12/18/84 MJR (2) I-18587

DR-0672-6

LBL-18571
UC-34D
CONF-8406190--

Proceedings of the Workshop on Electroweak Symmetry Breaking

June 4 - 22, 1984

October 1984

Lawrence Berkeley Laboratory
University of California
Berkeley, California 94720

Prepared for the U.S. Department of Energy under Contract DE-AC03-76SF00098
and the National Science Foundation under Research Grant No. PHY-84-01063

LBL-18571

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WORKSHOP ON ELECTROWEAK SYMMETRY BREAKING†

Held at Lawrence Berkeley Laboratory, June 4 - 22, 1984

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† This Workshop supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098 and in part by the National Science Foundation under Research Grant No. PHY-84-01063.

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PREFACE

A theoretical workshop on electroweak symmetry breaking at the Superconducting Supercollider was held at Lawrence Berkeley Laboratory, June 4 - 22, 1984. The purpose of the workshop was to focus theoretical attention on the ways in which experimentation at the SSC could reveal manifestations of the phenomenon responsible for electroweak symmetry breaking.

This issue represents, at present, the most compelling scientific argument for the need to explore the energy region to be made accessible by the SSC, and a major aim of the workshop was to involve a broad cross section of particle theorists in the ongoing process of sharpening the requirements for both accelerator and detector design that will ensure detection and identification of meaningful signals, whatever form the electroweak symmetry breaking phenomenon should actually take.

The workshop was scheduled so as to immediately precede the DPF Workshop at Snowmass, with sufficient overlap in participants that the results of the theoretical workshop could be efficiently communicated to the participants at Snowmass.

The first two days of the workshop were devoted to summaries of the conclusions of earlier working groups both in the U.S. and abroad. The remainder of the three week session consisted of a single seminar per day on a topical issue with the rest of the time being spent on individual work and discussions among participants, mainly through working groups that were organized around various alternative scenarios for electroweak symmetry breaking. The discussions were lively; many new ideas were generated and older ideas were pursued in some depth. We are grateful to all the participants who contributed to a lively and productive session.

We are indebted to Susan Fidelman for her superb and smooth management before, during and after the workshop, and to Brenda Allen-Clearlake for her invaluable assistance throughout. We further enjoyed the support and clerical assistance of Betty Sublett, Luanne Neumann and Mary Gorman, during the workshop and in preparing the proceedings. We wish to thank Orlando Alvarez for the social organization and all the members of the LBL theory group, most especially our willing graduate students, for their participation, support and material assistance.

We are grateful to Peggy Little for her assistancee with organization and to J.D. Jackson and David Shirley for the support of the Physics Division and of the Laboratory.

This Workshop was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098 and in part by the National Science Foundation under Research Grant No. PHY-84-01063

We hope that this workshop has achieved the objective of drawing the theoretical community into substantial involvement with the national SSC effort and that these proceedings will serve as a stimulus to further study within the high energy physics community.

Thomas Appelquist
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OVERVIEW OF THE LBL WORKSHOP

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This overview of the LBL Workshop on Electro-Weak Symmetry Breaking at the SSC is provided as a brief guide to the more detailed papers which follow. It is based on a talk given at the opening of the Snowmass workshop and is intended mainly for a non-theoretical audience. The comments expressed in this overview are mine and do not necessarily represent the view of the majority of participants.

The LBL workshop had a loose organizational structure; theorists often work best in a freewheeling environment. Several subgroups were formed to attack some of the more important theoretical problems relevant to physics at the SSC.

These groups discussed

1. Non-Standard Higgs
2. Intermediate Mass Higgs
3. Non-Standard Higgs Bosons
4. Strongly Interacting W's and Z's
5. Supersymmetry
6. Interpretation of Unusual Events from CERN and DESY
7. Compositeness

The reports of the the subgroup chairmen can be found following this article.

In addition, a number of seminars were held with two purposes. Firstly, talks were given at the opening of the workshop to acquaint participants with the current state of affairs. Only two of these talks are included in the Proceedings. The one by M.K. Gaillard which acts as an introduction to the workshop and sets out its goals and the one by J. Ellis on the status of European Facilities. Other talks were given by Jay Marx (reporting on the SSC Reference Design, Ref. 1), Stu Loken (reporting on the activities of the PSSC, Ref. 2), Bob Cahn (reporting on the Chicago Workshop, Ref. 3), Dave Jackson (reporting on the ICFA meeting held in Japan in late May, Ref. 4), and Chris Quigg (reporting on some conclusions of EHLQ, Ref. 5). These talks all cover material which is readily available elsewhere and are not included in these proceedings.

Secondly, seminars were held on topical subjects relevant to the SSC. These seminars are written up in these proceedings and will not be discussed further in this summary.

The group looking at the new and unusual events from CERN and DESY attempted to draw together the information from these odd events, summarized in Table 1 of the report (L. Hall, et al., Ref. 6). They looked at and evaluated the various theoretical proposals for these events. In particular, they tried to look for common features of the events and to search for explanations which are capable of dealing with more than one type of event.

I will first comment on the events from the UA1 and UA2 collaborations with missing transverse energy (Ref. 7). The invariant mass of the system with the large transverse momentum in these events is quite large, the transverse mass of the jets, leptons, "photons" and missing transverse momentum provides a lower bound on the invariant mass of the parton-parton system which generates the final state. This transverse mass is given for the UA1 events in Table 3 of the report which shows that it is of order 100 GeV. The UA2 events with leptons and jets have a total invariant mass of order 150 GeV (roughly the same as the jet-jet mass peak also claimed by UA2) if they are interpreted as final states of W plus jet. The typical mass of the parton-parton system in either case is of order 150 GeV. In order to have a chance of explaining the rate of these events, the cross-section in the parton-parton system must be due to strong interactions, and the particles in the final state will be strongly interacting. This simple observation follows from the typical values of parton momenta which are required and the behavior of the parton distribution functions. The "old" physics background to these events is from W/Z + Jet final states which has a partonic cross-section of order $\alpha_s\alpha$ and is consequently much too small to explain the signal given that the W/Z has to decay leptonically in order to generate the missing transverse momentum.

The only candidate proposed by theorists before the events were found is supersymmetry which generates events with jets and missing transverse momentum. The immediate difficulty is that one naively expects events with several jets, two from events with squark pairs and four from events with gluino pairs. Some of these jets will either be soft or will coalese so that the physical final state could consist of fewer than the expected number of jets. Figure 6 of Hall et al. shows this effect. However, a problem remains with the supersymmetric interpretation of the events; the multiplicity of particles within, and the invariant mass of, the jets in the UA1 sample is lower than is expected. Also, although the number of events is small, the transverse momentum distribution is rather flat in the UA1 case, another problem which is not obviously explained by supersymmetry.

All other candidate explanations were invented after the events were found; they are summarized in Table 8 of the subgroup report. None of them is very compelling, and some fail to explain the quantitative features of the events.

The events interpreted as Z decay to $e^+e^-\gamma$ or $\mu^+\mu^-\gamma$ (Ref. 8) are summarized in Table 2 of the subgroup report. The most natural explanation is in terms of Bremsstrahlung. it gives a Dalitz distribution which agrees with the data but unfortunately has a rate which is a factor of fifty too small. None of the explanations discussed by the group are particularly compelling and most fail to explain the Dalitz distribution.

The other odd events found at PETRA (Ref. 9) are at a rather small rate and it is not completely clear that they are inconsistant with background. There is only one event from CELLO of the type $\mu^+\mu^- + 2$ jets. Again they were not anticipated by theorists although a number of explanations were discussed. The conclusion of this group was that one should wait and see. If the UA1 events are real the next run should settle the issue and provide more detail which would enable some of the explanations to be excluded.

Higgs Rates and Signals

In the minimal Glashow-Salam-Weinberg[10] model there is only one particle remaining to be found; the scalar Higgs boson (H). As discussed by C. Quigg,[26] the only unknown quantity surrounding this particle is its mass, all its couplings are fixed. It couples to W^+W^- with strength $eM_W/\sin\theta_W$ to Z^0Z^0 with strength $2eM_Z/\sin 2\theta_W$ and to a fermion anti-fermion pair (each with mass m_f) with strength $em_f/2M_W \sin\theta_W$. The factor m_f/M_W ensures that its coupling to fermions is very weak (with the exception of the top quark and any new heavier generations). Consequently H cannot be produced directly in e^+e^- annihilation.

The Z^0Z^0H vertex can be exploited by using the decay $Z \to H\mu^+\mu^-$ (or He^+e^-) where the $\mu^+\mu^-$ (e^+e^-) pair came from the decay of a virtual Z^0.[11] The ratio of widths $\Gamma(Z^0 \to e^+e^-H)/\Gamma(Z^0 \to e^+e^-)$ is shown as a function of the Higgs mass in Fig. 11.[12] The Z^0 can be formed in e^+e^- and at a luminosity of 10^{30} cm^{-2} sec^{-1} a value of 10^{-3} for the ratio of widths corresponds to approximately one event every 10 days. The signal is very clear, one looks for a peak in the missing mass recoiling against the lepton pair. There is a background from the production of two heavy quarks where they both decay leptonically. By imposing isolation cuts on the leptons this background can be controlled. In the e^+e^- channel, there is an additional background at low m_H from $e^+e^- \to e^+e^-$ + hadrons which imposes a lower limit of about 8 GeV upon the Higgs mass that one can see.[12]

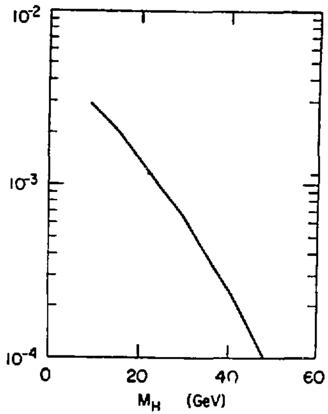

Fig. 11.
The ratio $\Gamma(Z \to e^+e^-H)/\Gamma(Z \to \mu^+\mu^-)$ vs m_H.

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One other area of theoretical speculation which was discussed by one of the groups concerns the issue of compositeness of quarks and leptons (Ref. 23). Motivated by the apparent arbitrariness of quark and lepton masses, some theorists have speculated that they may not be elementary but rather composed of some "preons". No semi-realistic model has yet emerged. Some signals for composite structure at the SSC are discussed in Ref. 5, but some work was begun on the question "What happens if the scale of compositness is 1 TeV and hence can be crossed at the SSC?" There was also some discussion of the use of polarization for disentangling the structure of any interactions caused by compositeness.

Finally, I would like to thank all those who helped to make this workshop so enjoyable and fruitful. My co-organizers Mary K. Gaillard and Tom Appelquist for their valuable work in planning and organization, the working group chairmen for summarizing clearly and concisely the work of their groups, Susan Fidelman for dealing with the day to day running of the workshop and for assisting in the preparation of these Proceedings, and finally to all the particpants for their enthusiasm and hard work.

This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098 and in part by the National Science Foundation under Research Grant No. PHY-84-01063.

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INTRODUCTION TO THE WORKSHOP: ELECTROWEAK SYMMETRY BREAKING AT THE TeV SCALE

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Introduction

As viewed from today's perspective, electroweak symmetry breaking is both the central issue to be addressed by physics in the TeV region, and the most compelling argument for the need to explore that region. While the picture may change considerably over the next decade, it seems reasonable to focus theoretical attention on this issue which is in fact very broad in terms of its possible ramifications. Such a concerted effort can help to sharpen the scientific case for the SSC and provide fresh theoretical input to the ongoing series of workshops and studies aimed at forming a concensus on a choice of SSC design parameters.

To set the mood of the workshop I will review briefly the physics to be explored prior to the SSC as well as the motivations for exploration of the TeV region for hard collisions. I will follow with an example of a possible scenario for the first manifestation of electroweak symmetry breaking at the SSC.

State of the Standard Model

In a sense we are reaching the end of an era in the study of electroweak interactions, which are by now well understood as being described by the Lagrangian of a renormalizable, spontaneously broken gauge theory. The list of successes and precise, quantitative predictions is impressive. The attempt to understand the four-fermion charged current interactions in terms of a renormalizable theory culminated in the prediction and subsequent observation of neutral current phenomena as well as of the W and Z bosons, with precise predictions for their masses and other properties. Within the same context, the presence of strangeness changing charged currents, together with the observed strong suppression of their neutral current counterparts, led to the prediction and subsequent observation of charmed particles with precisely defined weak couplings and approximate estimates of their masses and other properties. The discovery of the τ-lepton implied, again within the context of a renormalizable theory, the existence of the (t,b) quark doublet; indeed, the entire third family of quarks and leptons had been anticipated in attempts to incorporate CP violation into the theory. Of this family, the t quark still awaits confirmation, as does direct evidence for the ν_τ.

There are hints from CERN that we may already be embarking on a new era. Possible interpretations of the zoo of intriguing $Sp\bar{p}S$ events were the focus of one of the workshop study groups. Whether any of these events really reflects new physics, as opposed to the traditional hiccups which tend to accompany the opening up of a new domain of experimentation, should be settled by the coming generation of facilities: an upgraded $Sp\bar{p}S$, TeV I, SLC, LEP and HERA.

In any event these facilities will provide a thorough testing of the standard model, including precision measurements of the W and Z masses and widths. In particular, the parameter $\rho = m_W/m_Z$ is sensitive to some high mass phenomena through radiative corrections. The high yield of Z's at SLC and LEP will permit searches for rare decays. About 5000 $W \rightarrow t\bar{b}$ events should be produced at TeV I for an integrated luminosity of $10^{37}cm^{-2}$, which should allow a rough check of GIM-KM unitarity.

An important aspect of the standard electroweak theory which has not yet been tested is the complex of trilinear and quadrilinear self-couplings of gauge bosons. Measurements of $e^+e^- \rightarrow W^+W^-$ at LEP II and of $qq \rightarrow WW$, WZ and $W\gamma$ at pp colliders will provide rough checks of the three vector boson coupling strengths. For LEP running somewhat below the two-W threshold, the process $e^+e^- \rightarrow eW\nu$ should allow a similar rough check[1] of the magnetic moment of the W.

It is possible that the "observed" electroweak gauge group $SU(2)_L \times U(1)$ is embedded in a larger group: TeV I should be able to probe for additional heavy Z's with masses up to 500 GeV if they have couplings to quarks of standard strength. LEP can search for very heavy Z's through propagator effects, while HERA will be sensitive to heavy W's as well, and also to the presence of right handed couplings for charged current reactions.

At lower energies, the copious sources of kaons and/or B mesons to be provided by CESR, TeV II and the AGS will help to pin down the parameters of the KM matrix, and in particular those governing CP violation. Searches for rare decays will also provide probes of higher mass scales.

Shelly's Plumber

There is still, of course, an important missing link in our present picture: the Higgs particle(s) or some other manifestation of electroweak symmetry breaking.

If the standard Higgs has a mass $m_H < 2\, m_W$ its decay will probably be too indistinctive to allow detection at a hadron collider. A possible exception is the case in which the decay $H \rightarrow t\bar{t}$ is kinematically forbidden; then there is a small window of possible Higgs masses for which the decay $H \rightarrow W + f\bar{f}'$ may have a substantial branching ratio[2]: $B(H \rightarrow W + f\bar{f}') = (6 \text{ to } 60)\%$ for $m_H = (120 \text{ to } 150)$ GeV. Generally, a Higgs with mass below the two-W threshold can most easily be detected using missing mass techniques in e^+e^- annihilation. For
$m_H \leq 40$ GeV, a standard Higgs could be found at SLC or LEP in Z^0 decay, or, depending on the top quark mass (and possibly at TRISTAN), via the decay of toponium into $H + \gamma$. LEP II can probe for a Higgs with mass $m_H \leq (2E_{beam} - m_Z)$ via the process $e^+e^- \rightarrow Z^0 + H$.

In the event that such a "light" standard Higgs turns up at the next generation of facilities, will the final chapter of weak interactions come to a close? There is strong reason to suspect that the Higgs phenomenon represents only the tip of the iceberg, and that qualitatively new physics must be involved. The deeper issue, commonly known as the gauge hierarchy problem, is the puzzle as to why the W and Z masses are so small in the presence of large scale parameters such as the hypothesized grand unification scale or the Planck scale. In the context of a weakly coupled renormalizable theory, such "light" gauge bosons require similarly "light" scalar bosons, but scalar masses are highly unstable against radiative corrections.

There are of course, other hierarchy problems, in particular large ratios among fermion masses, which by rights should all be of the same order as the W and Z masses since they are governed by the same symmetry breaking scale parameter. This issue has received less attention, probably because we haven't yet understood how to sensibly formulate the question. In the case of the usual gauge hierarchy we know how to ask the question and even how to answer it. The three most popular answers are listed below.

Technicolor. A scalar particle may be kept light by a global chiral symmetry which is broken spontaneously by a condensate of massless technifermions, characteried by a scale parameter Λ_T

$$<\bar{\psi}_T\psi_T> \sim \Lambda_T^3,$$

which is the scale at which the presumed asymptotically free technicolor interactions become strong. If ψ_T is an electroweak gauge non-singlet, the condensate also breaks the electroweak symmetry, giving the observed W and Z masses for $\Lambda_T \sim (\sqrt{2}\, G_F)^{-\frac{1}{2}} \simeq 250$ GeV. The exactly massless goldstone bosons of spontaneously broken chiral SU(2) are eaten by the $W^\pm$ and Z to become their longitudinal components. This hypothesis predicts a rich spectrum of technihadrons with masses in the TeV region. For ordinary fermions to acquire masses, the theory must be extended in a way which generally leads to the prediction of additional pseudo-goldstone bosons that are considerably lighter. At present no phenomenologically viable (nor grand unifiable) model for technicolor exists, but the idea is sufficiently attractive to warrant attention.

Supersymmetry. Since chiral symmetries control fermion masses, scalar masses can be controlled if they are superpartners of chiral fermions. In practice, the gauge hierarchy is usually implemented in supersymmetric models by exploiting instead the "non-renormalization" property of supersymmetry which protects scalar masses against large radiative corrections. The electroweak breaking scale is related to the supersymmetry breaking scale which is generally adjusted by hand. Supersymmetry is motivated by other arguments as well, and may play a vital role in the ultimate connection between gravity and the observed gauge interactions. In this case, it may or may not also provide the mechanism for stabilizing the Higgs mass. If it does, as in most popular models, one expects to discover lots of superpartners of quarks and gauge bosons with masses below a TeV-except in the perhaps perverse but logically possible event that the Higgs mass is greater than a TeV and/or that the Higgs sector "sees" supersymmetry breaking only through radiative corrections, in which case many squarks, sleptons and gauginos could have masses larger than the Higgs mass by an order of magnitude or more.

Compositeness. A third possibility is that the standard model is in fact an effective theory for describing composite quark and lepton (and gauge?) fields which appear point-like at energies well below the inverse radius of compositeness. Perturbative calculations break down for virtual momenta higher than this inverse radius which provides an effective cut-off that stabilizes the Higgs mass, or, equivalently, its vacuum expectation value. Present data already suggest that the scale of compositeness exceeds a TeV; if it is indeed the Higgs mass stabilizer, it must not exceed a few TeV. Just as for supersymmetry, it is possible that ordinary particles are composite on a scale which is unrelated to the weak interactions. Signals of compositeness include new interactions: effective four fermion couplings with strength characterized by the squared radius of compositeness, and new particles: excited states of quarks and leptons, in particular color non-singlet quarks that may be quasi-stable.

Since none of the above models is sufficiently well constrained and/or well formulated to allow quantitative mass predictions, searches in any available mass range are of interest. The Sp$\bar{p}$S, TeV I, SLC, LEP and HERA complex of facilities should allow probes for supersymmetric particles up to mass scales of about 100 GeV, and for compositeness up to a scale of about 6 TeV The SSC will be able to push these scales up considerably,[3] and should be able to weed out technicolor if that is the mechanism which sets the electroweak breaking scale.

Physics at the TeV Scale

We saw that attempts to understand the relatively small scale of electroweak symmetry breaking tend to suggest the existence of new particles or new phenomena. There are various other hints from both particle physics and cosmology that new physics should appear at scales well below those associated with grand unification or gravity.

One is the non-observation of the decay $p \to \pi e$ with a partial life-time as estimated in the minimal SU(5) model. A possibility is that the unification idea is totally wrong, but then we must abandon our present understanding of the value of the weak mixing angle and the observed nucleon to photon density ratio. Furthermore, since the observed spectrum of fermions is indeed an SU(5) spectrum, it is difficult to imagine that the ultimate unification scheme does not embed SU(5) at some level. An alternative possibility is that the unification scale is much higher than standard model calculations predict. This has the possibly attractive feature that the unification scale and the Planck scale are essentially the same, and the danger that proton decay may be unobservable altogether, closing an important experimental window on unification. In any event the latter interpretation requires some new particles - if only extra generations of quarks and leptons - with masses above present laboratory sensitivity, but well below the unification scale.

Attempts to reconcile the density fluctuations required for galaxy formation with the observed degree of homogeneity in the microwave radiation background invoke particles which were thermally decoupled from photons by the time of galaxy formation; candidates have included massive neutrinos, gravitinos, axions, etc. Models for an inflationary scenario compatible with both astrophysical observation and particle phenomenology may require additional new fields. Finally, we are at present totally in the dark on the complex of issues including the spectrum of fermion masses, the Cabibbo-KM mixing angles and CP violation.

Whether any of the above issues is related to electroweak symmetry breaking is an open question, and it is not possible to pin down a mass scale at which their resolution should be revealed. However it seems likely that resolving the issue of electroweak symmetry breaking should point us in a clearer direction towards answers to some of the other questions.

Why is the TeV scale an immediate target? The mass of the $W^\pm$ was successfully predicted by a simple formula:

$$m_W = [\pi\alpha/\sqrt{2}\, G_F]^{\frac{1}{2}}, \qquad (1)$$

where α is the fine structure constant of QED and G_F is the Fermi constant. In the standard model the Higgs mass is predicted by a similarly simple formula:

$$m_H = [8\pi\alpha_H/\sqrt{2}\, G_F]^{\frac{1}{2}}, \qquad (2)$$

with the unfortunate difference that the "fine structure constant"

$$\alpha_H = \lambda^2/4\pi \qquad (3)$$

appearing in (2) is the expansion parameter for the perturbative theory of scalar interactions, and is itself unknown, except for the requirement that the observed vacuum be stable against radiative corrections, suggesting $m_H \geq 10$ GeV.

On the other hand, should the Higgs mass exceed a TeV, Eq. (2) implies that the parameter α_H exceeds unity and that perturbation theory is inapplicable to the scalar sector. One might worry that this would render fortuitous the successful predictions of the standard model, but it has been shown[4,5] that ordinary physics is highly screened from strong interaction effects in the scalar sector, just as, for example, anomalous magnetic moments of leptons are very insensitive to strong interactions in the hadron sector.

The relevance of all this to a supercollider is that a strongly interacting scalar sector cannot remain screened at very high energies. The reason is that α_H also governs the strength of the self couplings of the $W^\pm$ and Z through their longitudinal components, acquired by absorbing three of the scalars that together with the physical Higgs particle form a complex doublet of the weak SU(2) gauge group.

In the 1960's it was argued correctly, on the basis of unitarity of the S-marix and the observed fermi couplings, that exploration of energy scales up to 600 GeV would necessarily reveal either strong parity violation or qualitatively new physics associated with the underlying structure of the weak interactions. The latter we now recognize as the W and Z of the standard model, and the upcoming generation of experimental facilities is well adapted to study their properties as discussed above.

An analogous argument, based on unitarity and the observed electroweak couplings of the standard model, leads to the conclusion[4,6] that either W's and Z's will develop strong interactions at effective c.m. energies of a few TeV or qualitatively new physics, related to the mechanism for electroweak symmetry breaking, will emerge. The latter may or may not take the form of a standard model Higgs or one of the richer scenarios described above.

The questions addressed at the workshop included:

What form might the new physics take?
What might be its experimental signals?

The purpose of asking questions such as these is to sharpen the requirements on energy and luminosity, and suggest directions for detector development, with the aim of assuring maximum accessibility to the physics of the TeV region, whatever form it might take. The physics reach for various choices of machine parameters and for the standard "bellwether" scenarios has been extensively treated by EHLQ.[3] Our purpose here was not to rehash the bellwethers, but rather to generate new ideas and new perspectives on old ideas.

An Example: A Minimal Scenario

Suppose that the study of hard collisons up to the TeV scale for effective c.m. energies reveals neither a standard model Higgs nor any obvious variation thereof. Suppose further that experimental data continues to conform to the standard model, so that observed electroweak physics is described by the GWS Lagrangian which in a renormalizable gauge takes the form:

$$\mathcal{L} = \mathcal{L}_{\mathrm{gauge,\,matter}}(g, G_{\mathrm{Yuk}}, m_{W,Z}, m_f, \xi) - V(\phi, H), \qquad (4)$$

where the first term includes mass and kinetic energy terms and gauge and Yukawa couplings, and ξ is a gauge parameter. The second term is the scalar potential; in the standard model:

$$V(\phi,H) = m_H^2 H^2/2 + m_H^2 H(|\phi|^2 + H^2)/2v + m_H^2(|\phi|^2+H^2)^2/8v^2 \qquad (5)$$

where H is the physical Higgs particle, $\phi = (w^+, z, w^-)$ are the unphysical scalars absorbed as longitudinal components of (W^+, Z, W^-) in the unitary gauge, and $v = 250$ GeV is the usual scalar vacuum expectation value. The relevance of the potential (5) to TeV physics is that S-matrix elements with external w's and z's calculated from the Lagrangian (4) are equivalent[6-8] to the S-matrix elements for external longitudinally polarized W's and Z's, up to corrections of order m_w/E.

The potential (5) is characterized by a single unknown parameter, the physical Higgs mass m_H. As discussed above, for a Higgs mass of a TeV or more, (5) describes a strongly interacting system. However, one can try to exploit[9] the property[5] that the potential V is invariant under non linear transformations among scalars, whose generators satisfy the algebra of chiral $SU(2) \times SU(2)$. The first term in (4) contains the weak couplings of $w^\pm$ and z to fermions through scalar and pseudoscalar densities and to transversely polarized W's and Z's through vector and axial currents that are conserved up to corrections of order of the weak couplings and the W, Z squared mass. The situation is analogous to that of low energy hadron physics where an (approximately) chiral SU(2) invariant strongly interacting system couples to leptons through (partially) conserved axial and vector currents. Here the longitudinal vector bosons $W_L, Z_L \simeq w, z$ play the role of the pions of hadron physics. These general features are moreover not specific to the standard model; the situation in a minimal technicolor scenario is identical, and any scenario where no symmetry breaking phenomenon is manifested below the TeV scale is expected to display similar properties.

Ideally, then, one would like to understand the dynamics of a strongly coupled σ-model, just as for pion chiral dynamics. The strongly interacting limit of the minimal Higgs model has been analyzed for the presence of bound states or resonances[6,10] Regge poles,[11] or skyrmions.[12] Recently Einhorn[13] found that the leading N behavior in a 1/N expansion (here N=2!) for a chiral SU(N) scalar sector suggests that there must be a J=0 scalar state (which might as well be called the Higgs particle) with a mass of at most a few hundred GeV.

What I shall discuss here is a more modest approach adopted by Mike Chanowitz and myself:[9] given that the longitudinally polarized gauge bosons W_L, Z_L develop strong interactions, how can we experimentally study this strongly interacting system? First, we must produce a system of two or more W_L, Z_L, which is not entirely trivial, as W_L and Z_L couplings to quarks are suppressed up to corrections of order m_q/m_W or m_W/E_W. In addition we are working in a regime where perturbation theory is not applicable. However, the replacement $W_L, Z_L \rightarrow w, z + O(m_W/E)$ in S-matrix elements and the chiral symmetry of strong w, z interactions allow us to determine the W_L, Z_L couplings near threshold through soft pion theorems. The threshold behavior obtained in this way is given precisely[8] by the Born approximation to the GWS Lagrangrian (4), (5). The resulting amplitudes for multiple W_L and Z_L production are roughly characterized by a factor E/v for each emitted W_L or Z_L. In the limit $m_H \rightarrow \infty$, v is the only scale parameter of the system, so simple scaling arguments imply that the Born approximation must be valid for some energy range between multi-W production thresholds and that energy at which a damping scale (m_H?, Λ_T?) sets in to restore unitarity in the s-wave scattering channel, as it must. This scale will presumably be signaled by resonance production or similar phenomena. Whatever the dynamics of such a strongly coupled system should turn out to be, it should be characterized by events with a high multiplicity of W's and Z's.

We therefore considered various mechanisms for the production of a system of two or more W_L,Z_L and made multiplicity estimates based on the E/v scaling law, which is equivalent to the Born approximation that automatically satisfies the current algebra constraints. We considered two extreme cases: a) the Higgs mass sits at its "unitarity limit" value[6] of a TeV, where it becomes so broad that establishment of a resonance in the WW and ZZ systems may be problematic, and
b) $1\ \mathrm{TeV} \leq \sqrt{s_{WW}} << m_H$ (in this case tree unitarity breaks down in the s-wave channel for $\sqrt{s_{WW}} = 1.8$ TeV).

The most copious source of longitudinally polarized W's and Z's turns out to be the analogue of the Cahn-Dawson mechanism[14] for Higgs production (Fig. 1). At first sight, this would appear to give a negligible contribution because the W_L coupling to light quarks is suppressed by a factor m_W/E_W in amplitude at each qq'W vertex. However this factor is exactly compensated for by the fact that, as opposed to the case for transversely polarized vector bosons, longitudinal W emission does not vanish in the forward direction.[9] We estimated the total yield from this mechanism using parameterizations of the total $W_L W_L$ cross section adjusted to reproduce the correct threshold behavior, an asymptotic logarithmic energy dependence, with or without a broad (Higgs) resonance in the energy region accessible to the SSC. In all cases, the yield of events including a pair of Z_L's is expected to exceed[15] the Z-pair yield[3] from conventional gauge interactions for sufficiently high sub-energies for the vector boson system. In contrast, the light $q\bar{q}$ annihilation channel, which can produce a significant yield of pairs of longitudinally polarized bosons only in a pure J=1 state, is dominated by pair production of transversely polarized vector mesons. In either case the Z/W production ratio will be enhanced in the presence of important strong interaction effects.

Figure 1

As the potential (5) conserves "parity", with w and z defined as parity-odd, the mechanism of Fig. 1 will produce only even numbers of vector bosons. To lowest order in the weak gauge coupling constant, the dominant mechanism for production of an odd number of w, z (or W_L, Z_L) is that of Fig. 2. For case a), $m_H \simeq 1$ TeV, the cross section is dominated by on-shell Higgs production and decay; for case b), $1\ \mathrm{TeV} \leq \sqrt{s} << m_H$, the Born approximation cross section is constant, and, if extrapolated to asymptotic energies, would exceed the cross section for three W, Z production via conventional gauge interactions which must (with appropriate cuts on angular separation) scale as 1/s. Unfortunately, for the energy range accessible to the SSC the gauge background apparently[16] dominates three body production, presumably because of multiple polarization degrees of freedom in the final state.

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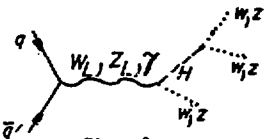

Figure 2

On the other hand, if the scalar system is indeed strongly interacting, events with four or more W's and Z's should significantly exceed the yield expected from gauge couplings alone. However, the multi-body event rates anticipated on the basis of the E/v rule are extremely small: for an integrated luminosity of 10^{40} cm^{-2} and a c.m. energy of 40 TeV, we found about 150 three body and 10-100 four body events using the prescriptions described above. In this case backgrounds present a major problem. While conventional gauge interactions[3] should not represent a prohibitive background for the total yield of multi W, Z events expected from the mechanism of Fig. 1, the anticipated two-jet QCD background is larger than the multi- W, Z signal by many orders of magnitude. Demanding one leptonic decay still leaves[15] an overwhelming background from W or Z plus high p_T jet. An important issue is thus whether the hadronic decays of W's and Z's can be distinguished from QCD jets. The wisdom which emerged from discussions at the $p\bar{p}$ workshop[15] is that a reduction factor of 1/7 in the background to signal ratio can be achieved by requiring a jet mass equal (within an appropriate definition) to the W,Z mass. This appears to be insufficient to extract two-body W and Z events for the yields estimated in Ref. 9. Demanding two leptonic decays (which excludes detection of two-W events) reduces the rates to a barely detectable level, even with a luminosity of 10^{33} cm^{-2}sec^{-1}. Therefore better methods for separating hadronic W and Z decays from QCD jets are highly desirable, not only for the scenario discussed here, but also if multi-W, Z events are to be used, for example, to test the standard model gauge couplings, or to search for a lighter ($2m_W < m_H <$ TeV) standard model Higgs, a technirho, etc. Sciulli[17] has considered the possibility of measuring the angular separation between individual particles in a jet. He concluded that the background reduction factor could be improved to about 1/25 in this way and suggested that a factor of 1/100 might be achievable. However this may be at the price of a severe reduction in solid angle: the total estimated[9] yield of multi W_L, Z_L events with invariant mass above 500 GeV is three to ten thousand for an integrated luminosity of 10^{40} cm^{-2}.

It might also be possible to extract a signal for strongly interacting vector bosons by less direct methods than identification of individual multi W_L, Z_L events, such as an anomalously high yield of W and Z leptonic decays and/or an anomalously high Z/W ratio in events with total transverse energy above, say, 500 GeV. If the muon angular distribution in $Z \rightarrow \mu\mu$ can be measured, it might further be possible to establish[15] an enhancement of longitudinally polarized Z's in this sample. These questions clearly require further study.

On the more theoretical side, a better understanding of the dynamics of a strongly interacting W_L, Z_L system, or plausible models of such a system, might give a better indication as to whether the low yield of events with multiplicity ≥ 3, estimated[9] by extrapolating the required threshold behavior (Adler zeros) is a fair guess or whether (hopefully?) it appreciably underestimates the multi-body event yield.

Conclusions

The questions raised above are intended to illustrate the way in which thinking about a specific scenario can raise further questions as part of an iterative process of providing not only input into the choice of SSC design parameters, but also directions for detector design, algorithms for data analysis, etc. Further study on the physics of strongly interacting W's and Z's did in fact go on at the workshop, as reported below along with the conclusions of other working groups. These included supersymmetry, compositeness, non standard Higgs particles, standard Higgs with mass $m_W < m_H < 2m_W$, mirror fermions and other exotics.

The physics to be revealed by exploring the TeV region at the SSC will undoubtedly bear little resemblance to anything discussed here, but hopefully exercises such as this and other workshops will leave us better prepared to exploit it.

Acknowledgments.

This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contracts DE-AC03-76SF00098 and the National Science Foundation under research grant PHY-84-06006.

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THE INTERMEDIATE MASS HIGGS

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Summary

The subgroup concentrated on two problems. The first was how to search for a standard neutral Higgs, H°, in the intermediate mass range, $2m_t < m_H < 2m_W$. The only method not obviously swamped by background is H° production in association with $W^\pm$, with $H^\circ \to t\bar{t}$ and $W \to \ell\nu$. However, this strategy puts great demands on the ability to identify t-quarks efficiently. The second problem addressed was the detectability of H° in the mass range 200 to 400 GeV. The question of signal versus background for a fusion-produced H° was reexamined. Heavy H° production in association with a Z° was also studied. The fusion process, with H° decay to W^+W^- or $Z^\circ Z^\circ$, remains by far the most promising way to look for a heavy Higgs scalar.

I. Introduction

Whether one believes or not in the validity of the minimal standard electroweak model with a single Higgs doublet,[2] the importance of confirming or excluding the existence of the physical scalar, H°, of that model can hardly be overstated. It is not surprising, therefore, that a great deal of attention has been given to the rates for signals and backgrounds for H° production. As we shall briefly review, most of that attention has been focussed on H° masses less than about $2m_t$ (~90 GeV?) or greater than $2m_W$. The intermediate mass Higgs (IMH), with $2m_t < m_H < 2m_W$, has received comparatively little study, presumably because the experimental chances looked slim, if not hopeless. The IMH subgroup's main task was to look carefully at this mass region and try to bring the situation there into sharper focus. That task was begun at Berkeley and vigorously continued by members of the Electroweak Subgroup at Snowmass. The IMH subgroup also extended their charge by looking again at the signals and backgrounds for H° with mass between 200 and 400 GeV.

Strategies for searching for both light ($m_H < 2m_t$) and heavy ($m_H > 2m_W$) Higgs have been discussed extensively in the literature. The light Higgs is primarily the property of e^+e^- colliders. (For a review, see Ref. 3.) The principal methods that have been proposed are e^+e^- annihilation to:

$$\Upsilon(b\bar{b}) \text{ or } \xi(t\bar{t}) \to H^\circ + \gamma \quad ,^4 \tag{1a}$$

$$Z_r \to Z_v + H^\circ;\ Z_v \to \ell^+\ell^- \quad ,^5 \tag{1b}$$

$$Z_v \to Z_r + H^\circ;\ Z_r \to \ell^+\ell^- \quad ,^5 \tag{1c}$$

The radiative decay from ξ should be useful for H° masses up to about $0.9m_\xi$. The decay of a real Z° into a virtual one plus H° is sensitive to $m_H < 45$ GeV, while the search at LEP II using virtual Z°'s can extend that limit to about 60 GeV (for $\sqrt{s}=175$ GeV and $L=2\times10^{31}\text{cm}^{-2}\text{sec}^{-1}$). The rates for all these electroweak processes are precisely known as a function of m_H because, in the minimal model, all of the H° couplings to gauge bosons and fermions are completely specified. If the recent report of the discovery of the top quark, with a mass of ~45 GeV,[6] holds up, the Wilczek process[4] may prove the most definitive way of searching for the light standard Higgs.

Once its mass exceeds $2m_Z$, the H° decays almost exclusively to W^+W^- and to $Z^\circ Z^\circ$ with the rates

$$\Gamma(H^\circ \to W^+W^-) = \frac{G_F m_H^3}{32\pi\sqrt{2}}(4-16m_W^2/m_H^2+12m_W^4/m_H^4)(1-4m_W^2/m_H^2)^{1/2} \tag{2a}$$

$$\Gamma(H^\circ \to Z^\circ Z^\circ) = \frac{G_F m_H^3}{64\pi\sqrt{2}}(4-16m_Z^2/m_H^2+12m_Z^4/m_H^4)(1-4m_Z^2/m_H^2)^{1/2} \tag{2b}$$

Only a multi-TeV hadron collider can access such a heavy Higgs. There, the primary production mechanisms are two-gluon fusion[7] and WW and ZZ fusion.[8] The signals and two-gauge-boson backgrounds have been discussed in detail by EHLQ,[9] who concluded that a 40 TeV collider with $\int L dt = 10^{40}\text{cm}^{-2}$ can reach up to $m_H=1$ TeV. Beyond that mass, the H° is expected to be so wide that it would be unrecognizable as a resonance in WW and ZZ production. At the same time, partial-wave unitarity breaks down in electroweak perturbation theory,[10] so that signals of the strongly-interacting Higgs sector should be seen.[11,12]

The two topics studied most intensively by the IMH group were:

(1) Methods of searching for the intermediate mass H°, with estimates of signal and background cross sections. The obvious method is two-gluon fusion of H°, followed by H° decay to $t\bar{t}$. The results in EHLQ, Chapter 4, indicate that the ordinary $t\bar{t}$ background swamps the H° signal. Thus, most attention was placed on the less background-dominated processes $p^\pm p \to H^\circ + W^\pm$ and $p^\pm p \to H^\circ + Z^\circ$,[13] with $H^\circ \to t\bar{t}$. The signal has a cross section of only O(1 pb), but the background (for the first process) is a factor of 2 or more smaller. In all of this, we took as a ground rule that t-quarks could be positively identified with finite efficiency. The discussion and (preliminary) results for H° production in association with weak gauge bosons are presented in Section II.

(2) More in-depth study of signal vs. background for H° with mass in the range 200-400 GeV. It was felt by some[14] that EHLQ had underestimated the real WW and ZZ background to $p^\pm p \to H^\circ \to WW$ or ZZ, so that this process was not suitable for finding a heavy Higgs in this lower mass range. Consequently, an attempt was made to estimate the background more conservatively, still assuming that real W- and Z-pairs constituted the only true background. Motivated by the possibility that the EHLQ estimates had been too liberal, members of the group also looked at the $H^\circ + W^\pm$ and $H^\circ + Z^\circ$ processes. The revised background estimates and results on associated H° production are given in Section III.

