5$ (95% CL). From this experiment again the solution $u_L \cdot u_R < 0$ is favoured. All the experimental data together determine uniquely the coupling constants u_L, d_L, u_R, d_R (α, β, γ, δ). It follows that the weak neutral current is dominantly isovector-axialvector like.

Using the experimental result from the scattering of longitudinally polarized electrons on deuterons[23] and the factorization relations (equ. 14) from the coupling constants α, β, γ, δ values for the leptonic coupling constants g_A and g_V can be derived.

From that experiment[23] the following combinations of coupling constants were determined

$$\tilde{\alpha} + \frac{\tilde{\gamma}}{3} = -0.60 \pm 0.16 \quad \text{and} \quad \tilde{\beta} + \frac{\tilde{\delta}}{3} = 0.31 \pm 0.51$$

Using the factorization relations one gets the following equation:

$$\underset{\substack{\Uparrow \\ (\nu_e)}}{\frac{g_V}{g_A}} = \underset{\substack{\Uparrow \\ (eq)}}{\frac{\tilde{\beta} + \frac{\tilde{\delta}}{3}}{\tilde{\alpha} + \tilde{\gamma}/3}} \cdot \underset{\substack{\Uparrow \\ (\nu_q)}}{\frac{\alpha + \gamma/3}{\beta + \delta/3}} \qquad (24)$$

which connects the neutrino - electron scattering with the electron-quark and neutrino-quark scattering. It turns out, that the allowed region in the g_A, g_V-plane coincides with the result obtained by the purely leptonic processes. The factorization relation is satisfied, if one selects the axialvector dominant solution.

5. Checks On Universality Of The Weak Neutral Coupling

Within the model SU(2) x U(1) the coupling strength of the electron, muon and τ-lepton to the Z-boson should be equal. On the quark side the coupling strength of the u, c and t-quark on the one side and of the d, s and b-quark on the other side should be the same. The following table shows the results on $\sin^2\theta_W$ and ρ for the various types of experiments, showing within the error that universality holds.

Coupling	ρ	$\sin^2\theta_W$	$\sin^2\theta_W(\rho = 1)$
$\nu_\mu q$(CHARM)[32)]	1.027 ± 0.023	0.247 ± 0.038	0.230 ± 0.023
eq[21)]	1.74 ± 0.36	$0.293^{+0.033}_{-0.100}$	0.223 ± 0.015
μq[22)]			0.23 ± 0.07
e^+e^-[30)]			0.259 ± 0.051
$\overset{(-)}{\nu}_\mu e$(CHARM)[26)]	1.12 ± 0.12		0.215 ± 0.040

The experiments at the electron-positron storage rings are testing the generation universality on the lepton side by comparing the results for the reactions $e^+e^- \to e^+e^-$, $e^+e^- \to \mu^+\mu^-$, $e^+e^- \to \tau^+\tau^-$. They are all in good agreement[30)]. Also the electron-quark and muon-quark scattering experiments are a check on the generation universality of the lepton side.

The following experimental results testing the generation universality on the quark side are available:

(i) The Charm-collaboration[33] quoted a value of $g_s^2 = s_L^2 + s_R^2$ which they determined from their y-distributions (fig. 13). They got the result

$$g_s^2 / g_d^2 = 1.39 \pm 0.43$$

i.e. the total coupling strength of the weak neutral current to the strange quark is consistent with being equal to that of the down quark.

(ii) The CDHS-Collaboration[37] reported a result on the J/ψ-production by the weak neutral current. They found in the invariant mass-spectrum of the two muons produced in the reaction $\nu_\mu Fe \rightarrow \mu^+\mu^- X$ (where only little hadronic energy is seen in the final state) a signal (peak) at 3.16 GeV containing 45 ± 13 events. They quote for the mass $M_\psi = 3.16 \pm 0.05$ GeV and a cross-section $\sigma_{diff}(\nu N \rightarrow \nu\psi X) = (4.2 \pm 1.5)\ 10^{-41} cm^2/N$. Their result agrees with the gluon fusion model, i.e. vector- and axialvector coupling is needed. The strength of the coupling of the c-quark to the Z-boson is consistent with the standard model[34]. ($g_c^2/g_u^2 = 2.1 \pm 1.0$)

(iii) At PETRA the cross-sections of the reactions $e^+e^- \rightarrow q\bar{q} \rightarrow$ hadrons are measured and compared with the point-like cross-section. Under the assumption, that all quarks (u, d, s, c, b) take part in the reaction and all quarks have the couplings predicted by the standard model, a value of $\sin^2\theta$ was deduced from JADE[38] of $\sin^2\theta = 0.22 \pm 0.08$ and MARK[39] of $\sin^2\theta = 0.27 \pm^{0.34}_{0.08}$.

The main result is, that at large Q^2 none of the quarks has an unexpected large coupling strength.

6. Conclusions:

(i) Even at higher energies and Q^2 the weak charged current is of (V - A) type: there are still 10 to 20% (V + A) contributions possible.

(ii) The weak neutral current is dominant of isovector-axialvector type. The experiments are in good agreement with the standard model. $SU(2)_L \times U(1)$.

(iii) The coupling of all leptons and quarks to the Z_o is of universal strength.

(iV) The factorization relations, which are only valid in a model with a single Z-boson, seem to hold, but there is a need of more accurate tests.

(V) There is a need of a more precise measurement of the mixing angle $\sin^2\Theta_W$ (possible discriminator for GUT).

Acknowledgements

I like to thank my colleagues from the CHARM-collaboration for their help. I would like to thank Mrs. M. Schmidt for an efficient and critical typing of the manuscript. I would like to express my thanks to Mrs. B. Lohl, who has prepared the drawings.

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Figure - Captions

Fig. 1: Time dependence of relative forward-backward positron asymmetry. The sinusoidal function is the best fit of eq. 3.

Fig. 2a: Asymmetry R_o as a function of the inelasticity

Fig. 2b: Expected polarization for different regions of y as a function of S and P.

Fig. 3: 90% confidence limits on neutrino beam polarization P and the value λ (equations 4-6).

Fig. 4a: Left-handed and right-handed structure functions in a particular Q^2 range.

Fig. 4b: Ratio $(d\sigma^{\bar{\nu}}/dy)/(d\sigma^{\nu}/dy)$ as a function of y for different x-regions.

Fig. 5: Experimental limits on the mixing angle Θ as a function of the mass-ratio.

Fig. 6: y-distribution for opposite-sign dimuon-events from the CDHS-Collaboration.

Fig. 7: Sakurai tetragon or "Mount Elbrus".

Fig. 8: The measured B asymmetry (equ. 16) at 120 GeV and 200 GeV vs $g(y)Q^2 = Q^2 \cdot ((1-(1-y)^2)/(1+(1-y)^2)$.

Fig. 9: Distribution of candidate events for the scattering of a) muon-neutrinos and b) muon-antineutrinos off electrons as a function of $E^2\Theta^2$.

Fig. 10: Cross-section ratio R as a function of $\sin^2\Theta_W$ compared to expectation for acceptance-corrected energy range (full line) and uncorrected energy interval (dashed line)

Fig. 11: Allowed regions in the g_A-g_V-coupling constant plane. (I like to thank Mr. W. Krenz, who updated the compilation of the data and prepared the figure)

Fig. 12: Comparison of various experiments on R_ν and $R_{\bar{\nu}}$ with the Glashow-Weinberg-Salam model.

Fig. 13: Differential cross-section $d\sigma/dy$ after resolution unfolding and acceptance correction.
(a) CC (b) NC

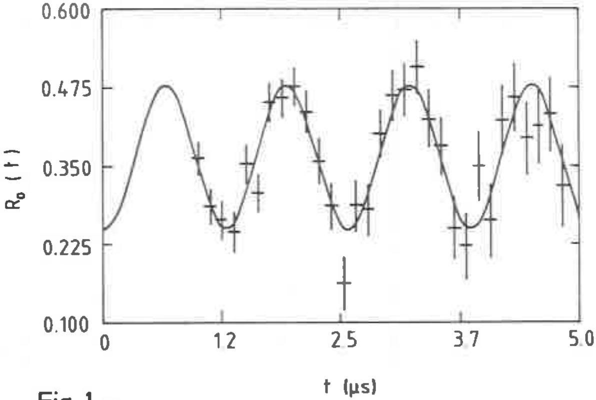


Fig.1

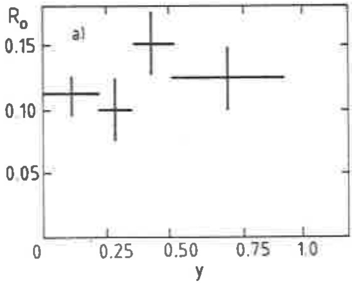


Fig.2a

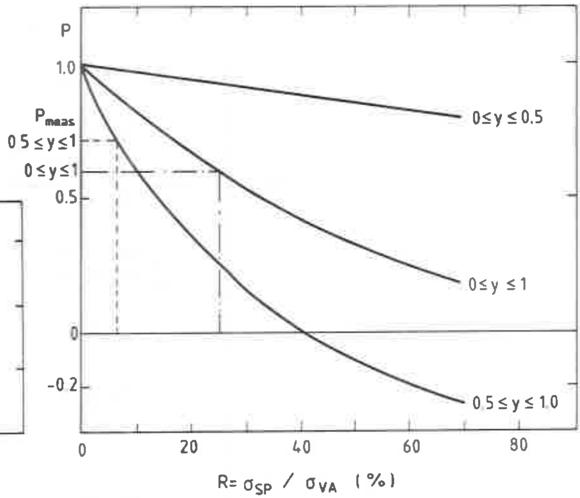


Fig.2b

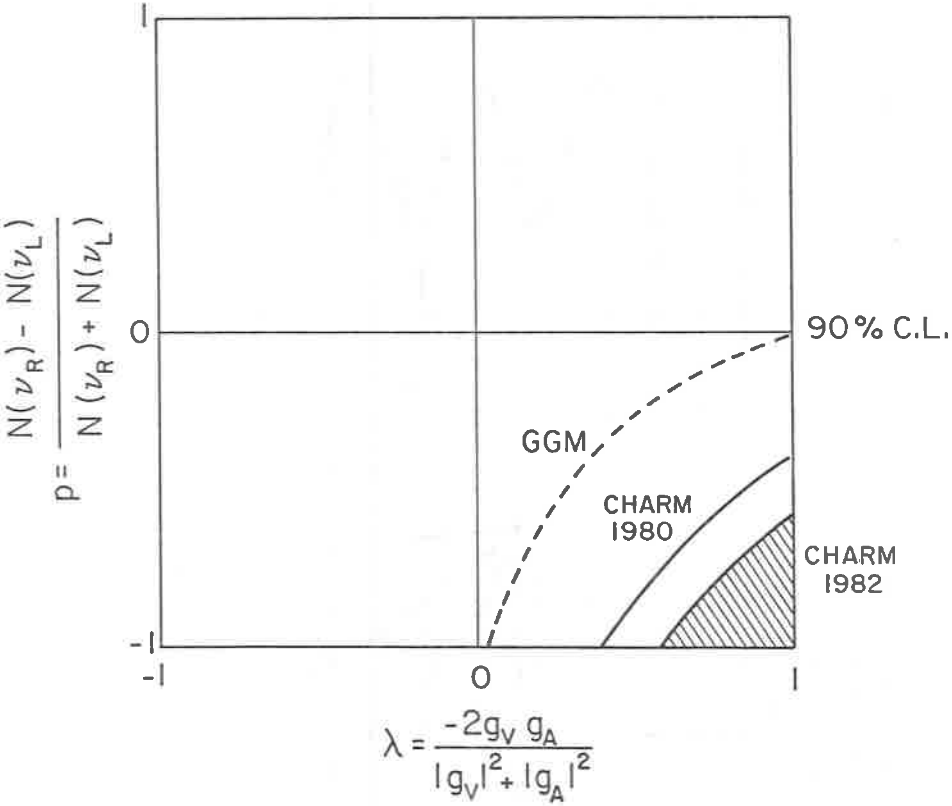

1
0
-1
-1
0
1
90% C.L.
GGM
CHARM
1980
CHARM
1982
$p = \frac{N(\nu_R) - N(\nu_L)}{N(\nu_R) + N(\nu_L)}$
$\lambda = \frac{-2g_V\, g_A}{|g_V|^2 + |g_A|^2}$

Fig. 3

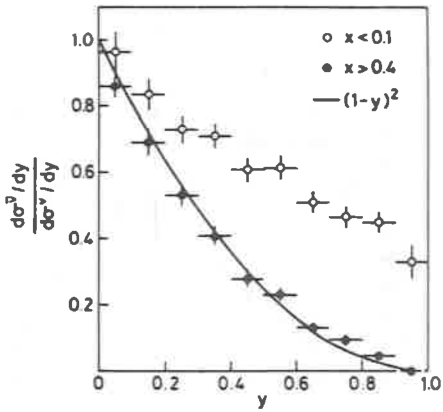


Fig. 4a

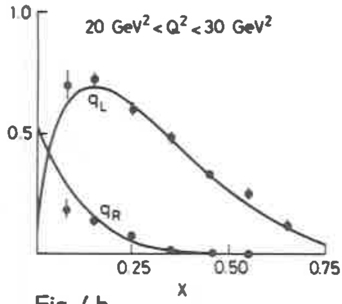


Fig. 4b

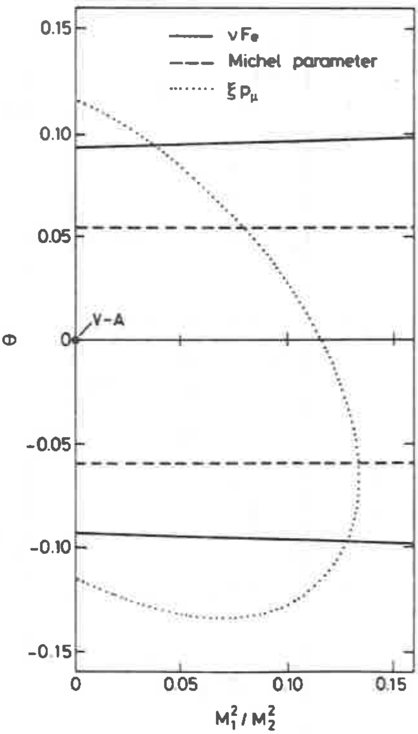


Fig. 5

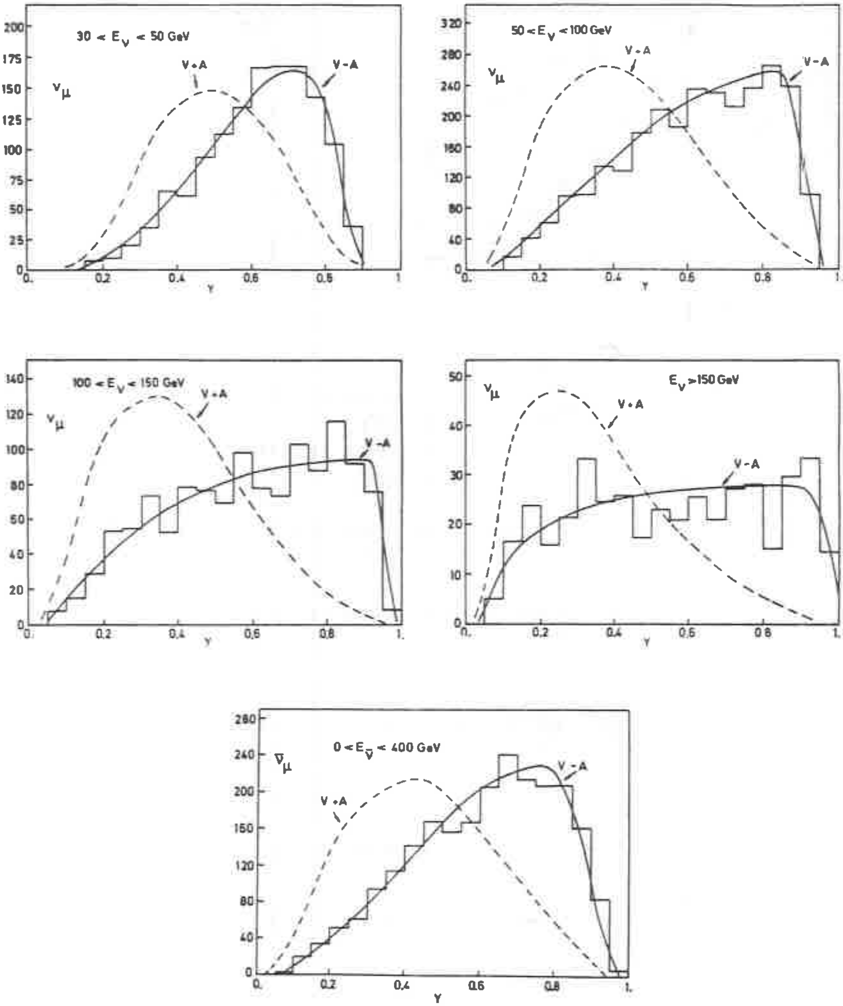
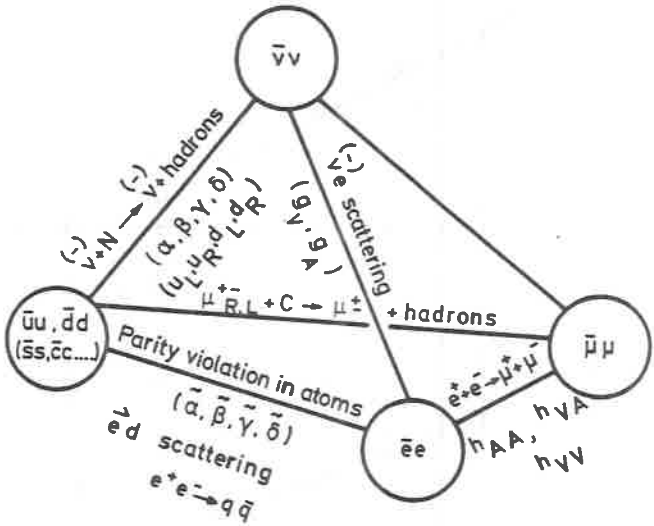

30 < E_ν < 50 GeV
ν_μ
V+A
V-A
50 < E_ν < 100 GeV
ν_μ
V+A
V-A
100 < E_ν < 150 GeV
ν_μ
V+A
V-A
E_ν > 150 GeV
ν_μ
V+A
V-A
0 < E_ν̄ < 400 GeV
ν̄_μ
V+A
V-A
Y

Fig. 6



$\bar{\nu}\nu$
$\overset{(-)}{\nu}+N \longrightarrow \overset{(-)}{\nu}+$ hadrons
$(\alpha, \beta, \gamma, \delta)$
(u_L, u_R, d_L, d_R)
(g_V, g_A)
$\overset{(-)}{\nu}e$ scattering
$\bar{u}u$, $\bar{d}d$
($\bar{s}s$, $\bar{c}c$....)
$\mu^{\pm}_{R,L} + C \to \mu^{\pm}$ + hadrons
$\bar{\mu}\mu$
Parity violation in atoms
$(\tilde{\alpha}, \tilde{\beta}, \tilde{\gamma}, \tilde{\delta})$
$\vec{e}d$ scattering
$e^+e^- \to q\bar{q}$
$e^+ + e^- \to \mu^+ + \mu^-$
h_{VA}
h_{AA},
h_{VV}
$\bar{e}e$

Fig. 7

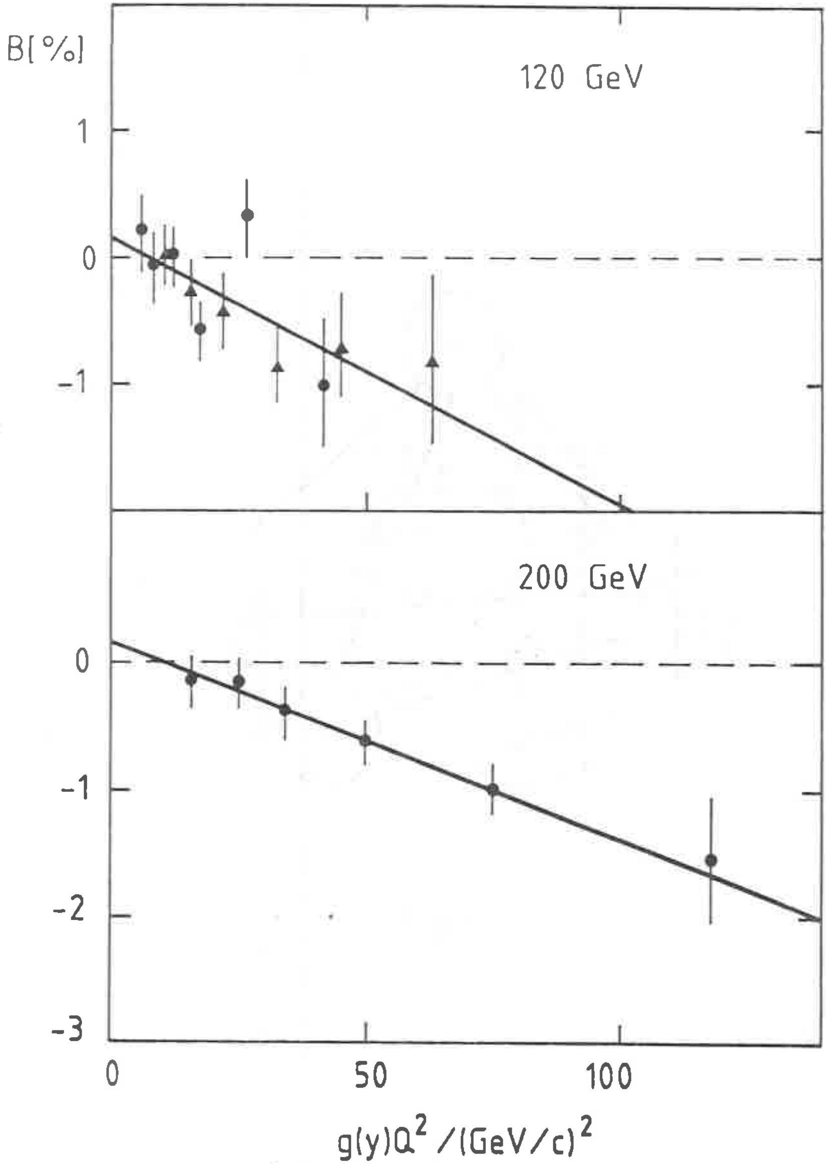
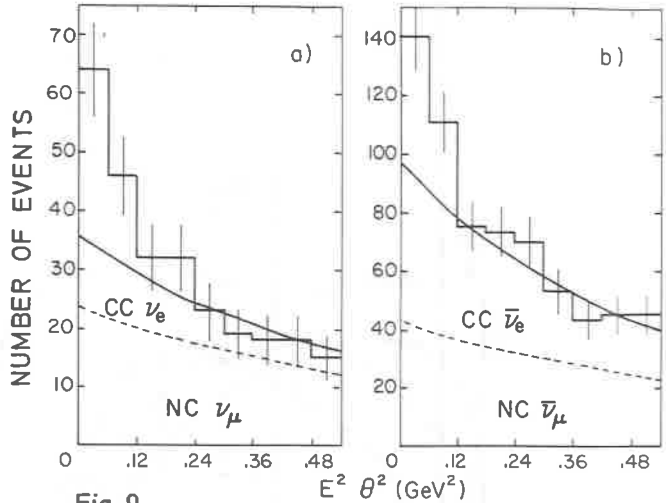
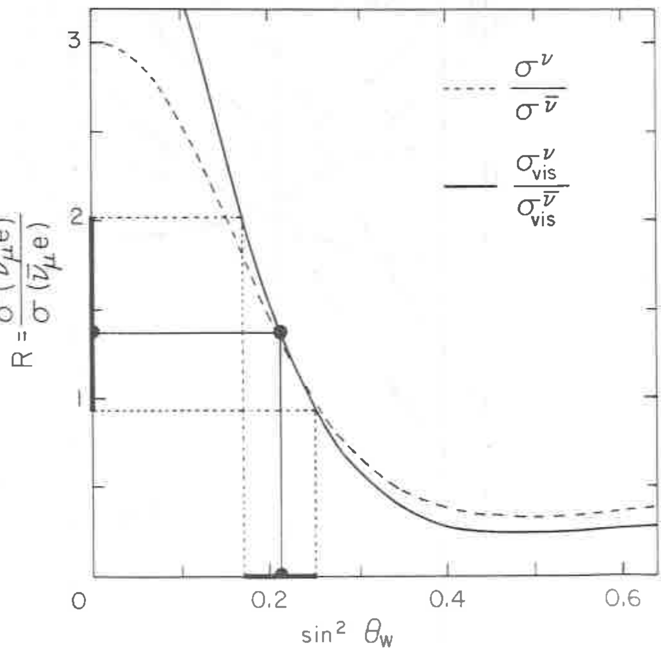


Fig.8



NUMBER OF EVENTS
a)
b)
CC ν_e
NC ν_μ
CC $\bar{\nu}_e$
NC $\bar{\nu}_\mu$
$E^2\ \theta^2$ (GeV^2)

Fig. 9



$R = \frac{\sigma(\nu_\mu e)}{\sigma(\bar{\nu}_\mu e)}$
$\frac{\sigma^{\nu}}{\sigma^{\bar{\nu}}}$
$\frac{\sigma^{\nu}_{vis}}{\sigma^{\bar{\nu}}_{vis}}$
$\sin^2\ \theta_W$

Fig. 10

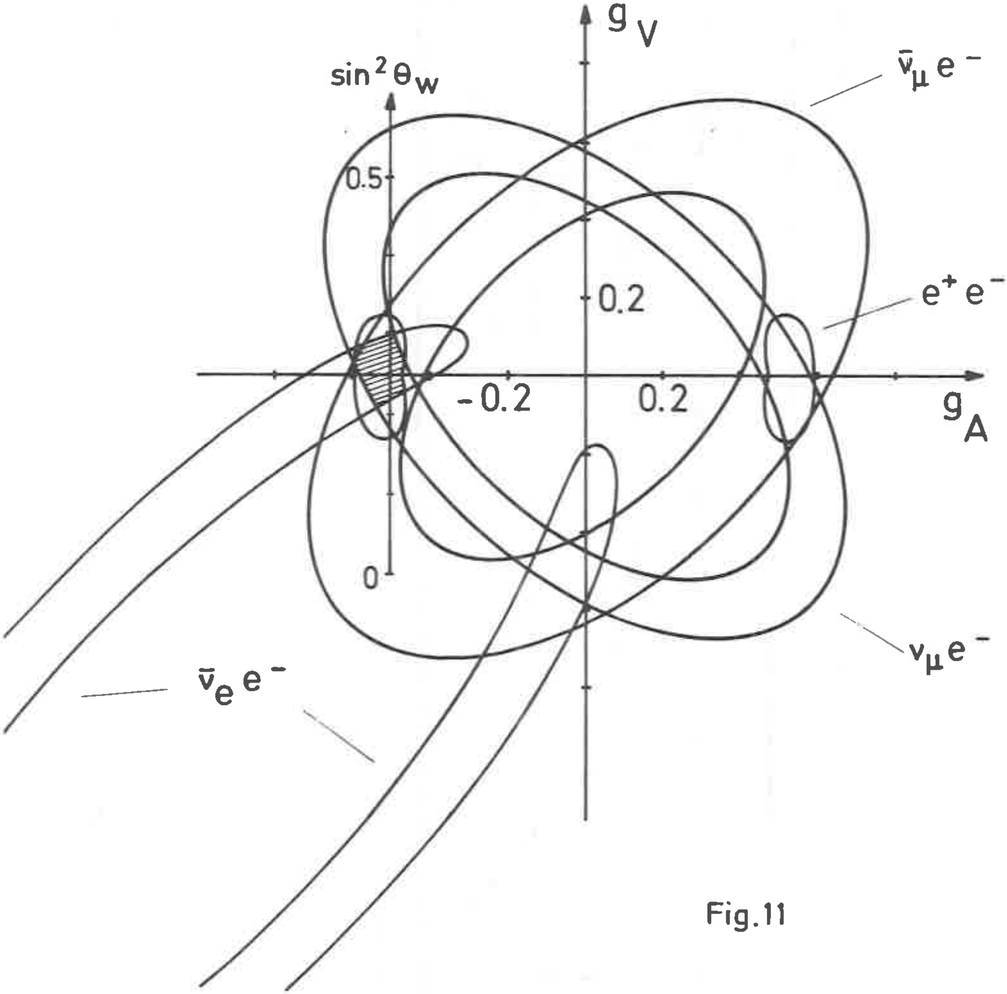


Fig.11

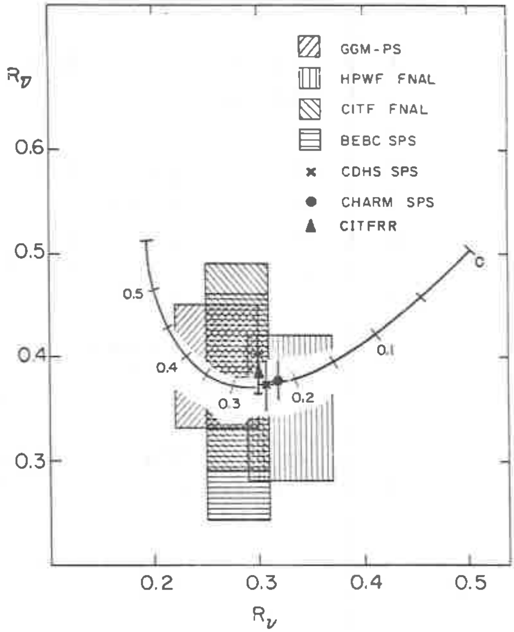

GGM-PS
HPWF FNAL
CITF FNAL
BEBC SPS
CDHS SPS
CHARM SPS
CITFRR
$R_{\bar{\nu}}$
R_ν
0.6
0.5
0.4
0.3
0.2
0.1
c

Fig. 12

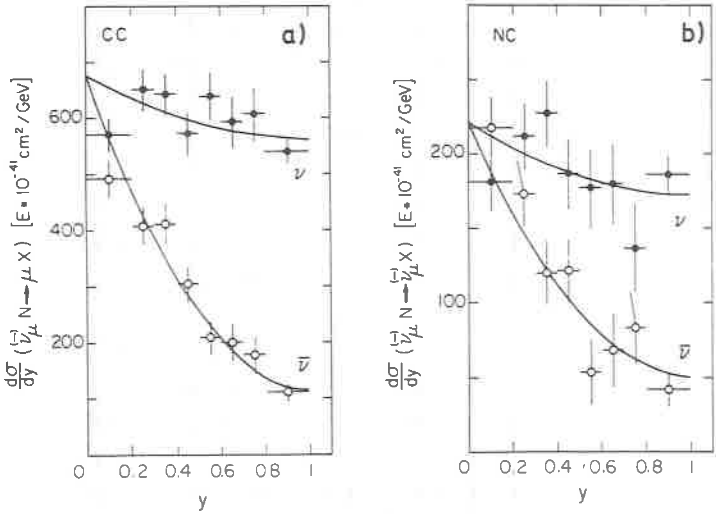

CC
a)
NC
b)
$\frac{d\sigma}{dy}(\overset{(-)}{\nu}_\mu N \to \mu X)$ [$E \cdot 10^{-41}$ cm^2/GeV]
$\frac{d\sigma}{dy}(\overset{(-)}{\nu}_\mu N \to \overset{(-)}{\nu}_\mu X)$ [$E \cdot 10^{-41}$ cm^2/GeV]
ν
$\bar{\nu}$
y

Fig. 13

RECENT RESULTS FROM THE UA1 COLLABORATION AT THE CERN PROTON-ANTIPROTON COLLIDER

Aachen-Annecy (LAPP)-Birmingham-CERN-Helsinki-
Queen Mary College, London-(Collège de France), Paris-Riverside-Roma-
Rutherford Appleton Lab.-Saclay (CEN)-Vienna

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1. Introduction

This workshop takes place at a critical time for the CERN SPS Collider community. In a few days, the UA1 and UA2 collaborations will begin an intensive two month period of data-taking, with expected peak luminosities between 10^{28} and 10^{29} $cm^{2}sec^{-1}$, and strong hopes to produce exciting new results in the field of electro-weak interactions. I will indicate the high level of preparation to study new phenomena by summarising results achieved so far by the UA1 collaboration, using a data sample taken at the end of 1981 at $\sqrt{s}$ = 540 GeV and corresponding to a total integrated luminosity of 22 μb^{-1}. I will discuss the results under three general headings which reflect the highly dynamic nature of this field.

"OLD" Collider Physics (up to summer 1982) covering basic studies of the general nature of $\bar{p}p$ interactions at $\sqrt{s}$ = 540 GeV. I apologise for focusing here, and elsewhere, on UA1 results ; in this case prolific contributions have been made by all collaborations, UA1, UA2, UA4 and UA5.

"NEW" Collider Physics (summer 1982) dealing with existence of jets and preliminary studies of their properties, first highlighted by new results from UA2 and UA1 at the XXI Paris conference in July and at the European Symposium at Santiago de Compostela in August.

"FUTURE" Collider Physics (from end 1982). At the time of writing, the future has already been reached, both for UA1 and UA2, and I can profitably conclude with a brief resumé of the first "electro-weak" results from UA1, derived from the 1982 data.

2. THE UA1 DETECTOR

The UA1 detector has been extensively described in existing literature[1)], and I recall only general features which are relevant to this presentation. The central region of the apparatus is designed to provide detection of W, Z bosons via different decay modes. The interaction area is enclosed in a transverse dipole magnet which can provide a uniform magnetic field of maximum strength 0.7 Tesla. Closest to the interaction is the central detector, a high resolution cylindrical drift chamber complex, 5.8 m long by 2.3 m diameter, which gives a detailed bubble-chamber quality image of each event, and allows momentum analysis of individual tracks with a typical accuracy of ± 20% for a 1 m track at p = 40 GeV/c. The central detector is surrounded in turn by fine-grain lead-scintillator electromagnetic calorimeters and by iron-scintillator hadron calorimeters, the latter incorporated in the return yoke of the magnet. This allows localisation and energy measurement of high energy electrons, identified by their characteristic energy deposition profile, with a typical energy resolution of $\Delta E/E = 0.15/\sqrt{E}$ GeV. High energy muons will penetrate beyond the iron of the magnet, and can be detected by 2 x 4 planes of drift tubes installed around the magnet on all 6 sides.

The calorimetry also permits measurement of global or local hadronic energy deposition, with a resolution $\Delta E/E = 0.8/\sqrt{E}$ GeV, and with a coverage extended down to 0.2^{0} with respect to each beam tube by additional electromagnetic and hadron calorimeters in the forward arms. This capability is important in two respects for W, Z searches, namely the missing transverse energy can be determined per event (characteristic of a neutrino) and high energy jets can be identified and measured (hadronic decay modes, QCD "background").

For the 1981 run, 4 pairs of small precision drift chambers and trigger scintillators were installed in the forward arms at ± 23 m from the interaction vertex to allow a specific analysis of elastic scattering events.

During data-taking, various trigger processors can select events which are primarily electron, muon or jet candidates ; at 10^{28}-10^{29} luminosities the data rate can be then controlled below 1 event/sec with a consequent dead time for rare processes < 10%.

The ensemble of elements for W, Z detection provides in UA1 a universal, near 4π, detector adapted to a wide range of physics topics. In addition to those discussed here, I mention in particular the study of heavy flavour decays, another important objective of future running periods.

3. "OLD" COLLIDER RESULTS

3.1. CHARGED PARTICLE MULTIPLICITIES

The charged particle multiplicity distribution has been studied in rapidity range $|\eta| < 3.5$ with a sample of 8000 minimum bias events recorded without magnetic field[2]. Beam gas events have been removed by timing and topological cuts, and corrections have been applied for track finding efficiency (96%) and geometrical acceptance (80%). Further corrections, estimated by Monte Carlo, have been included for γ-conversion, neutral particle decay and nuclear interaction losses.

The distribution of the average pseudo rapidity density, $dn/d\eta$, is shown in Figure 1, giving a measurement at $\eta = 0$ of 3.3 ± 0.2, compatible with earlier results from UA5 and UA1[3]. As illustrated in Figure 2, this represents a 70% rise from ISR results at $\sqrt{s}$ = 63

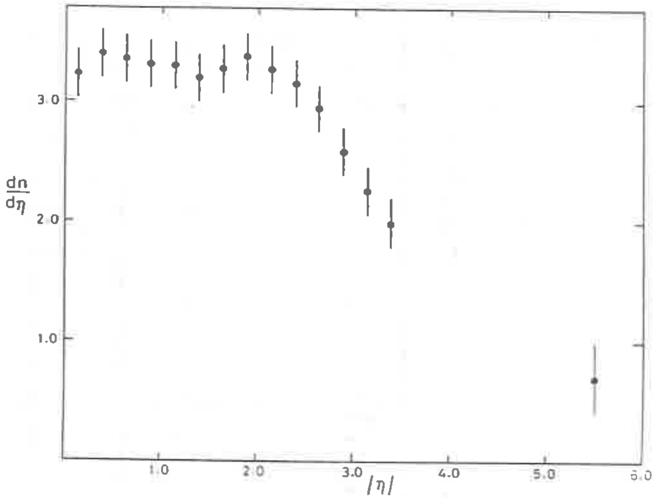


Figure 1. Charged Particle Pseudo-Rapidity Density vs Pseudo-Rapidity.

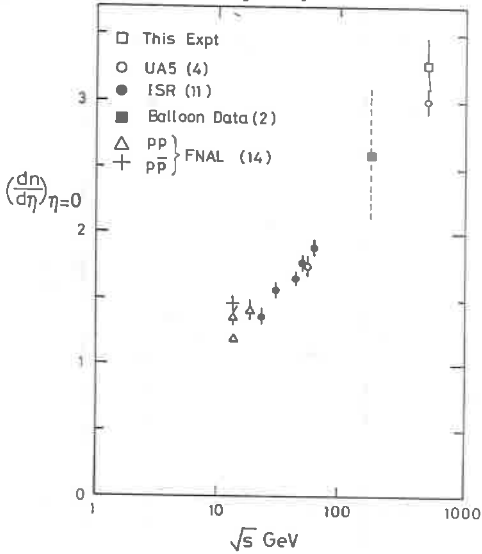


Figure 2. Charged Particle Pseudo-Rapidity Density vs CMS Energy.

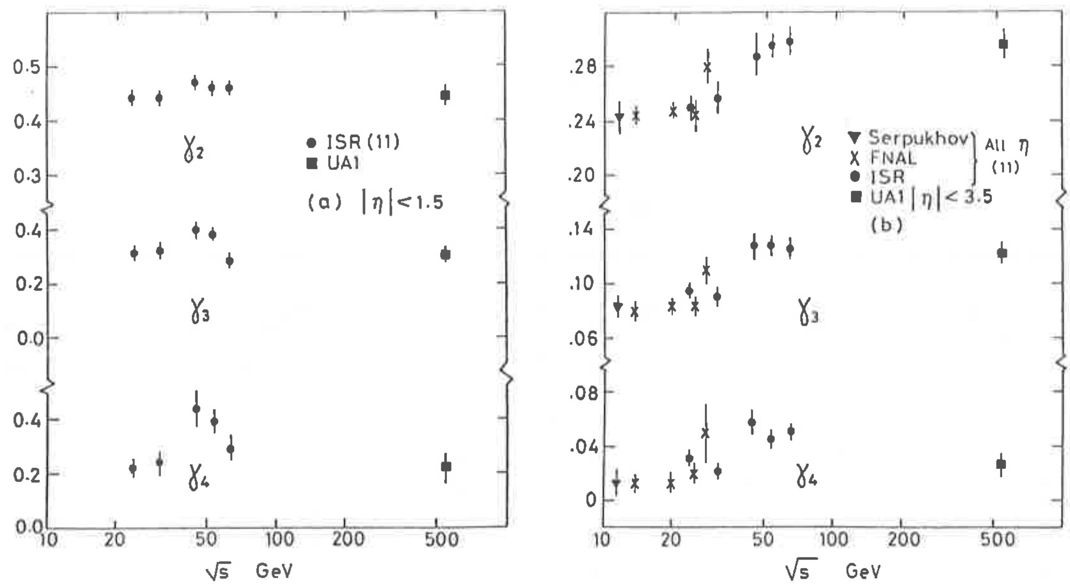


Figure 3. Moments of the Density Distribution of Charged Particles vs CMS Energy.

GeV.[4] The shape of the distribution shows a narrowing of about 2 units with respect to a simple extrapolation from ISR energies. Analysis of the moments of the distribution allows comparison with other energies. Figure 3 shows the energy dependence of the following three moments :

$$\gamma_2 = \frac{<(n-<n>)^2>}{<n>^2}$$

$$\gamma_3 = \frac{<(n-<n>)^3>}{<n>^3}$$

$$\gamma_4 = \frac{<(n-<n>)^4> - 3<(n-<n>)^2>^2}{<n>^4}$$

Little change has occurred between ISR and collider energies, i.e. the moments are independent of $\sqrt{s}$, implying that KNO scaling is approximately valid over this very large energy range.

3.2. SMALL ANGLE ELASTIC SCATTERING

$\bar{p}p$ elastic scattering has been studied in the range $0.14 < |-t| < 0.26$ GeV^2 using a sample of 5437 elastic trigger events obtained at a luminosity of $\sim 10^{26}$ $cm^2 sec^{-1}$ with normal beta operation of the SPS[5]. The detector was arranged to record elastic scattering events in the vertical plane, and approximately 20% of the triggers gave a clean coincidence of "up-down" or "down-up" single tracks in the respective forward arms. The remaining triggers contained large splashes in one arm, partly due to accidental coincidence of beam-gas secondaries and partly to inelastic beam-beam events.

The acceptance of the apparatus has been studied in detail by Monte Carlo simulation. The full acceptance regions of the up-down and down-up categories have been combined for the final t distribution shown in Figure 4 which is well described by the hypothesis $dN/dt \propto e^{bt}$ with $b = 13.3 \pm 1.5$ GeV^{-2}. The data is presented as a differential cross-section, involving scale factors with systematic uncertainties, the most important being 30% for the integrated luminosity measurement

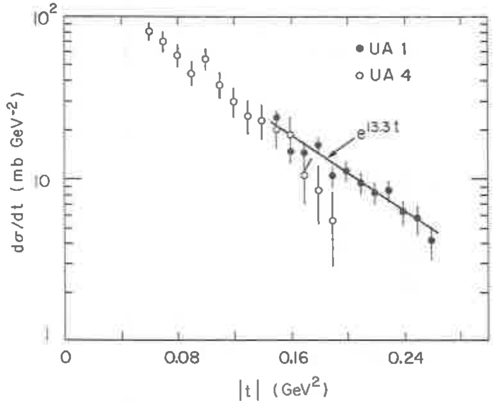


Figure 4. Elastic Differential Cross-Section.

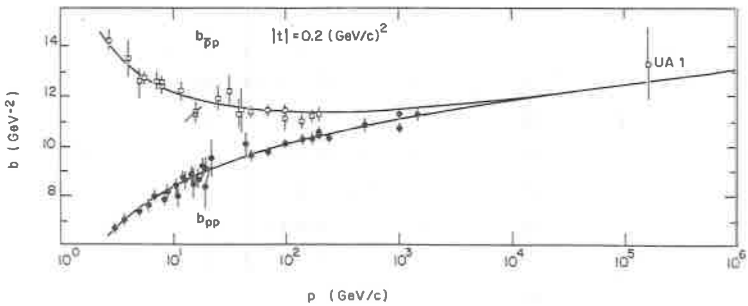


Figure 5. Elastic Slope Parameter vs CMS Energy.

of 2.6 x 10^{30} cm^{2}. Figure 4 also shows the UA4 data[6] over the range $0.05 < |-t| < 0.19$ GeV^{2}, which give a measured slope parameter $b = 17.2 \pm 1.0$ GeV^{-2}. The combined results suggest that b increases with decreasing t, as observed in pp scattering at the ISR[7].

A comparison with the compilation of Burq et al.[8] in Figure 5 indicates that the forward elastic $\bar{p}p$ peak is shrinking over the ISR-Collider energy range.

3.3. SEARCH FOR CENTAURO EVENTS

The expected characteristics of Centauro events, reported from Cosmic Ray experiments[9], are :

- high hadron multiplicity
- negligible fraction of electromagnetically showering particles
- high $<p_t>$ per particle (above 1 GeV/c).

It is natural to search for similar features in collider events where the equivalent energy on a stationary target is 155 TeV, always admitting a possible higher threshold (typical Centauro energies are estimated at around 1500 TeV).

Interpretation of Centauro events in the collider context is difficult due to various uncertainties. One concerns the K_γ factor ie. the fraction of hadronic energy converting electromagnetically, which depends on the pion/nucleon population. Another concerns the accuracy of the vertex height determination which directly affects measurement of $<p_t>$. The expected kinematics in the collider is also dependent on the underlying mechanism, being very different for the two popular hypothesis, namely collision with an air nucleon or decay of a massive object.

With such reservation in mind, we have looked for equivalent dramatic behaviour in a sample of 48000 minimum bias events[10], recorded with a magnetic field of 0.56 Tesla. A reasonable measure of the overall e.m./hadron composition is given in the central calorimeters by the respective energy depositions in the first 4 radiations lengths (e.m.) and beyond 12 radiation lengths (hadronic).

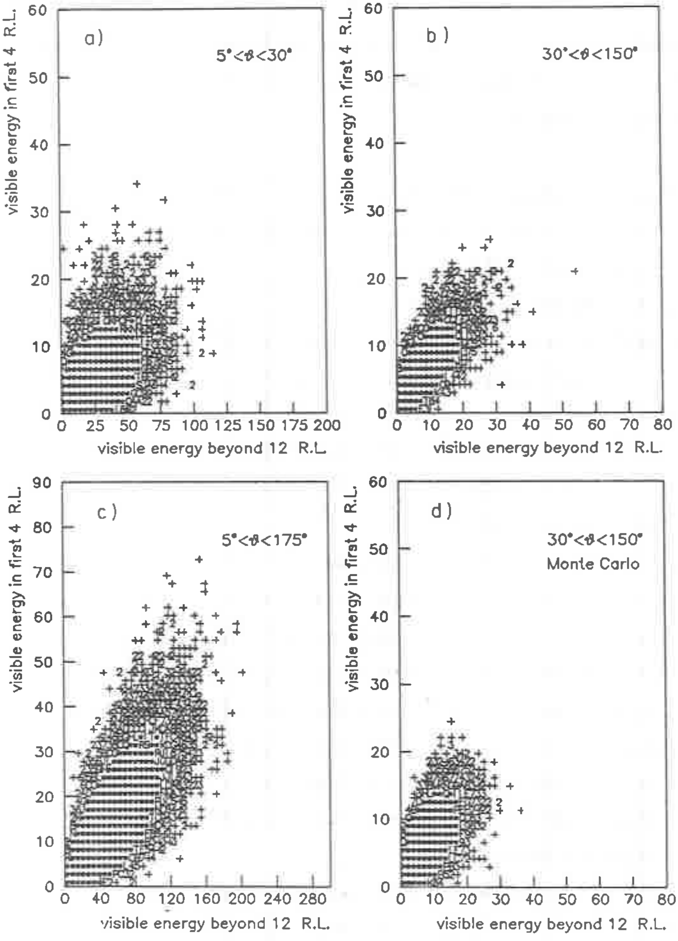


Figure 6. "Electromagnetic" vs "Hadronic" Energy.

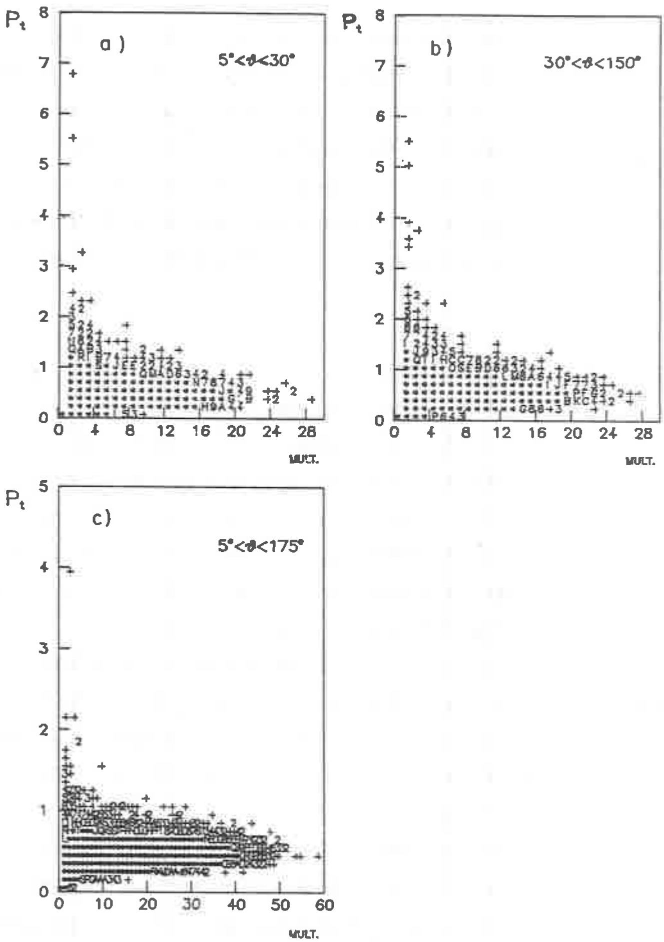


Figure 7. Mean Transverse Momentum vs Charged Multiplicity.

Figure 6 shows a smooth behaviour in different relevant angular regions, without any clearly anomalous events, and consistent with a Monte Carlo simulation. The same is true for the mean transverse energy of charged tracks as a function of their multiplicity, shown in Figure 7.

A more detailed study has been performed on Centauro 1, the best determined of the family of (currently) 6 events. Centauro 1 has been transformed to the collider cms under various K_γ and production mechanism hypotheses. There is no concentration of collider events in the predicted regions, which are nonetheless fully accessible.

With the qualifications mentioned above, we can give an upper limit for Centauro-like processes at 155 TeV equivalent energy of around 1 μb.

3.4. CHARGED PARTICLE p_t SPECTRA

The inclusive transverse momentum spectrum of charged particles in rapidity range $|y| < 2.5$ has been measured up to 10 GeV/c[11], using the sample of 48000 minimum bias events taken with a magnetic field of 0.56 Tesla. Track momenta were reconstructed in the central detector which has a typical resolution ∿ 280 μ in the drift plane and ∿ 2.3% of wire length for the transverse coordinate given by current division. The track finding efficiency was ∿ 97%. Corrections have been applied for geometrical acceptance and for smearing due to resolution. Contamination near p_t = 10 GeV/c is estimated as < 17% due to bad track reconstruction and < 30% due to particle decays. At this stage only about one third of the wires were equipped with electronics.

Figure 8 compares the UA1 inclusive spectrum with ISR results[12], showing an increase of about 3 orders of magnitude in cross-section at p_t= 10 GeV/c. Both the ISR and collider spectra are in good agreement with QCD predictions by Odorico[13]. The UA1 spectrum is also well described by a simple empirical from :

$$Ed^3\sigma/ = Ap_{t0}^n/(p_t + p_{t0})^n$$

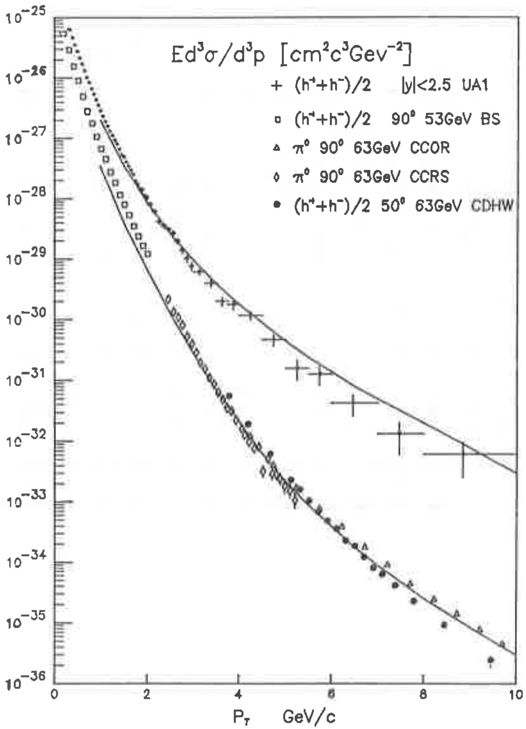


Figure 8.

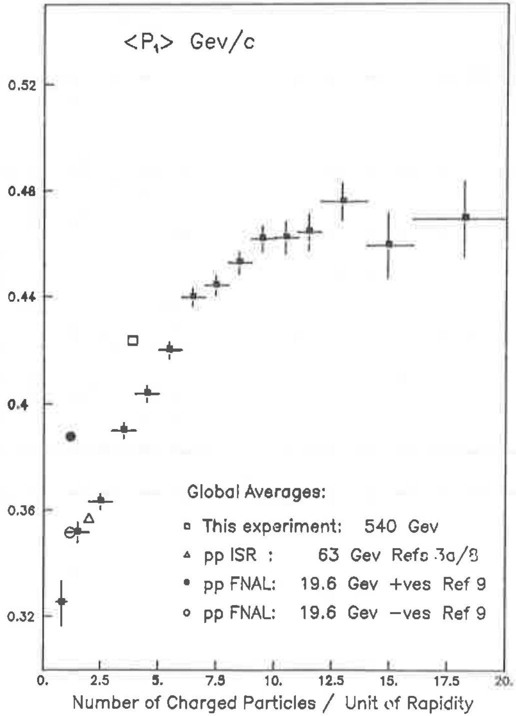


Figure 9.

with $p_{t0} \sim 1.3$ and $n \sim 9.1$

The spectrum is clearly dependent on charged track multiplicity in the same rapidity interval. This is illustrated in Figure 9 where the mean transverse momentum is plotted as a function of the charged particle density ; the $<p_t>$ increases from about 350 to 470 MeV/c, apparently saturating beyond $\sim$ 10 particles per unit of rapidity. This saturation may be kinematic, due to exhaustion of available energy, but may also be interpreted as a thermodynamic effect, involving a charge of state of hot hadronic matter, as suggested by Van Hove[14].

3.5. CORRELATIONS BETWEEN HIGH p_t CHARGED PARTICLES

Extensive studies at the CERN ISR[15] have demonstrated strong correlations in the production of high transverse momentum particles ; these correlations are interpreted as hard scattering of partons, where the scattered partons fragment into clusters of hadrons, or jets. Similar correlations have been observed in UA1 events[16], both in the minimum bias sample, and in a sample of 39000 events recorded with a central transverse energy trigger, $\Sigma E_t > 30$ GeV.

Selected events have been analysed which include a "software trigger" particle of transverse momentum $p_t > 4$ GeV/c. Using the trigger direction to define two regions, respectively "towards" ie. within azimuthal angle $|\Delta\phi| < 90^0$ of the trigger , and "away", Figure 10 shows the rapidity and azimuthal difference with respect to the trigger for other particles in each region. For "towards" particles, a clear trend is seen, increasing strongly with the p_t of the other particles, to cluster around the trigger in rapidity/azimuthal space. This is further demonstrated in Figure 11 where the p_t spectrum of particles near to the trigger is compared with the inclusive minimum bias distribution.

In the away region, the higher p_t particles tend to be produced at $\Delta\phi = \pm 180^0$, ie. coplanar, with respect to the trigger. There is no obvious correlation with the trigger rapidity. However

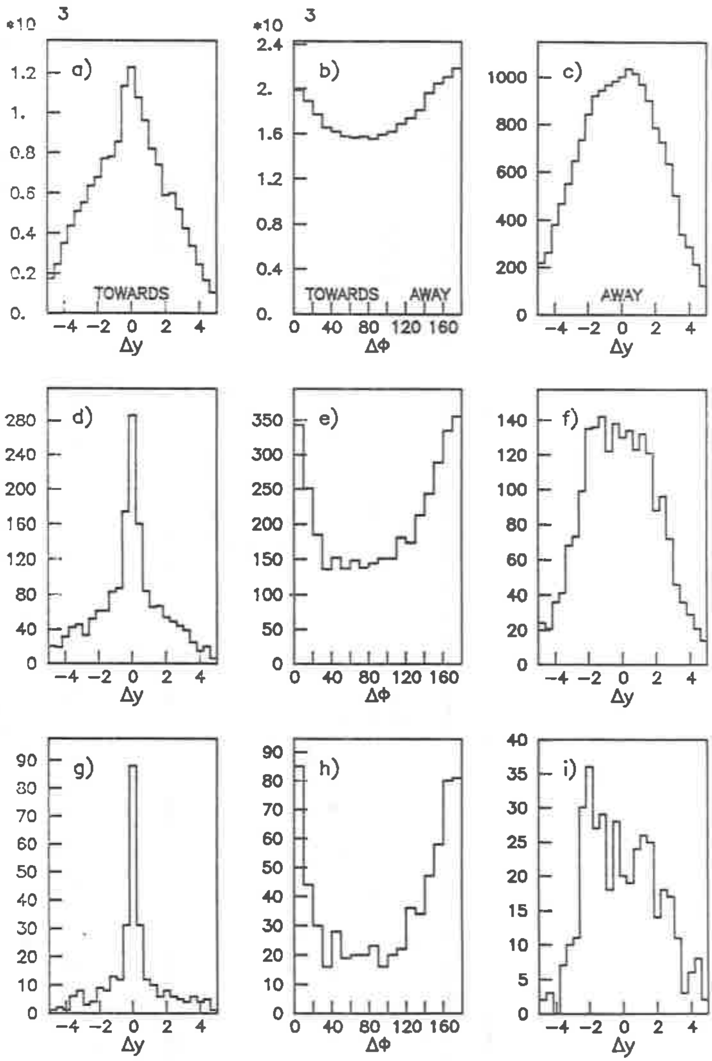


Figure 10. Rapidity-Azimuth Correlations of other Particles Relative to Trigger Track of p_t > 4 GeV.

p_t of other Particles > 0 GeV/c (a → c)

> 1 GeV/c (d → f)

> 2 GeV/c (g → i)

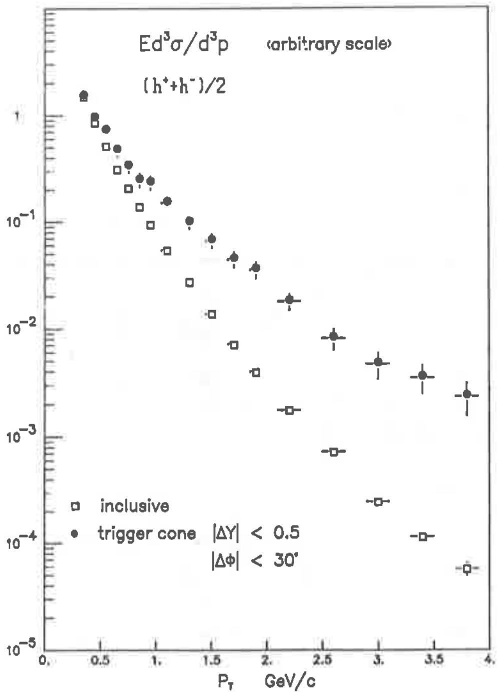


Figure 11.

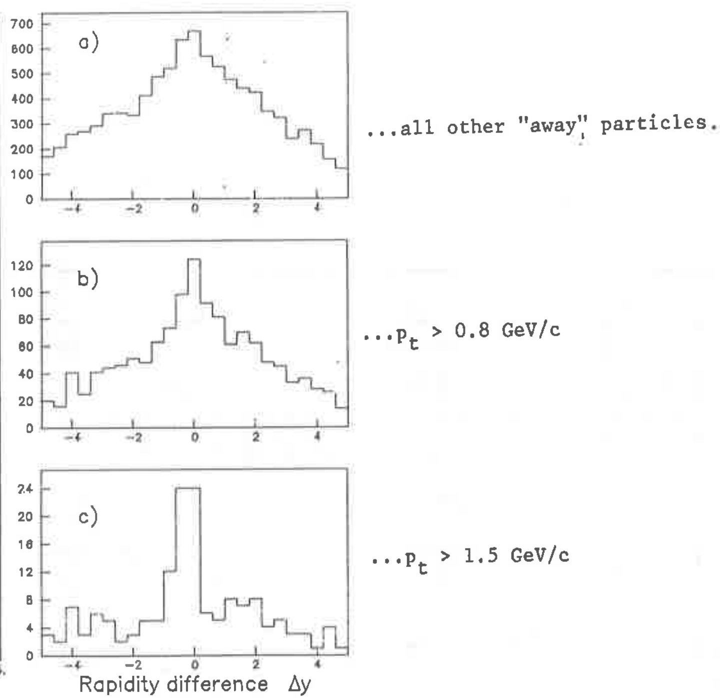


Figure 12. Rapidity Correlations w.r.t. highest p_t particle in the "away" region.

Figure 12 demonstrates a clustering of other high p_t particles around the highest p_t particle in the away region.

This observation of two coplanar cluster, uncorrelated in rapidity, and present in < 1% of minimum bias events, is the first indication that parton jets are present in collisions at the SPS collider.

4. "NEW" COLLIDER PHYSICS

At the time of this workshop, the most recent collider results concern the production of jets, and in particular the dominance of jet production at high transverse energies[17,18]. From the jet physics viewpoint, particle tracking and calorimetry are complementary ; however, for analysis of the initial low statistics data samples, separate analysis have been carried out, concentrating on different p_t regions and stressing different measurable characteristics.

4.1. LOW ENERGY JETS

Low energy jets, containing dominantly soft fragments which are spread by the magnetic field, are most accurately measured in the central detector (at least for charged tracks!). For the present analysis[19], a jet-finding algorithm has been developed to isolate clusters which have low relative internal transverse momenta. Initially considering each track as a cluster, the two clusters with the lowest relative p_t are merged and this procedure monotonously repeated until the relative p_t increases too steeply or crosses a threshold. Various fiducial cuts are then applied, and events are retained with at least one "central" cluster ($25^0 < \theta < 155^0$) with $p_t > 5$ GeV/c. The following table gives the percentage yields from different data samples recorded at 0.56 Tesla. The numbers are approximate due to uncorrected acceptance effects and inefficiencies , but show clearly the increased fraction of "jetty" events at higher global transverse energies, both for single and multi-cluster production.

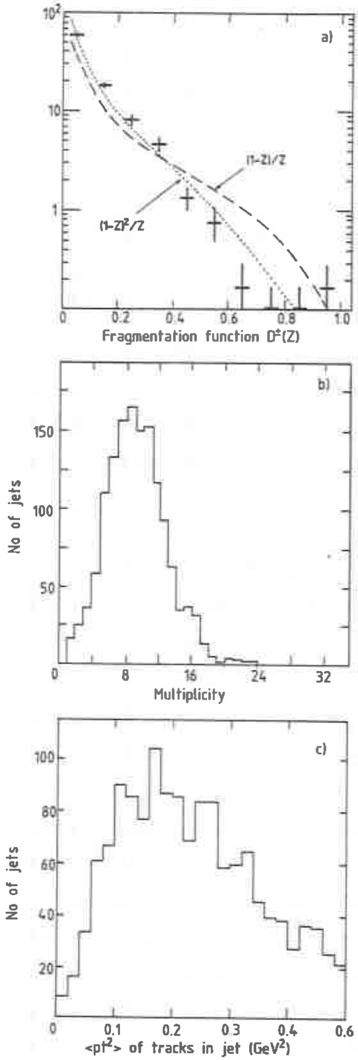


Figure 13. Properties of Charged Particle Jets ($25^0 < \theta < 155^0$).

TABLE I

Data sample	Number of Events	$\geq$ 1 cluster	3 clusters	> 4 cluster
Min. Bias	44887	2.1(± 0.1)	0.96(± .05)	0.83(±0.04)
High ΣE_t (> 30,40 GeV)	37062	12.4(± 0.2)	4.9(± 0.2)	6.1(± 0.2)
ΣE_t > 60 GeV in Δy = ± 1.5	436	51(± 4)	12(± 2)	34(± 3)

Events mixing and Monte Carlo techniques have been used to demonstrate that the measured jet signals are not due to chance fluctuations. In addition there is only a weak dependence on the parameters of the algorithm.

A first indication of internal jet properties is given in Figures 13 (a), (b) and (c), although it is stressed that these results are preliminary and much more sensitive to acceptance corrections and to algorithm parameters. The normalised fragmentation function $D^{\pm}(z) = 1/N\ dN/dz$ is compared with the functions $(1-z)/z$ and $(1-z)^2/z$, and shows a preference for the softer form. The mean multiplicity of charged particles within a cluster is ~ 8, and the mean internal p_t^2 is $\sim 0.2\ GeV^2$.

4.2. HIGH ENERGY JETS

High energy jets have been studied by detection of large local deposition of transverse energy in the central calorimeters[17]. The present analysis is restricted to the barrel calorimeters, covering pseudo rapidity range $|\eta| < 1.5$, but will be extended later to include the end-cap calorimeters $[1.5 < |\eta| < 2.6]$. The data samples were obtained by triggering on local transverse energy in the central calorimeters in excess of various thresholds ($\Sigma E_t > 20$, 30, 40 GeV), representing a total integrated luminosity of 22 μb^{-1}. The combined samples contain 6051 events with $\Sigma E_t > 40$ GeV, 279 with $\Sigma E_t > 20$ GeV and 5 with $\Sigma E_t > 100$ GeV.

Any jet study depends on initial definition of a jet, leading to a choice of algorithm and associated parameters which may influence quantitative results, especially concerning jet properties. A simple algorithm is however sufficient to demonstrate the presence of high energy jets at the collider. In the UA1 "window" algorithm, a jet is declared present in one half-shell of the barrel calorimeters if at least two-thirds of the transverse energy in that half shell is contained within 8 (out of 24) adjacent electromagnetic cells plus the matching hadronic ones. This allows events to be clearly categorised as 0-jet, 1-jet or 2-jet.

Figure 14 shows the fraction of such 2-jet events as a function of total transverse energy. The decreasing fraction up to 40 GeV is expected from multiplicity fluctuations, coupled with the dependence of ΣE_t on multiplicity, and has been reproduced by Monte Carlo simulation (solid curve). Above 40 GeV, the fraction increases due to onset of jet production which becomes dominant at the highest ΣE_t values. This conclusion has been confirmed by scanning of individual events on a graphical display system. For the five events with $\Sigma E_t > 100$ GeV, four of them are clear 2-jet events while the fifth contains a multi-jet structure of three (or four) jets. The 2-jet events are dominantly coplanar, as expected in a hard-scattering process, and are roughly balanced in p_t. Examples of different jet topologies are given in Figures 15 (a), (b) and (c) which show the struck calorimeter cells and the central detector track vectors above modest thresholds of order 1 to 2 GeV.

The inclusive cross-section for jet production around $\eta=0$ has been studied with both the window algorithm and with a more sophisticated "cluster" algorithm, to reveal possible sensitivity to jet definition. The cluster algorithm associates a vector to each struck calorimetric cell, and combines these vectors into clusters on the basis of their separation in $\eta\phi$ space. The inclusive cross-section

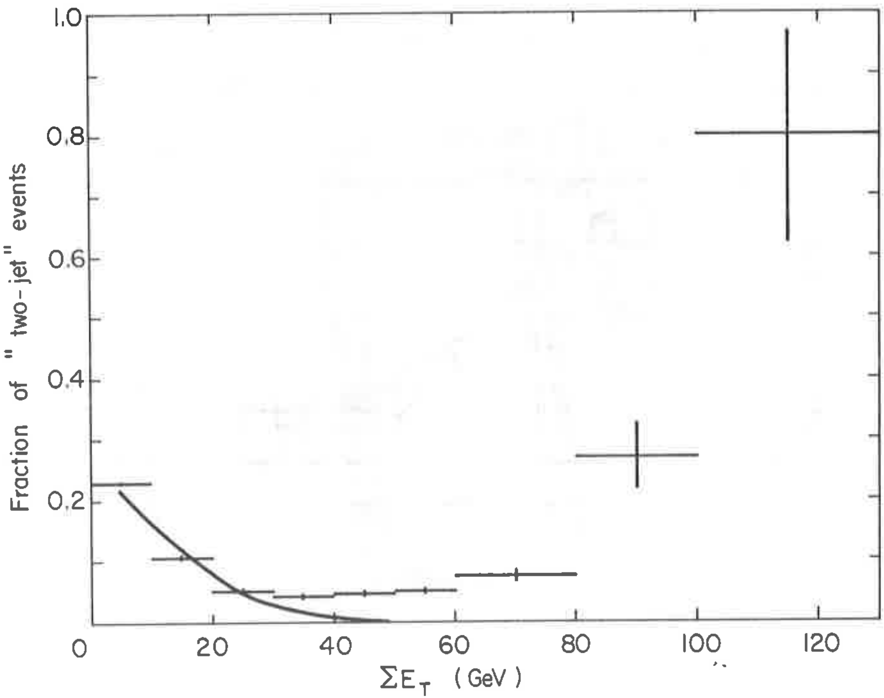


Figure 14. Fraction of 2-Jet Events vs Total Transverse Energy.

Figure 15 a) 2-JET Back to back

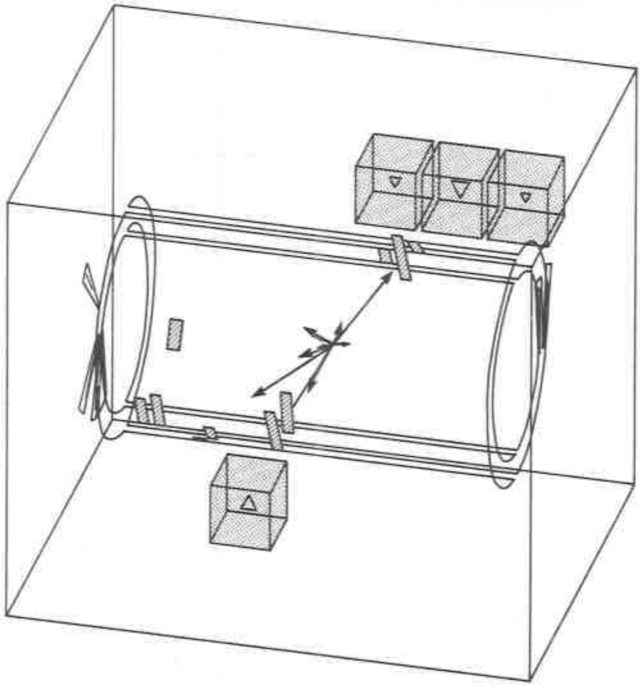


Figure 15 b) 2-JET Back-to-antiback

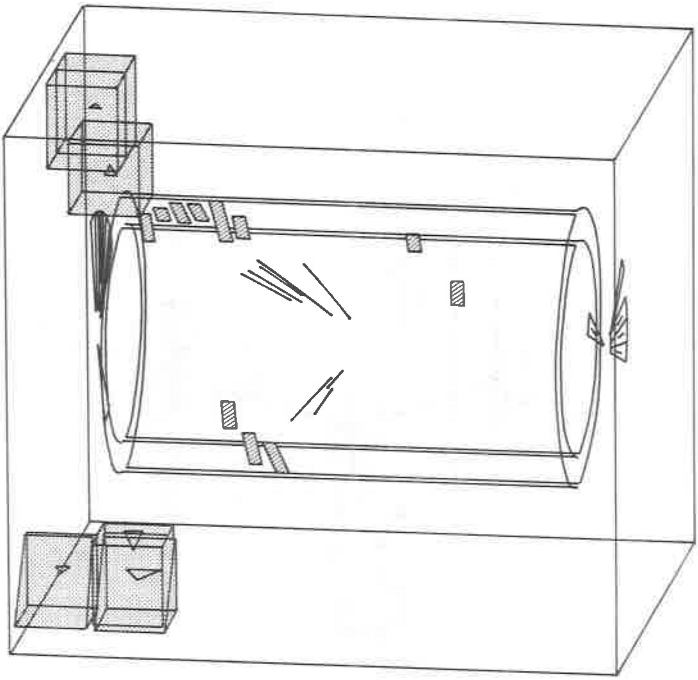
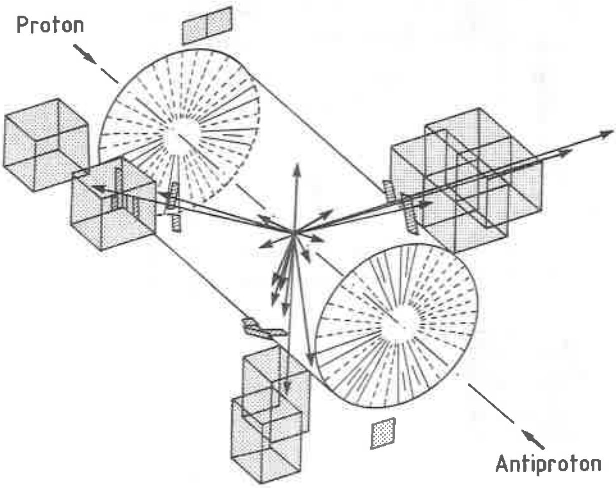


Figure 15 c) 3-JET



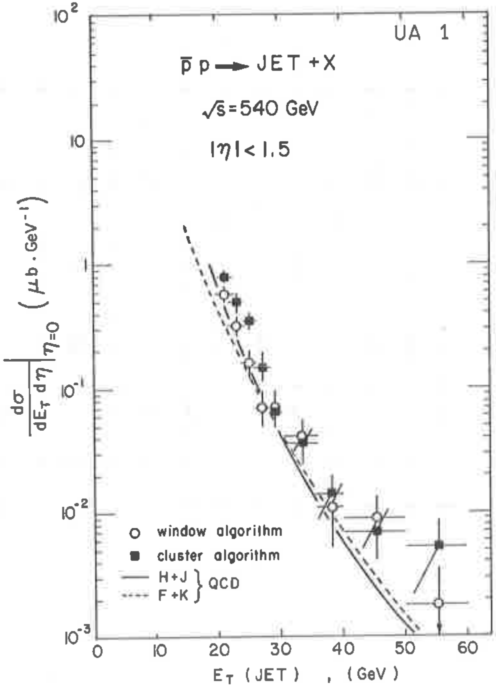


Figure 16.

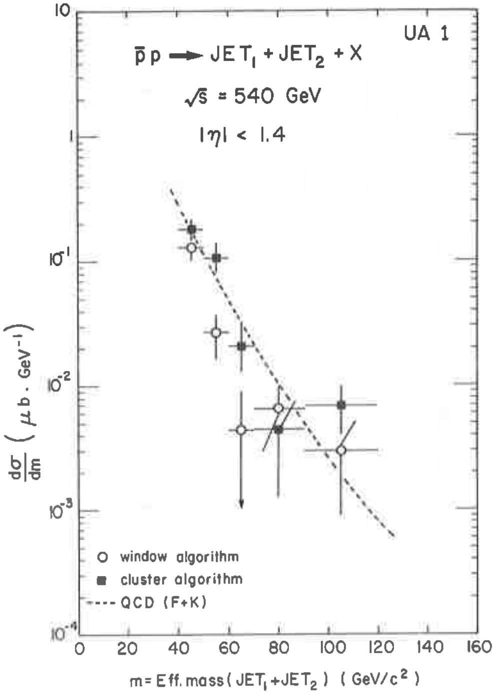


Figure 17.

for both algorithms is shown in Figure 6. The results agree reasonably well, with somewhat higher estimates from the cluster algorithm which tends to include more "soft" cells from the fringe area around a jet.

The results are also in good agreement with the predictions of two QCD-motivated models[20]. This is again true for the 2-jet mass spectrum, shown in Figure 17, which already extends to the region of W, Z decays. There is therefore a good prospect to detect these particles via their hadronic decay into two (or more) jets from later high statistics runs at the collider.

Addendum : At the time of writing this article, the jet cross-section has been measured with 1982 data for a total integrated luminosity of about 18 nb^{-1} and has been presented at the January workshop in Rome on proton-antiproton collisions[21]. The new spectrum is fully consistent with our previous one but remains high with respect to the published UA2 result[18].

5. "FUTURE" COLLIDER PHYSICS

(...with the benefit of hindsight!)

Future collider physics became present-day reality during November -December 1982, when UA1 collected data with a fully operational detector for a total integrated luminosity of 18 nb^{-1}, with peak rates of 5×10^{28} $cm^{2} sec^{-1}$. High energy electron candidates were isolated and studied during the run, leading to a presentation by Prof. C. Rubbia in January at the Rome workshop on proton-antiproton collisions and to the first publication[23] of events which are entirely consistent with the $W^{\pm}$-decay hypothesis. Independent evidence from the UA2 collaboration[24] was also presented at Rome.

The UA1 results are based on 5 events, plus one with an alternative interpretation, in which a high energy track in the central detector matches in position and energy with local energy deposition in the calorimeters with an electromagnetic profile. A parallel search, based on large missing transverse energy, yielded the same 5 events. To avoid background contamination from 2-jet events with precocious jet fragmentation, all events where the electron forms part of a jet, or is

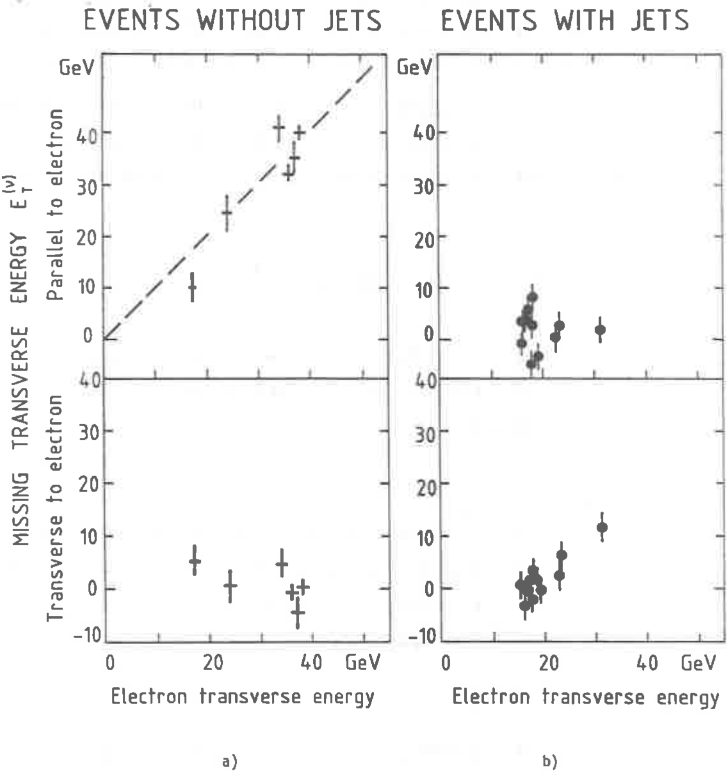


Figure 18. Correlation Between Missing Transverse Energy and Electron Transverse Energy for Selected Events (1982 data).

accompanied by a coplanar jet, have been scrupolously eliminated. Under these conditions no significant background source has been found.

The analysis also reveals a sample of 11 events with a high energy electron and an opposite jet. In Figure 18, we show the correlation of the missing transverse energy and the electron transverse energy for the electron-neutrino candidates and for the electron-plus-jet events.

The simultaneous presence of an electron and a neutrino which approximately balance in p_t suggests presence of the two-body decay $W^{\pm} \to e^{\pm} + \nu$. Assuming W-decay kinematics and V-A coupling, and correcting for the transverse motion of the W, the W-mass has then been calculated as :

$$m_W = 81^{+5}_{-5} \text{ GeV}/c^2$$

in excellent agreement with the Weinberg-Salam model[25]. The number of events observed, after correction for inefficiencies, is also consistent with predicted cross-sections.

6. Acknowledgements

I am extremely grateful to the organisers of the workshop, and to Prof. C. Rubbia, for providing the possibility to present UA1 collider results to an enthusiastic community. My very special thanks go M. Keller for her help in preparing this manuscript.

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Seminars

POLARIZATION EFFECTS IN e^+e^--ANNIHILATION AND RADIATIVE CORRECTIONS**

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** Talk presented by W. Hollik.

Abstract: The cross section for fermion pair production $e^+e^- \to \gamma, Z^0, \ldots \to f\bar{f}$, including transverse (natural) and longitudinal beam polarizations, and the final fermion polarization is calculated in the standard electroweak model and in extended models. The results contain the 1-loop electromagnetic corrections to γ and $Z^0_{(k)}$ exchange together with the soft photon bremsstrahlung thereby taking account of the resoncance character of the $Z^0_{(k)}$ boson(s)[1]. These QED corrections are the main parts of the complete electroweak 1-loop corrections. The virtual photon contributions consist of γ-self-energy, vertex corrections to γ and $Z^0_{(k)}$ vertices, $\gamma-\gamma$ and $\gamma - Z^0_{(k)}$ box diagrams.

The cross section for polarized $e^+e^- \to f\bar{f}$ has the general structure

$$d\sigma \sim (1 - P_L^+P_L^-)\,\sigma_U + (P_L^+ - P_L^-)\,\sigma_L + P_L^+P_L^-\tilde{\sigma}_L$$

$$+ P_T^+P_T^-\,(\sigma_T \cos 2\phi + \tilde{\sigma}_T \sin 2\phi) \quad ,$$

with

$$\sigma_{U,L,T} = \sum_\alpha \sigma^\alpha_{U,L,T} \quad , \quad \alpha = \gamma,\ \gamma Z_{(k)},\ Z_{(k)},\ (Z_k Z_l) \quad ,$$

longitudinal $(P_L^\pm)$ and transverse $(P_T^\pm)$ polatization degrees ($\tilde{\sigma}_L$ is only present in $e^+e^- \to e^+e^-$). The following results were presented:

1. Transverse polarization gives rise to an azimuthal asymmetry $A_T = \sigma_T/\sigma_U$. $1-\gamma$ exchange predicts $A_T = 1$ $(\theta = 90^0)$ for $f = \mu, \tau, q$. Radiative corrections to this QED part are negligibly small. Deviations from 1 are purely weak effects, quadratic in the Z^0 propagator. For $\sqrt{s} < M_{Z^0}$, where A_T changes its sign, radiative corrections to A_T become important. Above the resonance the radiative tail effect occurs, strongly dependent on $\Delta E/\sqrt{s}$ with ΔE = maximum bremsstrahlung energy. On the Z^0 resonance, A_T is practically insensitive to radiative corrections.

2. Longitudinal polarization: At PETRA energies an antiparallel spin configuration of e^+ and e^- increases the weak effect in Bhabha scattering[2] from ~3% (unpolarized beams) to ~10% . The polarization asymmetry $A_L = \sigma_L/\sigma_U$, however, is small except around the Z^0 .

In quark pair production[1], A_L is sizeable (20 - 50%) already

in the γZ interference region. Radiative corrections are also of importance and practically independent of $\Delta E/\sqrt{s}$ for $\sqrt{s} < M_{Z0}$. On resonance, where A_L is large, the corrections to A_L are very small.

3. Final state polarization P_f: For $f = \mu,\tau$ P_f is identical to A_L also in higher order, but $P_f \neq A_L$ for $q\bar{q}$ production. The radiative corrections to P_f are important in the γZ region; on resonance they are very small. Since P_f is 70 - 90%, the Z^0 would be a good source for polarized quarks, which can be used for investigations of the helicity transfer in fragmentation processes.

4. Alternative models: In models with more Z^0 bosons the polarization asymmetry A_L is different from the standard model prediction, also in the case $M_{Z0} \approx M_{Z^0_{(1)}}$ [3]. The inclusion of radiative corrections on the (first) Z^0 resonance does not shift the values of A_L, as obtained in lowest order, neither in the standard nor in more boson models.

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RADIATIVE CORRECTIONS TO HIGGS PRODUCTION BY e^+e^- IN THE GLASHOW-WEINBERG-SALAM MODEL

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* Report on work in collaboration with J. Fleischer

Abstract:

At least one neutral scalar Higg-boson (H) must "show up", if our present understanding of electro-weak interactions is correct. In the case the Higgs particle is not too heavy, $m_H \lesssim 100$ GeV, there is a good chance for the Higgs to be produced in future e^+e^- - and $pp(p\bar{p})$ - experiments. H is most abundantly produced in association with heavy particles and it decays into the heaviest particles which are energetically accessible.

The total Higgs-boson width varies widely with m_H [1]: $\Gamma_{H,tot}(m_H) = 5\times10^{-4}(10)$, 0.1(100), 2(200) GeV. Below the W^+W^--threshold ($2M_W$) the decay into heavy fermions predominates and the Higgs-boson is expected to be pretty stable.

Presumably, the best "clean" (non-hadronic) production process is [1] $e^+e^-\to Z^*\to ZH$ which above the threshold $\sqrt{s}\gtrsim M_Z+m_H$ highly dominates the process $e^+e^-\to f\bar{f}H$ observed actually [2].

In the Born approximation the ratio $R = \sigma_{ZH}/\sigma_{\mu\mu}$ as a function of the centre of mass energy $E = \sqrt{s}$ has a peak. The values $R(m_H, E_{peak}) = 1.23(10, 109)$, $0.74(24, 127)$ and 0.18 $(100, 243)$ show that an "observation" of H should be possible.

Since the energy involved in this process is rather high ($\sqrt{s}\gtrsim100$ GeV) the knowledge of radiative corrections might be important. We have analyzed in detail the electro-weak one-loop corrections including standard soft photon bremsstrahlung [3].

The corrected cross-section

$$\sigma = \sigma_o + \sigma_c = \sigma_o(1+C_F+C_{VB})\exp C_{IR}(x_r)$$

includes contributions from virtual fermions C_F, from virtual

vector-bosons C_{VB} including infrared regular terms from QED and soft photon bremsstrahlung.

$$C_{IR}(x_r) = -\frac{e^2}{2\pi^2}\{\ln x_r(1-\ln y) + \frac{1}{2}\ln y\};\ y = {}^{s}/m_e^2$$

is the standard QED infrared sensitive term, x_r the fraction of e^+e^--energy carried by soft photons. The total correction relative to σ_o is

$$\Delta C(x_r) = (1+C)e^{C_{IR}}-1 \quad \text{with} \quad C = C_F + C_{VB}$$

the "weak" correction. Whereas C is specific for the process $e^+e^- \to ZH$, the bremsstrahlung correction $C_{IR}(x_r)$ is universal for all processes $e^+e^- \to$ neutrals, depends sensitively on the "counter resolution" x_r, however.

Some of the results obtained (using standard values for the parameters) are given in Tab. 1.

$\sqrt{s}$ GeV	100	200	500	1000
ΔC	-6	-8.5	-11	-14
C	27.5	26.5	26	23.5
C_F	17	19	21	23.5
C_{VB}	10	7.5	5	~0

Tab. 1: Percentage corrections to σ_{oZH} for x_r=0.1.

The following general features have been found:

- The "weak" correction C, for $m_H < 2M_W$, depends weakly on m_H and slightly decreases with energy.
- For $m_H < 2M_W$, the fermion contributions dominate particularly at higher energies.
- C_F is independent on m_H and $C_{VB} \propto m_H^2$ for $m_H > 2M_W$.
- As expected the "weak" corrections increase the cross-section due to an increase of the effective couplings with energy. This constructive effect is essentially cancelled by the soft photon effects (for x_r = 0.1).

We notice that the "weak" corrections are not small and the cancellation by soft photon effects, which have been taken into account to all orders by exponentiation, need not be at work for higher order corrections.

There is one major problem: the problem of heavy fermions [4]. We assumed the existence of a t-quark with m_t = 35 GeV. The results are independent on m_t if $m_t < M_Z/2$. However, there might exist additional super-heavy fermions. Decoupling of heavy fermions fails in the GWS-model. Whereas the existence of fermion doublets with large mass splitting is restricted by low energy data (ρ-parameter) [5] heavy mass-degenerate doublets are not excluded.

For $m_f < M_Z/2$ the cross-section σ_{ZH} is essentially independent on m_f. However, for $m_f > M_Z/2$ we find a strong suppression of σ_{ZH} in both cases, for heavy single fermions and for heavy doublets. When $m_f > 300$ GeV the perturbation expansion breaks down. As an effect, in the domain $2m_f > \sqrt{s} > M_Z + m_H$, the production of the Higgs particle by $e^+e^- \to ZH$ could be heavily suppressed by (non-perturbative?) virtual heavy fermion effects.

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MASSIVE NEUTRINOS: NOVEL THEORETICAL AND EXPERIMENTAL IDEAS

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Abstract

Massive neutrinos, with masses of a few MeV to a GeV or so, will be discussed in this talk. We first describe the theoretical motivation for having such heavy neutrinos and then - a new experimental search for such neutrinos.

The present experimental limits on the three neutrino masses give rise to the following perplexing questions:

1. $m_{\nu_e} < 60$ eV : Why is the electron neutrino so much lighter than the electron?

2. $m_{\nu_\mu} < 500$ KeV : Since cosmology tells us that stable neutrinos must be lighter than 50 eV , do we have to conclude that the mass region between these two values, separated by four orders of magnitude, is absolutely excluded? Or could the muon neutrino be made unstable to avoid the cosmological constraint?

3. $m_{\nu_\tau} < 250$ MeV : This limit being so weak, how could one improve it to find out whether ν_τ is not heavier than a few MeV for instance?

By first attempting to answer the first two questions we are led to the introduction of right-handed neutrinos with masses of the order of a hundred MeV (part I of our talk); we then answer the third question by describing an experimental search for heavy neutrinos (part II).

I. Theoretical Motivation for $m_\nu \sim 0(100$ MeV$)$ [1)]

1. We adopt the beautiful idea of Gell-Mann, Ramond, Slansky and Yanagida which answers the first question by attributing the lightness of the left-handed neutrino to its Majorana character and to the existence of a very heavy right-handed partner. The two mass eigenvalues which emerge from this scheme are m_M and m_D^2/m_M where the Majorana mass m_M obtains its value from the GUT scale whereas the Dirac mass m_D is related to the corresponding charged fermion mass. The two mass eigenstates are almost pure right and left chirality states, each containing a tiny component m_D/m_M of opposite chirality. We will actually work within the left-right symmetric model[2)] in which parity restora-

tion occurs at an intermediate energy m_R (m_M may be smaller than m_R just as $m_e \ll M_W$).

2. To answer the second question we note that when both Dirac and Majorana masses are introduced for the neutrinos, the GIM cancellation breaks down and flavour changing neutral currents among neutrinos are allowed. This gives rise to the process $\nu_{\mu L} \to \nu_{eL} \nu_{eL} \bar{\nu}_{eL}$ with a rate proportional to $(m_D/m_M)^4$ which may avoid the cosmological constraint. The two inequalities for m_D/m_M, which follow from the requirement $50\ \text{eV} < m_{\nu_\mu} < 500\ \text{KeV}$, have a consistent solution which leads to the following values of the masses of the light and heavy muon neutrinos: $m_{\nu_{\mu L}} \sim 200 - 500\ \text{KeV}$, $m_{\nu_{\mu H}} \sim O(100\ \text{MeV})$. The heavy (right-handed) neutrino is expected to give rise to a secondary peak in $K \to \mu\nu$ decay[3] with a rate of the order of 10^{-3}. For other consequences see part II of this talk.

II. Experimental Search for Heavy Unstable Neutrinos[4]

3. If the τ neutrino is heavier than one MeV neutrino beams will contain a fraction of these neutrinos which while traveling in the beam may decay into final states with electrons and muons. The feasibility of the decay signals is demonstrated by showing their high sensitivity to neutrino mixing parameters as small as $10^{-6}(m_\mu/m_{\nu_\tau})^6$. The main conclusions of our study are the following:

a) The bound $|U_{e3}|^2 < 10^{-6}(m_\mu/m_3)^6$ may be derived from limits on $\nu_3 \to e^- e^+ \nu_e$ as possible background to $\nu_\mu e \to \nu_\mu e$.

b) The very conservative bound $|U_{e3}|^2 < 10^{-4}(m_\mu/m_3)^6$ is obtained from the number of prompt ν_e charged-current events observed in beam dump experiments. This bound may be improved by two orders of magnitude by concentrating on low E_{vis} events.

c) Dedicated neutrino decay experiments are suggested to considerably improve these limits. A currently running decay beam dump experiment is shown to have the potential of improving the bound down to $|U_{e3}|^2 < 10^{-8}(m_\mu/m_3)^6$.

d) Similar bounds may be obtained for $|U_{\mu 3}|^2$ if $m_\mu < m_3 < 250\ \text{MeV}$.

e) The natural "generation gap" assumption $|U_{e3}| < |U_{\mu 3}|$ attributes the excess of prompt muon events over prompt electron events observed in beam dump experiments to neutrino decay.

f) Similar, however somewhat different, limits may be derived for right-handed neutrinos and higher generation neutrinos. Neutrinos heavier than the kaon may be produced and decay in beam dump experiments.

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II. GRAND UNIFICATION AND SUPERSYMMETRY

GRAND UNIFICATION-PHENOMENOLOGICAL CONSEQUENCES

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Sometime ago Reinhart Kögerler called me and asked me to give a review talk on the phenomenological aspects of grand unification at this 1982 DESY Workshop. I agreed without hesitation. My immediate reaction was "what a wonderful task"; there was so much new material I could discuss. But when Reinhart enumerated the topics I was not supposed to touch (cosmology, neutrino masses and oscillations, supersymmetry, etc.), as they belonged to other speakers, I felt quite uneasy. I was asked to review mainly the present situation of nucleon instability within the framework of old fashioned grand unified theories[1,2] (GUT), a saga which most of you have heard many times before. So, I will not repeat it once more but shall mainly concentrate on recent results obtained within the last year. I will also discuss some other baryon and lepton number violating processes. The topics to be discussed are:

1. Why grand unification
2. Recent results on nucleon decay
   - 2a. Family mixing
   - 2b. Recent calculations of lifetime
   - 2c. Nucleon decay in supersymmetric GUT
3. Double beta processes
4. Other baryon/lepton number violating processes

1. Why grand unification

The standard electroweak-QCD theory (in spite of having provided a decisive step forward in our understanding of nature) is unlikely to be telling the whole story of nongravitational forces. Many fundamental questions fall outside its scope. Some such questions are:

Why are there quarks and leptons? Our gauge theories would have been simpler and more beautiful without them. Why three families? Why is charge quantized? Why is the symmetry pattern $U(1)_y \times SU(2)_L \times SU(3)_c$ which gives for example V-A charged currents, etc.? There are also a number of empirical facts which are not explained such as the pattern of masses and mixing angles, the strength of the coupling constants (g_1, g_2, g_3) and the value of $\sin^2\theta_W$.

GUT provide a framework for answering some of the above questions.[2] The quantization of charge, in these theories, is as natural as the quantization of angular momentum in quantum mechanics.

Also by unifying the three interactions the coupling constants (g_1, g_2, g_3) are related. The "observed" fractional quark charges, the value of $\sin^2\theta_W$ and the ratio of masses of b-quark and τ-lepton are explained in the minimal grand unified model, the SU(5) of Georgi and Glashow.

A very attractive feature of GUT is that they relate the quarks and leptons and thus lead to baryon and lepton number nonconservation which constitutes the most striking phenomenological consequence of GUT. This is gratifying because we know of no fundamental reason why baryon and lepton numbers should be good quantum numbers.

A famous German physicist who was once giving a talk on the physics of the past fifty years said: "Much has been learned during the past fifty years; many problems have been solved. But the present time seems to offer still more puzzels, and perhaps harder ones."

The physicist was Max Born and the year 1953. Now, thirty years later the statement seems to be just as valid. Grand unification also solves some problems but "seems to offer still more puzzels", such as the hierarchy problem and the monopole problem.

2. Recent results on nucleon decay

The present experimental situation[3)] in a nutshell is that the Kolar Gold Mines Collaboration reports[4)] three fully contained proton decay candidates and a lifetime $\tau_p = 8 \times 10^{30}$ years. The Mont Blanc experiment[5)] has a number of fully contained neutrino events and one possible proton decay candidate which could be $p \to K^0\mu^+$, $p \to K^*\nu$, $p \to 3\mu$, $p \to \rho^0\mu^+$ or with a small probability a neutrino event $\nu n \to \mu\pi n$, with an interacting pion. Assuming that the event is a proton decaying one finds $\tau_p = (0{,}8-1{,}6)r \times 10^{31}$yr, where r is the branching ratio. Other experiments in progress have not yet reached the necessary sensitivity level.

On the theoretical side the major question is what is the predicted value of the nucleon instability lifetime? I shall now discuss some new theoretical developments.

2a. Family mixing

Altschüler et al[6)] and Nandi et al[7)] have examined, in some detail the question of Cabibbo-like mixing angles in the nucleon decay hamiltonian of SU(5) as I shall now sketch.

Just as the interactions of quarks with $W^{\pm}$ are not diagonal in family, it is possible that the interactions of quarks and leptons with the gauge bosons $X^{(\pm 4/3)}$ and $Y^{(\pm 1/3)}$ of SU(5) are not diagonal in family. If one is unlucky the proton decay hamiltonian may prefer the transition $p\to\tau^+\pi^0$ which is however forbidden by the energy-momentum conservation law.

In SU(5) the interactions of the X and Y with the matter fields are shown in Fig. 1.

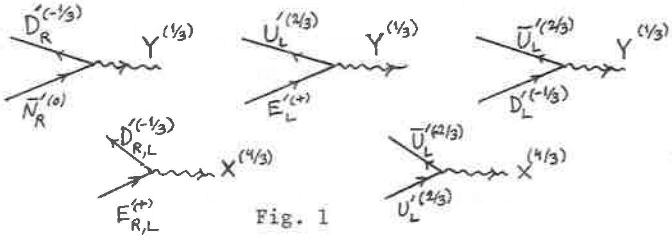


Fig. 1

In Fig. 1, D', U', etc. are related to the physical fields by unitary transformations,

$$D_H'^{(-1/3)} = \begin{pmatrix} d' \\ s' \\ b' \\ \vdots \end{pmatrix}_H = H(-1/3)\, D^{(-1/3)} = H(-1/3) \begin{pmatrix} d \\ s \\ b \\ \vdots \end{pmatrix}_H , \qquad H = L, R .$$

H(Q) is a unitary matrix, and Q denotes the electric charge. Note that the vertices in Fig. 1 are not the most general ones allowed by charge conservation. For example $E_R'^{(+)} \to U_R'^{(2/3)} + Y^{(1/3)}$ is missing because the E_R^+ and U_R' belong to different multiplets of SU(5). The vertices in Fig. 1 give

$$\begin{aligned}\mathcal{L} \sim & \left[\bar D_R' N_R'^{c} + \bar U_L' E_L'^{(+)} + \bar U_L'^{c} D_L'\right] Y + \left[\bar D_L' E_L'^{(+)} + \bar D_R' E_R'^{(+)} + \bar U_L'^{c} U_L'\right] X = \Big\{ \bar D_R R^+(-1/3) L^*(0) N_R^c + \\ & + \bar U_L L^+(2/3) L(+1) E_L^+ + \bar U_L^c R^T(2/3) L(-1/3) D_L \Big\} Y + \Big\{ \bar D_L L^+(-1/3) L(+1) E_L^{(+)} + \bar D_R R^+(-1/3) R(+1) E_R^{(+)} + \\ & + \bar U_L^c R^T(2/3) L(2/3) U_L \Big\} X ,\end{aligned} \tag{1}$$

where we have dropped the γ-couplings and the hermitian conjugate terms. R(L) projections are $(1\pm\gamma_5)/2$ and C denotes charge conjugation. Thus, for example

$$\begin{aligned} N_R'^{c} &= \frac{1+\gamma_5}{2} C \bar N'^{T} = \frac{1+\gamma_5}{2} C\gamma_0 N'^{*} = C\gamma_0 \left(\frac{1-\gamma_5}{2} N'\right)^{*} = C\gamma_0 N_L'^{*} = \\ &= C\gamma_0 L^*(0) N = L^*(0) N_R^c , \\ \bar U_L'^{c} &= \bar U'^{c} \frac{1+\gamma_5}{2} = -U'^{T} C^{-1} \frac{1+\gamma_5}{2} = \bar U_L^c L^T(2/3). \end{aligned}$$

We note also that

$$L(Q) = R^*(-Q)\,, \qquad L^+(2/3)\,L(-1/3) = V_c$$

where V_c is the Cabibbo (Kobayashi-Maskawa) matrix for two (three) families. The coupling of the charged leptons to W-boson defines the neutrino

$$\bar{E}'_L N'_L = \bar{E}\, L^+(-1)\, L(0)\, N_L \,.$$

Since the neutrino is massless in SU(5) we may put $L^+(-1)L(0)=1$. Then with the definitions

$$U_1 = R^+(-1/3)\,L^*(0)\,, \qquad U_2 = L^+(-1/3)\,L(+1)\,, \quad U_3 = R^T(2/3)\,L(-1/3)\,,$$

we have

$$L^+(2/3)\,L(+1) = V_c U_2\,, \quad R^T(2/3)\,L(2/3) = U_3 V_c^+ .$$

Putting these into eq. (1) and using $L(0)=L(-1)$ gives

$$\mathcal{L} \sim \{ \bar{D}_R U_1 N_R^c + \bar{U}_L V_c U_2 E_L^{(+)} + \bar{U}_L^c U_3 D_L \} Y + $$
$$+ \{ \bar{D}_R U_1 E_R^{(+)} + \bar{D}_L U_2 E_L^{(+)} + \bar{U}_L^c U_3 V_c^+ U_L \} X + h.c. \qquad (2)$$

The nucleon will be stable if the transitions involving the light quarks (u,d,s) should vanish, i.e.,

$$(U_3)_{11} = (U_3)_{12} = 0 \quad , \qquad (U_3 V_c^+)_{11} = 0 . \qquad (3)$$

For three families U and V are three by three unitary matrices, then the conditions (3) imply

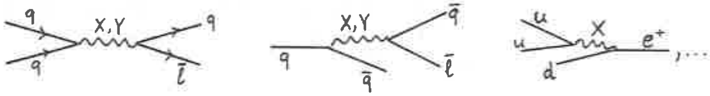
$$|(U_3)_{13}| = 1$$
$$(U_3 V_c^+)_{11} = (U_3)_{13} (V_c^+)_{31} = 0 \quad \rightarrow \quad (V_c^+)_{31} = 0$$

i.e., $V_{ub}=0$. The only point of the above exercise is that one cann't reliably predict the nucleon lifetime, unless one knows more about the pattern of the spontaneous symmetry breaking and family mixings.

2b. Recent calculations of lifetime

The possible mechanisms for nucleon decay in SU(5) are two quark fusion,[8] quark decay,[9] three quark fusion[9] mediated by X and Y exchange; some examples are shown in Fig. 2.

Higgs exchange contribution is expected to be small.[8]



The quark decay is much suppressed[9] due to the three body phase space. Previously it was estimated[9] that also the three quark fusion mechanism is unimportant because the nucleon has to be far off mass shell $q_N^2 \approx 0$. Recently some authors[10,11] have given estimates which are larger than one would have expected a priori. Note that the three quark fusion is relevant, for example, in the pole diagram (Fig. 3) calculations.

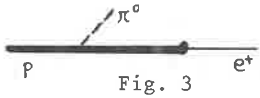


Fig. 3

Approaches used in calculating the nucleon lifetime are
- SU(6) quark model
- Bag models
- Current algebra
- Pole dominance model and dispersion relations.

There is a clear tendency seen in the more recent calculations: the estimated lifetime in the minimal SU(5) and with M_X, etc. as in Langacker's review article,[1] is now in a serious difficulty with experiments.

In the pole model, for example the main uncertainties are the 3 quark wavefunction at the origin and the value of the pion nucleon coupling constant $g(M_p^2, m_\pi^2, m_e^2)$. Berezinsky et al[11] write a dispersion relation for the coupling constant $g(M_p^2, m_\pi^2, t)$ in the variable t and saturate the imaginary part with the nucleon pole whereby the coupling constant falls off gently. They find $\tau(p \to e^+ \pi^0) \approx 5 \times 10^{28}$yr. The current algebra approach[12] gives a very similar result. Also a new bag calculation[13] by Thomas and McKellar using the so-called cloudy bag model finds τ_p is again smaller than the experimental limit. Actually, they find a large cancellation between the pole and non-pole contributions. Very recently Isgur and Wise[14], in a nonrelativistic quark model calculation, taking into account the pole and non-pole contributions again find that $\tau(p \to \pi^0 e^+)$ is about an order of magnitude shorter than the experimental limit.

Does all this mean that the minimal SU(5) is out? I believe that the answer is no. All the theoretical calculations done so far contain assumptions which, in spite of looking reasonable, may in fact be wrong.

As far as the neutron baryon number violating decay is concerned it was first considered by the author and Ynduráin (Ref. 9). In SU(5), the decay $n \to e^+ \pi^-$ has a nice signature. However the rate in the SU(6) quark model is expected to be somewhat slower than that of $p \to e^+ \pi^0$. Experimentalists don't seem to care much about it, perhaps because of π^- absorption in matter.

2c. Nucleon decay in supersymmetric GUT

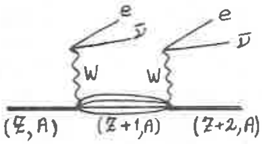
In supersymmetric (SUSY) grand unified theories (GUT) each spin $\frac{1}{2}$ fermion has two spin zero companions and each gauge boson is accompanied by a spin $\frac{1}{2}$ partner. In a nutshell what happens is that there are many more scalars and spin $\frac{1}{2}$ fermions around. The β-functions which govern the evolution of coupling constants become smaller. Thus the strong and the SU(2) coupling constants fall off more slowly than in GUT and meet later giving a much larger unification mass. The X and Y mediated decays become much less important. Generally in SUSY-GUT the nucleon instability is due to Higgs exchange. As the Higgs couplings are proportional to masses one usually has that transition to heavier fermions is favoured. Therefore the family mixing angles could be more important here than in GUT. The nucleon instability in SUSY-GUT has been studied by many authors[15]. The major decay modes are expected to be $p \to \bar{\nu} K^+$, $n \to \bar{\nu} K^0$, which if true is very unfortunate. However the SUSY-GUT models so far seem to have very serious difficulties and there is no reason to put much confidence in their predictions. For example Masiero et al[16] by introducing a 50-plet of Higgses can completely change the previous picture. The nucleon would then decay essentially as it does in GUT. But "who ordered the fifty?"

Now I will turn into some other lepton/baryon number instability phenomena. It is indeed essential to look for all possible manifestations of baryon and lepton number non-conservation. As one says in this country

"Nur die Fülle führt zur Klarheit
Und in Abgrund wohnt die Wahrheit".

3. Double beta processes

The double beta decay was invented[17] by Maria Goeppert Mayer in 1935, nearby half a century ago and only one year after the Fermi theory of beta decay was published. Her double beta process is shown in Fig. 4.



$$Z \to (Z+2) + 2e + 2\bar{\nu}.$$

Fig. 4

Of course she didn't have the W-boson, but that is irrelevant (the $W^{\pm}$ were introduced by Oskar Klein in 1938). A few years later it was realized by Furry[18] that if the neutrino is its own antiparticle (i.e. a Majorana neutrino) then the neutrinoless double beta process shown in Fig. 5 may take place.

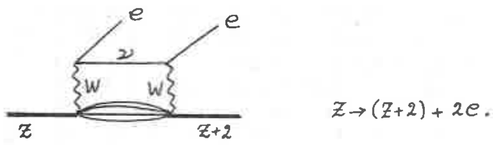


Fig. 5

Furry realized that this lepton number violating amplitude is much enhanced by phase space and that the neutrinoless process could have many orders of magnitude shorter lifetime than the lepton number conserving double neutrino process of Fig. 4. It is sad that in spite of nearly half a century of efforts there is as yet no convincing evidence for the double beta decay. Recently just because of GUT double beta decay has come into the focus. The experimental activity in the field has increased and theorists have looked for new mechanisms for it, for example exchange of exotic gauge bosons and Higgs particles, as shown in Fig. 6.

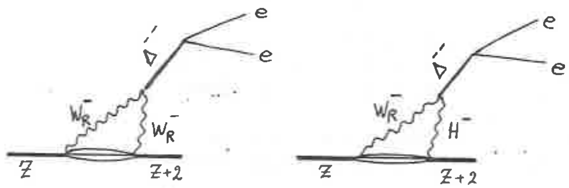


Fig. 6

Experiments on double beta decay are either a) geochemical where one looks for the daughter isotope (Z+2,A) in old rocks containing the mother isotope (Z,A). In order to see an effect the daughters should stay at "home" for a long time ($\sim 10^9$ years), which they don't seem to do. That is the major problem with these experiments. Also the (0ν)- and (2ν)-modes are not distinguished because the electrons are not detected and so one doesn't know if there is any missing energy due to neutrino emission. The second kind of experiments are b) direct or laboratory experiments, where the electrons are observed. The present experimental situation in a nutshell is as follows: A direct experiment[19)] reports the observation of $^{82}Se \to {}^{82}Kr + 2e^- +$ .. where the reaction is presumably the (2ν)-mode, as energy is missing. However the measured partial lifetime is an order of magnitude shorter than the total lifetime $^{82}Se \to {}^{82}Kr +$ .. measured in geochemical experiments. So something is wrong. In the Te-Xe system, as suggested by Pontecorvo[20)] one compares

$$^{130}Te \to {}^{130}Xe + 2e^- + \text{ .. and } {}^{128}Te \to {}^{128}Xe + 2e^- + \ldots$$

The rate for each is assumed to be the sum of the (2ν)- and (0ν)-modes. The ratio of the rates is very sensitive to the (0ν)-mode. For purely left-handed currents, for the (0ν)-mode (see Fig. 7)

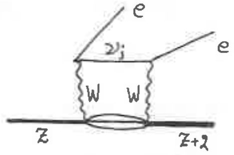


Fig. 7

the leptonic current $\bar{e}\gamma_\lambda(1-\gamma_5)\nu_e$ enters at two vertices. Now ν_e is assumed to be a linear combination of fundamental Majorana neutrino fields ν_j with well defined masses m_j, i.e.,

$$\nu_e = \sum_j U_j \nu_j \;, \qquad \sum_j |U_j|^2 = 1$$

Here U_j are in general complex constants. The diagram gives a factor

$$\Big(\sum_j m_j U_j^2\Big)\, \bar{e}(1-\gamma_5)e^c \equiv \langle m_\nu \rangle\, \bar{e}(1-\gamma_5)e^c,$$

where the quantity denoted by $\langle m_\nu \rangle$ is not the neutrino mass (it need not even be real) but is a kind of effective mass. The rate for (0ν)-mode is then proportional to $|\langle m_\nu \rangle|^2$. If one knew the nuclear matrix elements one could extract this neutrino mass from data. Until recently an experiment[21)] measuring the Te-Xe system and indicating the existence of the (0ν)-mode was used by several theorists[22)], who depending on their choice of nuclear matrix elements, were obtaining different values for the Majorana effective mass of the neutrino, ranging between 1 eV to 30 eV. There is now a contradicting experiment[23)] by T. Kirsten and his collaborators who find that there is no evidence for the (0ν)-mode. So one will get an upper limit on $\langle m_\nu \rangle$. Other double beta experiments give so far only lower limits on the lifetimes. For example, the Milano group[23)] is now at the level of 10^{22} years (Ge→Se) and plans to penetrate into 10^{23} years range in the near future.

By "crossing" the double beta reaction $Z \to (Z+2) + e^- + e^- + \ldots$ one obtains a number of associated reactions such as

$$\begin{aligned}
&e^- + e^- + (Z+2) \to Z \\
&e^- + e^- + (Z+2) \to Z + \nu_e + \nu_e \\
&e^- + (Z+2) \to Z + e^+ \\
\text{etc.} \qquad & \to Z + e^+ + \nu_e + \nu_e \,,
\end{aligned}$$

The first reaction, the double electron capture (with no neutrinos) is specially interesting as it may induce mixing between different atoms. The neutrinoless double beta capture was studied[24)] a long time ago.

A recent paper by Georgi et al[25] triggered our[26] interest to start a serious study of this topic. We are looking for cases in which mixing looks favorable. Unfortunately we have no results which I could report on.

4. Other baryon/lepton violating processes

In this section I would like to mention a number of other baryon number and/or lepton number violating processes, which have been considered so far. One such process is neutron - antineutron oscillation which is not allowed in the mimimal SU(5) by the B-L conservation rule but well may happen in some GUT, as was emphasized by Glashow.[27] One knows, from an experiment in progress at Grenoble, that the $n-\bar{n}$ oscillation time is longer than 10^5 sec. In SO(10) model the $n-\bar{n}$ is allowed but the estimate of $\tau(n-\bar{n})$ is very model dependent as multi Higgs exchange is expected to give the dominant contribution.[28]

Hydrogen - antihydrogen, $H-\bar{H}$, oscillation is allowed in the minimal SU(5) but is suppressed as it involves the exchange of more than one superheavy gauge boson (or Higgs).

The oscillation $H-\bar{H}$ has an interesting astrophysical signature: the produced $\bar{H}$ annihilates and gives energetic photons. Feinberg et al[29] using the effective hamiltonian

$$\mathcal{H} = c\,\frac{G_F}{\sqrt{2}}\;\bar{e}\gamma^{\mu}(1-\gamma_5)p\;\bar{e}\gamma_{\mu}(1-\gamma_5)p,$$

and from absence of the expected annihilation photons give $C \lesssim 10^{-21}$. These authors also quote upper limits on a number of other processes such as

$$p+n \to e^+ + \nu$$
$$n+n \to e^+ + e^-$$
$$p+p \to e^+ + e^+$$

Recently the process $n \to \bar{p}\, e^+ \nu_e$ has been considered by Gupta et al[30], who calculate the rate in a pole model (Fig. 9).