II. Searching for the Intermediate Mass Higgs[15]

In a hadron collider the dominant production mechanism for the standard neutral Higgs with $m_H < 200$ GeV is two-gluon fusion.[7,9] For $m_H < 2m_W$, the H° decays to the heaviest fermion-antifermion pair that is kinematically accessible, and it has long been believed that backgrounds dwarf the signal for any H° mass in this range.[9] The situation is epitomized in Chapter 4 and Figs. 4-51 (Fig. 157 in Rev. Mod. Phys.) and 4-52 (158) of EHLQ, for $m_t=30$ GeV. By definition, the IMH decays to $t\bar{t}$. The production cross section in a 40 TeV collider, with both t and $\bar{t}$ emitted in the rapidity interval (-1.5, 1.5), falls from 0.1 nb at $m_H=100$ GeV to 0.03 nb at $m_H=2m_W$. Even under the hopeful assumptions that t-quarks can be identified, that the only background is true $t\bar{t}$-jet production, and that the $t\bar{t}$-invariant mass resolution is 10 GeV, the background is at least 500 times the signal. For a t-quark mass of 45 GeV, the signal is still more than 200 times smaller than this most optimistic background.[16]

Therefore, it was decided to consider associated production of H° with a weak gauge boson, either $W^\pm$ or Z°. This is the analog in a hadron collider of the Bjorken process.[5,13] Although the cross sections are $O(\alpha^2)$, associated production of an intermediate mass H° has a much more favorable signal-to-background ratio than gluon fusion so long as one tags the outgoing W or Z and identifies top quarks efficiently. The differential cross sections for the $q\bar q$ annihilation subprocesses, averaged over colors, are

$$\frac{d\sigma}{d\Omega}(q_i\bar q_j \to H^\circ W^\pm) = \frac{\alpha^2|K_{ij}|^2}{48\sin^4\theta_W}\left(\frac{2k}{\sqrt{s}}\right)\left(\frac{1}{s-m_W^2}\right)^2\left(m_W^2+\frac{k^2}{2}\sin^2\theta\right) \quad (3a)$$

and

$$\frac{d\sigma}{d\Omega}(q_i\bar q_i \to H^\circ Z^\circ) = \frac{\alpha^2[(T_{3i}-Q_i\sin^2\theta_W)^2+(Q_i\sin^2\theta_W)^2]}{24\sin^4\theta_W\cos^4\theta_W}$$

$$\times \left(\frac{2k}{\sqrt{s}}\right)\left(\frac{1}{s-m_Z^2}\right)^2\left(m_Z^2+\frac{k^2}{2}\sin^2\theta\right) \quad (3b)$$

where k is the c.m. momentum of the final particles, θ is the c.m. scattering angle and K_{ij} is the appropriate Kobayashi-Maskawa matrix element ($K_{ij}=\delta_{ij}$ in calculations). The corresponding total cross sections are

$$\sigma(q_i\bar q_j \to H^\circ W^\pm) = \frac{\pi\alpha^2|K_{ij}|^2}{36\sin^4\theta_W}\left(\frac{2k}{\sqrt{s}}\right)\left(\frac{1}{s-m_W^2}\right)^2(k^2+3m_W^2) \quad (4a)$$

$$\sigma(q_i\bar q_i \to H^\circ Z^\circ) = \frac{\pi\alpha^2[(T_{3i}-Q_i\sin^2\theta_W)^2+(Q_i\sin^2\theta_W)^2]}{18\sin^4\theta_W\cos^4\theta_W}$$

$$\times \left(\frac{2k}{\sqrt{s}}\right)\left(\frac{1}{s-m_Z^2}\right)^2(k^2+3m_Z^2) \quad . \quad (4b)$$

The total, uncut, cross sections in a pp collider with $\sqrt{s}$=10, 20 and 40 TeV are shown in Fig. 1 for $H^\circ+W^\pm$ and in Fig. 2 for H°+Z°. [These are adapted from Figs. 4-53 (159) and 4-55 (161) of EHLQ.] The cross sections in a $p\bar p$ collider are indistinguishable. We see that, for the IMH mass range and $\sqrt{s}$=40 TeV, $\sigma(pp\to H^\circ+W^\pm) \approx 2\sigma(pp\to H^\circ+Z^\circ) \approx 10-1$ pb.

The signal for associated production is two top jets from H° plus the weak boson. To eliminate large multi-jet backgrounds, it was decided, and it probably is necessary, to tag the W or Z in its leptonic (e or μ) decay modes. It was assumed further that the t-quarks could be distinguished from other flavors and from gluon jets. This optimistic assumption received much study at Snowmass.[17,18] Then the background processes are $pp\to g+W^\pm$; $g\to t\bar t$ (Fig.3a), and $pp\to g+Z^\circ$; $g\to t\bar t$ and $pp\to gg\to t\bar t+Z^\circ$ (Fig. 3b). It takes no calculation to see that the signal-to-background ratio should be considerably better for $H^\circ+W^\pm$, and so all attention was focused on that reaction. Since there is one missing neutrino from the W, the t and $\bar t$ will have to be identified in their nonleptonic decay modes. This puts a special premium on the ability to distinguish t's from b's, not to mention lighter flavors and gluons. If this cannot be done efficiently, it is clear that the backgrounds are hopelessly large once again.

The signal and background cross sections for $pp\to H^\circ+W^+$ (only), $H^\circ\to t\bar t$ in a 40 TeV collider are shown in Fig. 4.[15] Here the top quark mass was taken to be 40 GeV. The cross sections with a W^- emitted are almost identical. Also, they are not significantly different if m_t=30 GeV or 50 GeV is used instead. In these calculations, the following cuts were imposed to enhance the signal and take account of the finite coverage of real detectors:

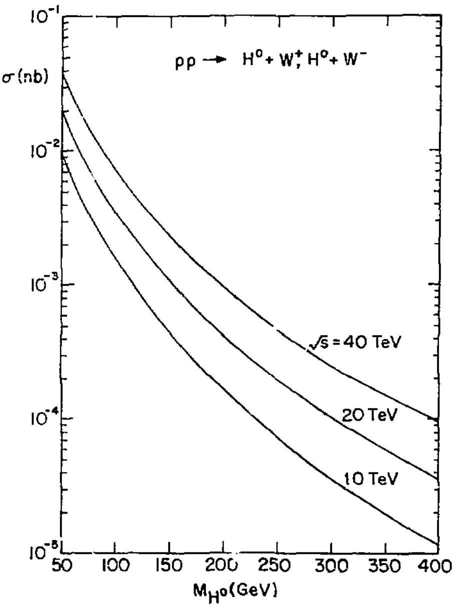

Fig. 1. Total cross sections in pp collisions for H° production in association with $W^\pm$. (From Ref. 9).

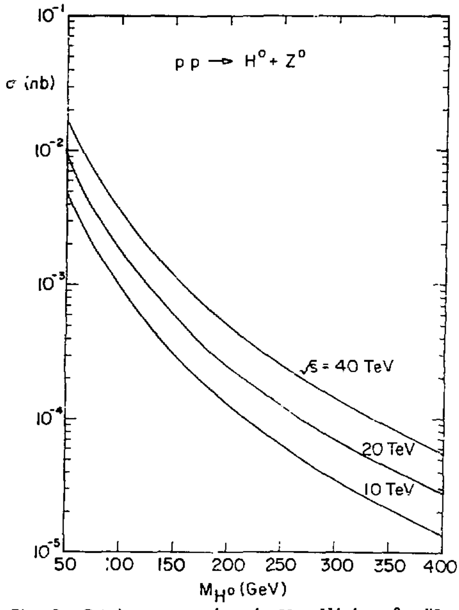

Fig. 2. Total cross sections in pp collisions for H° production in association with Z°. (From Ref. 9).

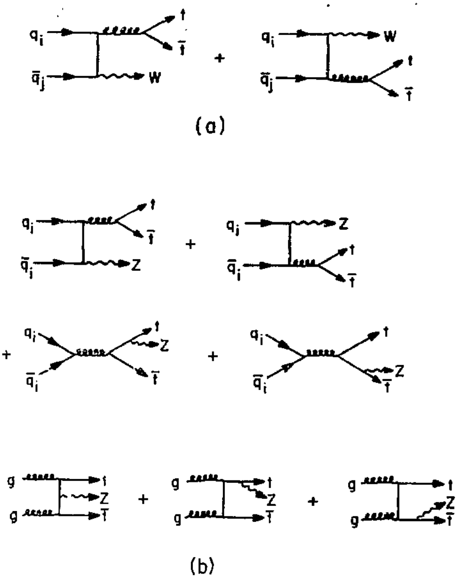

Fig. 3. Elementary subprocesses for background to (a) $p^{\pm}p\to H^\circ+W^{\pm}$, $H^\circ\to t\bar t$ and (b) $p^{\pm}p\to H^\circ+Z^\circ$, $H^\circ\to t\bar t$. Curly lines denote gluons.

$$p_T(W) > 40 \text{ GeV} \quad ; \qquad -2 < y_W, y_H < 2 \quad . \tag{5}$$

In Fig. 4 there is no additional cut on the t and $\bar t$ directions. The background was calculated assuming a $t\bar t$ invariant-mass-squared resolution of 0.1 m_H^2 . Also shown in Fig. 4 is the rate for the potential background, $pp\to \bar t bW^+$, divided by 100. This reaction proceeds mainly by two-gluon fusion of the final state. The cuts of Eq.(6) are put on the W and $\bar t$b-system. It is clear from this plot that the b/t discrimination must be better than ~10% if the signal is not to be swamped.

To get some idea of the demands on machine and detector that the curves in Fig. 4 imply, let us estimate the minimum effective integrated luminosity required to discover the IMH in associated production. For this, we assume that the b/t discrimination is good enough to ignore the $\bar t bW$ background. We require that there be at least 100 signal events in a standard 10^7 second run and that these represent at least ≈ 5 standard deviation excess over the $t\bar tW$ background. We assume that $W\to e\nu$ and $W\to\mu\nu$ are detected with 100% efficiency when $p_T(W) > 40$ GeV. Then, since we require that both t's decay nonleptonically, our "discovery criterion" amounts to

$$N_{sig} = 2B(W\to e\nu)(\epsilon_t B(t\to bq\bar q))^2\sigma(H^\circ W)\int Ldt > 100 \tag{6a}$$

$$N_{sig}/\sqrt{N_{bkd}} > 5 \tag{6b}$$

where

$$N_{bkd} = 2B(W\to e\nu)(\epsilon_t B(t\to bq\bar q))^2\sigma(t\bar tW)\int Ldt \tag{6c}$$

and ϵ_t is the t-quark detection efficiency. For the cross sections shown in Fig. 4, doubled to include W^-

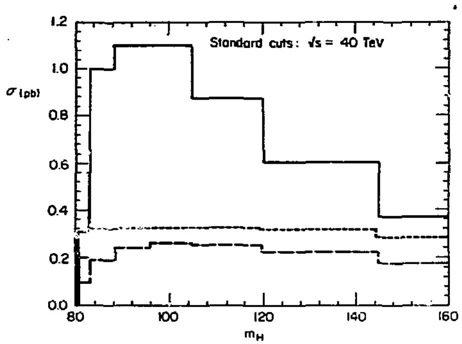

Fig. 4. Cross sections at √s=40 TeV for $pp\to H^\circ+W^+$ (solid line), $pp\to g+W^+$, $g\to t\bar t$ (long dashed line) and $pp\to \bar t b+W^+$, divided by 100 (short dashed line), from Ref. 15. m_t=40 GeV; cuts are given in Eq.(5).

as well as W^+, the requirement of 100 (or even 50 events) is the most stringent. We find that $\epsilon_t^2 \int Ldt$ must be at least 6×10^{38} if m_H=100 GeV and at least 2×10^{39} if m_H=160 GeV. Assuming integrated luminosity of 10^{40} cm^{-2}, these estimates imply ϵ_t=25-50%. These are extraordinarily, perhaps unrealistically, high detection efficiencies. However, it is so important to have the capability to confirm or exclude the existence of a Higgs scalar in the intermediate mass range that much more work on the problem of efficient t-identification at high luminosity is well-justified.

III. Higgs Detection for 200 GeV < m_H < 400 GeV[9,14,19]

The reach of a supercollider for a standard neutral Higgs with $m_H > 2m_W$ was determined by EHLQ as follows: They considered the processes $p^{\pm}p\to H^\circ\to W^+W^-$, $Z^\circ Z^\circ$ via both two-gluon and two-weak boson fusion. The background to these processes was assumed to be due only to true WW and ZZ production, respectively. To estimate the background, they assumed further that the weak boson invariant mass resolution was the greater of 10 GeV and the total H° width, Γ_H. The signal and background rates thus determined for a pp collider with √s=40 TeV are shown in Fig. 5 for WW and Fig. 6 for ZZ. Note that the rapidity of each weak boson is assumed to lie in the range (-2.5, 2.5). Finally, the collider reach as a function of machine energy and luminosity was determined by requiring that there be at least 5000 H°→WW or ZZ events and that the signal must stand above the background by at least 5 standard deviations. (The number of events required is large enough to guarantee that the H° could be discovered, if need be, in $H^\circ\to Z^\circ Z^\circ$ with both Z°'s decaying to e^+e^- or $\mu^+\mu^-$.) This discovery criterion implies that a 40 TeV collider with $\int Ldt=10^{40}$ can access the m_H range from $2m_W$ to 1 TeV.

The EGK collaboration[14] argued, perhaps justifiably, that EHLQ's estimate of the WW and ZZ background is too hopeful for low H° masses (see also Ref. 20). They proposed instead that the pair invariant mass resolution be taken to be the larger of $m_H/10$ and Γ_H. The two methods agree for m_H > 500 GeV. Below this mass, EGK found that the backgrounds exceeded the signals by a factor of 3-5.[21] Thus, they concluded that the fusion-produced H° would be very difficult to discover for m_H = 200 - 400 GeV.

In their calculations, EGK placed no restriction

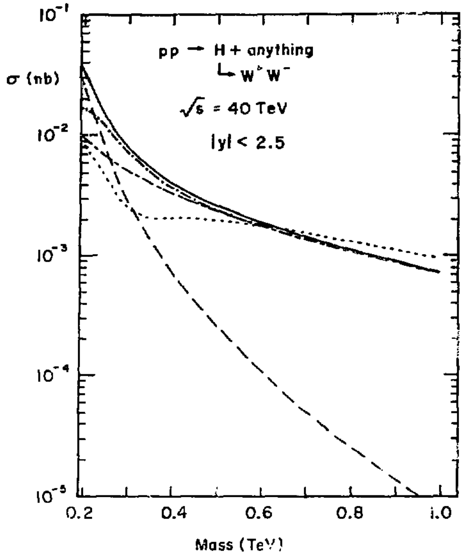

Fig. 5. Cross sections and backgrounds for pp→H°+anything, H°→W⁺W⁻, at √s=40 TeV. The dashed line is gg fusion of H°, the long-short dashed line is WW and ZZ fusion of H°, and the solid line is the total cross section. The dotted line is the WW background as computed by EHLQ (Ref. 9) and the dash-dotted line is the revised background suggested by EGK (Ref. 14).

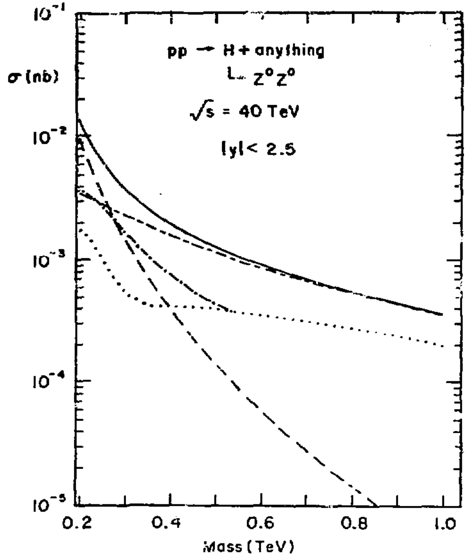

Fig. 6. Same as Fig. 5 for pp→H°+anything, H°→Z°Z°.

on the weak boson rapidities. This favors the background relative to the signal because the ordinary electroweak production of gauge boson pairs is peaked forward in their c.m. frame, while their production from H° decay is isotropic in that frame. The rapidity cuts imposed by EHLQ thus enhance significantly the signal-to-background ratio. The net result of EGK's poorer mass resolution and the cuts is shown for √s=40 TeV in Figs. 5 and 6. For H°→WW, the signal/background is about 1/1 over the entire mass range; for H°→ZZ, it is about 2/1. At √s=10 TeV, the revised signal/background is ~1/2-1/3 for H°→WW and ~1/1 for H°→ZZ over the accessible H° mass range. The corresponding ratios for a 20 TeV collider are ~1/1 and ~2/1. Given the discovery requirement of 5000 events, even the revised backgrounds are never a limiting factor. Thus EHLQ's conclusions on the reach of a supercollider as a function of luminosity and c.m. energy remain unchanged (see Fig. 4.50 in Ref. 9). In particular, assuming realistic W and Z detection efficiencies,[22] experiments at a 40 TeV collider with $\int L dt = 10^{40} cm^{-2}$ can discover a heavy Higgs with mass in the range $2m_W$-1000 GeV. At the same luminosity, a 10 TeV collider can reach only to m_H=350 GeV, a 20 TeV collider to 650 GeV.

As a hedge against the possibility that the signal for pp→H°+anything, $H^\circ \to W^+W^-$ or $Z^\circ Z^\circ$, would be "drowned" by the background for Higgs masses in the range 200-400 GeV, Ellis and Sharpe[10,23] investigated the associated production pp→H°+Z° with H°→WW or ZZ. The trigger topology is $Z^\circ \to e^+e^-$ or $\mu^+\mu^-$ plus four jets from the two additional weak bosons. One then has to plot the invariant mass of one pair of jets against that of the other pair, select events with $m_{12} \approx m_{34} \approx m_W$, m_Z, and then look for an enhancement in the four-jet mass.

As of the Berkeley Workshop, Ellis and Sharpe had not done detailed calculations of the signal and background with cuts. However, we can estimate the signal rates as follows: From Fig. 2,[9] the total H°+Z° cross section in a 40 TeV colider is seen to fall from 0.5 pb at m_H=200 GeV to 0.05 pb at m_H=400 GeV. Including the $Z^\circ \to e^+e^-$ branching ratio (~6%) and allowing no more than one W or Z decay involving a hard neutrino, we expect no more than 2 x 10^{-2} to 2 x 10^{-3} pb of signal in this mass range. This amounts to at most 200 (20) events for a 200 (400) GeV H° produced at 40 TeV with $\int L dt = 10^{40}$. The background is unknown.

By comparison, the cross section at √s=40 TeV for pp→H°→Z°Z°, with $|y_Z| < 2.5$, is seen from Fig. 6 [9] to fall from 10 pb at m_H=200 GeV to 2 pb at m_H=400. The background from continuum ZZ production is small. If it is necessary to tag one of the Z's in its e^+e^- or $\mu^+\mu^-$ decay modes, this amounts to between 12000 and 2400 events per year at $\int L dt = 10^{40}$. Inclusive H° production via the fusion mechanisms thus remains the key to discovering a heavy Higgs scalar.

Conclusions

Definitive tests of many of the scenarios for new physics in the 1-TeV energy regime will require efficient identification of heavy quarks and leptons and of the W and Z bosons. In particular, the search for an intermediate mass Higgs scalar apparently requires very high t-quark detection efficiency at high collider luminosity, of order 25-50% at $\int L dt = 10^{40}$ cm^{-2}. The situation is not hopeful. But the importance of being able to cover this Higgs mass range warrants continued study of this problem.

For a neutral Higgs heavier than $2m_W$ the two-gluon and two-weak boson fusion mechanisms give large production rates and the true WW and ZZ backgrounds are quite manageable. Searching for the heavy H° should be straightforward with only modest W and Z detection efficiencies. The rates for production of a heavy H° in association with a W or Z are small. If a heavy H° exists, the observation of this process will make for an interesting second-generation experiment at the SSC. Meanwhile, much more study of the signal and, especially, the background rates is needed.

Acknowledgements

Thanks are due to the organizers of the Berkeley Electroweak Workshop, Tom Appelquist, Ian Hinchliffe and Mary K. Gaillard, for providing the stimulating atmosphere in which the Subgroup's work was done. I also thank my other two collaborators, Estia Eichten and Chris Quigg for important contributions to the contents of this report. This work was supported in part by the U.S. Department of Energy under Contract No. EY-76-C-02-1545.

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NONSTANDARD HIGGS BOSONS

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Summary

We have considered the prospects for producing and detecting very massive (eg. 100 GeV - 1 TeV) nonstandard Higgs bosons at a high energy pp collider. By nonstandard Higgs bosons we mean the additional color-singlet spin-0 particles that occur in most extensions of the minimal (i.e. one Higgs doublet) $SU_2 \times U_1$ model. Our main conclusions are: (a) The physics of nonstandard neutral Higgs particles is likely to be very similar to that of *the neutral scalar in the minimal model.*
(b) Massive charged Higgs bosons will probably (ie. for most models and most parameter ranges) be very difficult to produce or detect, especially if their dominant decay is into $t\,\bar{b}$.
(c) The W^+ZH^-, $W^+W^+H^{--}$, and $ZH^0H^{0\prime}$ vertices are especially interesting.
(d) The identification of t and b quark jets, mass resolution for jet pairs, and lepton sign identification are extremely important.
(e) More detailed studies of various production mechanisms and backgrounds would be very useful.

Motivations[1,2]

In the minimal $SU_2 \times U_1$ model, spontaneous symmetry breaking is accomplished by a single Higgs doublet (implying one physical scalar particle). However, many modifications and extensions require additional elementary Higgs multiplets. These include most models incorporating supersymmetry, the Peccei-Quinn solution to the strong CP problem, spontaneous CP breaking, spontaneous lepton number non-conservation, family symmetries, etc. In addition, extra physical spin-0 particles emerge in most models without elementary Higgs fields, such as composite Higgs[3] models or technicolor.[4]

Existing Constraints

There are no significant direct contraints on neutral Higgs particles[1] unless they are lighter than $\simeq 1$ GeV. In particular, the standard Higgs cannot[1] be lighter than $m_K - m_\pi$. Experiments at PETRA require $m_{H^\pm} > 16$ GeV *for* singly charged Higgs bosons unless their couplings to $\tau\nu_\tau$ are anomalously small. Also, if the UA1 group has indeed observed the t quark (via its expected semileptonic weak decay) one has $m_{H^\pm} > m_t - m_b$ for most models.

An important indirect constraint is provided by the neutral current parameter $\rho \equiv m_W^2/m_Z^2 \cos^2\theta_W$, which is experimentally very close to unity. This strongly suggests that only Higgs doublets or singlets have significant vacuum expectation values (VEVs). Other possibilities include (a) a custodial SU_{2C} symmetry to relate the VEVs of other Higgs representations[3,5]. (b) fine tuning, or (c) the use of some specific Higgs representations with $I \geq 3$ which also give $\rho = 1$.

Another indirect constraint comes from flavor changing neutral current *(FCNC) limits*, the $K_L - K_S$ mass difference, etc. In particular, neutral Higgs bosons in general have dangerous flavor off-diagonal couplings. This is avoided in models with natural flavor conservation[6], in which the neutral member of only one Higgs doublet couples to the charge 2/3 quarks, and only one couples to the charge -1/3 quarks (in such models the Yukawa matrix for each charge quark is proportional to the mass matrix and is therefore flavor diagonal). In the minimal model the same doublet couples to both types of quarks and the Yukawa coupling to quark q_i is $h_i = gm_i/2M_W$, where $g = e/\sin\theta_W$ is the SU_2 gauge coupling. In a particularly simple two doublet model, the neutral member of the first (second) doublet couples only to the charge 2/3 (-1/3) quarks, so that the corresponding Yukawa couplings of the weak eigenstate Higgs fields are

$$h_i^1 = \lambda_1\, g m_{u_i}/2M_W$$

$$h_i^2 = \lambda_2\, g m_{d_i}/2M_W \quad , \tag{1}$$

where $\lambda_{1,2} > 1$. Both versions of the natural flavor conservation model can be generalized (and the Yukawa couplings enhanced) by adding additional Higgs multiplets with non-zero VEVs but no Yukawa couplings to quarks.

Another possible solution to the FCNC constraints is to suppress the unwanted amplitudes, either by fine-tuning parameters or by assuming that the relevant Higgs bosons are very massive (> several TeV).

Nonstandard Higgs bosons do not contribute significantly to g-2, nonleptonic parity violation, etc., for most reasonable mass and parameter ranges.[1,2]

There are a number of theoretical constraints on Higgs masses. In the minimal model vacuum stability requires[7,8] $m_{H^0} \gtrsim m_{CW}/\sqrt{2} \simeq 7$ GeV, where m_{CW} is the Coleman-Weinberg[9] mass, while cosmological arguments[7,1] *suggest* $m_{H^0} > m_{CW} \sim 10.3$ *GeV. Note that these bounds* are weakened by a factor $\left[1-(m_t/81.2\text{ GeV})^4\right]^{1/2}$ for a heavy t quark or by a fourth fermion family. These limits do <u>not</u> hold in nonstandard models (except[1] for bounds on a particular linear combination of $m^4_{H_i}$).

There are no rigorous upper bounds on Higgs masses. In the minimal model $m_{H^0} = \sqrt{2\lambda}\,\nu$, where $\nu = (\sqrt{2}\,G_F)^{-1/2}$ $\simeq 246$ GeV is the weak scale and λ is the Higgs quartic self-coupling. If $\lambda > 8\pi/3$ ($m_{H^0} \simeq 1$ TeV) one has a strongly coupled theory[10,11] and, in particular, tree unitarity[10] fails. However, this is not necessarily disastrous: the effects of the strongly coupled Higgs sector are screened at low energies[12] and there may be interesting consequences at high energy,[13,10] such as Higgs-Higgs bound states or multiple W and Z production. For nonstandard Higgs it can be shown[14] very generally that the lightest physical neutral Higgs particle must satisfy $m_{H^0} \lesssim \sqrt{2\lambda}\, v$, where λ is a particular quartic coupling.

Two Models

In the two-doublet model[2] with natural flavor conservation one has

$$\tilde{\psi}_1 \equiv \begin{pmatrix} \tilde{\psi}_1^{\,o} \\ \tilde{\psi}_1^{\,-} \end{pmatrix} \qquad \tilde{\psi}_2 \equiv \begin{pmatrix} \tilde{\psi}_2^{\,o} \\ \tilde{\psi}_2^{\,-} \end{pmatrix} \tag{2}$$

where $\tilde{\psi}_1^{\,o}$ ($\tilde{\psi}_2^{\,o}$) has Yukawa couplings to the charge $\frac{2}{3}$ ($-\frac{1}{3}$) quarks only. (The leptonic couplings must also be specified). This can come about because of some additional symmetry, such as supersymmetry or an additional global U_1 or discrete symmetry. The VEVs are $\langle \tilde{\psi}_i^{\,o} \rangle = \nu_i/\sqrt{2}$, with $(\nu_1^2 + \nu_2^2)^{\frac{1}{2}} = \nu = 246$ GeV $= 2M_W/g$. The physical (mass eigenstate) Higgs particles are

$$\psi^{\pm} = (\nu_2\, \tilde{\psi}_1^{\,\pm} - \nu_1\, \tilde{\psi}_2^{\,\pm})/\nu$$

$$\psi_1^{\,o} = \cos\theta \left[\frac{\tilde{\psi}_1^{\,o} + \tilde{\psi}_1^{\,o\dagger}}{\sqrt{2}} - \nu_1 \right] + \sin\theta \left[\frac{\tilde{\psi}_2^{\,o} + \tilde{\psi}_2^{\,o\dagger}}{\sqrt{2}} - \nu_2 \right]$$

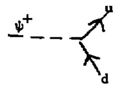
$$\psi_2^{\,o} = -\sin\theta \left[\frac{\tilde{\psi}_1^{\,o} + \tilde{\psi}_1^{\,o\dagger}}{\sqrt{2}} - \nu_1 \right] + \cos\theta \left[\frac{\tilde{\psi}_2^{\,o} + \tilde{\psi}_2^{\,o\dagger}}{\sqrt{2}} - \nu_2 \right]$$

$$\psi_3^{\,o} = i \left[\frac{\nu_2 (\tilde{\psi}_1^{\,o} - \tilde{\psi}_1^{\,o\dagger}) - \nu_1\, (\tilde{\psi}_2^{\,o} - \tilde{\psi}_2^{\,o\dagger})}{\sqrt{2}\, \nu} \right] \tag{3}$$

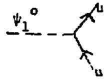
where θ is a mixing angle determined (like ν_1 and ν_2) by the parameters in the Higgs potential. $\tan\theta$ is typically of order[2] ν_1/ν_2 (times a ratio of quartic couplings). For simplicity, CP invariance has been assumed in (3), in which case $\psi_1^{\,o}$ and $\psi_2^{\,o}$ couple to fermions as scalars and $\psi_3^{\,o}$ as a pseudoscalar. (If this assumption is relaxed then $\psi_1^{\,o}$, $\psi_2^{\,o}$, and $\psi_3^{\,o}$ can all mix and the Yukawa couplings become mixtures of S and P). The Yukawa couplings are given in Table 1.

For $\nu_1 \ll \nu_2 \simeq \nu$ (or $\nu_2 \ll \nu_1 \simeq \nu$) the Yukawa couplings proportional to M^u (M^d) are enhanced. The existing constraints are not very stringent, however. For $\nu_1 < \nu_2$ one has (Abbott et al.[2]) $|\nu_2/\nu_1| < (2m_\psi/m_c)^{\frac{1}{2}}$ from the K_L-K_S mass difference, which is of O(10) for $m_{\psi^+} \simeq M_W$ (this may be strengthened somewhat for a large t quark mass[15]). Sher and Silverman[2] have recently argued that $|\nu_2/\nu_1| < m_{\psi^+}/\sqrt{m_t m_c}$ if the t quark decays normally, and have discussed (future) constraints from toponium. The limits on $\nu_2 < \nu_1$ are even weaker: $|\nu_1/\nu_2| < 10\ (m_{\psi^+}/m_s)^{\frac{1}{2}}$ from $D^o \nleftrightarrow \bar{D}^o$ (Abbott et al.[2]) and $< 4.0 (m_{\psi^+}/m_b)$ possible from b decay (Sher and Silverman[2]).

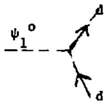
Another interesting model (Machacek, ref. 5), motivated by composite Higgs models[3], leads to doubly charged Higgs particles as well as new couplings. It has an ordinary Higgs doublet $X = (X^+ X^o)^T$ and complex



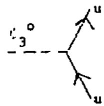
$$\frac{g}{2\sqrt{2}\, M_W} \left[M^u K \frac{\nu_2}{\nu_1}(1-\gamma_5) + K M^d \frac{\nu_1}{\nu_2}\,(1+\gamma_5) \right]$$



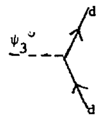
$$\frac{g}{2M_W} \qquad M^u\, \frac{\nu}{\nu_1} \cos\theta$$



$$\frac{g}{2M_W} \qquad M^d\, \frac{\nu}{\nu_2} \sin\theta$$



$$\frac{g}{2M_W} \qquad M^u\, \frac{\nu_2}{\nu_1}\, \gamma_5$$



$$\frac{g}{2M_W} \qquad M^d\, \frac{\nu_1}{\nu_2}\, \gamma_5$$

Table 1. Yukawa couplings to quarks in the two doublet model with natural flavor conservation. M^u and M^d are the diagonal mass matrices for the charge $\frac{2}{3}$ and charge $-\frac{1}{3}$ quarks, respectively, and K is the Kobayashi-Maskawa-Cabibbo matrix. The couplings for $\psi_2^{\,o}$ are obtained from those for $\psi_1^{\,o}$ by replacing $\cos\theta \to -\sin\theta$, $\sin\theta \to \cos\theta$.

(ψ) and real (ξ) Higgs triplets

$$\phi = \begin{pmatrix} \psi^o & \xi^+ & \psi^{++} \\ \psi^- & \xi^o & \psi^+ \\ \psi^{--} & \xi^- & \psi^{o*} \end{pmatrix}, \tag{4}$$

where $\psi^{--} = (\psi^{++})^\dagger$, etc. The triplets have no Yukawa couplings. One postulates that for some reason there is a custodial SU_{2C} symmetry manifested by $\langle\phi\rangle = $ diag (b b b), so that $\rho = 1$. Then $\nu = 246$ GeV $= \sqrt{8b^2 + a^2}$, where $\langle X^o \rangle = a/\sqrt{2}$. The physical Higgs states then consist of an SU_{2C} 5-plet (H_5^{++}, H_5^{+}, H_5^{o}, H_5^{-}, H_5^{--}), an SU_{2C} triplet (H_3^+, H_3^o, H_3^-), and two SU_{2C} singlets ($H_{1a}^{\,o}$, $H_{1b}^{\,o}$).

Couplings

The various models may be distinguished by the couplings that are allowed at tree level (or at one loop level for the gluon vertices). The allowed couplings for the minimal (one doublet) model, for two nonstandard models just described and for technipions in technicolor models (Lane, ref. 1) are indicated in Table 2. A noteworthy feature is the $ZH_S^0H_P^0$ vertex, which first appears in the two doublet model[16]. Also, the ZZH^0 and $W^+W^-H^o$ vertices are present in all

Vertex	Minimal model (One doublet)	Two doublet	Doublets and Triplets	Technipions
$\gamma, Z \to H^+H^-$	x	yes	yes	yes
$Z \to H_S^o H_P^o$	x	yes	yes	x
$W^+ \to H^+H^o$	x	yes	yes	yes
$ZZ \to H^o$	yes	yes	yes	suppressed
$W^+W^- \to H^o$	yes	yes	yes	x
$W^+Z \to H^+$	x	x	yes	x
$W^+W^+ \to H^{++}$	x	x	yes	x
$W^+W^-, ZZ \to H^oH^o$	yes	yes	yes	probably
$\gamma\gamma, W^+W^-, ZZ \to H^+H^-$	x	yes	yes	probably
$\gamma W, ZW \to H^{\pm}H^o$	x	yes	yes	probably
$gg \to H^o$	yes	yes	yes	yes
$gg \to H^oH^o$	yes	yes	yes	–
$gg \to H^+H^-$	x	yes	yes	–

Table 2. Vertices that are allowed at tree level (or one loop level for gluons). The subscripts S and P refer to scalar and pseudoscalar, respectively.

elementary Higgs models, but the W^+ZH^- vertex is absent in all models involving Higgs doublets only (this is most easily seen by going to the basis in which only one doublet has a nonzero VEV).

Production, Detection, Backgrounds

Standard Model

If the standard Higgs mass m_{H^o} is $\lesssim 100$ GeV, it can be produced by toponium $\to H^o\gamma$ (ref. 17), by $Z \to Z^*H^o$, or by $e^+e^- \to Z^* \to ZH^o$ (Z^* is a virtual Z), with $H^o \to t\bar{t}$ or $b\bar{b}$.

Higgs physics will be very difficult[18] if 100 GeV $< m_{H^o} < 2M_W$. The H^o may be produced by gg fusion (ref. 19), by WW or ZZ fusion (ref. 20), or by $gg \to H^o t\bar{t}$ (ref. 21), with a large cross section,[19-22] but the dominant $t\bar{t}$ or $b\bar{b}$ decays should be swamped by QCD backgrounds[18,21-22]. (The decay[23] $H^o \to WW^* \to Wf\bar{f}$ may be a possibility, however). The best possibility is probably $q\bar{q} \to W^* \to WH^o$ or $q\bar{q} \to Z^* \to ZH^o$; although the cross section is low (see Figure 1), the signal may significantly exceed the background with appropriate cuts.[18]

For $m_{H^o} > 2M_W$ one has the cleaner signature $H^o \to W^+W^-$ or ZZ (assuming that one can recognize W's and Z's). Cross sections for $gg \to H^o$, WW or $ZZ \to H^o$, $gg \to H^o t\bar{t}$, and $q\bar{q} \to WH$ are shown in Figure 1. Assuming that $\sigma \sim 10^{-4}$ nb is required for observability (10^3 events for $L = 10^{40}$ cm^{-2}, which allows the secondary W's to be observed by their leptonic decays), it should be possible[18] to observe the WH^o or $H^o t\bar{t}$ processes for $M_{H^o} \lesssim 400$ GeV (the $WW \to H^o$ process is defeated by $q\bar{q} \to W^+W^-$ background[18]). For 400 GeV to 1 TeV it may be possible to observe $WW \to H^o \to W^+W^-$, although background from $q\bar{q} \to W^+W^-$ is still a severe problem.

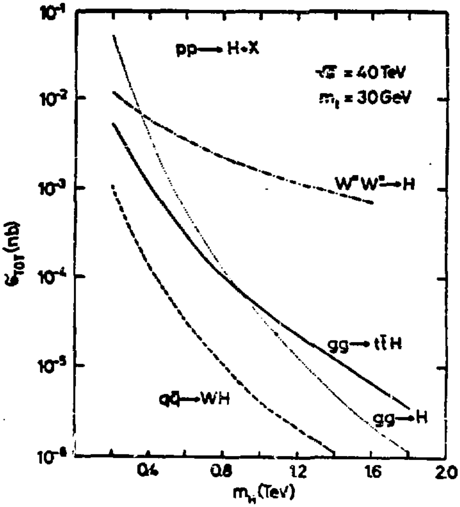


Figure 1. Cross sections for production of a conventional Higgs boson at $\sqrt{s} = 40$ TeV (from ref. 21).

Nonstandard Neutral Higgs

The production mechanisms and signatures for neutral Higgs are similar to those in the minimal model, with two exceptions.

(a) The $Z^oH_S^oH_P^o$ vertex[16] allows neutral Higgs production via the Drell-Yan process $pp \to Z^* \to H_S^oH_P^o$. This has been calculated for the case of $m_{H_S} = m_{H_P}$ by Lane[24], who finds a significant cross section for $m_H < 200$ GeV (Fig. 2).

(b) The $\bar{t}tH^o$ vertices may be enhanced by $|v_2/v_1| < (2\, m_{H^+}/m_c)^{1/2}$ leading to enhancements of the cross sections for $gg \to H^o t\bar{t}$ (Fig. 1 and ref. 21), $t\bar{t} \to H^o$ (Fig. 2), and $gg \to H^o$ (via a t or a new very heavy quark loop) by $(v/v_1)^2$ and to $gg \to H^oH^o$ by $(v/v_1)^4$. The latter two processes have been calculated in the limit $m_t \gg m_H$, $\sqrt{s}$ by Gavela and Gunion[25] with encouraging results (Fig. 2). (The $gg \to nH^o$ cross section is also enhanced by very heavy quarks in the minimal model[25].)

Charged Higgs

Sufficiently light charged Higgs bosons may be produced via $e^+e^- \to \gamma$, $Z \to H^+H^-$ or in $t \to bH^+$ decay, with $H^+ \to c\bar{s}$ or $t\bar{b}$, depending on the masses.

Heavier charged Higgs bosons may be very difficult to produce or identify. In models with Higgs doublets only the dominant decay modes are expected to be $H^+ \to t\bar{b}$ and/or $H^+ \to W^+H^o$ (if it is kinematically allowed). The former decay should generally be swamped by QCD backgrounds unless t and b jets can be identified with high efficiency. The W^+H^o mode is more promising, especially if the H^o is heavy enough to decay into W^+W^-.

Models with doublets and triplets have the

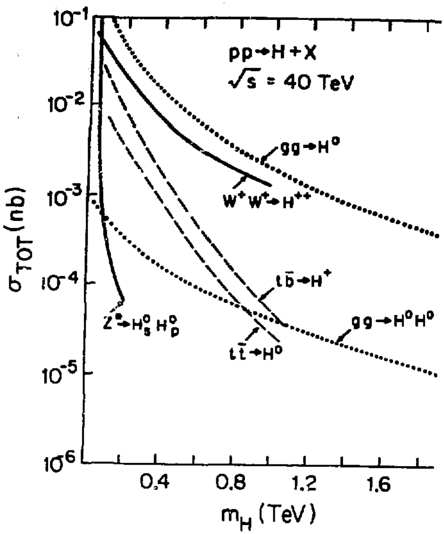


Figure 2. Cross sections for the production of nonstandard Higgs bosons. The Drell-Yan cross sections[24] for $Z^* \to H^+H^-$ and $W^* \to H^+H^o$ are similar to $Z^* \to H_s{}^oH_p{}^o$ and all assume equal mass scalars. The gg cross sections[25] assume $m_t > m_H$, $\sqrt{s}$ and should be multiplied by $(\nu/\nu_1)^2$ for H^o or $(\nu/\nu_1)^4$ for H^oH^o. The $t\bar{b} \to H^+$ and $t\bar{t} \to H^o$ cross sections assume $m_t = 30$ GeV and should be multiplied by $(\nu/\nu_1)^2$. The $W^+W^+ \to H^{++}$ cross section[5] should be multiplied by $b^2/(8b^2 + a^2)$.

advantage that the H^+ can decay into W^+Z (if kinematically allowed), or $H^{++} \to W^+W^+$. The H^{++} decay is one example of interesting physics for which good sign identification is needed.

Charged Higgs bosons may be produced by the Drell-Yan processes $pp \to Z^* \to H^+H^-$ or $pp \to W^* \to H^+H^o$. Unfortunately, the cross sections are small[24] above $m_H \sim 200$ GeV (Fig. 2) and QCD backgrounds(eg. to the $t\bar{b}t\bar{b}$ final state) are likely[21] to be severe (O(nb)). More detailed estimates of the backgrounds, of the possibility of identifying t and b jets, and of the Drell-Yan cross section for the case $m_{H_1} \neq m_{H_2}$ would be very helpful.

Other pair production mechanisms are $gg \to H^+H^-$, which is proportional to $(\nu_2/\nu_1)^4$ and is likely to be similar to $gg \to H^oH^o$ (ref. 25), and (W_LW_L or $Z_LZ_L \to H^+H^-$ or $W_L^+Z_L \to H^+H^o$ (L represents longitudinal). Byers[26] has estimated the cross section for $pp \to H^oH^o + X$ via $W_LW_L \to H^oH^o$ (which is expected to dominate) for the standard Higgs (the H^+H^- and H^+H^o cross sections are presumably similar[27]). Using the amplitudes of Lee, Quigg, and *Thacker*[10] *and the effective* W_LW_L luminosities of Dawson[28], she finds the discouragingly low value $\sigma \sim 10^{-6}$nb $(\sqrt{s}/40 \text{ TeV})^2$.

Charged Higgs bosons may also be produced singly. Machacek[5] has estimated the cross sections due to subprocesses such as $W^+Z \to H^+$ or $W^+W^+ \to H^{++}$ in the doublet/triplet model. The cross sections are typically of the same order as the $W^+W^- \to H^o$ cross section of the minimal model (Fig. 1), assuming that unknown suppression factors (ratios of VEVs) are not too small. The $pp \to H^{++} + X$ cross section is shown in Fig. 2.

There are no W^+ZH^- vertices in n doublet models, but charged Higgs bosons may be produced by direct $t\bar{b}$ fusion (the expected t and b distribution functions at $\sqrt{s} = 40$ TeV are expected to be small but not negligible[22]). In the two doublet model with $\nu_1 < \nu_2$ both the $t\bar{b} \to H^+$ and $t\bar{t} \to H^o$ vertices are enhanced by $|\nu_1/\nu_2| < \sqrt{2}m_{H^+}/m_c$. Estimates of the resulting cross sections, using the Eichten et al.[22] luminosity functions, are shown for $m_t = 30$ GeV in Fig. 2. Even for $\nu_2 = \nu_1$ the cross sections are reasonably large, suggesting that detection may be possible if $H^+ \to W^+H^o$. As usual, the backgrounds are severe if $H^+ \to t\bar{b}$. If one optimistically assumes that the t and b jets can be reliably identified the dominant background is from the QCD subprocess $t\bar{b} \to t\bar{b}$, for which the signal to background ratio is of order $10(\nu_2/\nu_1)^2$ $(m_{H^+}/100 \text{ GeV})(10 \text{ GeV}/\Delta)$, where Δ is the experimental mass resolution for the $t\bar{b}$ pair. (The background from the Drell-Yan process $u\bar{d} \to W^* \to t\bar{b}$ is an order of magnitude smaller for the Eichten et al[22] structure functions.) That is not by itself too discouraging. However, the cross section for $pp \to 2$ jets + X is of order 100 nb (Δ/10 GeV), so even small inefficiencies in t and b identification could be disastrous. A more detailed analysis of these issues would be desirable.

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STRONGLY INTERACTING W's AND Z's

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Introduction

The study focussed primarily on the dynamics of a strongly interacting W, Z (SIW) sector,[1] with the aim of sharpening predictions for total W, Z yield and W, Z multiplicities expected from WW fusion[1,2] for various scenarios. Specific issues raised in the context of the general problem of modeling SIW included the specificity of the technicolor (or, equivalently, QCD) model, whether or not a composite scalar model can be evaded, and whether the standard model necessarily implies an $I = J = 0$ state ($\simeq$ Higgs particle) that is relatively "light" ($M \lesssim$ hundreds of TeV). The consensus on the last issue was that existing arguments are inconclusive, and I shall not pursue it further in this report. While I shall briefly address compositeness and alternatives to the technicolor model, quantitative estimates will be of necessity based on technicolor or an extrapolation of pion data.

As discussed previously,[1] up to mass dependent effects, S-matrix elements with external longitudinally polarized W's and Z's (W_L, Z_L) are the same as S-matrix elements for their respective unphysical, or "eaten", spinless counterparts (w, z). In the strongly interacting limit where scalar self-couplings are much stronger than gauge couplings, the w, z system possesses an approximate chiral SU(2) symmetry analogous to that of pion chiral dynamics. Modeling the w, z system by scaling the pion system by the ratio

$$v/f_\pi = 250\ \text{GeV}/93\ \text{MeV}, \tag{1}$$

of the parameters that characterize spontaneous chiral symmetry breaking in each case is, by definition, equivalent to a technicolor model for electroweak symmetry breaking with an SU(3) technicolor group.

Another possible source of multi-W, Z production was suggested,[3] namely a strongly coupled Yukawa interaction that would arise in the context of the standard electroweak model if there were very heavy fermions. This might then provide multi-W, Z events via gluon fusion (Fig. 1). Calculations of these processes are in progress,[4,5] in this report I will discuss only the ideas involved.

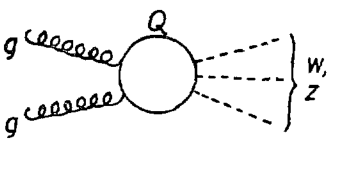

Fig 1 Possible gluon fusion mechanism for multi-W_L, Z_L production

Technirho Production

Modeling the SIW sector on the pion sector would suggest a $J = I = 1$ (W, W), (W, Z) resonant analogue of the $\rho(770)$: this is none other than the technirho for an SU(N) technicolor gauge interaction with $N = 3$. In addition to production[6] via resonance dominated $q\bar{q}$ annihilation, the ρ_T can be produced by the WW fusion process of Fig. 7, Ref. 1. The resulting differential cross section for $pp \to W^+W^- + X$ with a total c.m. energy of 40 TeV as a function of the W-pair invariant mass is shown in Fig. 2 for several values of N. The technirho parameters scale according to[7]

$$\begin{aligned} m_{\rho_T}(N) &= (3/N)^{1/2}\, v\, m_\rho/f_\pi = 2\ \text{TeV}\,(3/N)^{1/2}, \\ \Gamma_{\rho_T}(N) &= (3/N)^{3/2}\, v\, \Gamma_\rho/f_\pi = 0.5\ \text{TeV}\,(3/N)^{3/2} \end{aligned} \tag{2}$$

As N increases the resonance peak moves into a region of higher quark luminosity, but this advantage is eventually compensated for by a more rapidly decreasing width. Whether such scaling behavior, suggested by the large N limit of SU(N) gauge interactions, would be characteristic of a more general class of SIW models remains an open question.

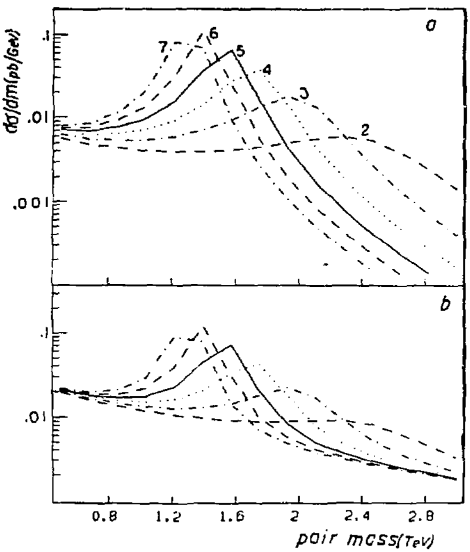

Fig. 2 Cross section for $pp \to W^+W^- + X$ through the W-fusion subprocess $q_1q_2 \to q_1'q_2' + W^+W^-$ via (a) technirho formation and (b) technirho formation plus continuum scattering in the Born approximation. The curves are labeled by the N of SU(N).

The results of Fig 2 are without rapidity cuts. Since qq interactions are as effective as $q\bar{q}$ collisions in WW fusion, and since for high invariant mass and a single partial wave the WW production is rather central, the effects of rapidity cuts should be fairly mild, reducing the signal by a factor of about 3 or less for $y < 1.5$. Comparison of Fig. 2 with Fig. 6-1 of EHLQ[6] suggests that the WW fusion process gives a contribution comparable to that of $q\bar{q}$ annihilation to production of the neutral technirho which EHLQ evaluated for $N = 4$. They found a higher yield for charged technirho production, while WZ fusion to $\rho_T^{\pm}$ should be less than WW fusion to ρ_T° because of the smaller Z-couplings to quarks. Aside from enhancement of a possibly detectable resonance peak in the WW channel, what one is getting is an appreciable excess (as compared with gauge interactions) of multi W, Z events in the high mass region. The contribution of Fig. 2(b) alone, which includes non resonant WW scattering, should give about 1500 W^+W^- events for an integrated luminosity of 10^{40} cm^{-2} in the W^+W^- invariant mass region above 500 GeV. Even without the resolution of a resonance peak, such an excess of W-pair events, if measurable, would signal new strong interactions.

S-Wave Scattering

It was suggested[8] that the s-wave amplitudes for W^+W^- scattering could be modeled by scaling up the measured s-wave $\pi^+\pi^-$ scattering amplitudes. Scaling with Eq. (1) is not entirely unambiguous near threshold, because threshold behavior involves two parameters: f_π or v which measures spontaneous chiral symmetry breaking, and m_π or $m_{W,Z}$ which measures its explicit breaking. Note that the ratios

$$m_\pi^2/f_\pi^2 = 2.27, \tag{3a}$$

$$m_W^2/v^2 = 0.11, \tag{3b}$$

are rather different. The masses are relevant for our purposes only as kinematic quantities that can play a role near threshold and also, for example, in governing phase space available for multi-particle production to be discussed below. In studying threshold behavior, scaling in momentum rather than energy may be the most reliable procedure.[8]

In the limit of a large Higgs mass, $m_H >> \sqrt{s}$, the s-wave born amplitude for $w^+w^- \to zz$ is

$$A_0(w^+w^- \to zz) = \sqrt{2}(A_{00} - A_{02})/3, \tag{4}$$

where A_{JI} is the partial wave amplitude for fixed spin J and "isospin" I:

$$\begin{aligned} A_{00} &= -(8k^2 + 7m^2)/2v^2, \\ A_{02} &= (2k^2 + m^2)/v^2, \end{aligned} \tag{5}$$

where $k = (s - 4m^2)^{1/2}/2$ is the scattering c.m momentum. With the substitutions $v \to f_\pi$ and $m \to m_\pi$ in Eq. (5), the A_{0J} are just the pion s-wave scattering amplitudes obtained[9] from PCAC and current algebra. In that case the terms linear in m_π^2 determine the scattering lengths.[9] In the w, z case m^2 is in fact a meaningless parameter since the substitution $W_L, Z_L \to w, z$ is valid only to order $m_{W,Z}/E_{W,Z}$, and the masses of the w and z are gauge dependent quantities ($m_w = m_W/\xi$). The ratio (3b) assures us that $m_{W,Z}^2$ corrections are negligible in the region $\sqrt{s} \gtrsim 500$ GeV $= 2v$ in which we are interested.

Unitarity of the S-matrix requires that in the region of negligible inelasticity, the partial wave scattering amplitudes are of the form:

$$A_{JI} = -16\pi E\, a_{JI}/k, \tag{6}$$

with

$$a_{JI} = \sin\delta_{JI}\exp(i\delta_{JI}), \quad E = \sqrt{s}/2. \tag{7}$$

Since A_{JI} does not diverge for $k \to 0$, we have

$$a_{JI} \xrightarrow[k\to 0]{} \delta_{JI} \propto k, \tag{8}$$

and the threshold behavior required by chiral symmetry is reproduced for any parameterization of δ_{JI} such that

$$\delta_{JI} \xrightarrow[k\to 0]{} (a_{JI})_{Born} = -k(A_{JI})_{Born}/16\pi E, \tag{9}$$

where the $(A_{JI})_{Born}$ for $J = 0$ are given in Eq. (5).

A standard unitarization procedure is the K-matrix formalism (we take $J = 0$) which defines the phase shift by:

$$\exp(2i\delta_K{}^I) = [1 + i(a_{0I})_{Born}]/[1 - i(a_{0I})_{Born}]. \tag{10}$$

Both unitarity and chiral symmetry will also be satisfied if we take instead the phase shifts

$$\delta_b{}^I = (a_{0I})_{Born}. \tag{11}$$

The s-wave intensity I for $\pi^+\pi^- \to \pi^\circ\pi^\circ$, defined as

$$I = |a_{00} - a_{02}|^2 = 9k^2\sigma(+- \to 00)/8\pi\ , \tag{12}$$

has been measured by Cason et al[10] (Fig. 3(b)). From this they extracted values of the $I = 0$ s-wave phase shift using as input a parameterization of the $I = 2$ s-wave phase shift:

$$\delta_{02} = -(k/1.1\ \text{GeV})/[1 + (k/1.12\ \text{GeV})^2]. \tag{13}$$

The resulting data points for δ_{00} are shown in Fig 3(a) along with the parameterizations (10) and (11) of the phase shifts. For comparison we also show a simple linear extrapolation from the current algebra values[9] of the scattering lengths:

$$\delta_L{}^I = -k[A_{0I}(k=0)]_{Born}/16\pi m_\pi. \tag{14}$$

That the data is better reproduced (see especially Fig. 3(b)) by the parameterizations (10) and/or (11) indicates that the k^2 terms in (5), that are the only ones relevant to the w, z system, are indeed accurately reproduced by the data.

We note from Fig. 3(a) that the "input" parameterization for δ_{02} is close to the K-matrix parameterization (10), while the extracted $I = 0$ phase shift agrees better with the simple Born parameterization of Eq. (11). We therefore include in the intensity parameterizations, Figs. 3(b) and 3(c), one using (11) for δ_{00} and (10) for δ_{02}. This appears to give a reasonable fit to the pion intensity, although the K-matrix prescription (10) may be better for relatively low k.

The relevant lesson is for the $ww \to zz$ intensity, shown in Fig. 3(c), where the Born amplitude and the intensity expected for a Higgs of mass 1 TeV are shown for comparison. Total rates for multi-W, Z production via the WW fusion process have been estimated[11] by interpolating between the Born approximation in the limit $m_H >> \sqrt{s}$ and a 1 TeV Higgs, giving 3,000 to 10,000 events for an integrated luminosity of 10^{40} cm^{-2}. The unitarized s-wave amplitudes shown in Fig. 3(c) suggest only a slightly reduced yield with respect to the Born approximation in the region $\sqrt{s} \sim$ TeV where quark luminosities are most significant. Taking into account contributions from other partial waves (e.g. $J = 1$ resonance production as discussed above), these estimates are probably not overly optimistic, but we have, unfortunately, found no reason to suspect that they are overly pessimistic.

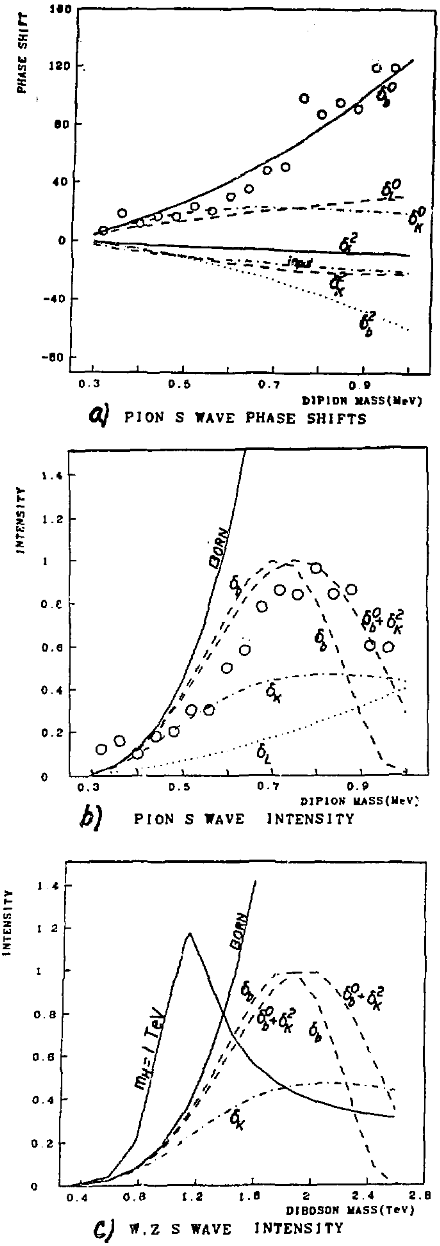

Fig. 3 Unitarized s-wave scattering intensities for $\pi^+\pi^- \to \pi^0\pi^0$ (b) and $W^+W^- \to ZZ$ (c) for various parameterizations of the phase shifts (a) consistant with low energy constraints. Circles are data points;[10] the $I = 0$ pion phase shift was obtained using an input parameterization for the $I = 2$ phase shift. Also shown are extrapolations of the Born, or soft pion, intensities.

Alternative Models of SIW

An alternative[12] to the QCD/technicolor model is the ultracolor of Georgi and collaborators: in the limit in which the scale where ultracolor becomes strong is close to the electroweak symmetry breaking scale (which, for viable models implies an additional strong, gauged axial U(1) coupling) ultracolor models resemble technicolor models, but with a richer spectrum of bosonic states. (An interesting feature of these models is that there are no baryonic states, so they are distinctly different from technicolor models.)

The class of ultracolor models that might provide viable models for SIW are those in which ultrafermions fall into real representations of the ultracolor gauge group. In the minimal model of this class, left handed fermions form a (2, 2) and a (1, 1) of the $SU(2) \times SU(2)$ of the (here composite) scalar sector. These five left handed fermionic degrees of freedom define a 5-plet of "ultraflavor" U(5). The fermion condensates of the strongly coupled, gauged SO(N) spontaneously break this U(5) flavor symmetry to an SO(5) flavor symmetry, implying the existence of 15 Goldstone bosons that transform under $SU(2) \times SU(2)$ as:

$$0^-: (2,2) + (1,1) + (3,3) + (1,1). \tag{15}$$

The above fermionic condensate, in contradistinction to conventional technicolor models, does not break the electroweak $SU(2)_L \times U(1)$ gauge symmetry. To achieve electroweak symmetry breaking an axial U(1) gauge interaction is introduced that explicitely breaks flavor U(5) to flavor U(4). The fermion condensate arising from the strong SO(N) gauge interactions now breaks flavor U(4) to flavor SO(4), leaving only 10 Goldstone bosons, namely the (3, 3) + (1, 1) of Eq. (15) above. This means that the (2, 2) + (1, 1) are not decoupled in *the zero-momentum limit. In particular, the (2, 2), which has the* electroweak quantum numbers of the conventional complex Higgs doublet, can acquire a negative squared mass and trigger the breaking of the electroweak gauge symmetry. This scenario has a well defined set of "low lying" resonances which is richer than that of minimal technicolor models. In addition to the (3, 3) + (1, 1) in (15) that might be relatively "light" ($m << $ TeV) because of their pseudogoldstone boson nature, there are 10 ground state spin one bosons:

$$1^-: (3,1) + (1,3) + (2,2), \tag{16}$$

that might have masses in the TeV region.

Since this scenario does not reduce simply to a v/f_π scaling of QCD, quantitative predictions of resonance parameters have not been attempted. The ultracolor alternative does however have in common with technicolor an underlying theory of fermions with strong gauge interactions, and it is anticipated that masses and widths should scale in a similar way with the number of gauged fermionic degrees of freedom.

Can we evade[13] fermions altogether as underlying constituents of a strongly interacting scalar sector? A pure scalar field theory is known to be inconsistant.[14] *On the other hand* no one would take seriously the notion that the scalars of SIW can be treated as an isolated system at arbitrary energies; the practical implications of difficulties of a pure scalar need not become manifest below the Planck scale.[15] Scalars can presumably be consistantly embedded within a supersymmetry/supergravity context, and pushing the scale of supersymmetry breaking up to the Planck mass poses no practical problem in this respect. Similarly scalars may be composite, but, again this could be relevant only at distances of the order of the

Planck length where perhaps fermions and even gauge bosons would also appear as composite. On the other hand, the results presented by Manton[16] may suggest that the elementary scalar sector of the standard electroweak model will cease to be a sensible description at a scale between 7 and 13 TeV.

The bottom line is that we have no real guidance. This makes experimental investigation of the TeV energy region all the more imperative and all the more exciting.

Multiplicity

For lack of a better guide we will proceed to further modeling with what we know how to do. Estimates[11] of multiplicities for an SIW sector have been made on the basis of massless phase space and assuming a factor $\sqrt{s}/v$ in amplitude, inspired by chiral symmetry, for each emitted W_L, Z_L. For processes involving couplings to "parity violating" weak external sources (W_T, Z_T, heavy quarks) this gives:

$$\sigma(n+1)/\sigma(n) = s/[16\pi^2 v^2 n(n-1)]. \tag{17}$$

For purely "strong" effects that govern the WW fusion process, "parity" is conserved, and only an even number of W_L, Z_L can be produced; this gives

$$\sigma(n+1)/\sigma(n-1) = (s/16\pi v^2)^2/[n(n-1)(n-2)]. \tag{18}$$

Ellis[17] has done a careful calculation of the "technirho" decay branching ratio using the constraints of current algebra and PCAC. The analogous calculation[18] for the ρ gave

$$\Gamma(\rho \to 4\pi^\pm)/\Gamma(\rho \to 2\pi^\pm) = 2 \times 10^{-5}. \tag{19}$$

Scaling this result according to $f_\pi \to v$, $m_\pi \to m_W$ and $m_\rho = 700$ MeV $\to M_{\rho_T} = 1800$ GeV (900 GeV), gives

$$\Gamma(\rho_T \to 4w)/\Gamma(\rho_T \to 2w) = 3.7(1.9) \times 10^{-3}, \tag{20}$$

to be compared with the prescription of Eq. (18) which gives $8.9\ (0.6) \times 10^{-3}$. Since the ρ_T decay involves p-waves, it is not surprising that phase space alone is inadequate, but the latter estimate is not off by an order of magnitude.

Ellis[17] applied the same techniques to calculate the $\sigma(2w \to 4w)/\sigma(2w \to 2w)$ cross section ratio in the Born approximation for $m_H >> s$. In this case there is a large $J = 0$ contribution, and the estimate using the prescription (18) is fairly accurate, as can be seen in Fig. 4. This unfortunately lends confidence to the estimates of Chanowitz and myself[11] who found 10 - 75 four-body W_L, Z_L events for $\int dt\, L = 10^{40}$ cm^{-2} with $\hat{s} \geq 0.5$ TeV from the WW fusion process for various parameterizations of the WW total cross section for $m_H \geq 1$ TeV.

At sufficiently high energies one expects the scaling law (18), or the chiral symmetric Born amplitudes, to break down. Again one can model[8,17] the SIW sector by scaling up the pion sector in the resonance region. This could underestimate SIW multiplicities, because, as indicated by the ratios (3), if we scale according to $s_W = (v/f_\pi)^2 s_\pi$, the available phase space for multi-W, Z production at s_W exceeds that for multipion production at s_π, because of accidents of mass values. Jaffe[8] pointed out however that 4π production does not become significant below multi-resonance thresholds, where the pion mass is itself insignificant. In other words the principle source of high multiplicity pion production appears to be resonance decay, with the primary interactions always being $2 \to 2$ scattering, or $1 \to 2$ decays. This feature may be specific to the underlying QCD structure and duality diagram prescription that it implies. So again the question arises: can we evade[13] an elementary fermion model?

Table 1 shows the number of 4-body W_L, Z_L events for $\int dt\, L = 10^{40}$ cm^{-2} expected from resonance production by W^+W^- fusion for various resonance parameters, assuming a product of branching ratios in the range

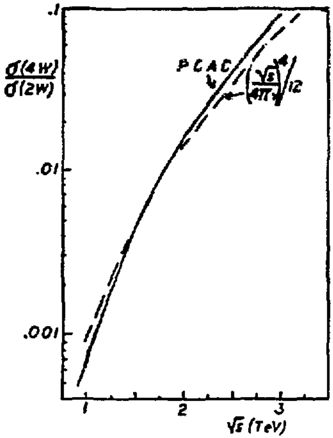

Fig. 4 Fraction of 4W production in W^+W^- scattering as predicted from chiral symmetry,[17] compared with estimates based on phase space and the $\sqrt{s}/v$ rule.[11]

$$0.1 \leq B(4w) \cdot B(2w) \leq 0.25 \tag{21}$$

The resonance parameters are obtained by scaling pion resonance parameters using the prescription (2) for $N = 3$ (QCD) and $N = 7$. Only if the scaled up version of the $\rho(7,\omega)$ has a substantial branching ratio into 4-body final states (not the case for the pion system, nor anticipated by the chiral symmetry estimate, Eq. (20)) can we expect more than a handful of 4-body events from these processes. Hopefully the real SIW will *not mimic* QCD *so closely.*

Table 1 Expected number of 4-body W_L, Z_L events from resonance production by WW fusion in pp collisions with $\int dt\, L = 10^{40}$ cm^{-2} and $E_{cm} = 40$ TeV, assuming Eq. (21).

Low energy model (MeV)	"Technicolor" extrapolation (GeV) $N = 3$	$N = 7$
$g\ 1^+(3^-)$	$g_T\ 1^+(3^-)$	$g_T\ 1^+(3^-)$
m = 1691	m = 4546	m = 2976
Γ = 200	Γ = 538	Γ = 201
$B_2 \cdot B_4 = 0.17$	< 0.3 event	5 - 12 events
$\rho'\ 1^+(1^-)$	$\rho_T'\ 1^+(1^-)$	$\rho_T'\ 1^+(1^-)$
m = 1600	m = 4301	2815
Γ = 300	Γ = 806	Γ = 301
$B_2 \cdot B_4 \leq 0.3$	1.5 to 4 events	8 to 20 events
$\rho\ 1^+(1^-)$	$\rho_T\ 1^+(1^-)$	$\rho_T\ 1^+(1^-)$
m = 770	m = 2070	m = 1355
Γ = 154	Γ = 414	Γ = 155
$B_2 \cdot B_4 = 0$	100 - 280 events	230 - 580 events

Strongly Interacting Yukawa Sector

If there are very heavy fermions that acquire masses via Yukawa couplings to the standard model Higgs doublet, then the physical Higgs and the longitudinally polarized vector bosons W_L, Z_L are strongly coupled to the heavy fermions. This has been suggested[3] as a possible source of multiple Higgs production, especially relevant to the case of an elusive "intermediate" mass ($40\ \mathrm{GeV} \lesssim m_H \lesssim 2m_W$) Higgs boson,[19] via gluon fusion through a heavy quark loop as in Fig. 1.

In the scenario considered here, $m_H \gtrsim 1$ TeV, the process of Fig. 1 could under certain circumstances provide an additional source of multi-W-Z events. Calculations of the general multi-boson loops are underway;[4,5] here I shall outline the physical principles involved. The strong sector of the theory is now defined by the scalar and heavy fermion sectors including Yukawa couplings. The scalar sector alone possesses as before a chiral SU(2) symmetry, and S-matrix elements for the "pseudo-scalars" $w^{\pm}$, z are equivalent to S-matrix elements for $W_L^{\pm}$, Z_L up to $O(m_W/E)$ corrections. To the extent that the heavy quarks are pair-wise degenerate, i.e. that their Yukawa couplings are invariant under suitably defined chiral SU(2) transformations, the full strongly interacting sector is chiral SU(2) invariant. As the Goldstone bosons of this spontaneously broken chiral symmetry, the $w^{\pm}$, z (or $W_L^{\pm}$, Z_L) must decouple at zero momentum. On the other hand, if there is substantial splitting for some quark doublet, this represents an explicit breaking of SU(2) × SU(2), and the Goldstone theorem need not apply.

To illustrate, consider gluon fusion to a single (real or virtual) z (or Z_L). Working in a renormalizable gauge the z couples to the heavy quark through a pure pseudoscalar coupling proportional to the third component of the quark weak isospin. The calculation of the gluon fusion process reduces to the Steinberger[20] calculation of $\pi^\circ \to \gamma\gamma$, giving the well known result that the amplitude drops rapidly to zero for $m_Q^2 < p_z^2/2$, and takes a non-zero value essentially independent of m_Q for $m_Q^2 \gtrsim p_z^2$. Since the two components of a quark weak isodoublet contribute with opposite signs, there is no net contribution unless they satisfy a mass relation $m_{Q1}^2 \leq p_z^2 \lesssim m_{Q2}^2$. Roughly, one obtains a contribution proportional to:

$$\Sigma_H I(Q_H) = \Sigma_H J(U_H) - \Sigma_H J(D_H). \qquad (22)$$

where U_H and D_H are "heavy" ($m_H^2 \gtrsim p_z^2$) quarks of charge 2/3 and $-1/3$ respectively and $I(U) = J(U)$ ($I(D) = -J(D)$) is the contribution to the amplitude of Fig. 1 from an external quark $Q = U$ or D.

Alternatively, one may transform the fields to obtain the non-linear σ-model formulation, in which case the w, z appear only with derivative couplings and couple to quarks through the derivative of an axial vector coupling. In this case w, z amplitudes vanish explicitly as $p/m_Q \to 0$ unless there are anomalies (or explicit chiral symmetry breaking). In this formulation, the calculation reduces to the more modern calculations[21] of $\pi^\circ \to \gamma\gamma$ (and is equivalent to a direct calculation of $gg \to Z_L$ in the unitary gauge). For "light" quarks (L) only the anomaly $A(Q)$ contributes, while for $m_Q^2 \gtrsim p_z^2$, the anomaly exactly cancels the mass dependent contribution $I(Q_H)$, giving a contribution to $gg \to z$ proportional to

$$\Sigma_L A(Q_L) = -\Sigma_L A(Q_H) = \Sigma_H I(Q_H), \qquad (23)$$

which is the same as (22); the first equality in (23) holds because the theory is by construction anomaly free when summed over all quarks.

We are really interested in multi W_L, Z_L production. Up to mass dependent effects, the production of an odd number of w, z is completely determined[22] by the axial anomalies, and, roughly, a non-zero amplitude should be found if some quark doublet satisfies $m_{Q1}^2 \lesssim s \lesssim m_{Q2}^2$, where $\sqrt{s}$ is the total cm energy of the di-gluon system. We already know from the strength of the neutral current couplings and the Z and W masses that quark doublet mass splittings cannot exceed a few hundred GeV. However, when mass dependent effects are included,[5] there may be some window for an observable effect.

Gluon fusion to a "parity-even" system of w, z and Higgs is governed by the trace anomaly rather than the axial anomaly, and the presence of a trace anomaly does not invalidate the soft meson theorems of chiral symmetry. Since gluon fusion into, say, a pair of z's is proportional to $(\tau_3)^2$ for each quark, no cancellation can occur between members of a doublet, so the decoupling theorem must hold for each quark loop separately, and it is probable that gluon fusion to 2n(w, z) alone will be suppressed as s/m_H^2 or s/m_Q^2.

It therefore seems likely that the gluon fusion process is most promising as a source of multiple "light" Higgs,[19] or, possibly for $2m_W < m_H << \mathrm{TeV}$, a source of multi-W events via Higgs decay. Results of explicit calculations[4,5] will give a more precise answer.

Conclusions

The lesson for SSC experimentation, is, as before,[1] that identification of W's and Z's is a crucial issue. Further questions that should be pursued include:

At what level of production can multi ($N \geq 3$) W, Z events be detected? This requires more serious study of multi-jet ($N_{jet} \geq 3$), W, Z + multi-jet, 2(W, Z) + jet, etc. backgrounds, as well as the effects of rapidity cuts on various classes of events.

Are the general features of a strongly interacting W, Z sector discernable without event-by-event identification? Signals include an enhanced W, Z yield, an enhanced Z/W ratio, and an enhanced component of longitudinally polarized Z's in the tail ($\sqrt{\hat{s}} \gtrsim 500$ GeV) of the effective center of mass spectrum. The question is whether these effects are sufficiently pronounced that deviations from purely gauge interaction effects could be extracted by comparing, say, events containing one or more leptonic decays with total transverse energy greater than or less than 500 GeV.

Acknowledgment

In addition to lively discussions among workshop participants, the ideas and results presented here benefited from discussions and consultations with I. Bars, P. Binétruy, S. Dawson and EHLQ

This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy Contract DE-AC03-76SF00098 and in part by the National Science Foundation under Research Grant PHY-84-06608

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SUPERSYMMETRY AT THE SSC

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Introduction

The topics studied by this working group can be grouped under four main headings:

(A) *Conventional SUSY phenomenology in which there is a multiplicatively conserved quantum number R and the lightest sparticle is a stable photino* $\tilde{\gamma}$ *which provides a missing* p_T *signature.*

(B) Alternative SUSY phenomenologies, including those with a different lightest sparticle such as the sneutrino $\tilde{\nu}$,[8] Higgsino $\tilde{H}$ or gluino $\tilde{g}$,[9] and models in which R parity is broken so that the photino becomes unstable.[10] We studied how much of the conventional missing p_T signature would survive in such models.

(C) What if some sparticle weighs less than 100 GeV and has been detected at the Sp$\bar{p}$S Collider or at the FNAL Tevatron Collider before the SSC comes into operation? We studied the distrubutions of light $\tilde{q}$ or $\tilde{g}$ partons in the proton,[11] and the problems of distinguishing the missing p_T signatures of light and heavy sparticles.

(D) SUSY Higgses, of which there must be two doublets whose masses are constrained.[12,13] Moreover, once two masses are known, all other physical Higgs masses and couplings are fixed.[14]

In what follows, topics (A) and (B) are reviewed in some detail, while topics (C) and (D) receive less attention.

Conventional SUSY Phenomenology

In most SUSY models there is a multiplicatively conserved quantum number $R = +1$ for all conventional particles (q, L, g, W, γ, H) and $R = -1$ for all their spartners ($\tilde{q}$, $\tilde{\ell}$, $\tilde{g}$, $\tilde{W}$, $\tilde{\gamma}$, $\tilde{H}$). R conservation would imply that sparticles are always produced in pairs (e.g. $pp \to \tilde{q}\bar{\tilde{q}}$ or $\tilde{q}\tilde{g}$ or $\tilde{g}\tilde{g}$ or ...), that every sparticle decays into another sparticle (e.g. $\tilde{q} \to q\tilde{\gamma}$), and therefore that the lightest sparticle is stable. There are cosmological reasons to expect this lightest sparticle to be neutral and not strongly interacting, with the favoured candidate being the $\tilde{\gamma}$ discussed here.[15] In such a SUSY model there is a missing p_T signature as photinos escape after being produced in the decays of heavier sparticles, e.g. $\tilde{q} \to q\tilde{\gamma}$, $\tilde{g} \to q\bar{q}\tilde{\gamma}$, $\tilde{W} \to f\bar{f}\tilde{\gamma}$, etc. Total cross sections are available for the pair-production of $\tilde{g}\tilde{g}$, $\tilde{q}\bar{\tilde{q}}$, $\tilde{q}\tilde{g}$, $\tilde{W}\tilde{g}$, $\tilde{\gamma}\tilde{g}$, and $\tilde{\ell}\tilde{\ell}$.[1,2,3,4,5,6] Detailed final state Monte Carlo calculations are available for the first two subprocesses,[3,4] and one in progress for the next three.[6] The results are not reproduced in detail here, since they are already available elsewhere.[4] Suffice to say that

$$\sigma_{\tilde{g}\tilde{g}} \geq \frac{1}{10}\ \text{pb} \quad \text{for} \quad m_{\tilde{g}} \lesssim \begin{cases} 1 \\ 1{\cdot}6 \\ 2{\cdot}4 \end{cases} \text{TeV at } E_{cm} = \begin{cases} 10 \\ 20 \\ 40 \end{cases} \text{TeV} \tag{1}$$

Conventionally, one expects the decay $\tilde{g} \to q\bar{q}\tilde{\gamma}$ to dominate, resulting[4] in a large missing p_T signature ($\not{p}_T = O(m_{\tilde{g}}/2)$) and ideally four final state jets. For $m_{\tilde{g}} = 1$ TeV and $E_{cm} = 20$ TeV we find that they emerge with rapidities y: $\langle y^2 \rangle^{1/2} \sim 1$, angular separations $\Delta\theta \sim 1.5$ radians, and the average of the minimum p_T jet being O(140) GeV. In the case of $\tilde{q}\bar{\tilde{q}}$ production, we find[4]

$$\sigma_{\tilde{q}\bar{\tilde{q}}} \geq \frac{1}{10}\ \text{pb} \quad \text{for} \quad m_{\tilde{q}} \lesssim \begin{cases} 0{\cdot}9 \\ 1{\cdot}4 \\ 2{\cdot}0 \end{cases} \text{TeV at } E_{cm} = \begin{cases} 10 \\ 20 \\ 40 \end{cases} \text{TeV} \tag{2}$$

and ideally 2(6)-jet final states if $\tilde{q} \to q\tilde{\gamma}$ ($\tilde{q} \to q\tilde{g}$, $\tilde{g} \to q\bar{q}\tilde{\gamma}$) decays dominate, in both cases with a large missing p_T signature. In the latter case, we find[4] that the jets emerge with $\langle y^2 \rangle^{1/2} \sim 1$, $\Delta\theta \sim 1$ radians and the average of the minimum jet p_T is 90 GeV, if we take $m_{\tilde{q}} = 1$ TeV and $m_{\tilde{g}} = 700$ GeV at $E_{cm} = 20$ TeV as an illustrative example. The observable cross-sections are not greatly reduced[4] if one restricts oneself to events with $\not{p}_T > 50$ GeV. Moreover, the missing p_T vector is generally isolated in azimuthal angle, having components p_{TT} transverse to the observable q and $\bar{q}$ jet axes which are much greater than 100 GeV on average. This makes heavy sparticles easy to distinguish from lighter sources of missing p_T, such as the t quark which gives $|p_{TT}| \leq m_t/2 \approx 20$ GeV. In the case of $\tilde{\ell}^+\tilde{\ell}^-$ production, the γ and Z^0 production diagrams give[2]

$$\sigma_{\tilde{\ell}^+\tilde{\ell}^-} \approx \begin{cases} 10^{-37}\text{cm}^2 \\ 10^{-38}\text{cm}^2 \end{cases} \quad \text{for} \quad m_{\tilde{\ell}^\pm} = \begin{cases} 180 \\ 400 \end{cases} \text{GeV}$$

$$\text{at } E_{cm} = 40\ \text{TeV}. \tag{3}$$

The expected decays $\tilde{\ell}^\pm \to \ell^\pm\tilde{\gamma}$ would give $\ell^+\ell^- + \not{p}_T$ events. Among the possible backgrounds are heavy lepton pair L^+L^- production, which has a cross-section $\geq 4 \times \sigma_{\tilde{\ell}^+\tilde{\ell}^-}$. Heavy leptons with masses $m_{L^\pm} \geq 100$ GeV are expected to decay into $\nu + W^\pm$: $W^\pm \to q\bar{q}$ or $\ell^\pm\nu$. *The latter mechanism yields* $\ell^+\ell'^-$ *pairs of uncorrelated* flavours and with relatively low invariant mass. Since each $W^\pm \to \ell^\pm\nu$ decay has an 8% branching ratio, this is not a substantial background. The cross-section for W^+W^- production is larger: $\sigma_{W^+W^-} \sim 10\ \sigma_{\tilde{\ell}^+\tilde{\ell}^-}$, but again the $W^\pm$ decays give $\ell^+\ell^-$ of uncorrelated flavours, *and the two branching ratio factors yield a* signal-to-background ratio of O(10). One can also produce $\tilde{\ell}^\pm\tilde{\nu}$ pairs via the $W^\pm$, but $\sigma_{\tilde{\ell}^\pm\tilde{\nu}} < 1/2\ \sigma_{\ell^\pm\nu}$, and this process seems unlikely to be observable if the $\tilde{\nu}$ decays invisibly.[16] Even for $\tilde{\nu} \to$ visible decay products there are many backgrounds.

Another possible production mechanism for heavy strongly interacting sparticles is diffractive,[4,7] with a $\tilde{g}\tilde{g}$ or $\tilde{q}\bar{\tilde{q}}$ pair produced in the forward direction with low p_T, while the opposite hemisphere is "quiet". One expects a cross-section for such a mechanism of the general form

$$\sigma = \frac{1}{m^n}\ f(m^2/s) \tag{4}$$

where m is the heavy particle mass. The form of the function $f(m^2/s)$ is unknown: the power n is expected to be ≥ 2, with recent opinion[17] favouring $n = 4$. The most optimistic extrapolation using present accelerator charm production data, $n = 2$, and working at fixed $m/\sqrt{s}$ would give

$$\left.\begin{array}{l} m_c = 1\cdot 5\ \text{GeV} \\ \sqrt{s} = 20\ \text{GeV} \\ \sigma \sim (1\ \text{to}\ 10)\mu\text{b} \end{array}\right\} \longrightarrow \left.\begin{array}{l} m_t = 40\ \text{GeV} \\ \sqrt{s} = 540\ \text{GeV} \\ \sigma = (1\ \text{to}\ 10)\text{nb} \end{array}\right\} \longrightarrow$$

$$\left.\begin{array}{l} m_{\tilde{g}} = 3\ \text{TeV} \\ \sqrt{s} = 40\ \text{TeV} \\ \sigma = (0.1\ \text{to}\ 1)\text{pb} \end{array}\right\} \qquad (5)$$

While the theoretical prognosis[17] for such an extrapolation is not encouraging, it can be checked by looking for diffractive t quark ($m_t = 40$ GeV) production at the CERN pp Collider with $\sqrt{s} = 540$ GeV. A possible way to search for diffractive gg pair production would be to look for events with two observed jets in the forward direction, accompanied by missing p_T. If one then computes the minimum mass m^* of a system decaying into the observed $\underline{p}_{jet_1}$, $\underline{p}_{jet_2}$ and $\not{p}_T$ for arbitrary p_L, one finds[4] that it is sharply peaked at the mass of the produced g, despite the combinatorial background from mixing the decay products of different gluinos. Moreover, it seems that there is no substantial background of missing p_T from unobserved forward-going particles. Searching for evidence of diffractively produced $\tilde{q}\bar{\tilde{q}}$ pairs would be even easier. one would look for a Jacobian peak at $|p_{T_{jet}}| = m_q/2$ from $\tilde{q} \to q\tilde{\gamma}$ decay. Unfortunately, the cross-section for diffractive production of TeV $\tilde{g}$ or $\tilde{q}$ pairs will be negligible if $n = 4$ in Eq. (4), as recently argued.[17]

Alternative SUSY Phenomenologies

The possibility discussed above is the mainstream SUSY phenomenology. During the Workshop there was considerable discussion of alternatives and their implications for the supercollider, addressing in particular the question whether the distinctive missing energy-momentum signature would be lost. The lightest sparticle may not be the photino $\tilde{\gamma}$, but is probably electromagnetically neutral and not strongly interacting.[15] Stable charged and/or strongly interacting heavy particles tend to condense into galaxies, stars, planets, etc. along with ordinary matter, and unsuccessful searches for superheavy isotopes of ordinary nuclei seem to exclude[9,18]

$$5\ \text{GeV} < m < 35\ \text{GeV},$$

$$\text{and possibly}\ 3\ \text{GeV} < m < 1000\ \text{GeV}\ . \qquad (6)$$

This leaves a window for light stable gluinos, if the stable color-neutral hadrons that they form are neutral (e.g. $\tilde{g}g$ or $\tilde{g}(u\bar{u} + d\bar{d})/\sqrt{2}$) rather than charged (e.g. $\tilde{g}ud$). Candidates for the lightest (neutral) sparticle therefore include the gravitino, $\tilde{\gamma}$, $\tilde{\nu}$, $\tilde{H}^0$ and $\tilde{g}$. Even if the gravitino is lighter than the others, in most "fashionable" supergravity models the lifetime $\tau(x \to x' + \text{gravitino})$ is much greater than 1 second, so that the next-to-lightest sparticle appears to be effectively stable as far as laboratory experiments are concerned

Cosmology imposes no lower limit on the sneutrino mass:[3] the large annihilation cross-section $\sigma_{\tilde{\nu}\tilde{\nu}}^{ann} \propto 1/m_W^2$ always suppresses the primordial ν abundance to acceptably low levels. In many models $m_{\tilde{\nu}} < m_{\tilde{\ell}}$, since

$$m_{\tilde{\nu}}^2 = O(m_{3/2}^2) - \frac{m_z^2}{2}\ \cos 2\alpha;\ \tan\alpha < 1 \qquad (7)$$

and in some models with a heavy t quark α is very small so that $\cos 2\alpha \to 1$. If $m_{\tilde{\nu}} < m_{\tilde{\gamma}}$ (the case of interest to us here) then one expects $\tilde{\gamma} \to \tilde{\nu}\nu$ decays (with $\tau_{\tilde{\gamma}} = O(10^{-16})\ \text{sec}/m_{\tilde{\gamma}}(\text{GeV}))$ to give <u>the same missing energy-momentum signature as in conventional SUSY</u>. If $m_{\tilde{\gamma}} < m_{\tilde{\nu}} < m_{\widetilde{other}}$ then

$$1 \ll \frac{\Gamma(\tilde{\nu} \to \nu\tilde{\gamma})}{\Gamma(\tilde{\nu} \to eud\tilde{g}\ \text{etc.})} \approx \frac{45}{2}\left(\frac{\alpha}{\alpha_s}\right)\left(\frac{m_W}{m_{\tilde{\nu}}}\right)^8 \qquad (8a)$$

for $m_{\tilde{u}} = m_{\tilde{\ell}} = m_{\tilde{W}}$ much greater than $m_{\tilde{\nu}}$. In this case, invisible neutral decays of the ν dominate, and it has a relatively short lifetime:

$$\tau_{\tilde{\nu}} = \frac{O(10^{-16})\text{sec}}{m_{\tilde{\nu}}(\text{GeV}) \times |F|^2}\ ;\quad |F| = |F(m^2/m_W^2)| = O(1)\ . \qquad (8b)$$

If $m_{\tilde{\nu}} > m_{\tilde{\ell}}, m_{\tilde{q}}$ etc., then "visible" decays $\tilde{\nu} \to eud\tilde{g}$, etc. may compete with, or dominate the "invisible" decay $\tilde{\nu} \to \nu\tilde{\gamma}$.