In this model, using the result $\tau(n-\bar{n}) \gtrsim 10^5$ sec., they estimate $\tau(n\to\bar{p}\, e\ \nu) > 2\times 10^{49}$ years, and improve by several orders of magnitude the limit given by Feinberg et al.[29]

There are also other processes such as muonium - antimuonium oscillation, $\mu\to e\gamma$, etc. which might happen in some GUT but I have no time to discuss them. The decay $\mu\to e\gamma$ is specially interesting as the limit will be improved in the near future.

5. Concluding remarks

One of the greatest virtues of GUT, in my opinion, is that they have triggered the interest of the physics community to seriously undertake a new kind of fundamental but very difficult experiments. Without GUT few would have had the desire or courage to do such null-looking experiments. For all we know the experiments may turn out to find nothing but limits. Nevertheless they will widen our horizon considerably. If we are lucky baryon and lepton number violation will be found with good enough rates for us to study. Such a state of affairs as J.J. Sakurai would have put it "would promise us many exciting developments which will keep us busy until the dawn of the 21st century".[31] We shall miss him, his joy and enthusiasm.

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THE FIRST MILLION YEARS

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ABSTRACT

. This talk offers a predestrian walk in the early universe. It is argued, that several puzzling astrophysical phenomena can be explained on a straightforward way by assuming a nonvanishing rest mass for the heaviest stable neutrino in the region of 30 eV .

Expanding Universe

In the past centuries research performed in physical laboratories clarified the fundamental equations of motion in form of differential equations:

$$\frac{dy}{dt} = F(y) \quad , \tag{1}$$

where F describes the interactions among the constituents of matter. In order to get the present state $y(t)$ of the universe, beside the differential equations (1) one needs the initial values

$$y(0) = a \tag{2}$$

as well. This could not be learned by experimenting in laboratories. One tried to borrow it from theology, from philosophical arguments. Scientists of the last century were worried about this alien element in science. To be able to explain the material world by itself, they tried to postulate a steady state for the universe, at least on large scale

$$\frac{dy}{dt} = 0 \quad .$$

In this way they hoped to deduce the specific steady state y_0 of the actual universe by solving the time independent equation

$$F(y_0) = 0 \quad .$$

Detailed investigations have shown, that this static world picture contains some inherent difficulties, which were expressed by three famous paradoxa of the history of science: the Olbers paradoxon (darkness of the night sky), Clausius paradoxon (heat death) and Seeliger paradoxon. The third paradoxon is connected with the mathematical fact, that the classical potential equation of gravity,

$$\nabla^2\phi = 4\pi\, G\, \rho$$

has a homogenous-isotropic-static solution ϕ only in the case of $\rho = 0$. Expressed on a more picturesque way: any mass M is the source of $4\pi\, GM$ gravitational lines of force. All existing masses M are positive, so these lines do not cancel each other. If gravitational force is present in space, the masses must accelerate along these lines. This is the theoretical background of the simple observation, that a stone cannot stand still steadily near to a large body. It must fall into, orbit around or run away from the other body. Only an empty universe can be static.

These paradoxa produced a lot of headache at the turn of the century. A shock treatment by history helped to break the mental barrier (Petrograd, 1922). Alex Fridman was the first to conclude from physical equations of motion, that our world cannot be and is not static. It must be either contracting or expanding.

For sake of simplicity we shall postulate the homogenity and isotropy of the world on large scale. (It will be shown later, that this assumption has been confirmed by astronomical observations to a high degree of accuracy.) In such a uniform, but unsteady world the galaxies move like grains on a stretching rubber sheet (Figure 1).

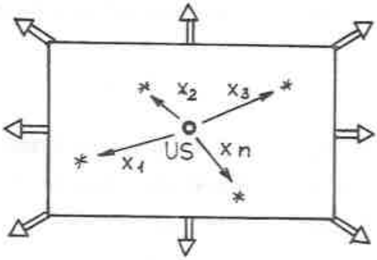

Figure 1.

If we are stretching the sheet uniformly, the coordinate of a specific grain can be written as

$$\underline{x}_n(t) = R(t) \cdot \underline{c}_n \quad . \tag{3}$$

Here $\underline{c}_n$ is a constant, characterizing the n-th galaxy and $R(t)$ describes the universal time dependence of the sheet, coming from the way we are stretching it. (If the distance $\underline{x}_1$ of the first galaxy doubles during a time interval $t_2 - t_1$, this is the consequence of $R(t_2) = 2R(t_1)$, so the distance of all the other galaxies will be doubled during the same time.) From eq. (3) one gets the velocity of the galaxy:

$$\underline{v}_n(t) = \dot{R}(t) \cdot \underline{c}_n \quad , \tag{4}$$

of by eliminating $\underline{c}_n$

$$\underline{v}_n(t) = H^{-1} \cdot \underline{x}_n \quad , \tag{5}$$

where

$$H = R(t) / \dot{R}(t) \tag{6}$$

is a factor of the dimension of time. The proportionality of the velocity of galaxies to their distances was discovered by Edwin Hubble (Mount Wilson, 1929). The general red shift in the spectra of galaxies indicated, that the galaxies were running away, each galaxy from any other. The material universe is in the state of general expansion. The best present value of the Hubble factor is

$$H^{-1} = \frac{15 \pm 3 \text{ km s}^{-1}}{10^6 \text{ light years}} \quad . \tag{7}$$

This number has an interesting interpretation. By rewriting Hubble's Law (5) in the form

$$\underline{x}_n = \underline{v}_n \cdot H$$

one can say, that a galaxy with a higher speed has a larger distance. Assuming a steady velocity, the present separation can be deduced from a state of zero separation, if the elapsed time is H . From eq. (7)

$$H = (21 \pm 4) \cdot 10^9 \text{ years} = 21 \pm 4 \text{ Gy} \quad . \tag{8}$$

This origin has the nickname "Big Bang". The actual time passed since the Big Bang is actually shorter, because the mutual gravitational attraction decreases the speed of galaxies with time:

$$t < H = 21 \pm 3 \text{ Gy} \quad .$$

A galaxy at a distance x_n is attracted by the mass enclosed by a sphere of radius x_n , with us at the centre (Figure 2).

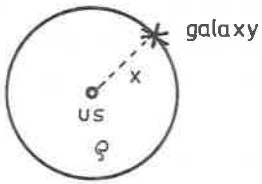

Figure 2

$$\ddot{x}_n = -\frac{G}{x_n^2}\left(\frac{4\pi}{3} x_n^3 \rho\right) \quad . \tag{9}$$

G is the Newtonian constant of gravity, ρ is the homogenous mass density of the universe. By substituting the expression (3) into the equation (9), one gets

$$\ddot{R}(t) = -\frac{4\pi G}{3} \rho(t) R(t) \quad . \tag{10}$$

Here the time dependence of $\rho(t)$ is related to that of $R(t)$ through the theorem of mass conservation:

$$M = \frac{4\pi}{3} R(t)^3 \rho(t) = \text{const.} \quad , \tag{11}$$

so

$$\ddot{R}(t) = - \frac{GM}{R(t)^2} \quad . \tag{12}$$

(The same equation is given also by the general relativity.) Eq. (12) can be integrated:

$$\frac{1}{2} \dot{R}(t)^2 - \frac{GM}{R(t)} = E = \text{const.} \quad . \tag{13}$$

This first order differential equation is soluble by the separation of variables. E.g., for $E = 0$ one gets

$$R(t) = (9GM\, t^2/2)^{1/3} \quad . \tag{14}$$

The corresponding time dependence of the density is

$$\rho(t) = \frac{3M}{4\pi\, R(t)^3} = \frac{1}{6\pi\, Gt^2} \quad . \tag{15}$$

(This is valid for a cool "dust": nonrelativistic mass conservation, negligible pressure.) From formula (14) one gets

$$H = R(t) / \dot{R}(t) = \frac{3}{2} t \quad , \tag{16}$$

so if the Hubble factor is taken from (8), the present age of the universe is

$$t = \frac{2}{3} H = 14 \pm 2 \text{ Gy} \quad , \tag{17}$$

assuming $E = 0$.

$E > 0$ gives faster expansion, $E < 0$ results in a turnover of expansion into collapse. In the latter case the world has only a finite life span. (And the space has a finite volume, according to general relativity. $R(t)$ is the radius of the homogenous curvature.) (See Fig. 3.) The universe is infinite in space and time, if

$$E = \frac{1}{2} \dot{R}^2 - \frac{4\pi G}{3} \rho\, R^2 \geqq 0 \quad , \tag{18}$$

i.e. if

$$\rho \leqq \rho_{crit} \quad ,$$

where the value of ρ_{crit} can be obtained from the measured red shift factor (7):

$$\rho_{crit} = \frac{3}{8\pi\ GH^2} \cong 10^{-26}\ \mathrm{kg/m^3} \quad , \tag{19}$$

corresponding to 10 protons per m^3. If ρ is smaller than this critical value, the kinetic energy dominates over gravity in eq. (18), so the expansion will last for ever. If ρ is larger than the critical value (19), then gravity dominates over motion, a total collapse will come within a finite time.

The optical mass density of galaxies is about 1.2% of the critical density. The ionized gas of the intergalactic space may add about twice as much. So the nuclear mass density ρ_{nucl} may be 3 of 4% of the critical value. If this is the only mass present, the universe has a positive energy constant, motion dominates over gravity, space will expand forever.

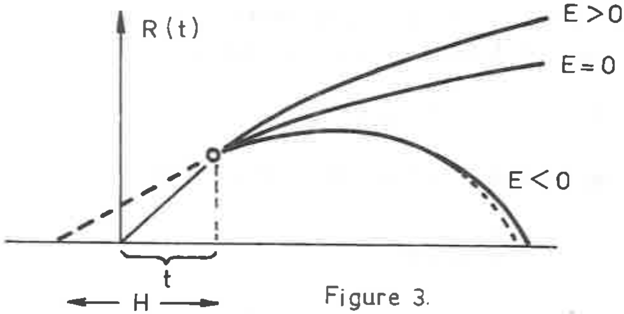

Figure 3.

(The actual mass density cannot be much higher than the critical density. A high density would give an age t too short to explain some facts of stellar dynamics and it would produce a considerable decrease of expansion speed with time. The expansion rate would be considerably smaller now than it was one billion years ago, so the observed red shift factor H^{-1} of nearby galaxies would be smaller that that of galaxies being one billion light years away. There is no indication for such a strong distance dependence, so $\rho < 2\,\rho_{crit}$.)

Hot Universe

If the overall density was extremely high in the early universe, the shortranged nuclear forces had to bind the nucleons to the almost stable medium-heavy nuclei. But the present universe is dominated by light elements (73% H, 26% He). The only explanation is, argued G. Gamow (1948), that the early universe was very hot, preventing the formation of composite nuclei.

Penzias and Wilson, radio engineers of the Bell Lab, observed (1964), that the universe is filled isotropically with an electromagnetic noise, corresponding to $T = 2.7$ K temperature. This means 550 million photons per m^3, compared to less that 1 proton or electron per m^3. This many photons cannot be of stellar origin. They originated in the early hot universe.

The present nuclear mass density is $\rho_{nucl}(t) = 10^{-28} kg/m^3$. The present radiation density, as calculated from the Stefan-Boltzmann law

$$\rho_{rad} = aT^4 / c^2 \quad , \tag{20}$$

is $10^{-31} kg/m^3$, negligible compared with the mass of stars and ion clouds. But this was not so all the time. Light means a sinus wave drawn on a stretching rubber sheet (Figure 4).

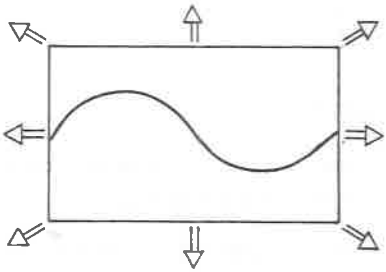

Figure 4.

Its wave length increases proportionally to the dilatation factor $R(t)$:

$$\lambda \sim R(t) \quad . \tag{21}$$

That means that the temperature of the thermal radiation drops (due to the scaling property, expressed by Wien's displacement law):

$$T \sim R(t)^{-1} \quad . \tag{22}$$

Taking this in eq. (20) into account, the time variation of the radiation density is given by

$$\rho_{rad}(t) \sim R(t)^{-4} \quad . \tag{23}$$

The time dependence of the stellar (nuclear) mass density is described by the nonrelativistic mass conservation (11):

$$\rho_{nucl} \sim R(t)^{-3} \quad . \tag{24}$$

Going back to smaller and smaller $R(t)$ values, we reach a time when radiation density dominated over nuclear mass density. In this early "radiation era" the density ρ of eq. (13) was proportional to R^{-4} and it was big. The constant term E was negligible compared to the density term:

$$\frac{\dot{R}^2}{R^2} = \frac{8\pi G}{3}\rho \quad . \tag{25}$$

If ρ is proportional to R^{-4}, then

$$\dot{\rho}/\rho = -4\dot{R}/R = -4(8\pi G\,\rho/3)^{1/2} \quad .$$

The equation

$$\rho^{-3/2}\dot{\rho} = (128\pi G/3)^{1/2}$$

can be integrated directly:

$$\rho(t) = \frac{3}{32\pi G\,t^2} \tag{26}$$

in the radiation-dominated universe. (Here the constant of the second integration is zero, if the time of Big Bang is $t = 0$.)

If the early universe contained n different types of radiation (n different sorts of relativistic particles), eq. (20) gives

$$\rho(t) = naT^4/c^2 \quad . \tag{27}$$

Combining eq. (27) with eq. (26), the temperature of the early universe turns out to be

$$T(t) = \left(\frac{3c^2}{32\pi G\,a\,n\,t^2}\right)^{1/4} \quad . \tag{28}$$

By substituting the values of fundamental constants, for $n \gtrsim 1$ one gets the orienting formula

$$kT \simeq \frac{1\ \text{MeV}}{\sqrt{t_{sec}}} \quad . \tag{29}$$

(Here k is the Boltzmann constant, $a = 8\pi^5 k^4/15c^3h^3$.) The thermal history of the universe is shown in Figure 5.

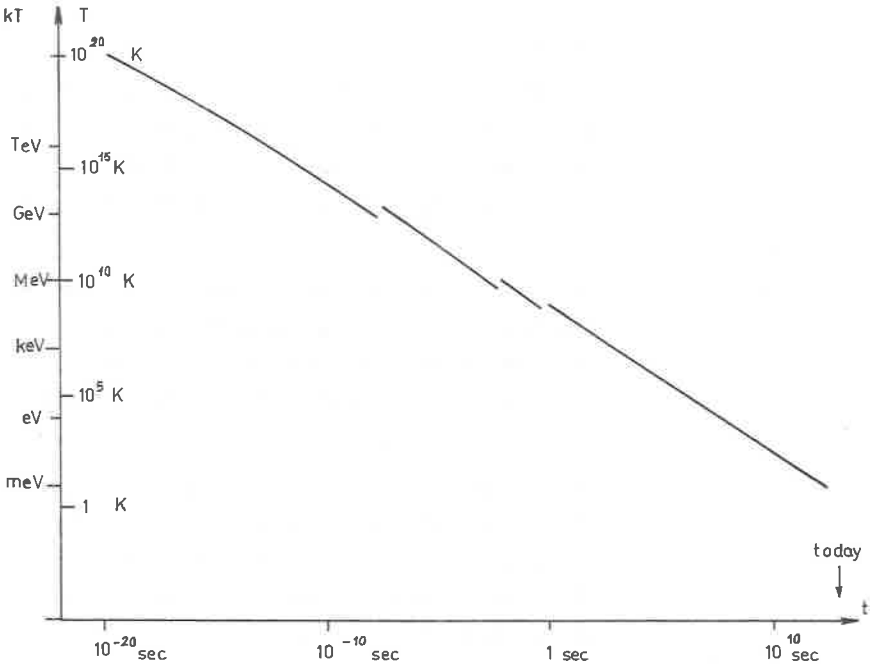

Figure 5

According to the formula (22) the rate of the geometric expansion changes as

$$R \sim t^{1/2} \quad .$$

Going back enough in time, one can reach any high values of density and temperature. The formulas (26) and (28) are free of arbitrary constants of integration. This answers the century-old question about the initial condition. The laws of matter, as we learned them by lab experiments, do not have singularity-free solutions at all. (Even in an inhomogenous-anisotropic universe the time axis of the solution cannot be extended into infinity on both directions, as proved by Hawking and Penrose.) The singularity, following from the very nature of these laws, offers a natural initial condition, which is rather free of arbitratiness.

The First Millisecond

When $t < 10^{-3}$ s , then $T > 0.3 \times 10^{12}$K and $kT > 30$ MeV . This early era was dominated by hadrons and their strong interactions. Under these conditions the physics is not well understood, which makes the era difficult and exciting. People call it "the poor man's high energy laboratory".

The universe started from a specific, but unknown microstate of high energy density. At such extrem densities the collisions were so frequent and vehement, that the universe reached its thermal equilibrium very fast (within Planck-time), forgetting the peculiarities of its initial microstate.

At such a high temperatures all kinds of particles were produced and destroyed immediately. While $mc^2 < kT$, the unstable particles were as abundant as the stabel ones. When the temperature dropped below mc^2/k , the unstable particles were not created any longer. In thermal equilibrium their concentrations vanishes as

$$\exp(-mc^2/kT) \ll 1 \quad , \quad \text{if} \quad kT < mc^2 \quad . \tag{30}$$

But the decay probability is a given number, so the particles need a certain time to be decayed. Till then an overcooled gas of the unstable particles is present: the universe departs from thermal equilibrium for a while. Later the unstable particles disappear by irreversible radioactive decay and so the thermal equilibrium gets restored.

This is, how heavy bosons (postulated by GUT) survived for a while, later they decayed charge-asymmetrically on an irreversible way, turning over the original quark-antiquark balance of the universe.

The quark excess is calculable: it is very small compared to the number of photons. (The quark excess ment a positive baryonic charge, which is carried by nucleons in our present universe.) The presence of protons, however, contradicts the condition of thermodynamical equilibrium. But the proton life time expected by GUT is so long, $\tau_p \simeq 10^{31}y$, that the present world still contains these transient particles. The number of protons is low compared to the number of photons:

$$N_p / N_\gamma = 10^{-9\pm1} \quad , \tag{31}$$

but due to their large mass they play now an important role in the gravitational history of the universe.

The jumps in the straight line of Figure 5 are due to changes of n . The superheavy bosons decayed. Quarks condensed into hadrons. Hadrons decayed into leptons. Electron-positron pairs annihilated into photons. Finally, only photons and neutrinos survived, with a slight proton-neutron-electron contamination.

The First Second

Between $t = 10^{-3}$ seconds and $t = 10^5$ years the average energy per particle decreased from $kT = 30$ MeV to $kT = 1$ eV . All the short lived particles had already disappeared. So this energy region is a world with well-known and simple linear physics. The conditions of thermal equilibrium make the reconstruction of this era straightforward and easy.

Below 30 MeV the existing particles were photons, electron pairs and neutrinos (with a negligible contribution of nucleons, due to the tiny charge asymmetry created in the previous era). All these particles were in thermal equilibrium with each other. All the neutrinos exchanged energy with the elctrons via weak neutral current interactions $(\nu_i + e^- \rightarrow e^- + \nu_i)$. But as density and energy decreased, the weak collisions became less and less frequent. When the temperature dropped below $kT = 10$ MeV , the average collision time for neutrinos grew larger than the age of the universe. This meant that each individual neutrino got decoupled from the other components of the plasma. The only change the neutrinos suffered was the increase of their wavelength (Figure 5). But due to the absence of any scale, the dilatation $\lambda \sim R(t)$ initiated a cooling $T \sim R(t)^{-1}$. In spite of having been decoupled the energy spectrum of neutrinos could be characterized by a temperature T_ν which was equal to the temperature of photons T_γ.

At the end of the first second the temperature cooled below 1 MeV. The elctron positron pairs annihilated into photons, but they were not regenerated, because the collision energies were not enough for pair creation any longer. The energy of the electron pairs heated the photon gas up with respect to the disconnected neutrino gas. For a more quantitative orientation, let us imagine an adiabatic transition of a photon-electron-positron gas with entropy

$$S_1 = S + S_{e^-} + S_{e^+} = \frac{4}{3}\, aT_1^3V + \frac{4}{3}\cdot\frac{7}{8}\, aT_1^3V + \frac{4}{3}\cdot\frac{7}{8}\, aT_1^3V$$

to a photon gas of equal entropy:

$$S_2 = \frac{4}{3}\, aT_2^3V \quad .$$

(The factor 7/8 takes into account the Pauli exclusion for fermions. Otherwise relativistic electrons behave rather similarly to photons.) Putting the two entropies equal, the photon temperature T_2 after annihilation can be expressed by the photon temperature T_1, which would be valid if the annihilation had not taken place:

$$T_2 = \left(\frac{11}{4}\right)^{1/3} T_1 \quad .$$

Since this transition the neutrinos are a bit cooler $(T_\nu = T_1)$ than photons $(T_\gamma = T_2)$:

$$T_\nu = \left(\frac{4}{11}\right)^{1/3} T_\gamma \quad . \tag{32}$$

The present proton temperature is measured: $T_\gamma = 2.7$ K. The corresponding neutrino temperature can be calculated from eq. (32): $T_\nu \simeq 2.0$ K. We have today $N_\gamma = 550$ photons per cm^3. The corresponding number of neutrinos is

$$N_\nu = r \times \frac{3}{4} \times \frac{4}{11} \times N = 150\ r \text{ neutrinos per cm}^3 \quad . \tag{33}$$

Here r is the number of the different types of neutrinos ($r = 3$ for ν_e, ν_μ, ν_τ) . 3/4 is due to Pauli exclusion. 4/11 reflects the temperature difference.

The First Three Minutes

In the thermal equilibrium the neutron to proton ratio is expected to be

$$N_n/N_p = \exp - \left(\frac{m_n c^2 - m_p c^2}{kT}\right) \simeq \begin{cases} 1 & \text{if} \quad kT > 1.26 \text{ MeV} \\ 0 & \text{if} \quad kT < 1.26 \text{ MeV} \end{cases} . \tag{34}$$

After the first second, when temperature dropped below 1 MeV , no more neutrons were formed in proton collisions, but the existing ones survived for about 15 minutes, due to their long life time. The presence of free neutrons in a cool environment meant again a transient deviation from thermal equilibrium, which made the formation of the first nuclei possible by capture of these overcooled neutrons:

$$n + p \rightarrow {}^2H,\ {}^3He,\ {}^4He,\ Li,\ Be,\ B \quad .$$

The fusion chain stopped after a few minutes because the free neutrons became either captured or decayed. The quantitative chemical composition of the produced material depended on the concentration of available nucleons among the photons. The present aboundance of ^2H , ^3He and Li , Be , B can be explained only by this early nuclear build-up. (These loosly bound nuclei could not be formed in stars, because they are excellent nuclear fuels and would burn fast to helium.) The observed $N(^2H)/N(^1H) \simeq 10^{-5}$ ratio gives $\rho_{nucl} = 3.5 \pm 2.0\%$ of ρ_{crit} , in excellent agreement with the direct astronomical evidence (Olive).

An other valuable information, coming from the observed deuteron concentration, is that the overall mass density (dominated by radiation) could not be very high, because large mass density would have resulted in a fast deceleration. If the speed of expansion had been too large, it would have left only a time too short for the formation of such a quantity of deuterons. The existing forms of radiation were photons and r types of neutrinos:

$$\rho = \frac{aT_\gamma^4}{c^2} + r\,\frac{7}{8}\,\frac{aT_\nu^4}{c^2} = [1 + \frac{7}{8}\,(\frac{4}{11})^{4/3} r]\,\frac{aT_\gamma^4}{c^2} \quad .$$

The quantity of existing deuterons limits the value of this radiation density: $r < 5$. If all the neutrinos $(\nu_e, \nu_\mu, \nu_\tau \ldots)$ are stable, one cannot expect more than one new neutrino type.

The First Million Years

The first 10^5 years were dominated by radiation:

$$\rho_{nucl} \ll \rho_\gamma + \rho_\nu \quad .$$

The electromagnetic interaction coupled the photons to the few electrons and ions, providing a uniform temperature T_γ . The decoupled neutrinos were characterized by a lower temperature T_ν . At about $t = 10^5$ y the temperature dropped to $kT_\gamma < 1$ eV , so the protons and other light nuclei became able to capture electrons, to form stable atoms. These neutral atoms were transparent for electromagnetic waves, so light had got decoupled from heavy particles. The photons of the radio noise, discovered by Penzias and Wilson, originated to this time. Since then they were distorted only by the geometric dilatation of wave length. The high isotropy of the present radiation $(\Delta T_\gamma / T_\gamma < 5.10^{-5})$ indicates that the plasma was also very homogenous at the time of atom formation!

At about the same time the fast decreasing radiation density $(\rho_{rad} \sim R(t)^{-4})$ got smaller than the nuclear density $(\rho_{nucl} \sim R(t)^{-3})$, because photon mass is devaluated by the cosmic red shift (21). From this time on the universe was dominated by nonrelativistic matter, resulting in a dust-like expansion $R \sim t^{2/3}$ instead of the earlier radiation-like expansion $R \sim t^{1/2}$. The main actor is the cool (a bit dirty) hydrogen gas. This is a new era, again with well-known physics, but it is far from being simple. The complication comes from the non-linearity of the hydrodynamical equations with gravitational interactions.

Before the neutral era any local inhomogenity (produced by statistical fluctuations) dissolved quickly because radiation transferred mass and energy from high density regions to places of low density. (The only surviving inhomogenities were larger than the horizont: time was then not yet long enough to level up on such large distances.) But with the advent of a neutral era the gravitational instability came into force. A mass concentration of radius L became stable and started growing by gravitational attraction, if its surface gravity exceeded the thermal motion (Jeans stability condition, Figure 6):

$$\frac{G}{L}\left(\frac{4\pi}{3} L^3 \rho\right) m_p > kT \qquad (35)$$

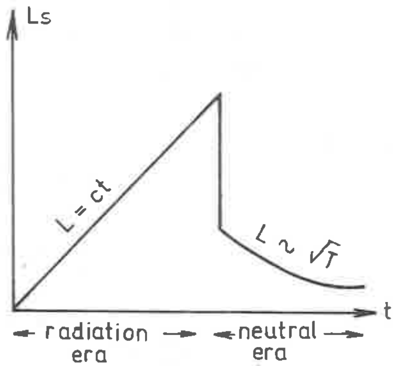

Figure 6.

The smallest of the stable masses was formed most frequently by statistical fluctuations. This critical size and mass is

$$L_{crit} = \left(\frac{3kT}{4\pi\, G\, \rho}\right)^{1/2} , \qquad M_{crit} = \frac{4\pi}{3} L^3 \rho = \left(\frac{3k^3 T^3}{4\pi\, G^3 \rho}\right)^{1/3} . \qquad (36)$$

ρ is the density and T is the temperature of the neutral gas. At the time of atom formation $T_{then} = 10^4 K$ and

$$\rho_{then} / \rho_{today} = (T_{then} / T_{today})^3 = (10^4 / 2.7)^3 \quad ,$$

so all quantities are known. One gets $M_{crit} = 10^5 M_{SUN}$. This is the size of globular clusters, it is much smaller than the size of galaxies. There is no way to understand, how such big objects like galaxies and clusters of galaxies could be formed in the neutral era.

An even greater discrepancy concerns the time scale of the evolution of inhomogenities. We have found that $\Delta\rho/\rho$ was $< 5.10^{-5}$ at the beginning of the neutral era. If $\rho = 3.5\%$ of ρ_{crit}, the inhomogenities could increase only by a factor 15 up to now. If $\rho = \rho_{crit}$, the inhomogenities could increase only 1000 fold. But to get dense astronomical objects, the inhomogenity has to reach $\Delta\rho/\rho \simeq 1$, where the nonlinear effects begin to work. To be short: the time of the neutral era seems to be too short for the formation of galaxies (Rosgacheva - Sunyayev, 1981, Figure 7). Galaxy formation supposes a higher nonrelativistic mass density: $\rho \simeq \rho_{crit}$! Atomic matter contributes only $3.5 \pm 2.0\%$ to this, so about 95 - 97% of the nonrelativistic mass of our universe is "missing".

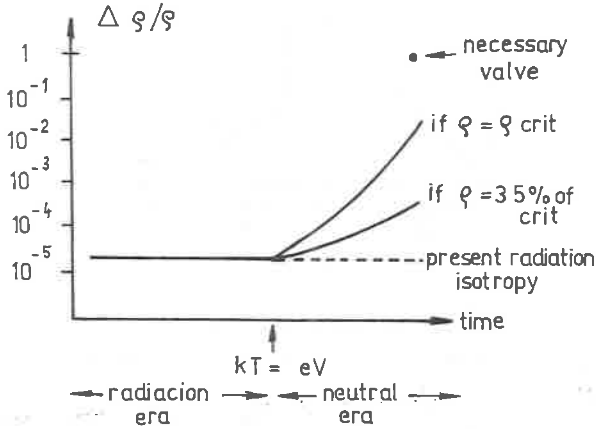

Figure 7.

Thermodynamics forces us to suppose that there are $N_\nu = 150$ million neutrinos per m^3 (of each type) in the universe, compared to N_p 1 proton per m^3. It has been suggested by A.S. Szalay and G. Marx (1976) that the missing mass can be explained by assuming a non-vanishing rest mass at least for one sort of neutrinos:

$$m_\nu = \frac{\rho_{missing}}{N_\nu} \simeq \frac{10^{-26} kg\ m^{-3}}{150 \cdot 10^6\ m^{-3}} = 30\ eV \quad .$$

This proposal created large interest when Lubimov et al. reported (1980) that the β spectrum of ^3H decay indicates an $\bar{\nu}_e$ rest mass of similar value $(16\ eV < m_\nu < 45\ eV)$.

Neutrino Superstars

The cooling process of decoupled particles is described by the law (21). The energy change of the photons is the following:

$$\varepsilon_\gamma(t) = h\ \nu(t) = \frac{hc}{\lambda(t)} \sim \frac{1}{R(t)} \sim T_\gamma(t) \quad .$$

From the decoupling $(T_{dec} = 10^4 K,\ kT_{dec} = 1\ eV)$ the photon energy was shifted to the present value

$$\varepsilon_\gamma(t) = 1eV \times \frac{T(t)}{T_{dec}(t)} = 1eV \times \frac{2.7K}{10^4 K} \simeq 2.7 \times 10^{-3} eV = 2.7meV \quad .$$

If the neutrinos posses a rest mass m_ν in the few eV region, when the temperature dropped below

$$kT_{NR} = \left(\frac{11}{4}\right)^{1/3} m_\nu\ c^2 \tag{37}$$

they become nonrelativistic.

$$\varepsilon_\nu(t) = \frac{p^2}{2m_\nu} = \frac{h^2/2m_\nu}{\lambda(t)^2} \sim \frac{1}{R(t)^2} \sim T_\gamma(t)^2 \quad , \tag{38}$$

so the present value of the average neutrino energy may be

$$\varepsilon_\nu(t) \simeq m_\nu\ c^2 \times \frac{T^2}{T_{NR}^2} = \left(\frac{4}{11}\right)^{2/3} \frac{k^2 T_\gamma^2}{m\ c^2} \quad . \tag{39}$$

E.g. for $m_\nu c^2 = 30$ eV one has now

$$\varepsilon_\nu(t) \simeq 3 \cdot 10^{-9} \text{ eV} = 3 \text{ neV} = kT_\nu \quad ,$$

corresponding to an effective neutrino temperature $T_\nu = 0.00004$ K . Today their average speed is $v_\nu = 6$ km/s . This is smaller than the escape velocity from the Earth! The relic ultracold neutrinos must be very sensitive to gravitational instability.

In the relativistic era the only mass accumulations were the statistical fluctuations of the size of/or larger than the horizon ($L = ct$) . The neutrinos became nonrelativistic when the temperature decreased below the value (37). The number of neutrinos is comparable to that of the photons, it may exceed the number of protons by a factor 10^8 or 10^9 . If $m_\nu c^2 > 10$ eV , in the nonrelativistic era they are the main source of gravity. (The contribution of the redshifted photons became negligible.) The beginning of the nonrelativistic era is characterized by the equation

$$kT_{NR} = k \left(\frac{3c^2}{32\pi \, Gant_{NR}^2}\right)^{1/4} = m_\nu c^2 \quad , \tag{40}$$

where $n = (11/4)^{4/3} + 3 \cdot (7/8) = 6.5$ takes into account the presence of photons $(T_\gamma = (11/4)^{1/3} \, T_\nu)$ and the Pauli exclusion for neutrinos $(7/8)$. Eq. (40) gives the time t_{NR} and the corresponding size L_o of the horizon.

$$L_o = c \, t_{NR} = \left(\frac{3k^4}{32\pi \, Ganm_\nu^4 \, c^6}\right)^{1/2} \quad . \tag{41}$$

This is the size of the neutrino inhomogenities inherited from the relativistic era. The question is, whether such an inhomogenity is stable or not? Let us quote the Jeans stability condition (35) for neutrinos:

$$G \left(\frac{4\pi}{3} L^3 \rho_\nu\right) \frac{m_\nu}{L} > kT_{NR} = m_\nu c^2 \quad .$$

This is satisfied if

$$L > \left(\frac{3c^2}{4\pi G \rho_\nu}\right)^{1/2} = \left(\frac{6k^4}{7\pi \, Gam_\nu^4 \, c^4}\right)^{1/2} \quad .$$

This stability condition is more or less fullfilled by the value (41). So the statistical neutrino concentrations, inherited from the relativistic era $(t < t_{NR})$ turn out to be stable in the nonrelativistic era $(t > t_{NR})$: they survive, start contracting and attract further neutrinos. These "neutrino superstars" were the first astronomical objects in the universe. They were created early in the radiation era.

The radiation era ended later $(kT_\gamma = 1\text{eV})$. The plasma had been homogenized by radiation $(\Delta\rho/\rho \simeq 10^{-4})$, but at the end of the radiation era the newly formed neutral atoms got decoupled from radiation and they found themselves in the inhomogenous gravitational field of neutrino superstars. In the gravitational valleys of the neutrino superstars huge hydrogen clouds were formed.

The original separation of neutrino superstars was given by eq. (41). This distance increased as a consequence of geometrical expansion:

$$L(t) = L_o \frac{R(t)}{R(t_{NR})} = L_o \frac{T_{NR}}{T(t)} = L_o \frac{m_\nu c^2}{kT_\gamma(t)} \left(\frac{11}{4}\right)^{1/3} \qquad (42)$$

$[T_\gamma(t) = 2.7\text{ K}$ today.] The separation $L(t)$ is inversely proportional to m_ν . The biggest objects of the present universe are the superclusters of galaxies. Their average separation is now 10^8 lightyears. If we assume, that the superclusters were born in the womb of neutrino superstars, we can put $L(t) = 10^8$ lightyears, and from this equality we can calculate the rest mass of the neutrino. We get $m_\nu c^2 \cong 30$ eV (A.S. Szalay - G. Marx, 1976). This rest mass makes the present density of the universe equal to the critical density, so it solves the time discrepancy of the galaxy formation as well (Figure 8).

The neutrinos of 30 eV rest mass may make about 95 - 97% of the mass in the universe. They became nonrelativistic at $T_\nu = 300000$ K , so the neutrino superstars were born when the universe was $t_{NR} = 1000$ years old. The young superstar filled the whole horizont, its radius was just $L_o = 1000$ lightyears. Since then the photon temperature decreased by a factor 10^{-5} , correspondingly each distance increased by a factor 10^5 . This gives 10^8 lightyears for the present separation of superstars. The superclusters were born much later $(t = 10^5 y)$ inside the superstars, thus the present distance of superclusters is also 10^8 lightyears.

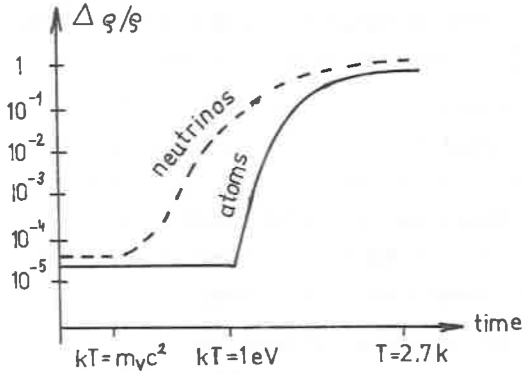

Figure 8.

Voids

Thermal pressure is isotropic. The interplay between pressure and gravity results in spherical objects. But neutrinos do not collide. They do not produce pressure.

Neutrino concentrations originated from statistical fluctuations. Let us follow the fate of a homogenous neutrino gas, which shows velocity inhomogenities. Perturbations smaller than the horizont $(L < ct)$ were eliminated in the relativistic era. The smallest surviving perturbation had a wave length $\lambda = ct$. Let us choose the x axis in the direction, in which the velocity perturbation has the largest amplitude (Figure 9). In a short time this velocity disturbance will make a flat pancake out of the homogenous neutrino cloud (Ya.B. Zeldovich, 1970).

As a matter of fact, an isotropic shrinking would be very exceptional: it supposes equal diagonal elements for the velocity fluctuation tensor (Figure 10). A general tensör (Figure 11) makes a flat sheet (or a rod) out of a cube.

The astronomical evidence confirms this prediction. Galaxies (indicators of the neutrino concentrations) prefer to be distributed in layers and filaments. On the sky we see just opposite of what was expected by the "hot hydrogen gas" hypothesis of the galaxy formation. Instead of bright spheres, separated by dark regions, we find huge empty voids, separated by dense layers and filaments of galaxies (Fig. 12). One can say: these voids are the largest and possibly oldest "things" in the universe (Kirschner et al., 1981; Gregory et al., 1982).

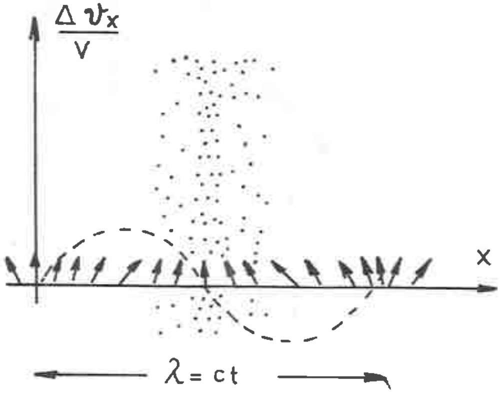

Figure 9.

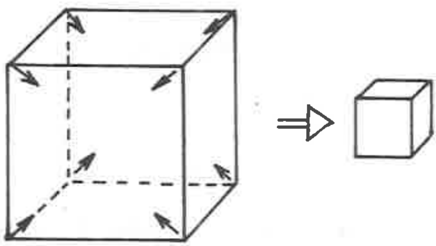

Figure 10.

The detailed evolution of the two component system (collision free neutrinos plus colliding atoms, held together by gravity) is a rather difficult dynamical problem (Bond - Szalay, 1981). This mixture will be granulated into galaxies and stars; to the end galaxies will concentrate in spherical clusters. But this relaxed state is not yet reached everywhere in the universe. The clear message is that the observed large scale structures are still young, in the state of development. They offer a fresh and attractive field for research.

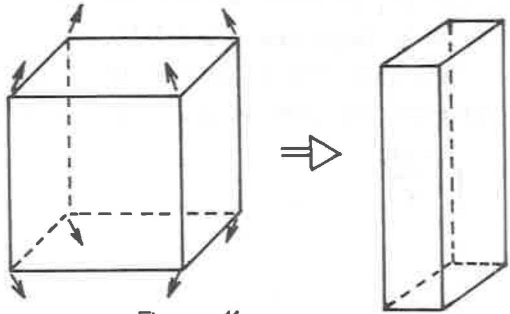

Figure 11.

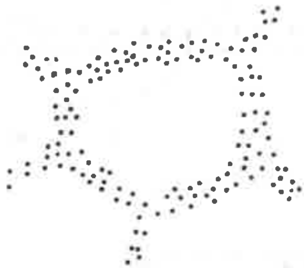

Figure 12.

Galaxies

Some clusters of galaxies look already to be stable relaxed formations, with an equipartition of energy among the member galaxies. The Coma cluster is a fine example for this. If it is stable, its dynamical energy has to be negative:

$$E = \Sigma \frac{1}{2} M v^2 - \frac{G}{R} (\Sigma M)^2 < 0$$

The value of the kinetic energy can be obtained from the observed red shift values, that of the gravitational energy from the counting of member galaxies. But this empirical estimation gives a positive value for E ! In order to explain the stability of the cluster, one has to increase its total mass by an order of magnitude! Well, this extra mass can be explained again by an invisible halo of massive neutrinos around the cluster (A.S. Szalay - G. Marx, 1976). With $m \simeq 30$ eV the neutrinos are able to stabilize the Coma cluster.

The discrepancy between dynamical mass (obtained from its gravitational effects) and optical mass (obtained by counting of stars and galaxies) is a well-known puzzle in galactic astronomy (Peebles, 1973). The discrepancy appears also in the case of single spiral galaxies, like in our Milky Way. By observing the Kepler period of bodies orbiting around the centre of mass at a distance r, one can calculate the total mass $M(r)$ of the galaxy within the radius r. One can extend the method to larger r distances, using satellite galaxies, tidal phenomena of twin galaxies etc. The surprising conclusion is that $M(r)$ increases linearly even beyond the optical boundary of the spiral galaxy (Figure 13, Peebles, 1973). This calls also for an invisible halo around the spiral galaxy, what is made of lighter particles and which produces a rather extended flat potential well, in which the stars of the galaxy are placed (Figure 14, Marx, 1964, 1967, 1976).

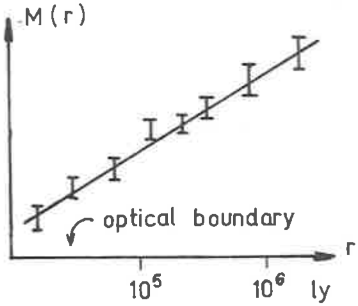

Figure 13.

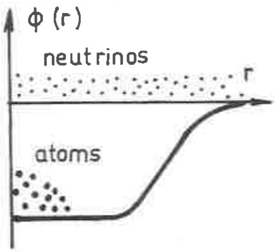

Figure 14.

On scales smaller than a galaxy (stars, globular star clusters) one finds high atom concentrations, which make the contribution of neutrinos negligible. The orbit of the Earth can be explained completely by taking only the attraction of the nuclei and electrons in the Sun into account.

Conclusion

The missing of the mass is a rather general phenomenon in extragalactic astronomy: it comes up at the level of cosmology, at clusters of galaxies and at giant spiral galaxies. These discrepancies can be explained by a simple hypothesis: a certain type of stable neutrinos has a rest mass of cca 30 eV (within a factor 2). This neutrino has to be weakly coupled to other sorts of matter. (Particles with fainter interactions, e.g. gravitons, decoupled much earlier, in the hadron era, so their number is too small to produce observable gravitational consequences.)

This astrophysical conclusion seems to be supported by the lab observations, which indicate a similar rest mass for ν_e. There are double β decay measurements ($^{128}Te/^{130}Te$ ratio in very old rocks), which indicate a Majorana mass for ν_e, but the experimental values are controversial (Hennecke, Kirsten). Lubimov et al. reported $16\ eV < m_\nu < 45\ eV$ limits for the ν_e mass form ^3H decay spectrum (1980). This experiment is being repeated in several laboratories, so we shall have a clear conclusion in one or two years.

The surprising convergence of astrophysical investigations and some laboratory experiments to $m_\nu \simeq 30$ eV is worth of sceptical interest. The unified theories may tolerate $m_\nu \neq 0$, but from the masses of the charged leptons ($m_e = 0.5$ MeV, $m_\mu = 105$ MeV, $m_\tau = 1800$ MeV) one would expect $m(\nu_e) < m(\nu_\mu) < m(\nu_\tau)$. But if both Lubimov and astrophysics are right, then ν_μ and ν_τ either have the same mass as ν_e, or they are lighter, perhaps they are unstable, not as simpleminded theoreticians would expect.

A possible "explanation" of the "convergence" to $m_\nu = 30$ eV can be offered by the old joke (Figure 15). Gravitational arguments can be used for neutrinos only if

$$m_\nu N_\nu > m_p N_p \quad .$$

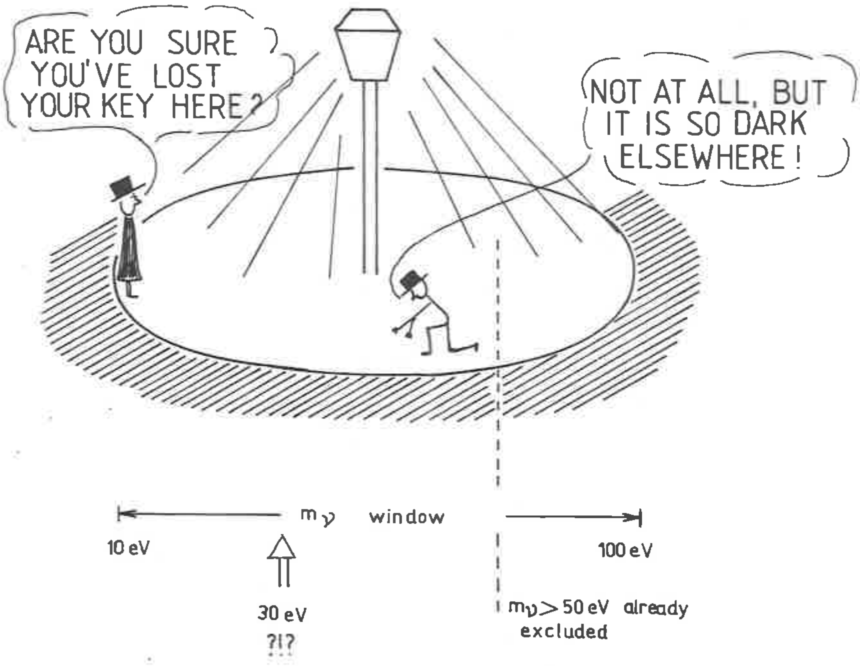

Figure 15.

Now $N_\nu = 10^8 N_p$, $m_p = 10^9$ eV , so astrophysical indications can work only if $m_\nu > 10$ eV . - β spectra are disturbed by complicated molecular fields in the eV region, so they have a similar sensitivity limitation: $m_\nu > 10$ eV . - It has been shown earlier, that cosmological evidence (lower limits on the age of the universe) allowes only the possibility of $m_\nu < 50$ eV (Zeldovich and Gerstein, 1966; G. Marx - A.S. Szalay, 1972; Cowsik - McLeland, 1972). Accurate β spectroscopy has put an upper limit $m_\nu < 55$ eV on the $\bar{\nu}_e$ mass (Berquist, 1979). So only a tight window $10 \text{ eV} < m_\nu < 50$ eV has been left open for an empirical approach to m_ν . If someone is about trying to explain any discrepancy by neutrino mass, the only allowed range is $m_\nu = 30 \pm 20$ eV . One will either find the neutrino mass indeed here within a few years, or one will conclude: $m_\nu < 10$ eV . In the latter case no further progress can be expected concerning m_ν in this century.

If we are lucky and the $m_\nu \simeq 30$ eV value will be confirmed by laboratory experiments, it will give answers to a series of questions concerning the origin of astronomical objects. But it will have a message also to particle physicists. It will say that beside the conventional weakly coupled lefthanded neutrinos we have also the righthanded ones among us, which are very peculiar objects. They seem to communicate with the rest of the world only through the tiny m_ν parameter. This small number brings us a piece of empirical information from a faraway continent of dreams, called "superweak interaction" or "Higgs sector" or "grand unification" or "supersymmetry". It may make the life of theoreticians more exciting again.

The Last Million Years

As it was mentioned above, neutrino superstars were the first things in the universe. What will be the last ones?

If GUT is right, nucleons (and planets, stars, galaxies) will decay in $10^{31\pm1}$ years. After them only photons and neutrinos will be left. But if some neutrinos have a rest mass, the annihilation

$$\nu + \bar{\nu} \to \gamma + \gamma$$

is an entropy generating reaction, so neutrino matter will change slowly into electromagnetic matter (producing photons with a sharp energy $h\nu = m_\nu c^2$).

This neutrino annihilation goes through the electromagnetic form factor of neutrinos. The diagrams of Figure 16 dominate at low energies.

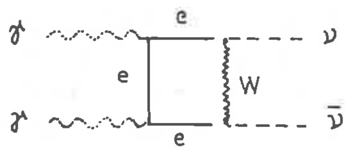

Figure 16.

The corresponding transition matrix elements will be proportional to α

(fine structure constant) and g (Fermi constant of weak interactions). The contribution of Figure 17 is 2 times that of Figure 16, so one has to pay attention only to the last mentioned diagram. In low energy limit it can be simplified to Figure 18.

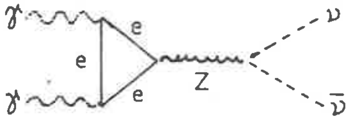

Figure 17.

Figure 18.

This triangle diagram is well known in literature. It gives an effective coupling $\sim \alpha\, g\, F^{\alpha\beta}\, \tilde{F}_{\alpha\beta}\, \overline{\Psi}_\nu\, \gamma_5\, \psi_\nu$, giving the annihilation cross section (Balog - Marx, 1982)

$$\sigma = \frac{\alpha^2}{72\pi} \cdot \frac{g^2 m_\nu^6}{m_e^2} \cdot \frac{c}{v_\nu}$$

for slow neutrinos, numerically

$$\sigma = 10^{-75}\ \mathrm{cm}^2 \cdot \left(\frac{m_\nu c^2}{30\ \mathrm{eV}}\right)^6 \cdot \frac{c}{v_\nu} \quad .$$

($\sigma \to 0$ if $m_\nu \to 0$, and σ is very small if m_ν is as small as 30 eV .) The average life time of a neutrino can be estimated from the formula

$$t = (N_\nu\, \sigma\, v_\nu)^{-1} \quad .$$

Today the average neutrino density is $N_\nu \simeq 150\ \mathrm{cm}^{-3}$, giving a life time of 10^{55} years, much longer, than that of protons. As time passes, however, the average neutrino density decreases as t^{-2} , so the actual life span of a lone neutrino may turn out to be considerably longer. But neutrino superstars are stable bound systems. They have a central density which may exceed the average value several thousand times. So the life time of neutrinos, making a superstar, can be estimated to be of the order of 10^{50} years. Neutrino annihilation is slow compared to proton decay (but it may be fast with respect to black hole decay). Neutrino superstars were the first macroscopic objects and they may be the last ones as well. They always remain invisible for us, but while they exist, they play a key role in making galaxies (and stars and planets and people) and in keeping them together. (The picture is valid if neutrinos have a rest mass of 30 eV indeed.)

After 10^{50}-odd years only zero rest mass particles (photons) will be left over in state of maximum disorder. The universe will return to the complete thermal equilibrium, to the state with maximum entropy once again. (It was in equilibrium already in the first 10^{-40} seconds.) Quarks and protons and neutrinos and stars and life is only a transient nonequilibrium phenomenon, forced to matter by the violent expansion of space. But this tendency for runaway expansion comes from the very structure of the laws, and scientists discovered these laws by investigating matter in terrestrial laboratories.

Is this a sad conclusion? I do not think so. A tragedy in 5 acts (like Hamlet) may be of higher value than an endless (consequently a bit boring) TV serial.

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PHENOMENOLOGY OF SUPERSYMMETRY : A SHORT SUMMARY

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ABSTRACT

We review the main consequences of supersymmetric theories, in particular the experimental constraints on a new neutral gauge boson.

In this paper we intend to present, briefly, some of the main consequences of supersymmetric theories. We refer the reader to more complete articles for additional information and references 1).

1. Supersymmetric theories

Supersymmetric theories associate bosons and fermions in multiplets. This leads one to introduce, next to ordinary leptons and quarks, spin-0 leptons and spin-0 quarks. At the same time the gauge bosons are associated with spin-1/2 particles sometimes called gauge fermions or gauginos : the photon is associated with a spin-1/2 photino, and the gluons with an octet of spin-1/2 gluinos. Moreover with a suitable choice of superfield representations one can construct theories in which the massive gauge bosons are associated, (together with massive heavy fermions) with the physical Higgs bosons ; i.e., charged and neutral Higgs bosons ($w^{\pm}$, z, ...) appear as extra spin-0 states of the massive gauge bosons ($W^{\pm}$, Z, ...).

A continuous or discrete R-invariance allows one to distinguish between ordinary particles (leptons, quarks, gauge and Higgs bosons), which have R = 0, and their superpartners (spin-0 leptons and quarks, photino, gluinos, etc.) which have R = ± 1.

If supersymmetry were conserved, leptons and quarks would be degenerated in mass with their spin-0 partners. Similarly the $W^{\pm}$, for example, would have the same mass m_W as the charged Higgs $w^{\pm}$ and <u>two</u> charged Dirac fermions.

In a globally supersymmetric theory it is necessary to extend the gauge group beyond SU(3) x SU(2) x U(1) (or SU(5)) if all spin-0 leptons and quarks are to receive large masses at the tree approximation. This leads us to consider SU(3) x SU(2) x U(1) x U(1) theories, in which the extra U(1) is responsible for spontaneous supersymmetry breaking and the generation of large masses for spin-0 leptons and quarks 2).

Alternately by using a different method of spontaneous supersymmetry breaking 3) relying on a particular set of interacting chiral superfields one may be able to avoid the introduction of the extra U(1) group and generate large masses for spin-0 leptons and quarks by radiative corrections at the several loop level 4).

Finally the super-Higgs mechanism of supergravity may also be used to generate large masses for spin-0 leptons and quarks 5). This leads to another class of models, such as those of refs. 6,7), with a heavy spin-3/2 gravitino of mass $m_{3/2} \gtrsim 20$ GeV/c^2, and interesting mass relations, such as 7)

$$m \text{ (spin-0 lepton or quark)} = \left| m_{3/2} \pm m_{\ell,q} \right| \qquad (1)$$

and

$$\begin{cases} m^2 (w^\pm) = m^2 (W^\pm) + 4\, m^2_{3/2} \\ m^2 (z) = m^2 (Z) + 4\, m^2_{3/2} \end{cases} \qquad (2)$$

The latter equations relate the $W^\pm$ and Z masses to the masses of the charged and neutral Higgs bosons $w^\pm$ and z which are associated with them under supersymmetry. At the same time the photino and gluinos can also have arbitrary masses, which may reasonably be expected to be of the order of $m_{3/2}$. This leads to the possibility, unpleasant from an experimentalist point of view, that for $m_{3/2} \gtrsim m_W$ all superpartners have large masses $\gtrsim m_W$. (Note, however, that $m_{3/2}$ cannot be much larger than m_W if no fine-tuning is to be performed.)

2. Spin-0 leptons and photinos

What do we know experimentally about the possible existence of the new particles ? Spin-0 leptons produced in e^+e^- annihilation 8) and decaying into a photino or a goldstino (the massless Goldstone fermion of supersymmetry, absorbed to give the extra polarization states of the massive spin-3/2 gravitino 9)) have been systematically searched for at PETRA and, more recently, at PEP 10). This gives the following results :

Table 1 : Limits on the masses of spin-0 leptons.

	Excluded mass range for spin-0 leptons		
	s_e	s_μ	s_τ
CELLO	2 → 16.8	3.3 → 16	m_τ → 3.8 ; 6 → 15.3
JADE	<16	-	4 → 13
MARK J	-	3 → 15	m_τ → 14
MARK II	-	-	m_τ → 9.9
PLUTO	<13	-	-
TASSO (2 spin-0 leptons)	<16.6	<16.4	-

These limits are close to the maximum beam energy available. It is, also, possible to produce a single spin-0 lepton, in association with its spin-1/2 partner and a photino or goldstino. This may allow one to improve the lower limit on the spin-0 electron masses 11). Of course in the future a good way to look for spin-0 particles would be to search for their pair-production in e^+e^- annihilation at the Z pole.

But there is a way to know, relatively soon, about spin-0 electrons and photinos, by looking for the process $e^+e^- \rightarrow$ photino antiphotino. The cross-section is significantly larger than for $e^+e^- \rightarrow \nu\bar{\nu}$, if spin-0 electrons are lighter than $\sim$ 40 GeV/c^2 and photinos not too heavy 12). By looking for single photons emitted in the process

$$e^+e^- \rightarrow \gamma \text{ photino antiphotino} \qquad (3)$$

(see Fig. 1) one may be able, if no signal is detected, to exclude the existence of light photinos coupled to spin-0 electrons lighter than $\sim$ 40 GeV/c^2.

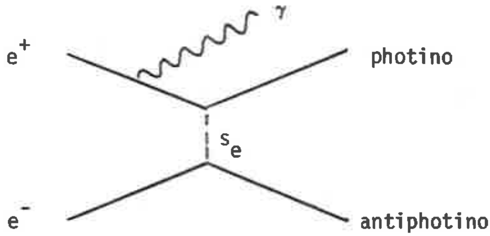

Fig. 1 : A diagram contributing to the radiative production of a photino pair in e^+e^- annihilation.