If $\tilde{H}^0$ (more accurately, a mass eigenstate mixture of $\tilde{W}^0$, $\tilde{B}^0$ and $\tilde{H}^0$) is the lightest sparticle, then one expects $\tilde{\gamma}$ decays of $\tilde{q}$, $\tilde{g}$ etc. to dominate:

$$\Gamma(\tilde{q} \to q\tilde{\gamma}) \gg \Gamma(\tilde{q} \to q\tilde{H}^0),\ \Gamma(\tilde{g} \to q\bar{q}\tilde{\gamma}) \gg \Gamma(\tilde{g} \to q\bar{q}\tilde{H}^0)\ . \qquad (9)$$

These would then be followed by $\tilde{\gamma}$ decays into the $\tilde{H}^0$:

$$\Gamma(\tilde{\gamma} \to \gamma\tilde{H}^0) \sim O\left(\frac{\alpha^3}{10^3}\frac{m_{\tilde{\gamma}}^3}{\tilde{m}^2}\right) \underset{?}{\gg} \Gamma(\tilde{\gamma} \to q\bar{q}\tilde{H}) \sim O\left(\frac{\alpha^2}{10^4}\left(\frac{m_q}{m_W}\right)^2\frac{m_{\tilde{\gamma}}^5}{\tilde{m}^4}\right). \qquad (10a)$$

Therefore one expects dominance of $\tilde{\gamma} \to \gamma\tilde{H}^0$ decays, so that the <u>missing energy-momentum is reduced by 50%</u>, and is <u>accompanied by "prompt" photons</u> ($\tau_{\tilde{\gamma}} < O(10^{-11})$sec) This should be quite a distinctive signature.

It is possible to have a light gluino ($m_{\tilde{g}} < m_{\tilde{\gamma}}$) in some supergravity models where $m_{\tilde{\gamma}}$ and $m_{\tilde{g}}$ arise from loop diagrams and $m_t < O(35)$ GeV.[9] Then the photino is unstable

$$\Gamma(\tilde{\gamma} \to \tilde{g}q\bar{q}) \sim \alpha_s\alpha\ O(10^{-3})\ \frac{m_{\tilde{\gamma}}^5}{m_{\tilde{q}}^4} \qquad (10b)$$

and one expects

$$\tau_{\tilde{\gamma}} \lesssim \tau(m_{\tilde{\gamma}} = 2\ \text{GeV},\ m_{\tilde{q}} = 250\ \text{GeV}) \sim 10^{-11}\ \text{sec}\ . \qquad (11)$$

As mentioned above, Eq. (6), stable particle searches require $m_{\tilde{g}} < 5(3)$ GeV, while the absence of a light ($\tilde{g}g$) hadron suggests $m_{\tilde{g}} > O(1)$ GeV. Between these two limits there is an allowable range where there are stable g hadrons which behave much like neutrons or K_L^0, <u>with no missing energy-momentum signature</u> Thus $\tilde{g}$ hadrons share with them the ability to deposit <u>energy in a hadronic calorimeter</u>, while leaving no charged track in front Therefore it is probably difficult to pick out a light, stable $\tilde{g}$ in high energy hadron-hadron collisions. The best place might be in $\chi_b \to \tilde{g}\tilde{g}X$ decay,[11] where $\tilde{g}$ final states have a branching ratio O(30)% if $m_{\tilde{g}} < 4$ GeV

Another possible scenario is that in which the lightest sparticle is not absolutely stable, because R parity is not multiplicatively conserved. R parity may either be broken explicitly[19] by Lagrangian terms which couple particles to sparticles (e.g. superpotential terms $\delta_i L_i H$ or scalar mass squared terms $\mu_i^2 \tilde{L}_i H$) and/or sponteneously[10] by $<0|\tilde{\nu}_i|0> \neq 0$ (e.g. $v_\tau = <0|\tilde{\nu}_\tau|0> \neq 0$). There are phenomenological upper limits on explicit R breaking terms: neutrino mass limits imply[19]

$$\mu_\tau < 100 \text{ GeV},\ \mu_\mu < 10 \text{ GeV},\ \mu_e < 1 \text{ GeV}, \qquad \text{and } \delta_e < 10^{-2} \text{ GeV} \tag{12}$$

while the upper limit on $\mu \to e\gamma$ decay grazes the limits, Eq. (12). In models with explicit R breaking, the $\tilde{\gamma}$ is unstable, with

$$\Gamma(\tilde{\gamma} \to b\bar{b}\nu)\ (\propto m_{\tilde{\gamma}}^3 m_b^2) > \Gamma(\tilde{\gamma} \to \gamma\nu)\ (\propto m_{\tilde{\gamma}}^5) \tag{13}$$

for $m_{\tilde{\gamma}} > O(10)$ GeV. In this case one has less missing E_T and a possible excess of heavy quarks, whereas a lighter $\tilde{\gamma}$ gives missing energy-momentum associated with a prompt γ.

We have already seen, Eq. (7), how $m^2_{\tilde{\nu}}$ may be less than $m^2_{\tilde{\ell}}$, and if $m^2_{\tilde{\nu}} < 0$, spontaneous R breaking through $<0|\tilde{\nu}_i|0> \neq 0$ ensues. A generic possibility, which appears the most likely,[10] is $<0|\tilde{\nu}_e|0> = <0|\tilde{\nu}_\mu|0> = 0$, $v_\tau = <0|\tilde{\nu}_\tau|0> = O(v, v')$ where $v, v' \equiv <0|H^0, H^{0\prime}|0>$. In this case there is a charged fermion mass matrix mixing the $\tilde{W}^\pm$, $\tilde{H}^\pm$ and $\tau^\pm$:

$$\begin{pmatrix} \tilde{W}^+ \\ \tilde{H}'^+ \\ \tau^+ \end{pmatrix}_L^{T} \begin{pmatrix} M_2 & g_2 v & g_2 v_\tau \\ g_2 v' & \epsilon & 0 \\ 0 & h_\tau v_\tau & h_\tau v \end{pmatrix} \begin{pmatrix} \tilde{W}^- \\ \tilde{H}^- \\ \tau^- \end{pmatrix}_L \tag{14}$$

where M_2 is the SUSY breaking SU(2) gaugino mass, ϵ is a HH' mixing term, and h_τ is the τ Yukawa coupling. In this case the physical charged mass eigenstates

$$\begin{aligned} &\text{"}\tilde{W}^+\text{"} \ni O\left(\frac{m_\tau}{m_W}\frac{v_\tau}{v}\right)\tau^+, \quad \text{"}\tilde{W}^-\text{"} \ni O\left(\frac{v_\tau}{v}\right)\tau^- \\ &\text{"}\tilde{H}^+\text{"} \not\ni \tau^+, \quad \text{"}\tilde{H}^-\text{"} \ni O\left(\frac{v_\tau}{v}\right)\tau^- \\ &\text{"}\tau^+\text{"} \ni O\left(\frac{m_\tau}{m_W}\frac{v_\tau}{v}\right)\tilde{W}^+, \quad \text{"}\tau^-\text{"} \ni O\left(\frac{v_\tau}{v}\right)\tilde{H}^- \end{aligned} \tag{15}$$

and there is analogous mixing for the $\tilde{\gamma}$, H^0 and $\tilde{\nu}_\tau$. There are two light neutral mass eigenstates, "$\tilde{\nu}_\tau$" and a H^0 state, in the limit $\epsilon \to 0$. The $e^+e^- \to \tau^+\tau^-$ forward-backward asymmetry is not substantially modified in such a model, since the only substantial mixing of the τ^- is with the H^- which has the same I_3 and Y, and hence the same Z^0 coupling as the unmixed τ^-. The τ lifetime is unaffected since the weak isospin partner of the τ^- mass eigenstate is essentially a combination of the "ν_τ" and H^0 states, both of which are possible decay products of the τ if ϵ is small. Since the neutral mass matrix has substantial entries mixing states with different I_3 and Y, there can in principle be substantial flavour non-diagonal couplings of the Z^0. In practice, the only substantial one is the $Z^0 - \tilde{\gamma} - \tilde{H}^0$ vertex, which can lead to a decay rate

$$\frac{\Gamma(Z^0 \to \tilde{\gamma}\tilde{H}^0)}{\Gamma(Z^0 \to \bar{\nu}_e\nu_e)} = O\left(m^2_{\tilde{\gamma}}/m^2_Z\right) . \tag{16}$$

In principle, the decays $\tilde{q} \to q\tau$ or $q\nu_\tau$ ($\tilde{g} \to qq\tau$ or $\tilde{g} \to \bar{q}q\nu_\tau$) are now possible because R-parity is broken, but in practice these novel decay modes are suppressed by $O(m_\tau \text{ or } m_q/m_W)^2$ relative to $\tilde{q} \to q\tilde{\gamma}$ ($\tilde{g} \to q\bar{q}\tilde{\gamma}$), as can be deduced from the last line of Eq. (15). One observable novelty is however the decay of the photino: there is no $W^\pm - \tilde{\gamma} - \tau^\mp$ or $Z^0 - \tilde{\gamma} - \nu_\tau$ coupling, but the $Z^0 - \tilde{\gamma} - \tilde{H}^0$ coupling mentioned above, Eq. (16), gives

$$\Gamma(\tilde{\gamma} \to \tilde{H}^0 X : X = e^+e^-, \mu^+\mu^-, \tau^+\tau^-, q\bar{q}, \bar{\nu}\nu) = O\left(G_F \frac{m^5_{\tilde{\gamma}}}{192\pi^3}\right) \times \left(m_{\tilde{\gamma}}/m_Z\right)^2 . \tag{17}$$

If $m_{\tilde{\gamma}} < O(10)$ GeV, the corresponding lifetime could be $\geq O(10^{-11})$ seconds, providing the useful signature of a separated decay vertex if X contains charged particles. If $m_{\tilde{\gamma}} \leq O(10)$ GeV, the decay $\tilde{\gamma} \to \gamma + \nu_\tau$ or $\tilde{H}^0$ may dominate over Eq. (17) providing a γ + missing energy signature. If X is neutral (i.e. $X = \bar{\nu}\nu$), then one has the missing energy signature of SUSY, albeit suppressed by the $\bar{\nu}\nu$ branching ratio relative to its magnitude due to $\tilde{q} \to q\gamma$ decays in conventional SUSY. We conclude that one could have large R and L_τ breaking with $v_\tau/v = O(1)$, that this would in general provide less missing E_T, but that there are possible signatures of a finite decay length for the photino, and/or $\tilde{\gamma} \to \gamma$ + missing energy decays.

What if...

Some sparticles are relatively light (e.g. $m_{\tilde{\gamma}}$ = 1 to 5 GeV,[3] $m_{\tilde{q},\tilde{g}}$ = O(40) GeV).[3] How would sparticle phenomenology at the SSC look in this case? Detailed calculations are available[11] of the evolution of supersymmetric parton densities in the nucleon. For example, one finds the following asymptotic momentum fractions at Q^2 much greater than $m_{\tilde{g}}^2$, $m_{\tilde{q}}^2$:

	Pure QCD	QCD + $\tilde{g}$	QCD + $\tilde{g}$ + $\tilde{q}$
$q + \bar{q}$	$\frac{3N_q}{3N_q + 16}(0{\cdot}53)$	$\frac{3N_q}{3N_q + 20}({\cdot}47)$	$\frac{3N_q}{5N_q + 20}({\cdot}36)$
g	$\frac{16}{3N_q + 16}(0{\cdot}47)$	$\frac{16}{3N_q + 20}({\cdot}42)$	$\frac{16}{5N_q + 20}({\cdot}32)$
$\tilde{g}$		$\frac{4}{3N_q + 20}({\cdot}11)$	$\frac{4}{5N_q + 20}({\cdot}08)$
$\tilde{q} + \tilde{\bar{q}}$			$\frac{2N_q}{5N_q + 20}({\cdot}24)$

corresponding to

$$\frac{\tilde{g} \text{ momentum}}{g \text{ momentum}} \to \frac{1}{4} , \qquad \frac{\tilde{q} \text{ momentum}}{q \text{ momentum}} \to \frac{2}{3} . \tag{18}$$

Thus there would be quite a large number of spartons in the proton at large Q^2. The cross-sections for all the $\tilde{X}X' \rightarrow \tilde{X}''X''$ parton-sparton scattering subprocesses have been calculated[20] in the limit s much greater than $m_{\tilde{X}}^2$, which is the domain in which the renormalization group calculations of the sparton densities are useful. These cross-sections are then comparable with conventional large p_T cross-sections, and hence should be readily observable. The missing p_T signature would be different from that due to heavy sparticle decay, since the missing p_T vector would no longer be isolated in azimuthal angle, but would instead be oriented almost parallel to one of the outgoing jet axes, with a momentum transverse to the jet axis

$$p_{TT} \lesssim (m_{\tilde{q}} \text{ or } m_{\tilde{g}}/2) \ . \tag{19}$$

Thus, one could hope to distinguish light sparticles from heavy sparticles.

SUSY Higgses

Supersymmetric models need at least two Higgs doublets,[12,13] and hence physical charged $H^{\pm}$ scalars, H_1^0 and H_2^0, and a pseudoscalar boson a. The quartic potential terms are fixed by SUSY, so the only free parameters are quadratic:

$$V(H) = m^2|H|^2 + m'^2|H'|^2 + m_3^2(HH' + \text{herm. conj.}) . \tag{20}$$

One combination of these three parameters is fixed by our knowledge of G_F (implying $v^2 + v'^2$ is known), while the other two parameters can be fixed by knowing two Higgs boson masses. For example β: $\tan\beta = v'/v$ is determined from

$$\tan^2 2\beta = \frac{\left(m_Z^2 - m_{H_1^0}^2\right)\left(m_Z^2 - m_{H_2^0}^2\right)}{m_{H_1^0}^2 \, m_{H_2^0}^2} \tag{21}$$

and there are restrictive mass formulae:

$$m_{H^{\pm}}^2 = m_a^2 + m_{W^{\pm}}^2 > m_{W^{\pm}}^2 \tag{22}$$

$$m_{H_1^0,H_2^0}^2 = \left(m_a^2 + m_Z^2 \pm \sqrt{(m_Z^2 - m_a^2)^2 + 4m_Z^2\, m_a^2 \cos^2 2\beta}\right)/2 \tag{23}$$

and one neutral Higgs is guaranteed to be light in many models[13] The couplings of all the Higgs bosons to $W^{\pm}$, Z^0 and fermions are fixed once one knows two Higgs masses and hence v'/v. Introducing θ:

$$\tan 2\theta \equiv \tan 2\beta \left(\frac{m_{H_1^0}^2 + m_{H_2^0}^2}{m_{H_1^0}^2 + m_{H_2^0}^2 + 2m_Z^2}\right) . \tag{24}$$

Relative to the conventional single $H^0 - W^{\pm}W^{\mp}$, $H^0Z^0Z^0$ and $H^0 f\bar{f}$ couplings one has[14]

$$g_{W^+H^-(H_1^0,H_2^0,a)} = \sin(\theta-\beta),\ \cos(\theta-\beta),\ 1 \tag{25a}$$

$$g_{Z^0a(H_1^0,H_2^0)} \propto \sin(\theta-\beta),\ \cos(\theta-\beta) \tag{25b}$$

$$g_{W^+W^-(H_1^0,H_2^0,a)} \propto \cos(\theta-\beta),\ \sin(\theta-\beta) \tag{25c}$$

$$g_{u\bar{u}(H_1^0,H_2^0,a)} \propto \sin\theta/\sin\beta,\ \cos\theta/\sin\beta,\ \cot\beta\gamma_5 \tag{25d}$$

$$g_{d\bar{d}(H_1^0,H_2^0,a)} \propto \cos\theta/\sin\beta,\ -\sin\theta/\cos\beta,\ \tan\beta\gamma_5 \tag{25e}$$

$$g_{d\bar{u}H^+} \propto (m_d\tan\beta + m_u\cot\beta) + (m_d\tan\beta - m_u\cot\beta)\gamma_5 \ . \tag{25f}$$

These couplings offer possibilities for enhancements in Higgs production rates above those expected for conventional single Higgs bosons. There could be additional enhancements in $gg \rightarrow H^0$ cross-sections from including virtual $\tilde{q}$ loops. Thus SUSY could have an important impact on Higgs phenomenology at the SSC

Conclusions

- Conventional SUSY should be easy to see because of its large missing p_T signature
- The usual missing p_T signature survives in many (less likely?) alternative SUSY scenarios, but one may have signatures of less missing p_T accompanied by γ, $\ell^+\ell^-$ or $q\bar{q}$.
- If either the $\tilde{q}$ or the $\tilde{g}$ is light: $m \approx 40$ GeV which is certainly not excluded by CERN collider data, then one expects copious sparticle production at a high energy supercollider
- SUSY needs at least two Higgs doublets and tightly constrains their possible masses and couplings, with interesting consequences for Higgs production at a supercollider

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INTERPRETATION OF UNUSUAL EVENTS FROM CERN AND DESY

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Abstract

Several classes of interesting and unusual events from the S$\bar{p}$pS and from PETRA are studied with two purposes in mind. Firstly, varieties of background within the standard SU(3)×SU(2)×U(1) model are described, together with estimates of the number of expected events. Secondly, a review of the recent explanations of the events involving new physics is given. Critical assessments of these proposals focus on the assumptions made, expected rates for the unusual events, and the ability to account for events of several categories.

I. Introduction

The CERN SPS $\bar{p}$p collider data taken up to 1983 have yielded more than 30 unexpected events in addition to those (W,Z,t candidates) anticipated. In e^+e^- interactions at the highest PETRA energies, unusual signatures also may be appearing. In this report we summarize consideration of these events by a working group of the LBL SSC Workshop on Electroweak Symmetry Breaking. For long-term planning, this exercise illustrates the surprises that arise when a set of detectors planned for one kind of physics encounters yet another. In the near term, we hope to aid critical assessment of the new physics interpretations of these events as further data are accumulated. To that end, we are preparing an expanded version of the present article [1].

Six classes of unusual events from the CERN SPS collider and two from PETRA were considered. We discuss the events themselves, and standard physics backgrounds to them, in Section II. Section III deals with proposals for explaining a class of radiative Z decays, while Section IV treats suggestions primarily motivated by events with large missing transverse momentum. Section V considers origins of dimuon events (in $\bar{p}$p collisions) and both 2μ and certain 1μ events in e^+e^- annihilations. Conclusions are drawn in Section VI.

II. Events and Background

We summarize the events to be discussed [2-8] in Table I. These have been reviewed in Ref. [10,11]. We have not included all reported interesting new signatures, such as the 3σ bump in the multijet invariant mass distribution around 150 GeV seen by UA2 [12].

A. Apparent $Z \to \ell^+\ell^-\gamma$ decays

Each of the UA1 samples of 4 $Z \to e^+e^-$ and 5 $Z \to \mu^+\mu^-$ decays contains an event with a hard photon, such that $M(\ell^+\ell^-\gamma) \approx M_Z$. The UA2 sample of 8 $Z \to e^+e^-$ decays also contains one such event (see Table 1).

An important conventional source of these events is QED internal Bremsstrahlung. No other explanation is capable of reproducing the strong observed clustering in the Dalitz plot [Fig. 1]. This clustering expresses the small angle between the photon and one of the leptons; one lepton-photon invariant mass $m_{\ell_1\gamma}$ in Table II is low. The lepton energies in the Z^0 rest frame are:

$$E_1 = \frac{M_Z}{2}\left(1 - \frac{m^2_{\ell_1\gamma}}{M_Z^2}\right) \tag{2.1}$$

External Bremsstrahlung (in which the lepton encounters material after being produced and then radiates) is unlikely since then $m_{\ell\gamma}$ would be extremely low.

The difficulty with an internal Bremsstrahlung explanation of the $\ell\bar{\ell}\gamma$ events is the high observed event rate. The UA2 collaboration has calculated the probability of observing an $e^+e^-\gamma$ event which is less likely than the event observed (and which leads to a signature of three separate electromagnetic energy depositions): $P(e^+e^-\gamma)$=1.0% per e^+e^- event. This would correspond to a probability of 8% for one such $e^+e^-\gamma$ event in the eight e^+e^- events observed. A parallel calculation based on a simulation program gives 13% for this last figure, or 25% if one adds together all configurations including ones in which one electron and the photon are not resolved.

TABLE 1
SALIENT FEATURES OF 8 CATEGORIES OF UNUSUAL EVENTS

EVENT	#	GROUP	REF	FEATURES	COMMENTS
$j\not{p}_T$ $jj\not{p}_T$ $>2j\not{p}_T$	6 0 1	UA1	2,3	Jets of low charged multiplicity and low invariant mass against large missing p_T	If p_T cut is relaxed to 4σ limit: $j\not{p}$ 17, $jj\not{p}_T$ 5, $>2j\not{p}_T$ 3
$ej(s)\not{p}_T$	4	UA2	4	A hard e isoloated from j(s)	In addition UA1 reduced W sample of 43 events contains 2 with $q_T(W)\geq$22GeV
$e^+e^-\gamma$ $\mu^+\mu^-\gamma$	2 1	UA1 UA2 UA1	5,6	EM shower isolated from lepton pair $m(\ell^+\ell^-\gamma) \sim M_Z$	UA1 sees no radiative W decays. UA2 has one $W \to e\nu\gamma$ with e,γ nearly collinear.
$\gamma\not{p}_T$	2	UA1	2,3	Eγ=53,54 GeV	One event may be $W \to e\gamma$ with missed charged track.
$\mu^+\mu^-j(s)$ $\mu^\pm\mu^\pm j(s)$	7 3	UA1	3	6 GeV <m(μμ)< 22 GeV Most events have j. Large abundance of K,Λ.	
Zj(s) $\hookrightarrow \ell^+\ell^-$	5	UA1	3,7	Hadronic activity associated with Z. 4 events consistent with m(Zj(s))∿160 GeV	More j, larger E_T, large n_{ch} than seen in W production, and expected from QCD.
$\mu^+\mu^-jj$	1	CELLO	8	$\sqrt{s}$=4.5 GeV; Little missing energy; All pair invariant masses large.	Mark J has similar events under analysis.
μj(s)	6	MARK J	9	$\sqrt{s}$=46.5 GeV; E_{vis} ~30 GeV; High sphericity.	Hadron distribution too coplanar for $t\bar{t}$ interpretation (2-3σ level).

j(s): hadron jet(s) (occasionally includes charged lepton).
$\not{p}_T$: a substantial imbalance in the observed momentum transverse to beam.
γ : large EM shower with no charged track pointing to it.

The absence of isolated hard photons in W decay [5] implies $P(e\nu\gamma/e\nu) \lesssim 1/50$ (for isolated γ). The UA2 collaboration observes one event consistent with $W \to e\nu\gamma$ in which separate showers for e and γ cannot be resolved. The probability for this event to be external Bremsstrahlung has been calculated to be 4.5% [6].

B. Events with jet(s) or isolated photon and missing p_T

The UA1 collaboration [2] has drawn attention to six events with a jet (A-F in Fig. 2), two with a photon (G,H), and one with 3 or more jets (Δ), opposite large missing p_T. The finite coverage of the UA2 detector prevents a similar statement from UA2, but one candidate for a photon opposite missing p_T has been reported [12]. The UA1 events A-H are summarized in Table III.

The background from QCD jets (with one jet missed) falls quickly with missing energy and is very small for $\Delta E_M \sim 35$ GeV. One monojet event (F) is consistent for expectations for $W \to \nu\tau$ with $\tau \to \nu j$, and will be ignored henceforth. The remaining 5 monojet events A-E have

TABLE II
RADIATIVE Z^0 DECAYS

	$e^+e^-\gamma$		$\mu^+\mu^-\gamma$
	UA1	UA2	UA1
$m(\ell^+\ell^-\gamma)$	98.7±5	90.6±1.9	$88.4^{+46.1}_{-15.2}$
$m(\ell^+\ell^-)$	42.7±2.4	50.4±1.7	$70.9^{+17.2}_{-12.4}$
$m(\ell\gamma)_{low}$	4.6±1.0	9.1±0.3	5.0 ± 0.4
$m(\ell\gamma)_{high}$	88.5±2.5	74.7±1.8	$52.5^{+27.5}_{-9.3}$
$\theta_{\ell\gamma}$	14.4±4.0°	25 ± 1°	7.9°
E_γ	38.8±1.5	24.4±1.0	28.3 ± 3

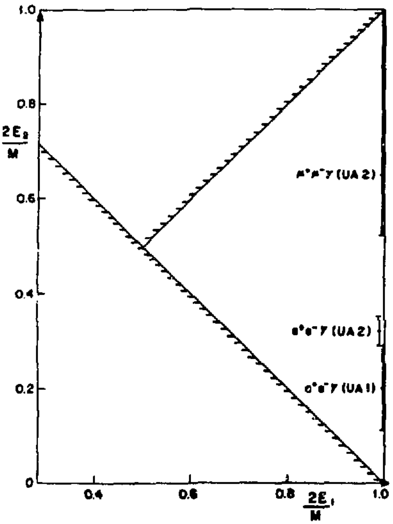

Fig. 1. Dalitz plot for $\ell\bar{\ell}\gamma$ events. E_2 is the energy of the lepton with smaller angular separation from the photon. M is the mean invariant mass of the events.

TABLE III

Properties of events with jet or isolated photon and missing p_T. "Charged tracks" denote those with p_T 0.5>GeV/c.

	Event (Ref. 5)	j or γ E_T (GeV)	ΔE_M (GeV)	M_T(j or γ, ΔE_M) (GeV/c²)	COMMENTS
j	A	25(71^a)	24±4.8 (66±8^a)	130±16	a) Value if hard muon included in jet.
	B	48	59±7	106±12	Three charged tracks m_{eff}=0.79±0.12 GeV/c²
	C	52	46±8	97±17	E_T^{em}=44.4 GeV. One visible charged track.
	D	43	42±6	85±12	Four charged tracks. (two in E^0), m_{eff}= 3.24±0.38 GeV/c².
	E	46	41±7	87±14	Unreconstructed tracks.
	F	39	34±7	73±14	Possible $W \to \tau\nu$
γ	G	46	40±6	84±6	Possible $W \to e\nu$, e track missed.
	H	54	40±4	93±5	

been claimed inconsistent with this interpretation, though it is possible that the τ background was underestimated in Ref. [2]. If the $\tau \to (\geq 4\pi)\nu_\tau$ modes are sufficiently important, a background to events B-E of nearly 2 events of $W \to \tau\nu$ was estimated in Ref. [13]. Without a contribution from $\tau \to (\geq 4\pi)\nu_\tau$, however, the estimate drops to 0.6 event. The background to the events G, H is estimated to be negligible.

The single-shower event G has an azimuth angle $\phi \sim 0$ corresponding to an insensitive area of the central detector: the event may be $W \to e\nu$. The shower in event H, by contrast, occurs in a region where a charged track would be hard to miss.

The monojet events and those involving W + jet(s) (to be discussed below) have an $O(\alpha_s)$ background consisting of hard gluon Bremsstrahlung with $Z(\to \nu\bar{\nu}$ for monojets) or $W(\to \bar{\nu}e$ for $ej\not{p}_T)$. The transverse momentum distribution of W's in Drell-Yan production has been evaluated [13]; it is quite hard. A similar naive estimate gives a QCD monojet background of ~ 1 expected event for $q_T > 25$ GeV (6 observed) and $\sim 1/10$ expected event for $q_T > 50$ GeV (2 observed). These estimates are borne out by more complete calculations [14]. For the monoshower event, the background would be a hard photon Bremsstrahlung together with $Z \to \nu\bar{\nu}$. This is down by a further α/α_s, giving 0.01 expected event with $q_T > 50$ GeV.

We conclude that events A-E and H of Table III cannot be easily dismissed. With the exception of the open

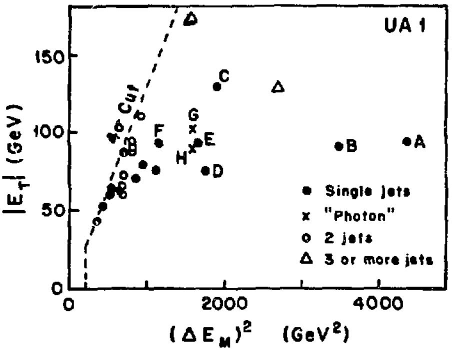

Fig. 2. Events with jet(s) or isolated photon and missing p_T. The dashed line corresponds to $p_T > 4\sigma$, $\sigma = 0.7\sqrt{|E_T|}$. Here $|E_T|$ is the scalar sum of the transverse energy in the detector. From Ref. [5].

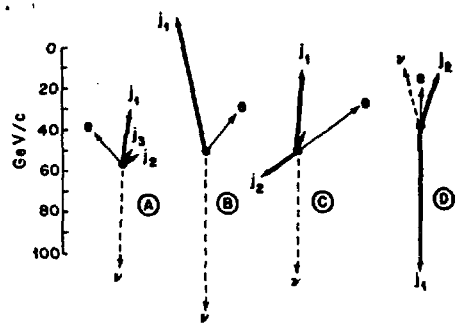

Fig. 3. UA2 events with e + jet(s) + missing p_T, viewed transversely [4].

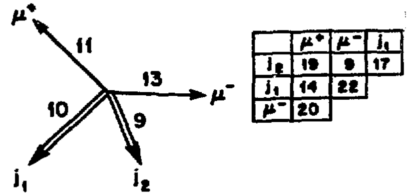

Fig. 4. The dimuon CELLO event [8] at $\sqrt{s}$ = 43.45 GeV. Pair invariant masses and energies are shown in GeV.

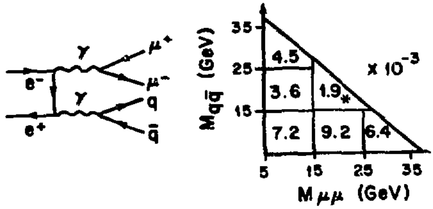

Fig. 5. The $O(\alpha^4)$ contributions to $\mu^+\mu^-jj$ together with the expected number of events plotted in bins of pair invariant masses (from [8]). Observed event is marked by a star.

question of some $W\to\tau\nu$ contamination in events B-E, the backgrounds to these events are all at least an order of magnitude below the observed rate.

C. Apparent W + hard jet(s) events

The UA2 group observes four events with e+jet(s)+(missing transverse momentum), shown in Fig. 3 [4]. One of them [D] could be a heavy quark pair, followed by semileptonic decay of one of the quarks. This is not so for events A-C, for which the missing p_T lies opposite the electron direction. In fact these events are consistent with W+jet(s) followed by $W\to e\nu$.

On the basis of calculated W transverse momentum spectra [13], UA2 would expect ~1 event with q_T>25 GeV (they have 3: A-C) and ~0.1 event with q_T>50 GeV (event B). This suggests that event A could be QCD background, as noted in Ref. [4]. The UA1 collaboration would expect 1.5 of their sample of 43 "clean" $W\to e\nu$ events to have q_T>25 GeV. In fact, they have two events with $22<q_T<24$ GeV [15], but none higher.

D. "Noisy" Z events

The UA1 collaboration has observed that Z production is accompanied by substantial jet activity. Of their sample of 4 $Z\to e^+e^-$ and 5 $Z\to\mu^+\mu^-$ decays, the fractions with (0,1,2,3) jets are (33%,11%,22%,33%). By contrast, their 68 $W\to e\nu$ candidates are accompanied by (0,1,2,3) jets (69%,24%,4.4%,2.9%) of the time [3,7]. The jets occurring with W production are found to agree with QCD expectations, so it is the high activity Z events which appear anomalous. (A signal of equal magnitude in W production could not be ruled out with present statistics, however.)

The calculations made for high-q_T production of Z referred to earlier are relevant here as well. Since one expects $B(Z\to\ell^+\ell^-)/B(Z\to\text{all }\nu\bar{\nu})\gtrsim 1/6$, the backgrounds are expected to be 1/6 of those to monojets.

E. Low-mass dimuons, sometimes with jets

The ability to identify muons has permitted the UA1 group to study a sample of 10 $\mu\mu$+jet(s) events, 7 with $\mu^+\mu^-$ and 3 with $\mu^\pm\mu^\pm$, having $m_{\mu\mu}$ between 6 and 22 GeV/c^2. [Other $\mu\mu$ events are consistent with Z production.] These low-mass $\mu\mu$ events are characterized by high occurrences of strange particles, and vary greatly in their jet activity and invariant masses. Many could be due to processes of the standard model only partially understood, such as gluon fragmentation to $c\bar{c}$, $b\bar{b}$, ... [16]. Heavy (*b*) quark pair production followed by semileptonic decays of both quarks may also play a role [17]. In this connection two of the three same-sign $\mu\mu$ events may be due to $B_s^0-\bar{B}_s^0$ mixing [17]. However [10], neither $b\bar{b}$ nor $c\bar{c}$ production mechanisms fit the kinematics of several of the events.

One $\mu^-\mu^-$ event could be due to $W^-\to\bar{t}b$, $\bar{t}\to\mu^-+\ldots$, $b\to\mu^-+\ldots$. (Some jets in this event must then not have come from the W). Similarly, it is not excluded that one or more $\mu^+\mu^-$ events come from $W\to t\bar{b}$ or hadronic $t\bar{t}$ production [18]. The recent announcement of events compatible with $W\to t\bar{b}$, $t\to b\ell\nu$ [19] should allow more precise calculation of rates in dilepton channels. The UA1 detector will resume running in September, 1984 with enhanced muon detection capability, and the t signal will certainly be searched for in the dilepton channel [20].

F. CELLO $\mu\mu$ + 2 Jet event

Fig. 4 shows a sketch of an interesting event seen in e^+e^- interactions at $\sqrt{s}$ = 43.45 GeV [8]. The event cannot be interpreted as the semileptonic decay of t and $\bar{t}$ quarks because there is insufficient rise in R and no $t\bar{t}$ onium resonance. The data cannot rule out a fourth Q=-1/3 quark, but its semileptonic decay would be expected to yield much larger missing energy and momentum than seen.

This event does have a possible QED explanation. A dominant graph and its expected contribution are shown in Fig. 5. The CELLO collaboration claim a background of ~10^{-3} expected events [8] from such standard radiative processes. This could be increased by as much as an order of magnitude if the background is obtained by integrating over all phase space in which the squared matrix element is smaller than its value near the observed event. In addition, it is tempting to ask what the probability is that the many PEP and PETRA detectors might have observed such an event so near the kinematic boundary. Since the process in Fig. 5 has a low threshold, and since these detectors have accumulated a great deal of integrated luminosity, the event then might appear not nearly as peculiar. Mark J has also seen events with μ pairs and jets, but the background analysis has not yet been completed [21].

G. Mark-J μ + (planar topology) events

The Mark-J group studied events of the form $e^+e^-\to\mu$+ (hadrons). When a cut on events with thrust<0.8 was applied, one expected a background of 1.1 events at the highest energy ($\sqrt{s}\geq 46.5$ GeV) by extrapolating from lower-energy data. Instead, 7 events were seen [22]. The hadronic activity in these events is predominantly confined to a plane which, however, does not contain the muon.

The study of these events is continuing, plagued by the difficulties of running PETRA at such high energies.

III. NEW PHYSICS IN $\ell\bar{\ell}\gamma$ EVENTS?

A. Excited leptons

The decay $Z\to\bar{\ell}\ell^*$, $\ell^*\to\ell\gamma$ is expected in some schemes of composite quarks and leptons. Several authors [23, 24] have ascribed the observed $\ell\bar{\ell}\gamma$ events to this process. The excited lepton is produced and decays via a transition magnetic moment operator of the form

$$\frac{1}{\Lambda}\,\bar{\ell}^*\sigma^{\mu\nu}\ell F_{\mu\nu} \qquad (3.1)$$

A small scale $\Lambda\lesssim 100$ GeV is needed to obtain sufficient rate. This is uncomfortably low in view of limits on other operators [25,26]. To be consistent with g-2 [27-29], further chiral constraints on couplings are necessary. The process $W\to\nu\ell^*$, $\ell^*\to\ell\gamma$ can be suppressed (forbidden) by making m_{ℓ^*} near (above) M_W.

The excited lepton scenario has severe difficulties with the Dalitz plot. Both the high and low ($\ell\gamma$) invariant masses differ by $\sim 3\sigma$ for the $e^+e^-\gamma$ events. ℓ^* is assumed to correspond to the high value as the low-mass ℓ^* would have been seen at PEP or PETRA. It is then very improbable that the γ should be correlated with the prompt lepton to give such low invariant mass values. The operator (3.1) leads to an essentially flat distribution in this mass [24].

B. Scalar boson in Z decay

An alternative explanation of $\ell^+\ell^-\gamma$ events via Z decays involves the chain [30-34]

$$Z \to X\,\gamma, \quad X \to \ell^+\ell^- \qquad (3.2)$$

where X is a scalar or pseudoscalar boson. The observed $\ell^+\ell^-$ invariant masses are barely compatible with one another and with limits ($M_X>47$ GeV) from Bhabha scattering [34,35]. A band at fixed $\ell^+\ell^-$ mass is expected in the Dalitz plot. A persistent feature of such schemes is the prediction of a large $X\to\gamma\gamma$ width, leading to the decay $Z\to 3\gamma$.

C. Scalar state $\to \ell\bar{\ell}\gamma$

It is possible that scalars expected in composite or technicolor schemes are heavier than the Z. In this case they could still yield the observed $\ell\bar{\ell}\gamma$ events provided their mass is less than ~ 100 GeV [36]. In this scheme the scalar (or pseudoscalar) boson R is taken to have large couplings to fermions only if these couplings are chirally invariant. A class of dimension 7,9,... operators then arises which leads to matrix elements for $R\to\ell^+\ell^-\gamma$ which vanish in the soft photon limit and peak along the edges of the Dalitz plot. The peaking is not sharp enough to predict the observed distribution, but it is a step in the right direction. Operators in which ℓ is replaced by ν also occur, leading to mono-shower events.

The scale factor needed to obtain suitable production rates is, as usual, uncomfortably [26] low: $\Lambda\sim 100$ GeV.

In both X and R boson mechanisms, $q\bar{q}\gamma$ events are expected with $M(q\bar{q}\gamma)\approx M_Z$. Present data cannot exclude such events [3].

D. Z mixing with exotic quarkonium

The Z would appear to have anomalous decay modes if it was degenerate and mixed with some other state. One such model [37] envisages the Z mixing with an excited 1^{--} onium state of a quark with exotic color. This is assumed to decay to the lowest 1^{--} state of mass ~ 50 GeV via emission of a hard photon to a 0^{++} state, followed by a soft (unobserved) photon. The lowest 1^{--} state will occasionally decay to $\ell^+\ell^-$ giving the observed signature. A sufficient rate requires essentially complete mixing of the states, with a 2% branching ratio for the $\ell^+\ell^-\gamma$ decay chain.

This scheme has much in common with the X boson idea mentioned above. Moreover, it requires the binding of quarks with higher color representations to produce an extraordinary spectrum of states, with the lowest P state nearly degenerate with 1S and the Z nearly degenerate with 2S.

E. Composite W,Z

If the W and Z are composite [27,33,38-40], operators of the form $(G/\Lambda^2)\,F^{\mu\nu}Z_\mu\Box Z_\nu$ could occur. The branching ratio for radiative decays is then [38]

$$\frac{\Gamma(Z\to\ell^+\ell^-\gamma)}{\Gamma(Z\to\ell^+\ell^-)} \sim 10^{-4}\,G^2\,\frac{M_Z^4}{\Lambda^4} \qquad (3.3)$$

This is sufficient only if $G\gg 1$ for $\Lambda\lesssim M_Z$. One would then expect to see $Z\to q\bar{q}\gamma$ [41] or $Z\to q\bar{q}g$ (i.e., jjj) [39]. If Λ is so low [41], however, one would expect a momentum dependence of vector boson masses, deviation of ρ from unity, and W radiative decays at a large rate.