In a similar way one can study the process $e^+e^- \rightarrow$ gravitino anti-photino, and

$$e^+e^- \rightarrow \gamma \quad \text{gravitino antiphotino} \qquad (4)$$

(see Fig. 2). The corresponding cross-sections are proportional to $G_{Newton} / m_{3/2}^2$. An experiment looking for single photons in e^+e^- annihilation may lead to an improved lower limit $\sim 10^{-6}$ eV/c^2 on the mass $m_{3/2}$ of the spin 3/2 gravitino 12).

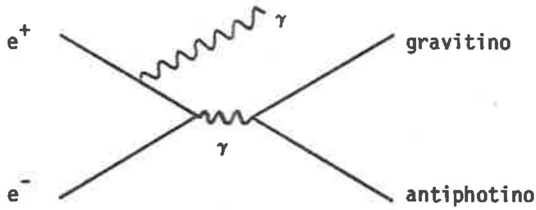

Fig. 2 : A diagram contributing to the radiative production of a gravitino-antiphotino pair in e^+e^- annihilation. (The one-photon exchange leads to a local four-fermion interaction proportional to $(G_{Newton}\,\alpha)^{1/2} / m_{3/2}$.)

3. Gluinos and R-hadrons

The octet of spin-1/2 fermions - gluinos - associated with the gluons might a priori be massless, light, or heavy. These fermions could combine with quarks, antiquarks, or gluons to give new color-singlet states called R-hadrons 13). R-hadrons could be unstable and decay into ordinary hadrons by emitting an unobserved photino or goldstino (gravitino). Limits on the existence of such particles have been obtained from calorimeter and beam dump experiments 13,14). The reinteraction cross-section of photinos with matter 15) is proportional to $m_{s_q}^{-4}$, in which m_{s_q} is the mass of the spin-0 quarks exchanged in the process. As a result the limit obtained, from beam dump experiments, on R-hadron production cross-sections, and therefore on R-hadron masses, depends on m_{s_q}. By subtracting about 1 GeV one gets the corresponding limits on gluino masses. The Fermilab experiment 14) implies :

$$\text{if} \quad m_{s_q} < 40 \ \text{GeV} \qquad m_R > 5 \ \text{GeV} \quad \text{or} \quad m_{\text{gluino}} > (\sim 4 \ \text{GeV})$$

$$\text{if} \quad m_{s_q} < 100 \ \text{GeV} \qquad m_R > 3.5 \ \text{GeV} \quad \text{or} \quad m_{\text{gluino}} > (\sim 2.5 \ \text{GeV}) \tag{5}$$

Gluinos, therefore, may still be relatively light. Their effects on nucleon structure functions and deep inelastic scattering have been studied in detail in Ref. 16).

4. Searching for a new neutral gauge boson

In theories in which the gauge group is extended to SU(3) x SU(2) x U(1) x U(1) the existence of a new neutral gauge boson U may manifest itself in a number of processes 1,17,18) : neutral current effects in neutrino scatterings, e^+e^- annihilation (asymmetry, etc.), parity-violation in the electron-nucleon interaction (SLAC and atomic physics experiments). If the new boson is sufficiently light it could appear as a narrow resonance in e^+e^- annihilation, with a typical electronic width

$$\Gamma_{ee} \sim 100 \ \text{eV} \ m_U^3 \ (\text{GeV}/c^2) \tag{6}$$

or, also, be produced in the decays of kaons, ψ's, Υ's, etc.

a) Why the U might be light

A large number of experimental constraints now exist, which restrict the possible existence of such a particle. A potentially useful remark is that U-exchange amplitudes are proportional to

$$G_F \quad \frac{m_U^2}{m_U^2 + q^2} \tag{7}$$

As a result U-exchanges do not affect the neutral current phenomenology if the U is sufficiently light, i.e. if m_U is smaller than the momentum transfer q in the experiment considered. This leads us to consider with special attention the possibility that the U may be light.

b) Radiative decays of the ψ and the Υ

A sufficiently light U could be produced in the radiative decays of the ψ and the Υ. One then runs the risk of a conflict with the results of the experiments which have been looking for single photons emitted in the decays

$$\psi \longrightarrow \gamma + \text{"nothing"}$$
$$\Upsilon \longrightarrow \gamma + \text{"nothing"} \qquad (8)$$

and have excluded the existence of a standard axion 19). In our case however the situation is slightly different 18) : only a fraction of the events in which U is produced may result in events of the type (8), due to the various decay modes of the U particle. While a very light U would be quasistable, a heavier U can decay into $\nu\bar{\nu}$, e^+e^-, $\mu^+\mu^-$ or $q\bar{q}$ pairs, etc. One has

$$B\,(\psi \to \gamma + U\,(\to \text{"nothing"})) \times B(\Upsilon \to \gamma + U \to \text{"nothing"}))$$
$$\simeq (5 \text{ or } 6)\ 10^{-5} \times (2 \text{ or } 3)\ 10^{-4} \times \quad B\,(U \to \text{"nothing"})^2 \qquad (9)$$

The experimental results on ψ and Υ decays 19)

$$\begin{cases} B\,(\psi \to \gamma + \text{unobserved neutrals lighter than 1 GeV/c}^2) < 1.4\ 10^{-5} & (10) \\ B\,(\Upsilon \to \gamma + \text{unobserved neutrals lighter than 4 GeV/c}^2) < .5\ 10^{-4} & (11) \end{cases}$$

lead to a constraint on the product (9). They allow for the existence of a light U, provided its branching ratio into unobserved neutrinos satisfies

$$B\,(U \to \nu\bar{\nu}) \lesssim 25\ \% \qquad (12)$$

Equation (12) may be satisfied as soon as m_U is somewhat larger than 2 m_e, provided the neutrino decay modes of the U are somewhat inhibited. (Remember, in addition, that m_U between $\sim 2\ m_e$ and 7 MeV/c^2 is forbidden by beam dump experiments, while a U lighter than $\sim$ 300 MeV/c^2 might possibly have been produced in the decay $K^+ \to \pi^+ U$ 18).) Since the U couplings to left-handed and right-handed fermion fields are proportional to ($\cos\varphi$ -1) and ($\cos\varphi$ +1), respectively, the neutrino decay modes $\nu_e\bar{\nu}_e$, $\nu_\mu\bar{\nu}_\mu$, $\nu_\tau\bar{\nu}_\tau$ are sufficiently inhibited, compared with the electron decay modes, as soon as

$$\cos\varphi_{\text{lepton}} \gtrsim \frac{1}{2} \qquad (13)$$

Alternately in the absence of such constraints on $\cos\varphi$ the U should be sufficiently heavy to have sizeable $\mu^+\mu^-$ and hadronic decay modes. We shall see that atomic physics experiments imply that the hadronic part of the U current is mostly axial ($\cos\varphi_{quark} \simeq 0$). The hadronic decays of the U would then be those of a 1^{++} particle, forbidden by parity and angular momentum to decay into two pions only. Moreover an isoscalar U would have even G-parity and is forbidden to decay into three pions ; therefore there is a threshold at $m_U = 4\, m_\pi$ for the hadronic decay modes of the U. In that case m_U cannot be much smaller than 1 GeV/c^2.

c) The "invisible U"

It is, however, very important to note that such constraints get relaxed if one introduces, in the theory, additional Higgs singlets, which have the effect of making the U almost "invisible" : the couplings of the U to leptons and quarks are multiplied by a factor $r < 1$. U-exchange amplitudes, as well as U-production cross sections and decay rates, are then multiplied by a factor $r^2 < 1$, which can be small 17,18). As an example the ψ and Υ decay experiments now imply only that, for $m_U <$ 1 GeV/c^2, one should have, instead of (12), the new inequality :

$$r^2 \; B\,(U \rightarrow \nu\bar{\nu}) \lesssim 25\ \% \qquad (14)$$

d) Asymmetry in e^+e^- annihilation

Let us now consider the asymmetry in $e^+e^- \rightarrow \mu^+\mu^-$ (or $\tau^+\tau^-$). If no extra Higgs singlet is introduced (which would make the U effects invisible) and in the absence of Z-U mixing we find the following (simplified) expression for the asymmetry 1) :

$$A \simeq -\frac{3}{2}\,\frac{G_F\, s}{8\pi\sqrt{2}\,\alpha}\left[\frac{m_Z^2}{m_Z^2 - s} + \frac{m_U^2}{m_U^2 - s}\right] \qquad (15)$$

The asymmetry measured at PETRA 20) (at $\sqrt{s} \simeq 34.4$ GeV)

$$A_\mu \simeq -(10.5\ \% \pm 1.2\ \%) \qquad (16)$$

in $e^+e^- \rightarrow \mu^+\mu^-$ agrees well with the expected value in the standard model ($\simeq -9.3$ %), but not with about twice this value. Under the above hypothesis, the experimental result $A_\mu > -7$ % implies

$$m_U < 17\ \mathrm{GeV}/c^2 \qquad (17)$$

In addition, searches for narrow resonances in e^+e^- annihilations exclude the existence of the U in most of the 1 to 8 GeV/c^2 mass interval 18), as well as in the Υ region.

If we confine ourselves, for simplicity, to models with no extra Higgs singlet and no Z-U mixing, we find that the U may exist in the (7 MeV/c^2 $\rightarrow$ 1 GeV/c^2) or (8 $\rightarrow$ 17) GeV/c^2 mass intervals.

e) Atomic physics and constraints on the U current

There are, in addition, other constraints that the couplings of such a U particle should satisfy. Remember, at first, that the axial couplings of the U are universal while its vector couplings are parametrized proportionately to a factor $\cos\varphi$. Neutrino scattering experiments imply that the coupling of a (not too light) U particle to ν_μ should be depressed, i.e. $\cos\varphi_\mu \gtrsim 1/2$. Moreover the U couplings to quarks should be mostly axial, in such a way that U-exchanges do not lead to unacceptable parity violation effects in atomic physics 1,17).

The cesium experiment 21) implies that the weak charge of the nucleus, $Q_W(Z, N)$, is given by 22) :

$$Q_W = -\,57.1 \pm 9.4(\text{stat.}) \pm (\sim 4.7)(\text{syst.}) \qquad (18)$$

One should also include an estimation of the theoretical uncertainty from atomic physics effects ($\lesssim 8.5$), before comparing (18) (see ref.22) with the expected value of Q_W in the standard model

$$Q_W^{st} \simeq -68.6 \pm 3 \qquad (19)$$

For a (not excessively light) U the expression of the weak charge is 17)

$$Q_W(Z,N) = \left[Z\left(1-4\sin^2\theta\right)-N\right] + 3(Z+N)\cos\varphi_q = Q_W^{st} + Q_W^{U} \qquad (20)$$

The extra contribution Q_W^U must satisfy, at the 90% confidence level

$$Q_W^U = 11.5 \pm 31.7 \qquad (21)$$

With Z + N = 133 for the cesium this gives

$$-\,.05 < \cos\varphi_q < .11 \qquad (22)$$

Under the above hypothesis the cesium experiment constrains $\cos\varphi_q$ to be very small : the hadronic part in the U current should be mostly axial. (This implies, as a byproduct, that the two spin-0 quarks s_q and t_q associated, respectively, with the left-handed and right-handed spin-$\frac{1}{2}$ quark fields should have nearly equal masses.) One can then see that the SLAC experiment 17,23) leads to no useful constraint on the factor $\cos\varphi_e$ which parametrizes the vector electronic part in the U current.

We have focused here, for a greater clarity of the discussion, on the simple situation with no Z-U mixing and no extra Higgs singlet, which favors the 7 MeV/c^2 to 1 GeV/c^2 and (8 to 17) GeV/c^2 mass intervals for the new particle. But more generally any mass is possible for the U, and one should remain particularly attentive to deviations from the standard model phenomenology, which might signal the existence of the new particle.

5. Constraints on the gravitino mass

Finally we briefly remind the reader that a number of constraints exist, from particle physics and astrophysics, on the mass of the spin-3/2 gravitino. They are discussed in detail in Refs. 1). A crucial fact is that, for a light gravitino the $\pm$ 3/2 polarization states are almost non-interacting, while the $\pm$ 1/2 polarization states (which behave like the goldstino of global supersymmetry) have interactions of effective strength fixed by the ratio $G_{Newton}/m^2_{3/2}$ 9,15). As a result, gravitinos, depending on their mass, may have decoupled sufficiently early, and be less numerous than neutrinos. By studying the effect of gravitinos on the abundance of primordial helium, and on the total energy density of the universe one finds that for quasistable gravitinos the preferred mass interval is

$$10^{-2}\ \mathrm{eV/c^2} < m_{3/2} < 100\ \mathrm{keV/c^2} \tag{23}$$

On the other hand gravitinos may also be heavy, provided they decay sufficiently early. This leads to the alternate possibility 1,24)

$$m_{3/2} \gtrsim 10^4\ \mathrm{GeV/c^2} \tag{24}$$

These bounds, however, should not be taken too seriously. In particular, the attractive possibility $m_{3/2} \sim m_W$ should not be disregarded.

As we have seen supersymmetric theories offer a rich array of new phenomena. If supersymmetry is an invariance of the world it could be discovered relatively soon. There exist, however, models with no new neutral gauge boson, in which all superpartners have large masses comparable to m_W. In such cases, supersymmetry could still remain well hidden for some time.

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SUPERSYMMETRIC GRAND UNIFICATION

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1. Introduction to SUSY-GUTs

There are several reasons for studying supersymmetric models, perhaps the best of which is "because its there". In fact it has been shown [1] that it is the only possible type of symmetry not yet used in addition to the Lorentz group and internal symmetries based on Lie groups (which commute with the Lorentz group). Supersymmetry relates fermions and bosons and may provide a natural home for the elusive scalars needed to spontaneously create gauge symmetries. It relates matter and radiation, a massive supermultiplet containing spin components (the gauge fields) together with spin ½ fields and spin 0 Higgs scalar. It is presently the only candidate for unifying particle physics and gravity (local supersymmtrey generates gravitational interactions), although the largest viable group with a nontrivial connection between space-time and internal symmetries has a local gauge symmetry So(8) which is too small to accommodate the standard model[2]. As such it is probably relevant only as a theory for constituent fields [3] (at the Planck scale?).

Recently there has been renewed interest in a more modest class of supersymmetric models, namely those based on a direct product of the gauge group G with a global supersymmetry. The reason for this interest is that it may solve the "hierarchy" problem[4] and if so the supersymmetry must be good at low energies O(1 TeV) (low relative to the Planck scale of 10^{18} GeV!). As the hierarchy problem is at the centre of the recent studies of low energy supersymmetric models I will spend some time reviewing it.

The standard model[5),6] of the strong, weak and electromagnetic interactions, based on the local gauge symmetry SU(3)xSU(2)xU(1) is renormalisable and (provided the Higgs boson, H, has mass $\lesssim$1 TeV) perturbatively unitary[7]. Thus radiative corrections are (apart from renormalisation terms) finite and calculable and scattering amplitudes have a good high energy behaviour. Since the model also agrees remarkably well with experiment why do we need to go beyond it? Admittedly there are 18 parameters which, one believes, ultimately will be related, as, for example, are the three gauge couplings in grand unified theories (GUTs), but this is no reason why the SU(3)xSU(2)xU(1) model should not be a good effective theory valid until the grand unification scale M_W ($\simeq 10^{15}$ GeV). However there is a flaw in this picture - the flaw is known as the "hierarchy problem".

Once we accept that ultimately the parameters in the standard model will be related (as in a GUT) we can treat the radiative corrections to them as meaningful and not just to be absorbed in a counter term. For example the graphs in Fig 1 contribute to the Mass M of the Higgs boson H. This (running) mass depends on the scale μ at which it is measured and calculation of the graphs of Fig 1 gives the result

$$M^2(\mu) = M^2(M_X) + \sum_i c_i \alpha_i M_X^2 \tag{1.1}$$

where c_i and α_i are coefficients and couplings depending on M_X. For the necessary result[7] that M^2 (1 TeV) $\lesssim$ 1 TeV we see that this may only be achieved by an unnatural cancellation of the large terms on the RHS of eq(1.1). It is unnatural in the sense that the parameters of the microscopic scale M_X combine to give a low mass ($\approx$ massless) scalar only at the macroscopic scale μ. ie simplicity is only apparent at the macroscopic scale[8].

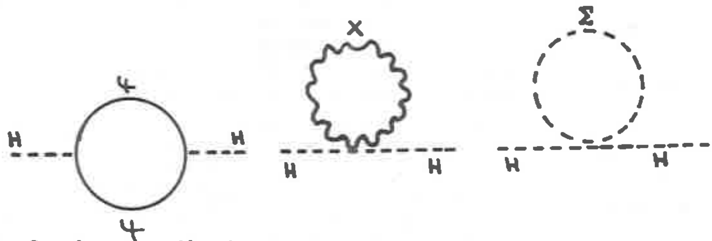


Fig 1. Graphs contributing to the Higgs scalar mass H. Σ and X are the additional Higgs and gauge bosons of the GUT and ψ are the fermions coupling to H.

This means that the standard model will break down at a scale ≈1 TeV, a scale which will be probed by the next generation of experiments! What could this new physics be?

There are three obvious possibilities. One is that the perturbative analysis leading to eq(1.1) fails due to an interaction becoming strong. If this happens only in the scalar sector the GUT predictions may still be valid for gauge and Yukawa couplings will not be greatly affected. In this case the departure from the standard model will be relatively difficult to see until the scalar mode is measured. For example $e^+e^- \to W_L^+W_L^-$ will have strong interactions in the final state as longitudinal W^S are generated by Higgs scalars (clearly this is not a first generation experiment!)

A second possibility is that there are no elementary scalars, as in technicolour models[9] or in other composite models. In this case the calculation leading to eq(1.1) fails for $M_X \approx$ scalar binding energy Provided this energy is $\lesssim \frac{1\ \text{TeV}}{\alpha}$ there will be no hierarchy problem.

Finally it may be that $(\Sigma c_i \alpha_i) = 0$. However it is not enough to arrange this only in leading order perturbation theory for then the next order will require $M_x \lesssim \frac{1\ \text{TeV}}{\alpha^2}$ and so on. If we are to arrange this cancellation involving both fermion and boson loops as in Fig 1 we need a symmetry and the only known symmetry that can do this is supersymmetry[10]. Supersymmetry achieves this[11] by associating all scalars with fermion partners so that the scalar mass is equal to the fermion mass. Since fermion masses can be forbidden by chiral symmetries the associated scalar mass too may be forbidden by a chiral symmetry. As a result in a supersymmetric theory there are cancellations between fermion and boson graphs of fig(1) and eq(1.1) is replaced by

$$M^2(\mu^2) = M^2(M_x^2)\left\{ \frac{\bar{g}^2\ (M_x^2)}{\bar{g}^2\ (\mu^2)} \right\}^{\gamma} \tag{1.2}$$

so that $M^2(\mu^2)$ vanishes if $M^2(M_x^2)$ does. Thus there is no strong constraint on M^2, and the supersymmetric model may be valid up to the GUT scale. Supersymmetry breaking in the gauge nonsinglet sector cannot occur much above 1 TeV, otherwise this evasion of the hierarchy problem will fail.

2.1 A supersymmetric standard model

The simplest sypersymmetric model which can be constructed is a direct product of the internal symmetry gauge group with a global (N = 1) supersymmetry[11,12,13,14]. The basic building blocks are massless supersymmetry multiplets[(15)] of the chiral or vector type as shown in Table (1).

Table 1
Fundamental massless supersymmetric multiplets in N = 1 supersymmetry

Chiral	$\begin{pmatrix}\psi\\ \phi\end{pmatrix}$	2 component majorana fermion 2 real scalar fields (≡ 1 complex)
Vector	$\begin{pmatrix}V_\mu\\ \psi\end{pmatrix}$	2 component massless vector 2 component majorana fermion

Thus the supersymmetric version of the standard model contains at least twice the number of particles needed in the nonsupersymmetric version. The minimal SU(3)xSU(2)xU(1) model can be made supersymmetric simply by assigning the usual states to supermultiplets of the type given in Table 2[12,13].

Table 2
Multiplet structure of the minimal
Supersymmetric SU(3) x SU(2) x U(1) model

Vector Supermultiplets			Spin J
V_G	$g^{a=1\ldots 8}$	Gluons	1
	$\tilde{g}^{a=1\ldots 8}$	Gluinos	½
V_W	$W^{\pm}_{\mu}$, Z_{μ}	W, Z bosons	1
	$\tilde{W}^{\pm}$, $\tilde{Z}$	Winos, Zinos	½
V_γ	A_μ	Photon	1
	$\tilde{A}$	Photino	½

Chiral supermultiplets			Spin J
$S_q, T_{\tilde{q}}$	q_L, q_R	Quarks	½
	$\tilde{q}_L, \tilde{q}_R$	Scalar quarks	0
$S_\ell, S_{\tilde{\ell}}$	ℓ_L, ℓ_R	Leptons	½
	$\tilde{\ell}_L, \tilde{\ell}_R$	Scalar leptons	0
S,T	ψ'_L, ψ''_L	Fermionic Higgs	½
	ϕ', ϕ''	Higgs doublets	0

This leads to the multiplet structure of Table (2). In this table the new superpartners carry the same SU(3)xSU(2)xU(1) quantum numbers as their conventional partners. Note that there are two Higgs doublets (plus their supersymmetric partners) rather than the single one usually included in the standard model. This proves to be necessary if there is to be a reasonable mass spectrum[13]. In addition there will be a massless goldstone fermion, the goldstino, the goldstone particle resulting from the breaking of global supersymmetry. The coupling of supersymmetric particles is simply related to that of their conventional partners. For example the photino coupling follows from the photon coupling to charged particles just by the replacement of the photon by the photino and one of the charged particles by its superpartner -eg see Fig (2).

Fig. 2
Photon and photino coupling to charged fermions and their scalar partners

The other fermions in vector supermultiplets ('inos) have couplings similarly related to the gauge couplings of their partners. The scalar quarks and leptons have gauge couplings with the vector supermultiplets, as in Fig(2), and also couplings to other chiral supermultiplets related to the usual Yukawa couplings necessary to give masses to the quarks and leptons.

The goldstino couples to the partners in a supermultiplet with a strength given by[15]

$$e_g = \pm \frac{\Delta m^2}{d^2} \tag{2.1}$$

where Δm^2 is the mass difference of the superpartners and d is a dimensional quantity related to the scale of supersymmetry breaking ($\approx$ 1 TeV), and the sign depends on the chirality of the fermion considered.

2.2 An SU(5) SUSY-GUT

In the previous section we discussed the spectrum of states necessary to build a supersymmetric (SUSY) version of the standard model. The masses of these states are not predicted until we have a model for supersymmetry breaking, but if the solution to the heirarchy problem is to be preserved they should not have masses greater than $O\left(\frac{300\ \text{GeV}}{\alpha}\right)$. In fact the breaking of supersymmetry proves to be the most difficult part of constructing a viable model. At the tree level supersymmetry breaking in a supersymmetric version of the standard model splits the scalar from fermion masses in a chiral supermultiplet but still leaves one scalar state lighter than its fermion partner. Since we would have seen a scalar electron lighter than the electron this is untenable. One possibility[16] is to extend the standard model to include an extra U(1) factor. This of course means that one cannot then grand-unify in the minimal SU(5). A second possibility[12,13,14,17)] is that radiative corrections to masses may generate an acceptable mass spectrum. Remarkably even in the minimal SUSY-SU(5), this possibly works, generating radiatively reasonably large masses for superpartners, and also triggering SU(2)xU(1) breaking and thus relating M_W to M_{SUSY}, the scale of supersymmetry breaking. We will discuss in detail only this second possibility here. The simplest SUSY-GUT is based on the gauge group SU(5)x[N=1 global supersymmetry]. The supermultiplet structure is analagous to Table 5 with now an SU(5) adjoint representation, Σ, of SU(5) as a vector supermultiplet of gauge bosons and associated fermions, together with three families of ($\bar{5}$ + 10) chiral supermultiplets, ψ_i^α and $\chi_{i\gamma\delta}$ respectively, containing the quarks and leptons together with their scalar superpartners. Finally we have H_1 and H_2, two chiral supermultiplets containing the Higgs fields and transforming as 5 and $\bar{5}$ respectively. The colour triplet and antitriplet components of H_1 and H_2 mediate proton decay and must have a very large mass ($\gtrsim O(10^{10}$ GeV)). The doublets however must be relatively light ($O(10^2$ GeV)) in order to generate SU(2)xU(1) breaking. One of the main differences in constructing a viable model is to generate this split multiplet pattern in a natural way. The fermion partners must also be split (super-symmetry breaking is only of order 10^3 GeV in this sector) and consequently they cannot be the usual quarks or leptons. For this reason we have had to introduce two completely new chiral supermultiplets.

The SU(3)xSU(2)xU(1) content of these SU(5) representations is as given in Table 3. In addition, if we want to break supersymmetry explicitly we need additional SU(5) singlet chiral superfields A,B,C and two further chiral superfields ϕ_1 and ϕ_2 transforming as 5 and $\bar{5}$ under SU(5) respectively.

Rôle	Notation	SU(5) Representation Content
Matter	ψ_a^α, $\chi_{a\ \alpha\beta}$	N_G x ($\underline{\bar{5}}$ + $\underline{10}$)
Higgs	$H_{1\alpha}$, H_2^α Σ_α^β Z	($\underline{5}$+$\underline{\bar{5}}$) $\underline{24}$ $\underline{1}$
O'Raifeartaigh Sector	A,B,C	3 x ($\underline{1}$)
Coupling to Gauge Sector	ϕ_{1_α}, ϕ_2^β	($\underline{5}$+$\underline{5}$)

Table 3: Chiral supermultiplets used in model-building

To specify the model completely it is necessary to write down the Yukawa and scalar couplings. These are related by supersymmetry and it is most convenient to describe them via the superpotential P. P is a gauge invariant function of dimension ≤ 3 constructed from the chiral superfields of the model (but not their complex conjugates). Then the Yukawa and scalar couplings are given in terms of P by

$$\mathcal{L}^{\text{Yukawa}} = \sum_{i,j} \frac{\partial^2 P}{\partial\phi_i \partial\phi_j} \psi_i \psi_j \tag{2.2}$$

$$\mathcal{L}^{\text{scalar}} = \left| \sum_i \frac{\partial P}{\partial\phi_i} \right|^2 = \sum_i F_i^* F_i \tag{2.3}$$

Here ψ, ϕ refer to the (LH) fermion and scalar components of the chiral supermultiplets respectively and the sums over i,j run over all the chiral supermultiplets. The F_i are the auxiliary fields.

For the minimal set of supermultiplets of Table 3 P is taken to be

$$P = P_{OR} + P_{GD} + P_{5M} + P_{5Z}$$

where

$$P_{OR} = \lambda_1 ABM + \lambda_2(A^2 - M^2)C \tag{2.4}$$

$$P_{GD} = \lambda_3 \phi_1{}^a \phi_{2a} A + \lambda_4 AM^2 \tag{2.5}$$

$$P_{5M} = M_{ij}^{(\alpha)} \phi_i^\alpha \chi_{\alpha\beta} H_2^\beta + M_{ij}^{(u)} \varepsilon^{\alpha\beta\gamma\delta\rho} \chi_{i_{\alpha\beta}} \chi_{i_{\gamma\delta}} H_{1\rho} \tag{2.6}$$

$$P_{5Z} = H_1(\gamma_1 Z + \gamma_{24}\Sigma)H_2 + \frac{\beta_3}{3}\Sigma^3 + \frac{\beta_2}{2}\mu\Sigma^2 \tag{2.7}$$

Supersymmetry is broken at the tree level through P_{OR}. P_{GD} couples this supersymmetry breaking to the gauge nonsinglet fields ϕ_1, ϕ_2 and includes a soft term breaking residual R invariance and allowing the gauginos to acquire radiative masses. P_{5M} generates quark and lepton masses once H_1 and H_2 develop vacuum expectation values (vevs). P_{5Z} is responsible for breaking SU(5) and splitting the Higgs multiplets H_1 and H_2.

2.3 The pattern of spontaneous symmetry breaking

In order to discuss the predictions of the model introduced above it is necessary to consider in detail the breaking of the underlying symmetries. This happens in several stages.

At a scale M, which could be as large as M_{Planck}, P_{OR} induces spontaneous breakdown of the N = 1 global supersymmetry. This is clear for it is a feature of supersymmetry that the potential V is positive semi-definite and supersymmetry is broken if and only if $V > 0$. Using

(2.3) we see P_{OR} contributes the terms to V

$$V_{OR} = \left|\lambda_1 AM\right|^2 + \left|\lambda_2(A^2 - M^2)\right| + \left|\lambda_1 BM+2\lambda AC\right|^2 \tag{2.8}$$

For no value of A does (5.10) vanish and so supersymmetry is broken. Without considering radiative corrections the supersymmetry breaking is confined to the supermultiplet A. The remaining fields at tree level are still invariant under SU(5)x[N = 1 global supersymmetry].

We turn now to the effect of P_{5Z}. It is easy to check that the potential coming from P_{5Z} has (at least[18]) three degenerate minima with $V_{5Z} = 0$ (remember supersymmetry is unbroken in this sector). These minima correspond to the gauge symmetries SU(5), SU(4)xU(1) and SU(3)xSU(2)xU(1) remaining unbroken. We assume that nature selects the SU(3)xSU(2)xU(1) minima and that Σ develops the corresponding large vev breaking SU(5) at a scale which is proportional to μ. (Whether this minima is chosen depends on radiative corrections and temperature dependent effects that we do not discuss here[19]). We thus have, at tree level, a SU(3)xSU(2)xU(1)x[N = 1 global supersymmetry] model as introduced in section (5.1).

The breaking of supersymmetry in this SUSY standard model proceeds only through radiative corrections. The graphs of Fig 3 generate a gluino mass of order

$$m_{\tilde{g}} = \frac{\alpha_s}{3\pi}\,\frac{\lambda_3^2\lambda_4}{16\pi^2}\,\ln\left(\frac{m_{\psi A}{}^2}{m_{\phi_A}{}^2}\right)\;A \tag{2.9}$$

where ψ and ϕ represent the fermion and scalar components of A.

However the superpartners of the gluinos, the gluons, are prevented by SU(3)c gauge symmetry from acquiring a mass radiatively. Similarly the Winos and Binos may acquire radiative masses.

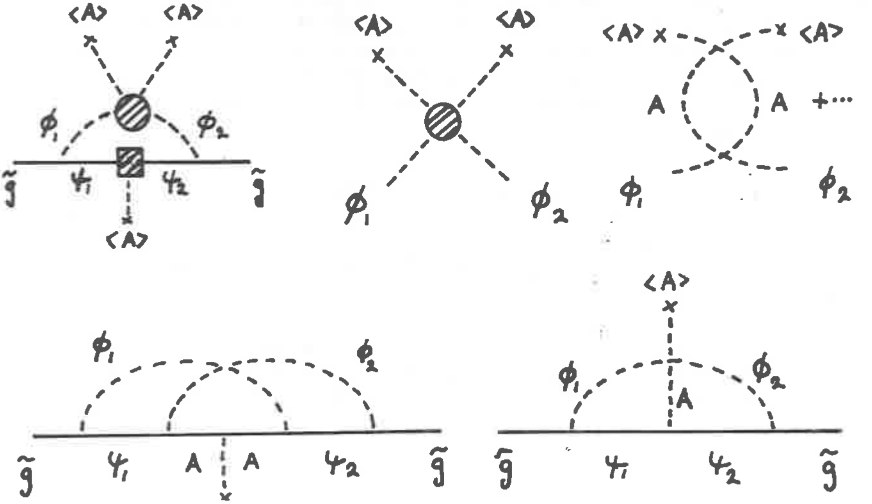


Fig 3. Graphs contributing to gluino in two loop order.

Once the 'inos acquire mass they can generate masses for the scalars via the graphs of Fig 4a; because there is no longer perfect cancellation between fermion and boson contribution. This gives

$$m_{\tilde q}^2 = \frac{4\alpha_s}{3\pi} \ln\left(\frac{m_\phi^2}{m_{\tilde g}^2}\right) m_{\tilde g}^2 \tag{2.10}$$

and $$m_{\tilde l}^2 = \delta m^2_{T,S} \approx \frac{3\alpha_2}{4\pi} \ln\left(\frac{m_\phi^2}{m_{\tilde W}^2}\right) m_{\tilde W}^2 \tag{2.11}$$

While the contribution (2.11) to the light Higgs and slepton masses is positive, there are also negative contributions to the Higgses (but not the sleptons) from the squark loops in Fig 4b which are of order

$$\delta m_{T,S}^2 = -\frac{3}{8\pi^2}\left|h_{U,D}\right|^2 \ln\left(\frac{m_\phi^2}{m_{\tilde q}^2}\right) m_{\tilde q}^2 \tag{2.12}$$

where $h_{U,D}$ are the quark-Higgs Yukawa couplings to up and down quarks respectively.

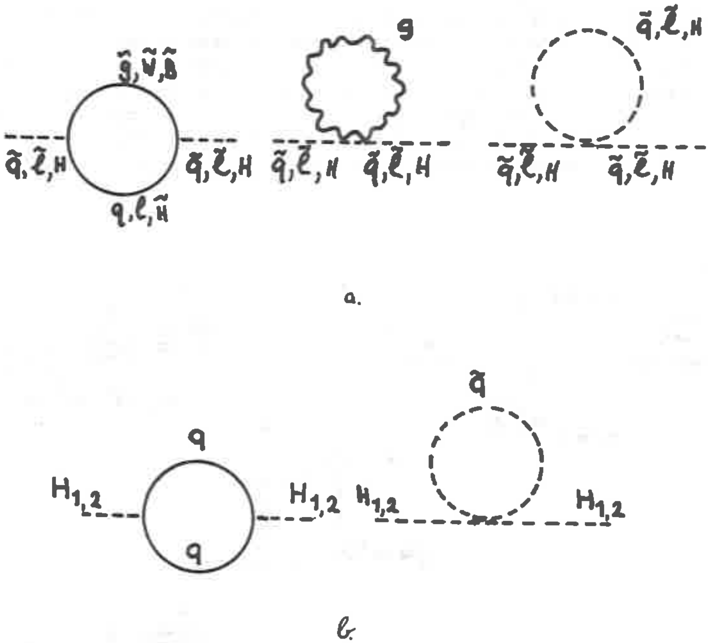


Fig 4. Graphs contributing to scalar masses.

It is easy to check that the negative contribution of eq(2.12) overwhelms the positive piece of (2.11) and triggers spontaneous symmetry breaking of SU(2)xU(1) by the Higgses (and not the sleptons) if we choose

$$h_{U_i}^{2} \gtrsim O\left(\frac{1}{2\ln\frac{m_\phi^2}{m_{\tilde W}^2}}\right) \quad \text{ie } m_t \gtrsim 25 \text{ GeV} \tag{2.13}$$

Once this happens the piece of P_{5Z} in eq(2.7) involving H_1 and H_2, automatically gives the triplet but not the doublet components a large mass as is required. This happens because on minimising the scalar potential Z (the "sliding singlet") develops a vev to make

$$|H_1 \ (\gamma_1 Z + \gamma_{24} \ \Sigma_{55})|^2 + |(\gamma_1 Z + \gamma_{24} \ \Sigma_{55} \ ) \ H_2|^2 = 0 \tag{2.14}$$

This is just the required condition.

To summarise, scalar partners of quarks and leptons and fermionic partners of gauge bosons are heavier than their conventional partners and thus need not have been seen yet. In addition the supersymmetry breaking scale is related via radiative corrections to the SU(2)xU(1) breaking scale: there is no need to input by hand a second scale. It is also easy to introduce a term that automatically splits the Higgs multiplets, as is required if proton decay is not to proceed too quickly. The final mass spectrum of the model is summarised in Table 4 with a more visible version in Fig 5. It is remarkable that radiative corrections have automatically generated a reasonable spectrum.

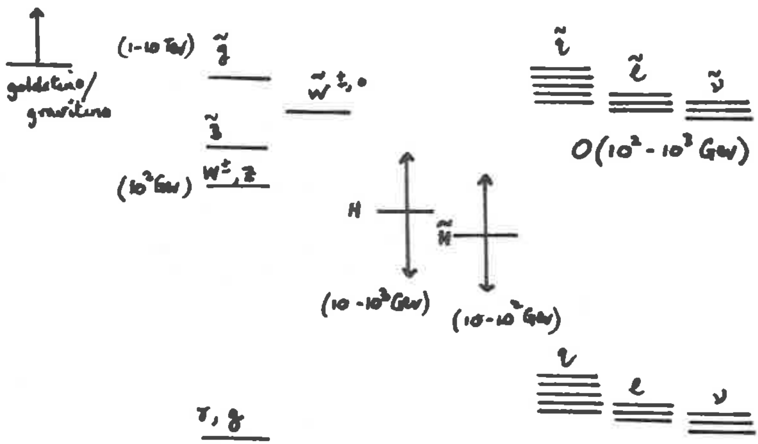


Fig 5. Low Mass spectrum for SU(5) SUSY-GUT.

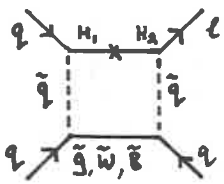
Fermions	Mass	Characteristics
$\tilde{W}^{\pm} + \dots$	$m_{\tilde{W}} + \dots$	Mainly $I = 1$, vector-like neutral weak couplings
$\tilde{H}^{\pm} + \dots$	$\frac{g_2^2 v_1 v_2}{m_{\tilde{W}}}$	Mainly $I = \frac{1}{2}$, vector-like neutral weak couplings decaying into $\tilde{H}^0 + (\ell\nu$ or $\bar{q}q)$
$\tilde{W}^3 + \dots$	$m_{\tilde{W}} + \dots$	Mainly $I = 1$ wino
$\tilde{B}^0 + \dots$	$\frac{5}{3}\left(\frac{\alpha_1}{\alpha_2}\right) m_{\tilde{W}} + \dots$	Mainly $I = 0$ bino
$\tilde{A}^{0\prime} \equiv \frac{v_1\tilde{H}_1^0 - v_2\tilde{H}_2^0}{v} + \dots$	$\frac{g_2^2 v^2}{5 m_{\tilde{W}}} + \dots$	Mainly $I = \frac{1}{2}$ shiggs decaying into $\tilde{H}^{\pm} + (\ell\nu$ or $\bar{q}q)$
$\tilde{S}^0 \equiv \frac{v_2\tilde{H}_1^0 + v_1\tilde{H}_2^0}{v}$	$= 0$	Mainly $I = \frac{1}{2}$ shiggs: end of all supersymmetric decay chains
OR		
$\tilde{S}^0_{\pm\prime} \equiv \frac{1}{\sqrt{2}}(\tilde{S}^0 \pm \tilde{Y}^0)$	$\frac{\lambda v}{2} \pm \dots$	Mixture of $I = \frac{1}{2}$ and $I = 0$, lighter one stable, heavier one quasistable decaying into lighter one $+(\ell^+\ell^-$ or $\bar{q}q)$
Bosons	**$(\text{Mass})^2$**	**Characteristics**
$\frac{v_1 H_2^0 + v_2 H_1^0}{v} + \dots$	$4\lambda^2 v_2^2 + \dots$	$I = \frac{1}{2}$, mainly coupled to charge $-\frac{1}{3}$ quarks
$\frac{v_1 H_1^0 - v_2 H_2^0}{v}$	$\frac{g_2^2 + g_1^2}{2} v^2 + \dots$	$I = \frac{1}{2}$, mainly coupled to charge $+\frac{2}{3}$ quarks
$a \equiv \frac{\eta_1 v_2 - \eta_2 v_1}{v}$	few GeV?	Pseudoscalar coupling, larger for charge $-\frac{1}{3}$ quarks
$\frac{v_1 H_2^{\pm} - v_2 H_1^{\pm}}{v}$	$\frac{g_2^2}{2} v^2 + \dots$	Conventional charged Higgs boson
Y?	$O(\lambda^2 v^2)$?	$I = 0$ scalar and pseudoscalar, very weakly coupled to light quarks

Table 4: Broken Supersymmetric Spectroscopy

3. Phenomenology of SUSY-GUTs

The spectrum of Fig 5 illustrates the expected mass pattern in a minimal SU(5) class of supersymmetric models in which radiative corrections are responsible for splitting the gauge non-singlet supermultiplets and in which R invariance is broken, allowing the gauginos to become heavy. Other schemes, based on tree level breaking, have been considered[20] although they are more difficult to grand unify. The major difference is they have a light goldstino and/or photino, giving rise to a different decay pattern for the new supersymmetric states. This affects the pattern of decay of supersymmetric states and consequently the phenomenological signals. For the models with a light photino these signals have been extensively studied[20]. Here we will mention some of the novel features of the more recent models with the spectra of Fig 5. First, however, there are several predictions insensitive to the low energy structure of the theory. In particular the predictions for $\sin\theta$ and M are changed from the usual grand unified predictions due to the new supersymmetric states (masses 1 TeV) which contribute to the β function. This gives [21),26)] $M_X = \begin{cases} 1 \\ 4 \end{cases} \times 10^{16}$ GeV and $\sin^2\theta_W = \begin{cases} .236 \\ .229 \end{cases}$ for $\Lambda_{\overline{MS}} = \begin{cases} 100 \text{ MeV} \\ 300 \text{ MeV} \end{cases}$, with the minimal SU(5) supermultiplet structure of Table 3 assuming light states are ($\lesssim 1$ TeV) and heavy states have a mass $\approx M_X$. M_X is about 20 times the non-supersymmetric value mainly because the adjoint of fermions introduced in the vector supermultiplet make the β functions smaller in magnitude, slowing up the evolution of the couplings. It may be seen that $\sin^2\theta_W$ is increased from the SU(5) value of .205 and is in poorer agreement with the current best (radiatively corrected) value for $\sin^2\theta_W$ of[22] .210 ± .005. This result is somewhat sensiitive to the masses assumed for the heavy states; for example the colour triplets of H_1 and H_2 had a mass $\gtrsim 10^{10}$ GeV, then[23] $\sin^2\theta = \begin{cases} .225 \\ .218 \end{cases}$

With such a high value for M_X it might be assumed that proton decay will be negligible. This is not the case for new processes occur involving the new superpartners which, in the minimal SU(5) theory, mediate proton decay[24] (see Fig 6).



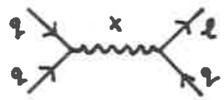


Fig 6. (a) Fermionic higgs contributions to proton decay
(b) Vector boson contribution to proton decay.

These give

$$\tau_p \propto M^2_{\tilde{H}_{1,2}} \; M^2_{\tilde{W},\tilde{B},\tilde{g}} \times \text{Yukawa couplings.}$$

Because $M^2_{\tilde{W},\tilde{B},\tilde{g}} \lesssim 1$ TeV and $M^2_{H_{1,2}} \approx M^2_x$ this rate is potentially much less than the usual SU(5) process of Fig 6(b) which gives $\tau_p \propto M_x^4$. In fact, because of the Yukawa couplings, the rate turns out to be $\approx 10^{30}$ years[25]. However, because the fermionic Higgs of Fig 1(a) preferentially couple to heavy states the dominant decay mode is $P \to \bar{\nu} K^+$ in contrast to the $\pi^\circ e^+$ mode in usual SU(5). Unfortunately this is not an unambiguous prediction of SUSY-GUTs. For example in some models the graph of Fig 1a does not exist and the dominant diagram may be through colour triplet Higgs exchange[26] with dominant decay modes $P \to \bar{\nu}_\mu K^+$, $\mu^+ K^\circ$ or even[23,26] through the original SU(5) diagram of Fig 6(b).

Another sensitive testing ground for unified models is in rare K decays. The most sensitive process is K°-$\bar{K}^\circ$, contributing to the $K_L K_S$ mass difference. In addition to the usual graph of Fig 7(a), there are also scalar quark contributions coming from Fig 7(b).

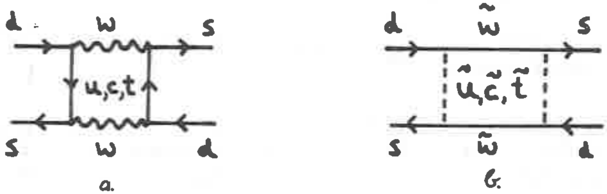


Fig 7. Graphs contributing to ΔS=2 processes.

Evaluating these graphs gives the constraint on the scalar quark masses[26]

$$\frac{m^2_{\tilde{c}} - m^2_{\tilde{u}}}{m^2_{\tilde{c}}} \lesssim 10^{-3}$$

This is a very strong constraint on the mass difference of scalar quarks, but it is satisfied in models of the type in section 2 in which scalar quarks acquire mass radiatively. For them there is a large flavour independent term coming from gauge couplings and a small flavour dependent term coming from Yukawa couplings. Thus the scalar quark mass differences are naturally of order of the quark mass differences .

We turn now to the specific properties and signals of the supersymmetric model of section 2[19,27]. There is a conserved quantum number of the new supersymmetric states introduced in Table 3 which means that these states are only produced in pairs and that they cannot decay entirely into conventional states. Thus all new symmetric states will ultimately decay into the lightest such state; in the case of the minimal model of section 2 it is mainly the fermionic partner of the Higgs (mass $\gtrsim$ 20 GeV)(cf Table 4). In models with the "sliding singlet" Z, there are two such states $\tilde{S}_{\pm}$ and these may be produced via the Z vector boson and may give rise to the characteristic signal of Fig 8.

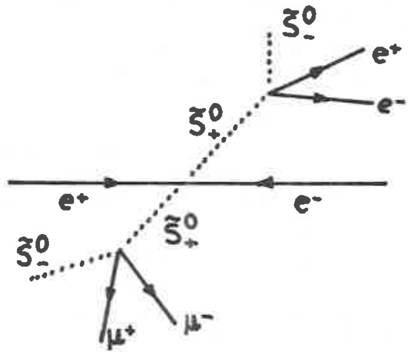


Fig 8. Characteristic production and decay pattern of the lightest new fermion states $\tilde{S}^0_{\pm}$

The other Higgs fermion states, the $\tilde{H}^{\pm}$, are also relatively light, will be produced in e^+e^- annihilation and will have characteristic neutral current couplings which will be able to distinguish them from a new family of heavy lepton or from a Wino (see Table 5).

Particle	g_V	g_A
e^-, μ^-, τ^-	$-\frac{1}{2}+2\sin^2\theta_W \approx 0$	$+\frac{1}{2}$
$\tilde{W}^- + \varepsilon\tilde{H}^-$	$-2-\frac{1}{2}\left(\frac{v^2}{\alpha_2^2 M_0^2}\right)g_2^2 + 2\sin^2\theta_W \approx -\frac{3}{2}$	$\frac{1}{2}\left(\frac{v_1^2-v_2^2}{\alpha_2^2 M_0^2}\right)g_2^2 \approx 0$
$\tilde{H}^- - \varepsilon\tilde{W}^-$	$-1-\left(\frac{v^2}{\alpha_2^2 M_0^2}\right)g_2^2 + 2\sin^2\theta_W \approx -\frac{1}{2}$	$\left(\frac{v_1^2-v_2^2}{\alpha_2^2 M_0^2}\right)g_2^2 \approx 0$

Table 5: Neutral Current Couplings of Charged Fermions

Another characteristic of supersymmetric models is that there are two multiplets of Higgs doublets together with Higgs singlets and consequently the Higgs scalar spectrum is quite rich. Table 4 summarises the expected properties and masses of this sector.

4. Mass scales in SUSY-GUTs

One of the reasons for studying SUSY-GUTs was the hope they would solve the hierarchy problem. In section 2 we found that they indeed allowed scalars to remain massless even in the presence of a superheavy scale M_x. However there remains the question of why the bosonic mass scales are what they are. Ideally one would like to explain the relative sizes of M_{Planck}, M_x, M_{SUSY} and M_W.

4.1 M_{SUSY}?

In section 2 we discussed a mechanism by which a large value for the supersymmetry breaking scale M_{SUSY} would still give a small value for M_W, the difference in scales being given by a high power of the coupling. It is tempting to try to choose $M_{SUSY} \simeq M_{Planck} \simeq M_x$ as the single mass scale in the theory. However this is not possible without including the effects of gravity, for gravitational corrections are suppressed only by powers of (M_{SUSY}/M_{Planck}). Recently the coupling to gravity of N = 1 supersymmetric models has been worked out[28]. At tree level the scalar potential is modified, eq(2.3) becoming

$$\mathcal{L}_{scalar} = e^{\sum_i |\phi|^2 K} \left\{ \left| \sum_i \frac{\partial P}{\partial \phi_i} + K\phi_i^* P \right|^2 - 3K|P|^2 \right\} \tag{4.1}$$

where $K = 1/M^2_{Planck}$. In the limit K = 0 eq(3.1) reduces to the usual N = 1 global superpotential. The form of eq(3.1) gives the canonical form for the scalar kinetic energy term. Relaxing this condition would allow more general arguments for the exponential in eq(3.1). It is clear from eq(3.1) that there are new contributions to scalar quark masses at O(K), whose natural order of magnitude is

$$m_\phi^2 \approx K^2 |P|^2 \tag{4.2}$$

If we insist the two terms in eq(3.1) cancel, giving zero cosmological constant, then $|P|^2$ is related to the supersymmetry breaking scale parameter $M_{SUSY} \equiv \sum_i \left| \frac{\partial P}{\partial \phi_i} + K\phi_i^* P \right| \times e^{\sum_i |\phi_i|^2 K}$. Thus the condition that the Higgs doublet should be less than 1 TeV translates to the condition that $\frac{M_{SUSY}^2}{M_{Planck}} \lesssim 1$ TeV ie $M_{SUSY} \lesssim 10^{10}$ GeV[29]. It is possible to avoid this bound by arranging that this mass squared term be negative, driving a Goldstone mode in which scalars remain massless at tree level because they are (pseudo) Goldstone bosons. However even in this case radiative corrections of the type shown in Fig 9 will generate a scalar mass because the gravitino acquires a mass

$$\tilde{m}_G = \exp(\sum_i \phi_i^2 |P| K \tag{4.3}$$

and spoiling the cancellation between these graphs. Estimates of these graphs[30] give $M_{SUSY} \lesssim 10^{13}$ GeV.

Fig 9. Graphs contributing to scalar masses. G and $G_{3/2}$ are the graviton and gravitino respectively.

Thus it appears that M_{SUSY} cannot be as high as M_{Planck}. Does this mean a new hierarchy problem has arisen? Probably not for there are several possible ways to get a small value for M_{SUSY}.

(i) It should not be forgotten that nonperturbative supersymmetry breakdown has not been ruled out for the models of interest with complex fermion representations. If this happens we would expect $M_{SUSY} \propto e^{-1/g^2} M_{Planck}$. However the lack of definite evidence for or against has caused people to look elsewhere for more tractable methods.

(ii) Since M_{SUSY} is related to a vacuum expectation value of a potential at its minimum it is not surprising if M_{SUSY} should be less than any of the mass scales M' in the potential. In fact it is easy to construct examples[31] with $M_{SUSY} \propto \left(\frac{M'}{M_{Planck}}\right)^n$ giving an arbitrarily small value for M_{SUSY}.

(iii) Another possibility is that large vacuum expectation values ($O(M_{Planck})$) in eq(3.1) may mean that gravitational effects give large suppressions in the effective potential[32]. As a toy example choose P so that V vanishes at the minimum. This may be done for a single chiral field ϕ by setting

$$\frac{\partial P}{\partial \phi} + K\phi^* P = \sqrt{(3K)}\ P \tag{4.4}$$

This cannot be solved for an analytic function P but instead use the equation

$$\frac{\partial P}{\partial \phi} + K\phi P = \sqrt{(3K)}\ P \tag{4.5}$$

This gives

$$P(\phi) = M^3 \exp\left[-\sqrt{3K}\,\phi - K\phi^2/2\right] \tag{4.6}$$

which, at the minimum has $\phi_I = 0$, and $V = 0$ for any value of ϕ. In this case the supersymmetry breaking term is

$$\begin{aligned} e^{|\phi|^2 K}\left|\frac{\partial P}{\partial \phi} + K\phi^* P\right| &= 3e^{-K\phi^2 - 2\mathrm{Re}(\sqrt{3K}\,\phi + K\phi^2/2)} \\ &= 3\, e^{-2\sqrt{3K}\,\phi_R}\, M^3 (M/M_{Planck})^2 \\ &\equiv M^3_{SUSY} \end{aligned} \tag{4.7}$$

even with $M = M_{Planck}$, for large ϕ_R the value of M_{SUSY} may be very small.

$\underline{M_W?}$

The simplest solution for M_W is $M_W \simeq M_{SUSY}$. However other possibilities have been discussed. The first[19], discussed in section 2, is that $M_W = f(\frac{\lambda}{4\pi})\, M_{SUSY}$, and that a large number of loops combine to give a small ratio for M_W/M_{SUSY}. An alternative scheme,[33] the geometric hierarchy scheme, generates M_W through radiative corrections too, but in this case the equivalent graphs to Fig (4) involve massive (gauge non singlet) states with mass $\approx M_x$. Then it is easy to see on dimensional grounds $M_W = \frac{M^2_{SUSY}}{M_x}$. Unfortunately in this class of model the sliding singlet trick of section 2 fails due to radiative corrections and models constructed so far have to adjust parameters to arrange the necessary splitting in the Higgs multiplets.

$\underline{M_x?}$

The simplest idea is to have $M_x = O(\alpha M_{Planck})$, or $O(\frac{1}{\alpha} M_{SUSY})$[19]. Both of these possibilities look feasible and reasonable. Another suggestion due to Witten [34] is that M_x (and also M_{Planck}!) may be related to $M_{SUSY} \approx M_W$ by radiative corrections generating an inverse heirarchy

$$M_x = e^{c/g^2} M_{SUSY}$$

where c is a coefficient determined by calculation of (perturbative) radiative corrections. Unfortunately the original idea does not appear to work for with realistic values for c and g a hierarchy of only of order $e^{c/g^2} \approx 10$ is natural[35]. However it may be possible to resurrect the idea of a low fundamental scale for M_{SUSY} in a different guise[36], although in these schemes there is no understanding of M_{Planck} relative to M_{SUSY}.

5. Conclusions

Global supersymmetry solves the hierarchy problem in a way which requires many new states, However it is expected these states should be heavy and would not yet have been observed, although not so heavy that they will elude detection soon if they exist. Viable models have been constructed and there are plausible explanations for the relative magnitude of the bosonic mass scales. The inclusion of gravity may render the models more elegant in the symmetry breaking sector and allows for the cancellation of the cosmological constant; although its cancellation is unnatural and understanding this problem remains a fundamental stumbling block. However it is an exciting fact that we may, for the first time in particle physics have phenomenological reasons for including gravity and for moving towards a truly unified theory.

Of course the whole construction is pure theory at present; encouragingly these models are testable and experiment will decide.

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SEARCH FOR NEW PARTICLES IN e^+e^- ANNIHILATION

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ABSTRACT

Extensive searches for new particles have been performed at e^+e^- storage rings. No new leptons or quark flavours have been found up to centre of mass energies of 36.75 GeV. Also the search for new scalars has been negative. In particular, supersymmetric leptons must have masses larger than 16 GeV and charged technipions and "standard" axions can be excluded.

1. Introduction

As the title indicates, this is a report on the discoveries which have not (yet) been made at e^+e^- storage rings. Motivations to search for new particles come from many sources. Even in the "standard" GWS-GIM-model[1,2] several particles are still missing. Numerous alternatives and extension to this model call for a large variety of other or additional particles:

- Fermions

  Besides the top quark and the τ neutrino, for which direct evidence is still missing, further families of quarks and leptons (sequential leptons) may still exist. More exotic leptons (stable, excited, or neutral) are predicted in specific models and fractional charged quarks (and monopoles) are under heated discussion again. Moreover, supersymmetric theories predict spin 0 partners of quarks and leptons (squarks, sleptons).

- Vector Bosons

  The mediators of the weak interaction, $W^\pm$ and Z^0, are not the only particles on the shopping list for new vector bosons. Again, in particular supersymmetry, predicts fermion partners. Some of them, photinos and gravitinos, will be discussed.

- Scalar Bosons

  Whereas the standard model asks for one single physical scalar (Higgs) boson H^0, also charged Higgses, axions and technipions are predicted in other models.

Many of the new particles suggested above are directly or indirectly accessible in e^+e^- reactions. In e^+e^- annihilation the direct production of charged (e_q), pointlike particles is given by ($\beta = v/c$, s = centre of mass energy squared, $\sigma_{QED} = 4\pi\alpha^2/3s$

$$\sigma = \sigma_{QED} \cdot e_q^2 \frac{\beta(3-\beta^2)}{2} \qquad \text{spin } 1/2 \qquad (1)$$

$$\sigma = \sigma_{QED} \cdot e_q^2 \; 1/4 \; \beta^3 \qquad \text{spin } 0 \qquad (2)$$

i.e. the production of scalars is 4 times smaller and suppressed by an additional threshold factor with respect to the production of spin 1/2 fermions (see Fig. 1). We will start with the discussion of fermions and then proceed to scalar particles.