Perhaps the worst feature of this scheme is that it prefers large invariant masses for both ($\gamma\ell$) pairs, rather than one large and one small. This has been emphasized in Ref. 42 by comparing the Dalitz plot distributions of the data with that of Bremsstrahlung, the X boson, excited leptons, and a composite Z. Bremsstrahlung does the best, and a composite Z the worst.

A model with an effective $Z\gamma\gamma$ vertex [43] also has been proposed. It has the same difficulties with the Dalitz plot distribution.

F. WW bound states

It has been proposed [41] that for very heavy Higgs mass the resulting forces between longitudinal W bosons are strong enough to bind them into a state of mass ~ 90 GeV. This state would then decay to a virtual $Z(\to e^+e^-)+\gamma$, in the manner of the R boson discussed earlier. The resulting eight-fermion operators coming from strong W-W interactions are conjectured to be responsible for same-sign multimuon events in neutrino scattering. One would also expect the 90 GeV state to decay to virtual $Z(\to\nu\bar{\nu})+\gamma$, giving the monoshower event(s), and to be produced in e^+e^- interactions via virtual Z exchange in association with a photon [44].

We believe the production rate for such a bound state is far too small to be relevant for the unusual events. For comparison, a 100 GeV Higgs boson is produced via W fusion with a quark subprocess cross section of less than 10^{-4} nb [45]. Quark luminosity factors reduce the $\bar{p}p$ cross section still more.

IV. NEW PHYSICS IN EVENTS WITH MISSING p_T?

A. Remarks on the 160 GeV mass region

Many of the unusual events we discuss seem to point to a common origin in the mass range of 160 GeV. (See in particular Ref. 7.) These include monojets (if interpreted as $j+Z$, $Z\to\nu\bar{\nu}$); monoshowers (if $\gamma+Z$, $Z\to\nu\bar{\nu}$), "noisy" Z events, and W + jet(s) events. This mass range will be more efficiently studied by raising the SPS energy ($\sqrt{s}$ = 630 GeV in the forthcoming run), and in particular at the Tevatron ($\sqrt{s}$ > 1.6 TeV). Meanwhile a cautionary note is that the selection of events containing W or Z (or their analysis as such), combined with cuts on a steeply falling p_T(jet) spectrum, can conspire to produce a peak.

B. Higgs bosons

It has been proposed that many unusual SPS events ($ej(s)\not{p}_T$, $j\not{p}_T$, $\gamma\not{p}_T$) come from decay of a 160 GeV Higgs particle [46]. The cross section must be enhanced by $\sim 10^6$ with respect to naive estimates in order to obtain a sufficient event rate ($\sim 10^3$ produced at CERN). This enhancement makes the Higgs boson so broad ($\Gamma \gtrsim 200$ GeV) that a peak is unlikely, and production violates unitarity [47,48].

C. New gauge interaction

As an example, we consider the case of "odor" [49], a proposed interaction with $\Lambda_O \sim \Lambda_{QCD}$ and with the lightest odor quark having a mass of ~ 75 GeV. A spectrum of $O\bar{O}$ ("odoronium") bound states between 150 and 300 GeV is then expected. $G\bar{O}$ production leads almost exclusively to odoronium. It is necessary in this scheme for the $O\bar{O}$ cross section to be ~ 1 nb. (A perturbative estimate falls short by about a factor of 100, for color triplet O quarks.) The observed events are then ascribed to specific products of $O\bar{O}$ annihilation, e.g.

$$O\bar{O}(1^-) \to Z\,H \;\; (Z \to \nu\bar{\nu}) \quad : \quad j\not{p}_T \tag{4.1}$$

$$O\bar{O}(0^-) \quad Z\,\gamma \;\; (Z \to \nu\bar{\nu}) \quad : \quad \gamma\not{p}_T \tag{4.2}$$

However, many other decay modes are expected, and it is not clear they all occur with consistent branching ratios. [See Table IV] Notable is the prediction [50] that approximately $10 e^+e^-$ and $\mu^+\mu^-$ events would be expected from the 1^- decay. The jet-jet bump seen by UA2 [12] around 150 GeV should also appear in the 3j spectrum. It is not clear whether odor gluon (G) emission is visible; odor gluons should form odor glueballs (GG or GGG) which are invisible except via energy and momentum balance.

D. Color octet mesons

In the previous two sections we saw that attempts to produce a 160 GeV state have generally led to insufficient rates, especially for a Higgs boson. It has been suggested that these problems can be overcome by producing a mesonic state, predominantly via $q\bar{q}$ fusion, which is a color octet [51]. This idea takes advantage of the large $q\bar{q}$ differential luminosity (6 times larger for $u\bar{u}$ than for gg at $\sqrt{s} = 540$ GeV and M=160 GeV), and allows for decay channels involving a weak boson (such as gW, gZ, gγ) at a rate down by only one power of α/α_s compared to the leading decay channel. Furthermore, for a given partial width to $q\bar{q}$, the production cross section for a color octet meson M_8 is 8 times larger than for a color singlet.

In Table V we show the number of events expected at the CERN collider if $\sigma(M_8)=\sigma(W)$. Drell-Yan production rates for the W and Z are shown for comparison. From the last three lines one finds 4 monojets, 2-3 $ej\not{p}_T$ events, and a 15% rate for jetty Z production. However, there are also an order of magnitude too many events in the jj bump, and 100 dramatic γj events. These latter should be searched for.

TABLE IV

1^- DECAY TO	EXPECTED EVENTS	POSSIBLE SIGNATURE
ggg	45	jjj bump at m=150 GeV
GGG	45	odor glueballs: invisible
G^2g^3	1/2	j(s) $\not{p}_T$
G^3g^2	1/2	
ZH, $Z\to\nu\bar{\nu}$	5	j $\not{p}_T$
ZGG, $Z\to\ell^+\ell^-$	1	$\ell^+\ell^-$ $\not{p}_T$
Zgg, $Z\to\ell^+\ell^-$	1	$\ell^+\ell^-$ j(s)
ZGG, $Z\to q\bar{q}$	10	jj$\not{p}_T$
γH, $H\to b\bar{b}$	5(unless $m_H<2m_b$)	γj
γgg	5	γ j
e^+e^-, $\mu^+\mu^-$	> 5 each	high inv. mass $\ell^+\ell^-$
0^- DECAY TO	**EXPECTED EVENTS**	**POSSIBLE SIGNATURE**
gg	15	jj bump at m=150 MeV
GG	15	odor glueballs: invisible
g^2G^2	1/7	jj $\not{p}_T$
Zγ, $Z\to\ell^+\ell^-$	1/20	$\ell^+\ell^-\gamma$ with $m(\ell^+\ell^-\gamma)$=150 GeV
HGG $H\to b\bar{b}$	1	j $\not{p}_T$
Zγ, $Z\to\nu\bar{\nu}$	1/3	γ$\not{p}_T$

$\ell^+\ell^-$ includes e^+e^- and $\mu^+\mu^-$ contributions.

TABLE V. Decays of color octet mesons (based on Table of [51])

DECAY MODE	# EVENTS	DECAY MODE	# EVENTS
$M_8\to q\bar{q}$	500	$W\to\bar{q}q$	500
$M_8\to gW$	37	$Z\to q\bar{q}$	100
$M_8\to gZ$	23		
$M_8\to g\gamma$	100		
$M_8\to gZ$, $Z\to\nu\bar{\nu}$	4		
$M_8\to gW$, $W\to\nu e$	2.2	$W\to\nu e$	35
$M_8\to gZ$, $Z\to e^+e^-$	.7	$Z\to e^+e^-$	4

TABLE VI. Decay modes of an excited quark [from Ref. 56]

MODE	SIGNATURE	$\Lambda=\Lambda'=150$ GeV # EVENTS	$\Lambda=50,\Lambda'=15$ GeV # EVENTS
$q^*\to qg$	jj bump at 150 GeV	10	80
$q^*\to\gamma q$	jγ bump at 150 GeV	.2	20
$q^*\to qW$, $W\to e\bar{\nu}$	$j\not{p}_T e$	.04	4
$q^*\to qZ$, $Z\to\nu\bar{\nu}$	$j\not{p}_T$	.04	4

E. Neutral leptons

The previous three explanations have assumed $\not{p}_T$ to be $Z\to\nu\bar{\nu}$, where ν is a conventional neutrino. It is also possible that $\not{p}_T$ could be carried off by a heavy neutral lepton ν_H of a few GeV mass, such as a fourth sequential neutrino ν_4 [52,53] or a mirror neutrino ν_M [52-54], produced in the decay $Z\to\nu_H\bar{\nu}_H$.

A sequential neutrino ν_4 typically has neutral-current decays suppressed via the GIM mechanism. To give missing p_T, it must decay outside the detector. Its mass must be chosen very carefully for this to be reasonable. It could give monoshower events by occasionally decaying to $\gamma+\nu_i$ ($i=e,\mu,\tau$). In that case one would also expect $\gamma\gamma\not{p}_T$ events, with $M(\gamma\gamma\not{p}_T) \le M_Z$. In many cases the neutrino ν_4 must decay inside the detector, giving rise to monojets with charged tracks originating some distance away from the interaction point.

For a mirror neutrino ν_M, the GIM mechanism is frustrated, and one expects [54,55] $B(\nu_M\to\nu\nu\bar{\nu})$=0.1, $B(\nu_M\to\nu$+hadrons)=0.2, $B(\nu_M\to\ell\ell'\nu)$=0.2, $B(\nu_M\to\ell$+hadrons)=0.5, where ℓ is a charged lepton. The missing p_T signature is then expected to come from $\nu_M\to\nu\nu\bar{\nu}$. The monojets come from such decays as $\nu_M\to\ell$+hadrons. The low charge multiplicity in the observed monojets and their low effective mass (when all tracks are reconstructed) argues for $M(\nu_M) \lesssim$ (few GeV). The mono-shower events correspond in this scheme to a less likely decay such as $\nu_M\to\nu$+(all neutral hadrons).

In future runs the sequential [52] and mirror [54] schemes may be differentiated. The sequential neutrino is expected to have a sizeable radiative decay, so that γj, γℓ events should be seen. The charged tracks should in general originate a detectable distance from the beam pipe, reflecting the finite lifetime needed to account for events with missing transverse momentum. The scheme based on neutrinos with a $\nu\nu\bar{\nu}$ decay mode can account for monojets without finite lifetime effects, but predicts jj as well as $j\not{p}_T$ events in which j has a high lepton content. Both schemes have difficulty in accounting for the most spectacular $j\not{p}_T$ event ("A" of UA1) unless another Z is postulated [54].

F. Excited quarks

A model [56] which could account for anomalous events at the CERN collider postulates an excited quark q^* belonging to a color 3, flavor doublet, with charges 2/3 and -1/3, and $M(q^*) \sim 150$ GeV. The q-q^*-gluon coupling is assumed to be of the anomalous moment type with scale Λ, and q-q^*-(W or B) couplings are also assumed to exist (B is the boson of electroweak $U(1)_Y$) with scale Λ'. A natural choice for these scale factors would be a compositeness scale, which might also be $M(q^*)$. q^* is produced by quark-gluon fusion and decays via $q^*\to qg, q\gamma, qZ, qW$. The expected number of events for $\int\mathcal{L}dt=137nb^{-1}$ are shown in Table VI [56].

Although there are signatures for $j\not{p}_T$, $j\not{p}_Te$, and jj events expected, there are rate problems if $\Lambda=\Lambda'=150$ GeV. It would seem necessary to lower $\Lambda=\Lambda'$ to 15 GeV to obtain a sufficient rate for $j\not{p}_T$ and $j\not{p}_Te$, but then the jj rate would be too high. Moreover, other high dimension operators with the same scale (e.g., $\ell\ell q\bar{q}$) are excluded by present data [26]. The result of taking $\Lambda\neq\Lambda'$ is also shown in Table VI. The number of $j\gamma$ events is still large (shown for $Q(q^*)=2/3$). It would be reduced to 5 events if $Q(q^*)=-1/3$. Again (as for color octet mesons) the γj signature appears worth looking for.

G. Supersymmetry [57,58]

A missing transverse energy signature has been recognized as a signal for the production of supersymmetric partners of the known particles, both at e^+e^- colliders [59] and at $\bar{p}p$ colliders [60,61].

The most favored supersymmetric phenomenologies have an unbroken R parity ensuring the stability of the lightest superpartner, taken to be the photino $\tilde{\gamma}$. This neutral particle, which interacts with strengths similar to that of the neutrino, is expected to carry off missing transverse energy. However, the observed events with large $\not{p}_T$ do not involve the predicted broad jets [60] opposite this momentum which would result from $W\to\tilde{w}\tilde{\gamma}$, $\tilde{w}\to q\bar{q}\gamma$. Instead, the jets appear narrow. Any supersymmetric scenario must cope with this feature. We are aware of five variants, involving production of $\tilde{g}\tilde{g}$, $\tilde{q}\tilde{q}^\dagger$, $\tilde{q}\tilde{g}$, $\tilde{w}\tilde{g}$ and even $\tilde{\gamma}\tilde{\gamma}$, upon which we now comment.

1,2. Gluino, squark pairs. The CERN collider can produce light gluino pairs copiously [62], and isolated $\not{p}_T$ signatures have been recognized as useful for gluino searches [63]. Gluino [64,65] and squark [66] pair production in fact yields monojet events under the UA1 event selection criteria, as a result of loss of soft jets or coalescence of two jets. The $\not{p}_T$ spectrum resulting from $\tilde{q}\tilde{q}^\dagger$ production ($\tilde{q}\to q\tilde{\gamma}$, $\tilde{q}^\dagger\to q^\dagger\tilde{\gamma}$) is harder than that from $\tilde{g}\tilde{g}$ production ($\tilde{g}\to q\bar{q}\tilde{\gamma}$) [65], and the jets are narrower, so $\tilde{q}\tilde{q}^\dagger$ is preferred. However, since most models do not give squarks much lighter than the gluino [66,67], it is likely that both mechanisms contribute.

The predicted $j\not{p}_T$ and $jj\not{p}_T$ rates for various squark masses, assuming the UA1 selection criteria, are shown in Fig. 6. Squark masses much below 35 GeV are ruled

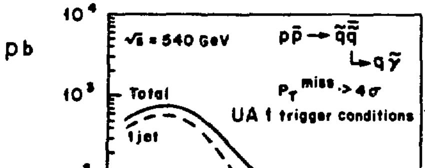

Fig. 6. The total and topological cross-section for $\tilde{q}\bar{\tilde{q}}$ production followed by $\tilde{q}\to q\tilde{\gamma}$ decay giving one- or two-jet final states with $p_T^{miss}>4$, and fulfilling the UA1 trigger conditions.

out by the observed monojet rate, and similar limits apply to the gluino mass. The spectacular monojet event "A" of UA1, and possibly the event "B", are left unexplained for a squark or gluino mass of 40 GeV, which otherwise fits the observed p_T distributions.

3. Singly produced squarks. The mechanism $\tilde{g}q\to\tilde{q}$ could lead to single squark production. One would require a gluino of about 25 GeV (a lighter $\tilde{g}$ seems ruled out by $jj\not{p}_T$ data [67]). The signature would be $\tilde{q}\to q\tilde{\gamma}$ [68]. This mechanism could account for the observed number of monojets, even if $m_{\tilde{q}}\sim150$ GeV [69]. The decay $\tilde{q}\to q\tilde{g}$ can provide a jj bump at 150 GeV, as seen by UA2 [12], and $\tilde{q}\to\tilde{w}q$, $\tilde{w}\to W\tilde{\gamma}$, $W\to e\nu$ yields $ej\not{p}_T$ and a possible explanation of the UA2 event "B" of Fig. 3. A related mechanism, $\tilde{g}q\to\tilde{\gamma}q$, has also been proposed as a source of monojets [70]. However, all these mechanisms require far more $\tilde{g}$ in the proton than one might expect for heavy quarks (such as t [71]).

4. Heavy squark, light gluino [72]. If the gluino were only slightly more massive than the photino it could escape the detector before decaying to $q\bar{q}\tilde{\gamma}$, thus frustrating the 25 GeV lower bound on its mass. In this case squarks could be produced singly even if their mass were as large as 100 GeV: $gq\to\tilde{g}\tilde{q}$. Presumably the monojet signature would come from $\tilde{q}\to q\tilde{g}$, where the energetic gluino is missed.

5. Photino pairs. One model with broken R-parity envisions production of a pair of 5-8 GeV $\tilde{\gamma}$'s, followed by $\tilde{\gamma}\to\tau\bar{\tau}\nu_\tau$ [73]. The monojets occur when one photino produces a fast ν_τ and a $\tau^+\tau^-$ pair with $p_T<10$ GeV, while the other throws the $\tau^+\tau^-$ forward. The monoshower event is viewed as a fluctuation to zero observed charged multiplicity of the $\tau\bar{\tau}$ decay products.

6. $\tilde{w}\tilde{g}$ production. The processes $q\bar{q}\to\tilde{w}\tilde{g}$ and $qg\to\tilde{w}\tilde{q}$ [74] can give $ej(s)\not{p}_T$ signatures. Rates for these processes are generically down by at least an order of magnitude compared to $\tilde{g}\tilde{g}$, $\tilde{q}\tilde{q}^\dagger$, and $\tilde{g}\tilde{q}$ production [58]. Optimal values for $\tilde{g}$, $\tilde{q}$, and $\tilde{w}$ masses could enhance the signal [75].

H. Heavy quark

We note that one UA2 event of $e2j\not{p}_T$ ("C" of Fig. 4) is barely compatible with heavy quark pair production, $p\bar{p}\to Q\bar{Q}+\ldots$, $Q\to W+q$, $W\to e\nu$; $\bar{Q}\to W+\bar{q}$, $W\to q\bar{q}$ [76]. Here q has to be light ($m_q\lesssim$few GeV), which may not be favored if Q is the lightest member of a fourth generation [77].

V. MECHANISMS FOR $\mu\mu$ + jet(s) AND μ + (PLANAR EVENTS)

A. Leptoquarks

It has been suggested [78] that the CELLO event [8] is due to the reaction $e^+e^-\to L\bar{L}$, where L is a leptoquark which decays to μ + jet. This possibility can explain the apparent back-to-back nature of each μ + jet in the event. However, it entails large cross section for leptoquark pair production at the CERN collider. The observed 2μ + jet(s) signal mentioned in §II.E can be used either to bound leptoquark pair production, or to provide confirmation of the hypothesis.

B. Neutral heavy leptons

One possibility suggested for the CELLO event is the production of a pair of neutral leptons $\nu_H\bar{\nu}_H$, either via a virtual Z^0 [8,79] or via a new, weakly coupled "Z_x" in the 50-70 GeV range [79]. In the latter case, ν_H could be a right-handed neutrino "N" of the type described in Refs. [54,55]. The decays $Z^0\to\nu_H\bar{\nu}_H$ or $Z_x\to N\bar{N}$ should then be observable at the CERN collider. A Z_x (50-70 GeV) should also have an observable e^+e^- decay mode, and should affect electroweak asymmetries at PETRA. Both signatures will be visible in forthcoming improvements of present data.

A neutral heavy lepton, produced in pairs, also could be responsible for the Mark-J events discussed in §II.G. One neutral lepton would decay to μ+(hadrons) and the other to (say) $\ell\ell'\nu$ or ν+(hadrons).

C. Heavy quarks

The Mark-J events have some properties in common with semileptonic decays of heavy quarks. Possibly related features are that (i) the largest fluctuation in R occurs at $\sqrt{s}$=44 GeV, and cannot exclude the 1S bound state of a Q=-1/3 quark and its antiquark; (ii) R is sufficiently poorly measured above 45 GeV that one cannot exclude the threshold for a Q=-1/3 quark. Further study of the $\sqrt{s}$=44 GeV region is in progress, and will probably be able to settle the question of whether a new quark is responsible for the events in the near future.

VI. CONCLUSIONS

A summary of our findings is presented in Table VII and VIII. A glance at Table VII is quite rewarding. Theorists have found the kinematics of the $\ell^+\ell^-\gamma$ events essentially impossible to explain except as a statistical fluctuation of bremsstrahlung, which thus remains the most likely source.

It is easier to invent explanations for events with large $\not{p}_T$, but few ideas apply to several event categories at once, and few explain the event topologies and rates in a natural way.

Many explanations are based on new physics in the 160 GeV mass range, with characteristic γj and jj peaks expected at this mass. All would benefit from improved statistics, which are eagerly awaited.

ACKNOWLEDGEMENTS

We would like to thank T. Appelquist, M. K. Gaillard, and I. Hinchliffe for organizing an enjoyable workshop at LBL where this work was begun. We thank all participants of the CERN Events working group, which includes B. Barbieri, I. Bars, M. Bowick, J. Ellis, K. Ellis, E. Farhi, B. Gavela, H. Georgi, J. Gunion, H. Haber, R. Holdom, P. Langacker, M. Machacek, P. Roy, S. Sharpe and M. Suzuki. In addition we thank J. Branson, R. Cashmore, G. Kane, Z. Kunszt, R. Peccei, R. Schwitters and members of the UA1 and UA2 collaborations for comments and suggestions.

TABLE VII

NEW PHYSICS EXPLANATION	REF	III $\ell^+\ell^-\gamma$	IV $\gamma\not{p}_T$
$Z\to\ell\ell^*$, $\ell^*\to\ell\gamma$	23,24	☼ ■	-
$Z\to X\gamma$, $X\to\ell^+\ell^-$	30-34	☼ ■	✓
$R\to\ell^+\ell^-\gamma$	36	☼✓	✓
$Z\to Q\bar{Q}$	37	☼ ■	✓
COMPOSITE W AND Z	29,33 38-40	☼ ■	✓
WW BOUND STATES	41	☼ ■	✓

Comments: (1) all schemes fail to account for the Dalitz plot distribution; however $R\to\ell^+\ell^-\gamma$ does the best in this regard.
(2) None of these schemes has an obvious mechanism for explaining any of the other events in Table I.

Symbols:
☼ primary motivation
✓ qualitative explanation
■ fails to give quantitative explanation
✸ Quantitative explanation
\- No obvious explanation

L.J.H. and J.L.R. thank the CERN Theory Division for their hospitality while this work was being completed.

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+Supported in part by the National Science Foundation under Grant No. PHY-82-15249.
++Supported in part by the Department of Energy under contract No. DE-AC02-76ER03069.
*Supported in part by the Department of Energy under contract No. DE-AC02-82ER40073.

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TABLE VIII

NEW PHYSICS EXPLANATION	REF.	I $j\not{p}_T$	II $ej(s)\not{p}_T$	III $\ell^+\ell^-\gamma$	IV $\gamma\not{p}_T$	V $\mu\mu j(s)$	VI $Zj(s)$	VII CELLO $\mu^+\mu^-jj$	VIII MARK I $\mu j(s)$	COMMENTS
1.0 GeV Higgs	46	☼ ■	☼✓	-	☼✓	-	✓	-	-	No clear $j\not{p}_T$ signal. Rate assumed, and conflicts with unitarity.
ODUR	49	☼✓	-	-	✓	-	☼✓	-	-	Can explain UA2 jj bump at 150 GeV. Predicts too many e^+e^-, $\mu^+\mu^-$ events at 150 GeV.
COLORED OCTET MESONS	51	☼✓	✓	-	✓	-	✓	-	-	Can explain UA2 jj bump at 150 GeV. Rate good. Predicts $j\gamma$ events at 150 GeV.
ν_4	52	• ■	-	•	☼✓	-	-	-	-	•These signatures exist, but would not look like observed events.
OTHER NEUTRAL HEAVY LEPTONS	53 54	☼✓ •	-	-	✓*	•	-	✓	-	*could be j with n_{ch}=0 •Monojet A and .. require Z' boson
EXCITED QUARKS	56	☼✓	✓	-	-	-	-	-	-	Predicts jj bump at M(q*).
SUPER-SYMMETRY	68-70	✸ •	✓	-	✓*	-	-	-	-	*could be j with n_{ch}=0. •Monojet A unlikely.

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LOW ENERGY SIGNALS OF COMPOSITE MODELS*

Compositeness Study Group at the SSC Theoretical Workshop

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SUMMARY

Some signals of compositeness that represent deviations from the standard model at low energies are discussed. Emphasis is given to exotic composites, strong P,C violation beyond the weak interactions and small deviations in relations among the parameters of the standard model. Such effects may be detected at energies obtainable at CERN, LEP and the SSC.

EXOTIC COMPOSITES

If quarks and leptons are composites of preons with a scale Λ_p, their low mass ($m \ll \Lambda_p$) can be understood only if there is a preonic chiral flavor symmetry group G_F that is not spontaneously broken in the vacuum. G_F must include $SU(3)_c \times SU(2)_w \times U(1)_Y \times$ Baryon Number $\times$ Lepton number $\times$ family quantum numbers. If all forces, except the strong precolor forces, are neglected at the scale Λ_p, G_F may appear to be a much larger symmetry. Since the formation of the bound states at Λ_p has nothing to do with the other smaller forces, (e.g.QCD), it is a good approximation to use the enlarged (globally conserved) G_F to classify all massless or massive composites in irreducible representations of G_F.

The massless composites, classified as {R} under G_F, must satisfy certain consistency conditions[1] if the symmetry G_F is to remain unbroken. General unique solutions have been given to a set of very restrictive conditions[2]. Thus, it is now known that there exists a classification of possible precolor groups and precolor representations (preons) that yield a predetermined set of massless composite states {R}. Potentially realistic examples can be found among these models.

The next stage is to analyze the $SU(3) \times SU(2) \times U(1)$ content of these representations to determine the quark and lepton structure. There may be more than one way of embedding $SU(3) \times SU(2) \times U(1)$ in G_F; each embedding may give a different structure. One may find along with the massless quarks and leptons that there are massless exotic composites in the sense that they carry high color or high weak-isospin or high hypercharge (or electric charge):

$$
\begin{aligned}
\{R\} \rightarrow\ & \text{leptons} = \{(1, I)_y\} \\
& + \text{quarks} = \{(3, I)_y\} \\
& + \text{High color} + \{(6 \text{ or } 8 \text{ or } 10 \text{ etc.}, I)_y\}
\end{aligned}
$$

where the isospin I and hypercharge Y may take non-exotic or exotic values. The emergence of exotics is not necessary in every model (see e.g. ref. 2) but they may occur quite naturally in many models.

When the symmetry is broken from G_F to $SU(3) \times SU(2) \times U(1) \times$ Baryon no. $\times$ Lepton no., the exotics must become massive while the ordinary families are still massless (exact SU(2)). Since by definition of G_F, this breaking is assummed to occur at a scale μ considerably smaller than Λ_p, $\mu \ll \Lambda_p$, the exotics are much lighter than the heavy composites. The study of exotics is therefore interesting since they may provide some clues[2] of compositeness at energies much smaller than Λ_p. Furthermore, it is conceivable that high color exotics may play a role in electroweak symmetry breaking[3] at 250 GeV and generate masses for the quarks and leptons in a composite model[4]. The abnormally energetic events seen at UA1 and UA2 at CERN may be associated with exotics.

These remarks provide some motivation for studying exotics at this workshop. We report here two possible occurances of exotics 1) Exotics produced freely, 2) Exotics within non-exotic bound states. We concentrate on high color exotics since their production cross section is large in pp reactions already at CERN energies.

In the reaction

$$pp \text{ or } p\bar{p} \rightarrow Q\bar{Q} + \text{anything}$$

The pair of high color exotics ($Q\bar{Q}$) may be produced freely or in a bound state, depending on the strength of the QCD color force that acts on them. The bound state may be in the form of -onium (like charmonium) or in the form of a pseudo-goldstone (η'-like object) if high color plays a role in electroweak symmetry breakdown at 250 GeV. The observation in the form of a bound state is possible only if the lifetime of the composite state is shorter than the lifetime of the individual exotic fermions.

The dominant parton subprocess in the production of high color exotics is gluon fusion $g + g \rightarrow Q + \bar{Q}$. the production cross section is enhanced by color factors relative to the production of heavy triplet quarks, bound or free. The production of a bound state $X = Q\bar{Q}$ is given by

$$\sigma(p\bar{p} \xrightarrow{+} X + \text{any}) = \frac{\pi^2 \Gamma_2}{8M_X^3} \left(\frac{\tau dL}{d\tau}\right)\bigg|_{\tau = M_X/s} \left[q_r^2 \frac{3}{d_r}\right]$$

where Γ_2 is the decay rate of $X \rightarrow gg$ if Q is a color triplet, $\tau\, dL/d\tau$ is the parton effective luminosity factor, as given in ref. 5., and the last bracket is a color factor for the color representation r: d_r is the dimension of the representation and q_r is related to the quadratic casimir operator normalized to 1 for a triplet

d_r	3	6	8	10	. . .
q_r	1	5	6	15	. . .

For $M_X = 150$ GeV, we estimate $\Gamma(X \rightarrow gg) = 1-10$ MeV for either -onium type bound state or goldstone type bound state. At CERN energies, $\sqrt{s} \sim 540$ GeV, the luminosity factor gives a cross section

$$\sigma \sim (10^{-7} - 10^{-6})\text{nb}$$

which is too small to produce any appreciable number of events. Thus, surprisingly such bound states may hide quite well at CERN. However, at the SSC, for $M_X \sim 1$ TeV, Γ_X is estimated to be ≈ 300 MeV, and at $\sqrt{s} = 40$ TeV the much larger luminosity factor produces

$$\sigma \sim 10^{-2} \text{ nb}$$

The analysis of the signals may be done as described in ref. 6, for a similar bound state of two gluinos. Such a state may produce detectable signals at the SSC.

The free production and subsequent decay of certain exotics can produce a much larger number of events so that their detection is enhanced by many orders of magnitude. See, for example, estimates of gluino production and detection at CERN[7] or the SSC energies[5]. Composite exotics with certain properties that can produce energetic jet signals plus large missing energy with cross section of the size alleged to be detected at CERN can exist in composite models. A minimal example of a model containing a zero-charge color nonet $L_8 + L_1$ is given in ref. (4). In the low energy theory one finds effective couplings of the form

$$L_{eff} = \bar{L}_8 i \not{D} L_8 + \bar{L}_1 i \not{\partial} L_1 + \frac{g}{\Lambda_p} \bar{L}_8 F\mu\nu \sigma^{\mu\nu} L_1$$

where $D L_8 = \partial L_8 - ig [A, L_8]$ couples the gluon to a pair of octets and g/Λ_p describes the magnetic coupling for the gluon + octet + singlet. Since the octet and singlet carry the same global quantum numbers (they come from a nonet) the magnetic coupling is not suppressed by factors of $(mass/\Lambda_p)$. thus, the octet decays dominantly to the singlet plus a gluon with a decay rate

$$\Gamma_8 = \frac{1}{\pi} \alpha_{QCD} M_8^3/\Lambda_p^2$$

This may be small, but since it is nearly 100% branching ratio, it is only the magnitude of the production rate that determines the number of events in the final state. The production cross section for the octet is large at the CERN (and higher) energies, as estimated[7] for production of 40 GeV gluinos in supersymmetric theories, and can produce many events with jets + missing energy of the type and numbers alleged to be seen at CERN. Thus

$$\text{Gluon} + \text{Gluon} \to \text{octet} + \text{octet}$$
$$\text{octet} \to \text{jet} + \text{missing}$$
$$\text{octet} \to \text{jet} + \text{missing}$$

produces signals of the type

proton + (anti-) proton→jet+jet+missing+anything

with the characteristics of the CERN events at UA1 and UA2. If one of the octets has a slow forward momentum, the momentum cuts will make it appear as if there was only 1 energetic jet in the final state. thus the mono-jet and two-jet events with approximately correct size cross sections may be explainable by composite exotics.

There may exist other exotics with interesting properties. For example, a sextet plus a triplet with the same global U(1) quantum numbers will have effective interactions of the type described above. The decay rate for sextet → triplet + gluon is similar in magnitude to the octet described above. [If the triplet does not share the same global quantum number as the sextet, then Γ_6 is suppressed by a factor $(m/\Lambda_p)^2$]. The production and decay of a pair of heavy sextets will now produce 2 + 2 = 4 energetic jets in the final state without any large missing energy. This kind of signal may help or destroy certain models.

An excited triplet quark, which also has unsuppressed magnetic couplings with a gluon and a quark, would differ from the properties of the sextet described above. For one thing, the mass of Q^* is likely to be of order Λ_p. But, if somehow its mass were low, it will be expected to have effective magnetic couplings $\bar{Q}^* \sigma_{\mu\nu} q F\mu\nu$ with $F_{\mu\nu}$ = photon, W, Z in addition to gluon. Thus, from the ratio of the coupling constants for

$$Q^* \to q+g : Q^* \to q+\gamma : Q^* \to q+W : Q^* \to q+Z,$$

taken proportional to the strenths of the gauge coupling constants, one may estimate the ratio of number of events for

$$\text{jet+jet} : \text{jet+}\gamma : \text{jet+}e\nu : \text{jet+}\nu\bar{\nu} : \text{jet+}e\bar{e},$$

with a result

$$10 : 1 : 0.3 : 0.7 : 0.1 \quad .$$

Unfortunately, this pattern does not appear to correspond to present observations. Furthermore, the absolute number of events would be too small if the magnetic coupling is of order $1/\Lambda_p$. (The sextet described above, if charged or if it has weak isospin, would also produce signals of the type discussed here.)

It appears that further study of exotic composites, with possible applications at CERN or the SSC, is warranted.

STRONG P,C VIOLATION

In the models with purely fermionic preons it appears impossible to preserve global flavor chiral symmetry if the gauge precolor interaction is vector-like[2]. By contrast, in left-right asymmetric chiral theories the necessary chiral preservation conditions can be satisfied [1,2]. Thus, we are biased to believe that a preon theory is likely to violate parity (P) and charge conjugation (C) (but perhaps not PC) by large amounts at the scale Λ_p, since the classification of composites and the dominant strong interactions (precolor) are expected to be left-right asymmetric. (Theories with scalar preons may avoid this conclusion). If we were doing experiments at energies near Λ_p this fact, if true, would be readily apparent. However, at low energies $E \ll \Lambda_p$, the effective Lagrangian contains symmetry violating effects in terms proportional to $1/\Lambda_p$ (such as 4-fermi terms), which are not easily detectable until the energies are sufficiently high. As emphasized since the early days [1,2], it is very interesting to test this distinguishing aspect of preon models by looking for P or C violating effects that increase with energy and eventually surpass in magnitude the parity violation of the weak interactions.

In this workshop we considered the possibility of polarized proton beams to help determine the chirality (and hence the P,C properties) of the contact 4 fermi interactions with strength $1/\Lambda_p^2$. The initial discussions revealed mainly the difficulties: At the SSC energies mainly the sea partons dominate[3]. Therefore we do not expect much effect from the initial polarization of the proton which is mostly shared by the valance quarks. Furthermore, in the final state we cannot distinguish a quark jet from an anti-quark jet. This prevents distinguishing final helicities! In addition, since the expected asymmetries are small at $E \ll \Lambda_p$, the parton distributions must be better understood to clearly disentangle the effect.

Another possible test of P,C violation in contact terms involves polarized electron-electron or electron-positron scattering. The available energies and resulting effect are of course much smaller, but the signal would be cleaner. One could measure forward-backward asymmetries in $ee \to ee$ or $e\bar{e} \to e\bar{e}$ with initially polarized electrons. Some estimates have been done by Peskin[9]. A more interesting test would be a measurement of $e\bar{e} \to \tau\bar{\tau}$ with initially polarized beams and observation of the polarizations of stopped τ, $\bar{\tau}$. A deviation from the helicity amplitudes and rates of the standard model may indicate the presence of contact terms with certain helicity properties. General 4-fermi helicity amplitudes and cross sections that can be applied to this problem are defined in ref. 10.

In addition to separate P,C violation it is interesting to consider (PC-T) violation in the contact terms, although there is no compelling theoretical reason to expect it. A possible effect involves the quantity

$$\frac{d\sigma}{dq\,dy}\bigg|_{y=0}(p\bar{p} \to e\bar{\nu}) - \frac{d\sigma}{dq\,dy}\bigg|_{y=0}(\bar{p}p \to e\nu) \equiv A$$

A signal above the standard model background would result if the 4-fermi strength is $\lambda^2/2\Lambda^2$, with $\lambda^2 = 4\pi$, $\Lambda_p \sim 5$ TeV, and $\sqrt{q^2}$ is larger than 500 GeV. For example, at $\sqrt{q^2} = 1.5$ TeV the quantity A would be 5 times larger than the Drell-Yan background. The effect gets larger as $\sqrt{q^2}$ increases.

By far the surest way to observe the possible strong P,C violation is to do experiments at large energies $E \geq \Lambda_p$. If Λ_p is not large relative to the SSC center of mass parton-parton energies (see ref. 4) the best way to distinguish viable models may be through asymmetries. Estimates of asymmetries have not yet been done for such energies, but a formalism that provides the methods and some cross section estimates at $E \geq \Lambda_p$ has been developed.[10]

SMALL DEVIATIONS IN THE STANDARD MODEL

Most discussions on tests of compositeness at low energies concentrate on new phenomena due to the non-renormalizable pieces in the effective Lagrangian. However, because of the electroweak symmetry breaking, the effects of compositeness may trickle down to the dimension ≤ 4 operators in the standard model. These effects which boil down to shifts in the W, Z masses, the ρ parameter, or the Z couplings, may be measurably large in precision experiments as discussed in more detail refs. 11, 12.

As an example consider the γ-Z kinetic mixing of ref. 12. This arises from an SU(2) X U(1) invariant term in the effective Lagrangian

$$-\frac{1}{2}\,\xi\, g'/\Lambda^2 \;\; B_{\mu\nu}\, \mathrm{Tr}\left[M^\dagger\, [D_\mu\, D_\nu]\, M\, \tau_3\right]$$

where M is the 2 x 2 matrix of $[SU(2) \times U(1)]/U(1)$ Goldstone boson fields (effective Higgs), $D_\mu M = \partial_\mu M - ig\, W_\mu(\tau/2)\, M + ig'\, M(\tau_3/2)\, B_\mu$, and $B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$ is the field for B_μ. Note the presence of τ_3 that breaks the right handed isospin: This is assummed to arise from an underlying preon theory or extended technicolor theory) with the abiltiy to generate up-down mass differences. Electroweak spontaneous breakdown allows us to replace M by the vacuum expectation value v, thus generating B-Z or γ-Z mixing in the form

$$-\frac{1}{2}\,\xi' \mathrm{Cos}\theta\; A_{\mu\nu}\; Z^{\mu\nu}$$

where $A_{\mu\nu}$ and $Z_{\mu\nu}$ are the photon and Z field strengths, and tan $\theta = g'/g$ is the Weinberg angle. the parameter ξ' may be estimated by relating it to 4-fermi interactions that violate the right-handed isospin. It's magnitude may be as large as $\xi' \approx 0.01$, unless it is suppressed by a factor $(v/m_Q)^2$. m_Q is the dynamically generated mass of a high-color or technicolor quark.

It can be shown that this mixing ξ' leads to a redefinition of the measured Weinberg angle $\tilde{S}$

$$\tilde{S}^2 = \mathrm{Sin}^2\,\theta + \xi' \mathrm{Sin}\theta\, \mathrm{Cos}^2\theta \;.$$

and to a modification of the relationship between the measured masses of W, Z and the Weinberg angle $\tilde{S}$, and electric charge:

$$m_W = \frac{1}{2}\left(\frac{ev}{\tilde{S}}\right)\left[1 + \frac{\xi'}{2\tilde{S}}(1-\tilde{S}^2) + \ldots\right]$$

$$m_Z = \frac{1}{2}\,\frac{ev}{\tilde{S}\sqrt{1-\tilde{S}^2}}\left[1 - \xi'/2\tilde{S} + \ldots\right]$$

This causes a shift in m_W or in m_Z up to 1% if not suppressed by v/m_Q. However, since the effective interaction above is only an example, a more general structure can lead to a larger shift in m_Z, m_W, according to the more general formulas in ref. 12. It is hard to obtain a precision measurements of m_W with present techniques. However, measurements at the Z resonance could reveal shifts in m_Z, as well as modifications in the couplings of the Z, as in eq. 46 of ref. 12, thus signaling deviations from the standard model due to the presence of a higher scale.

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EUROPEAN FACILITIES AND PLANS

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Summary

The capabilities and time schedules of present and planned accelerators in Europe are reviewed. The history of the Large Hadron Collider (LHC) project is recalled, and the results of machine feasibility studies are summarized. It seems possible to build in the LEP tunnel a pp collider with $\sqrt{s} = 10$ to 18 TeV and a luminosity of 10^{33} cm^{-2} sec^{-1}. Results from the Lausanne Workshop on LHC physics are reviewed, and some aspects of Higgs and supersymmetric particle production are discussed.

Present and Planned Accelerators

Sp$\bar{\text{p}}$S Collider. In the past this machine has operated at a centre-of-mass energy $\sqrt{s} = 546$ GeV at an instantaneous luminosity up to 2×10^{29} cm^{-2} sec^{-1}, and has accumulated an integrated luminosity of about 10^2 nb^{-1}. Plans for this year include a test for an exploratory run in a pulsed mode, with the beam energy oscillating between 120 and 450 GeV on a time-scale measured in tens of seconds. It is hoped to achieve next year an instantaneous luminosity of order 10^{25} cm^{-2} at $\sqrt{s} = 900$ GeV, which should enable searches for Centauro events or any other novel phenomena with a threshold below $\sqrt{s} = 1$ TeV. Also planned for this year is a conventional physics run at $\sqrt{s} = 620$ GeV for about 12 weeks between September and December 1984, during which, with the aid of some improvement in the instantaneous luminosity, one should be able to increase the present integrated luminosity by a factor of 3 to 5. In the future, it is planned to upgrade the Sp$\bar{\text{p}}$S Collider with a new $\bar{p}$ accumulator ring called ACOL, which should increase the peak luminosity by O(10) from 1987 onwards. One may hope to accumulate an integrated luminosity of $O(10^3)$ nb^{-1} before then, and perhaps $O(10^4)$ nb^{-1} afterwards. Plans for detector upgrades include a microvertex detector for UA1 during 1984, and subsequently improved calorimetry. UA2 plans include Ring Imaging Cerenkov (RICH) in the forward direction during 1984, and improved solid angle coverage for their calorimetry in 1987.