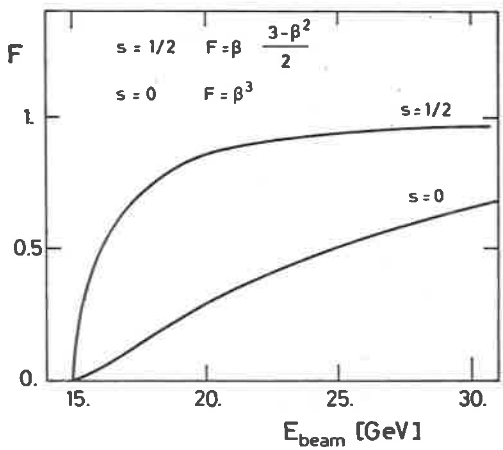


Fig. 1

Threshold behaviour of pointlike scalars and spin 1/2 fermions in e^+e^- annihilation[3c)]

This report will include searches carried out by the MAC, MARK II, PEP12 and PEP14 experiments at PEP, the CELLO, JADE, MARK J, PLUTO, and TASSO experiments at PETRA, the CLEO and CUSB experiments at CESR, and the Crystal Ball and MARK II experiments at SPEAR. Similar reports have been published previously[3)].

2. Search for New Fermions

2.1 The "Old" Sequential Heavy Lepton τ

Before we talk about new lepton families, let us first have a quick look at the last one[4)].

The τ lepton has so far not shown any deviation from the straight-forward predictions for a sequential heavy lepton. Latest results from high energies show pointlike behaviour

$$\sigma_{\tau\tau} = \sigma_{QED}\ (1.03 \pm 0.05 \pm 0.07) \qquad \text{CELLO}^{5)}$$

$$\sigma_{\tau\tau} = \sigma_{QED}\ (0.97 \pm 0.05 \pm 0.06) \qquad \text{MARK II}^{6)}$$

and a charge asymmetry in the angular distribution

$$A = -8.0 \pm 2.3 \qquad \text{PETRA average}^{7)}$$

in good agreement with the prediction of the standard model

$$A = -9.3 \pm 0.2 \qquad \text{Theory}$$

As reported in Paris, several groups have now been able to measure the lifetime of the τ lepton. In particular, the beautiful measurement of the MARK II group yields

$$\tau_\tau = (3.31 \pm 0.57 \pm 0.6) \cdot 10^{-13}\,\text{s} \qquad \text{MARK II}^{8)}$$

in agreement with the value expected from e-μ-τ universality

$$\tau_\tau = (2.8 \pm 0.2) \cdot 10^{-13}\ \text{s} \qquad \text{Theory}$$

Apart from the value this measurement has in itself, it bears heavily on the question of whether or not the τ neutrino exists[4b].

Already before the new measurements on the τ lifetime were performed in 1982, it was known that an extra neutrino was involved in τ decays and that this neutrino was different from ν_μ, $\bar{\nu}_\mu$ and $\bar{\nu}_e$[4c]. However, the one case $\nu_\tau = \nu_e$ could not be excluded. Now, arguments similar to those which already led to the exclusion of $\nu_\tau = \nu_\mu$ can also be applied for ν_e.

Let us consider the coupling strength $\varepsilon(\tau - \nu_e)$ of a hypothetical τ-ν_e vertex. If $\varepsilon(\tau-\nu_e) \neq 0$, this would lead to a production of τ's in ν_e interactions. In a beam dump experiment, the BEBC group has determined an upper limit[9])

$$\varepsilon(\tau-\nu_e) \quad < \quad 0.35 \text{ of universal coupling strength} \quad (90\% \text{ C.L.})$$

from the comparison of charged/neutral current ν_e interactions. On the other hand, the MARK II lifetime value constrains the value from below

$$\varepsilon(\tau-\nu_e) \quad > \quad 0.75 \text{ of universal coupling strength} \quad (90\% \text{ C.L.}) \quad .$$

This contradiction indicates that ν_τ cannot be identical to ν_e. Thus, by exclusion of all alternatives, we have indirect experimental evidence for a new neutrino ν_τ in τ decays.

2.2 The Next Sequential Heavy Lepton L

The production and decay of a hypothetical new heavy lepton L is shown in Figs. 2a and b. As an example the predicted branching ratios for a mass of 16 GeV are given in Table 1 [10]). The signatures which

Table 1: Branching ratios for a hypothetical sequential heavy lepton L^- of mass M_L = 16 GeV [10])

$$BR(L \to \ell\nu\nu) \quad \simeq \quad 10.7\% \quad (\ell = e,\mu,\tau)$$

$$BR(L \to \nu u\bar{d}) \quad \simeq \quad BR(L \to \nu c\bar{s}) \quad \simeq \quad 32\%$$

$$BR(L \to \nu u\bar{s}) \quad \simeq \quad BR(L \to \nu c\bar{d}) \quad \simeq \quad 2\%$$

can be used are similar to those for τ detection (Figs. 2c,d,e). Signature (c) - a μ acoplanar with a jet of invariant mass larger than the τ mass - has been used by most groups to search for a new lepton. Signature (d) - two acoplanar jets of invariant mass larger than the τ mass - has been used by JADE, the signature (e) -acoplanar e μ pair -

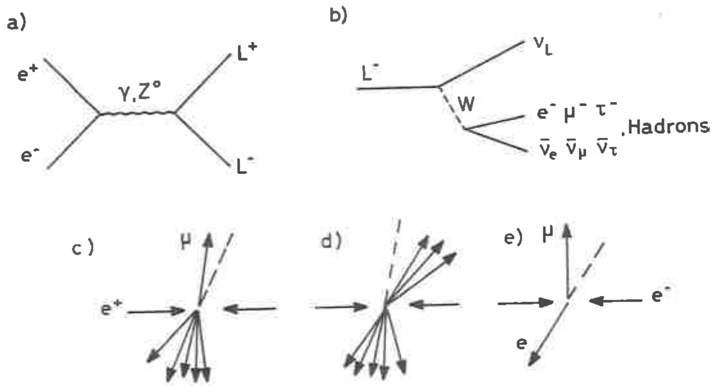


Fig. 2: Production (a) and decay (b) of sequential heavy leptons and different experimental signatures (c,d,e).

by the MARK II group. The results are summarized in Table 2.

Table 2: Experimental lower limits (95% C.L.) on the mass M_L of a new sequential heavy lepton L

experiment	mass limit	Ref.
CELLO	$M_L > 16.3$ GeV	11
JADE	$M_L > 18.1$ GeV	3a
MAC	$M_L > 14.5$ GeV	12,13
MARK J	$M_L > 16.0$ GeV	14
MARK II	$M_L > 13.8$ GeV	12
PLUTO	$M_L > 14.5$ GeV	15
TASSO	$M_L > 15.5$ GeV	16

The cases of ortho ($\nu_L = \nu_e,\nu_\mu,\nu_\tau$) or para ($\nu_L = \bar{\nu}_e,\bar{\nu}_\mu,\bar{\nu}_\tau$) leptons have not been evaluated in detail, but would lead to similar limits.

2.3 Stable Charged and Heavy Neutral Leptons

New lepton families may exist, in which the neutral partner is heavier than the charged one[17]:

$$\begin{pmatrix} L^o \\ L^- \end{pmatrix} \quad \text{with} \quad M_{L^o} > M_{L^-} \tag{3}$$

This would lead to lepton signatures much different from the sequential case discussed above.

Stable Charged Leptons

If the new lepton family has its own conserved quantum number, the charged partners are stable and behave like additional "μ pairs" with low momenta

$$e^+e^- \rightarrow L^+L^- \qquad L^{\pm} \text{ stable} \tag{4}$$

The MARK J and JADE groups have searched for such signals. Fig. 3

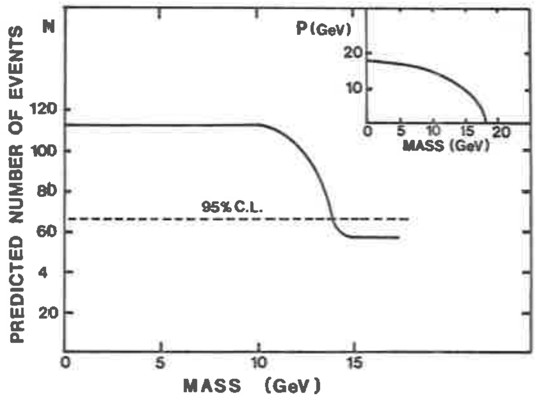


Fig. 3

Experimental limit on the production of stable charged leptons (MARK J)[19] shows the upper limit deduced from μ pair production in MARK J compared to the expected number of events. The JADE group has looked for excessive stable charged particle pair production (see below). The experimental limits of the two groups are (95% C.L.)

$$M_{L^\pm} > 12 \text{ GeV} \qquad \text{JADE}^{18)}$$

(for charges $\geq 2/3$)

$$M_{L^\pm} > 14 \text{ GeV} \qquad \text{MARK J}^{19)}$$

It should be kept in mind that these limits only hold for $M_{L^0} > M_{L^-}$ [4a]. If both masses would be similar with the charged lepton slightly heavier, L^- decays of the type

$$L^- \rightarrow L^0 + e^- + \bar{\nu}_e \qquad \text{etc.} \tag{5}$$

would be very difficult to detect. Because of the low electron energy and large missing mass they would probably be indistinguishible from 2γ events.

Neutral Heavy Leptons

For the heavy neutral lepton in a hypothetical new family the following phenomenlogical cases have to be considered ($\ell = e,\mu,\tau$) (Ref. 4a)

- New conserved lepton number, charged (L^-) or neutral (ν_L) partner

$$L^0 \rightarrow \nu_L + \ell^+ + \ell^- \quad \text{or} \quad \nu_L + \text{hadrons}^0 \tag{6a}$$

$$L^0 \rightarrow L^- + \ell^+ + \nu_\ell \quad \text{or} \quad L^- + \text{hadrons}^+ \tag{6b}$$

- No new conserved lepton number, "old" lepton partners (ℓ^-,ν_ℓ)

$$L^0 \rightarrow \nu_\ell + \ell^+ + \ell^- \quad \text{or} \quad \nu_\ell + \text{hadrons}^0 \tag{6c}$$

$$L^o \rightarrow \ell^- + \nu_\ell + \ell^+ \quad \text{or} \quad \overline{\ell^- + \text{hadrons}^+} \qquad (6d)$$

Since the production of L^o can only occur through weak interactions (Fig. 4), cross sections will be low. The cross section of pair production via Z^o (Fig. 4a) is estimated to be of the order of[20]

$$\sigma(e^+e^- \rightarrow L^o\overline{L^o}) \simeq 0.016\ \sigma_{QED} \qquad (\sqrt{s} = 40\ \text{GeV})$$

at PEP/PETRA energies, corresponding to about 60 events/100 pb^{-1}. Note that the cross section ratio increases like $\sigma/\sigma_{QED} \sim s^2$.

For the particular case of ℓ being an electron in (6d), the heavy neutral electron E^o can also be produced via $W^\pm$ exchange (Fig. 4b) with a substantially higher cross section

$$\sigma(e^+e^- \rightarrow \bar{E}^o\ \nu_e) \approx 0.1\ \sigma_{QED}$$

at PEP/PETRA energies. This possibility has been studied by the JADE group[3a,c] for the case, where $\bar{E}^o$ decays into an electron and hadrons. The experimental signature and the expected cross section are given in Fig. 5. No event has been observed, which corresponds to the following 95% C.L. lower limits on the mass of E^o in case of V+A or V-A interaction[21].

$$M(E^o) > 24.5\ \text{GeV} \quad (V{+}A) \qquad \text{JADE}$$
$$M(E^o) > 22.5\ \text{GeV} \quad (V{-}A) \qquad \text{JADE}$$

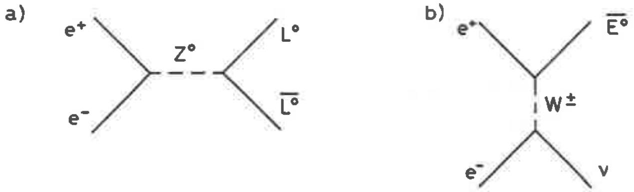
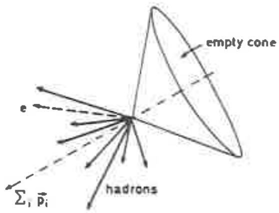


Fig. 4
Production of neutral leptons in e^+e^- annihilation



Event signature for
neutral heavy electron

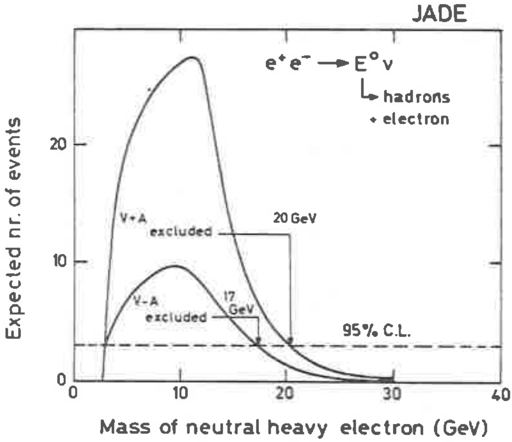


Fig. 5

Experimental limits on the production of neutral electrons (JADE) (Ref. 3a)

2.4 Excited Charged Leptons

If leptons were composite particles, one would expect to observe finite structures and excited states. As far as structures are concerned, no deviation from pointlike behaviour of e, μ, and τ has been observed up to 0 $(100\ \text{GeV})^{-1}$.

Direct Production of μ^*

There are two reactions through which excited leptons (heavier particles with the same quantum numbers as the corresponding leptons) could be produced directly in e^+e^- annihilation (Fig. 6)

$$e^+e^- \quad \to \quad \mu^*\mu^* \tag{7}$$

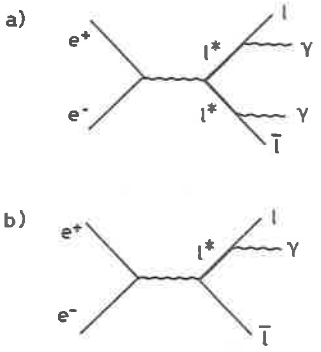


Fig. 6
Production of excited charged leptons

$$e^+e^- \rightarrow \mu^*\mu \tag{8}$$

Whereas the cross section for (7) in the case of a pointlike μ^* with a mass less than the beam energy would be given by (1), reaction (8) would require an unconventional current of the type[22] (coupling strength λ)

$$e\lambda\ \psi_{\mu^*}\ \sigma_{\beta\alpha}\ \psi_\mu \tag{9}$$

The cross section then reads (M_{μ^*} = mass of μ^*)

$$\sigma = \frac{8\pi\alpha^2}{3}\lambda^2\ (1-M^2_{\mu^*}/s)^2\ (1+2\ M^2_{\mu^*}/s) \tag{10}$$

If μ^* decays rapidly

$$\mu^* \rightarrow \mu + \gamma$$

one would observe a signal

$$e^+e^- \rightarrow \mu^+\mu^-\gamma\ (\gamma) \tag{11}$$

which has to be separated from the radiative QED background. Mass limits on M_{μ^*} are of course restricted to less than the beam energy in reaction (7), whereas higher values ($\lesssim \sqrt{s}$) can be reached in reaction (8).

Searches for excited muons μ^* have been performed by the CELLO, JADE, MARK J, and MAC collaborations. For masses below beam energy, mass limits can be deduced from the observed limits on excess events of type (11) compared to the expected cross section (1). The experimental limits are summarized in Table 3.

Table 3: Experimental limits on μ^* masses from reaction $e^+e^- \to \mu^*\mu^*$

Experiment	Mass Limit (GeV) (95% C.L.)	Ref.
CELLO	> 16.9	11
JADE	> 10	3a
MARK J	> 10	19
MAC	> 14	12,13

In the case of reaction (8), an excited muon would show up as a peak in the invariant mass distribution of the $\mu\gamma$ system. As an example, the data of the MAC group[13)] are shown in Fig. 7. No such signal has been observed. From a comparison with the QED expectation, upper limits on the observed cross section for reaction (8) can be obtained. By means of the expected cross section (10) this can be transformed into limits on the coupling constant λ as a function of M_{μ^*}. Figs. 7 and 8 show two recent results from the MAC[13)] and MARK J[19)] groups. Similar limits were obtained by the CELLO and JADE collaborations[3b)]. Fig. 8 also shows the constraints on λ which can be obtained from the anomalous magnetic moment of the muon[23)].

At SPEAR the MARK II collaboration has also searched for signals of excited electrons and muons. They can put stingent limits on the production of e^* and μ^* up to masses of about 3 GeV. The results are shown in Fig. 9 [24)].

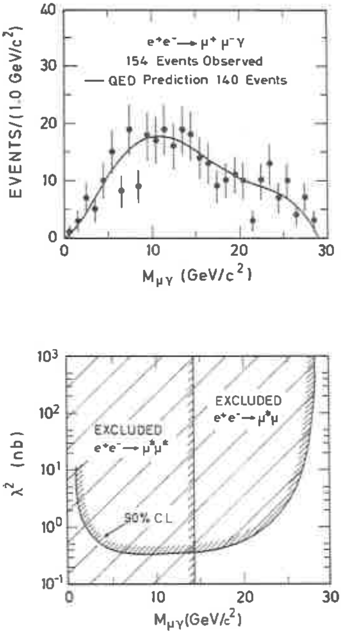


Fig. 7

Experimental limits on the production of excited muons μ^* (MAC)[13)]

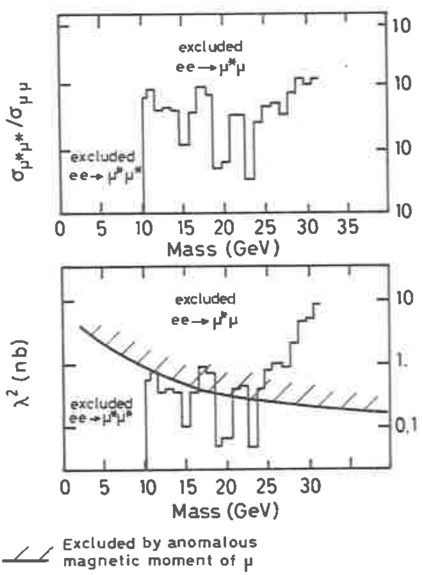


Fig. 8

Experimental limits on the production of excited muons μ^* (MARK J)[19)] together with limits from the anomalous magnetic moment of the muon[23)]

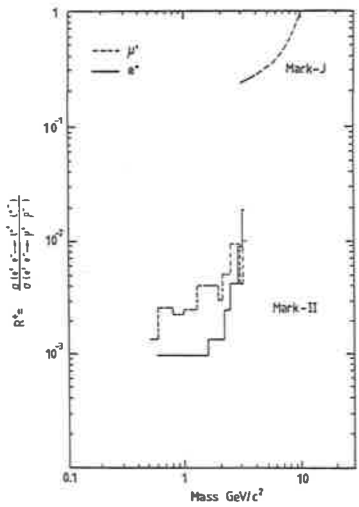


Fig. 9

Experimental limits on the production of excited muons μ^* and e^* (MARK II)[24]

Virtual Excited Electrons

If currents of the type (9) would exist, they would interfere with the electron exchange diagram of the reaction

$$e^+e^- \rightarrow \gamma\gamma \tag{12}$$

in the way indicated in Fig. 10. This modifies the cross section for (12)

$$\sigma = \sigma_{QED} \left(1 + (s^2/\Lambda^4) \sin^2\theta\right) \tag{13}$$

where $\Lambda = M_{e*} \sqrt{\lambda'}$ and λ' is the relative coupling strength of current (9).

Many groups have looked for deviations from QED in reaction (12) for which radiative corrections have been calculated up to $O(\alpha^3)$ and which is unaffected in lowest order by weak interactions. Figs. 11 and 12 show two recent results of the MARK II and CELLO collaborations. Both agree well with the QED predictions. According to (13) this can

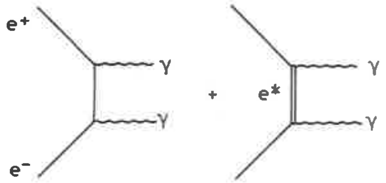


Fig. 10

Possible contributions to $e^+e^- \to \gamma\gamma$

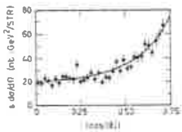


Fig. 11

Invariant differential cross section for the reaction $e^+e^- \to \gamma\gamma$. Data are compared to the QED prediction to O (α^3) (MARK II)[12)]

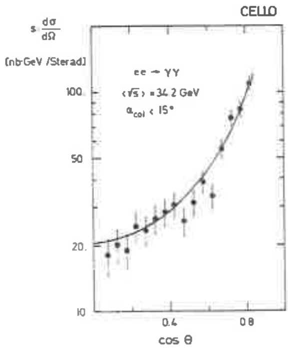


Fig. 12

Invariant differential cross section for the reaction $e^+e^- \to \gamma\gamma$. Data are corrected for radiative effects to O (α^3) and compared to lowest order QED (CELLO)[25)]

be transformed into a limit on $M_{e^*} \cdot \sqrt{\lambda'}$. Table 4 summarizes the results obtained by different groups at PEP and PETRA:

Table 4: Experimental Limits on $M_{e*} \cdot \sqrt{\lambda'}$ (95% C.L.)

Experiment	$M_{e*} \cdot \sqrt{\lambda'}$	Ref.
CELLO	> 59 GeV	25
JADE	> 47 GeV	26
MARK J	> 58 GeV	19, 27
MARK II	> 50 GeV	12
PLUTO	> 46 GeV	15,28
TASSO	> 34 GeV	29

2.5 Search for New Quark Flavours

One of the main objectives of the high energy e^+e^- machines is the search for new quark flavours. I will only very briefly sketch the extensive experimental program which has been carried out at highest PETRA energies to search for thresholds or resonances in the hadronic cross section[3a,b,c].

Search for Thresholds

Variations in the total cross section, event shapes and inclusive lepton rates would be indicative of new quark thresholds. No such changes have been observed up to center of mass energies of $\sqrt{s}$ = 36.75 GeV. This excludes new quark flavours with charge 2/3 (from total cross section) or 1/3 (from event shape analysis).

Search for Narrow Resonances

The highest energy region 33 GeV < $\sqrt{s}$ < 36.75 GeV has been scanned in steps of 20 MeV to look for narrow resonances. No resonances have been found. Quantitatively, an upper limit of

$$B_h \cdot \Gamma_{ee} \quad < \quad 0.65 \text{ keV} \qquad (90\% \text{ C.L.})$$

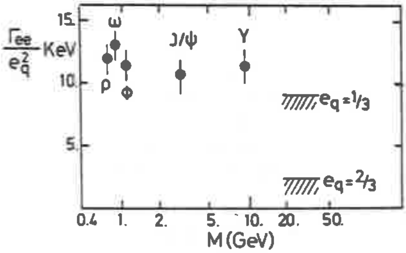


Fig. 13

The reduced leptonic width Γ_{ee}/q_e^2 of the known vector mesons compared to the experimental limit at highest PETRA energies[3c)]

can be deduced from the combined PETRA data where B_h is the hadronic branching ratio and Γ_{ee} the leptonic width of the resonance. If we assume B_h = 0.7 this gives an upper limit of

$$\Gamma_{ee} \quad < \quad 0.93 \text{ keV} \qquad (90\% \text{ C.L.})$$

In Fig. 13, this limit is compared to the width of known narrow vector mesons in terms of Γ_{ee}/e_q^2, e_q being the quark charge. If we assume Γ_{ee}/e_q^2 to be about constant, as previous data suggest, we can safely exclude a $t\bar{t}$ bound state with e_t = 2/3.

2.6 Search for Free Fractionally Charged Quarks

Searches for fractional charges have been performed in several experiments at PEP and PETRA. They all rely on a measurement in the specific ionisation

$$dE/dx \sim e_q^2 \cdot f(\beta) \qquad (14)$$

The JADE group measures dE/dx and the reduced momentum p/e_q in a "jet" chamber[18)]. The PEP14 collaboration has no magnetic field and uses dE/dx and TOF measurements in scintillation counters[30,31)]. In addition, tracks are defined in proportional chambers. The MARK II group combines the dE/dx and TOF information from scintillation counters with a p/e_p measurement in the drift chambers[32)].

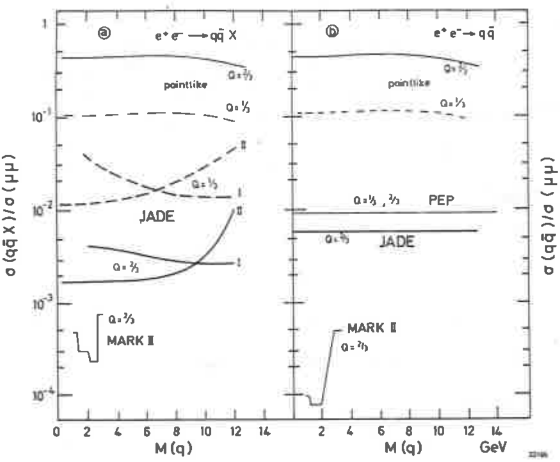


Fig. 14

Experimental upper limits on the cross section for inclusive and exclusive quark production as a function of the quark mass (JADE, MARK II, PEP14)[3a)]

(a) limits on inclusive production
(b) limits on exclusive production

All three groups have searched for the exclusive production of quarks

$$e^+e^- \rightarrow q\bar{q} \tag{15}$$

JADE and MARK II have also determined limits on the inclusive production

$$e^+e^- \rightarrow q\bar{q}X \tag{16}$$

No signals have been seen. The upper limits relative to the pointlike cross section $\sigma_{\mu\mu}$ as a function of the quark mass are shown in Fig. 14 (Ref. 3a).

2.7 Search for Monopoles

Monopoles (which, strictly speaking, do not fit in here, but are somewhat related to the question of fractional charges) were originally proposed by Dirac to symmetrize the Maxwell equations and to quantize the elctric charge[32)]. These Dirac Monopoles have a magnetic charge g which is related to the electric charge e by

$$g/e = n/2\alpha \sim 68.5\ n \ ; \quad n = 1,2,... \tag{17}$$

The mass of these monopoles M_M is predicted to be (M_p = proton mass)

$$M_M \sim 2.56\ n^2 \cdot M_p \tag{18}$$

Later, T'Hooft and Polyakov[34)] pointed out that monopoles appear as soliton solutions of the classical field equations and must exist in unified gauge theories. Their mass is much higher than that of Dirac monopoles, $M_M \sim M_G/\alpha$, where M_G is a typical mass of the gauge, e.g. 100 GeV for electroweak theories or 10^{14} GeV for unified theories. Clearly, this kind of monopole is inaccessible in e^+e^- interactions and our search is restricted to Dirac monopoles (18).

To search for monopoles in e^+e^- reactions, we make use of the high specific ionisation which is expected from (17)

$$dE/dx \sim (g/e)^2\ (dE/dx)_{min} \quad \text{for } \beta \sim 1 \tag{19}$$

This is about 4700 times the ionisation of a minimum ionising particle for $n = 1$, and rises quadratically with n.

The experimental procedure to look for these highly ionising particles is quite simple and has been applied at PEP and PETRA. Plastic foils are wrapped around the vacuum pipe near an interaction point and later analyzed for traces of heavily ionising particles. Searches in the PEP16 experiment have yielded an upper limit for monopole production[30)]

$$\sigma(e^+e^- \to \text{monopoles} + X) < 0.007 \cdot \sigma_{point}$$

for $M_M < 14$ GeV and n=1. For $n > 2$ these simple experiments are insensitive because the particles are absorbed in the beam pipe. Therefore, foils have been put into the vacuum of the beam pipe in the MARK J detector at PETRA. No results have been reported yet.

3. Search for New Scalars

Motivation to look for new scalar particles has recently come from two sources:

a) To explain symmetry breaking, models of increasing complexity have been proposed which require one or several (many) new scalar particles[35).

b) Supersymmetry associates a new scalar particle to each fermion[36,3d).

As already mentioned above, the search for scalars in e^+e^- is rendered more difficult compared to fermion searches by a factor 1/4 in the asymptotic cross section and a slowly rising threshold function $\sim \beta^3$ (see Fig. 1).

3.1 Search for Supersymmetric Particles

In supersymmetry, each fermion with spin J has one boson counterpart with spin J ± 1/2 and vice versa. Thus, to each particle, there exists a supersymmetric partner ("sparticle")[3d). Table 5 contains a few examples which will be relevant to the further discussion.

Table 5: Supersymmetric Partners

Particle	J	Sparticle	J ± 1/2
lepton l	1/2	slepton s_ℓ t_ℓ	0
photon γ	1	photino λ_γ	1/2
graviton	2	gravitino λ_g	3/2

Some properties of these hypothetical particles are illustrated in Fig. 15. The sleptons s_ℓ, t_ℓ (one for each helicity state of the lep-

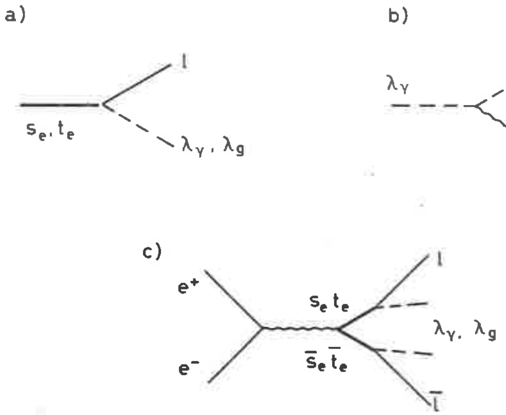


Fig. 15

Production and decay of supersymmetric particles

a) Supersymmetric lepton decay,

b) Model-dependent photino decay,

c) production and decay of supersymmetric leptons in e^+e^- annihilation

ton 1) decay rapidly into the related leptons and a photino λ_γ or gravitino λ_g (which could also be a goldstino if gravity were not included in the theory). The photino may be stable or may decay into a photon and a gravitino (Fig. 15b). This process is strongly model-dependent and will be discussed later.

3.1.1 Search for Direct Production of Supersymmetric Leptons

Pointlike sleptons are pair produced with the cross section (2) in e^+e^- annihilation and decay rapidly according to Fig. 15. Thus the signature is (l = e, μ, τ)

$$e^+e^- \rightarrow l^+l^- \quad \text{(acollinear)} \tag{20}$$

Although this is a simple and clean signature in principle, background, in particular from radiative QED processes,has to be eliminated carefully. Many groups at PEP/PETRA have searched for sleptons. As examples, the results of the CELLO group on supersymmetric electron (t_ℓ, s_ℓ) and muon (t_μ, s_μ) searches and of the MARK J collaboration on a supersymmetric tau (t_τ, s_τ) search are shown in Figs. 16 and 17. It should be noted that the additional graph of λ_γ exchange contributes to the supersymmetric production and increases the expected rate considerably (Fig. 16a).

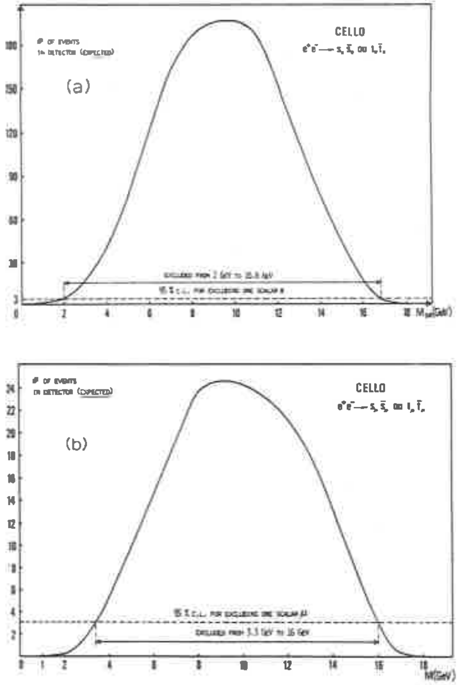


Fig. 16

Experimental limits on supersymmetric electrons and muons (CELLO)[37)]

(a) suypersymmetric electrons

(b) supersymmetric muons

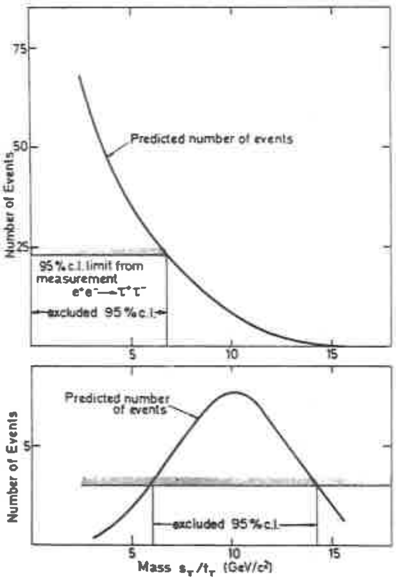


Fig. 17

Experimental limits on supersymmetric taus (MARK J)[39)]

This improves the experimental limits which can be obtained. The signature used by the MARK J group to search for supersymmetric taus is similar to (2c) in the case of the sequential lepton search and the hyperpion search discussed below. The results of the experimental searches at PEP/PETRA are summarized in Table 6. The value derived from the anomalous magnetic moment of the muon is also included.

Table 6: Experimental Lower Limits on the Masses of Supersymmetric Leptons in GeV (95% C.L.)

	$s_\ell t_\ell$	$s_\mu t_\mu$	$s_\tau t_\tau$ *	Ref.
CELLO	16.8	16	15.3	37
JADE	16		14	38
MARK J		15	14	14,39
MARK II			9.9	40
TASSO	16.6	16.4		41
MAC		13	13	13
g-2		13		42

* mostly from hyperpion search (c.f. Fig. 24)

The experiments quoted usually also give upper limits for the lower masses. These values have not been included in the table since the low mass region is already excluded by precision measurements on electron and muon pair production, in a way similar to that indicated in Fig. 17a for the case of τ's.

To improve the supersymmetric mass limit, a study of the single production shown in Fig. 18 was suggested[43)]. The electron would stay invisible in the beam pipe and a single supersymmetric electron would emerge under large angle and decay into an electron and a photino.

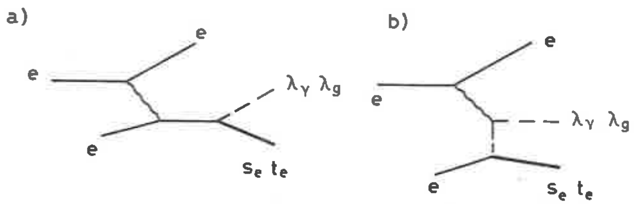


Fig. 18: Single production of a supersymmetric electron in e^+e^--collisions

Thus the experiment would have to look for single electrons at large angle. The authors claim that the supersymmetric electron mass limit could be pushed to about 1.25 times the beam energy.

3.1.2 Search for Virtual Supersymmetric Electrons and Search for Photinos

One may try to study the supersymmetric analogue to reaction (12) as given in Fig. 19a

$$e^+e^- \rightarrow \lambda_\gamma\bar{\lambda}_\gamma \tag{21}$$

We have to distinguish the following two alternatives:

λ_γ, λ_g are stable

As λ_γ and λ_g are supposed to interact only very weakly, they will be invisible if they do not decay in the detector. Therefore, the only way to study this process experimentally would be through initial state radiation (Fig. 19b) as suggested by Fayet[44]. The expected rates of

$$\sigma(\lambda_\gamma\bar{\lambda}_\gamma\gamma) \sim 0.3 \text{ pb} \quad \text{for} \quad M_{s_e} = 20 \text{ GeV}$$
$$\sim 0.1 \text{ pb} \quad \text{for} \quad M_{s_e} = 40 \text{ GeV}$$

would lead to a signal of $10 \div 30$ events for 100 pb^{-1}. In view of the

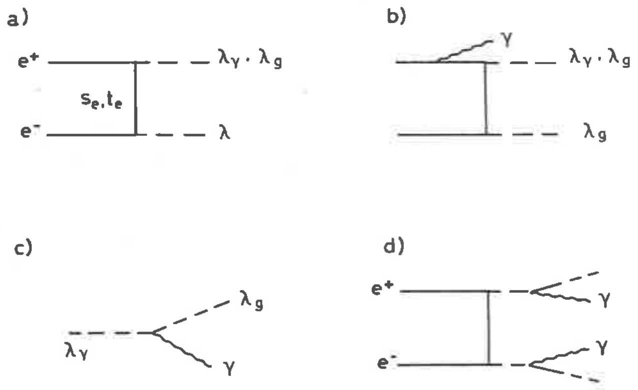


Fig. 19: Production and decay of photinos in e^+e^- annihilation

experimental problems of triggering on a single photon of 2 to 3 GeV and fighting the background from radiative QED processes, this is certainly a very low signal. It requires a very good trigger and complete solid angle coverage for electromagnetic showers.

λ_γ is unstable

If the photino is massive, with a mass larger than the gravitino mass, it will decay according to Fig. 19c with a lifetime[45)]

$$\tau_{\lambda_\gamma} = \frac{8\pi d^2}{M_{\lambda_\gamma}{}^5} \tag{22}$$

where M_{λ_γ} is the photino mass and d is a parameter characterizing the symmetry breaking, which is unknown. If τ_{λ_γ} is such that λ_γ decays inside the detector, the process (21) would lead to two photons in the final state (Fig. 19d) with the experimental signature

$$e^+e^- \rightarrow \gamma\gamma + \text{missing energy} \tag{23}$$

The CELLO group has searched for events of this type[46)]. Fig. 20

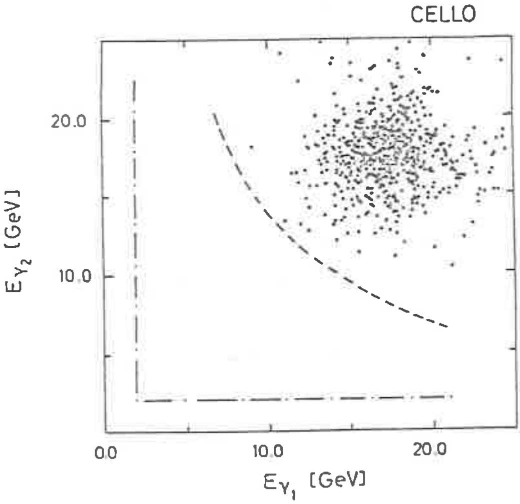


Fig. 20

Correlation of the two photon energies in the reaction $e^+e^- \to \gamma\gamma$ (CELLO) (Ref. 46)

shows the correlation of the energies of the two photons in reaction (23). All events pass the missing energy cut indicated in the figure. Fig. 21 shows the constraints which can be derived on the mass of the photino and the scale parameter d assuming m_{s_e} = 40 GeV, together with limits from cosmological arguments[56)] and search for the decay $J/\psi \to$ nothing [42)].

3.2 Search for Higgs Particles

3.2.1 Search for the Neutral Higgs Particle H^o

In the Standard Model with one Higgs doublet only, one physical particle, the neutral Higgs H^o, emerges after symmetry breaking.

Since H^o couples preferentially to large mass fermions, the best way to look for it in e^+e^- annihilation would be in the decay of heavy onium O [47)] (Fig. 22a). The relative width is estimated to be (m_o, m_H = onium and Higgs mass)

$$\frac{\Gamma(O \to H^o\gamma)}{\Gamma(O \to e^+e^-)} \simeq \frac{G_F\, m_o^2}{4\sqrt{2}\,\pi\,\alpha}\left(1 - \frac{m_H^2}{m_o^2}\right) \qquad (24)$$

which leads to the following branching ratios[48]

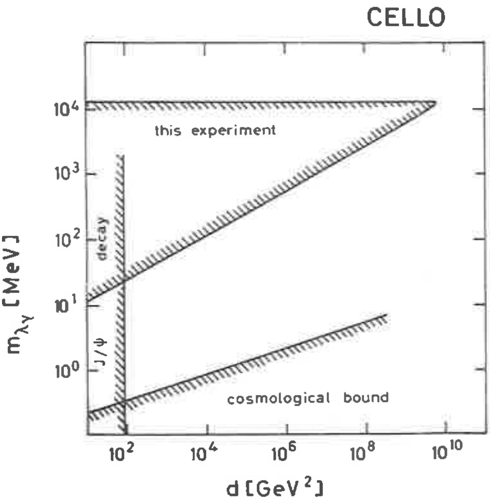


Fig. 21

Experimental bounds on the mass of the photino as a function of the scale parameter d of the supersymmetry breaking (CELLO)[46]

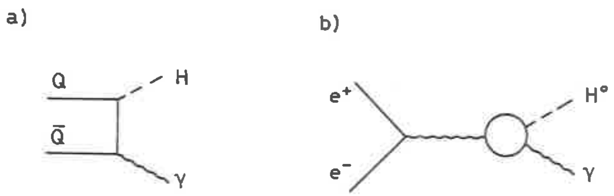


Fig. 22

Production of neutral Higgs particles in e^+e^- reactions

a) Decay of heavy onium

b) Production in the continuum

$$BR(J/\psi \to H^o\gamma) \simeq 1/2 \cdot 10^{-4}$$

$$BR(\Upsilon \to H^o\gamma) \simeq 2 \cdot 10^{-4}$$

$$BR(t\bar{t}(^3S_1) \to H^o) \simeq 10^{-2} \qquad (M_{t\bar{t}} = 40 \text{ GeV})$$

An upper limit on the J/ψ branching ratio of 10^{-3} is quoted by the Crystal Ball group[48].

Even if the mass of H^o were sufficently low, it would be very difficult to detect H^o on the J/ψ or Υ because of the low branching ratios. Toponium would certainly be an ideal place to look for H^o.

Another possible way of producing H^o in e^+e^- annihilation is shown in Fig. 22b. This production in the continuum is however highly suppressed since it involves heavy quark and vector boson loops. The cross section of[49)]

$$\sigma(e^+e^- \rightarrow H^o\gamma) \simeq 10^{-6}\ \sigma_{QED}$$

is hopelessly small. Fig. 23 shows the radiative background

$$e^+e^- \rightarrow \gamma + q\bar{q} \rightarrow \gamma + \text{hadrons} \qquad (25)$$

From the absence of any structure or deviation from the expectation for (25) an upper limit of $\sigma(e^+e^- \rightarrow H^o\gamma) \lesssim 10^{-2}\ \sigma_{QED}$ can be deduced (Ref. 3c,50).

3.2.2 Search for Charged Higgs Particles and Hyperpions

There are several reasons why the Higgs sector may be more complicated than proposed in the Standard Model e.g.[48)]

- CP violation may be due to Higgs fields
- asymptotically the electro-weak interaction is left-right symmetric
- Higgs particles have supersymmetric partners
- the strong CP problem is solved by Higgs particles.

In all these cases, charged Higgs particles $H^{\pm}$ will occur.

Moreover, in dynamical schemes of symmetry breaking like extended technicolour (ETC) many Goldstone bosons become physical particles and acquire mass (pseudo-Goldstone bosons, PGB). Some of these lie in the mass range accessible to high energy e^+e^- machines. In particular, in a minimal ETC, the colour singlet 0^- PGB's acquire their mass from electroweak interaction. The mass of the charged states, $P^{\pm}$ is then

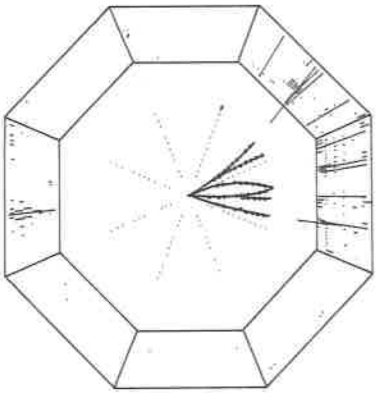
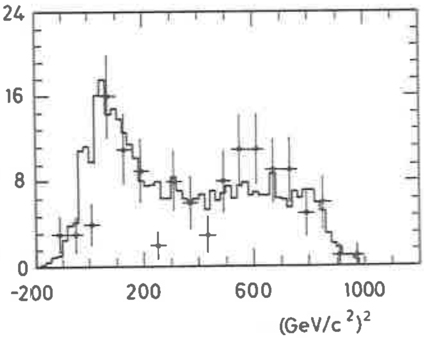


Fig. 23

Search for neutral Higgs in the reaction $e^+e^- \to \gamma$ + hadrons

a) hadronic event with an isolated photon as seen in the CELLO detector[50)]

b) Photon recoil mass spectrum (JADE) of events with an isolated photon[3c)]



predicted to be in the range[35)]

$$5 \text{ GeV} \lesssim M_{p^\pm} \lesssim 14 \text{ GeV} \qquad (26)$$

Extensive searches for charged Higgs particles or hyperpions have been carried out at PEP and PETRA. Both types of particle have the same production and decay properties

$$
\begin{array}{l}
e^+e^- \rightarrow H^+H^- \\
\qquad \text{or } P^+P^- \text{ decay into heavy fermions} \\
\qquad\quad \rightarrow \bar{c}s,\ \bar{c}b,\ \tau^-\nu \\
\qquad\quad \rightarrow c\bar{s},\ c\bar{b},\ \tau^+\nu
\end{array}
\tag{27}
$$

The relative branching ratios for hadronic and leptonic decays are given by

$$
\frac{BR(H^+ \rightarrow c\bar{s})}{BR(H^+ \rightarrow \tau^+\nu)} = \frac{m_c^2}{m_\tau^2}\,\chi \tag{28}
$$

where χ is, a priori, an unknown parameter. In different models, χ may vary over a wide range (χ = 3 for Higgs, 1/3 or 27 for PGB's)[51]. We shall, therefore, consider χ as a free parameter and discuss the following cases (Fig. 24):

a) $\chi \ll 1$: $H^\pm \rightarrow \tau^\pm\nu$ Predominant

This situation is illustrated in Fig. 24a. Both H decay into τ and ν. The events look identical to the ones expected for supersymmetric τ's. If λ_γ is undetectable, $H^\pm$ decaying with 100% branching ratic into $\tau\nu$ is indistinguishable from s_τ or t_τ. Therefore, the limits derived for $H^\pm$ production can also be applied to s_τ (special case of $B_{\tau\nu}$ = 100% in Fig. 25, c.f. Table 6).

The CELLO[37] and MARK J [39] groups have used the decay $\tau \rightarrow$ 1 prong - signature (1) in Fig. 24a - to search for $H^\pm$ signals: an acoplanar pair of particles with missing energy. Signature (2) implies τ decays into all channels: 1 prong acoplanar with 3 or 5 prongs in the opposite hemisphere. The JADE collaboration used this signature in the combinations (1 prong - 3 prong) and (1 prong - 5 prong) whereas the MARK II group looked for the (1 prong - 3 prong) combination. The mass of the 3 or 5 prong system is restricted to less than the τ mass (< 2 GeV).

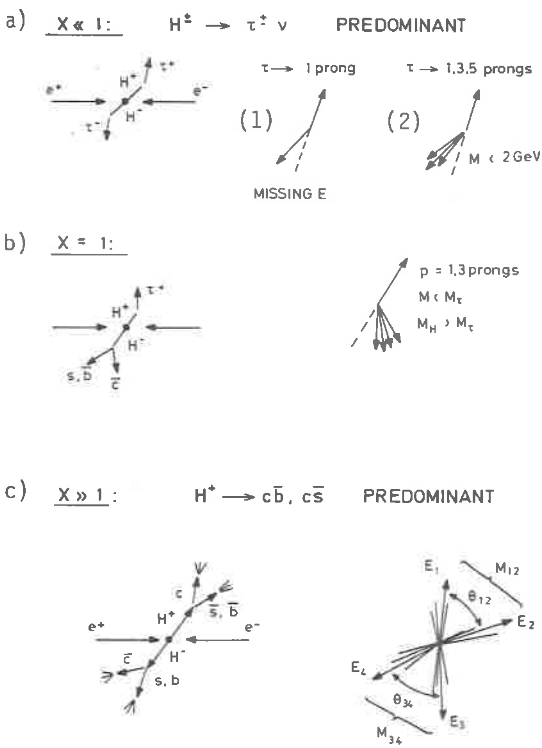


Fig. 24

Experimental search for charged Higgs particles and hyperpions

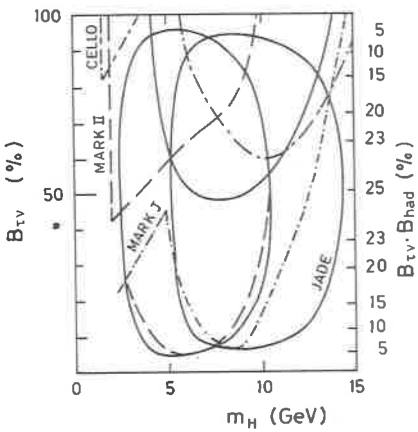


Fig. 25

Experimental bounds on the mass of the Higgs or hyperpions $H^{\pm}$ as a function of the leptonic and hadronic branching ratios. The areas included by the curves are excluded with 90% C.L.

b) $\chi = 1$

In this case, both the leptonic and hadronic decay modes occur with similar branching ratios. The situation is described in Fig. 24b.

Four groups have looked for events characterized by one or three prongs (originating from the τ) acoplanar with a hadronic jet of mass $M_H > M_\tau$ (c.f. Fig. 24b)

group	M_H	$\tau \to P$	M_P	Ref.
JADE	> 2.5 GeV	1,3 prongs + γ's	< 2 GeV	38
MARK II	> 2 GeV	1 prong + γ's	< 2 GeV	40
MARK J	> 2 GeV	μ	-	39
MAC	> 2 GeV	μ	-	13

The results of searches a) and b) are summarized in Fig. 25 which shows the experimental boundaries for $H^\pm$ production as a function of mass and branching ratio. Combining the results of all groups, the mass range from about 3 to 13 GeV can be excluded except for very low leptonic branching ratios of about 10%. This interesting area is again shown in Fig. 26.

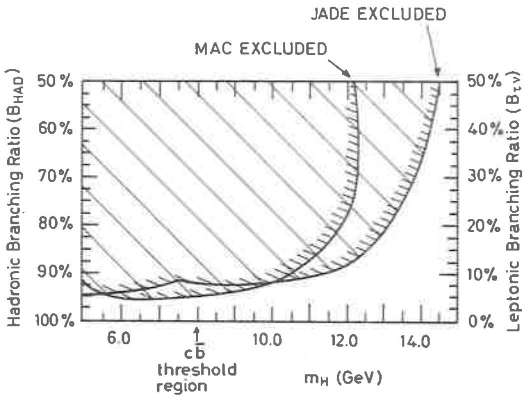


Fig. 26

Experimental bounds like in Fig. 25 (JADE, MAC)

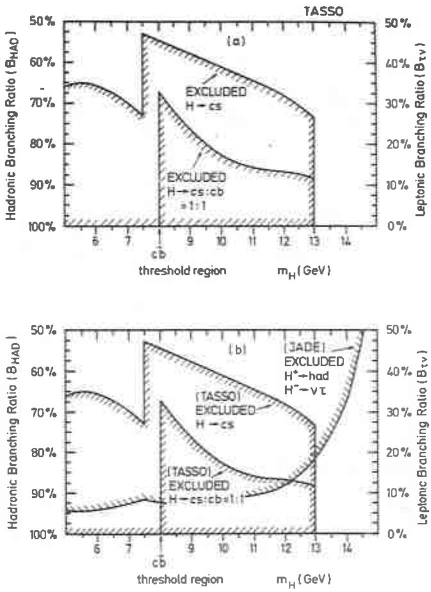


Fig. 27

Experimental bounds on the hyperpion mass (TASSO)[52]

c) $\chi >> 1$, $H^+ \to c\bar{b}$, $c\bar{s}$ predominant

To fill in the gap in Fig. 25, the TASSO collaboration has studied $H^{\pm}$ decays into hadrons according to Fig. 24c [52]. 75 pb^{-1} with about 20 000 hadronic events at $\sqrt{s}$ = 34.5 GeV energy were analysed. The experimental procedure is illustrated in Fig. 24c. The group has searched for four jet events and determined the jet energies, E_i, and the opening angles, θ_{ij}, between jets. Cuts are then applied according to the constraint that two pairs of jets have to form systems of equal invariant mass with beam energy. Two events survive in the 5 to 7.5 GeV mass range, no candidates are found between 7.5 and 13 GeV. The resulting limit on the leptonic branching ratio as a function of $H^{\pm}$ mass is shown in Fig. 27. Together with the limits from Fig. 26, this excludes charged Higgs particles or technicolour particles in the mass range 5 to 13 GeV. Similar results are quoted by the MAC group[13].

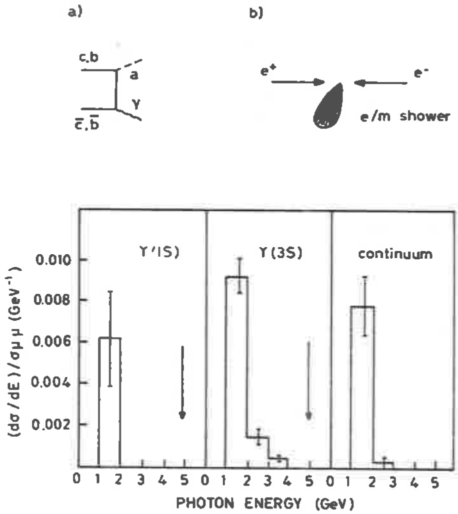


Fig. 28

Search for the axion at the Y resonances. Events with one single photon of beam energy ($E_\gamma > 5$ GeV). Arrows indicate, where an axion signal would be expected (CUSB)[55]

3.2.3 Search for the Axion

To explain the absence of P and T violation in QCD, it has been proposed to introduce two Higgs doublets ϕ_1 and ϕ_2. In such a model, a light neutral pseudoscalar particle appears, the axion a [53]. Similar to H^o, heavy quarkonia decay into an axion and a photon (Fig. 28a). The branching ratio for the decay of J/ψ and Y can be precisely predicted up to a parameter $X = \langle\phi_2\rangle/\langle\phi_1\rangle$ which is the ratio of the expectation values of the two Higgs fields. Since c and b have opposite weak isospin, the two branching ratios have different proportionality to X

$$\begin{aligned} BR(J/\psi &\rightarrow \gamma a) \sim X^2 \\ BR(Y &\rightarrow \gamma a \sim 1/X^2 \end{aligned} \tag{29}$$

so that the product of the two is a precisely predictible number

$$BR(J/\psi \to \gamma a) \times BR(\Upsilon \to \gamma a) = B^{\Upsilon}_{\mu\mu} \cdot B^{J/\psi}_{\mu\mu} \; \frac{G_F^2 m_c^2 m_b^2}{2\pi^2\alpha^2}$$

$$= (1.6 \pm 0.3) \times 10^{-8} \; . \qquad (30)$$

The experimental signature for reaction (29) is rather clear (Fig. 28b)

$$e^+e^- \to J/\psi \;\text{ or }\; \Upsilon \to \gamma + \text{NOTHING} \qquad (31)$$

if the axion decays outside the detector (M_a << 10 MeV).

The Crystal Ball group has published a negative result for a search at the J/ψ resonance[54)]

$$BR(J/\psi \to \gamma a) < 1.4 \cdot 10^{-5} \; . \qquad (32)$$

Recently, the CUSB and CLEO groups have looked for a signal at the $\Upsilon(1s)$ and $\Upsilon(3s)$ resonances. Fig. 28 shows the CUSB result. No events have been found (arrows in Fig. 28) which implies the following limits[55)]

$$\begin{array}{llll} \text{CUSB:} & BR(\Upsilon(1s) \to \gamma a) & < 3.5 \cdot 10^{-4} & (90\% \text{ C.L.}) \\ \text{CUSB:} & BR(\Upsilon(3s) \to \gamma a) & < 1.2 \cdot 10^{-4} & (90\% \text{ C.L.}) \\ \text{CLEO:} & BR(\Upsilon(1s) \to \gamma a) & < 9.0 \cdot 10^{-4} & (90\% \text{ C.L.}) \end{array} \qquad (33)$$

The combined result of (32) and (33) is then

$$BR(\Upsilon \to \gamma a) \quad BR(J/\psi \to \gamma a) < 0.6 \times 10^{-9} \quad (90\% \text{ C.L.})$$

which is more than an order of magnitude lower than the expectation (30). Thus the "standard" axion is ruled out.

4. Summary

We can summarize the search for new particles as follows

Fermions

- There is now (indirect) experimental evidence that the "old" sequential heavy lepton τ has its own neutrino.
- No new leptons (sequential, stable, neutral, excited) have shown up in the PEP/PETRA energy range.
- The mass limit on toponium is larger than 36.75 GeV.
- No free quarks or monopoles have been seen at PEP/PETRA.

Scalars

- If they exist, supersymmetric leptons have masses larger than about 16 GeV.
- Charged Higgs particles are excluded in the mass range from 5 GeV to 13 GeV.
- Charged technipions do not exist in the proposed mass range from 5 GeV to 13 GeV.
- The "standard" axion does not exist.

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PRELIMINARY RESULTS FROM THE NUSEX EXPERIMENT ON NUCLEON STABILITY

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This paper will be mainly devoted to the NUSEX experiment on nucleon stability which started operating a few months ago at the Mont Blanc Laboratory; recently published data from the Kolar Gold Field experiment will be also discussed.

The NUSEX experiment [1)]

The laboratory. Our detector is installed in the "Laboratorio di Cosmogeofisica del CNR", in the Mont Blanc Tunnel connecting Italy to France. The overburden of the laboratory has a complicated shape, due to the structure of the mountain-chain, but it is always thicker than $4.5 \cdot 10^3$ m of water equivalent (Fig. 1), allowing a great reduction in atmospheric muon flux.

The detector. It consists of a uniform sandwich of iron plates 1 cm thick interleaved with layers of limited stramer tubes (Fig.2). The total volume is 3.5 x 3.5 x 3.5 m^3 corresponding to a total mass of $\sim$150 tons ($\sim$ 134 tons of iron and $\sim$ 16 tons of PVC). Resistive cathode limited streamer tubes are the basic elements of the sensitive part of the detector. These tubes, developed in Frascati by Battistoni et al.[2)], show a large high voltage plateau, large output signals, good efficiency and moreover do not need an external trigger and allow a very simple method for bidimensional read-out of multitube layers. Resistive cathodes are transparent to the pulsed e.m. field associated with any pulsed current generated around the sense wire; thus pick-up electrodes can be placed outside the tubes and pulses collected by these electrodes. In our case both coordinates of the streamer are read by means of external strips (Fig. 3).

In total there are 43.000 limited streamer tubes and $\sim$82.000 read-out channels.

Signals from read-out strips are fed in amplifying, discriminations and recording circuits on Le Croy cards (32 channels/card) and finally collected by a PDP 11/60 computer.

Trigger. At present the basic element for triggering the data acquisition is a prompt OR among signals from strips of a whole plane. The counting rate of a single plane is $\sim$ 350 Hz. "Good triggers" are selected requesting the following patterns

- an AND among four consecutive planes
- an AND among three adjoining planes and two other consecutive planes
- other configurations of greater complexity.

Trigger rate is about 6 h^{-1}, and is dominated by radioactivity coincidences.

Starting from the beginning of August a new piece of data has been added to the previous ones. Whenever a "good trigger" is recorded, a clock for each plane is started, and it stops only either when the first hit is recorded from that plane or when it is reset by the computer. This procedure allows to store the relative time delay of hits and eventually the one that comes from the decay of a muon stopping inside the detector.

Efficiency for μ^- is almost 0 (due to the μ^- being absorbed in Fe nuclei before decay), while it is $\simeq$ 35% for μ^+; time resolution is $\simeq$ 100 nsec.

In this way we can say in 17% of the cases which way the track goes, and its charge.

Tests. During '81 a test module with the same features (1 cm iron plates, 3.5 m long tubes, same electronics ...) as the Mont Blanc detector, but smaller dimensions, was exposed to e, π and ν beams at CERN.

The ν data are of special interest, in fact the unavoidable background that can mimic genuine nucleon decay is due to atmospheric neutrino interactions.

Atmospheric and accelerator ν's have similar spectra (Fig.4) except in the very low and very high energy regions, but accelerator beams contain ν_μ only and have a single direction, while both ν_μ's and ν_e's are present in cosmic radiation and enter isotropically the detector.

Our test detector was exposed to an unfocused ν beam from 10 GeV protons, in two geometrical conditions: at 0° (i.e. with beam orthogonal to the iron plates) and 45°. In total $\sim$ 400 events (50% at 0° and 50% at 45°) were collected. They have been classified according Table I, but a new and more refined analysis is in progress.

TABLE I

Neutrino events from test runs	
1 prong	209 event's
2 prongs	98 "
≥ 3 prongs	32 "
1 prong + e.m. shower	19
2 prongs + " "	6
neutral current candidates	17

Results. The detector started operation on May '82 for preliminary tests, with 2/3 of the volume active. Up to now useful data have been collected during two periods: June and July (∿ 1050 hours) with 84% of the resistive tubes in operation, August and September (∿ 1160 hours) with the detector fully active. The two runs correspond to ∿ 34.5 tons x x year; collected data are synthetically reported in Table II.

TABLE II

single muons	2490 (∿1.18 μ/h)
muons stopping in the detector	21
parallel muon	34
fully contained events	4

Examples of these events are shown in Fig. 5; azimuthal angle distribution for single μ's, shown in Fig. 6, reflects the shape of the mountain covering the detector. Atmospheric muons not only provide an useful and continuous monitoring of the detector, but they also are interesting by themselves. However discussion will be restricted to fully confined events which are directly connected with the problem of nucleon stability.

Ev. 9-532 (Fig.7a). It is a single track event; if the particle is a μ, its total energy is E_μ=320 ± 15 MeV.

Ev. 122-526 (Fig.7b). It is interpreted as a single μ accompanied by soft electrons; its total energy is E_μ=370 ± 20 MeV.

These two events are probably due to ν_μ or $\bar{\nu}_\mu$ interactions.

Ev. 40-524 (Fig.7c). It consists of two prongs plus an electromagnetic shower; the interaction point and the direction of the primary particle are unmistakable; two interpretations are possible:

$$\nu_\mu(\bar{\nu}_\mu) + N \rightarrow N' + \mu + \pi^{ch} + \pi^0 \quad \text{or}$$
$$\nu_e(\bar{\nu}_e) + N \rightarrow N' + e + \pi^{ch} + \pi^{ch}$$

The computed ν energy and momentum are 1.3 ± 0.2 GeV and 1.1 ± 0.2 GeV; these figures support the ν interpretation, being nearly equal.

Ev. 19-503 (Fig.7d). It is a 3 prong event, of $\sim$ 1 GeV total energy and approximately balanced in momentum. Let us consider first the proton decay hypothesis.

a) $p \to 3\mu$. All the track are assumed to be muons; the total energy is 0.90 ± 0.15 GeV and the total momentum 0.3 ± 0.1 GeV.

b) $p \to K^0 + \mu^+$. Track AB is interpreted as a μ while AC and AD should be pions from $K^0{}_s$ decay; the $\pi\pi$ invariant mass is 0.55 ± 0.08 GeV and the μ total energy 0.38 ± 0.15 GeV. This last value must be compared with the predicted one (for a proton decay at rest) of 0.34 GeV. Total visible energy is 1.0 ± .2 GeV and momentum is 0.4 ± 0.2 GeV/c.

c) Other possibilities like $p \to K^*\nu$ $(K^* \to K^0_s\pi)$ or $p \to \mu\rho^0$ $(\rho^0 \to \pi^+\pi^-)$ are kinematically consistent with the above mentioned hypothesis.

Now the crucial point is a correct estimate of the ν background. We could envisage only three possibilities:

i) A ν interaction in A producing three charged particles. Visible energy (1.0 ± 0.2 GeV) and total momentum (0.4 ± 0.2 GeV) are largely inconsistent. Furthermore, in our ν sample collected in the test runs, only one of the 3-prong events shows an angle between the μ candidate and any one of the two other tracks greater than 140°, while track AB and AD form an angle of 160° ± 7°. We conclude that the probability that the event be neutrino induced is $\lesssim$ 1%.

ii) A neutrino interaction in D producing a μ (DB) and a π(DAC) which scatters in A. The angle between the particles must be < 10°. In the neutrino test run, we did not observe any two-prong events with an opening angle less than 15°. We conclude that the probability of such a hypothesis is $\lesssim$ 1%.

iii) A neutrino interaction producing a single pion; the probability of this hypothesis is negligible in comparison with (i) and (ii).

In conclusion, the event is kinematically compatible with a genuine proton decay, if one takes into account the error measurements and Fermi motion in nuclei.

The probability that it is due to a ν (or a n) interaction is small ($\sim$ 0.05 event expected), but not completely negligible and partially based on the ν test run data which refer to ν impinging on the detector at 0° and 45° only.

Nucleon mean life is given by

$$\tau_{nucleon} = \frac{\text{mass x time x efficiency}}{\text{n. of events}} \quad ;$$

therefore assuming an overall efficiency of 0.5, 1 event corresponds to 10^{31} y's. This figure can be considered as a lower limit or as a rough estimate of the mean life, according to the preferred interpretation of the candidate event. This limit (or value) is indeed for $B\tau$, when B is the branching ratio for decay modes in "visible" particles (μ, $\pi^{\pm}$, π^0, K ... but not ν's).

Comments on the Kolar Gold Field data [3)]

I was requested by the organizers to comment the recently published results obtained by the Indo-Japanese collaboration in more than one year of measurement with a detector installed in the very deep mine of Kolar Gold Field (KGF). It is not an easy task, because it is always difficult to collect and analyze all the details of a detector and to envisage all possible interpretations of few and "delicate" events like those we are considering.

Certainly the pioneering work of this collaboration must be acknowledged but, at least in my opinion, their positive conclusion about the existence of the nucleon instability is weakly supported by their experimental data.

Table III summarizes (and compares) the main feature of the KGF and Mont Blanc experiments and it can help in understanding some of the considerations which follow.

TABLE III

	KGF	NUSEX
overburden	7000 m w.e.	4500 m w.e.
μ flux	$\sim$ 1.8 μ/d	$\sim$ 1.2 μ/h
mass	140 t	150 t
iron thickness	(1.2+0.46)x2 cm per view	1 cm
n. of detectors	1600 proportional detectors, 10x10 cm^2 cross section measure of the deposited energy in counters (±30%)	$\sim$ 42000 limited streamer tubes, 1x1 cm^2 cross section
calibration	energy deposition in tubes with radioactive sources	tests to e, π and ν beams
ν induced confined events (+)	0 in 461 d's	3 in $\sim$2000 h's
p decay candidate confined	3 " "	1 " "

(+) The ν flux is different at different latitudes.

The 3 fully confined events are shown in Fig.8. Two general remarks can be done: i) the granularity of the detector does not allow an easy identification of the topology of the events (i.e. interaction or decay point, number and nature of emitted particles) because of the very low (a few units) number of hits per view; ii) the absence of fully contained ν induced interactions.

Let us now consider the 3 events in some more details (Fig.8a,b,c).

Ev. 867. The authors interpret this event as due to a proton decay in a π^{ch} and a ν ($p \to \pi^{ch} \nu$), because the total pion energy is $\sim$435 (±?) MeV, in agreement with the decay kinematics. I would remark that this interpretation is based on a single hit (hit nearby point C in 6M view). How large is the probability that this track is due to a μ^- produced in a ν_μ elastic interaction and captured at the end of its range with emission of a soft γ or n? In any case I believe that the observation of a single track carrying $\leq$ 1/2 of the total energy liberated in a nucleon decay cannot be assumed as a proof of the existence of this process.

Ev. 587. It has been interpreted as $p \to e\pi$ decay; possible background reactions are:

$$\text{(i)} \quad \nu_e(\bar{\nu}_e) + N \to e^-(e^+) + N'$$

$$\text{(ii)} \quad \nu(\bar{\nu}) + N \to \nu(\bar{\nu}) + \pi^0 + N'$$

The number of events due to reactions (i) and (ii), as evaluated by the authors, is (in total) $\leq$ 0.5 event, then not negligible. The ν hypothesis is disliked on the basis of Monte Carlo calculations on electromagnetic showers. In view of uncertainties in the ν flux and cross sections and in M.C. computations without experimental checks on the very same detector, it seems to me quite difficult to reject the hypothesis that this event is due to a ν interaction.

Ev. 877. Two interpretations are given: $n \to e^+\pi^{ch}$ and $p \to \mu K^0_s$. The second one is very hard to be demonstrated; in fact the two pions from K^0_s decay give 3 hits in total in one view (6M view) and one pion leaves no track in the other view.

First hypothesis is compatible with the observed pattern of the event. I would only remark that: track BA is not so clearly due to a showering particle (in fact the authors also interpret it as a charged π) and the interaction point is somewhat undefined. Moreover other interpretations could be exploited; for instance, tracks CB and BA form a single track due to a μ of $\sim$ 600 - 700 MeV, the extra hit in 6M view being due to some soft γ-ray.

All this criticism is, of course, due to the relevance of the problem and to the difficulties of such a kind of experiment where background is not completely known and statistics is very low.

Conclusions

I belive that present situation[4] can be summarized in the following way:

- the KGF data cannot prove unambiguously the existence of the nucleon decay at a level of $\tau_p \gtrsim 10^{31}$ y's, because the detector granularity does not allow a clear identification of the event topology;

- the Mont Blanc event is a reasonable good candidate, but the hypothesis that it is due to a ν interaction cannot be completely excluded; so it is appropriate to wait for more statistics;

- if the $e^+\pi^0$ KGF candidate is a genuine p-decay, the IMB experiment, at present in operation, would give a clear answer in a short time unless background problems due to cosmic rays in their quite shallow laboratory will result too severe.

- the Freyus experiment, which would become operating next year, would give valuable information on processes like $p \to \mu k^0$.

- present data seem to justify the construction of a new and improved (in respect to the Nusex detector) digital calorimeter of large mass (> 1000 tons) possibly in a short time.

Acknowledgements

I should like to thank the members of the NUSEX collaboration, in particular Profs. E. Fiorini, E. Iarocci, P. Picchi, A. Pullia and L. Zanotti.

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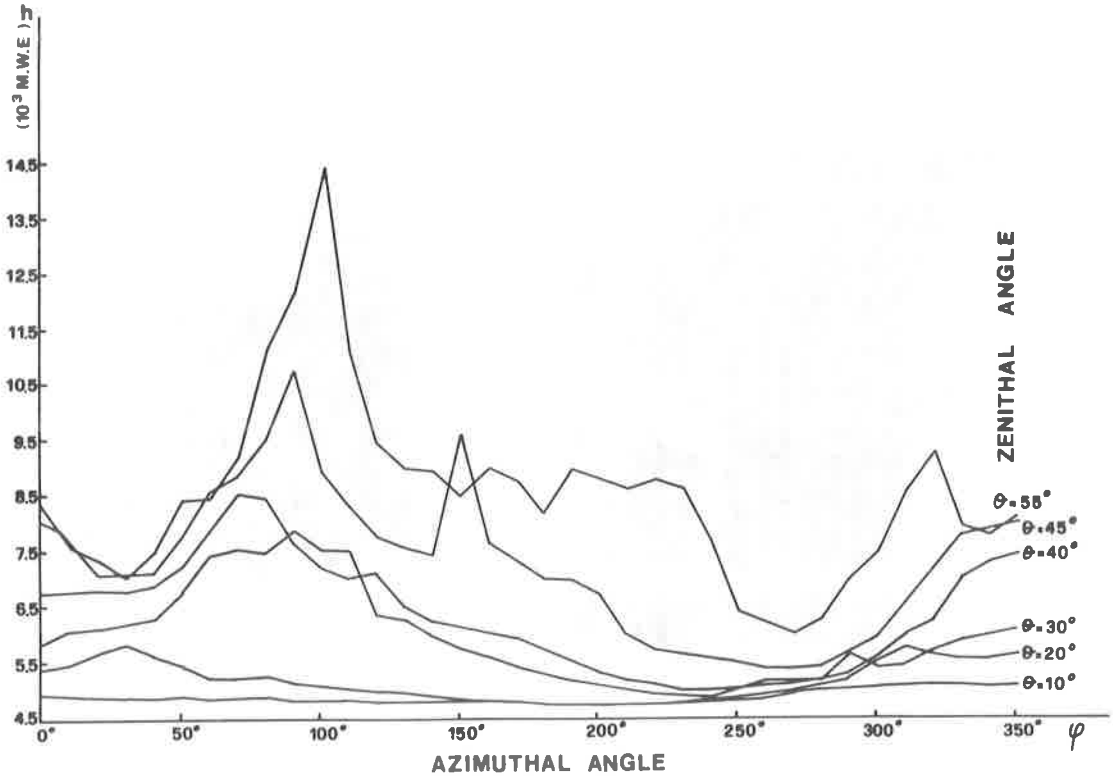


Fig. 1. Overburding as a function of zenithal and azimuthal angles.

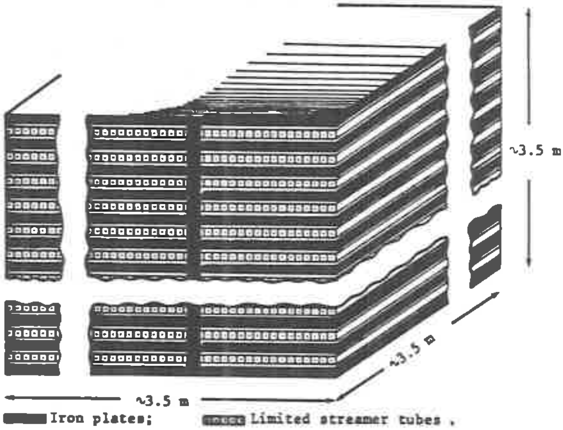


Fig. 2. Sketch of the detector.

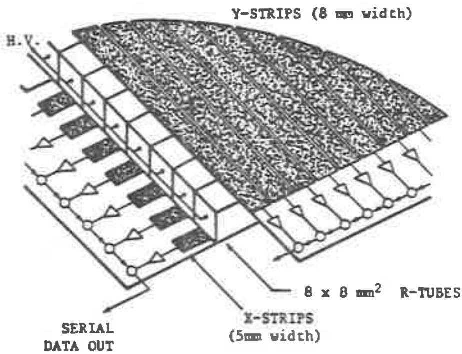


Fig. 3. Read-out.

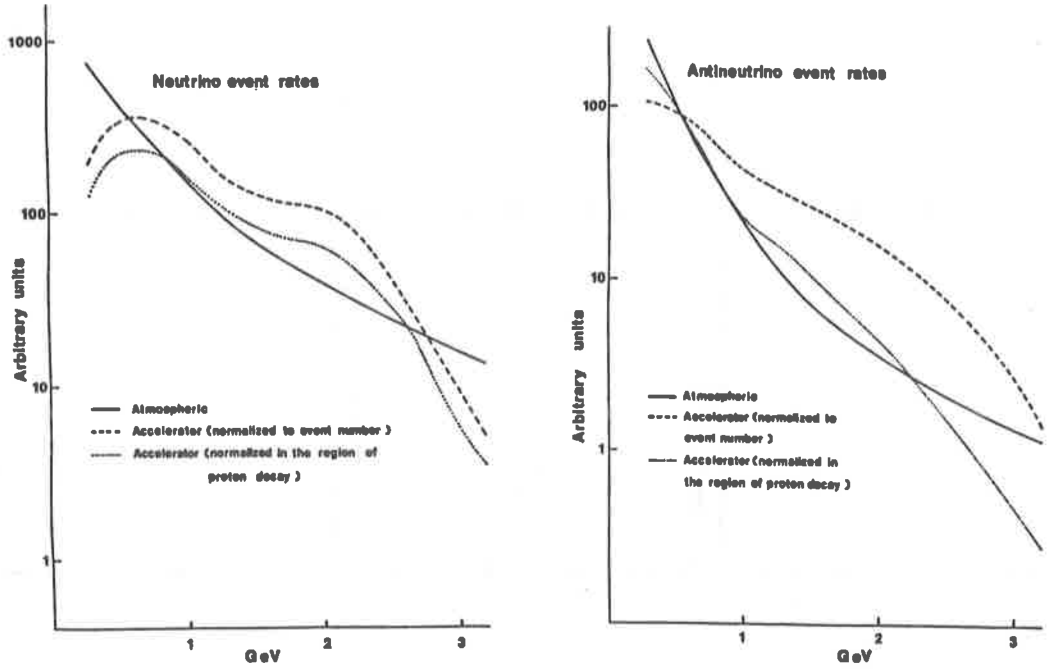
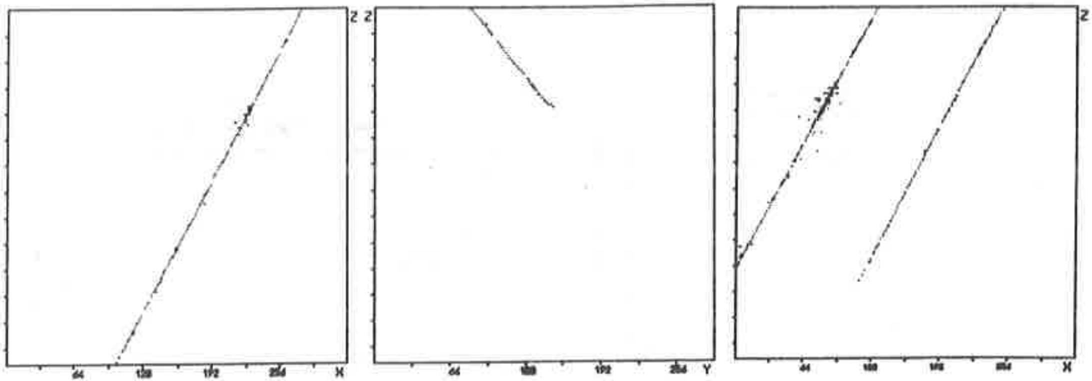


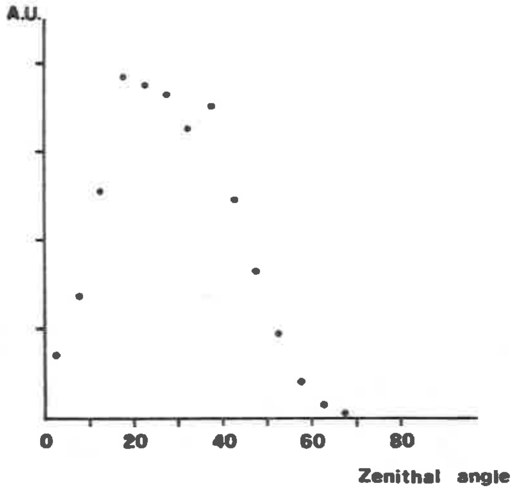
Fig. 4. Event rates from atmospheric and accelerator ν and $\overline{\nu}$.



a) b) c)

Fig. 5. a) single μ; b) stopping μ; c) parallel μ's

a)



b)

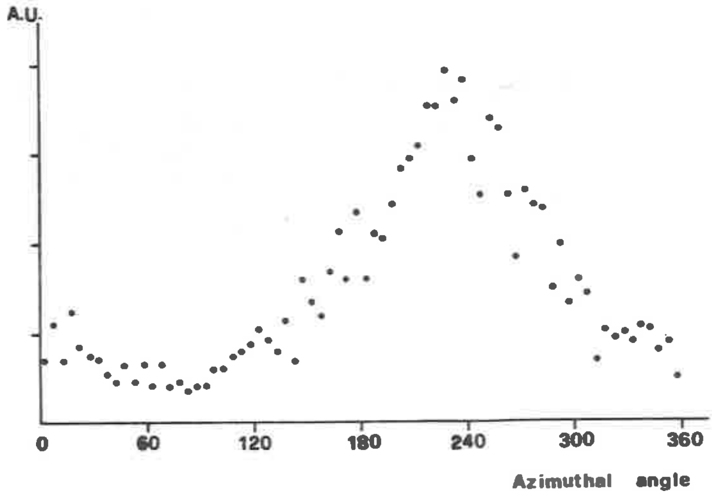


Fig. 6. Single muons; a) zenithal distribution;
b) azimuthal distribution.

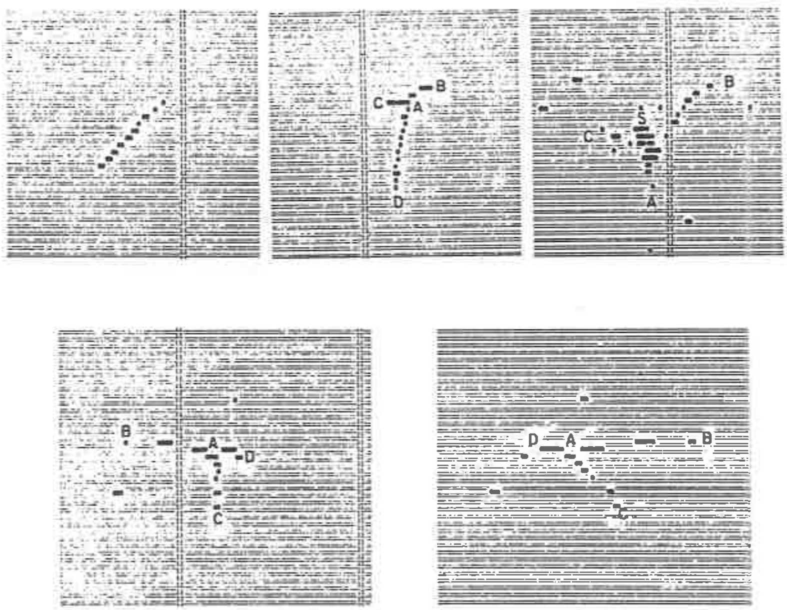


Fig. 7. Confined events (see text).

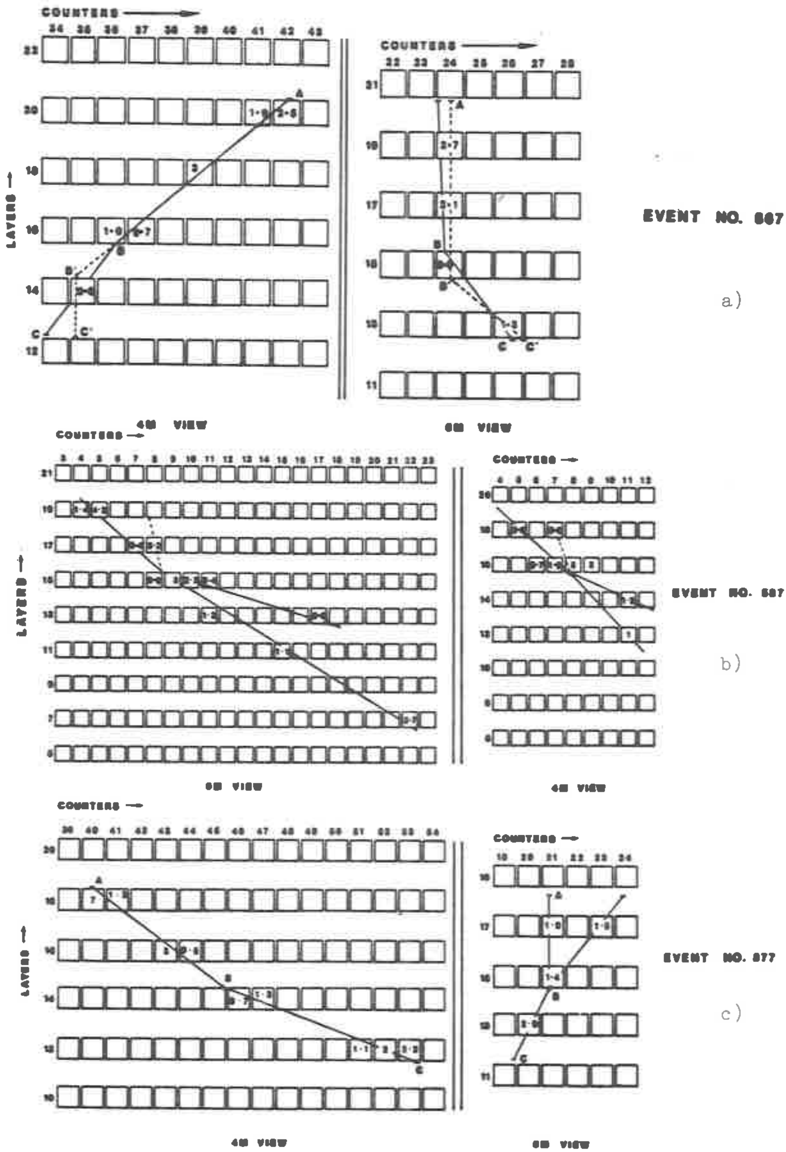


Fig. 8. Confined events observed in the KGF experiment (see text).

Seminars

ELECTROWEAK GAUGE THEORIES WITH EXTRA U(1)'s IN SO(10) AND E_6

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Abstract

We have examined the possibility of extending the standard $SU(2)_L \times U(1)$ electroweak theory by adding only $U(1)$ factors within the context of $SO(10)$ and E_6 grand unification.[1,2] Besides considering the left-right symmetric breakdown of $SO(10)$ through an $SU(4) \times SU(2)_L \times SU(2)_R$ stage which can lead to an $SU(2)_L \times U(1)_R \times U(1)_{B-L}$ electroweak group we consider the breakdown $SO(10) \to SU(5) \times U(1)_\chi$ (via an adjoint 45 Higgs) which leads to an $SU(2)_L \times U(1)_Y \times U(1)_\chi$ gauge group[3] and no new physics in the "desert" except one additional $U(1)$, i.e., a second Z . This $U(1)_\chi$ is generated by the linear combination $\sqrt{2/5}\, T_{3R} - \sqrt{3/5}\, T_{B-L}$. Breaking this symmetry by giving vacuum expectation values (VEVs) to the neutral members of $SO(10)$ $\underline{16}$'s we obtain the usual charged current interactions and the neutral current interaction

$$H^{NC}_{SO(10)} = \frac{4G_F}{\sqrt{2}} \left[(J_{3L} - \sin^2\theta_W\, J_{em})^2 + \frac{1}{R} (J_{3R} - \frac{3}{5} \cos^2\theta_W\, J_{em})^2 \right]$$

where $R = V^2/v^2$ is a ratio of VEVs. A similar expression is obtained for the left-right breakdown with the substitutions $\sin^2\theta_W \to e^2/g_L^2$ and $3\cos^2\theta_W/5 \to e^2/g_R^2$ by a suitable rotation of the diagonal generators.[4] This interaction is identical to the standard model one for neutrinos. Parity violation experiments in atomic physics and measurements of the forward-backward asymmetry in $e^+e^- \to \ell^+\ell^-$ limit R to be bigger than 10 . This forces the lighter Z_1 (now almost the Z^0) to satisfy $1.0 \geq M_1/M_{Z^0} \gtrsim 0.98$ while the heavier Z_2 (now nearly the Z_χ boson) must have $M_2/M_{Z^0} \gtrsim 2.5 - 3.0$.

A similar analysis is performed for the "minimal" breakdown of $E_6 \to SO(10) \times U(1)_\psi \to SU(5) \times U(1)_\chi \times U(1)_\psi$ which allows three low mass Z's. The neutral currents can also be arranged to coincide with the standard model predictions for neutrinos and similar limits are obtained on the masses of the light Z and the two additional heavy Z's.

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CP VIOLATION AND R INVARIANCE IN SUPERSYMMETRIC MODELS OF STRONG AND ELECTROWEAK INTERACTIONS

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Abstract

We investigate[1] the general structure of the minimal supersymmetric extension of the standard model of strong and electroweak interactions. Certain terms in the effective Lagrangian, which may yield too large a neutron electric dipole moment, are not allowed by R invariance. The hypothesis of a partially conserved R current suppresses also the transition $\mu \to e\gamma$ as well as dimension-5, B- and L-nonconserving operators; furthermore, it has interesting implications for the strong CP problem. The interplay of soft supersymmetry breaking terms and gauge interactions allows for a generation of fermion masses without Yukawa couplings. The mechanism necessarily renders the neutrino massless.

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U(n) SUSY GUTS

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There are three fundamental problems in the minimal supersymmetric SU(5) model of Dimopoulos-Georgi-Sakai[1]: i) An unnatural fine tuning of parameters is necessary to adjust the masses of Higss doublets to zero. This problem has been handled by Ibañez, Ross, Nanopoùlos and Tamvakis[2] by the introduction of a gauge singlet superfield decoupled from the rest of fields and, therefore, with an arbitrary v.e.v. However, this method changes the fine tuning of a mass parameter by that of a v.e.v., which is not much more natural. ii) There is not any invariance that prevents a violation of the baryon number at $d < 6$. In particular, $d = 5$ operators give a short proton lifetime $\sim m_X^2$ (instead of m_X^4). A solution to this problem has been advocated by Ellis, Nanopoulos and Rudaz[3] showing there are suppression factors coming from renormalization effects. However, they come into troubles when four (or more) light higgses are used. Anyway, it should be desirable to get rid of $d = 4, 5$ baryon number violating operators. iii) SUSY has to be broken softly and, therefore, it is not possible to give a mass to the higgsinos.

Problems ii) and iii) have been already considered by Weinberg[4] in the supersymmetric standard model. He comes to the conclusion of the necessity of introducing an additional $\tilde{U}(1)$ preventing the presence, in the superpotential, of terms violating baryon number. All three problems can be solved in a $U(n)$ SUSY GUT. Since $U(n)$ is not semi-simple two additional problems do arise: a) In addition to $(SU(n))^3$-anomalies, there are $(SU(n))^2\tilde{U}(1)$ and $\tilde{U}(1)^3$-anomalies which must be eliminated, otherwise the theory is not renormalizable. b) A $\xi\tilde{D}$-term is perturbatively generated, i.e. $\xi \sim m_{Pl}^2$, unless $\mathrm{tr}Y = 0$[5] or $\tilde{U}(1)$ is unified in a semi-simple group[6].

We shall study the symmetry breaking pattern $U(5+n)\times SUSY \to U(5)\times SUSY \to \tilde{U}(1)\times SU(3)\times SU(2)\times U(1)\times SUSY$, where the first breaking is accomplished via the Higgs superfield $\Phi^{r_1\ldots r_n}$, $\Phi_r^{(i)}$, $i=1,\ldots,n$, and the second breaking through the adjoint representation Σ. The condition $\tilde{D} = 0$ will lead to the canonical assignment of charges: for a tensor with m upper and m' lower indices, $\tilde{Y} = m - m'$. The model we propose is given by the superpotential $f = \lambda\,\mathrm{tr}\,\Sigma^3 + m\,\mathrm{tr}\,\Sigma^2 + m'H''\Sigma H$

where Higgs chiral superfields in the fundamental (H) and anti-fundamental (H') representations, suitable for the electroweak breaking, are introduced. SUSY conditions $f_\Sigma = f_H = f_{H'} = 0$ lead to massless WS doublets <u>without any fine tuning</u> between m and m' [7]. The problem of anomalies can be solved as follows. Let f_L be the set of N_k chiral supermultiplets $\Phi^{r_1 \ldots r_k}$ $(k = 1,\ldots,n)$. The condition of cancellation of anomalies implies $N_i = \sum_4^n a_k^{(i)} N_k$ $(i = 1,2,3)$ with $a_k^{(1)} = -(k-3)(k-2)\,\Gamma(n-1)/2\Gamma(k-1)\,\Gamma(n-k)$, $a_k^{(2)} = (k-3)\,\Gamma(n-2)/\Gamma(k-2)\,\Gamma(n-k)$ and $a_k^{(3)} = -\Gamma(n-3)/\Gamma(k-3)\,\Gamma(n-k)$. Thus, for any set $(N_4,\ldots,N_n)$ of integer numbers we get an anomaly-free set. Let us note that the $\xi\tilde{D}$-term is not perturbatively generated since $\mathrm{tr}\tilde{Y} = 0$ for the above solution.

The above ideas can be particularized for the "minimal" U(5) SUSY GUT. Absence of anomalies leads to p light (ordinary) families $\underline{10}(2)+\underline{\bar{5}}(-6)$ and p heavy families $\underline{10}(-3)+\underline{\bar{5}}(-1)+\underline{1}(-5)$. Four higgses are needed with non canonical assignment of $\tilde{Y}$ charges, but free of anomalies and with $\mathrm{tr}\tilde{Y} = 0$: $\underline{5}(6)+\underline{\bar{5}}(4)$, which give a mass to heavy fermions, and $\underline{5}(-4)+\underline{\bar{5}}(-6)$ giving a mass to the ordinary fermions. B and L violating $d = 5$ operators due to the exchange of massive Higgs triplets do not appear because $\tilde{U}(1)$ conservation. To draw a realistic model, electroweak and SUSY breaking patterns must be described. The problem of SUSY breaking is the hardest one since a Fayet-Iliopoulos scheme, $\tilde{D} \neq 0$, is very difficult to achieve since squarks and sleptons can acquire non vanishing v.e.v.'s leading to a supersymmetric vacuum breaking colour and/or lepton number. A way out is to introduce $N = 1$ supergravity coupled to the SUSY GUT, and a Higgs superpotential like $f = \rho\underline{5}(6)\,\underline{\bar{5}}(-6)+\mu\underline{5}(-4)\,\underline{\bar{5}}(4)+\Delta$. Nevertheless, this superpotential is forbidden by R-invariance[8] unless $\xi = 0$, which is technically permitted since $\mathrm{tr}\tilde{Y} = 0$ for canonical solutions. This would achieve a possibly realistic local U(5) SUSY GUT.

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A COMPLETE HIGGS STRUCTURE FOR SUSY-SU(5)

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Abstract

We construct a complete supersymmetric SU(5) model where all terms consistent with R-invariance and gauge symmetry are included in the superpotential. For all values of the tree parameters the model exhibits a spontaneous breakdown (S.B.) of supersymmetry (Susy), R-invariance, and SU(5) with unique breaking patter $SU(5) \to SU(3) \times SU(2) \times U_Y(1)$. Within a broad range of the tree parameters radiative effects trigger a second breakdown $SU(2) \times U_Y(1) \to U_Q(1)$ at a scale which is hierarchically small. Apart from heavy particles and the phenomenologically required light particles the model contains an invisible Goldstone fermion stemming from the S.B. of Susy and an equally invisible axion related to the S.B. of the anomalous T-invariance. As a by-product the strong CP-problem is solved. Flavour changing neutral currents are suppressed by a super-GIM-mechanism.

The superpotential W is composed out of the following parts:

1. An O'Raifeartaigh part[1] to obtain S.B. of Susy and SU(5)

$$W_{OR} = X(g \operatorname{tr} A^2 - m^2) + f \operatorname{tr} YA^2 \quad .$$

The complex fields A and Y transform as $\underline{24}$ under SU(5), X is a singlet. The minimum of the potential $V = \sum_{\text{all fields}} |\partial W/\partial\phi|^2 +$ gauge part is uniquely acquired for $\langle A \rangle \propto \operatorname{diag}(2,2,2,-3,-3)$, the magnitude of $\langle Y \rangle \propto \langle A \rangle$ is fixed by radiative corrections[2].

2. A hierarchy-generating part with a sliding singlet Z

$$W_H = H'(\gamma_{24}\Sigma + \gamma_1 Z)\, H + h^i_j\, H'\, M_i\, M'^j + f^{ij}\, H\, M_i\, M_j$$

with Higgs fields H' and H ($\underline{\bar{5}}$ and $\underline{5}$ under SU(5)) to break $SU(2) \times U_Y(1) \to U_Q(1)$ and give masses to quarks and leptons which are contained in M_i and M'^j ($\underline{10}$ and $\underline{\bar{5}}$ under SU(5); $i,j = 1,2,3...?$ is a family index). Σ is a $\underline{24}$ under SU(5) and gets vev $\langle\Sigma\rangle \propto \operatorname{diag}(2,2,2,-3,-3)$. The vev $\langle z \rangle$ is undetermined at tree level and slides so as to make the doublet components of H' and H massless. It thereby generates a natural hierarchy without finetuning[3].

3. A piece to copy $\langle A \rangle$ on $\langle \Sigma \rangle$:

$$W_{copy} = \mathrm{tr}\ \chi(\sigma\ \phi_{1+r/2}\ \Sigma + \alpha\ \phi_r\ A) \quad .$$

χ is an additional $\underline{24}$, the fields ϕ are singlets with R-charge given by the subscript. (Each term of W has R-charge 2.)

4. A completion of the O'Raifeartaigh part to give non-vanishing vevs to the singlet fields

$$W_{singlet} = X(\delta_1\ \phi_r\ \phi_{-r} - m^2) + \delta_2\ \phi_{-r}^2\ \phi_{2+2r} + \delta_3\ \phi_r^2\ \phi_{2-2r} \cdot$$

(The term $-m^2X$ is already contained in the first part W_{OR} .)

5. A piece to fix the value r of the singlet R-charges

$$W_{Fix} = \varepsilon_1\ \phi_{-r}\ \phi_{1+r/2}^2 + \varepsilon_2\ \phi_{1+r/2}\ \phi_{1/2-r/4}^2$$

$$+ \varepsilon_3\ \phi_{1/2-r/4}\ \phi_{3/4+r/8}^2 + \varepsilon_4\ \phi_{3/4+r/8}\ \mathrm{tr}\ \Sigma \cdot A \quad .$$

The value of $r = 18/5$ turns out to be sufficiently complicated to exclude additional terms not listed above. Without arbitrarily finetuning terms to zero we can therefore keep the singlet $\langle Z \rangle$ sliding and obtain a natural hierarchy with correct SU(5) breaking pattern. Granted the structure of $W_{OR} + W_H$, the model has the minimal structure sufficient for the solution of the hierarchy problem.

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$\sin^2\theta_W$ AND THE PROTON LIFETIME IN UNEXTENDED TECHNICOLOR MODELS

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Abstract: Recently Georgi and Glashow[1] proposed the so called unextended technicolor models (uTc). uTc models assume the technicolor theory to be a fixed point theory. The electroweak and strong interactions are unified in a simple group (in our case SU(5)). In contradistinction to extended technicolor models, no unification of technicolor with strong or electroweak interactions is performed. Fermions get their masses through heavy scalars $(M_S \sim M_G)$. Georgi[2] advocates that, in the neighbourhood of the fixed point, this strongly interacting theory can be rewritten as an effective, asymptotic free theory.

Assuming such an effective theory, we derive $\sin^2\theta_W$ and the proton lifetime in the framework of SU(5) unification. Following Weinberg[3], we define below each threshold M_k a new model wherein only particles with a mass less than M_k contribute to the running coupling constant. At each threshold the β function is approximated by a step function and the running coupling constant is continuous. We obtain the formulae:

$$\sin^2\theta_W\,(M_W) = E^{-1}\,\{\,(B_2 - B_3) + \frac{\alpha C^2}{\alpha_3}(B_1 - B_2) - C^2((B_1 - B_2)\,A_{32} + (B_3 - B_2)\,A_{21})\,\} \quad ,$$

$$\frac{1}{2\pi}\ln\frac{M_G}{M_W} = E^{-1}\,(\frac{1}{\alpha} - \frac{1+C^2}{\alpha_3} - (C^2A_{13} + A_{23})) \quad ,$$

$$A_{ij} = (2\pi)^{-1}\sum_{k=1}^{N-1}((b_i^k - b_j^k) - (b_i^{k+1} - b_j^{k+1})) \quad , \quad b_i^N = B_i \quad .$$

b_i are the coefficients of the different β functions.

$E = ((B_2 - B_3) + C^2(B_1 - B_3))$.

The different thresholds are defined by the technicolor scale (1 TeV) and the masses of the low-lying bound-states. In all our numerical results, we use the input parameters: $\alpha_{QCD} = 0.12$ and $\alpha_{em} = (128.2)^{-1}$.

We consider two types of effective models:

Model A: The technifermions carry SU(5) quantum numbers, an example being the one family model of Dimopoulos[4]. The technifermions

are a copy of the ordinary fermions. Below the technicolor scale, the model is defined by the bound states (Table 1). In this model the technifermions contribute equally to each β function, therefore $\sin^2\theta_W$ and M_G are independent of the dimension of the technifermion representation.

Complex representation: $\sin^2\theta_W = 0.205$ $M_G = 1.5 \cdot 10^{15}$ GeV .

Real representation: $\sin^2\theta_W = 0.202$ $M_G = 5.0 \cdot 10^{15}$ GeV .

α_G depends on the dimension of the technifermion representation. A proton lifetime of 10^{30} years restricts the dimension to be smaller than 4 . This model is ruled out because of new proton decay channels (through scalars).

Model B: The technifermions carry only electroweak quantum numbers. - The one doublet model[4]. - If the technifermion representation is complex, the model has no Pgb's. In the case of real representation, the model has 6 Pgb's. The results depend on the dimension of the technifermion representations.

		d	3	4	5	6	
Complex representation:	$\sin^2\theta_W$	=	0.214,	0.220,	0.223,	0.225	
	M_G	=	3.3,	2.3,	1.6,	1.1	10^{14} GeV
Real representation:	$\sin^2\theta_W$	=	0.215,	0.221,	0.224,	0.227	
	M_G	=	2.6,	2.0,	1.3,	0.9	10^{14} GeV

In this model α_G is independent of d . It is in agreement with experimental data on $\sin^2\theta_W$ and the proton lifetime.

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Table 1

Complex representation:

SU(3)	$SU(2)_L$	U(1)	M	b_1	b_2	b_3
8 8	3 1	0 0	270 GeV	0	8/3	4
3 3 + their antiparticles	3 1	-2/3 -2/3	160 GeV	$32C^2/9$	2	4/3
1 1	3 1	0 0	M_W	0	1/3	0

Real representation: The 60 Pgb's of the complex representation plus

SU(3)	$SU(2)_L$	U(1)	M	b_1	b_2	b_3
6 + their antiparticles	3	1/3	270 GeV	$4C^2/3$	4	5
3 3 3 + their antiparticles	3 1 1	-1/3 -1/3 -1/3	160 GeV	$10C^2/9$	2	5/3
1 + their antiparticles	3	-1	M_W	$2C^2$	2/3	0

A REALISTIC QUARK-LEPTON MASS SPECTRUM IN A MODEL WITH EXTENDED TECHNICOLOR

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Abstract: In the frame of an SU(8) grand unified model incorporating an $SU(3)_{TC}$ technicolor subgroup we show that with certain assumptions concerning the technicolor dynamics a realistic fit of the quark-lepton mass spectrum can be achieved[1]. The minimal anomaly free fermion representation containing 3 families is

$$3[\bar{1}] + 2[2] + [\bar{3}] \quad . \tag{1}$$

Above [m] are the antisymmetric tensors of rank m which we denote by $\psi_A^{(k)}$, $\chi^{(k)AB}$, φ_{ABC}, with k being a family index.

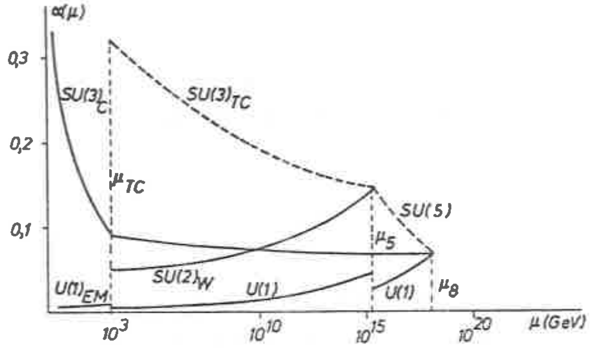

Fig. 1 shows the symmetry breaking scheme and the energy dependence of the running coupling constants for different subgroups.

Our main dynamical assumptions are:

1. The symmetry breaking at the scales μ_8 and μ_5 is due to elementary Higgs fields.
2. The extended technicolor (E.T.C.) mass scales μ_8 and μ_5 can be taken arbitrarily large due to a mechanism proposed by Holdom[2].
3. The global symmetry due to the repetition of irreducible representations in (1) is explicitly broken by Yukawa couplings to scalar fields with masses of $O(\mu_5)$:

$$L_Y = [A_1\, \chi^{(1)AB}\, \chi^{(1)CD} + A_2\, \chi^{(2)AB}\, \chi^{(2)CD}]\, H_{ABCD} + V\, \psi_A^{(1)}\, \varphi_{BCD}\, H^{*ABCD} \tag{2}$$

4. The technifermions condensates at a mass scale $\mu_{TC} \simeq 1$ TeV break

the Weinberg-Salam group. The subgroup alignment is determined by other perturbing interactions: $U(1)_{e.m.}$, $SU(3)_c$, and E.T.C. interactions.

Assuming that the E.T.C. interaction (2) dominates, one obtains in first approximation the following mass pattern: The t, c and b quarks and the τ lepton become massive with $m_\tau = m_b$. Also one obtains the relation $m_t > 2m_b/U_{ub}$ between t, b masses and U_{ub} Kobayashi-Maskawa matrix element[3] which implies $m_t > 70$ GeV. Other E.T.C. interactions induce small masses to u, s, d quarks and μ and e leptons. The mass relation $m_d \simeq m_s U_{ub}$ results in a natural way. Also the ν_τ neutrino becomes massive. $\nu_e - \nu_\mu$ mixing occurs only as a second-order effect.

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III. COMPOSITE MODELS OF LEPTONS AND QUARKS

ELECTROWEAK INTERACTIONS IN COMPOSITE MODELS -
PHENOMENOLOGICAL ASPECTS

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1. Introduction

Recently, the opinion has been advocated[1], that leptons and quarks may not be elementary but composed of more fundamental subconstituents ("preons"). The general properties of such bound states together with the necessary ingredients of the corresponding models will be described in the talks by Harari[2] and Barbieri[3]. Therefore I can restrict myself to some short introductory remarks concerning the main arguments for and features of the subconstituent picture.

Among the numerous arguments, which have been advocated in favour of composite leptons and quarks, the following two are most convincing:

a) The generation pattern among the "fundamental" fermions (leptons and quarks): We have at least three groups (generations) of fermions, which show similar behaviour under $SU(3)_c \times SU(2) \times U(1)$. They differ only by their masses and (mass-like) mixing angles. Futhermore, within each generation, quarks and leptons are very similar. Their electric charges cancel exactly. This generation pattern can be naturally explained by assuming that quarks and leptons are composed of subconstituents and that the more massive generations represent higher excitations.

b) The singular nature of weak interactions as compared to strong and electromagnetic interactions: Within the standard scheme ($SU(3) \times SU(2) \times U(1)$) and within grand unified extensions, the weak part of $SU(2) \times U(1)$ is the only spontaneously broken local gauge symmetry, all other local gauge symmetries being exact. On the other hand, all <u>global</u> symmetries seem to be broken (with the possible exception of $B-L$). Since global (gauge) symmetries have a rather vague physical significance (they only show up in selection rules, but their charges are not sources of a field), there is no convincing argument,why they should be strictly valid. On the other side,one would expect - at least from aesthetical reasons - that all local gauge symmetries are exact. Subconstituent models open the possibility for such a situation[4], as we shall discuss in a moment. There,the presently observed weak processes ought not to be interpreted as manifestations of a local SU(2)-gauge symmetry but are rather connected with a (broken) global symmetry. Within that picture,one can easily avoid the existence of fundamental (Higgs) scalars,which constitute another unwanted feature of the Glashow-Weinberg-Salam-model.

In the present talk I will exclusively dwell on the second aspect of subconstituent models, mainly because there is some progress on that field,whereas no really convincing model has been constructed,which could describe the generation pattern in terms of the underlying subconstituent dynamics[5].

2. General Dynamical Features

Very little is known about the forces which bind the subconstituents inside leptons and quarks. Therefore people have tried to look for ingredients of a theory,which are suggested either by comparison with phenomenology or by our knowledge about QCD. The following features are usually believed:

- Subconstituents are bound by a superstrong force, based on some unbroken nonabelian local gauge symmetry. The corresponding gauge group G_H (which for simplicity is assumed to be some $SU(n)$ or a product of SU(n)-groups) acts on an internal degree of freedom called hypercolor. The theory is asymptotically free and shows confinement at some scale. Since the confinement radius has to be smaller than 10^{-17}cm ,one expects that the mass scale of the hypercolor dynamics Λ_H is larger than 300 GeV .

Physical particles (quarks, leptons, intermediate vector bosons, eventually Higgs mesons) are hypercolor singlets. Some of these singlet states (leptons, quarks) are light ($m << \Lambda_H$), others may have masses of the order of Λ_H (f.i. intermediate vector bosons or Higgses). A convenient way to explain the existence of almost massless fermionic states is provided by the assumption that chiral symmetry, which is present on the subconstituent level, (partially) survives the confinement mechanism and thereby guarantees massless fermions[6]. On the other hand, it is not clear, how such simple arguments can account for the gigantic range of the physical lepton and quark masses, which extends over 9 orders of magnitudes.

Between G_H-singlets a residual hypercolor force (in addition to electromagnetic and color forces) is acting much in the same way as there are residual color interactions between color-singlet hadrons. At low q^2 (i.e. for distances much larger than Λ_H^{-1}), these van der Waals like forces represent weak interactions.

If intermediate vector bosons exist, they are also composed of subconstituents; the same is true eventually for Higgs-mesons. On the other hand, gluons and the photon are considered to be elementary. Therefore, the full local gauge group contains at least $G_H \times SU(3)_C \times U(1)_{EM}$ as a subgroup.

The dynamics of physical particles at length scales much larger than Λ_H^{-1} (where the internal structure of quarks and leptons is ineffective) can be described by an effective Lagrangian containing the G_H-singlet fields as (quasi-) local quantities. This Lagrangian includes all ordinary QED- and QCD-terms; additional unusual terms (anomalous magnetic moments of leptons and quarks[7]), flavour-changing terms[8], baryon number violating[8] terms) are of higher dimension and should consequently be suppressed by appropriate powers of Λ_H^{-1}. We will neglect those contributions in the following.

Having now sketched the general frame for the subconstituent dynamics, we go over to a more detailed discussion of the expected structure of weak interactions within that scheme.

As we have stated already, weak interactions, as we observe them today, are interpreted as the low q^2 $(q^2 \ll \Lambda_H^2)$ residual hypercolor interactions between (hypercolor singlet) leptons and/or quarks. An important open question concerns the relative magnitude of this residual force scale (given approximately by $G_F^{-1/2} \simeq 300$ GeV or by m_W) and the underlying hypercolor force scale Λ_H . Although it may well be that $\Lambda_H \gg m_W$, we will take in the following the extreme attitude that $\Lambda_H \sim O(G_F^{-1/2}) \sim O(m_W)$, since that assumption will yield a very simple and appealing explanation for the weak interaction scale.

The residual force may be described by the exchange of (G_H-singlet, composite) intermediate vector bosons W^+, W^-, Z^0 . A typical case is depicted in Fig. 1.

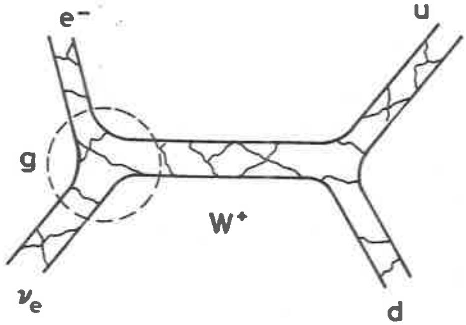


Figure 1

The low q^2 behaviour of the diagram in Fig. 1 is given by g^2/m_W^2 , which due to phenomenological reasons has to be equal to $8G_F/\sqrt{2}$:

$$\frac{g^2}{m_W^2} = \frac{8G_F}{\sqrt{2}} \quad . \tag{1a}$$

From this we get

$$m_W = \sqrt{\frac{g^2}{8G_F/\sqrt{2}}} = g \cdot 123 \text{ GeV} \quad . \tag{1b}$$

Since the coupling of the W's to leptons or quarks is of the type of a strong coupling, the corresponding coupling constants are not bound to be small, but rather may be of order one. If we tentatively put $g \gtrsim 1$ in eq. (1), we would get

$$m_W \gtrsim 123 \text{ GeV} \quad .$$

Thus, there is no a priori reason for the intermediate boson masses to coincide with the values predicted by the G-W-S-model. On the other hand those specific values are not at all ruled out and we will show later, how they appear naturally as a consequence of rather general dynamical assumptions.

To account for the phenomenology of charged current processes, one has to assume that the effective Lagrangian (of composite fields) shows (at least when electromagnetic interactions are switched off) a global $SU(2)_W$ symmetry ("weak isospin") with the usual doublet assignment

$$\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L , \ldots \begin{pmatrix} u \\ d_c \end{pmatrix}_L , \ldots$$

Within most of the models, which are presently on the market, this global SU(2) symmetry is already established on the fundamental level. Then, in addition to W^+, W^- and $W^{(3)}$, which are members of a SU(2)-triplet, there should also exist an electrically neutral SU(2)-singlet vector boson $W^{(0)}$ coupling to some isoscalar current. Since the isoscalar contribution to the (axial) neutral current is measured to be small (if existing at all), we have to assume that the mass of the $W^{(0)}$ is much higher than the W-masses.

Since all intermediate vector bosons are compound objects, we expect that in addition to the lowest lying ones $(\vec{W}, W^{(0)})$ there should also exist higher excited states $(\vec{W}', W^{(0)\prime}, \vec{W}'', \ldots)$, which again couple to quarks and leptons. On the other hand, the excitation energy should be of the order of Λ_H, therefore these excited states will lie fairly high in mass $(m_{W'} \geq 300 \text{ GeV})$ and their contribution to the residual interactions should be small.

The photon couples directly and pointlike to the subconstituents, but since leptons and quarks are composite states, their coupling to the photon is of a more complicated nature (like the hadron photon coupling) and has to be described by formfactors. We expect that the q^2-dependence of those lepton- or quark-formfactors is somehow determined by the hypercolor scale Λ_H (which is near to m_W in our case). In addition, as long as $W^{(3)}$ and $W^{(0)}$ are composed of charged subconstituents, there should be a transition coupling of the photon both to $W^{(3)}$ and to $W^{(0)}$ (see Fig. 2).