LEP. Access pits are being dug, as well as the part of the tunnel under the Jura mountains. The first e^+e^- collisions are expected towards the end of 1988. In its first phase (LEP I) the centre-of-mass energy $\sqrt{s} \leq 103$ GeV with an instantaneous luminosity $\leq 10^{31}$ cm^{-2} sec^{-1}. The cost of LEP I is about 4.5×10^8\$. Probably it will be run for at least two years at the Z° peak, gathering up to 10^7 Z° decays. Also, the latest data on the t quark make it seem quite likely that toponium is too heavy to be produced at previous lower energy e^+e^- machines, and that LEP may provide the first detailed exploration of the toponium system, something which might take a year or more. Beyond LEP I there are options for increasing the available $\sqrt{s}$, called generically LEP II. One possibility is to fill up the first two planned RF galleries with superconducting RF. This would provide $\sqrt{s}$ up to 140 GeV at luminosities up to 10^{32} cm^{-2} sec^{-1}. There is also provision for excavating two other RF galleries and filling them with superconducting RF which should give $\sqrt{s} = 180$ GeV at a luminosity of 10^{32} cm^{-2} sec^{-1}. The cost of LEP II is estimated to be about 10^8\$, 1 1/2 times the part of the annual CERN budget available for LEP construction. The additional RF for LEP II could be installed during shutdowns of LEP I, which is not expected to run for more than about 2000 hours per year. It is not unreasonable to foresee LEP II starting to operate above the W^+W^- threshold in 1992, and then running there for at least two years. Four detectors are currently under construction for LEP. Among these are ALEPH and DELPHI which feature TPC's and other slow drift devices, and OPAL and L3 which do not have such slow drift devices, and could perhaps be used in a modified form at the LHC.

HERA. This ep collider was finally approved in April 1984, and civil engineering is expected to start soon. It is planned to collide 30 GeV $e^\pm$ with 820 GeV protons for a total center-of-mass energy of 314 GeV, with a luminosity around 10^{32} cm^{-2} sec^{-1}. A simpler, cheaper and higher field magnet design is now being prepared, which might enable higher proton energies of O(1) TeV to be achieved. It is expected to have electrons in their ring by early 1988, and the protons in mid 1989, with ep collisions in mid 1990. Intensive e^+e^- physics is not a priority for the machine, and going to 40 GeV per beam in the $e^\pm$ rings would require superconducting RF: this makes it unlikely that HERA will be suitable for a toponium factory. As for possible subsequent developments of HERA, the internal diameter of the tunnel is large enough to contain another ring. There are no prizes for guessing what that ring might be! The present cost of the HERA projected is estimated at around 2.5×10^8 \$. About a third of this is being covered by contributions of labour, materials and components by countries outside West Germany, including Italy, Canada, France, the Netherlands, Great Britain and Israel. This is a novel model for international participation in accelerator construction, which might be interesting for future projects such as the SSC. No experiments for HERA have yet been chosen, and it is expected that decisions will be made in late 1985.

History of the Large Hadron Collider (LHC) Idea. There were many informal discussions of a possible hadron collider in the LEP tunnel even before LEP was officially approved. Indeed, the possible subsequent installation of a hadron ring or rings was taken into consideration when choosing the LEP tunnel radius (as large as possible, since it might be the last piece of real estate available for accelerator construction in Europe, as well as to get to the highest possible e^+e^- energies) and diameter (large enough to leave room for hadron magnets). Discussion of an LHC was greatly stimulated by the recent success of the Sp$\bar{\text{p}}$S Collider. Serious work was, however, triggered by the initiation of the SSC project in the United States. There were prepatory meetings at CERN in December 1983 and February 1984, in which about 40 leading European physicists participated. Also during the winter 1983/1984 there was an LHC machine study group at CERN which produced the feasibility study[1] reviewed in the next section. Then, in March 1984, about 150 physicists met and worked on LHC physics at Lausanne for four days, with a subsequent two day open presentation at CERN.[2] All this activity was summarized in presentations[3] at the ICFA meeting at KEK near Tokyo in May 1984.

LHC Machine Studies

Possible Options. Although I am a total ignoramus about accelerators, I have been asked to review here the "feasibility study of possible options" for a Large Hadron Collider in the LEP tunnel made by the CERN Machine Group, so I will do the best that I can. The three main options considered have been:

(A) $p\bar{p}$ in a single beam channel, with either present day or futuristic high-field magnets. Such a device would yield relatively low luminosity and need a relatively large aperture so as to allow separation of the beams (Fig. 1a).

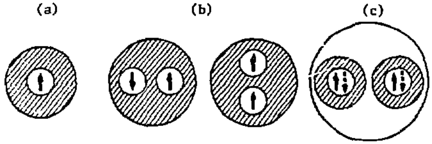

Figure 1 Possible Options for LHC Beam Channels

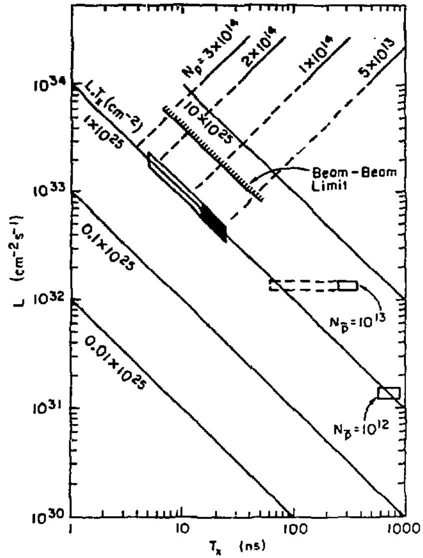

Figure 2 Luminosity and Bunch Separation

(B) Two beam channels with a common magnetic circuit (Fig. 1b). The space available in the LEP tunnel can only accommodate one cryostat to contain both beam channels. The most interesting possibility is to put them side by side, which permits a high luminosity pp machine with many bunches. If one uses present-day lower-field magnets one can have the magnetic fields in the two channels either antiparallel (allowing pp in separated channels or $p\bar{p}$ in the same channel) or parallel (allowing $p\bar{p}$ in separated channels). If one goes to high-field magnets only the antiparallel possibility (permitting pp) is available.

(C) Two beam channels with independent magnetic circuits (Fig. 1c). In this case one is restricted to moderate magnetic fields, but either pp or $p\bar{p}$ could be realized. Among these various options, (A) has received some attention but most interest has been focused on (B) which is the most demanding technologically as well as making the most physics available.[1] The attention paid to the different options in this brief review will reflect these priorities.

Luminosity and Bunch Separation. Most high energy hadron-hadron colliders use bunched beams: let us call the separation between successive beam crossings T_X , and the average number of collisions per beam crossing n . It is clear that if n is kept fixed, an increase in the luminosity L can only be obtained at the price of decreasing T_X , as seen in Fig. 2. Feasibility studies[2] of experimentation at a Large Hadron Collider suggest that it may be difficult to imagine working with n larger than 1, while one can do experiments with $T_X \geqslant 25$ ns. If these are strict limits, then the luminosity of a collider in the LEP tunnel is restricted to $L \leqslant 4 \times 10^{32}$ cm^{-2} s^{-1}. However, if higher values of n are tolerable, one can increase L to 1.5×10^{33} cm^{-2} s^{-1} (corresponding to $n \sim 4$) before hitting the beam-beam tune-shift limit of 0.0025 reached at the $Sp\bar{p}S$ collider. Such a luminosity is feasible with the pp option, but would require between 3000 and 4000 bunches. Matching the RF in the SPS and the LHC is only possible for a few choices of bunch number, and the only choice in this range is 3564. With $N_p = 5 \times 10^{13}$ protons the stored energy would be 70 MJ , a reasonable number, while there should be no problems with beam-beam effects as long as $n \lesssim 1$ ($\Delta Q \leqslant 0.0013$). The luminosities in $p\bar{p}$ machines are restricted by the number of antiprotons available. Possible choices of T_X and n for $N_{\bar{p}} = 10^{12}$ and 10^{13} are shown in Fig. 2: luminosities in the range $L = 1.5 \times 10^{31}$ cm^{-2} s^{-1} to 1.5×10^{32} cm^{-2} s^{-1} would be possible.

pp Option. General parameters[1] of the LHC pp option are set out in the Table. There would be eight

Table: General Parameters and Performance

General Parameters

Collider Type in LEP	Proton-Proton	
Separation Between Orbits (mm)	165-180	
Number of Bunches	3564	
Bunch Spacing (ns)	25	
Number of Crossing Points	8	
Beta Value at Crossing Point (m)	1	
Normalized Emittance $4\pi\gamma\sigma^2/\beta$ (μm)	5π	
Full Bunch Length (m)	0.31	
Full Crossing Angle (μrad)	96	
Lattice Period Length (m)	79	158
Lattice Phase Advance	$\pi/3$	$\pi/2$
Dipole Magnetic Field (T)	10	10
Operating Beam Energy (TeV)	8.14	8.99

Performance

<n> at σ = 100 (mb)	1	4
Luminosity ($cm^{-2}\ s^{-1}$)	4×10^{32}	1.5×10^{33}
Number of Particles/Bunch	1.34×10^{10}	2.6×10^{10}
Circulating Current (mA)	86	167
Beam-Beam Tune Shift	0.0013	0.0025
Beam Stored Energy (MJ)	63	121
RMS Beam Radius (μm) *	12	
Beam Life-Time (h) **	42	21

* at interaction point for $\beta^* = 1$ m
** particle loss due to beam-beam collisions

crossing points, though only six of these might be available for experiments, with two (or possibly just one) being reserved for dumping the beam. The length of the ring is fixed by the size of the LEP tunnel: 26.658 km. Each straight section would be 0.490 km long and the average bending radius would be 3.494 km. Thus the maximum beam energy is fixed by the attainable magnetic fields, which has been limited to B = 10T in the studies made. In this case the maximum beam energy is between 8 and 9 TeV, the latter requiring a longer lattice period which would not normally be the preference of the machine physicists, though experimentalists would obviously prefer the highest possible beam energy. The beam crossing angle would be 96 μrad.

Magnet System. Existing magnet technology, as developed by FNAL in particular and also BNL, DESY, KEK and UNK, can reach $B \leq 6$ or 7T. Getting to 10T would be a technological challenge which will be discussed in a moment. General features of the magnetic system which would be largely independent of the maximum field value are the following. For reasons of space and economy interest centres on a two in one design with a common magnet yoke and cryostat. For simplicity, economy and the highest possible beam energy one would like the dipoles to be as long as possible. However, the geometry of the LEP access pits restrict one to dipoles of length ≤ 12 m. The area available above the LEP magnets is about 0.8 m horizontally by 1.1 m vertically. A design[1] for the magnet and its cryostat which fits within this cross-section while having two side by side beam channels is shown in Fig. 3.

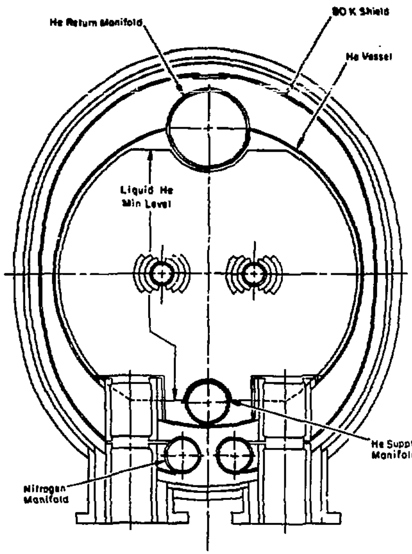

Figure 3 Possible Two-Channel Magnet Configuration

For reasons of space the vacuum chambers would be cold. The inner diameter of the superconducting coils would be 35 to 50 mm (the smaller the better from the point of view of the magnet designers). The ramping time needed to get from the field of 0.5 T required at the injection energy of 450 GeV from the SPS, to the maximum of 10 T at 9 TeV would be about 10 minutes. Magnet system parameters[1] for different period lengths are shown in Table 2.

10 T Magnets. To get B = 10 T in the magnet aperture one needs B = 11 T in the conductor itself. The maximum current required is j_c = 1300 A/mm^2, corresponding to $\langle j_c \rangle$ = 300 A/mm^2 after dilution by copper, a filling factor and allowing for a 20% safety factor. There are two candidate superconductors for attaining such a performance: Nb_3 Sn at 4.2° K and Nb Ti + Ta at 1.8° K, as seen in Fig. 4. It seems that Nb_3 Sn may have the greater potential for future development, in the sense that the attainable j_c falls more slowly as B is increased. So far, for reasons of time, LHC machine studies[1] have only included Nb_3 Sn at 4.2°K: it is planned to compare later with $Nb_3\ T_i + T_a$ at 1.8° K. Small quantities of Nb_3 Sn superconductor have been made in industry, but there is a long way to go to reach the thousands of tons needed for an LHC magnet system. Moreover, the mechanical properties of Nb_3Sn cause engineering problems for magnet fabrication. Therefore considerable research and development work will be needed before a 10 T magnet system can become a reality. As discussed earlier, there is time in the European

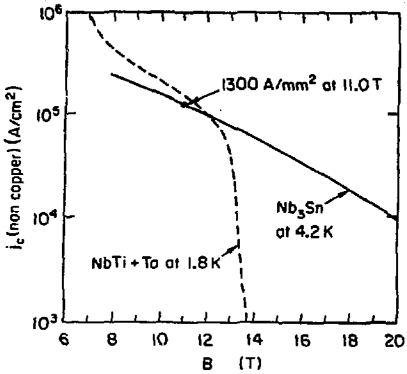

Figure 4 Performances of Candidate Superconductors

schedule for such an R and D programme to be undertaken, and several European laboratories are interested in participating in it.

Other Subjects Studied. The purpose of this section is just to inform you that a serious study has been made of most machine systems: please see the LHC machine study group report[1] for details. Practicable designs for quadrupoles and sextupoles have been made. The scaling of the magnets for different apertures between 35 and 50 mm has been studied. A 4.5° K cryogenic system has been designed—one interesting feature is the 1.5° slope of the tunnel which creates as pressure difference of 1.4 bar across the ring for liquid helium—annoying but not impossible to live with. The vacuum system has been studied. For the RF system it is proposed to use a frequency of 400.8 MHz , just twice the SPS RF frequency, which gives flexibility in the choice of bunch number. The SPS will be used to inject bunches of 10^{11} particles at 450 GeV . Since the SPS already provides up to 3×10^{13} protons per pulse and has a repetition rate of 10s , it will quickly provide the 5×10^{13} protons needed for LHC in the pp option. Since the ratio of LHC to SPS circumferences is 27/7 , the entire LHC ring can be filled with four SPS pulses. Two possibilities for the transfer lines have been investigated. One involves two 2 km transfer lines with large bending radii for which conventional magnets could be used, while the other involves two 1 km transfer lines with smaller radii which would need superconducting magnets. The shorter transfer lines would need a 12% slope to get down to the LEP/LHC tunnel, which is lower than the SPS tunnel, but the problems this raises for the transfer line cryogenic system seem to be soluble. Among other topics studied which seem to raise no particular problems are the beam dumps and radiation protection.

$p\bar{p}$ Option. In this case the attainable luminosity is limited by the $\bar{p}$ accumulation rate A . The rate required is $\geqslant N_{\bar{p}}$ divided by the luminosity decay time T_L . For T_L = 20 hours and $N_{\bar{p}} = 10^{12}$ one needs $A \geqslant 5 \times 10^{10}$ per hour. This is about an order of magnitude larger than that attained with the present CERN $\bar{p}$ source, but is expected to be reached by the FNAL Tevatron I and CERN ACOL sources. As seen in Fig. 2, with $N_{\bar{p}} = 10^{12}$ one can attain $L \leqslant 1.5 \times 10^{31}$ cm^{-2} s^{-1} with 108 bunches corresponding to T_X = 825 ns. A more sophisticated source would probably be designed to reach $A = 5 \times 10^{11}$/hour , permitting $N_{\bar{p}} = 10^{13}$ and $L \leqslant 1.5 \times 10^{32}$ cm^{-2} s^{-1} . However, one would either require a complicated separation scheme to avoid interactions at about 2000 unwanted crossing points (each separator being about 40 m long) or have many events per bunch crossing (unacceptable to the experimentalists?) or some judicious combination of these two extremes. These difficulties would not arise with a low luminosity $p\bar{p}$ machine, or one with two beam channels.

Conclusions on LHC Design. 1) One could build a 10^{33} cm^{-2} s^{-1} luminosity pp machine in the LEP tunnel which would yield $\sqrt{s}$ = 18 TeV with 10 T magnets.

2) To achieve this, one needs a vigorous programme for the development of materials and techniques needed to make such magnets.

3) One would easily reach $\sqrt{s}$ = 10 to 13 TeV with six to seven T magnets after a shorter development programme.

4) All other components and systems are feasible with present technology.

What would be the cost of the LHC? The cost of the magnets, scaled up from HERA, would be about 400 M \$, namely less than the total LEP cost and about 1/6 of the recently projected cost of the SSC. A simple-minded theorist's way of understanding this ratio is as follows. A factor of 1/3 comes from the smaller size of the LEP ring, and is reflected in the correspondingly lower energy of the LHC. Since calculations [4,5] suggest that the physics reach of such hadron-hadron colliders grows roughly as $s^{1/4}$ for fixed L , it becomes relevant to ask how much an extra factor of 1.5 to 2 in physics reach is worth. Another factor of 1/2 in cost comes from the existence of a tunnel and laboratory infrastructure at CERN. FNAL can provide the latter in the U.S., but not the former, so that siting the SSC at or near FNAL would not save such a large factor in cost.

A final comment concerns a possible ep collider. This would be available for free in the LEP tunnel: by colliding LEP and the LHC one could get $\sqrt{s} \gtrsim 2$ TeV for ep collisions. This possibility has been discussed enthusiastically by theorists,[6] but no detailed machine study has yet been undertaken.

Lausanne Workshop

Questions. As mentioned earlier, this meeting involved 150 physicists meeting at Lausanne for four days, followed by a two day public presentation at CERN. Needless to say, most of the work was done beforehand by the different working groups. The big questions to be addressed by the physicists at Lausanne were:

a) Do you want $\sqrt{s}$ = 18 TeV , or would you be satisfied with 10 TeV ?

b) Can you work with $L = 10^{33}$ cm^{-2} s^{-1} , or are you limited to $L \leqslant 10^{32}$ cm^{-2} s^{-1} ?

c) Would you prefer $n \sim 1$ and $T_X \sim 25$ ns , or would you prefer $n \sim 6$ and $T_X \sim 150$ ns ?

d) How interesting do you find $p\bar{p}$ at $L = 10^{32}\ cm^{-2}\ s^{-1}$?

The experimental physicists seeking answers to these questions were divided into 8 working groups, whose subjects and conveners are listed in ref. 6. To aid (?) them in their deliberations there were the seven plenary talks by theorists listed in ref. 7, and a continuous theoretical discussion session held in parallel to the experimental working groups.

Summary of Experimental Conclusions

Jets. A 1 TeV jet is expected to have a multiplicity or order 100, with 1/2 of its energy as carried by particles with energy fractions $z < 0.04$, and 1/2 of the jet energy dispersed outside a 5° cone about the jet axis. These expectations must be kept in mind when considering jet energy resolution and separation in angle. In order to determine transverse momentum balance and hence be able to look for missing transverse energy one needs to be able to measure muon momenta. Going from $n = 1$ event per crossing to $O(5)$ has little effect on such physics desiderata as the mass resolution in dijet combinations. It was not obvious that a magnetic field was needed for jet physics. The working group felt that it could operate with bunch separations as low as $T_X = 10$ ns .

Electrons and Photons. It was thought that calorimetry was the only feasible technique, and that a "traditional" resolution $\Delta E/E \sim 0.15/E^{1/2}$ would be adequate. No obstacle was seen to working with $n > 1$ and $T_X = 25$ ns. Most of the above remarks apply to the detection of isolated e and γ . The only conceivable technique for detecting electrons inside a hadronic jet was thought to be transition radiation detection, but this was not pursued in detail. This group also did not feel a strong need for a magnetic field.

Muons. This group felt able to identify muons with p_T larger than a few GeV , and to measure the sign up to momenta of 2 TeV . A preference was expressed for $n = 1$, while $T_X = 25$ ns was found acceptable.

Tracking and Vertex Detector. These groups wanted $n = 1$. They thought that drift chambers would be adequate for tracking. Silicon strips, CCD's, scintillating fibres and Josephson junction devices were considered as possible vertex detectors. CCD's were preferred on grounds of versatility, but would need to be speeded up to operate at the LHC with high luminosity. Concern was expressed about radiation damage to vertex detectors, which might have a lifetime of only about 1 year if $L \sim 10^{33}\ cm^{-2}\ s^{-1}$.

Triggering, Data Acquisition and Processing. Philosophies for these tasks were developed and no essential difficulties were forseen.

Based on these studies, a summary[2] of the big answers to the big questions posed is as follows.

a) No dramatic threshold is expected, but everybody prefers $\sqrt{s} = 18$ TeV .

b) It is possible to do experiments with $L \sim 10^{33}\ cm^{-2}\ s^{-1}$.

c) While much physics could be done with $n > 1$, $n \sim 1$ is preferred for tracking in particular, while nobody objected strongly to $T_X \sim 25$ ns .

d) No-one saw an overwhelming physics case for $p\bar{p}$ collisions at $L \sim 10^{32}\ cm^{-2}\ s^{-1}$.

Summary of Theoretical Discussions. For reasons of time, space and to avoid duplication, I will only summarize here some studies which are of particular interest to the focus of this Workshop, namely electroweak symmetry breaking, Higgses etc. This area was studied before, during and after the Lausanne meeting by Graciela Gelmini, Heniek Kowalski and myself.[5] Our most detailed investigations were of Weinberg-Salam elementary Higgs production, signatures and backgrounds, and of supersymmetric particles, mainly gluinos and squarks.

We found[5] that the detection of H produced by gg or WW collisions was difficult for $2m_W < m_H \lesssim O(400)$ GeV because of the large WW pair production background. One reaction which seems to give a more favourable signal-to-background ratio, though a smaller total cross-section, is $q\bar{q} \to W^* \to H + W$ production, where there is an additional W in the final state to tag the event. Another possibly interesting reaction might be gg or $q\bar{q} \to \bar{t}t + H$, although the background from $\bar{t}t + WW$ has not yet been evaluated and it is not clear how efficiently the trigger $\bar{t}t$ pair could be tagged. We also thought about the detection of Higgses with $100\ GeV < m_H < 2\, m_W$, in which range their dominant decay mode would be $\bar{t}t$. The backgrounds in $gg \to WW \to H$ and gg or $q\bar{q} \to \bar{t}t + H$ seem to be overwhelming, but the background from $q\bar{q} \to \bar{t}t + W$ to the reaction $q\bar{q} \to W^* \to H + W$ seems to be no larger than the signal, and might enable it to be detected.

As concerns supersymmetry, in addition to computing the total cross-sections for $\tilde{g}\tilde{g}$ and $\tilde{q}\bar{\tilde{q}}$ production, we also investigated their missing p_T decay signatures using a Monte Carlo programme based on $\tilde{g} \to q\bar{q}\tilde{\gamma}$ and $\tilde{q} \to q\tilde{g}$ or $q\tilde{\gamma}$ decay. We found that $\tilde{g}$ or $\tilde{q}$ with masses up to $O(2)$ TeV could be detected with missing p_T of several hundred GeV and several well-separated hadronic jets (four from $\tilde{g}\tilde{g}$, $\tilde{g} \to q\bar{q}\tilde{\gamma}$, six from $\tilde{q}\bar{\tilde{q}}$, $\tilde{q} \to q\tilde{g}$, $\tilde{g} \to q\bar{q}\tilde{\gamma}$, two from $\tilde{q}\bar{\tilde{q}}$, $\tilde{q} \to q\tilde{\gamma}$) each with $p_T \gtrsim O(100)$ GeV . The azimuthal angle of the missing p_T vector is not strongly correlated with the observed jet azimuthal angles, enabling heavy sparticles to be clearly separated from semileptonic quark decay backgrounds which give a charged lepton and missing neutrino p_T vector strongly correlated with one of the observed jets. We concluded $\tilde{g}$ or $\tilde{q}$ should be quite easy to detect.

For more details of our work, see ref. 5. For subsequent discussions of the signatures of intermediate mass Higgses and of supersymmetry, see the corresponding study group summaries[8] in these Proceedings.

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A SADDLE-POINT SOLUTION IN THE WEINBERG-SALAM THEORY*

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Summary

We briefly discuss some properties of a new saddle-point solution (energy $\sim$ 10 TeV) in the standard theory of the electroweak interactions.

Recently it was shown that the field configuration space of the classical Weinberg-Salam theory[1] without fermions has noncontractible loops passing through the vacuum configuration.[2] This makes it likely that there is a static classical solution of the field equations, which is a saddle-point of the energy functional, and therefore unstable. The solution would be the maximal energy configuration on some noncontractible loop, and all other loops homotopic to this one would pass through configurations of equal or greater energy. It would therefore be at the top of the energy barrier for going from the vacuum to the vacuum along a topologically nontrivial path. This energy barrier has a definite height because the Weinberg-Salam theory has a mass scale (the Higgs vacuum expectation value), in contrast to the case of a pure gauge theory such as Quantum Chromodynamics.

One believes that there is such a saddle-point solution by analogy with work of Taubes,[3] who has shown rigorously that in a slightly different context, namely the zero-monopole sector of the SO(3) gauge theory with an adjoint Higgs field and vanishing Higgs potential, one can apply Morse theory arguments to relate topological information about the space of field configurations to the existence of stationary points of the energy functional. Forgács and Horváth[4] have reviewed a number of other field theories in one, two, and three spatial dimensions, where explicitly known saddle-point solutions are related to the topology of the field configuration space.

We have coined the word "sphaleron"[5] to describe any classical solution of this type in a relativistic field theory. A sphaleron, being static and localized in space, is particle-like, but since it is unstable, we do not want to call it a soliton. Unlike a soliton, a sphaleron almost certainly does not correspond to a stable particle state in the quantum theory.

One can find a good approximation to a sphaleron in the Weinberg-Salam theory, which has gauge group $SU(2) \times U(1)$, by expanding to lowest nontrivial order in the weak mixing angle θ_W. In the limit that θ_W vanishes the U(1) field decouples and may consistently be set to zero in the field equations. There remains an SU(2) theory with a doublet Higgs field. Dashen, Hasslacher, and Neveu (DHN)[6] discovered some time ago a static solution in this theory, which was later rediscovered by Boguta.[7] The solution has finite energy, and the energy density is localized and spherically symmetric. The fields, strictly speaking, are only axially symmetric. Burzlaff[8] proved recently that this solution rigorously exists, and also proved that it is unstable, by presenting a one-parameter family of field configurations, among which it is the configuration of maximal energy. In fact, there is a noncontractible loop in the configuration space passing through the vacuum and the DHN sphaleron, on which the sphaleron is the configuration of maximal energy.

We have estimated numerically the energy of the SU(2) sphaleron for the whole range of values of the quartic Higgs coupling λ (neither Dashen et al. nor Boguta had done this). We find that the sphaleron energy E_S increases from 7.6 TeV for $\lambda = 0$ to 13.5 TeV for $\lambda = \infty$ (see the Appendix for details).

When $\theta_W \neq 0$ the presence of the U(1) field makes it impossible for the solution to remain spherically symmetric in any sense. However, it is quite easy to find the changes in the sphaleron's properties to leading order in θ_W. To first order, the SU(2) gauge field and Higgs field remain unchanged, but they produce a U(1) current density which is axially symmetric and which acts as a source for the U(1) gauge field. By calculating the asymptotic form of this U(1) field, we find that the sphaleron has a magnetic dipole moment of strength $\sim 0.3\ \mathrm{GeV}^{-1}$. This is about 80 times larger than the magnetic moment of the W-boson, which is e/M_W.[9] The energy of the sphaleron decreases by $\sim$ 1% relative to the energy when $\theta_W = 0$.

Another interesting property of the sphaleron is that it has a baryon number and a lepton number of 1/2. This is a direct consequence of the anomalies in the fermionic currents of the Weinberg-Salam theory, whose significance was first discussed by 't Hooft.[10] These fractional charges are also related to the existence of a zero binding energy solution to the Dirac equation in the SU(2) sphaleron background. Consider for simplicity the SU(2) part of the Weinberg-Salam theory with only leptons, that is, with an SU(2) doublet $\Psi_L = (e_L^-, \nu_L)$ and with an SU(2) singlet e_R^-, and let these leptons be massless (no Yukawa term in the Lagrangian). In the sphaleron background there is a normalizable zero energy solution to the classical Dirac equations: Ψ_L as given in Ref. 11 and $e_R = 0$. Following Jackiw and Rebbi[12] this then implies that the sphaleron has lepton number 1/2, which agrees with the result derived from the anomaly equation. A similar analysis holds for the Weinberg-Salam theory with only quark fields. Thus in the full theory with one generation of quarks and leptons the sphaleron also has baryon number 1/2.

It has been claimed by 't Hooft[10] that tunneling between topologically distinct vacua is negligible in the weak interactions because the Euclidean action is so large, namely $> 8\pi^2/g^2$. But this argument only applies to a virtual quantum process. If there were sufficient real energy available, more than the sphaleron energy E_S, the tunneling process might be enhanced. There seem to be two situations, at least, where this could happen. The first is in high energy collisions of particles from a very powerful accelerator, and the signature of the process would be the violation of baryon and lepton number conservation. The other situation is for a system at very high temperature ($kT > E_S$). Thermal fluctuations might then produce the baryon number violating process via the sphaleron at a substantial rate. This could be important in the Universe at early times ($t \sim 10^{-15}$s, $dT \sim 10$ TeV), where these processes could be related to the baryon number as observed today, cf. Ref. 13. But, before one can address these problems, a better understanding of the role of the sphaleron and other solutions is required.

*This summary is based on a paper submitted for publication in Physical Review D (preprint NSF-ITP-84-57) and was presented at this workshop by NSM.

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In conclusion, it is clear that already the standard Weinberg-Salam theory for the electroweak interactions contains some interesting non-perturbative structure.

Acknowledgments

We thank John Boguta for bringing Ref. 11 to our attention. This research is supported in part by the National Science Foundation under Grant No. PHY77-27084, supplemented by funds from the National Aeronautics and Space Administration. One of us (FRK) is supported by the Netherlands Organization for the Advancement of Pure Research (Z.W.O.).

Appendix

We will present here a table with numerical results for the energy of the SU(2) sphaleron (zeroth order in θ_W) and the magnetic dipole moment to order θ_W. Let us first establish our notation: v is the Higgs vacuum expectation value; λ is the quartic Higgs coupling; g and g' are the SU(2) and U(1) gauge coupling constants, respectively; the weak mixing angle θ_W is related to the electric charge e by $e = g \sin \theta_W = g' \cos \theta_W$. To be definite we take the following values: W-boson mass $M_W = \frac{1}{2} gv = 80$ GeV, $e^2/4\pi = 1/137$, and $\sin^2\theta_W = 0.23$.

The field configuration of the SU(2) sphaleron is determined by two radial functions. The differential equations for these functions, which follow from the field equations, can be integrated numerically. In the table below we give the resulting values for (i) the energy in units $4\pi v/g = 5.0$ TeV, and (ii) the magnetic dipole moment in units of $\frac{2\pi}{3} \frac{g'}{g^3 v} = 2.99\ e/M_W$. For comparison we also give the variational estimates of these quantities given in our original paper, to which we refer for further details.

λ/g^2	energy		magnetic moment	
	num.	var.	num.	var.
0	1.52	1.57	28	19.1
1	2.07	2.10	18	21.3
∞	2.70	2.72	16	19.5

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Dupe

THE THIN END OF THE WEDGE
OR
HOW SUGRA GUT-DRIVEN RADIATIVE ELECTROWEAK BREAKING CAN ACCOMMODATE A LIGHT TOP ONLY THROUGH A LIGHT HIGGS

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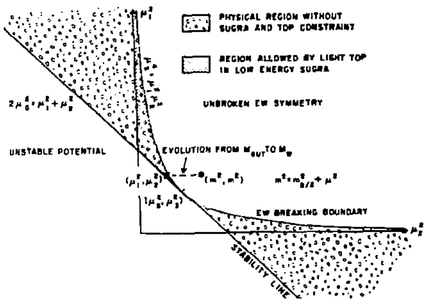

A top mass in the range 20–55 GeV imposes tight bounds on the light Higgs masses in the minimal SUSY-extended electroweak theory with soft SUSY-breaking and radiative EW-breaking induced at the same scale by an underlying N=1 SUGRA GUT. In particular, the lightest Higgs must weigh quite a bit less than $M_Z/2$, being therefore detectable at SLC and LEP. This will be a test for such theories.

1. SUSY 3-2-1 Theory

Consider the minimal (softly broken) SUSY extension[1] of the standard 3-2-1 theory retaining only[2] the heaviest quark/lepton generation. It has a spectrum of seven left-chiral superfields. In a transparent $z(R_3,R_2,Y_1)$ notation, they are:

$$\hat{Q}(3,2,1/3), \hat{T}^C(\bar{3},1,-4/3), \hat{B}^C(\bar{3},1,2/3), \hat{L}(1,2,-1),$$

$$\hat{\tau}^C(1,1,2), \hat{H}_T(1,2,1), \hat{H}_B(1,2,-1).$$

The first five from the left are R-parity-odd matter superfields while the remaining two are R-parity-even Higgs superfields. The most general R-parity invariant and 3-2-1 symmetric superpotential is

$$f_\ell = \lambda_T \hat{Q}\hat{T}^C\hat{H}_T + \lambda_B \hat{Q}\hat{B}^C\hat{H}_B + \lambda_\tau \hat{L}\hat{\tau}^C\hat{H}_B + \mu\hat{H}_T\hat{H}_B \,. \qquad (1)$$

Here λ's are dimensionless Yukawa coupling strengths and μ is a dimensional coefficient. The subscript ℓ in the LHS refers to our consideration as yet of only light superfields.

In this theory the Higgs potential, including the most general soft SUSY-breaking terms characterized by a mass-scale $O(M_W)$, has the following decomposition into quartic and quadratic parts:

$$V^H = V_4^H + V_2^H \,, \qquad (2a)$$

$$V_4^H = 1/2\, g_2^2\, (H_B^\dagger\, \vec{\tau}/2\, H_B + H_T^\dagger\, \vec{\tau}/2\, H_T)^2$$

$$+ 1/2\, g_1^2\, (1/2|H_T|^2 - 1/2\, |H_B|^2)^2, \qquad (2b)$$

$$V_2^H = \mu_1^2\, |H_B|^2 + \mu_2^2\, |H_T|^2 - \mu_3^2\, (H_B H_T + \text{h.c.}) \,. \qquad (2c)$$

Here $g_{2,1}$ are the SU(2), U(1) gauge coupling strengths and the Higgs doublets have been contracted by an ε-matrix. The dimensional coefficients $\mu_{1,2,3}$ are $O(M_W)$. One can choose a phase convention such that μ_3^2 is real and positive. Upon spontaneous EW breakdown, only the neutral Higgs components acquire VEVs

$$\langle H_B \rangle = \begin{pmatrix} h_1 \\ 0 \end{pmatrix}, \quad \langle H_T \rangle = \begin{pmatrix} 0 \\ h_2 \end{pmatrix}, \qquad (3)$$

where (ignoring CP-violation in the fermionic sector) $h_{1,2}$ can be taken to be real. The physical weak vector boson masses are then given by

$$M_{W,Z}^2 = 1/2\, (g_2^2,\, g_1^2 + g_2^2)(h_1^2 + h_2^2). \qquad (4)$$

Near the minimum, V^H looks like

$$V^H(h_1,h_2) = 1/8\, (g_1^2 + g_2^2)(h_1^2 - h_2^2)^2$$

$$+ (h_1 \;\; h_2) \begin{pmatrix} \mu_1^2 & -\mu_3^2 \\ -\mu_3^2 & \mu_2^2 \end{pmatrix} \begin{pmatrix} h_1 \\ h_2 \end{pmatrix} \,. \qquad (5)$$

Stability along the direction $h_1 = h_2$ implies

$$\mu_3^2 \leq 1/2\, (\mu_1^2 + \mu_2^2), \qquad (6)$$

while a necessary condition for EW symmetry breakdown is that the determinant of the matrix in (5) be negative, i.e. that

$$\mu_1^2 \mu_2^2 < \mu_3^4. \qquad (7)$$

The boundaries corresponding to (6) and (7) are given in the figure on the abstract. The minimization of $V^H(h_1,h_2)$ leads to two further relations (with $\cot\theta \equiv h_2/h_1$):

$$\frac{2h_1h_2}{h_1^2+h_2^2} \equiv \sin 2\theta = \frac{2\mu_3^2}{\mu_1^2+\mu_2^2} \,, \qquad (8a)$$

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$$\mu_1^2 \frac{h_1^2}{h_1^2-h_2^2} - \mu_2^2 \frac{h_2^2}{h_2^2-h_1^2} = 1/2\, M_Z^2. \tag{8b}$$

The physical Higgs spectrum has five particles $H^{\pm}$ and $H^{\circ}_{a,b,c}$. Of the neutrals $H^{\circ}_{a,b}$ emerge from the real parts of the neutral Higgs fields and act as scalars with respect to fermionic interactions while H°_c involves an imaginary part and acts as a pseudoscalar. The physical Higgs masses, calculated from (2), are

$$m_c^2 = \mu_1^2 + \mu_2^2\ , \tag{9a}$$

$$2m^2_{a,b} = m_c^2 + M_Z^2 \pm \{(m_c^2+M_Z^2)^2$$

$$- 4m_c^2 M_Z^2 \cos^2 2\theta]^{1/2} \geq\ ,\ \leq 2M_Z^2, \tag{9b}$$

$$m_{\pm}^2 = m_c^2 + M_W^2 \geq M_W^2\ . \tag{9c}$$

Thus[1] the charged Higgses have to weigh more than the W while the two neutral scalars must be on either side of the Z in mass. Further, it follows that $m_b \leq m_c \leq m_a$. It is also useful to write a mass-relation linking m_a and m_b but without involving m_c:

$$m_a^2 = M_Z^2\ (1-m_b^2/M_Z^2)\ (1 - \frac{1}{\cos^2 2\theta}\, m_b^2/M_Z^2)^{-1}. \tag{10}$$

Stronger bounds can be put on $m_{a,b}$ if $\mu^2_{1,2}$ are known to have the same sign. (As discussed later, this is what is forced by a light top in the SUGRA GUT-driven radiative EW breakdown scenario.) In that case one can take square roots of both sides of (7). Thus defining $\delta \equiv h_2/h_1 - 1 = \cot\theta - 1$, i.e. $\cos^2 2\theta = [1 - (1 + \delta + \delta^2/2)^{-1}]^2$, one can rewrite (6) and (7) as

$$|\mu_1\mu_2| \leq \mu_3^2 = 1/2\ (\mu_1^2 + \mu_2^2)\ \frac{\left(1 + \frac{\delta^2}{1+\delta}\right)^{-1}}{\sin 2\theta} \tag{11}$$

The quadratic inequality (11) on δ implies

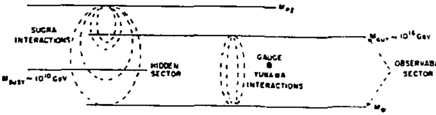
$$-\alpha\ (1-\alpha) \leq \delta \leq 1 - \alpha, \tag{12a}$$

$$\alpha \equiv \min.\ \left|\frac{\mu_1}{\mu_2}\right|,\ \left|\frac{\mu_2}{\mu_1}\right|\ . \tag{12b}$$

Since $0 \leq \alpha \leq 1$, the lowest that the lower bound in (12a) can go down to is $-1/4$ while the upper bound can come up to, at most, 1. Thus, the two VEVs $h_{1,2}$ now have to be roughly of the same order of magnitude. One can easily lower the upper bound on m_b somewhat below M_Z (e.g. $m_b \leq M_Z \cos 2\theta \leq 3/5\ M_Z$), but far stronger bounds are obtained below by restricting δ more stringently by use of the underlying SUGRA dynamics.

2. Minimal Low Energy Supergravity[3] and Radiative[4] EW Breaking

We[5] proceed in a framework in which the above 3-2-1 theory with a minimal softly broken supersymmetric extension is in fact the low-energy remnant of a deeper underlying theory: an N = 1 SUGRA coupled[3] to a GUT. We assume not any specific model for the GUT part but only the unification scale $M_{GUT} \sim 10^{16}$ GeV, the unified fine structure constant $\alpha_{GUT} \sim 1/24$ and the absence of significant new thresholds below M_{GUT}. There are superheavy (vector and chiral) superfields $\hat{V}^s, \hat{z}^s$ as well as light ones $\hat{V}^{\ell}, \hat{z}^{\ell}$ and these define the observable sector where gauge and Yukawa forces operate. In contrast, there is a gauge singlet hidden sector consisting of chiral superfields generically described by z^h at a typical scale $M_{SUSY} \sim 10^{10}$ GeV. This sector has only supergravity interactions though the latter cover all superfields and span all scales. The situation is pictorially represented in the following figure.



We work in the class of theories where the total superpotential can be additively split between contributions from the observable and the hidden sectors:

$$f(\hat{z}) = f_{obs.}(\hat{z}^{\ell}, \hat{z}^{s}) + f_{hid.}(\hat{z}^{h}). \tag{13}$$

Local SUSY is taken to be spontaneously broken at $M_{SUSY} \sim 10^{10}$ GeV by a nonzero Kähler covariant derivative VEV in the hidden sector. The embryotic goldstino gets eaten up by the gravitino through the super-Higgs mechanism, consequent to which the hidden sector superfields decouple from the rest of the theory leaving a physical gravitino of mass $\sim \sqrt{1/3}\ \kappa\ \langle F(z^h)\rangle \sim \sqrt{8\pi/3} \times M_{SUSY}/M_{P\ell}$. Here κ is the rationalized gravitational constant and $F(z^h)$ is the auxiliary component of $\hat{z}^h$; $\langle F(z^h)\rangle \sim M^2_{SUSY}$ and $\langle z^h\rangle \sim M_{P\ell}$. The gravitino mass, characteristic of the scale of the residual soft SUSY-breaking terms after the passage to lower energies, is identified with the weak scale and is taken to be $O(M_W)$.