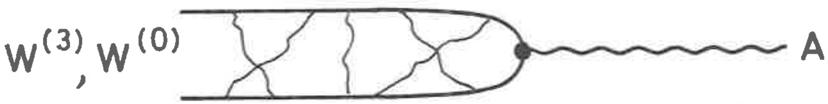
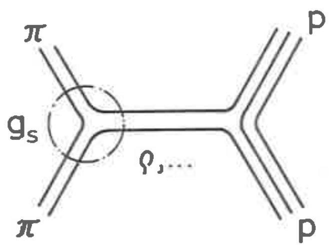
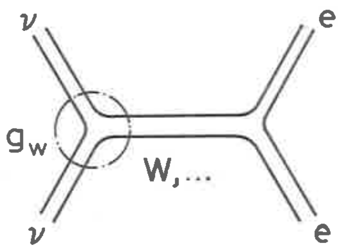


Figure 2

This generates a mixing mechanism between γ, $W^{(3)}$ and $W^{(0)}$, which will be important for the weak neutral current structure.

Having summarized the main expectations as to the weak interactions within subconstituent models, we recognize a striking parallelity between hypercolordynamics and weak interactions on the one hand, and colordynamics and strong hadronic interactions on the other, which is designated in the following schematic way:

Colordynamics (QCD)	Hypercolordynamics (QHD)
<u>Strong</u> interactions between color singlet hadrons are manifestations of residual QCD-interactions. For example:	<u>Weak</u> interactions between hypercolor singlet leptons and quarks are manifestations of residual QHD-interactions. For example:
	

Colordynamics (QCD)	Hypercolordynamics (QHD)
Effective couplings of the exchanged bosons are invariant under strong isospin (SU(2)).	Effective couplings of the exchanged bosons are invariant under global weak isospin ($SU(2)_W$).
There are SU(2)-triplet ($\rho^{\pm}$, ρ^{0}) and SU(2)-singlet (ω) contributions.	There are $SU(2)_W$-triplet ($W^{\pm}$, Z_0) and $SU(2)_W$-singlet ($W^{(0)}$) contributions.
Isospin is broken by $\rho^0-\gamma$ and $\omega-\gamma$ mixing.	$SU(2)_W$ is broken by $W^{(3)}-\gamma$ and $W^{(0)}-\gamma$ mixing.
The energy scale of QCD is in the GeV-region.	The energy scale of QHD is in the TeV-region.

We should add a short historical remark here. In 1960, J.J. Sakurai had proposed a non-Abelian gauge model for strong hadronic interactions with ρ and ω (yet undetected at that time) playing the role of gauge vector bosons. Later, it became clear that hadronic interactions should only be viewed as an indirect manifestation of the interquark forces, which themselves are of a gauge type (QCD). Twenty years later, a similar development seems to occur with weak interactions. Whereas Glashow, Weinberg and Salam have tried to explain weak interactions as generated by a local gauge theory, it may happen, that they also are the residue of a local gauge interaction acting on a deeper lying (subconstituent) level.

In spite of the numerous parallelities between hadronic and weak interactions one should not overlook the drastic differences between both. History does not really repeat itself! Among the apparent new (unorthodox) features of the residual (weak) interactions between leptons and quarks are:

- Parity-non-invariance of the effective Lagrangian: Where does parity violation come from? One could either imagine that parity is violated genuinly[10] (i.e. that left- and righthanded subconstituents transform differently under G_H),or that the fundamental Lagrangian is left-right-symmetric with parity being spontaneously broken (f.i. by right- and lefthanded fermion-fermion condensates with different vacuum expectation values).

V-A-structure: Weak interactions seem to be exclusively described by vector and axialvector exchange. Why are there no (pseudo-) scalar or tensor exchanges,which on the other hand contribute so drastically in hadronic interactions? One could imagine several possible answers to this question. Maybe that the vector channel is much more attractive than scalar (S) of tensor (T) channels within QHD. It is also possible that there are cancellations between S- and T-contribution and/or between different parity contributions (parity is not a good quantum number, after all). The most simple explanation would be, of course, that the scalar and tensor bosons have much higher masses than the vector bosons, or that their fermionic couplings go proportional to the fermion masses and are therefore small.

Isovector nature of the (axial) neutral current: Why is the contribution of the $W^{(0)}$ so small (if existing at all) as compared to the isovector ($W^{(3)}$) contribution? If that stems from a high $W^{(0)}$-mass: Why is it such high? A possible explanation could come from a comparison with the π-η-system, where the isoscalar state (mainly η) is also approximately three times more massive than the pion. Maybe that this comes from a large (hyper-) gluon annihilation part within the isoscalar bound state.

We see that the whole picture has to contain a lot of novel features,in order to provide for a correct description of weak interactions. And it will be a big challenge to construct specific models, which can account for those new requirements.

3. Mixing and W-Dominance[11]

The expected features of weak interactions which we have described before, may seem convincing - but do they really provide us a framework for the presently experienced weak phenomenology, which has been so well reproduced by the Glashow-Weinberg-Salam model?

In answering this question we fortunately can use the results of a small group of physicists[12], who - contrary to the overwhelming majority - have looked for alternatives to the G-W-S-description. Their works showed, that both charged current and neutral current phenomenology (at low q^2) can be successfully described on the basis of global SU(2) invariance of lefthanded current couplings, once mixing between the photon and the neutral carriers of weak interactions (Z_0,...) is assumed to occur. Since global SU(2)-symmetry is a (natural) assumption within, and γ-Z_0-mixing is even a necessary consequence of subconstituent models, we can hope that the main features of weak interactions will be reproduced correctly within the scheme proposed before.

a) General Mixing Effects

Let us shortly show how mixing accounts for the structure of neutral current phenomenology. For simplicity, we thereby take into account only one (isovector) neutral intermediate vector boson, which we call $W^{(3)}$ (which after diagonalization will show up as the physical particle Z_0). The mixing between the fields $W_\mu^{(3)}$ and A'_μ (which is the primordial photon before diagonalization) can be described by the following term in the effective Lagrangian[13]:

$$L_{mix} = -\frac{\lambda_W}{4}\left(W_{\mu\nu}^{(3)}\, F'^{\mu\nu} + F'_{\mu\nu}\, W^{(3)\mu\nu}\right) \tag{2}$$

$$\left(W_{\mu\nu}^{(3)} = \partial_\mu W_\nu^{(3)} - \partial_\nu W_\mu^{(3)} \quad ; \quad F'_{\mu\nu} = \partial_\mu A'_\nu - \partial_\nu A'_\mu\right)$$

which yields for each $W^{(3)} - A'_\mu$ vertex a contribution $\lambda_W q^2$ (see Fig.3).

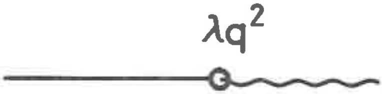


Figure 3

Such a transition coupling has the following consequences:

- It changes the $W^{(3)}$-propagator in the manner indicated in Fig. 4.

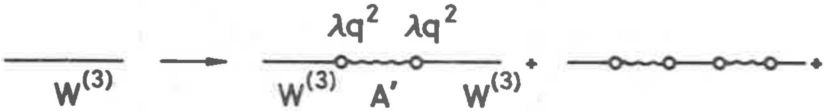


Figure 4

In formulae this reads like

$$\frac{-1}{q^2 - m_W^2} \to \frac{-1}{q^2 - m_W^2} \left[1 + \frac{1}{q^2} \frac{(\lambda q^2)^2}{q^2 - m_W^2} + \ldots\right]$$

$$= \frac{-1}{q^2 - m_W^2} \frac{1}{1 - \frac{\lambda^2 q^2}{q^2 - m_W^2}} = \frac{-1}{q^2(1-\lambda^2) - m_W^2} = \frac{-1/1-\lambda^2}{q^2 - \frac{m_W^2}{1-\lambda^2}} \quad .$$

The full propagator has a pole at $q^2 = m_W^2 / 1 - \lambda^2$, which defines the mass of the physical Z_0-particle, M_W :

$$M_W^2 = \frac{m_W^2}{1-\lambda^2} \geq m_W^2 \quad . \tag{3}$$

The mass of the Z_0 is thus shifted by the mixing mechanism to a value higher than the primordial mass m_W (which is equal to the mass of the charged intermediate boson $W^\pm$ because of the SU(2)-symmetry of the non-mixing parts of the effective Lagrangian). Of course, the exact value of M_W is not predictable as long as λ is unspecified. But we can derive a bound on λ from (3)[14]:

$$\lambda^2 < 1 \quad .$$

- A further consequence of mixing is, that NC neutrino processes are in lowest order (of weak coupling and mixing) described by two diagrams (see Fig. 5) instead of single Z_0-exchange.

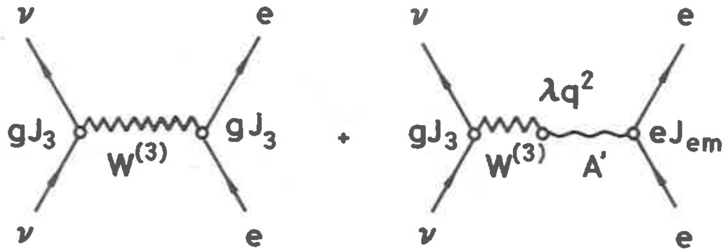


Figure 5

The resulting matrixelement therefore takes the form (at low q^2):

$$T_{fi} = g\, J_3^{(\nu)} \frac{-1}{q^2 - m_W^2}\, g\, J_3^{(e)} + g\, J_3^{(\nu)} \frac{-1}{q^2 - m_W^2}\, \lambda q^2 \frac{-1}{q^2}\, e\, J_{em}^{(e)}$$

$$= g\, J_3^{(\nu)} \frac{-1}{q^2 - m_W^2}\, g\, (J_3^{(e)} - \lambda \frac{e}{g} J_{em}^{(e)})$$

$$\xrightarrow[q^2 \to 0]{} \frac{g^2}{m_W^2} J_3^{(\nu)} (J_3^{(e)} - \lambda \frac{e}{g} J_{em}^{(e)}) \quad .$$

This is equal to the Weinberg-Salam result for the same quantity, if

$$\frac{g}{m_W^2} = \frac{8G}{\sqrt{2}} \quad , \tag{5}$$

and if

$$\lambda \frac{e}{g} = \sin^2\theta_W \quad . \tag{6}$$

The detailed treatment done in ref. 11) shows that these conditions remain true when mixing is taken into accout to all orders.

- Eqs. (5) and (6) imply a connection between the charged boson mass m_W and λ :

$$m_W = \sqrt{\frac{\pi\alpha}{\sqrt{2}\, G_F}}\; \frac{1}{\sin\theta_W}\; \frac{\lambda}{\sin\theta_W} \quad . \tag{7}$$

The first part of this product is equal to the prediction for the $W^\pm$-

mass within the G-W-S-model:

$$m_W^{(WS)} = \sqrt{\frac{\pi\alpha}{\sqrt{2}\,G_F}}\;\frac{1}{\sin\theta_W} = 37.3\text{ GeV}\,\frac{1}{\sin\theta_W} = 77.9\text{ GeV} \tag{8}$$

where a value of $\sin^2\theta_W = 0.23$ has been used in (8). Therefore (7) can be written as

$$m_W = m_W^{(WS)}\,\frac{\lambda}{\sin\theta_W} = 162\text{ GeV}\cdot\lambda \quad . \tag{9}$$

Together with (4), that yields the upper bound[14)]

$$m_W \leq 162\text{ GeV} \quad . \tag{10}$$

As a result of this schematic view one realizes that mixing can account for the rough features of NC phenomenology, although mixing alone (with g and λ unspecified) does not specify the values for the W-masses. On the other hand, the mass values predicted by the G-W-S-model can be easily reproduced within the mixing scheme, if appropriate values of g and λ are postulated. In fact, the so-called "unification condition"[12)]

$$\lambda = \frac{e}{g}$$

is sufficient to guarantee W-S-mass values. But this condition does not follow from the pure mixing postulat. Therefore, there is no convincing reason to assume its validity at the present state. We will see that (11) is a natural consequence within the subconstituent scheme.

b) Effects of Compositeness, W-Dominance

Now, let us concentrate our attention to the physics of composite quarks and leptons. Mixing is automatically guaranteed there, but in a more extendend version, namely mixing of the photon with $W^{(3)}$ and with $W^{(0)}$. The latter is described by a Lagrangian analogous to (2) but with a different mixing coupling λ_Y .

What will radically change now, is the structure of the photonic coupling of the fermions, since they have an internal structure, thus, these couplings have to be described by means of form factors. Without detailed knowledge of the hypercolordynamics we are not able to calculate the leptons- and/or quark-form factors. In such a situation it is tempting to adopt an approach, which has been successful in a similar physical situation: We know from the electromagnetic interactions of ordinary hadrons that hadronic form factors can be described by saturating the hadron-photon vertex by those vector mesons, which couple to the hadrons in a Zweig-allowed way (ρ and ω for the nucleon form factor). In the same spirit we treat here the coupling of the photon to leptons and quarks by saturating the vertex function by those vector-bosons, which couple directly to the fermions, i.e. with $W^{(3)}$, $W^{(0)}$, $W^{(3)\prime}$,... Graphically that is pictured in Fig. 6.

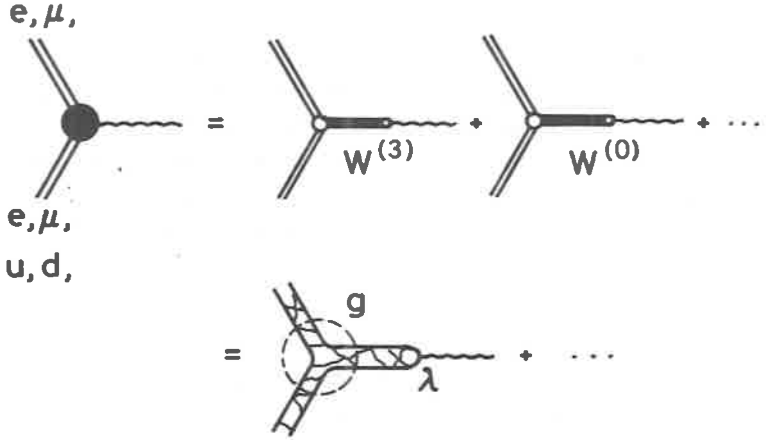


Figure 6

As a consequence, there is no direct coupling of the electromagnetic field to leptons or quarks. On the other hand, the mixing Lagrangian (2) alone cannot represent the whole amount of photonic interactions, since it vanishes for real $(q^2 = 0)$ photons. Due to the work of Kroll, Lee and Zumino[15] we know how the Lagrangian has to be completed in order to yield the correct electromagnetic interaction within the vector dominance picture. The corresponding full Lagrangian for our case (two vector bosons) can be found in ref. 11)[16]. In order to utilize this Lagrangian for calculating weak interaction matrix elements one could now start to diagonalize the theory (i.e. its inverse propagator matrix). This has been done in ref. 11), and the result of the exact diagonalization procedure is discussed there. For our demonstrative purpose within that talk, we adopt a more transparent way of diagonalization by using a perturbative approach. Thereby we treat the weak couplings

$$g_W \vec{J}_\mu \vec{W}^\mu + g_Y J_{o\mu} W^{(o)\mu} \tag{12}$$

in first order, but the electromagnetic mixing couplings λ_W and λ_Y to all orders. This is necessary since we will (a posteriori) find rather large values for the corresponding coupling constants.

In order to simplify the situation and to obtain more specific predictions, we neglect, in the first instance, the higher excited intermediate vector bosons and take into account only the lowest lying ones ($\vec{W}$ and $W^{(o)}$). Since the corresponding mass splitting should be of the order of Λ_H (i.e. at least 300 GeV), that may not be a too bad approximation for describing low energy pheonomena. But we cannot be sure about it (there might be excitations of different sort, like those which constitute the higher fermionic generations) and therefore, we will discuss the corrections, steming from higher states, in some detail later.

For the moment, we stick to three isovector bosons $(W^+, W^-, W^{(3)})$ and one isoscalar boson $(W^{(o)})$. A typical lowest order charged current matrix element is depicted in Fig. 7.

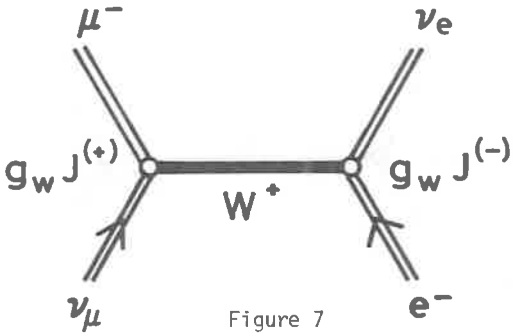


Figure 7

The corresponding effective four-fermion Lagrangian takes the form

$$-2L^{cc}_{eff} = g_W^2 \frac{-1}{q^2 - m_W^2} J_\mu^{(+)} J^{(-)\mu} \xrightarrow[q^2 \to 0]{} \frac{g_W^2}{m_W^2} J_\mu^{(+)} J^{(-)\mu} \quad . \tag{13}$$

Comparison with low q^2 phenomenology yields the already well known calibration relation

$$\frac{g_W^2}{m_W^2} = \frac{8G_F}{\sqrt{2}} \quad \text{with} \quad G_F = 1.166 \times 10^{-5}\ \text{GeV}^{-2} \quad . \tag{14}$$

As a next step we consider a typical neutral current process. The contributing diagrams (in the indicated order of perturbation expansion) are shown in Fig. 8.

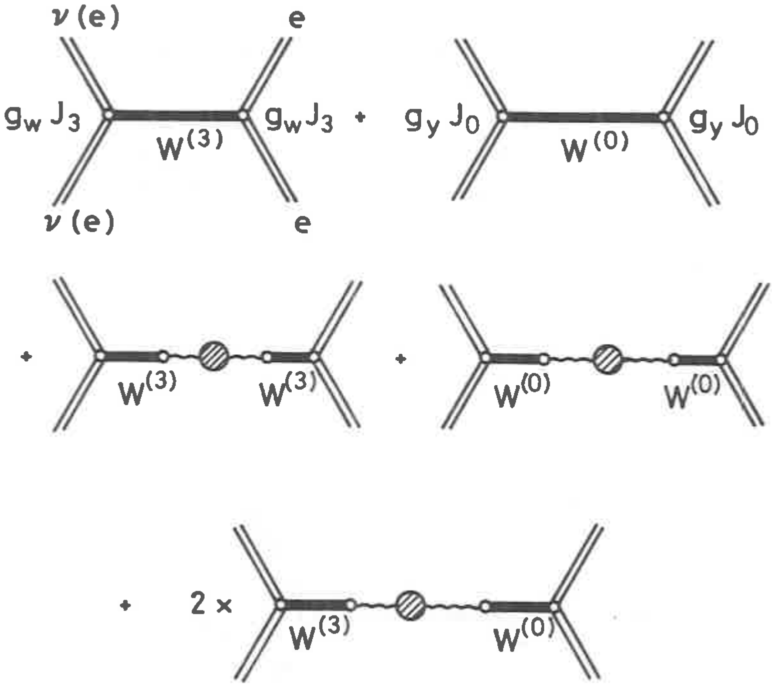


Figure 8

The shaded blob denotes the full (all orders of mixing) propagator of the photon $(D(q^2))$, which will be discussed in a minute. The diagrams in Fig. 8 correspond to an effective four fermion interaction Lagrangian (for weak neutral current and electromagnetic interactions) of the following form[17)]

$$2L_{eff}^{NC+EM} = J_3 \frac{g_W^2}{q^2 - m_W^2} J_3 + J_0 \frac{g_Y^2}{q^2 - m_Y^2} J_0$$
$$-(J_3\, g_W\, \lambda_W \frac{m_W^2}{q^2 - m_W^2} + J_0\, g_Y\, \lambda_Y \frac{m_Y^2}{q^2 - m_Y^2})^2\, D(q^2) \quad . \qquad (15)$$

An expression for the full photon propagator can also be found by diagrammatic means. Fig. 9 indicates the series of contributing diagrams.

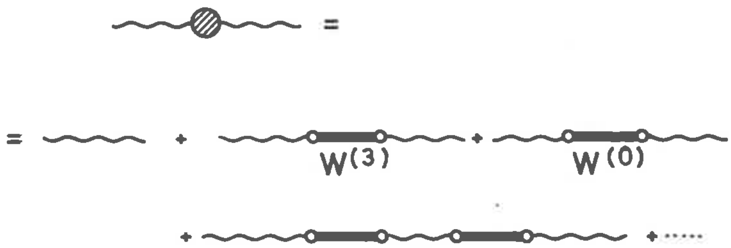


Figure 9

We find

$$D(q^2) = \frac{-1}{q^2} (1 + \sum_{i=W,Y} \frac{\lambda_i^2\, q^2}{q^2 - m_i^2} + (\sum_i \frac{\lambda_i^2\, q^2}{q^2 - m_i^2})^2 + \ldots)$$
$$= -\frac{1}{q^2} \frac{1}{1 - \sum_i \frac{\lambda_i^2\, q^2}{q^2 - m_i^2}} \quad . \qquad (16)$$

The poles of $D(q^2)$ represent the physical (diagonalized) particles. One pole evidently lies at $q^2 = 0$, which corresponds to the physical photon. The locations of the other two poles define the masses of the (physical) isovector and isoscalar bosons Z_0 and Y, respectively. The corresponding mass values are given by rather complicated expressions which can be found in ref. 11. A simple approximation formula is obtained in the following way: The poles of $D(q^2)$ are determined by the zeros of

$$1 - \Sigma \frac{\lambda_i^2 \ q^2}{q^2 - m_i^2} \quad .$$

For those values of q^2, which lie near to a pole $(q^2 \simeq M_i)$, we can neglect the second part of the sum, because the physical masses M_i will be near to the bare masses m_i, as long as the mixing couplings λ_i are small. We can therefore approximately write

$$1 - \frac{\lambda_i^2 \ M_i^2}{M_i^2 - m_i^2} \simeq 0 \quad , \qquad i = W \ \text{ or } \ Y \quad ,$$

which gives

$$M_i^2 \simeq \frac{m_i^2}{1 - \lambda_i^2} \qquad i = W,Y \quad . \tag{17}$$

Formula (17) is correct up to terms of order λ_i^4.

Now, let us analyse in more detail the content of the effective Lagrangian (15). For comparison with weak interaction phenomenology we need the form of L_{eff}^{NC+EM} at low q^2. Therefore we expand everything in powers of q^2. For $D(q^2)$, which has a pole at $q^2 = 0$, we get

$$D(q^2) = -\frac{1}{q^2} + \sum_i \frac{\lambda_i^2}{m_i^2} + O(q^2) \quad . \tag{18}$$

Consequently, the Lagrangian at $q^2 \sim 0$ contains terms proportional to $1/q^2$ and terms of order $(q^2)^0$. The contribution of the former is

$$2L_{eff}^{NC+EM} = \frac{1}{q^2} \left(J_3 \ g_W \ \lambda_W + J_0 \ g_Y \ \lambda_Y\right)^2 + O((q^2)^0) \quad . \tag{19}$$

This expression has to represent the photon exchange contribution, therefore $J_3 g_W \lambda_W + J_0 g_Y \lambda_Y$ must be equal to eJ_{em}, which implies

$$g_W \ \lambda_W = e = g_Y \ \lambda_Y \quad . \tag{20}$$

We call this relation "saturation condition", because it follows directly from saturating the electromagnetic form factors (of leptons and quarks) by the lowest lying vector mesons. The first part of (20) is identical to the unification condition within general mixing schemes,[12] where it had to be put in by hand, in order to establish the W-S mass relations. Here it is a consequence of the subconstituent dynamics.[18)]

Using (20) and taking into account also the terms $O((q^2)^0)$ we obtain the following low-q^2-form of the effective Lagrangian

$$2\, L_{eff}^{NC+EM} = \frac{e^2}{q^2} J_{em}^2 - \frac{g_W^2}{m_W^2} [(J_3 - \lambda_W^2 J_{em})^2 + + \frac{m_W^2}{m_Y^2} \frac{g_Y^2}{g_W^2} (J_Y - \lambda_Y^2 J_{em})^2] + O(q^2) \quad . \tag{21}$$

The first part within square brackets represents the W-S-form of the NC-interactions. The second part gives an (weak) isoscalar contribution, which according to experimental indications has to be less than 10% of the isovector one. This suppression is guaranteed if

$$m_Y \geq 3 m_W \quad .$$

On the basis of this assumption we can neglect the second term in (21). The remainder of (21) reproduces correctly the low-q^2-behaviour of the W-S-neutral current Lagrangian if we put

$$\lambda_W^2 \equiv \sin^2\theta_W \simeq 0.23 \quad . \tag{22}$$

Note that due to eq. (14) we also get the correct normalization of the neutral current relative to the charged current Lagrangian $(\rho = 1)$. Equations (22) and (20) imply

$$\lambda_W = \frac{e}{g_W} = \sin\theta_W \quad . \tag{23}$$

Since (23) represents the unification condition, we should not be surprised that it also yields the Weinberg mass formula. In fact, (9) together with (23) gives exactly

$$m_W = m^{(WS)} \frac{\lambda}{\sin\theta_W} = m^{(WS)} = 78 \text{ GeV} \quad . \tag{24}$$

What about the Z_o-mass? According to (17) we get

$$m_{Z_o} = M_W^2 = \frac{1}{1 - \lambda_W^2} m_W^2 + O(\lambda_W^4, \lambda_Y^4) \quad .$$

Together with (22),this yields

$$m^2_{Z_o} \simeq \frac{m^2_W}{\cos^2\theta_W} \quad , \tag{25}$$

which is (in this approximation) identical to Weinberg's mass relation.

In summarizing,we see that within a picture of composite quarks and leptons one should not only expect the correct low-q^2-structure of CC- and NC-interactions but also mass values for $W^\pm$ and Z_o, which are very near to the canonical values.

c) Extendend W-Dominance

The main dynamical input, which leads to that rather unexpected behaviour, is the assumption of single-W-dominance of the form factors. Therefore, this assumption has been the subject of some debate, which concentrated on the question of a possibly large non-resonance contribution. It has been argued[19)] on the basis of asymptotic freedom QCD sum rules[20)] that it is only for those models, where the sum of the squared electric charges of subconstituents exceeds that of the quarks and leptons, that one can approximately neglect the continuum contribution. On the other hand, within the spirit of (generalized) vector dominance (which has not yet been derived from QCD) one would expect that continuum effects are accounted for by including higher excited vector mesons. Therefore, it seems interesting to study the influence of higher excitations (which have been neglected up to now) on our predictions.

A first approach into this direction has been done by Kuroda and Schildknecht[21)]. They take into account a series of vector meson resonances $(W_i^{(3)}, W_i^{(0)}, i = 1,2,\ldots)$. Then, eqs. (14) and (20) generalize to

$$\frac{8G_F}{\sqrt{2}} = \sum_i \frac{g^2_{W_i}}{m^2_{W_i}} \equiv G_W \tag{26}$$

and

$$\sum_i \lambda_{W_i} g_{W_i} = e = \sum_i \lambda_{\gamma_i} g_{\gamma_i} \quad , \tag{27}$$

respectively, and the low-q^2-behaviour of the effective four fermion weak-NC-Lagrangian can be written as

$$2\, L_{eff}^{NC+EM} = \frac{e^2}{q^2} J_{em}^2 - \rho \frac{8G_F}{\sqrt{2}} \left((J_3 - \sin^2\theta_W\, J_{em})^2 + C \cdot J_{em}^2 \right) \tag{28}$$

where

$$\rho = 1 + \frac{\Sigma \frac{g_{Y_i}^2}{m_{Y_i}^2}}{\Sigma \frac{g_{W_i}^2}{m_{W_i}^2}} \quad , \tag{29}$$

and $\sin^2\theta_W$ and C are complicated expressions of bare masses and coupling constants. (Note that $C = 0$ if only lowest lying vector mesons are contributing.) Thus, again, NC-phenomenology is well reproduced, if ρ, C, $\sin^2\theta_W$ are compatible with experimental results, which indicate

$$\sin^2\theta_W = 0.23 \tag{30}$$

$$\rho = 1 \pm 0.03 \qquad\qquad C \leq 0.015 \quad .$$

It is clear that there are too many parameters now, in order to fix the values of m_W and m_Z uniquely. To investigate the accessible regions for those quantities, the authors of ref. 21) took into account only the first excitation of $W^{(3)}$ $(W^{(3)\prime})$ thereby assuming that isoscalar excited states are very massive (which must be the case in order to keep ρ sufficiently near to 1). Furthermore, they assume that $\lambda_Y \simeq \lambda_W$, and that $\lambda_{W'}$ and λ_W are connected by the duality relation $\lambda_{W'} / \lambda_W = m_W / m_{W'}$. This reduces the number of unknown masses and coupling constants to 5 $(m_W, \rho_W, \lambda_W, m_{W'}, \rho_{W'})$, which are connected by the three relations (26), (27) and (30). Consequently two parameters (f.i. m_W and $r = m_{W'} / m_W$) remain undetermined. There are nevertheless restrictive bounds on the possible combinations of m_W and r, either steming from the required positivity of $1 - \lambda_W^2 - \lambda_{W'}^2$ or from the experimental bound on C.

The resulting allowed regions for m_W and r are depicted in Fig. 10. Also included there is a bound steming from the additional dynamical hypothesis that the (quasi-hadronic) coupling of W to fermions is larger than the W'-coupling. We see that m_W is expected to lie between 65 and 105 GeV (from $C \leq 0.01$), or even between 68 and 94 GeV, an absolute upper bound being

$$m_W \leq 135 \text{ GeV} \quad . \tag{31}$$

Thus, the single W-dominance prediction seems to be rather stable against contributions from higher excited vector bosons.

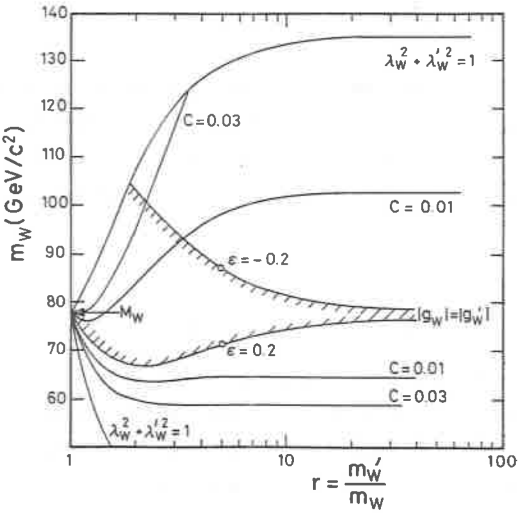


Figure 10

Let us finally add two remarks as to the consistency of the whole picture.

- We have mentioned that the electromagnetic mixing coupling λ_W^2 is numerically (approximately, in the case of higher excitations) equal to $\sin^2\theta_W$, i.e.

$$\lambda_W^2 \simeq 0.23 \quad .$$

This value seems to be astonishingly large for an electromagnetic coupling constant, especially if one compares it with the analogous ρ-γ-coupling

$$\lambda_\rho^2 \simeq 0.008 \quad .$$

Therefore, an independent additional information on the size of λ_W would be welcome. It comes from duality sum rule[22], which connects the high-q^2-behaviour of the vacuum polarization (governed by the charge square Q_S^2 of the subconstituents) with its low energy behaviour. Once the latter is saturated by the lowest lying vector boson $(W^{(3)})$ one gets

$$\lambda_W^2 = \frac{\alpha}{3\pi} N_S \, Q_S^2 \, \frac{\Delta m^2}{M_W^2} \quad , \tag{32}$$

where N_S denotes the number of (color and hypercolor) degrees of freedom of the subsonstituents. For instance, in the case of the so-called haplon model[23] ($N_S = 9$ and $Q_S^2 = 1/2$) , λ_W^2 can be as large as 0.23 provided that the mass-splitting Δm^2 is of the order of 0.4 TeV . That would imply $m_{W'} \simeq 600$ GeV , which is inside the expected range, if $\Lambda_H \simeq 300$ GeV . We realize that due to the high mass scale of hypercolor dynamics, λ_W can easily take a value much higher than expected from low energy hadronic phenomenology. Of course, there again rises the question of possible continuum contributions to the sum rule which has been investigated (within QCD-sum rules) in ref. 19).

- Universality. A crucial problem in all subconstituent models is connected with the fact that different generations (of leptons, at least) behave identically under weak interaction. Within the previously described picture that would imply

$$g_{Wee} = g_{W\mu\mu} = g_{W\tau\tau} \tag{33}$$

and similar relations for the g_Y's . Why should these couplings be equal, though e, μ, τ (and u, c, t) are states with different preon contents or with different dynamics?

Recently, an answer to this question has been proposed[24], which utilizes once more the above described similarity between hadronic and weak interactions. The argument roughly goes like follows. We know from ordinary hadronic phenomenology that the ρ couples with the same strength to all hadronic states of equal isospin (f.i. to K mesons, D mesons, B mesons). We also know that this follows from the assumption of local (strong) SU(2)-current algebra (which is realized by the quark currents $J_\mu(x) = \bar{q}(x)\,\gamma_\mu \frac{\vec{\tau}}{2}\,q(x)$, $q = \binom{u}{d}$) and from ρ-dominance of the (strong isovector) form factors.

Let us transfer this idea to the case of composite leptons and quarks. We assume that the weak-isospin charges, which generate $SU(2)_W$, can be constructed as integrals over local charge densities $\vec{\mathfrak{F}}^W_\mu(x)$:

$$\vec{I}_W = \int d^3\vec{x}\, \vec{\mathfrak{F}}^W_0(x)$$

and that those local quantities obey an equal time current algebra

$$[\mathfrak{F}^W_{0,i}(x), \mathfrak{F}^W_{0,j}(y)]_{x_0=y_0} = i\, \varepsilon_{ijk} \mathfrak{F}^W_{0,k}(x)\, \delta^3(\vec{x}-\vec{y}) \quad . \tag{34}$$

This assumption would be a trivial one, if leptons and quarks would be point-like and elementary (then the local weak currents would be the usual bilinears in the quark and lepton fields), but it becomes highly nontrivial for composite fields. Nevertheless, there might be the possibility that $SU(2)_W$ is already established on the subconstituent level,and then the weak local current algebra could be realized by means of currents which are bilinear in subconstituent fields. This is, e.g., the case in the so-called haplon model. Once (34) is assumed,we sandwich the equivalent version of the commutation relation (34)

$$[\mathfrak{F}^W_{0,1+i2}(x), \mathfrak{F}^W_{0,1-i2}(y)]_{x_0=y_0} = 2\mathfrak{F}^W_{0,3}(x)\, \delta^3(\vec{x}-\vec{y}) \tag{34'}$$

by isodoublet states $|A\rangle$ (all lefthanded leptons and quarks are members of isodoublets!) and integrate over $d^3\vec{y}$. The right hand side can be expressed by means of a weak form factor $F_A(q^2)$, and we get

$$F_A(0) = \langle I^W_3 \rangle_A \quad . \tag{35}$$

Thus, the form factors (at $q^2 = 0$) of all lefthanded doublet states are normalized universally.

Within the spirit of the fundamental assumption about weak interactions it is suggestive to assume that the spectral function for the weak $SU(2)_W$-currents can - at low q^2 - be saturated by the lowest-lying vector boson pole (W) . Then we get

$$F_A(q^2) = \frac{m^2_W}{f_W}\, g_{WAA} \frac{1}{m^2_W - q^2} \tag{36}$$

where f_W is the W-decay constant. For $q^2 = 0$, (35) and (36) imply

$$g_{WAA} = f_W \langle I_3^W \rangle_A \quad . \tag{37}$$

We recognize that g_{WAA} is the same for all isodoublet states. This guarantees universality of weak (CC and NC) interactions.

4. Experimental Signatures of Compositeness

An important issue within subconstituent models is to single out experimentally accessible indications of a composite structure of leptons and quarks. It is easy to list a whole number of spectacular phenomena, which would occur at utopically high energies $(E \gg \Lambda_H)$. For instance, radial and orbital excitations of leptons and quarks should appear, thus prolifering the number of (quasi fundamental) fermions. At even higher energies, we expect that the hypercolor degree of freedom itself starts to get melted up. Then, the difference between leptons and quarks will become uneffective, and leptons will be copiously produced by processes at such high energies.

Here we don't want to investigate utopia, but look for indications, which could show up at energies which are accessible at present day or in the near future (LEP). In the early stage of the subconstituent philosophy it has been believed, that a spectacular consequence of the subconstituent picture of weak interactions would be that the masses of intermediate vector bosons should deviate grossly from the values predicted by the Glashow-Weinberg-Salam model. Yet, the more detailed investigations, which we have sketched before, show that this need not be the case. Even more: Depending on the W-excitation spectrum, we expect that those masses lie very near (less than 20% deviations) to the canonical value. Thus, masses are not a good criterion as to compositeness.

On the other hand, if we stick to the simple model presented below, one should not overlook the fact, that it coincides with the standard picture (WS) only at low q^2 . At $q^2 \geq m_W^2$ it leads to a behaviour, which is apparently different from the G-W-S-predictions (for details see ref. 11)):

- For $q^2 \sim M_Z^2$ both weak couplings and weak mixing angles should deviate from the values at low q^2 (which are the fixed values of GWS). This would also influence the asymmetry around M_Z.
- At $q^2 > M_Z^2$, isoscalar contributions to NC processes will become sizable. In connection with it, we expect the appearance of currents with a different relative weight of vector and axialvector contributions.
- It might also happen that (pseudo-)scalar and/or tensor exchanges become important at some (completely unknown) energy.
- For energies high enough, higher excited intermediate vector bosons should come into the game.

Further indications of the composite nature of intermediate bosons have been recently studied by Renard[25]. He recognized that a composite Z_0 should show a decay pattern different to that of an elementary (G-W-S) boson. Firstly, multifermionic (multileptonic) final states will be more important. Secondly, some specific decay modes, which are rare modes within the G-W-S-scheme, will be enhanced by three or four orders of magnitude. For instance, the mode

$$Z_0 \rightarrow 3 \text{ gluons} \quad , \tag{38}$$

which shows up as a 3-jet-final state, could have a branching fraction as large as 10% for a composite Z_0, whereas in the G-W-S-model it proceeds via quark loops and is suppressed down to $10^{-3}\%$ due to the additional quark propagators. (38) is an especially interesting mode, since it is also sensitive to the color assignement of subconstituents. Other crucial decay modes are

$$Z_0 \rightarrow \gamma\gamma\gamma \quad , \qquad Z_0 \rightarrow gg\gamma \tag{39}$$

Thus, a detailed study of the Z_0 could yield some indications of a composite structure of particles, which originally have been believed to be elementary.

5. Conclusions

In this talk we have discussed a possible interpretation of electroweak interactions within the subconstituent picture of matter. Based on the assumption of global $SU(2)_W$-symmetry we have shown that both charged current and neutral current phenomenology at low q^2 can be easily accounted for, mainly because mixing of the photon with (composite) intermediate vector bosons is necessarily implied by the composite structure of the latter.

To describe the electromagnetic coupling of (composite) leptons and quarks we have adopted the view of W-dominance. In its most radical version (dominance of the lowest-lying bosons) it leads to mass values for $W^{\pm}$ and Z^0 identical to the (uncorrected) values predicted by the G-W-S-model. Inclusion of higher excitations would only imply corrections of 15 - 20% to this prediction. Thus, contrary to the original guesses, the intermediate boson masses are expected to lie near to the canonical values, even within subconstituent models.

Within this picture, the weak mixing angle ($\sin^2\theta_W$) is interpreted as an electromagnetic effect, its relatively large size being connected with the large mass scale of hypercolor dynamics.

Although the predictions of the new scheme coincide with the G-W-S-results at low q^2, at higher momentum transfers ($q^2 \geq m_W^2$) drastic departures from the canonical picture are expected. Clearly, a deeper knowledge about the subconstituent dynamics is needed in order to get more detailed predictions on the high energy behaviour.

Acknowledgements

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$$L_{mix} = -\frac{1}{2}\, m^2_{\gamma W}\, (W^{(3)}_\mu A'^\mu + A'_\mu W^{(3)\mu}) \quad ,$$

but in this case there must also exist a specifically chosen mass term for the (unphysical) photon, in order to restore gauge invariance of the whole Lagrangian.

14) One can easily convince oneself about the following fact: If there are two intermediate vector bosons, f.i. one coupling to left-handed current and the other to right-handed current, then there are twice as much contributions to the full propagator and the bound (4) takes the form

$$2\lambda^2 < 1 \quad \text{or, equivalently,} \quad \lambda < \frac{1}{\sqrt{2}} \quad .$$

Consequently, (10) changes to

$$m_W \leq \frac{1}{\sqrt{2}} \, 162 \text{ GeV} = 127 \text{ GeV} \quad .$$

This is the main result of Barbieri and Mohapatra, ref. 4).

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16) For completeness, we mention the main parts of this Lagrangian. They describe the couplings of the various isovector $(\vec{W}_i)$ and isoscalar $(W_i^{(0)})$ bosons:

$$L_{IVB} = - \sum_i g_{W_i} \vec{W}_{i_\mu} \vec{J}^\mu - \sum_i g_{Y_i} W^{(0)}_{i_\mu} J^{(0)\mu}$$

and the photonic couplings (à la Kroll-Lee-Zumino):

$$L_A = \sum_i \left(- \frac{\lambda_{W_i}}{4} (W^{(3)}_{i_{\mu\nu}} F'^{\mu\nu} + F_{\mu\nu} W_i^{(3)\mu\nu}) - \lambda_{W_i} g_{W_i} A_\mu J^{(3)\mu}\right)$$
$$+ (W^{(3)} \rightarrow W^{(0)}) \quad .$$

Here, $\vec{J}$ denotes the usual left-handed weak isospin current of leptons and quarks, whereas $J^{(0)}$ represents the isoscalar current, which consists of left- and right-handed pieces, in order to guarantee the parity conserving photon coupling.

17) The combinations

$$g_i \lambda_i \frac{m_i^2}{q^2 - m_i^2} \quad ,$$

which appear on the second line, are a result of both the (current) mixing $W^{(i)}$-γ-vertex and the contact term (see footnote 16)):

$$g_i \frac{-1}{q^2 - m_i^2} \lambda_i q^2 + g_i \lambda_i = - g_i \lambda_i \frac{m_i^2}{q^2 - m_i^2} \quad .$$

18) It is worth noting that the unification condition implies some further interesting properties[12], like asymptotic $(q^2 \rightarrow \infty)$ global $SU(2) \times U(1)$ symmetry, tree unitarity (orders) and canonical magnetic moments of W^+ and W^-. Whether it also guarantees renormalizability of the effective Lagrangian is not yet clear.

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COMPOSITE QUARKS, LEPTONS AND WEAK GAUGE BOSONS

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The motivations for looking into possible substructures of some or all the "elementary" particles of the standard model are not lacking (1). The computability of the basic parameters, the rationale of the quantum numbers and the economy in the fundamental degrees of freedom can be quoted among these motivations. They are generally shared by alternative (or complementary) attempts to go beyond the standard model with a different focus put on the various cases. Grand unification, as an example enphasizes the rationale of the quantum numbers, e.g. charge quantization, but may have little to improve on the economy of the basic fields and on computability.

Clearly an increasing degree of radicality is involved in considering as composite particles:

a) the Higgs scalar;

b) quarks and leptons;

c) W and Z;

d) photons and gluons.

That the Higgs particles come first in this list is quite natural. Technicolor is an interesting idea whose problems (read extended technicolor) may find a solution by going to step b).

To go further, a picture of composite gauge vector bosons is not without problems, to say it mildly. This is why in the above list W and Z come after quarks and leptons. On the other hand, are they gauge particles with the same status of the gluons and, even more, of the photon? As especially emphasized by Pati and Salam, a composite picture for the gauge bosons, including the photon and the gluons, is advocated by the requirement of economy in the basic degrees of freedom.

Quite in general it must be said that there is no sharp boundary between the various possibilities, since one can have compositeness scales which are vastly different from each other and stay in different ratios to the masses of the particles. Let me consider the case of composite vector bosons. Can a field theory model be conceived of where vector bound states exist with a mass M negligible relative to the inverse size (the compositeness scale Λ_H)? If this is the case, arguments based on the high energy behaviour of the effective interactions between these vector bosons require them to be gauge interactions, spontaneously broken if $M \neq 0$. So for all energies $E \ll \Lambda_H$, such a model would not be distinguishable from a model of fundamental vector bosons. However, do we know of anything else but gauge invariance itself that may possibly give rise to light vector bosons? [+] Indeed the same considerations of the high energy behaviour may require that the basic lagrangian be gauge invariant, with fundamental, rather then composite, vector bosons if renormalizability is demanded. On the other hand, for the W and the Z one can entertain the possibility that their mass is not much smaller than their inverse size ($\Lambda_H \lesssim$ 1 Tev). In such a case, one would have a situation at the border between a ρ-like

+) In these respects supersymmetry may provide a completely new and favorable situation.

picture and a gauge picture for the weak vector bosons.

The possibility is considered that a compositeness scale of quarks, leptons and weak gauge bosons be related to (or some what bigger than) the Fermi scale $G_F^{-1/2}$. This is an interesting possibility per se, - it might give a more physical significance to the Fermi scale than in the fundamental Higgs picture, - even though it is not clear at the momentum how it could meet the motivations already mentioned.

Some explicit models exist in the literature[1),2)]. They are generally based on a preonic lagrangian with a gauged symmetry including colour, electromagnetism and the binding hypercolour interaction. The chiral global symmetry is assumed to remain at last partly unbroken to ensure the presence of massless (light) fermionic composites. This unbroken symmetry also includes the $SU(2)_L$ of the weak interactions, with electromagnetism turned off, and may contain the Pati Salam SU(4) - symmetry when colour too is neglected. Many of the proposed models have fermionic as well as scalar preons, which might call for a supersymmetric generalization[3)].

Within this general context a few selected topic represent potential problems of the nearby compositeness picture. In examining them, I expressely take the view of the ρ-like picture, but I keep in mind a situation at the bolder with the gauge picture as I have explained.

i) Unwanted effective interactions.

A too low value of the compositeness scale brings forth unobserved effective interactions between the composite fermions.

Otherwise stated, to have precisely the interactions mediated by the weak composite vector bosons emerging from all the expected effective low energy interactions is difficult. At least the effect of a weak isoscalar vector exchange should show up at some level[4]. The present limit for the strenght of the weak isoscalar current-current interactions relative to the standard isovector term is $G_F^{(0)}/G_F \lesssim 5\%$.

ii) The value of "sen$^2\theta_W$" and of the W and Z masses.

Here is another issue which strongly calls for an anomalously light W-boson. Infact the observed structure of the neutral current interactions is accounted for by an admixture between the elementary photon and the composite W-boson. As such the effective mixing parameter, analogous to sen$^2\theta_W$ of the standard model, is of order α and therefore expected to be too small[4]. This unless the proximity of the W to the photon mass on the scale of the binding interactions produces a dynamical enhancement of this mixing parameter[5]. In principle an ideal way to study this problem is offered by the framework of the asymptotic freedom sum rules applied to the binding "hypercolour" gauge theory. In analogy with what is done in the QCD case for the ρ-photon mixing, by introducing the W-photon mixing parameter through

$$\mathcal{L}^{(mix)} = \lambda W_{\mu\nu} F^{\mu\nu} , \qquad (1)$$

one can have[6]

$$\lambda^2 \simeq \frac{\alpha}{3\pi} \left(\Sigma Q^2\Big|_{preons} - \Sigma Q^2\Big|_{\substack{quarks\ and\\ leptons}} \right) \left(\frac{M}{M_W} \right)^2 , \qquad (2)$$

where Q is the charge in units of e, and M is the mass where the perturbative preonic continuum sets in for example in $e^+ e^-$ annihilation. Preesumably $M \simeq (5\div 10)\Lambda_H$, so that $M_W \lesssim \Lambda_H$ might give the wanted enhancement. Eq.(2) is meaningful only for a total squared charge over the preons bigger than that over the composite quarks and leptons. Otherwise the quark and lepton continuum itself saturates or even oversaturates the sum rule with no room left for a significant contribution of the W-resonance.

Needless to say, none of these arguments says what M_W should precisely be, nor what its relation with the Z-mass should be. No more insight is gained other then the limits given by the general analysis of Bjorken[4] and Hung and Sakurai[7]. Only the following comment can be made related to what I said in the introduction. The lighter the W needs to be relative to the compositeness scale, the more the ρ -like picture goes into the gauge picture, as require by consistency. Then only relatively small deviations may be expected from the values of M_W and M_Z predicted by the standard model.

iii) Quark-lepton and interfamily universality.

In the extreme ρ-like picture of the weak interactions, the CVC concept is lost[+]. This poses the problem of quark lepton universality in an apparently severe way. If however the preonic theory possesses a Pati-Salam symmetry the situation can be remedied. This is because the violations of quark lepton universality relative to those existing in the standard model are only due to gluon exchanges occurring within the size $\sim 1/\Lambda_H$ of the composite quarks and leptons.

Let us look more closely at the phenomenological situation.

+) The presence of a conserved preonic $SU(2)_L$ current cannot be related to quark-lepton universality unless additional hypotesis are made[8].

Because of Cabibbo-like mixings a test of quark-lepton universality cannot be disentangled from interfamily universality. In terms of the Kobayashi Maskawa parameters, the present knowledge is summarized by saying that in the standard model, the relation

$$|u_{ud}|^2 + |u_{us}|^2 + |u_{ub}|^2 + \ldots = 1$$

puts an upper limit

$$|u_{ub}|^2 + \ldots \lesssim 1\,\%$$

on the couplings of the u-quark to the b-quark and the possible heavier quarks[9]. It is worth noticing that this analysis includes a radiative correction to $|u_{ud}|^2$ -related essentially to $\cos\theta_c$ - which significantly reduces the zeroth order value of about 3%. On the other hand the source of this anomalously large correction can be traced bach to a logarithmic divergence of the pure four fermion interaction, cut off in the standard model by the intermediate boson mass.

In a composite model with an approximate $SU(4)_{PS}$ symmetry the right hand side receives a correction of relative order $\alpha_s(\Lambda_H)/\pi$, namely 2 - 3% for $\Lambda_H = 10^2 \div 10^3$ GeV, due to gluon exchanges within the size $1/\Lambda_H$ of the composite quarks and leptons. This calls for an effect of violation of quark lepton universality which is of the order of the present uncertainty. On the other hand in a composite model too, on top of the effective four fermion interaction there will be a radiative correction effect similar to the one occorring in the standard model and cut off at the compositeness size. In other words, the logarithmic correction alone cannot be considered as a really distinctive test of the standard model.

iv) Flavour changing neutral currents

In the previous discussion I have tacitly assumed in the composite model under consideration the validity of interfamily universality. Such an assumption goes together with the fact that one does not know yet how to describe families, except by introducing the corresponding explicit degrees of freedom in the preonic lagrangian+). In this last case interfamily universality is enforced by a symmetry of the basic lagrangian, which of course has to be slightly broken. This happens explicitly for example in the Abbott Farhi model.

Assuming then the family universality, the GIM mechanism to suppress the flavour changing neutral currents may not be a distinctive feature of the standard model. For that it is important to recognize that the standard radiative correction diagram, giving rise for example to a $\Delta S=2$ effective lagrangian (Fig.1a), can be contracted, in the limit of infinite W-boson mass, to the still finite diagram of fig.1b). This is because of the occurrence of the GIM cancellation.

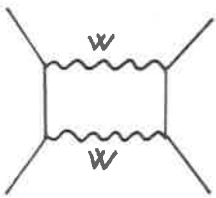

Fig.1a)

Fig.1b)

Diagrams giving rise to the effective $\Delta S=2$ lagrangian.

+) Here again supersymmetry may have something new to say[3].

In the last diagram only the point-like four fermion interactions appear and the typical momentum running in the loop is of the order of the relevant quark masses. Now the diagram of fig 1b) would also be the appropriate description of the $\Delta S = 2$ effective lagrangian in a composite model with a family universal four-fermion interaction. A GIM cancellation still takes place and no reference is made to the details of the model at energies of order Λ_H.

Why to contrast the neat picture of the standard electroweak model with a "scenario" which is still, at best, in the clouds of an unknown strong interaction dynamics and with so many potential problems? Admittedly all this seems a rather marginal point of view. However the idea that the Fermi scale be somehow related to a new compositeness scale has its appeal. True enough, the standard description of the Fermi scale, with a fundamental scalar field acquiring a non vanishing vacuum expectation value is much more concrete. But it has drawbacks too, wich infact motivate a lot of research. In any case its seems fair to let experiments decide.

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Composite Models for Quarks and Leptons *

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Abstract

We discuss the motivation for constructing composite models for quarks and leptons, the hopes we have for a successful model and the difficulties encountered, so far, in this field. This paper corresponds to the contents of lectures given at the SLAC Summer Institute (August 1982), at the DESY Workshop on "Electroweak Interactions at High Energies" (September 1982) and at the Solvay Conference at the University of Texas, Austin, Texas (November 1982).

1. Foreword

The possibility that quarks and leptons are composite, plays a peculiar role in present-day particle physics. On one hand, it is the most natural extrapolation of the development of modern physics and the least imaginative proposition for extending our theoretical ideas beyond those of the "standard model" of electromagnetic, weak and color interactions. On the other hand, any attempt to construct an explicit composite model immediately faces serious difficulties, necessitating assumptions and ideas which are at least unusual, possibly revolutionary. Thus, we are dealing with an approach which is, paradoxically, extremely speculative and somewhat unimaginative at the same time. In this lecture we review the present status of this field, emphasizing the hopes as well as the difficulties. We do not discuss specific models in any detail.

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2. Compositeness and the Fundamental Interactions.

Every level of compositeness in the history of modern physics, led to a major revision of the list of fundamental forces. The understanding of atomic structure showed that Van-der-Waals forces are residual electromagnetic effects and are therefore not a separate fundamental interaction. The substructure of the nucleus revealed the existence of a new short range force (the Strong or Nuclear or Hadronic Force) and led to the identification of an even shorter-range force (the Weak force). The hadron substructure uncovered the Color force and showed that the Nuclear or Hadronic forces are residual color effects. It is almost certain that a possible substructure of quarks and leptons, if found, will reveal one or more new forces. It is also possible that it may demote one or more of the existing "fundamental interactions" into the role of a "residual interaction".

New fundamental particles may be found in the future and may have escaped detection in the past in several different ways:

(i) The production of high mass free particles requires experiments at sufficiently high energies. The W and Z bosons and the toponium bound states have presumably escaped detection in such a way.

(ii) Low mass (or even massless) particles may escape detection if they interact weakly with all visible matter. That is how the neutrino eluded experiments for several decades. We may have a variety of light neutral particles, especially Goldstone or pseudo-Goldstone-bosons escaping us now in such a way.

(iii) Confined particles, even if they have small effective masses, require experimental probes with energies far exceeding their confinement energy scale. Only such probes can reach into small distances in which the confined particle can be indirectly observed. Thus, many Gev's were needed for indirect evidence for light quarks and massless gluons, where the confinement scale is $\Lambda_c \sim 100$ MeV If there is a second confining force with a scale Λ_H, we will need energies which substantially exceed Λ_H in order to "observe" the new confined states.

In close correspondence with these three methods, there are at least three ways of discovering a new fundamental interaction:

(i) It may be a short range interaction corresponding to distance scales and energy scales beyond those presently available. That is how the strong nuclear force remained unknown until the 1930's.

(ii) It may be an extremely weak force, not necessarily with a short range. Such a force can be uncovered only by increasing the sensitivity and accuracy of low energy experiments.

(iii) It may be a confining force, mediated by a massless boson but possessing a confinement scale beyond present energies.

Any of these possibilities may "hide" additional new interactions which could play an important role in quark and lepton compositeness. There may also be a fourth, hitherto unknown, method by which new particles and new interactions may escape detection. After all, had we discussed this subject 15 years ago, we would have listed only the first two items in the above lists of possibilities. We would not have anticipated the possiblity that an extremely strong force with massless force-carriers (Color) could have escaped detection. Today, we may be equally blind. The binding force of the constituents of quarks and leptons in perhaps eluding us in a new clever way.

3. Why Do We Wish to Go Beyond the Standard Model?

The standard model of electromagnetic, weak and color forces is based on a renormalizable local gauge theory, where $SU(3)_c xSU(2)xU(1)$ is the basic gauge group, and its spontaneous breaking via a Higgs mechanism leaves an unbroken $SU(3)_c xU(1)_{EM}$ gauge group. Ignoring temporarily the Higgs sector, the model involves massless quarks, leptons, photon, gluons and weak bosons with only three independent parameters representing the coupling constants of the three interactions. The Higgs sector induces all the necessary masses, increasing the number of arbitrary parameters to twenty or more.

There are several reasons which lead us to believe that the standard model, in spite of its elegance, self-consistency and experimental success, cannot be the final answer. Every one of these reasons is, to some degree, a matter of taste. However, the emerging overall picture convinces us that there must be some new fundamental physics beyond the standard model.

Why do we wish to go beyond the standard model?

(i) Too many parameters. It is unlikely that the laws of physics contain over twenty independent parameters and that the various Cabibbo angles and quark or lepton masses are as fundamental as the fine structure constant. A theory beyond the standard model may enable us to calculate most of these parameters, starting from a small number of constants. Such a calculation will hopefully explain the peculiar mass spectrum of the observed quarks and leptons.

(ii) Generation Puzzle. The standard model contains no clues explaining the existence of several generations of quarks and leptons. There is no good reason for having three generations (or any other specific number of generations). We do not know what distinguishes the generations from each other. A new quantum number or a new property which "labels" the generations can emerge only from theories which go beyond the standard model. Any hope of calculating the mass matrices for the quarks and leptons must involve an understanding of the distinction between fermions of different generations.

(iii) Pattern within one generations. The mysterious triplication of generations enhances the significance of another puzzle, namely- the pattern observed within one generation. The standard model does not explain why quark and lepton charges are quantized in a related way. It does not explain why the color and electric charge are correlated (integer charge always comes with color singlets; noninteger charge with color triplets). It also does not explain why the quarks and leptons possess identical SU(2) properties (left-handed doublets; right-handed singlets) and why all integer multiples of $Q=\frac{1}{3}$ between $Q=-1$ and $Q=+1$ appear but no $|Q|>1$ fermions seem to exist. All of these features would have required explanation even if we had only one generation. Since the same pattern occurs three times, there must be a particularly good reason for repeatedly having that specific pattern within the generation.

(iv) Unification. One other motivation for going beyond the standard model is the obvious hope of unifying the three interactions. We must remember that the SU(2)xU(1) gauge theory provides us with a beautiful and important connection between the electromagnetic and weak interactions but it does not fully unify them. We still have an independent coupling constant for each interaction.

(v) Hierarchy problem and fine tuning. These may count as two problems or one problem depending on one's point of view. In any way of looking at it, we have here a mismatch between different energy scales. Assuming that there is an important energy scale beyond the standard model (be it the Planck mass or a somewhat lower energy scale of some other type of new physics), it is difficult to understand how particles with masses corresponding to the low energy scales of the standard model can survive enormous self-energy corrections. Vector bosons and fermions may be protected from such corrections by a gauge symmetry or a chiral symmetry, respectively. Scalars are not protected, in general.

The above five arguments are not compelling in any rigorous sense. However, after considering them, it is difficult to avoid the conclusion that some deeper theory must lie beyond the standard model, settling at least some of the above issues.

4. Avenues Leading Beyond the Standard Model.

We have listed five reasons to go beyond the standard model: (i) Too many parameters; (ii) Generation puzzle; (iii) Pattern within one Generation; (iv) Unification; (v) Hierarchy problem and fine tunning. At least five different classes of approaches have been proposed for handling these issues. There is no one-to-one correspondence between the five problems and the five classes of models. Let us briefly review each approach:

(a) "Horizontal Symmetries": These are alleged new symmetries[1] connecting the different generations. Such symmetries cannot settle any issue except, possibly, the generation puzzle (ii). We return to some of these ideas in section 9. No convincing horizontal symmetry scheme has been proposed, so far.

(b) "Technicolor". The fine tuning problem (v) may be resolved or at least postponed, to higher energies, by postulating that the Higgs scalars are condensates of new fundamental "technifermions" bound by a new confining "technicolor" interaction[2]. This approach does not shed any light on problems (i)-(iv). In fact, it normally leads to additional particles and additional free parameters. The technicolor idea is a limited version of compositeness, in which scalars and, consequently, the longitudinal components of massive gauge bosons, are composite objects. Quarks, leptons, transverse gauge bosons and technifermions are all fundamental. In order to produce quark and lepton masses, technicolor schemes must be extended to include new gauge bosons which connect quarks to techniquarks[3]. Such bosons usually play a role similar to that of Horizontal gauge bosons, thus incorporating the Horizontal symmetry approach into the technicolor scheme. Here, again, no satisfactory model is presently available.

(c) Grand Unification. This approach satisfactorily settles points (iii) and (iv). The structure of one generation is beautifully accounted for in any scheme[4] based on O(10) or any of its candidate subgroups (SU(5) or SU(4)xSU(2)xSU(2)). The three interactions are clearly unified. However, the number of free parameters is increased beyond those of the standard model, no explanation is given to the generation puzzle and the hierarchy problem remains unanswered. In addition, we must face an energy "desert" spanning twelve orders of magnitude and the controversial possibility of heavy magnetic monopoles. Proton decay is the earliest crucial test of grand unified theories.

There are ambitious attempts[5)] to combine Grand unification, technicolor and horizontal symmetries in all-encompassing schemes based on large groups such as SU(7), SO(14) etc. No convincing model has been found.

(d) Supersymmetry. Supersymmetry, besides being an attractive mathematical "toy", provides a hope for settling the hierarchy problem[6)]. Unfortunately, the particle spectrum is doubled, introducing alleged supersymmetric partners for all gauge bosons, quarks , leptons and scalars. Thus we have more particles and more parameters, without solving any of the problems (i)-(iv). Grand unification can be combined with supersymmetry, thus combining possible solutions to points (iii)(iv)(v), but the price is, again, a significantly more complicated spectrum of particles as well as additional theoretical problems.

(e) Composite Models. Here we assume that quarks, leptons, scalars and possibly some gauge bosons are composite objects of some new fundamental constituents. Here, again, no satisfactory model is available. In the following sections (6-12) we discuss the hopes and difficulties of such models, vis-a-vis the various reasons for going beyond the standard model.

But, before we move on to our discussion of composite models, we must discuss the question of experimental tests.

5. Experiments Beyond the Standard Model

Theories which go beyond the standard model, like all other theories in physics, must pass two types of tests: self-consistency and agreement with experiment. The requirement of theoretical self-consistency is not as simple as it sounds. We must remember that in all previous stages in the understanding of the structure of matter, from the Bohr atom to the quark model, the original version of the theory had many correct ingredients but suffered from serious theoretical inconsistencies. In all cases, experimental clues played a crucial role in the acceptance of the correct ideas. The satisfactory theoretical self-consistency came only gradually, with modifications which were partly discovered by pure reasoning and partly through new experimental facts. Our main difficulty today is the total lack of any experimental facts which might force us to go beyond the standard model. To rely entirely on arguments of theoretical self-consistency is dangerous. It is not clear that such arguments would have allowed the quark model to be developed!

We therefore believe that experimental clues are crucial. In order to review such clues we must first discuss the relationship between the hypothetical new theory which goes beyond the standard model, and the Lagrangian of the standard model. In all cases we do not wish to discard the standard model. We would like to keep it as a good aproximation of the ultimate theory, valid at energies well below the new high energy scale. Schematically, we may consider the following situation: We have a new theory at small distances and high energies. Its Lagrangian, L_{NEW} is useful for describing high energy phenomena. It hopefully corresponds to a renormalizable theory which is fully self-consistent. At lower energies we may have an effective Lagrangian L_{EFF} which, in principle, can be derived from the fundamental high energy Lagrangian. It is presumably approximately equal to the standard-model Lagrangian L_{SM}.

$$L_{NEW} \rightarrow L_{EFF} \cong L_{SM}$$

When we search for experimental tests of the new theory, we may look for two general classes of tests:

(i) Tests of L_{NEW}. These are, necessarily, high energy tests involving future high-energy accelerators. If the energy scale of L_{NEW} is Λ_H, we may look for particles of mass Λ_H (e.g. Horizontal gauge bosons in Horizontal symmetry schemes, Monopoles in grand unified theories, Technihadrons in Technicolor models, etc.). We may also look for light particles which are confined within small confinement radii corresponding to a high energy scale Λ_H (e.g. preons in some composite models, techniquarks etc).

(ii) Tests of the small difference between L_{EFF} and L_{SM}. The low-energy phenomenology of the new scheme is, presumably, almost identical to that of the standard model. The small differences between the two may actually provide us with the first clues for physics beyond the standard model. Such clues may come from a variety of terms in L_{EFF}. A few examples:

(a) L_{EFF} may contain high-dimension four-fermion operators such as $uude^-$ or $\bar{\mu}e\bar{e}e$. These must be preceded by a coefficient of order Λ_H^{-2}. Such terms induce transitions like $p \to e^+\pi^o$ or $\mu \to 3e$, respectively. These are low-energy processes which, if observed, would necessitate some physics beyond the standard model.

(b) L_{EFF} may show slight deviations from certain coupling-constant relations of the standard model. For instance, the coupling constants $g_{We\nu}$, $g_{W\mu\nu}$ which are identical in the standard model ("universality") may turn out to differ by terms of order m_ℓ/Λ_H as a result of some e-μ difference which is revealed only at the scale of Λ_H.

(c) L_{EFF} may contain weakly coupled light Goldstone bosons with Yukawa couplings such as $\frac{m_e}{\Lambda}\chi\bar{e}e$. Such bosons are difficult to detect. If they exist, they could provide hints for some new physics.

The above classes of experimental tests are relevant to any theory which goes beyond the standard model. They are also true for composite models of quarks and leptons. The right-hand side of figure 1 indicates the general orders of magnitude of the new energy scale, corresponding to various theoretical approaches beyond the standard model. On the left-hand side of the same figure we show some of the experimental bounds which are relevant to composite models of quarks and leptons. We now turn to a discussion of these bounds.

6. Experimental Constraints on Composite Models of Quarks and Leptons

Any attempt to construct a composite model for quarks and leptons, must take into account several experimental constraints:

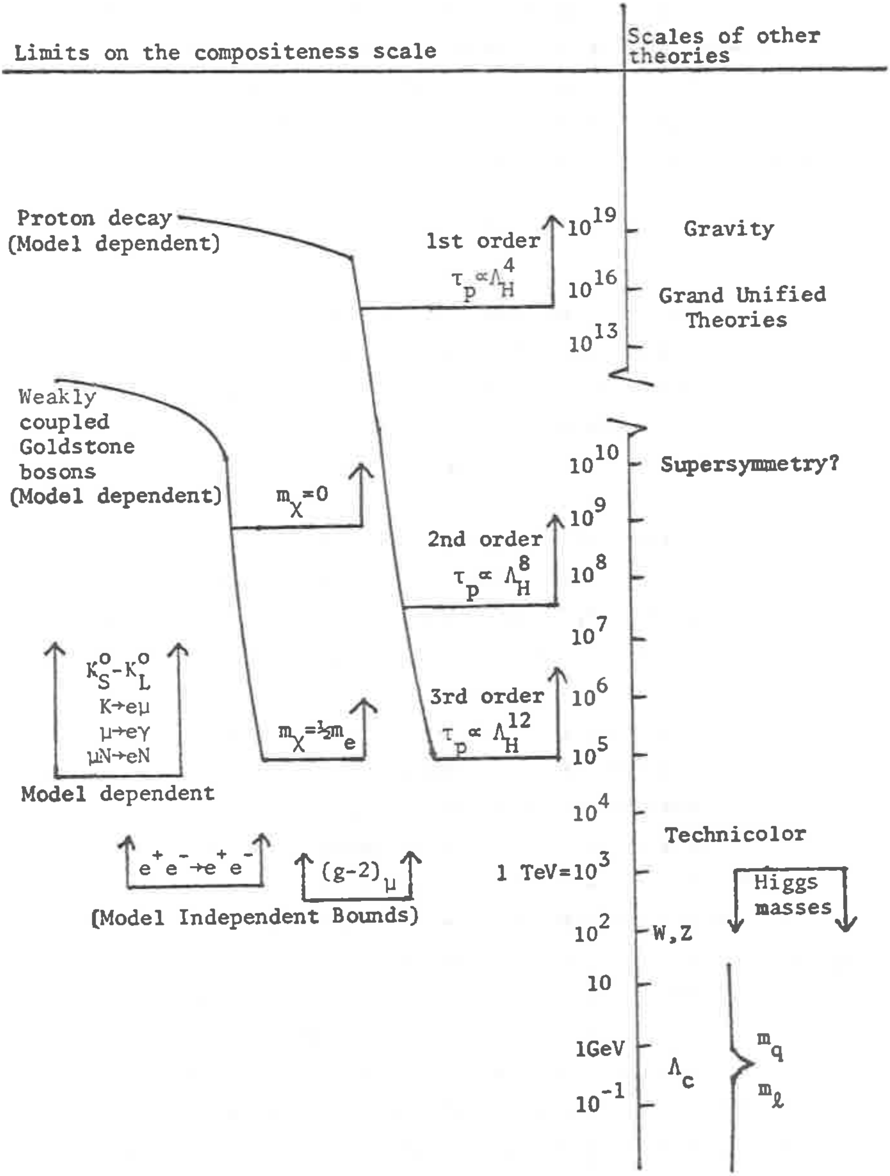

Figure 1: Experimental limits on the energy scales of theories beyond the standard model.

(i) The anomalous magnetic moments of the electron and muon provide us with an (almost) model-independent constraint[7]. It has been argued that if a composite structure at a scale Λ_H leads to deviations from the QED predictions for g-2, and if the composite model has a chiral symmetry, we expect:

$$\delta(g-2)_\ell \sim 0\left(\frac{m_\ell}{\Lambda_H}\right)^2$$

In the case of the muon, present g-2 experiments yield $\Lambda_H > 500$GeV. If we relax the chiral-symmetry assumption, the bound is much more severe. The 500GeV bound for the muon is essentially model-independent. However, it is only an order of magnitude estimate, and factors of π could easily change it in either direction.

(ii) If electrons are composite, we expect L_{EFF} to contain a four-fermion term of the form $\bar{e}e\bar{e}e$. Such a term would contribute to the cross-section for $e^+e^- \to e^+e^-$. The present agreement between this cross-section and QED places a new model-independent bound[8] on Λ_H. Here, again, one can only estimate the order of magnitude, obtaining $\Lambda_H > 700$GeV. Similar related estimates may be obtained for $e^+e^- \to \mu^+\mu^-$ and for neutral current neutrino reactions[8].

(iii) The absence of the decays $\mu \to e\gamma$, $\mu \to 3e$, $K \to \mu e$, $K \to \pi\mu e$ and the reaction $\mu N \to eN$ provide us with model-<u>dependent</u> bounds on compositeness. If the muon and the electron are prevented from transforming into each other by some selection rule or by a strong suppression-factor which depends on their detailed internal structure, no useful bounds can be derived from the present experimental limits. However, it is possible that the muon and the electron can easily convert into each other, the transition being supressed only by the physical dimension of the composite system. In such a case we would expect L_{EFF} to include effective four-fermion terms such as $\frac{1}{\Lambda^2}\bar{\mu}e\bar{e}e$ or $\frac{1}{\Lambda^2}\bar{s}d\bar{e}\mu$. These would enable us to derive limits of the order of $\Lambda_H \sim 10$-100TeV for various processes. We emphasize, however, that this case is very different from the previous item. No selection rule can forbid an $\bar{e}e\bar{e}e$ effective interaction, but it is perfectly reasonable to expect a small suppression factor based on some selection rule, appearing in an expression of the form $\frac{\varepsilon}{\Lambda^2}$ $\bar{\mu}e\bar{e}e$. If Λ is sufficiently small, the bound on Λ_H may become totally useless.

(iv) A related experimental bound follows from the well-measured K_S-K_L mass difference. The observed value can be accounted for by the standard model. A composite model for quarks might allow an effective term of the form $\frac{\varepsilon}{\Lambda^2}\bar{s}d\bar{s}d$ contributing to ΔM. Here, again, the constraints are model dependent. If ε is sufficiently small, due to a selection rule based on the different internal structure of s and d, no useful bound can be deduced. If $\varepsilon \sim O(1)$, we may obtain a very strong bound around 1000 TeV. However, it is somewhat unlikely that any composite model would allow $\varepsilon \sim O(1)$.

(v) Many composite models[9] may involve massless Goldstone bosons or extremely light pseudo-Goldstone bosons χ, whose Yakawa couplings to ordinary quarks and leptons are of order $\frac{m_\ell}{\Lambda_H}$ or $\frac{m_q}{\Lambda_H}$. Such bosons can easily escape detection in terrestrial experiments. However, the process $\gamma+e \to \chi+e$ which must occur frequently in stars places a limit[10] on the $\chi\bar{e}e$ coupling, and through it - on the compositeness scale. From the known limit on the allowed energy loss of red giant stars, we obtain: for a massless Goldstone boson - $\Lambda_H>10^9$GeV; for $m_\chi=\frac{1}{2}m_e$, $\Lambda>10^5$GeV; for $m_\chi=m_e$ no useful limit is obtained in this way. Thus the constraints are extremely sensitive to the boson mass.

(vi) Proton decay provides another crucial, but model-dependent, test for the compositeness scale. If no selection rules or supression factors exist, we find the usual result:

$$\tau_p \sim \frac{\Lambda_H^4}{m_p^5}$$

leading to $\Lambda_H>10^{15}$GeV. However, proton decay is actually forbidden in some composite models. In other models it may proceed in second order[11] (giving $\Lambda_H>10^7$ GeV) or in third order[12] (giving $\Lambda_H>10^5$GeV)

The overall picture is the following: At present, there is no experimental evidence for quark or lepton compositeness. Model-independent bounds tell us that $\Lambda_H > 0.5$ TeV or $r < 4.10^{-17}$cm. Any specific model must be compared with a variety of model-dependent tests. For instance, anyone who wishes to suggest that quark and lepton compositeness will be revealed already at energies around, say, 10-1000 TeV, must provide strong suppression factors for proton decay, and the K_S-K_L mass difference, as well as avoid massless Goldstone bosons which couple to electrons and up and down quarks. Additional experimental tests which one must consider involve $\sin^2\Theta_W$, the W-Z mass ratio, the W magnetic moment, the possible existence of "right-handed" weak bosons, etc.

7. Requirements from an Ideal Composite Model

What do we hope to achieve by constructing a successful composite model of quarks and leptons?

(i) Such a model should include a few species of fundamental objects interacting with each other through few types of fundamental interactions. The total number of parameters is presumably extremely small: several coupling constants and possibly (but not necessarily) a few mass parameters. All masses of the composite quarks and leptons should, in principle, be calculable from the parameters of the fundamental theory, in the same way that all hadronic masses and coupling constants are, in principle, calculable from the QCD coupling and a few quark masses.

(ii) The pattern of quarks and leptons within one generation should be fully explained in terms of the features of the fundamental fermions. For instance, if both quarks and leptons are composites of the same set of fundamental fermions, their charge quantization must clearly be related. The peculiar relation between electric charge and color may simply emerge from the color and charge of the fermions. The restrictions on $|Q|$ may be related to the number of constituents within a composite quark or lepton, in the same way that the limitations on the strangeness or isospin of hadrons follow from the number of valence quarks in a hadron.

(iii) The different generations may be excitations of a composite system, similar to excited atoms, nuclei or hadrons. The type of excitation in each case must be different, however.

(iv) The scalar particles, as well as the quarks and leptons, are presumably composites of the new fundamental fermions. Hopefully, no fundamental scalar particles are necessary. The fundamental fermions may be massless or may have explicit mass terms, but need not gain masses through symmetry breaking. The problem of fine tuning may thus be avoided.

(v) Other features are left open. Color, Electromagnetism and the weak interactions may all exist in the underlying theory. Alternatively, one or more of these interactions may turn out to be a residual force, appearing only in L_{EFF}. Additional color-like or other types of forces may be needed in order to bind the new fermions inside the quarks and leptons. The underlying theory may be left-right symmetric, with Parity being spontaneously broken at the composite level. Alternatively, the fundamental theory may already include explicit parity violation.

It is not at all clear that a composite model with all the above desired features can be constructed, but it is certainly worth exploring. So far no one has come close.

Among the various problems which face model builders, we choose to discuss four in some detail:

(a) The problem of scales. Quark and lepton masses are much smaller than any possible compositeness scale. This is the most difficult problem for all composite models, and it has several interesting aspects, which we discuss in section 8.

(b) The generation puzzle. If quarks and leptons are composite, what kind of quantum number distinguishes among generations and what kind of excitation can yield a higher-generation quark or lepton? We discuss this issue in section 9.

(c) Structure within one Generation. Can we select a simple set of fundamental constituents, such that the entire pattern within one generation will be fully accounted for in a natural way? We propose an answer in section 10.

(d) Possible Compositeness of Gauge Bosons. In addition to the quarks, leptons and scalar Higgs particles some of the gauge bosons of the standard model may be composite. If they are, the corresponding interactions become residual and do not appear in the fundamental high-energy Lagrangian. Among the various gauge bosons, the most likely candidates for compositeness are the W and Z bosons. We discuss their possible composite nature in section 11.

8. The Problem of Energy Scales

We have already explained why composite quarks and leptons must be approximately massless with respect to their compositeness scale Λ_H. Such masslessness must emerge from a symmetry principle. The simplest symmetry which may prevent a fermion from acquiring a mass is a chiral symmetry[13]. We may therefore wish to look for a composite model with a chiral symmetry.

The chiral symmetry is essentially automatic if the fundamental fermions appearing in L_{NEW} are massless. However, the existence of a chiral symmetry in the fundamental Langrangian does not necessarily guarantee its preservation at the composite level. The chiral symmetry may be broken spontaneously, leaving no reason for massless fermions at the composite level.

Thus the necessary logical sequence of assumptions is as follows:

(i) The fundamental Lagrangian contains massless fermions and therefore possesses a chiral symmetry.

(ii) The full chiral symmetry or, at least, a chiral subsymmetry remains unbroken at the composite level.

(iii) The chiral symmetry of the effective Langragian containing the composite fermions, prevents the latter from gaining a mass. We have composite massless fermions.

Three questions imediately arise:

(a) If the new fundamental fermions are massless, why don't we observe them?

(b) What is the interaction which binds the fundamental fermions inside the composite quarks and leptons?

(c) If both fundamental and composite fermions are massless, what provides us with the necessary "compositeness scale" Λ_H?

All three questions can be immediately answered by one postulate[13], if we assume a new color-like force ("hypercolor") with a scale parameter Λ_H. All fundamental fermions carry hypercolor. They are confined by hypercolor forces of characteristic scale Λ_H into hypercolor-singlet composite fermions with an effective radius $r \sim \Lambda_H^{-1}$. The confined fundamental fermions cannot be experimentally observed. The binding and the scale are provided by the hypercolor gauge force.

The above scenario is an attractive framework for the construction of a composite model. However, it is crucial that the chiral symmetry or at least a chiral subgroup must remain unbroken at the composite level. This is not apriori impossible but it differs from the observed pattern of chiral symmetry breaking in QCD. No composite massless fermions emerge in QCD. The hypercolor situation must, for some reason, be different!

We now face a dilemma which stems from the following statements:

(i) We believe that in two-flavor massless QCD, the chiral symmetry is <u>completely broken</u>. No chiral subgroup remains intact.

(ii) If we neglect all interactions except hypercolor (all other interactions are probably much weaker at the Λ_H-scale), a hypercolor model with K fundamental massless fermions is <u>isomorphic</u> to K-flavor massless QCD.

(iii) In order to have massless composite fermions, some chiral symmetry should remain <u>unbroken</u> in the hypercolor case.

(iv) In no case do we have a <u>full dynamical understanding</u> of chiral symmetry and its breaking.

It is hard to reconcile statements (i), (ii), (iii), but no negative proof can be given. What are the logical possibilities?

(a) A resonable attitude, advocated by some theorists, is simply to declare that (i), (ii) and (iii) are inconsistent. In that case one should not continue to pursue our discussion beyond this point and the hypercolor idea should be abandoned. Perhaps this is true. Perhaps not.

(b) One way out is to consider a composite model in which left- and right-handed fermions have different transformation properties under the gauge group. Such a model is not isomorphic to QCD and statement (ii) does not apply to it. In such a model an $\bar{f}f$ condensate cannot break the chiral symmetry without breaking the original gauge symmetry. Two options are open: Either there is no condensation or the gauge symmetry breaks itself into a smaller subgroup. The first possibility has been studied by various authors and no realistic model was found. The second possibility is the interesting "tumbling" approach[14]. Here, again, no realistic model was found.However, the left-right symmetric classification may still be the correct solution.

(c) It is possible that the pattern of chiral symmetry-breaking depends on the number of flavors K. This could happen at least in two ways. There may be a phase transition at some K-value, K>3, leading to a different pattern for QCD and for a hypercolor theory with K>3 flavors of fundamental fermions. It is also possible that the general $SU(K)_L xSU(K)_R$ chiral symmetry always breaks, leaving a small conserved chiral subgroup which is trivial for K=2 but is nontrivial for large K. An example could be a discrete Z_K chiral group. A chiral Z_2 cannot protect any fermion from acquiring a mass. A chiral Z_4 or Z_6 can do it. There is no dynamical reason to expect any of these speculations to be true, but there are no complete arguments against them.

(d) Another possible speculation is that the presence of the color or electroweak interactions somehow influences the pattern of chiral symmetry breaking in a hypercolor scheme. This is the most obvious difference betweeen the hypercolor case and QCD. The simplest attitude would be to treat color and electroweak interactions as minor perturbations which cannot substantially change anything. However, subtle effects may occur. For instance, imagine a situation in which the chiral symmetry can break via $\bar{f}f$ or $\bar{f}f\bar{f}f$ condensates, the potential having two similar minima. A small perturbation could conceivably change the balance between the two minima, making the $\bar{f}f\bar{f}f$ condensate the likely one. At this point we may also add that the usual $N_c \to \infty$ argument for the breaking of chiral symmetry in QCD[15] does not necessarily remain valid if N_c/N_f is held fixed. In some composite models, such a fixed ratio may be a necessary requirement.

The above discussion can be summarized very simply: One can speculate about scenarios which provide the required pattern of chiral symmetry breaking for a composite model. All such scenarios are not supported by any decent dynamical arguments, but they cannot be ruled out.

Even if we succeed in producing a composite model with a chiral symmetry which is not completely broken, we still have to worry about the anomaly-matching condition[13], to which we now turn.

Let us assume that we have constructed a composite model of quarks and leptons based on an $SU(N)_H$ hypercolor gauge group and containing K fundamental massless fermions, all assigned to the N-dimensional representation of $SU(N)_H$. The underlying Lagrangian automatically possesses a global $SU(K)_L x SU(K)_R x U(1)$ symmetry. The U(1) factor is a vector charge counting the number of fermions. An additional axial U(1) factor is broken by instanton terms .

The model contains "flavor" triangle anomalies corresponding to products of three SU(K) currents or to products of two SU(K) currents and the U(1) current. Such anomalies are perfectly legitimate, since the $SU(K)_L x SU(K)_R x U(1)$ symmetry is not gauged. However, in the zero momentum limit, a given anomalous term can be exactly calculated both from the underlying theory and from the low-energy effective theory containing trhe composite particles. The results must be the same, thus imposing a severe constraint on the spectrum of composite particles.

If, in the underlying level, the anomaly does not vanish, the theory must produce massless composite particles[13)16)]. We may consider three logical possibilities:

(i) The chiral symmetry is not broken at all. There are composite massless fermions. Their contribution to each anomaly must be exactly equal to that of the fundamental massless fermions. Thus a severe constraint is imposed, connecting the fundamental fermions to the composite fermions. This is the famous 't Hooft condition[13)].

(ii) The chiral symmetry is completely broken. No chiral subsymmetry remains. The only massless composite particles are Goldstone bosons. Their contribution to the anomaly is equal to that of the fundamental fermions. However, since the Goldstone bosons have unknown couplings, the anomaly constraint can only be used in order to compute these couplings, leading to equations similar to the Goldberger-Trieman relations.

(iii) The chiral symmetry is broken, but a chiral subgroup remains conserved. The chiral subgroup may be continuous or discrete. In this case, massless Goldstone bosons must exist but massless fermions may also exist. The combined contributions of the massless composite bosons and fermions must balance the anomaly of the underlying theory.

The anomaly constraint is particularly powerful in case (i). It is not very useful in case (ii), but we are interested in massless composite fermions, and they do not occur in that case. Case (iii) allows composite massless fermions, and the anomaly constraint is somewhat less powerful.

We suspect that case (iii) is the most likely candidate for a realistic composite model. In particular, we may consider the interesting possibility of a continuous chiral symmetry in the original Lagrangian, broken into a discrete chiral symmetry at the composite level. Such situations arise naturally in simple unrealistic "toy" models[17)].

In case (iii) the massless composite fermions are accompanied by massless Goldstone bosons. Such bosons appear in a wide variety of composite models. They may escape detection because their Yukawa couplings to quarks and leptons are of order m_q/Λ_H and m_ℓ/Λ_H, respectively. We have mentioned the resulting experimental constraint in section 7.

We have gone here through an elaborate maze of difficulties, all stemming from the fundamental mismatch between the compositeness scale Λ_H and the masses of the composite objects. If the chiral symmetry hypothesis, together with a new hypercolor force, will not solve the problem, what other options do we have? The most likely possibility is some new fundamental force with some new features, not resembling any of the known interactions. Various ideas in that direction have been considered, including magnetic monopoles [18)], dimensional compactification[19)], nonlocal theories[20)], and quarks and leptons as massless supersymmetric Goldstone fermions[21)]. It is difficult to believe that the correct theory can be found without some experimental hints.

9. The Generation Puzzle

The existence of quark and lepton generations is, perhaps, the most striking experimental fact which guides us beyond the standard model. Let us first discuss the general problem, then turn to the possible description of generations within composite models of quarks and leptons.

We have three identical generations of quarks and leptons. The standard model does not contain any quantum number which distinguishes among the generations. Yet, we suspect that such a quantum number must exist. Three classes of solutions have been considered for a generation-labelling quantum number. In all cases we are looking for a symmetry which is already spontaneously broken at the stage of creating the fermion masses. The existence of Cabibbo mixing tells us that any "generation number" cannot remain exactly conserved.

The three possibilities are:

(i) A discrete generation label. A discrete symmetry is introduced, such that each generation obtains a different eingenvalue under the symmetry operation. It is necessary that, say, e, μ and τ have different eingenvalues. It is not necessary that e and u have the same eingenvalues, although it would be more elegant if they do. The scalar particles must have well-defined transformation properties under the discrete symmetry and the allowed Yukawa couplings are severely restricted by the symmetry. The mass matrix for analogous

states in different generations contains matrix elements contibuted by different scalar fields. If scalar fields with a non-vanishing generation number obtain vacuum expectation values, the discrete symmetry is broken and Cabibbo mixing is introduced. There is no real theoretical difficulty in describing the generations using a discrete symmetry. The only drawback of such an approach is the fact that all such discrete symmetries appear completely arbitrary and artificial.

(ii) A Continuous Global Symmetry. A variation on the same theme would be a continuous global symmetry under which each generation obtains a different eingenvalue. Here we face a serious difficulty: If the continuous symmetry is spontaneously broken, an unwanted Goldstone boson appears. Here, again, the ad hoc nature of the symmetry is usually unattractive.

(iii) A Gauged Generation Label. A third possibility which avoids the dangerous Goldstone boson is to consider an extra "horizontal" gauge group under which different generations form a gauge multiplet. The simplest example is a U(1) gauge symmetry but larger groups can be considered. The complications are: A severe anomaly constraint; the existence of a new gauge boson (or bosons); the danger of flavor changing neutral currents associated with "horizontal" gauge bosons.

Of the three possibilities, the discrete one is the only one which leads to no great difficulties. If we could find a discrete symmetry which is "natural" in the sense that its existence is caused or guaranteed by some other feature of the overall theory, it would be a likely candidate for a generation labeling scheme.

Another important property of a generation-labeling quantum number is its space-time nature.

Let us consider an operator under which e^{o} ,μ^{o} and τ^{o} possess the quantum numbers X_e, X_μ, X_τ. Here e^{o} is a massless electron appearing in the standard model Lagrangian. If the symmetry is vectorial, $X(e_L)=X(e_R)$ etc. If it is axial, $X(e_L)=-X(e_R)$ etc In the first case a scalar field with X=0 can induce diagonal mass terms for e , μ and τ.

The necessary X-values for scalar fields which contribute to mass-matrix elements are:

$$\begin{pmatrix} 0 & X_e-X_\mu & X_e-X_\tau \\ X_\mu-X_e & 0 & X_\mu-X_\tau \\ X_\tau-X_e & X_\tau-X_\mu & 0 \end{pmatrix}$$

On the other hand, if X is an axial quantum number, the three diagonal mass-matrix elements must be contributed by three different scalar fields.

The necessary values are:

$$\begin{pmatrix} 2X_e & X_e+X_\mu & X_e+X_\tau \\ X_e+X_\mu & 2X_\mu & X_\mu+X_\tau \\ X_e+X_\tau & X_\mu+X_\tau & 2X_\tau \end{pmatrix}$$

In view of the different scales of the masses of different generations, we believe that the axial option is preferable[22)]. In grand unified theories such as 0(10) only axial quantum numbers are possible, since e_L^- and e_L^+ belong to the same multiplet and must have the same eigenvalue for a given generation-labeling operator. Consequently, e_L^- and e_R^- must have opposite eigenvalues.

We conclude that, on quite general grounds, an attractive generation- labeling symmetry would be an <u>axial</u> <u>discrete</u> symmetry, provided that it is not artificially concocted. In composite models we do not have an arbitrary freedom for inventing such symmetries. The fundamental Lagrangian in such models is fully specified and all symmetries at the composite level must follow in one way or another from the properties of the theory.

Corresponding quarks and leptons in different generations must have the same $SU(3)_c xSU(2)xU(1)$ quantum numbers. They differ by some "generation number". All generations are approximately massless in comparison with the compositeness scale. Hence, they cannot be obtained by radial or orbital excitations of the first-generation "ground state". A possible excitation of a composite massless fermion which may lead to a different composite massless fermion is an excitation by one or more pairs of fermionic constituents. In a given composite model we should therefore investigate the possibility of constructing a system of preons and antipreons forming a scalar under the Lorentz group as well as under $SU(3)_c xSU(2)xU(1)$, but possessing a nonvanishing value of some "generation number". Such a system could be the difference between corresponding composite fermions in different generations.

An interesting possibility[23)]: In hypercolor composite models with K massless constituent fermions, we have a global $SU(K)_L xSU(K)_R xU(1)$ symmetry. An additional axial U(1) factor is broken. Hovewer, a discrete axial Z_{2K} symmetry always remain unbroken. Such a symmetry may serve as an adequate candidate for a generation number. It is an axial, discrete symmetry and it is not artificial at all. It exists in the theory, "waiting" to be used.

10. Structure Within One Generation

Each generation of quarks and leptons contains eight types of states. We list them in table 1, arranged in descending order of their electric charges.

An inspection of the table reveals a few features which cannot be explained within the standard model:

(i) The electric charges of the quarks and the leptons are quantized in a related way. Thus $Q(u)=\frac{2}{3}Q(e^+)$ and the hydrogen atom is exactly neutral. This is not at all guarranteed if the SU(2) and U(1) gauge interactions are unrelated.

(ii) The sum of the electric charges of all fermions vanishes. This is the famous condition for the vanishing of the triangle

	Color	Q	B-L
e^+	1	1	1
u	3	$\frac{2}{3}$	$\frac{1}{3}$
$\bar{d}$	$\bar{3}$	$\frac{1}{3}$	$-\frac{1}{3}$
ν_e	1	0	-1
$\bar{\nu}_e$	1	0	1
d	3	$-\frac{1}{3}$	$\frac{1}{3}$
$\bar{u}$	$\bar{3}$	$-\frac{2}{3}$	$-\frac{1}{3}$
e^-	1	-1	-1

Table 1: Fermions and antifermions of the first generation

anomaly in SU(2)xU(1). It is the only ingredient of the standard model which explicitly connects quarks and leptons and which tells us that a model with quarks and no leptons (or vice versa) is not renormalizable. A fermion (antifermion) is defined as a left-handed doublet (singlet) of the SU(2) gauge group in the standard model.

(iii) There are certain color-charge combinations which exist (and repeat themselves in higher generations). Other combinations do not exist. We have surprising correlations. For instance, 3Q (mod 3) is identical to the color triality, although no relation between charge and color is implied by the standard model.

(iv) The electric charge is limited by $|Q|\leq 1$.

The above regularities cannot be accidental. They must be explained by some theoretical structure which goes beyond the standard model, either by embedding the three different gauge groups in a larger simple group or by constructing all quarks and leptons from more fundamental constituents.

In a grand unified O(10) theory all of these regularities are beautifully accounted for by the structure of the group and its 16-dimensional spinor representation. The related charge quantization of quarks and leptons is guaranteed by the relationship between the SU(2) and U(1) couplings. The absence of anomalies is automatic in O(10). The color-charge correlations are dictated by the specific way in which $SU(3)_c xU(1)_{EM}$ is embedded in O(10). The $|Q|<1$ limitation is a property of the 16-dimensional multiplet. In a SU(5) grand-unified theory all of these features are also explained.

In a composite model one would hope to explain the pattern within one generation by a set of simple rules based on the properties of the fundamental fermions. Such rules should presumably be analogous to the quark model rules which neatly explain the repeated appearance of decuplets, octets and singlets of the flavor-SU(3), with no other representation appearing.

A particularly satisfactory explanation of all features of one generation is given in the rishon model[24], based on a hypercolor SU(3)-group. There we postulate two types of fermions: The T-rishon in a $(3,3)_{1/3}$ of $SU(3)_H xSU(3)_c xU(1)_{EM}$ and the V-rishon in a $(3,\bar{3})_0$. The structure of one generation is given by the hypercolor-singlet lowest-color states of three rishons and three antirishons. These states are listed in Table 2.

All four features which we mentioned at the beginning of this section, can be neatly explained in such a model:

(i) All electric charges are due to the T-rishon or $\bar{T}$-antirishon. Hence, quark and lepton charges obey simple ratios. A Hydrogen atom contains $e+u+u+d \equiv 4T+4\bar{T}+2V+2\bar{V}$. Its neutrality is trivially understood.

(ii) The quarks and leptons in one generation are $3(u+d)+e+\nu_e \equiv 6(\bar{T}+T+\bar{V}+V)$ Hence, their sum of electric charges (or any other additive quantum number) vanishes, and the standard-model anomaly cancellation is simply understood.

	Color	Q	B-L	Rishon Combination
e^+	1	1	1	TTT
u	3	$\frac{2}{3}$	$\frac{1}{3}$	TTV
$\bar{d}$	$\bar{3}$	$\frac{1}{3}$	$-\frac{1}{3}$	TVV
ν_e	1	0	-1	VVV
$\bar{\nu}_e$	1	0	1	$\bar{V}\bar{V}\bar{V}$
d	3	$-\frac{1}{3}$	$\frac{1}{3}$	$\bar{T}\bar{V}\bar{V}$
$\bar{u}$	$\bar{3}$	$-\frac{2}{3}$	$-\frac{1}{3}$	$\bar{T}\bar{T}\bar{V}$
e^-	1	-1	-1	$\bar{T}\bar{T}\bar{T}$

Table 2: Rishon model assignments of first generation fermions and antifermions.

(iii) The color-charge correlation is automatic. Replacing a T by a V corresponds to $\Delta Q=-\frac{1}{3}, \Delta(\text{triality})=-1$. Hence, the equality between 3Q(mod 3) and the color triality.

(iv) All quarks and leptons are three-rishon states, all combinations appear and the fundamental charge is $\frac{1}{3}$. Hence, the observed electric charges must correspond to all integer multiples of Q ranging between Q=+1 and Q=-1. $|Q|>1$ values cannot be obtained from three rishons or antirishons.

The rishon model, like all other composite models, suffers from several difficulties which we have discussed in detail elsewhere[24]. Its success in accounting for the structure within one generation is, however, impressive.

11. Composite Weak Bosons and Residual Weak Interaction

In composite models of quarks and leptons we usually face at least four types of gauge bosons: Hypergluons, gluons, photon, and weak bosons. Which of these bosons must be elementary? Can some of them be composite?

There is a certain confusion in discussing the possibility of composite gauge bosons. There are theories in which a certain fermion-antifermion pair of fields may appear to have some or all of the properties of a gauge boson. In some sense this is a composite gauge boson, but it appears in the same fundamental Lagrangian with the fermion fields and all other fields of the theory. Such a possibility is very interesting and some toy-models incorporating it have been constructed.

A different concept of a composite gauge boson is this: It does not appear at all in the fundamental Lagrangian of the theory (in the same way that other composite objects do not appear there). It does appear in the low-energy effective Lagrangian together with all other composite particles. Here we would like to study whether some of the gauge bosons may appear as composites in this sense.

In a hypercolor composite model, the hypergluon must clearly appear as a fundamental massless gauge field in the underlying Lagrangian. It will not appear at all in the low-energy Lagrangian. What about the gluon, photon and weak bosons? Consider first the massless gauge bosons (gluon and photon). If an exactly massless gauge boson appears in the low-energy effective Lagrangian, the Lagrangian must be exactly gauge invariant under the corresponding gauge group. This gauge invariance cannot be broken by higher dimension terms which are proportional to Λ_H^{-N}(N positive). If no small corrections of any kind are allowed to break the exact gauge invariance of the effective Lagrangian, it is essentially unavoidable that the original underlying Lagrangian also possesses the same local gauge symmetry. But in that case, it would probably contain the corresponding massless gauge bosons as fundamental fields. We therefore suspect that the gluon and the photon are not composite. They have the same status as the hypergluon in the underlying Lagrangian which must now be gauge invariant at least under $SU(N)_H \times SU(3)_c \times U(1)_{EM}$.

The above argument does not necessarily apply to the massive W and Z weak bosons. The weak gauge symmetry of the effective Lagrangian could be an approximate symmetry, broken by higher dimension terms which vanish as $\Lambda_H \to \infty$. It is conceivable that this approximate gauge symmetry is not fully present at the underlying level. In fact, the longitudinal components of W and Z are "born" from the scalar fields which are probably formed as condensates of the fundamental fermions. In some sense, at least the longitudinal W and Z must be composite in such a scheme.

The possibility of composite W and Z which do not appear in the underlying Lagrangian is extremely interesting. It leads to exciting consequences but also to serious difficulties.

Let us consider a process such as neutrino-neutrino scattering, in a model in which neutrinos are composite. The neutrino carries no hypercolor, color or electric charge. However, it must contain objects which possess at least hypercolor, possibly even color. At short distances of order Λ_H^{-1}, the two scattered neutrinos must experience a short range residual hypercolor force. Even if the nature of the binding force of the constituents is different, and even if it is not a color-like interaction, we still expect a short-range residual interaction between two composite neutrions. One way of parametrizing this short-range force, is in terms of the exchange of the lightest bosonic bound states of the same fundamental constituents. If W and Z are such composite states, their exchange may control the longest-distance component of the residual short-range force. In that case, the conventional weak interactions are identified as a residual effect[25)] of the original hypercolor force (or any other fundamental binding force inside the neutrino). The weak interaction are then eliminated from the list of fundamental interactions in the same way that hadronic interactions are residual color forces.

If the photon is fundamental but W and Z are composite, how do we understand the "unified" electroweak theory of the standard model?

In the standard model, the electromagnetic and weak interactions are not fully unified. Their relative strength remains an arbitrary parameter, related to $\sin^2\theta_W$. The standard model provides us with a clear mechanism for γ-Z^o mixing which could be somewhat analogous to the γ-ρ^o, γ-ω^o and γ-ϕ^o mixing of the old "vector dominance" idea. A major difference between the two situations stems from the different order of magnitude of the ρ^o and the Z^o direct couplings. Experimentally, $g^2_{\gamma\rho} \sim 1/300$ while $g^2_{\gamma Z} \sim 1/4$. How can we explain such a difference?

Many authors have discussed this issue during the last year[26]. Their consensus in that there is no difficulty in obtaining a $g_{\gamma Z}$ of the correct order of magnitude, provided that the spacing between Z^o and any higher composite boson is of the order of a TeV or so. That sets another bound or the compositeness scale Λ_H, for the case of a composite Z.

In fact, we may consider two extreme possibilities in theories with composite W and Z bosons:

(i) The compositeness energy scale is relatively low, say, between 1 TeV and 10 TeV. In such a case we may hope to observe experimental deviations from the standard model predictions for the properties of W and Z. Such a deviation could be seen in the W/Z mass ratio, W magnetic moment, small violations of universality etc. For $\Lambda_H \sim 1$ TeV, the effects could probably be detected within the next decade (but the model should cope with all the constraints of figure 1, a highly nontrivial requirement!)

(ii) The compositeness scale is $\Lambda_H >> M_W$ (say, above 100 TeV). In such a case L_{EFF} should be extremely close to L_{SM} and no experimental effects can be observed in the near future. In that case, however, we must face a new puzzle: If $\Lambda_H >> M_W$, what symmetry principle protects M_W and M_Z from obtaining higher order mass corrections which would lift them up to the order of magnitude of Λ_H? So far no one has proposed a convincing reason for a small mass of a composite W or Z boson. In the absence of such a reason, the possibility $\Lambda_H >> M_W$ appears to be unlikely.

Needless to say, there is a continuum of possible Λ_H-values. The lower Λ_H is, the sooner we can detect deviations from the standard model . For smaller Λ_H it is easier to understand the value of M_W but it is more difficult to construct a model which survives all the tests of figure 1.

We therefore conclude that the possibility of composite W and Z bosons is interesting, but serious difficulties exist, especially for large Λ_H-values.

12. Summary and Outlook

We conclude with four statements:

(i) At present, there is no experimental evidence for quark or lepton compositeness. Both high energy and low energy experiments should continue to search for such evidence. Experimentalists should probably ignore specific theoretical composite models and concentrate on pushing the various experimental limits until some new effects are discovered.

(ii) There is strong circumstantial evidence for the compositeness of quarks and leptons. In view of the lack of experimental clues, it is perhaps too early to demand a serious self-consistent theoretical model.

(iii) There is no satisfactory theoretical model of composite quarks and leptons. However, the models proposed so far contain many interesting new ideas. Each of these ideas should be investigated on its own merit, regardless of the detailed model which may have led to it. Several correct ingredients of the correct theory may already be with us now.

(iv) We have hardly began to investigate the subject of compositeness. It is almost unavoidable that the next decade or two will bring new experiments, new theoretical ideas as well as new difficulties for quark and lepton compositeness. The subject will certainly stay with us for a long time. It is not at all clear whether by the end of the century (and the millenium) we will know whether the electron is composite.

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Seminars

A CONFINING MODEL OF THE WEAK INTERACTIONS WITHOUT FUNDAMENTAL SCALARS

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Abstract

A strong-coupling, confining $SU(2)_L \times SU(2)_R$ gauge model for the weak interactions is proposed and analysed. The general, underlying idea is that weak interactions as measured at presently available energies are residual interactions among composite quarks and leptons much in the same way as the conventional strong interactions among composite hadrons. The model under consideration represents the simplest extension of the generic Abbott and Farhi model[1] dispensing with unwanted fundamental scalars and allowing for both, left-handed and right-handed composite quarks and leptons. The $SU(2)_L \times SU(2)_R$ scenario, with two different confinement scales $\Lambda_L \sim G_F^{-1/2}$ and $\Lambda_R > \Lambda_L$, is investigated within 't Hooft's framework[2]: The massless fundamental fermions (preons) introduce a chiral symmetry into the original gauge Lagrangian, which in the present case is taken to be $SU(4)_L \times SU(4)_R$, corresponding to three colors and electro-magnetic charge in the global limit $\alpha_{QCD}/\alpha_{L,R}$, $\alpha_{QED}/\alpha_{L,R} \to 0$. Following 't Hooft, this chiral symmetry is assumed to (partially) "survive" $SU(2)_L \times SU(2)_R$ confinement in the Wigner-Weyl mode, such that a set of fermionic bound states, the quarks and leptons, is kept "naturally" massless via 't Hooft's anomaly matching conditions. It turns out that these conditions single out one "ground-state" family of massless composite quarks and leptons for $1 \leq \Lambda_R/\Lambda_L < (\Lambda_R/\Lambda_L)_{crit}$, with both left-handed and right-handed ones having a radius $O(1/\Lambda_L)$. The spectrum is left-right symmetric, no exotics appear and none of the composite leptons involves colored preons, which is reassuring.

In order to gain more insight into the effective interactions for momenta $\Lambda_L < p \lesssim \Lambda_R$ and $p < \Lambda_L$, a two-step confinement analysis is performed for the idealized case $\Lambda_L \ll \Lambda_R$. Thus, in a first step, for the momentum range $\Lambda_L \ll p \lesssim \Lambda_R$, the global $SU(2)_L$ limit $\alpha_L/\alpha_R \to 0$ may be considered. The relevant chiral symmetry then becomes $SU(6) \supset SU(2)_L \times SU(4)_R$. Using the general anomaly conditions for the $SU(6)^3$ anomalies and the fact that all $SU(2)_R$ singlet composites must be mesons, it is shown that the $SU(6)$ chiral symmetry must be spontaneously broken for $p < \Lambda_R$. The associated Goldstone bosons enter

the effective Lagrangian for $\Lambda_L < p < \Lambda_R$ as new composite scalar fields and realize the SU(6) symmetry non-linearly. On this level, the important global SU(2) symmetry of weak isospin arises (approximately) in the Goldstone boson sector. It is well known[3] how to fix the form of the effective weak interaction Lagrangian at present energies $p < \Lambda_L$. After switching on the $SU(2)_L$ gauge forces in the second step the Abbott and Farhi Lagrangian is essentially recovered, however, with the (composite) Goldstone boson multiplett playing the role of the scalar doublet. Finally, it is argued that the weak interactions of the right-handed composite quarks and leptons are suppressed for $p < \Lambda_L$ by powers of Λ_L/Λ_R relative to those of the left-handed ones.

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CHIRAL SYMMETRY BREAKING AND CONDENSATES IN THE RISHON MODEL

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Abstract: There has to date been a good deal of theoretical work on models of composite quarks and leptons, although so far they have no experimental support. Despite this effort no elegant and realistic model has been shown to satisfy all the constraints. For example consider the Rishon model[1] which is based on the local gauge group $SU(3)_H \times SU(3)_C \times U(1)_{EM}$ with spin 1/2 Rishons $t = (\underline{3}, \underline{3}, 1/3)$ and $V = (\underline{3}, \underline{\overline{3}}, 0)$. As emphasised by Okun[2] this model apparently fails to satisfy 't Hooft's constraints[3] or the experimental bound on the proton lifetime $\tau_p \gtrsim 10^{30}$yr. Progress had been made on the latter problem with the proton decay rate sufficiently suppressed for a compositeness scale $\Lambda_H \gtrsim 10^7$ GeV [1]. In addition to this I would like to report on recent progress I have made with the former problem[4].

The first of Okun's problems is that in the limit $g_C \to 0$, $\alpha \to 0$ the Rishon model is analogous to six-flavour QCD, having a global symmetry $G_F = SU(6) \times SU(6) \times U(1)$. The Rishon model then faces the dilemma that if the chiral symmetry is spontaneously broken as in QCD in the diagonal mode $G_F \to SU(6) \times U(1)$ then no light fermions can result, whereas if it is spontaneously unbroken then the resulting anomaly equations possess no solution[3]. The only way I know of to overcome this dilemma is if the chiral symmetry is partially spontaneously broken in a new mode $G_F \to H_F$. But this suggestion opens up a Pandora's Box of questions such as: What is the value of H_F and the condensate? Why does the Rishon model differ from QCD? And what are the experimental implications? Let us see if we can't catch some of these devils.

By imposing certain basic requirements on H_F (such as H_F must contain $SU(3)_C \times U(1)_{EM}$ with the physical Rishons and must allow the anomaly constraint to be satisfied) we quickly arrive at a list of possible values of H_F. We can then supplement this by a corresponding list of candidate condensates which contain the singlet of H_F. In this way we are led to the following table of candidates (see table and ref. 4).

To select the true condensate and pattern of chiral symmetry breaking from the table we may employ a simple plausibility argument based on one gluon exchange. We can show that in fact the most attractive channel for massless QCD is a six fermion colour singlet and not the usual bilinear colour singlet channel (see ref. 4 for details). Assuming (a) Electromagnetism is preserved, and (b) Left-right symmetry is spontaneously violated in the usual way, leads to the rishon condensates, $\langle(V_R V_R)^3\rangle \ll \langle(V_R V_R)^2(V_L V_L)\rangle \ll \langle(V_L V_L)^2(V_R V_R)\rangle \ll \langle(V_L V_L)^3\rangle \sim \Lambda_H^3$. At energy $\sim\Lambda_H$ the condensate is then $\langle(V_L V_L)^3\rangle$ which corresponds to the last entry in the table. According to an argument due to Albright, Schrempp and Schrempp, if the top quark mass has exceeded a certain critical value then real life six flavour QCD would refuse to behave in this way but would undergo a phase transition into the diagonal mode, as observed experimentally.

It can be shown that the condensate $\langle(V_L V_L)^3\rangle$ and corresponding pattern of chiral symmetry breaking (see table) leads to an anomaly equation with the solution of three generations of massless quarks and leptons[4]. Further colour states of quarks in $\underline{6}$, $\underline{15}$ and leptons in $\underline{8}$, $\underline{10}$ are also produced but are expected to gain large dynamical masses of 10 - 100 GeV from the QCD vacuum. The phenomenology of these exotics is discussed in ref. 5. Goldstone bosons are also produced in abundance, although there seems little hope of detecting them unless Λ_H is very small[4].

Table

Candidate Condensate	H_F
$\langle(\bar{\Psi}\Psi)^2\rangle$	$SU(4)\times SU(4)\times U(1)_R$
$\langle(t_L\bar{t}_L V_L\bar{V}_L)\rangle$	$SU(3)\times SU(3)\times U(1)_R\times U(1)_{B-L}\times U(1)_Y$
$\langle(t_L\bar{t}_L)^2+(V_L\bar{V}_L)^2\rangle$	$SU(3)_C\times U(1)_R\times U(1)_{B-L}\times Z_4^Y$
$\langle(t_L\bar{t}_L)^3+(V_L\bar{V}_L)^3\rangle$	$SU(3)_C\times U(1)_R\times U(1)_{B-L}\times Z_6^Y$
$\langle(V_L V_L)^3\rangle$	$SU(3)\times SU(3)\times U(1)_{EM}\times Z_6^V\times Z_6^Y$

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STRUCTURE OF QUARK MASS MATRICES IN THE RISHON MODEL

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Abstract: The family structure of quarks and leptons is still awaiting a natural explanation. Neither the low-energy gauge theories such as the standard model nor grand unified theories have been able to answer the question: what distinguishes the generations? Phenomenologically, the only difference seems to reside in the fermion masses and weak mixing angles. Therefore, any explanation of the generation puzzle has to yield at least some insight into the structure of fermion mass matrices.

It is an obvious conjecture that a solution of the generation problem requires a more fundamental theory than presently known. An interesting proposal in this direction has recently been put forward by Harari and Seiberg[1)] in the framework of composite models. The global symmetry of a subconstituent model based on a confining hypercolour gauge theory usually contains a discrete symmetry which is the remnant of an axial $U(1)$ symmetry broken by the hypercolour anomaly. If the generations are postulated to differ by constituent pair excitations such a discrete symmetry could provide a generation labelling quantum number. This symmetry will then be spontaneously broken at the same scale as the low energy gauge group.

In a recent work[2)] I considered the rishon model[1)] as a specific example of such a mechanism to examine possible traces of the discrete horizontal symmetry (the cyclic group Z_{12} in this case) in the quark mass matrices. The rishon model is supposed to lead to a left-right symmetric low-energy gauge theory. Assuming that only a few effective scalar fields dominate the quark mass matrices upon spontaneous symmetry breaking one may look for constraints between quark masses and mixing angles.

Among the several possibilities[2)] one finds in this manner a single model with a definite prediction for the Cabibbo angle in terms of quark masses. There are two effective scalar fields in this case. The resulting quark mass matrices resemble those of a model encountered earlier in the context of a general classification of horizontal symmetries[3)]. Although the Kobayashi-Maskawa mixing angle Θ_2 is only weakly constrained by an approximate upper bound the model predicts with an accuracy of better than 9%

$$\sin \Theta_1 \simeq \sqrt{m_d / m_s}$$

in excellent agreement with experiment. The additional relation

$$\sin \Theta_3 \simeq \frac{m_s}{m_b} \sin \Theta_2$$

implies a very small mixing angle Θ_3 and will serve as the crucial test of the model.

All CP violating amplitudes due to the left-handed mixing matrix are proportional to $\sin \Theta_3$. If all other sources of CP violation are sufficiently suppressed by a large enough mass of the right-handed vector bosons the predicted magnitude of Θ_3 gives a natural explanation of the smallness of CP violation.

In conclusion, the hypothesis that generations differ by constituent pair excitations and may be distinguished by a discrete axial symmetry seems promising and encourages further study of composite models.

References

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