The term in the Lagrangian $\mathscr{L}$(N = 1 SUGRA + matter), that is most relevant to us, is the scalar potential given by

$$V = e^{\kappa^2 d(z^N, z^*_M)}\ \{F_N (d^{-1})^N_M\ F^{*M}$$

$$- 3\kappa^2 |f|^2\} + 1/2\ \mathrm{Re}\ f^{-1}_{ab}\ D_a D_b. \tag{14}$$

Here $d(z,z^*)$ is the Kähler potential which is a real function defined in a Kähler manifold of chiral scalar fields z^N and their complex conjugates $\bar{z}_N$. F_N is the Kähler covariant derivative of the superpotential f, i.e. $F_N = f_N + \kappa^2 d_N f$ with $f_N \equiv \partial f/\partial z^N$, $d_N = \partial d/\partial z^N$. Further, $(d^{-1})^N_M$ is the inverse of the Kähler metric tensor $d^M_N \equiv \partial^2 d/\partial z^N \partial z^*_M$. Moreover, f_{ab} is an analytic function of z^N which transforms as a symmetric product of two adjoint representations of the GUT gauge group. (It occurs in the kinetic energy term $\mathscr{L}_{KE} = f_{ab}\hat{W}_a\hat{W}_b$, $\hat{W}^a$ being the spinor superfield for the gauge field strength). Finally, we have $D_a = d_N\ (gT)^N_{aM}\ z^M$, where $(gT)_a$ refers to a generator of the gauge group multiplied by the corresponding coupling. (In the $\kappa \to 0$ limit the RHS of (14) goes into the familiar $f_N(z) f^{*N}(z) + 1/2\ D_a D_a$ form.) We shall not assume any particularly simple structure for $d(z,z^*)$ and $f(\hat{z})$ since specific tree-level forms are likely[6] to be destroyed by gravitational radiative corrections.

The passage[7] to energy scales much below $M_{P\ell}$ and just below M_{GUT} can be achieved by expanding in powers of the gravitino mass divided by $M_{P\ell}$ and integrating out the superheavy and hidden sector superfields. One then recovers the light superpotential f_ℓ of (1) plus the following soft SUSY-breaking terms:

$$\mathscr{L}_{SOFT} = 1/2 \sum_a M_a \bar{\lambda}_a \lambda_a - m_{3/2} \, (Af_\ell^{(3)} + Bf_\ell^{(2)} + \text{h.c.})$$

$$- m_{3/2}^2 \sum_\ell |z_\ell|^2 - 1/2 \, D_a D_a . \qquad (15)$$

In (15) $f_\ell^{(3)}$ and $f_\ell^{(2)}$ are those parts of f_ℓ which are respectively cubic and quadratic in the light sector scalar fields. Further, $m_{3/2}$ is a mass of the order of the gravitino mass, being scaled from it by a factor $\exp(\kappa^2 \langle d(z,z^*) - \sum_N z_N^* z^N \rangle)$; thus $m_{3/2}$ equals the gravitino mass exactly when the Kähler metric is flat. A and B are complex coefficients[7] of order unity which are given by complicated formulas involving the Kähler potential, the hidden sector superpotential, their derivatives etc. evaluated at the minimum of the potential (e.g. in the case of a flat Kähler metric $A = \kappa^2 m_{3/2}^{-1} \sum_h \langle z^h \rangle \langle F(z^h) \rangle$ and $B = 1$). Finally, M_a — the Majorana masses of the gaugino fields λ_a — are the eigenvalues of the matrix $\sum_h \langle \partial \mathrm{Re} f_{ab} / \partial z^h \rangle \langle F(z^h) \rangle$. It is reasonable to assume the equality of the 3,2,1 gaugino masses at the GUT scale $M_{1,2,3} = M$. This follows if we seriously take the unitary symmetry of transformations[6,7] $\hat{z} \to Uz$ among chiral superfields which supergravity interactions (ignoring gauge and Yukawa forces) respect. Then f_{ab} must have the form $(1+Y)\,\delta_{ab}$ with Y being a unitary singlet function of z^N such as $M_{P\ell}^{-3} d_{LMN} z^L z^M z^N$. Allowing for the spontaneous breakdown of this unitary symmetry in the hidden sector, we then have $M_{1,2,3} = \sum_h \langle \partial \mathrm{Re} Y / \partial z^h \rangle \langle F(z^h) \rangle$. If Y is of the order of $(z^h/M_{P\ell})^3$, then $M = O(m_{3/2})$.

We now have the Higgs potential of (2) but with $g_{1,2}$ and $\mu_{1,2,3}$ evaluated at M_{GUT}. While evolving downwards in mass-scale, $\hat{\alpha}_i \equiv (g_i/4\pi)^2$ and μ_i^2 become functions of $t = \ln M_{GUT}^2/Q^2$ where Q is the runnng scale so that the evolution from M_{GUT} to M_W is really from t = 0 to t = 67. Thus

$$\mu_1^2(0) = \mu_2^2(0) = m_{3/2}^2 + \mu^2, \qquad (16a)$$

$$\mu_3^2 = -2\mu B m_{3/2}, \qquad (16b)$$

$$\hat{\alpha}_3(0) = \hat{\alpha}_2(0) = \hat{\alpha}_1(0) = 1/24 \;, \qquad (16c)$$

where (16a,b) follow from (1) and (15). Similarly, the full scalar potential for the squark $(\tilde{Q})$-slepton $(\tilde{L})$ sector may be obtained from (1) and (15) with quadratic, cubic and quartic parts in these fields. For us it suffices to display only the first two parts:

$$V_2^{\tilde{Q},\tilde{L}} = m_Q^2 |\tilde{Q}|^2 + m_T^2 |\tilde{T}^C|^2 + m_B^2 |\tilde{B}^C|^2 + m_L^2 |\tilde{L}^C|^2 + m_E^2 |\tilde{E}^C|^2,$$

$$V_3^{\tilde{Q},\tilde{L}} = m_{3/2} \, (A_T \lambda_T \tilde{Q}\tilde{U}^C H_B + A_B \lambda_B \tilde{Q}\tilde{D}^C H_T + A_L \lambda_L \tilde{L}\tilde{E}^C H_T) + \mu \, (\lambda_T H_B^* \tilde{Q}\tilde{U}^C + \lambda_B H_T^* \tilde{Q}\tilde{D}^C + \lambda_L H_T^* \tilde{L}\tilde{E}^C) + \text{h.c.} \qquad (17)$$

In (17), m_ℓ^2, A_ℓ and λ_ℓ ($\equiv 4\pi\sqrt{Y_\ell}$) are functions of t and the latter two can be chosen to be real and positive by appropriately redefining phases and ignoring CP-violation. Further, (15) yields the boundary conditions

$$m_Q(0) = m_T(0) = m_B(0) = m_L(0) = m_{3/2},$$

$$A_T(0) = A_B(0) = A_L(0) = A. \qquad (18)$$

It has been shown[8] that A must be in the range $0 \le A \le 3$ for the charge and color invariance of the vacuum to be preserved, i.e. for A outside this range some squark or charged scalar develops a nonzero VEV.

The relevant 3-2-1 renormalizable group (RG) equations[9] are coupled nonlinear matrix differential equations that are not analytically integrable in the most general case. The system can be solved numerically but then much of the physics underlying the final result becomes obscure. It is therefore attractive to make the following reasonable approximations that make these equations analytically tractable. 1) With $h_{1,2}$ comparable in magnitude, all Y's scale according to the fermion $(\text{mass})^2$ so that only Y_T can be retained in the RG eqns. 2) We can drop μ from the equations although keeping μ_3. Arguments have been given[4] about the expected smallness of μ because of radiative origin while B could be quite large. Our conclusions should be insensitive to the value of μ unless it is large; but a large μ is unacceptable since that will impede SU(2)xU(1) braking. 3) The terms $\hat{\alpha}_1 M_1^2$ and $\hat{\alpha}_2 M_2^2$ are neglected in comparison with $\hat{\alpha}_3 M_3^2$. If the tree-level gaugino masses are all equal at M_{GUT}, then — since they scale like the $\hat{\alpha}$'s — for any significant value of t, the third term will totally dominate[9] over the first two. These approxiations enable us to write the following analytic relations (N.B. now we have included three generations of quarks and leptons):

$$\frac{\hat{\alpha}_a(t)}{\hat{\alpha}_a(0)} = \frac{M_a(t)}{M_a(0)} = \left(1 + \hat{\alpha}_a(0) b_a t\right)^{-1},$$

$$b_1 = \frac{33}{5},\; b_2 = 1,\; b_3 = -3 \;, \qquad (19a)$$

$$\mu_1^2(t) = \mu_1^2(0) = m_{3/2}^2 = \mu_2^2(0), \qquad (19b)$$

$$\mu_3^2(t) = \mu_3^2(0) \; \left(1 + \frac{t}{96\pi}\right)^{3/2} \left(1 + \frac{11t}{160\pi}\right)^{1/22} \left[1 + 6F(t)Y_T(0)\right]^{-1/4} , \qquad (19c)$$

$$Y_T(t) = Y_T(0)E(t)\left[1 + 6F(t)Y_T(0)\right]^{-1}. \qquad (19d)$$

Here $E(t) = (1+t/32\pi)^{16/9} \, (1-t/96\pi)^{-3} \, (1+11t/160\pi)^{13/99}$ and $F(t) = \int_0 dt' \, E(t')$ so that $E(67) \simeq 14.5$ and $F(67) \simeq 290$. However, the μ_2^2 equation cannot be integrated and $\mu_2^2(t)$ extracted unless M_3 is put equal to zero. We can nonetheless exploit the sign of the gluino mass term and derive even for a nonnegligible gluino mass the inequality

$$\mu_2^2(t) < (\mu_2^2(t))_{M_3=0} = 3/2 \, m_{3/2}^2 \Big(-1/3 + \{1+6F(t)Y_T(0)\}^{-2}\{1+6F(t)Y_T(0)(1-1/3A^2)\}\Big). \qquad (20)$$

3. Bounds

(19) and (20) clairfy the connection between the strength of the top Yukawa coupling $Y_T(0)$ — related via (19d) to $Y_T(67)$ and the top mass $m_T = 4\pi h_2 \sqrt{Y_T(67)}$ — and the existence of tighter bounds on the Higgs masses. Since $E(t)$ and $F(t)$ are monotonically increasing functions of t, a large $Y_T(0)$ makes μ_3^2 decrease significantly as t goes from zero to 67 and makes $\mu_2^2(67)$ end up being negative. In this situation the shaded region of the figure in the abstruct shrinks to virtually nothing and, with the stability line having shifted towards the origin, the cross ends up being to the left of the μ_1^2 axis. The square root of (7) cannot be taken in this case so that there is no upper bound on m_b tighter than M_Z. In practice, this occurs for m_T somewhat larger than 60 GeV. However, for a lower value of m_T — as suggested by the current UA1 experiment — a totally different situation obtains. Now μ_3^2 does not change much with t, the numerator and denominator in (19c) nearly balancing each other. μ_2^2 decreases as t goes from zero to 67 but not a lot; it is barely able to carry itself across the hyperbola into the shaded region — the thin end of the wedge!

Part of (4) can be rewritten as $1 = 8\, G_F (\sqrt{2})^{-1} \rho(\delta)\, h_2^2$ with $\rho(\delta) = (1+\delta)(1+\delta+\delta^2/2)^{-1}$ and h_2 can be eliminated by introducing $Y_T(0)$ and m_T. We then have

$$1740Y_T(0)\left(2.3\times10^4(1+\delta)^2(1+\delta+\delta^2/2)^{-1}\right.$$
$$\left.(M_{proton}/m_T)^2-1\right) = 1. \qquad (21)$$

Now (11), (19b), (20) and (21) together imply

$$\frac{\delta^2}{1+\delta}\,\frac{h(\delta,A)}{\left(1-h(\delta,A)\right)^2} < 1, \qquad (22a)$$

$$h(\delta,A) \equiv \left(3/2\ \left\{1+(1-A^2/3)\rho(\delta)\right\}\right.$$
$$\left.\left\{1+\rho(\delta)\right\}^{-2} - 1/2\right)^{1/2}. \qquad (22b)$$

Eqs. (22) impose a kind of bootstrap condition on δ as follows. For any A in the range $0 \le A \le 3$, there exists an upper bound $\delta_{max}(A)$ on δ such that, if δ exceeds $\delta_{max}(A)$, the inequality (22a) gets violated. This $\delta_{max}(A)$ is a monotonically increasing function of A so that $\delta_{max}(3)$ is an absolute upper bound on δ — leading via (10) to the upper bound

$$m_b \le M_Z\ \left(1-\left\{1+\delta_{max}(3) + 1/2\ \delta^2_{max}(3)\right\}^{-1}\right)^2. \qquad (23)$$

The values of $\delta_{max}(A)$ and correspondingly of $(m_b/M_Z)_{max}$ have been displayed against various values of A in the range 0-3 in the following table for a characteristic m_T of 40 GeV. The lower entry in the right most column varies from 0.09 to 0.41 as m_T goes from 20 GeV to 55 GeV. Thus, taking m_T < 55 GeV, we find that m_b < 38.3 GeV which is quite a bit less than 1/2 M_Z. The remarkable thing about this result is that it is independent of the precise values of $m_{3/2}$ and A. There are additional results depending on their values, e.g. $m^2_{3/2} + M_0^2 \le m_a^2 \le 2m^2_{3/2} + M_0^2$ and $m^2_{3/2} \le m_c^2 < 2m^2_{3/2}$. The lower bound on m_a can be moved up above M_Z in an $m_{3/2}$-dependent but A-independent way while some regions in the m_a^2, m_c^2 plot can be excluded in an $m_{3/2}$-independent but A-dependent way. The details are to be found in Ref. 5. An H° with a mass < 38.3 GeV should certainly be discoverable in the Bjorken process $Z \to H\mu\mu$ and also probably in the Wilczek process toponium $\to H\gamma$ both at SLC and LEP. A definite negative result on its existence would rule out this entire clas of theories.

TABLE

A	0.1	1.0	1.5	2.0	2.5	3.0
δ_{max}	.061	.081	.106	.143	.193	.258
$(m_b/M_Z)_{max}$	.059	.078	.100	.133	.175	.226

4. Acknowledgments

I thank Partha Majumdar for a fruitful collaboration and Riccardo Barbieri, Mike Dine, John Ellis, Howie Haber, Larry Hall, Joe Polchinski and Marc Sher for helpful discussions. I am indebted to Mary K. Gaillard and Ian Hinchliffe for their hospitality at this stimulating workshop and to Mike Creutz and Bill Marciano for the hospitality at Brookhaven.

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6. Postscript

I have just come across a preprint by H.P. Nilles and M. Nusbaumer (Université de Geneve Report No. UGVA-DPT 1984/05-431) who have independently arrived at somewhat similar conclusions.

THE PRODUCTION AND DECAY OF HEAVY GAUGE BOSONS IN $p\bar{p}$ AND pp COLLISIONS

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Summary

We have studied the prospects for producing and detecting massive (e.g. 1 TeV) new gauge bosons at pp and $p\bar{p}$ colliders, including charged W bosons with right-handed couplings and several types of neutral bosons. We especially emphasize that forward-backward asymmetries of decay leptons, which can occur in both pp and $p\bar{p}$ collisions, may be a very useful probe of the gauge boson couplings.

Conclusions

Let us start by giving our major conclusions. Further details may be found in ref. 1-3.

(a) At a $\sqrt{s}$ = 40 TeV pp collider with an integral luminosity L of $10^{40}cm^{-2}$ it should be possible to detect a right-handed charged W_R by its leptonic decays if $M_{W_R} \lesssim 8$ TeV. We assume that detection of a W_R^+ is possible if there are at least 10 events each of $pp \to W_R^+ \to e^+N$ and $pp \to W_R^+ \to \mu^+N$.

(b) Under similar assumptions the Z_χ (the additional neutral boson in SO_{10}) can be detected by its e^+e^- or $\mu^+\mu^-$ decays if $M_{Z_\chi} \lesssim 6$ TeV.

(c) For the purposes considered here a pp collider with $L = 10^{40}cm^{-2}$ is slightly better than a $p\bar{p}$ collider with the same energy and $L = 10^{39}cm^{-2}$.

(d) For a boson mass $M_B \lesssim 1$ TeV a useful diagnostic signature for both pp and $p\bar{p}$ collisions is the forward-backward asymmetry of the emitted lepton:

$$A_{FB}(y) = \left[\left(\frac{d\sigma}{dy}\right)_{z^*>0} - \left(\frac{d\sigma}{dy}\right)_{z^*<0}\right] \Big/ \frac{d\sigma}{dy} \quad , \tag{1}$$

where z^* is the cosine of the lepton angle with respect to the beam direction in the gauge boson rest frame and y is the gauge boson rapidity.

(e) For $M_B \lesssim$ 1-5 TeV global asymmetry variables are useful. For pp a promising variable is $\langle E_{\ell^-}\rangle/\langle E_{\ell^+}\rangle$, where $\langle E_\ell\rangle$ are the average lepton energies in $W'^{\pm} \to \ell^\pm X$ or $Z' \to \ell^+\ell^-$.

(f) The secondary decays $N \to e^\pm + X$, where N (the SU_{2R} partner of e_R^-) is produced in $W_R^\pm \to e^\pm \overset{(-)}{N}$ or $Z_\chi \to N\bar{N}$, may occur essentially instantaneously, a finite distance from the production vertex, or even outside the detector (depending on the model). Observation of such decays could be very useful both for reconstructing the W_R and for determining the nature (eg. Majorana or Dirac) of the N (see ref. 1).

(g) Lepton sign identification is extremely important.

Typical Bosons

An important candidate for an additional charged boson is the $W_R^\pm$ of $SU_{2L} \times SU_{2R} \times U_1$ models which couples to the right-handed doublets $(u\ d)_R$ and $(N_e\ e^-)_R$. The new neutrino N_e may be Dirac or Majorana, but we usually assume that it is lighter than the W_R. We assume $g_R \simeq g_L$ where g_L and g_R are the SU_{2L} and SU_{2R} gauge couplings (for example, $g_R^2 < 1.1\ g_L^2$ for $M_{W_R} < 1$ TeV if $SU_{2L} \times SU_{2R} \times U_1$ is embedded in SO_{10}, although that model gives an unacceptable value for $\sin^2\theta_W$). We allow M_{W_R} to be arbitrary, although plausible arguments[4] based on the K_L-K_S mass difference strongly suggest $M_{W_R} >$ (1-2) TeV (or > 5 TeV if QCD corrections are included).[5]

An important possible neutral boson is the Z_χ which occurs in $SO_{10} \to SU_5 \times U_1$ or under more general circumstances.[1] The normalized X charge is

$$X = \frac{1}{\sqrt{10}}\left[5\ I_{3R} + 3(I_{3L} - Q)\right] \quad , \tag{2}$$

so that Z_χ couples with strength

$$\begin{array}{llll} \frac{3}{2\sqrt{10}}\ g_\chi & \text{to} & \bar{d}_L, e_L^-, \nu_{eL} & (5^* \text{ of } SU_5) \\ \frac{-1}{2\sqrt{10}}\ g_\chi & \text{to} & d_L, u_L, \bar{u}_L, e_L^+ & (10 \text{ of } SU_5) \\ \frac{-5}{2\sqrt{10}}\ g_\chi & \text{to} & \bar{N}_L & (1 \text{ of } SU_5) \ . \end{array} \tag{3}$$

We see that the Z_χ couplings to d and e^- are mainly V + A and V - A, respectively, while the couplings to u are purely axial and relatively weak. We assume $g_\chi = \sqrt{\frac{5}{3}}\ g' = \sqrt{\frac{5}{3}}\ g_L \tan\theta_W$, which will occur if $U_{1\chi}$ and U_{1Y} split off from an underlying grand unified group at the same scale.

Another interesting neutral boson is the Z_ψ, which occurs in $E_6 \to SO_{10} \times U_{1\psi}$. The couplings of Z_ψ to all of the ordinary fermions $[(u, \bar{u}, d, \bar{d}, e^\pm, \nu_e, \bar{N}_e)_L$ and higher families, which occur in the 16 of $SO_{10}]$ are the same $(g_\psi/\sqrt{24})$, with larger couplings to exotic heavy fermions (which occur in the 10 and 1 of SO_{10}). Hence, the couplings to ordinary fermions are purely axial. We assume $g_\psi = g_\chi = \sqrt{\frac{5}{3}}\ g'$.

We neglect any mixing between the $W_R^\pm$ and $W_L^\pm$ or between the Z_χ or Z_ψ and the Z. It can be shown[2] under quite general conditions that such mixings are negligible for heavy ($M \gtrsim 1$ TeV) bosons because of the success of the $SU_2 \times U_1$ predictions for M_W and M_Z.

Other neutral bosons are discusssed in ref. 1.

Production

We assume the Drell-Yan cross section

$$\frac{d\sigma}{dy} = \frac{4\pi^2}{3M^3}\, x_1 x_2 \sum_{i,j} f_i^{\ A}(x_1)\ f_j^{\ B}(x_2)\Gamma_{ij} \tag{4}$$

for the production of a heavy gauge boson of mass M and rapidity y in an AB collision at energy $\sqrt{s}$. Here $x_1 = \frac{M}{\sqrt{s}}\, e^{+y}$, $x_2 = \frac{M}{\sqrt{s}}\, e^{-y}$, $f_i^{\ A}(x_1)$ is the distribution function for quark or antiquark i in hadron A, and Γ_{ij} is the partial width of the boson into $q_i\bar{q}_j$ or $\bar{q}_i q_j$.

We use the Q^2 dependent structure functions of Eichten, et al,[6] based on CDHS data and numerically extrapolated to large Q^2 using the Altarelli-Parisi equations with $\Lambda = 0.29$ GeV.

Total production cross sections and rapidity distributions for various bosons as a function of M and $\sqrt{s}$ are given in ref. 1 and 6. Here we limit ourselves to Table 1, in which is shown the maximum mass of W_R^+, W_R^-, and Z_χ attainable at specific cross section levels. For example, for $M_{R^+} = 8.6$ TeV and $\sqrt{s} = 40$ TeV, one has $B\sigma = 10^{-39}\text{cm}^2$ for pp (B is the branching ratio into $e^+\nu_e$). Hence, one expects 10 events each of $W_R^+ \to e^+\nu_e$ and $\mu^+\nu_\mu$ if $L = 10^{40}\text{cm}^{-2}$.

	$\sqrt{s}$ (TeV)	pp $B\sigma = 10^{-39}\text{cm}^2$	pp 10^{-38}cm^2	$\bar{p}p$ 10^{-38}cm^2	$\bar{p}p$ 10^{-37}cm^2
W_R^+	10	3.2	2.3	2.9	1.9
	20	5.3	3.8	4.4	2.7
	40	8.6	5.8	6.5	3.7
W_R^-	10	2.7	2.0	same	
	20	4.5	3.1	as	
	40	7.3	4.8	W_R^+	
Z_χ	10	2.2	1.5	1.9	1.1
	20	3.5	2.3	2.8	1.5
	40	5.6	3.3	3.8	2.0

Table 1. Maximum $W_R^\pm$ and Z_χ masses in TeV/c^2 attainable at specific cross section levels. B is the leptonic branching ratio into $W_R^+ \to e^\pm \overset{(-)}{N}_e$ (or $\mu^\pm \overset{(-)}{N}_\mu$) or $Z_\chi \to e^+e^-$ (or $\mu^+\mu^-$).

Signatures

Assuming that a heavy boson has been observed one would like to extract as much information as possible about its couplings. Total rates and branching ratios will probably be very difficult (we are assuming that heavy bosons can only be identified by their leptonic decay modes). More promising are the decays into heavy neutral leptons (N) or exotic fermions (see ref. 1).

Here we consider the characteristic forward-backward asymmetries of the emitted leptons in the boson rest frame. (Recall that such asymmetries were useful in establishing the spin and $|V| = |A|$ nature of the W boson).

In the $q_i\bar{q}_j$ center of mass, the differential cross section for $q_i\bar{q}_j \to \ell_i\bar{\ell}_j$ is proportional to

$$(L_q^2L_\ell^2 + R_q^2R_\ell^2) \quad (1 + z^*)^2$$
$$+ \quad (L_q^2R_\ell^2 + R_q^2L_\ell^2) \quad (1 - z^*)^2 \tag{5}$$

where L_q and R_q are the couplings of M to the left and right-handed quarks, respectively, and similarly for L_ℓ and R_ℓ. Clearly, a measurement of the forward-backward asymmetry can yield useful information concerning the quark and lepton couplings. In practice, one must multiply this expression by the appropriate parton distributions and sum over parton types. For $pp \to W^+_{L,R} + X$, with $W^+_{L,R} \to e^+\nu(N)$, for example, one has

$$A_{FB}(y) = \frac{3}{4}\ \frac{\bar{D}(x_1)U(x_2) - U(x_1)\bar{D}(x_2)}{\bar{D}(x_1)U(x_2) + U(x_1)\bar{D}(x_2)} \tag{6}$$

(note that $V \pm A$ lead to the same asymmetry). Although this vanishes at $y = 0$ (for pp), $A_{FB} \to -\frac{3}{4}$ for large y. This is because $x_1 >> x_2$ for large y, so that the scattering is mainly from a valence quark (sea antiquark) in the first (second) proton.

A_{FB} is shown for the ordinary Z and the Z_χ in Figures (1) and (2), respectively, for both pp and $p\bar{p}$. (The asymmetries vanish for Z_ψ which has purely axial couplings). Notice that the pp asymmetries quickly approach those for $p\bar{p}$ for $y > 0$. Also note that the asymmetries become smaller for smaller M or larger $\sqrt{s}$ (for which sea - sea collisions are more important). In particular, the Z^0 asymmetries are positive for $y > 0$ but very small (because the electron coupling becomes purely axial for $\sin^2\theta_W = \frac{1}{4}$), while the Z_χ asymmetries are negative and larger.

Clearly, $A_{FB}(y)$ would be a good diagnostic tool if adequate statistics are available to determine it. G. Gollin[7] has estimated that the Z_χ asymmetries could be reliably determined at a 40 TeV collider if $M_{Z_\chi} < 1$ TeV.

For larger masses, one must resort to some sort of y average to get enough events. For $p\bar{p}$ a useful variable would be the overall forward-backward asymmetry

$$A_{FB} = \int dy \left[\left(\frac{d\sigma}{dy}\right)_{z^*>0} - \left(\frac{d\sigma}{dy}\right)_{z^*<0}\right] \Big/ \sigma \tag{7}$$

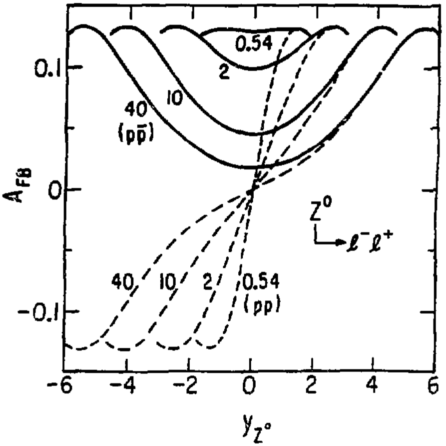

Figure 1. A_{FB} of ℓ^- vs. y_{Z^0} in pp (dashed line) and $p\bar{p}$ (solid lines) $\to Z^0 \to \ell^-\ell^+$. Curves are labeled by total c.m. energy in TeV.

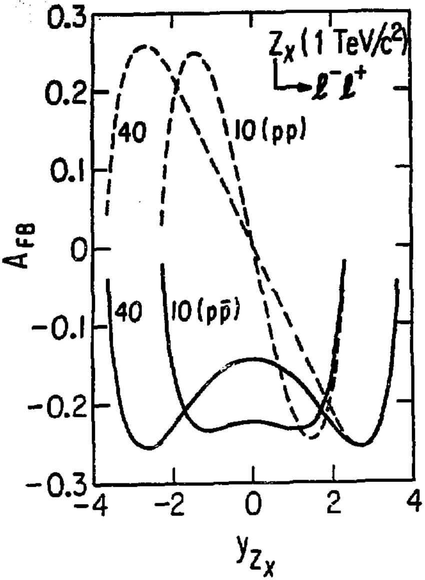

Figure 2. A_{FB} of ℓ^- vs. y_{Z_χ} for $M_{Z_\chi} = 1$ TeV/c^2, vs. y_{Z_χ}. (Other labels as in Figure 1).

For pp useful variables would be the average lepton energies $\langle E_{\ell^-}\rangle$ and $\langle E_{\ell^+}\rangle$. For $Z' \to \ell^+\ell^-$ any difference is due entirely to the FB asymmetries. For a comparison of $W_R^+ \to \ell^+$ vs. $W_R^- \to \ell^-$, part of the difference is due to different y distributions for $W_R^{\pm}$, but most of it is from the asymmetries. Predicted average energies for $W_R^{\pm}$ and Z_χ decays are shown in Figure 3. We have estimated that adequate statistics should be available to derive useful information from the global variables if the boson masses do not exceed (1-5) TeV.

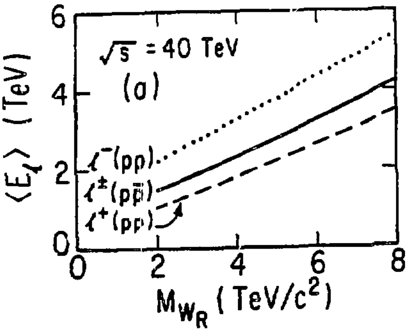

Figure 3a.

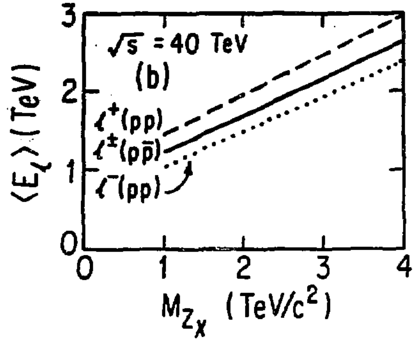

Figure 3b. Comparison of average lepton energies at $\sqrt{s} = 40$ TeV in $p\bar{p}$ collisions (solid lines) and pp collisions (ℓ^+; dashed lines; ℓ^-; dotted lines) for a) W_R decay, b) Z_χ decay.

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COMPOSITE HIGGS BOSONS

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I discuss the possibility that the Higgs boson may exist but be a bound state. This possibility is intermediate between technicolor models and models with a fundamental Higgs and it leads to a rich and interesting phenomenology at SSC energies.

I want to discuss the idea that the Higgs meson of the SU(2) x U(1) model may exist but be a composite state of strongly interacting fermions. Often, when *physicists hear this idea for the first time* their reaction is "Oh, you mean TECHNICOLOR." Well, I don't mean technicolor. What's more, I believe that this reaction is the result of a widespread muddling of the issues involved in SU(2) x U(1) symmetry breaking. My first goal will be to explain as carefully as I can what I think is going on. Then I will try to convince you that the composite Higgs idea is very rich and interesting and deserves some attention.

Part of the problem that I have in explaining the composite Higgs (CH) idea is that Peter Higgs' name has come to be associated with two completely different things: the Higgs mechanism; and the Higgs boson. The Higgs mechanism is the process in which a spontaneously broken gauge symmetry trades Goldstone bosons for gauge boson masses. This process is common to the simplest SU(2) x U(1) model and the most complicated technicolor models. It involves the Goldstone bosons and the gauge bosons and their interactions and nothing else. The Higgs mechanism depends on the fact that gauge (and sometimes global) symmetries have been spontaneously broken, producing Goldstone bosons with the appropriate properties, but it does not depend on any details of the symmetry breaking mechanism. The Higgs boson in the SU(2) x U(1) model, on the other hand, is a remnant of the symmetry breaking mechanism. It is there precisely when the SU(2) x U(1) symmetry is broken by the vacuum expectation value (VEV) of an SU(2) doublet scalar field, the Higgs doublet. The Higgs boson is the non-Goldstone part of the Higgs doublet.

Unitarity requires that some remnants of the spontaneous symmetry breaking mechanism appear at energies of a few TeV or less. But it does not tell us exactly where they will appear, or exactly what the remnants are. When the symmetry breaking mechanism involves other scalar multiplets than the doublet, the remnants of the symmetry breaking are more complicated than the single Higgs boson and should not be called by the same name (although they usually are!) When the symmetry breaking mechanism is dynamical, in a technicolor model, the remnants are not bosons at all, but instead are the techniquarks, with dynamical masses which violate the symmetry. If the technicolor is confining, these are bound into technihadrons, some of which may be scalars, but *none of them is the Higgs* boson. There is absolutely no sense in which any of them were ever part of an SU(2) x U(1) doublet with the Goldstone bosons. Indeed, there is no readon that I know of why there must be any massive scalar bound states at all. One can imagine, for example, a nonconfining technicolor interaction which breaks the chiral symmetry and binds the Goldstone bosons but does not produce any other bound states.

I hope that I have convinced you that the Higgs boson is an object with very definite properties and that it may or may not exist, depending on the structure of the SU(2) x U(1) symmetry breaking interactions. *If so, we can proceed to discuss what* I call *the composite* Higgs idea. What would it mean for the Higgs boson to exist, but be composite? Simply that in high energy scattering experiments we should see an object with the properties we would expect for the Higgs boson. In particular, at energies large compared to its mass and the gauge boson masses, it should behave like part of an SU(2) x U(1) doublet with the longitudinal polarization states of the heavy gauge bosons, which (at these high energies) are the Goldstone bosons in disguise. This implies, for example, that it has all the usual couplings to quarks and leptons. But at some even higher energy scale Λ, this Higgs boson must reveal itself as a bound state of strongly interacting particles, with the usual properties expected of a bound state, such as form factor effects and other nonrenormalizable interactions, strong interactions, etc.

Of course, if the Higgs boson is composite and it is in an SU(2) x U(1) multiplet with the Goldstone bosons, then the Goldstone bosons must be composite as well. Thus the heavy gauge bosons will exhibit strong *interaction properties* (at least for their *longitudinal* parts) at the same energy scale, Λ. Let us call the putative strong interaction which binds the Higgs and the Goldstone bosons ULTRACOLOR, to distinguish it from technicolor. This scenario is intermediate, in a sense, between the fundamental Higgs model and technicolor. The ultracolor dynamics does not directly break $SU(2) \times U(1)$, but it does bind the Higgs doublet whose VEV breaks the symmetry at a smaller scale. Thus this mechanism should perhaps be investigated simply because it represents another possibility. We might also expect that if the confinement scale is large, a model of this kind will have some of the advantages of a technicolor model without such severe flavor changing neutral current problems. But, as we will see, there are even better reasons to study composite Higgs models. When implemented in the most natural way, the idea is both rich and predictive.

If we said only that the Higgs was composite at *some large scale* Λ, *it would not gain us much information.* Nor would it be very plausible. In general, we would expect the ultracolor confining interactions to produce bound states with mass of order Λ, while our Higgs is by assumption much lighter. To go further, we must explain why the Higgs mass m_H is much smaller than Λ. We know of two principles which can keep spinless bosons light: supersymmetry and the Goldstone mechanism. I am alergic to supersymmetry, so I will discuss only the Goldstone mechanism.

It is obvious that the Higgs boson cannot be a Goldstone boson. True Goldstone bosons have *only* derivative interactions. They cannot have gauge interactions nor can they have a potential which leads to a VEV. Thus they are useless for our purposes. But the Higgs could be a pseudo-Goldstone boson in the following sense. It could be that the ultracolor interactions have, in the absence of any other interactions, a set of spontaneously broken global symmetries which would make the Higgs a Goldstone boson, but that these global symmetries (GS) are broken by the $SU(2) \times U(1)$ gauge interactions and other interactions weaker than ultracolor. These weaker interactions could then, perhaps, generate the $SU(2) \times U(1)$ breaking VEV for the composite Higgs field. Specifically, I will assume that ultracolor is an unbroken gauge dynamics, like QCD, that the particles which feel the ultracolor interaction are fermions, and that the spontaneously broken global symmetries are chiral. *We know from our* experience with QCD that composite pseudo-Goldstone bosons can actually be produced in this situation. The weaker interactions can be of three kinds: ultrafermion masses; nonrenormalizable interactions resulting from unspecified physics above the scale Λ which I will call extended ultracolor (EUC); and gauge interactions such as $SU(2) \times U(1)$.

The first prerequisite for a composite Higgs model is thus an ultracolor dynamics which need not break $SU(2) \times U(1)$ but which can produce pseudo-Goldstone bosons which transform like an $SU(2) \times U(1)$ doublet. A simple example of such a model is based on an analogy with QCD. Suppose that there are three Dirac ultrafermions which transform under a complex representation of the ultracolor gauge group. These ultrafermions are like the three light quarks in QCD. The composite pseudo-Goldstone bosons in QCD include the K meson, an isospin doublet with charges 1 and 0. Thus if we treat weak $SU(2)$ in our ultracolor model just like isospin in QCD, and weak $U(1)$ like hypercharge (appropriately normalized), the analog of the K meson will be a good candidate for the Higgs doublet. In this model, therefore, our three Dirac ultrafermions transform as a doublet plus a singlet under $SU(2) \times U(1)$.

For example, if we have a strong $SU(N)_{UC}$ gauge group and consider ultrafermions transforming under $SU(N)_{UC} \times SU(2) \times U(1) \times U(1)_A$ as:

$$Q_L = \begin{cases} (N,2)_{a,1+c} \\ (N,1)_{b,-2+d} \end{cases} \qquad Q_R = \begin{cases} (\bar{N},2)_{-a,1-c} \\ (\bar{N},1)_{-b,-2-d} \end{cases} \tag{1}$$

For the time being we will not gauge the $U(1)_A$. This theory possesses an approximate $SU(3) \times SU(3)$ symmetry which is spontaneously broken to $SU(3)$ by the ultrafermion condensate at the scale Λ. Note that the $U(1)_A$ breaks at this scale.

The pseudo-Goldstone bosons (PGB's) that appear are an $SU(2)$ triplet, a complex $SU(2)$ doublet, and an $SU(2)$ singlet--the ultrapions, -kaons, and -eta respectively. (When extra particles are added to render the $U(1)_A$ anomaly free and it is gauged, then the ultraeta will be eaten.) Note that the ultrakaons have the right quantum numbers to be the Higgs doublet. In the composite Higgs scenario the ultrafermion condensate is misaligned slightly with the $SU(2) \times U(1)$ preserving direction, which is equivalent to giving the Higgs a VEV $v \ll \Lambda$.

We can conveniently describe the Higgs potential in terms of a nonlinear sigma model. Let $\Sigma = \exp[2i\pi^a(x)T^a/f]$, where the π^a are the PGB fields and the T^a are the $SU(3)$ generators. The parameter f is analogous to $f_\pi = 93$ MeV in QCD, but here is of order Λ. The VEV of Σ is then the orientation of the condensate, in the basis we have chosen. In particular

$\langle\Sigma\rangle = I$ is $SU(2)\times U(1)$ preserving.

The alignment problem reduces to minimizing $V(\Sigma)$. Since we do not have bare ultrafermion masses, $V(\Sigma)$ will not contain any pieces transforming under the global $SU(3)\times SU(3)$ symmetry like $(3,\bar{3})$.[f1] The leading contribution to $V(\Sigma)$ in coupling constants squared and powers of $\Lambda^2/\Lambda^2_{EUC}$ will transform under $SU(3)\times SU(3)$ as an $(8,8)$. Thus the most general form for $V(\Sigma)$ consistent with $SU(2)\times U(1)$ is, to leading order:

$$V(\Sigma) = c_1 f^4\,\mathrm{Tr}\,\Sigma T_8 \Sigma^\dagger T_8 + c_2 f^4 \sum_{k=4}^{7} \mathrm{Tr}\,\Sigma T_\alpha \Sigma^\dagger T_\alpha + c_3 f^4 \sum_{a=1}^{3} \mathrm{Tr}\,\Sigma T_a \Sigma^\dagger T_a \tag{2}$$

c_1 and c_3 get contributions of $\mathcal{O}(g^2)$ from $SU(2)\times U(1)$ interactions, while c_2 must be due to nonrenormalizable interactions and be of $\mathcal{O}(\Lambda^2/\Lambda^2_{EUC})$. The c_2 term is the sole piece in $V(\Sigma)$ to give the ultraeta a mass, and must vanish if $U(1)_A$ is a good symmetry.

The Yukawa couplings which give rise to quark and lepton masses must be generated by the EUC interactions and are $\mathcal{O}(\Lambda^2/\Lambda^2_{EUC})$. The fact that the quarks and leptons we know are much lighter than the W suggests that $\Lambda^2/\Lambda^2_{EUC} \ll g$. Thus it is arguable that c_1 and c_3 are primarily due to $SU(2)\times U(1)$ interactions, in which case they can be calculated in terms of the $\pi^+-\pi^0$ mass difference in QCD due to electromagnetic interactions, and that c_2 is much smaller or zero. To see what this means for the self interaction of the K^0--i.e., the Higgs boson--we write

$$\Sigma = \exp \frac{2i}{f}\begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & K^0/2 \\ 0 & K^0/2 & 0 \end{pmatrix} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c & is \\ 0 & is & c \end{pmatrix} \tag{3}$$

where $c = \cos K_0/f$, $s = \sin K_0/f$. Then the terms in $V(\Sigma)$ of Eq. (2) become

$$\sum_{a=1}^{3} \mathrm{Tr}\,\Sigma T_a \Sigma^\dagger T_a = \frac{1}{4}\cos^2 K_0/f + \cos K_0/f \tag{4}$$

$$\sum_{a=4}^{7} \mathrm{Tr}\,\Sigma T_\alpha \Sigma^\dagger T_\alpha = \cos^2 K_0/f + \cos K_0/f$$

$$\mathrm{Tr}\,\Sigma T_8 \Sigma^\dagger T_8 = \frac{3}{4}\cos^2 K_0/f.$$

We see that $SU(2)\times U(1)$ gauge interactions cannot account for the whole Higgs potential--for c_1 and c_3 will have the same sign (negative [5]) and c_2 will be negligible. $V(\Sigma)$ is then minimized by the $SU(2)\times U(1)$ preserving vacuum with $\cos K_0/f \approx 1$.

Thus we follow the route taken in ref. [3] and gauge the $U(1)_A$, after cancelling its anomalies with spectators. When we do this, c_2 is forced to be strictly zero, and the ultra eta is eaten by the $U(1)_A$ gauge boson at the Λ_{UC} scale. More importantly, the $U(1)_A$ interactions give a <u>positive</u> contribution to c_1, and can cause the vacuum to misalign. It is now simple to write down the Higgs potential

$$V(K^0) = -\frac{c_3 f^4}{2\cos v/f}\left(\cos K_0/f - \cos v/f\right)^2 \tag{5}$$

(where $v = M_W\, 2\sin\theta/e \simeq 250$ GeV) where we used the fact that the linear $\cos K_0/f$ term comes solely from the c_3 piece of $V(\Sigma)$, as seen in (4). We may now compute the Higgs mass:

$$m^2_{K^0} = -(c_3 f^2 \sin^2 v/f)/\cos v/f = -c_3 v^2 [1 + 1/6(v/f)^2 - \ldots]. \tag{6}$$

We can relate c_3 to the $\pi^+-\pi^0$ mass splitting, which is primarily due to one photon exchange. This process is accounted for in the QCD chiral Lagrangian by the term

$$\mathscr{L} = \ldots + \frac{1}{2}\,\Delta m^2 f_\pi^2\,\mathrm{Tr}\,\Sigma Q \Sigma^\dagger Q \tag{7}$$

where Δm^2 is the observed $\pi^+-\pi^0$ mass difference to within about 5% ($\pi^0-\eta$ mixing and long distance effects are primarily responsible for this discrepancy.) Comparing this expression with Eq. (2) we may determine c_3 to be

$$c_3 = -\frac{\Delta m^2}{2\sin^2\theta f_\pi^2}\left[\frac{3}{N}\right] \tag{8}$$

where the factor $(3/N)$ comes from large N arguments.[f2] Using this result in Eq. (6) then gives us an expression for the Higgs mass:[6]

$$m_H^2 \simeq \left(\frac{m^2_{\pi^+} - m^2_{\pi^0}}{2\pi\alpha f_\pi^2}\right) M_W^2 (3/N) = [1.7\, M_W \sqrt{3/N}]^2,\quad N \gtrsim 3. \tag{9}$$

A crucial ingredient which went into the above calculation was the gauging of $U(1)_A$. In order for this to work, ordinary leptons and quarks must carry $U(1)_A$ charges which allow them to couple to the ultrakaon doublet. Furthermore, any additional spectators needed to cancel the $U(1)_A$ anomalies must be real under $SU(2)\times U(1)$ and have $U(1)_A$ charges which allow them to couple to the large $SU(2)\times U(1)$ preserving condensate, so that they acquire heavy masses, of order $\Lambda^3/\Lambda^2_{EUC}$. (If the anomalies were cancelled by particles which transform under $SU(N)_{UC}$, this would introduce an additional global $U(1)$ and a corresponding light $SU(2)\times U(1)$ singlet PGB). These additional fermions and the new $U(1)_A$ interactions of ordinary matter may be of experimental interest, depending on the model dependent ratios M_W/Λ and Λ/Λ_{EUC}.

There is no reason, in principle, why Λ cannot be very large. On the other hand, as Λ becomes much larger than the $SU(2)\times U(1)$ breaking scale v, the composite Higgs idea becomes less attractive for two reasons. The first is theoretical. The parameters of the GS breaking interactions must be tuned to an accuracy of $\sim v^2/\Lambda^2$. The second is practical. If Λ is too large, the differences between CH and FH models become very difficult to see at accessible energies. Thus the most attractive version of a CH model is one in which Λ is not much larger than v. We want to raise Λ as little as possible, consistent with FCNC constraints. If the EUC interactions are well chosen (as in the ETC models of ref. [7]), these constraints may not be very severe.

There is, however, another strong constraint on Λ in CH models like the model of references [3,6]. The trouble with this model is that it does not have a custodial $SU(2)_C$ symmetry. A custodial $SU(2)_C$ symmetry is an approximate global symmetry (broken by the U(1) gauge interactions) which enforces the relation[8]

$$M_W^2 = M_Z^2 \cos^2\theta. \qquad (10)$$

It is the existence of such a symmetry which ensures, *in appropriately constructed TC models, that the SU(2)x* U(1) breaking has the right form to agree with experiment.

A TC model without an $SU(2)_C$ symmetry is presumably just wrong, for although we do not know how to calculate the corrections to (10) in such theories, there is no reason to believe that they will be small. But in CH theories, the $SU(2)_C$ symmetry is not obviously necessary. When $\Lambda \gg v$, the theory is essentially a theory with an almost fundamental Higgs doublet, so (10) will be very nearly satisfied. But when Λ gets close to v, the compositeness effects will get more important and we may expect significant corrections to (10).

In the $SU(N)_{UC}$ model, in lowest order in chiral perturbation theory, the gauge boson masses come from the kinetic term

$$\frac{f^2}{4}\,\mathrm{tr}(D^\mu\Sigma D_\mu\Sigma^\dagger) \qquad (11)$$

where D^μ is the $SU(2)\times U(1)\times U(1)_A$ gauge covariant derivative,

$$D^\mu\Sigma = \partial^\mu\Sigma + i\, g_2[\vec{T}\cdot\vec{W}^\mu,\Sigma] \qquad (12)$$

$$+ i\,\frac{g_1}{\sqrt{3}}\,[T_8X^\mu,\Sigma] + i g_A Y^\mu(T_A^L\Sigma-\Sigma\, T_A^R)$$

where $\vec{W}^\mu(\vec{T})$, $X^\mu(T_8/\sqrt{3})$ and Y^μ (T_A^L and ^{R}) are the SU(2), U(1), and $U(1)_A$ gauge bosons (generators).

From (3), (11), and (12) we can read off the lowest order contributions to the weak gauge boson masses:[9]

$$M_Z^2\cos^2\theta = \frac{g_2^2 f^2 s^2}{8}$$
$$M_W^2 = \frac{g_2^2 f^2}{8}[s^2 + (1-c)^2]. \qquad (13)$$

(We have ignored Z-Y mixing--it will be absent if the coupling is pure axial, $T_A^L = -T_A^R$.[10] The second term in M_W^2 us clearly a correction to (10). It behaves like a VEV of a real SU(2) triplet Higgs multiplet, although here it arises from the nonrenormalizable terms in (4) required by the chiral symmetry. There will also be contributions from nonleading terms in chiral perturbation theory, but we expect these to be suppressed compared to (13). Thus for small s, we can replace (10) with

$$M_W^2 \simeq M_Z^2\cos^2\theta(1-v^2/4f^2). \qquad (14)$$

If we require that the ρ parameter in neutrino interaction be $1\pm.01$, we must have

$$f > 5v. \qquad (15)$$

It is not obvious that the constraint (15) is serious (since FCNC's must be suppressed anyway) but it is at least interesting to ask whether there are composite Higgs models which have an $SU(2)_C$ symmetry, and in which, therefore, the corrections to (10) are small. In fact, in ref. [2], we discussed such a model in which the ultraquarks transformed as 2N's of Sp(2N), one SU(2) doublet and two singlets. In this most economical model, the GS is SU(4) which is spontaneously broken down to Sp(4), leaving 5 PCB's, the 4 real components of the CH and a neutral singlet. Here there is an approximate global $SU(2)_G$ symmetry under which the two SU(2) singlet fields have a doublet. The Higgs transforms under $SU(2)\times SU(2)_G$ as a (2,2). The Higgs VEV breaks the symmetry down to the $SU(2)_C$ generated by the diagonal sum of weak SU(2) and $SU(2)_G$ *generators. Thus* (1) *is satisfied in lowest* order in chiral perturbation theory, and corrections are of order α.[8]

Unfortunately, in this model, the mechanism of refs. [3,6] for stabilizing the SU(2) invariant vacuum with the SU(2) gauge interactions does not work. An ultrafermion mass term is needed in order for the SU(2) breaking to be small compared to the UC confinement scale.

The SU(N) model of refs. [3,6] and the Sp(2N) model of ref. [2] are the simplest CH models in which the ultrafermion transform according to complex and pseudoreal representations respectively. The other simple possibility is to try a real representation. If the ultraquarks are real under UC, the GS breaking condensate is symmetric in flavor. This implies that there must be two SU(2) doublet pairs of ultrafermions in order to allow an SU(2) preserving condensate. To produce a composite Higgs, we need a singlet as well. Thus the simplest such model, there are five ultrafermions, two doublets and one singlet. For simplicity we will assume that there are N's under an UC SO(N).[9]

We organize the LH uq fields into a five component column vector,

$$\Psi = \begin{pmatrix}\psi_1\\ \psi_2\\ s\end{pmatrix}, \tag{16}$$

where ψ_j are the $SU(2)_L$ doublets. The GAS is SU(5) which is spontaneously broken down to an SO(5), producing 14 PGBs. It is convenient to describe the GS generators in a hybrid notation in which matrices which act on only the first four components of Ψ are written as tensor products of two sets of Pauli matrices, $\vec{\sigma}$, which acts on the ψ's and $\vec{\tau}$ which acts on the subscripts. Thus the SU(2) generators are $\vec{T} = \vec{\sigma}/2$ and the U(1) weak generator can be taken to be $S = -\tau_3/2$. The GS which leads to the custodial $SU(2)_C$ is $SU(2)_G$ generated by $\vec{T}_G = -\vec{\tau}^T/2$. If the $U(1)_A$ is chosen to commute with $SU(2)_G$, it is

$$T_A = \frac{1}{2\sqrt{5}}\begin{pmatrix}I & 0\\ 0 & -4\end{pmatrix} \tag{17}$$

where I stands for the 4x4 unit matrix. Note that $\mathrm{tr}\ T_A^2 = 1$.

Under $SU(2)\times SU(2)$, the s field is a singlet and the ψ's transform as a (2,2). Indeed the ψ's can be combined into a 2x2 matrix field

$$\psi = (\psi_1,\psi_2) \tag{18}$$

which behaves like a σ-model multiplet. Under an SU(2) transformation V_W and an $SU(2)_G$ transformation V_G,

$$\psi \rightarrow V_W\,\psi\,V_G^{\dagger}. \tag{19}$$

The composite Higgs, which is a bound state of ψ and s transforms the same way.

The breaking of the SU(5) GS down to SO(5) and the associated PGB's can be described by a symmetric unitary matrix

$$U = U^T, \quad U^{\dagger} = U^{-1}. \tag{20}$$

Under a SU(5) transformation V, $U \rightarrow V\,U\,V^T$. For convenience, we will define a unitary matrix Σ by

$$\Sigma = U\Delta, \Delta = \begin{bmatrix}\sigma_2\tau_2 & 0\\ 0 & 1\end{bmatrix}, \tag{21}$$

because $\Sigma=1$ corresponds to an SU(2) invariant vacuum. Σ satisfies

$$\Sigma = \Delta\Sigma^T\Delta = e^{2iX_\alpha\Pi_\alpha/f},$$
$$x_\alpha = \Delta X_\alpha^T\Delta, \quad \mathrm{tr}\ X_\alpha X_\beta = \delta_{\alpha\beta}. \tag{22}$$

This defines the PGB fields Π_α in terms of the SU(5) generators X_α which are broken in the SU(2) invariant vacuum (the unbroken generators satisfy $T_a = -\Delta T_a^T\Delta$), and the dimensional parameter f which is of order Λ.

We can write the PGB fields as follows

$$2\,X_\alpha\Pi_\alpha = \sigma_a\tau_b^T\rho_{ab} + 2T_A\xi + \begin{pmatrix}0 & \phi\\ \phi^{+} & 0\end{pmatrix}$$

where

$$\phi = \begin{pmatrix}\tilde{\phi}\\ \phi\end{pmatrix}, \ \tilde{\phi} = i\sigma_2\phi, \ (\tilde{\phi},\phi) = h + i\vec{\sigma}\cdot\vec{\pi}. \tag{24}$$

Evidently, ϕ is a Higgs doublet which transforms as a (2,2) under $SU(2)\times SU(2)_G$. h is the neutral Higgs field. ρ_{ab} transforms as a (3,3) and ξ is a singlet (eaten by the $U(1)_A$).

The kinetic energy term for the PGB is contained in the invariant

$$\frac{f^2}{8}\,\mathrm{tr}[(D^{\mu}\Sigma)(D_{\mu}\Sigma)^{+}] \tag{25}$$

where

$$\begin{aligned}D^{\mu}\Sigma &= \partial^{\mu}\Sigma + i\,g_2\,W^{\mu}[\vec{T},\Sigma]\\ &\quad + i\,g_1\,X^{\mu}[S,\Sigma] + ig_A Y^{\mu}\{T_A,\Sigma\}.\end{aligned} \tag{26}$$

In addition, as in the SU(N) model, SU(2), U(1), and $U(1)_A$ gauge bosons produce GS breaking nonderivative interactions which generate a potential for the PGB fields of the form (where y is a constant of order 1)

$$\begin{aligned}v(\pi) &= -y\,f^4 g_2^2\ \mathrm{tr}(\vec{T}\Sigma\vec{T}\Sigma^{+})\\ &\quad -y\,f^4 g_1^2\ \mathrm{tr}(S\Sigma S\Sigma^{+}) + f^4\,g_A^2\ \mathrm{tr}(T_A\Sigma T_A\Sigma^{+}).\end{aligned} \tag{27}$$

Note that the g_2^2 term stabilizes the $SU(2)\times U(1)$ invariant vacuum, giving positive mass-squared to the ρ_{ab} and ϕ while the g_A^2 term does not contribute to the ρ_{ab} mass, but destabilizes the invariant vacuum by giving a negative mass-squared term to the ϕ. Thus we expect to be able to find a range of parameters for which ρ_{ab} are heavy and have no VEV, while ϕ has a small negative mass-squared. For the potential to yield a useful CH model, it must also produce a positive $\lambda\phi^4$ coupling so that the ϕ VEV will be small. To see that this scenario, we can simply assume that $h \neq 0$ while all other PGB fields have zero VEV. Then

$$\Sigma = \begin{bmatrix} \frac{1+c}{2} & 0 & 0 & \frac{c-1}{2} & is/\sqrt{2} \\ 0 & 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 \\ \frac{c-1}{2} & 0 & 0 & \frac{1+c}{2} & is/\sqrt{2} \\ is/\sqrt{2} & 0 & 0 & is/\sqrt{2} & 1-2s^2 \end{bmatrix} \tag{28}$$

where

$$s = \sin(\sqrt{2}\ h/f),\ c = \cos(\sqrt{2}\ h/f). \tag{29}$$

The potential $v(\pi)$ can then be written as

$$v(\pi) = \frac{5}{4}\ y\ f^4 g_A^2 (\cos(h/f) - \cos\nu_0)^2 + \text{const}. \tag{30}$$

where

$$5\ g_A^2 \cos\nu_0 = 3g_2^2 + g_1^2. \tag{31}$$

If $5g_A^2 > 3g_2^2 + g_1^2$, $v(\pi)$ is minimized for $h = \nu_0 f$. Clearly for $5g_A^2 \approx 3g_2^2 + g_1^2$, the h VEV is small compared to f (and Λ). Thus the gauge interactions can generate an appropriate composite Higgs potential, just as in the SU(2) model.

Here, however, the $SU(2)_C$ symmetry is preserved because h is an $SU(2)_C$ singlet. From (18-23), we find

$$M_W^2 = M_Z^2 \cos^2\theta = g_2^2 f^2 (1-\cos\nu_0)/2 \tag{32}$$

which satisfies (10) as expected.

To calculate the Higgs mass in the $SO(N)_{UC}$ theory we need to know y in (27). Now we cannot relate this coefficient directly to the $\pi^+ - \pi^0$ mass difference because the SO(N) is a different strong gauge group. I don't really know how to do it, but I have a guess. The idea is to consider the SO(N) as a subgroup of SU(N) for which we know the answer. Then the SU(N) condensate looks just like a special alignment of the SO(N) alignment. My assumption is that the extra generators in SU(N) which are not in SO(N) are not *involved in an important way in the formation of the* condensate. Their job is simply to align it in the SU(N) preserving direction. This assumption allows us to normalize the f constants in the SO(N) condensate properly for comparison with an SU(N) theory with the same confinement scale. In fact, this is why we chose the normalization (22) of the generators in Σ.

If we now further assume that the effect of a weak gauge group on the condensate is the same in SU(N) (because the SU(N)/SO(N) gauge bosons are just aligning the condensate), we can calculate y. In SO(N),

$$y_N \approx \frac{1}{2}\ x_N, \tag{33}$$

$$x_N \approx \frac{c_3}{g_2^2} = -\frac{\Delta m^2}{2e^2 f_\pi^2} \cdot \frac{3}{N}. \tag{33}$$

One reason that I like this guess is that identical reasoning can be applied to the SU(N) subgroup of SU(2N) (under which the 2N transforms as N+N). It yields the familiar large N results,

$$\begin{aligned} f_N &\approx f_{2N}/\sqrt{2} \\ x_N &\approx \sqrt{2}\ x_{2N}. \end{aligned} \tag{34}$$

Note also that this guess can be checked in lattice gauge theory calculations.[11] At any rate, for what it is worth, (33) gives

$$m_H \approx 5.4\ M_W \sqrt{3/N}. \tag{35}$$

Several aspects of this type of composite Higgs model may be interesting from the point of view of SSC physics. The most obvious is that the Higgs mass is in an interesting range. If $\Lambda >> v$, this is really the only observable consequence. But for Λ not much larger than v, there are others. The axial U(1) gauge boson has mass $\approx g_A f$. In fact, this gives the strongest constraint on f in this model. The quarks and leptons must transform nontrivially under the $U(1)_A$ in order for the EUC interactions which produce their masses to be invariant. Thus $U(1)_A$ gauge boson exchange produces new neutral current interactions, weaker than SU(2) x U(1) neutral currents by a factor of $\approx v^2/f^2$. A *detailed analysis of neutral current constraint leads* to the bound

$$f \gtrsim 2v. \tag{36}$$

But still the $U(1)_A$ gauge boson could show up with a mass of few hundred GeV, in which case it would be copiously produced at the SSC and might even be observable at the FNAL collider.

Also, with masses of order gf are the other PGB's, which transform as 5, 3, and 1 under the custodial SU(2). The most spectacular of these are the doubly charged scalars in the 5. Unfortunately, in this theory the coupling of these scalars to $W^\pm W^\pm$ (which would show up at the SSC) is small because it is produced only by anomalies.

All the rest of the ultracolor physics is at still higher energies. The non-Goldstone ultraquark bound states should show up with masses of about $4\pi f$.[13] The EUC physics is at still higher scales and will probably best be studied in rare processes (ie. $K_L \to \mu e$).

Footnotes

1. In the models of reference [2] bare ultrafermion masses were needed to stabilize the vacuum. We

were dissatisfied with this mechanism as it seemed difficult to implement dynamically. We were pleased to learn from Tom Banks of the SU(3) x SU(3) model where ultrafermion masses were unnecessary.[3]

2. The order α_W contribution to the vacuum energy contians an ultrafermion loop and is hence of order N. Thus $c_3 f_2^4 = \mathcal{O}(N)$ which means that c_3 is $\mathcal{O}(1/N)$, because f_2^2 is O(N).

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ISOSPIN SYMMETRY BREAKING WITH A DYNAMICAL HIGGS SECTOR

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Abstract

The breaking of isospin symmetry in the electroweak theory with a dynamical Higgs sector is analyzed. An effective low energy theory is first constructed and used to discuss the natural size of the breaking in various amplitudes, given that the breaking is very large in the quark mass matrix. Special attention is paid to the parameter $\rho \equiv M_W^2/M_Z^2 \cos^2\theta_W$. Technicolor models are then investigated. There it is noted that $\rho - 1$ will naturally exhibit a sensitive linear dependence on the fermion-doublet mass splittings.

The observation of the charged and neutral weak bosons[1,2] reinforces our belief in the standard model of color-electroweak interactions. The final missing link is an understanding of the mechanism of electroweak symmetry breaking. It is attractive to assume that the breaking arises spontaneously, and within this framework two distinct possibilities emerge. One is that of the light, elementary Higgs field whose weak self-interactions are arranged to give it a non-zero vacuum expection value. The problems with this scheme have been stressed repeatedly. One shortcoming of special interest here is that as long as the Higgs fields are truly elementary, the model offers no understanding of the origin of fermion masses--they are merely parametrized by the Yukawa couplings of the Higgs field to the fermions.

The other possibility is that the spontaneous symmetry breakdown is due to the existence of new matter and strong forces at a mass scale around 1 TeV. An example are the technicolor models, in which the Higgs sector is a set of bound states and resonances formed from new, strongly interacting fermions.[3] Within this class of models, the masses of ordinary fermions must come from some direct interactions between the ordinary fermions and the new matter. In technicolor models, this must look like a direct four-fermion interaction or, at a deeper level, perhaps arise from the exchange of very massive ($\gtrsim$ 10 TeV) extended-technicolor bosons. While no completely realistic model of this sort has been constructed, it remains an attractive idea and one that at least offers the possibility of a deeper understanding of the origin of fermion masses.

In either of these cases, the observed strong isospin violation in the fermion mass matrix must be built into, if not explained by, the Lagrangian. It is important then to ask how this breaking infects other sectors of the theory through higher order corrections. In particular, one may expect corrections to the gauge boson mass matrix and to the prediction

$$\rho \equiv M_W^2/M_Z^2\cos^2\theta_W = 1 , \tag{1}$$

which is known to hold experimentally to within a few percent.[4] The relation $\rho = 1$ will be satisfied if the global symmetry of the Higgs sector is $SU(2)_L \times SU(2)_R$, which then spontaneously breaks to $SU(2)_{L+R}$.[5] The explicit breaking of the $SU(2)_R$, required to get the right fermion mass matrix, will necessarily give corrections to this relation. In this paper, we briefly review this problem in the elementary Higgs case and then discuss these corrections in the presence of a heavy, strongly interacting Higgs sector and, in particular, in technicolor models. There we find qualitatively new effects, perhaps at the edge of the experimental upper bounds.

In the elementary-Higgs model the SU(2) symmetry necessary to produce $\rho = 1$ arises as an accidental global symmetry of the potential of a complex scalar doublet. Writing this doublet as

$$\Phi = \begin{pmatrix} \phi^0 \\ \phi^- \end{pmatrix} = \begin{pmatrix} \sigma_0 + i\pi_3 \\ -\pi_1 + i\pi_2 \end{pmatrix} \tag{2}$$

we see that $V = V(\phi^+\phi)$ is $O(4) \approx SU(2) \times SU(2)$ symmetric. When $\langle\phi^0\rangle = \langle\sigma_0\rangle \neq 0$, this O(4) breaks to $O(3) \approx SU(2)$. The original O(4) symmetry is larger than the $SU(2)_L \times U(1)_Y$ symmetry needed to consistently couple the Higgs sector to the gauged weak interaction sector, and this extra symmetry is the origin of the relation $\rho = 1$.

It is clear that both gauge interactions and isospin breaking in the fermion mass matrix violate the SU(2) symmetry. Thus $\Delta\rho \equiv \rho - 1$ receives contributions of order α_W from radiative gauge-boson exchange diagrams.[6]

To see the consequence of isospin breaking in the fermion mass matrix we write out the Yukawa couplings for a generic family of fermions

$$\mathcal{L}_Y = Y_u\,(u_L^+ \; d_L^+)\begin{pmatrix} \phi^0 \\ \phi^- \end{pmatrix} u_R \; + \tag{3a}$$
$$Y_d\,(u_L^+ \; d_L^+)\begin{pmatrix} -\phi^{-*} \\ \phi^{0*} \end{pmatrix} d_R \; + \; h.c$$

Fig. 1. Fermion loop contribution to the gauge boson propagator.

where u denotes charge $^2/_3$ quarks and d charge $-^1/_3$ quarks. For $Y_u = Y_d = Y$ this takes the $SU(2)_L \times SU(2)_R$ symmetric form

$$Y(u_L^+ \; d_L^+)\begin{pmatrix}\phi^0 & -\phi^{-*}\\ \phi^- & \phi^{0*}\end{pmatrix}\begin{pmatrix}u_R\\ d_R\end{pmatrix} + \text{h.c} \tag{3b}$$

$$= Y q_L^+ M q_R + \text{h.c}$$

When $Y_u \neq Y_d$ the $SU(2)_L \times SU(2)_R$ symmetry is explicitly broken--the strength of the breaking should go smoothly to zero as Y_u approaches Y_d. In fact the contribution of fermion loops to the gauge-boson two-point function (Fig. 1) gives rise to [7]

$$\Delta\rho = \xi \frac{1}{16\pi}\left(\frac{\alpha_w}{M_w^2}\right)\left[\frac{2 m_u^2 m_d^2}{m_u^2 - m_d^2}\ln\frac{m_d^2}{m_u^2} + m_u^2 + m_d^2\right] \tag{4}$$

where ξ is a color factor— 1 for leptons and 3 for quarks. For small $m_u - m_d$ this reduces to

$$\Delta\rho = \xi\frac{\alpha_w}{12\pi}\left(\frac{m_u - m_d}{M_w}\right)^2 \tag{5}$$

The mass splitting $\Delta m = m_u - m_d$ appears quadratically since two mass insertions are required by helicity conservation. The experimental constraint

$$\Delta\rho \lesssim 0.05 \tag{6}$$

then requires $m_Q \lesssim 400$ GeV, for a heavy quark with a massless partner.

One may also derive (4) by considering solely Goldstone-boson dynamics.[8] By current conservation the gauge-boson two-point function has the transverse form

$$\Pi_{\mu\nu}(q^2) = \left(g_{\mu\nu} - \frac{q_\mu q_\nu}{q^2}\right)\Pi(q^2) \tag{7}$$

The $q_\mu q_\nu$ part is generated by the Goldstone-boson diagrams of Fig. 2 and yields the result of Eq. (4).

Fig. 2. Fermion loop contribution via the Goldstone boson propagator.

The arbitrary structure of the fermion mass matrix in the elementary-Higgs model thus allows very strong isospin violation without upsetting the relation $\rho \approx 1$.

If the Higgs sector is strongly interacting, the full theory at $E < 1$ TeV is conveniently described by the gauged nonlinear sigma model.[9] Corrections to the leading low energy behavior can then be summarized in the form of operators of higher and higher dimension which exhibit the $SU(2)_L \times SU(2)_R$ symmetry of the Higgs sector together with $SU(2)_R$ breaking effects. By listing and then estimating the natural size of these operators, the sensitivity of low energy measurements, such as the ρ parameter, to the 1 TeV dynamics and to the breaking of $SU(2)_R$ to $U(1)$ can be completely described.[10]

The coupling of the Goldstone fields to the gauge fields is given by

$$\mathcal{L}_{NL} = \frac{f^2}{4}\mathrm{Tr}\,(D_\mu U)^+(D^\mu U) \tag{8}$$

where $f \equiv \langle\sigma_0\rangle = 250$ GeV; where $U \equiv M/f$ obeys the nonlinear constraint $UU^+ = U^+U = 1$ and where $D_\mu U = \partial_\mu U + i\, {}^g/_2\, \vec{\tau}\cdot\vec{W}_\mu U - i\, {}^{g'}/_2\, B_\mu U\tau_3$. The fermion masses are again described by the Yukawa couplings $\mathcal{L}_y$ (Eq. 3). Among the new operators consistent with $SU(2)_L \times U(1)$ symmetry, there is only one of the same dimension as $\mathcal{L}_{NL}$ (dimension two with U dimensionless). It can be written in the form

$$\mathcal{L}_2 = a\frac{f^2}{4}\left[\mathrm{Tr}\,\tau_3 U^+ D_\mu U\right]^2 \tag{9}$$

where a is a dimensionless parameter.

The operator $\mathcal{L}_2$ is the only operator, in addition to $\mathcal{L}_{NL}$, that contributes to the gauge boson propagators at $q^2 = 0$ and, therefore, it completely determines $\Delta\rho$ as measured in low energy neutrino scattering experiments. A simple computation reveals that

$$\Delta\rho = -2a \tag{10}$$

A minimal size for a can be estimated by noting that $\mathcal{L}_2$ will be induced by radiative corrections involving the U(1) gauge field. It is found[10] that this source of $SU(2)_R$ breaking gives rise to $a = O\,(g^2/4\pi^2 \tan^2\theta_W)$, a value well within the experimental bound. The $SU(2)_R$ breaking in the Yukawa interactions (Eq. 3) will also induce $\mathcal{L}_2$. The lowest order computation is the same as in the elementary-Higgs theory and has already been described (Eqs. 4 and 5).

This analysis can be extended to higher order in the loop expansion, such as the contribution to $\mathcal{L}_1$ shown in Fig. 3, and to higher-dimension operators. The result, taking into account the fact that the *loop-expansion rapidly breaks down due to the strong* Higgs interactions, is the following: If there is no other source of $SU(2)_R$ symmetry breaking, beyond the Yukawa coupling $\mathcal{L}_Y$ and the U(1) gauge field coupling, then even in the presence of a strongly *interacting Higgs sector with a 1 TeV mass scale*, $\mathcal{L}_2$ and $\Delta\rho$ will not exceed the result (Eq. 5) which depends quadratically on the breaking parameter. This analysis will be presented in detail and extended to higher-dimension operators and other physical effects in a forthcoming paper.[11]

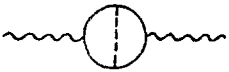

Fig. 3. Higher order contribution to $\Delta\rho$.

We now analyze the case where the Higgs sector is formed from (new) strongly interacting fermions—technifermions. The SU(2) isospin symmetry which guarantees $\rho = 1$ at the tree-level is the diagonal subgroup of the $SU(2)_L \times SU(2)_R$ chiral symmetry group of the technifermion sector.

Without further interactions the chiral symmetry of the ordinary fermions is unbroken—they are massless. The most economical way to give them mass is to couple them to technifermions via an effective four-fermi interaction of the form

$$\mathcal{L}_{fT} = \frac{G_1}{2}\bar{f}_L(\bar{T}T)f_R + \frac{G_2}{2}\bar{f}_L(\bar{T}T)\tau_3 f_R + h.c \tag{11}$$

The technifermions form a bound state at some scale Λ_{TC}

$$\langle \bar{T}_L T_R \rangle \sim \Lambda_{TC}^3 \tag{12}$$

Since isospin violation in the strong technicolor interaction would spread directly to the ρ parameter

$$\rho - 1 = \mathcal{O}(\alpha_{TC}) = \mathcal{O}(1) , \tag{13}$$

we assume that the only source of isospin violation, other than the $U(1)_Y$ gauge field coupling, is the coupling $\mathcal{L}_{fT}$. Then, a simple one-loop estimate gives

$$\Delta m = m_u - m_d \cong \frac{G_2}{8\pi^2}\Lambda_{TC}^3 \tag{14}$$

We now describe how this isospin violation affects the operator $\mathcal{L}_2$ and the shift of the ρ parameter. The first observation is that there will naturally exist a variety of other four-fermion interactions with coupling strengths on the order of G_1 and G_2. These will, for example, be induced by iterations of the interaction $\mathcal{L}_{fT}$. If the integrations are not cut off at energies below the unitarity bound, these new interactions will naturally be of the same strength as $\mathcal{L}_{fT}$. It might, of course, be that the iterations are damped below the unitarity bound. In extended technicolor (ETC) models, for example, $\mathcal{L}_{fT}$ could be due to an ETC boson exchange with a small dimensionless coupling constant. The higher order iterations will then be small. However, it is natural in these models for the additional four-fermion interactions to arise from the tree-level exchange of an ETC boson. Unless some suppression mechanism is present, this will be of strength comparable to $\mathcal{L}_{fT}$. Among the four-fermion interactions to be expected from these considerations are the following, involving only right-handed fermions:

$$\mathcal{L}'_{fT} = G'_2\,\bar{f}_R\gamma^\mu\tau_3 T_R\,\bar{T}_R\gamma_\mu\tau_3 f_R + h.c \tag{15}$$

$$\mathcal{L}_{TT} = G_3\,\bar{T}_R\gamma^\mu\tau_3 T_R\,\bar{T}_R\gamma_\mu\tau_3 T_R \tag{16}$$

These have been selected simply because they contain the "largest amount" of $SU(2)_R$ symmetry breaking, and they contribute most directly to $\Delta\rho$.

In order to estimate the contribution of either $\mathcal{L}_{fT}'$ or $\mathcal{L}_{TT}$ to $\Delta\rho$, it is simplest to imagine turning off the U(1) gauge coupling g'. In this limit, $\theta_W = 0$ and ρ is simply M_W^2/M_Z^2. Turning g' back on will then give higher order corrections to the dominant contributions.

Consider first the interaction $\mathcal{L}_{fT}'$. It can contribute to the W and Z masses through the graph shown in Fig. 4. The fact that $\mathcal{L}_{fT}'$ involves only

Fig. 4. The contributions of $\mathcal{L}_{fT}'$ to $\Delta\rho$.

right handed fermions allows the presence of the two τ_3's in the trace around the figure-eight loop and therefore a non-vanishing contribution to $\rho = M_W/M_Z$. (Recall that we are here defining mass and the ρ parameter as the zero-momentum limit of the inverse propagators.) The problem here is that, with $g' = 0$, the right-handed fermions do not couple to the gauge bosons without the helicity flip provided by mass insertions on the fermion lines. With the necessary insertions on the light fermion lines, an extra suppression factor of m_f^2/Λ_{TC}^2 will be introduced, giving a contribution to $\Delta\rho$ even smaller than that in the one-loop estimate (Eq. 5).

It is the operator $\mathcal{L}_{TT}$ that gives the largest contribution to $\Delta\rho$. Since it is the product of two $SU(2)_R$ violating currents, the figure-eight trace in Fig. 5 will contribute directly to $\Delta\rho$. Furthermore, since the technifermion mass is expected to be of order Λ_{TC} ($\simeq$ 1 TeV), no price is paid to convert the right-handed technifermion to the left-handed one necessary to couple to the gauge fields.

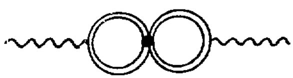

Fig. 5. The contributions of $\mathcal{L}_{TT}$ to $\Delta\rho$.

Because of strong technicolor interactions, it is only possible to estimate the contribution of Fig. 5. We find

$$\Delta\rho = \frac{1}{8\pi^2}\left(\frac{G_3}{8\pi^2}\Lambda_{TC}^2\right) \tag{17}$$

where the Λ_{TC} factors arise since the electroweak spontaneous symmetry breaking is restored for $q^2 > \Lambda_{TC}^2$. (In the nonlinear sigma model description of the low energy (< 1 TeV) theory, Fig. 5 can be regarded as giving rise to the operator $\mathcal{L}_2$ (Eq. 9).) To see the significance of this result, we rewrite it in terms of mass splittings. Using Eq. (14),

$$\Delta\rho = \frac{1}{8\pi^2}\frac{G_3}{G_2}\left(\frac{m_t - m_b}{\Lambda_{TC}}\right) = \frac{1}{8\pi^2}\frac{G_3}{G_2}\, g_w\left(\frac{\Delta m}{M_w}\right) \tag{18}$$

This is linear in the ordinary fermion mass splittings and the weak gauge coupling. The origin of the linear dependence on mass splittings we have found is the new scale $G^{-1/2}$ associated with the four-Fermi interactions that generate mass.

The numerical factors in this expression are only rough estimates. Nevertheless, it is perhaps worthwhile to compare its size to the result (Eq. (4)) and to the experimental bound. For $m_t \gg m_b$ and for $\xi = 3$, the expression in Eq. (4) becomes $(3/16\pi^2)(\Delta m^2/f^2)$. Clearly for small enough Δm, the linear expression (Eq. (18)) will dominate the quadratic one. If the numerical factor in Eq. (18) can be trusted, the two expressions become comparable when Δm is of order f, corresponding to a $\Delta\rho$ on the order of a few percent. This lies just within the current experimental bound. If this bound can be reduced, then the linear expression (Eq. (18)) will play the dominant role in constraining the mass splitting Δm.

Although the linear dependence is only obtained indirectly, through the effect of the four technifermion operator (Eq. (16)), we have argued that it is likely to be a generic property of techni-color theories. It is, in fact, not unlikely that this kind of constraint is a general feature of any theory in which electroweak symmetry breaking is due to some new matter with a mass scale of a few TeV. There must always be some $SU(2)_R$ violating interactions to give the correct fermion mass matrix and these could well feed back into $\Delta\rho$ in the manner discussed here.

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THE QCD PARTON MODEL AT COLLIDER ENERGIES

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Summary

Using the example of vector boson production, the application of the QCD improved parton model at collider energies is reviewed. The reliability of the extrapolation to SSC energies is assessed. Predictions at $\sqrt{S}$ = 0.54 TeV are compared with data.

Predictions for the interactions of hadrons in the TeV range are usually made using the parton model, suitably modified to include the effects due to QCD[1]. The model has been remarkably successful in analysis of experiments at fixed target energies, but present colliders test the model in a new energy regime, which will be further extended by the projected super colliders. This extension of the kinematic range raises certain theoretical issues which are addressed here, and elsewhere in these proceedings[2]. It is also of interest to compare the predictions of the model with data at $\sqrt{S}$ = 0.54 TeV, in order to assess the accuracy of projections to super-collider energies. This program is carried out in this paper.

Schematically, the parton model cross-section may be written as

$$\sigma(P_i) = \sum_{j,k}\int dx_1 dx_2\, f_j(x_1,Q^2)\, f_k(x_2,Q^2)\, \sigma_{jk}(x_i) \qquad (1)$$

where f_j are the parton distributions and j,k run over parton species. The QCD parton model contains three ingredients. These are,

a) the specification of distributions of quarks, antiquarks and gluons inside the colliding hadrons.

b) the extrapolation of the parton distributions to the higher energies relevant for collider experiments.

c) the calculations of parton cross-sections which, when combined with the parton distributions, fix the overall hadronic cross-section.

The first topic, the measurement of the parton densities will only be mentioned briefly. The principal source of information on these distributions comes from deep-inelastic lepton hadron scattering. For a review of the experimental problems in these determinations we refer the reader to ref.(3). The shape of the valence quark distributions is well determined. The uncertainties in the measurement of the antiquark distributions are somewhat larger, but the distributions themselves are smaller at fixed target energies. The shape of the gluon distribution, which is determined from scaling violations in deep-inelastic scattering, is correlated with the measured value of the scale breaking parameter Λ.

Setting aside the question of the experimental determination of the parton distributions, we now discuss the extrapolation to collider energies. In general the parton distribution functions are required at values of x and Q^2 which are outside the range measured in deep-inelastic scattering. The particular values depend on the transverse energy or mass of the object being produced. A W-boson produced in proton anti-proton collisions at $\sqrt{S}$ = 0.54 TeV is most likely to have come from a pair of partons having a fraction x = 0.15 of the hadrons' longitudinal momentum. Values of x which are higher or lower are probed if the W is produced in the forward or backward direction. At $\sqrt{S}$ = 40 TeV the typical value of x has become x = 2.0×10^{-3}, although in the measurable rapidity range, one is sensitive to values as small as x = 10^{-4}. For the production of hypothetical heavier particles, say of mass Q, the values of x are larger but the values of Q^2, at which the distribution is needed are also larger. We are therefore interested in a range such that,

$$Q/x < E \qquad (2)$$

where E is the total centre of mass energy of the collider.

The extrapolation to the values of x and Q^2 required is performed using the Altarelli-Parisi equation.[4]

$$\frac{d}{d(\ln Q^2)}\begin{pmatrix} q(x) \\ G(x)\end{pmatrix} = \frac{\alpha_s(Q^2)}{2\pi}\int_x^1 \frac{dz}{z}\begin{pmatrix} P_{qq}(z) & P_{qG}(z) \\ P_{Gq}(z) & P_{GG}(z)\end{pmatrix}\begin{pmatrix} q(\frac{x}{z}) \\ G(\frac{x}{z})\end{pmatrix} \qquad (3)$$

The functions P are the evolution kernels which are calculated as a perturbation series in the strong coupling constant. Normally the equations are used including only the first order evolution kernel, although the second order terms[5] and certain terms of even higher orders have also been calculated.[6] As the evolution proceeds uncertainties in the sea and gluon distribution functions tend to diminish. This is shown in Fig. (1) for the case of the gluon distribution function at Q^2 = 4 GeV^2 and Q^2 = 2000 GeV^2. The curves which are very different at low Q^2 approach one another at high Q^2. These curves were obtained using the two parameterisations of Duke and Owens[7] which evolve with different values of Λ. Part of the reason why different starting distributions, (compatible with data), give similar results after evolution is that Eq. (3) is driven by the hardest term on the right hand side, which is the well measured valence distribution.

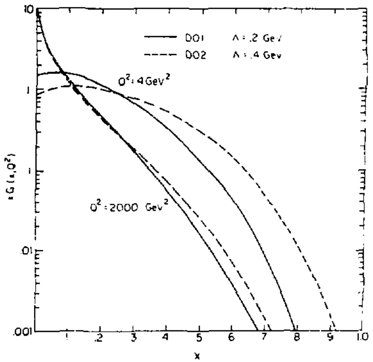

Fig. 1
Two parametrisations for the gluon distribution function at Q^2 = 4 GeV^2 and Q^2 = 2,000 GeV^2.

The extrapolation using the first-order Altarelli-Parisi kernels is expected to be acceptable throughout the range explored at super-collider energies. A possible source of danger is the low x region, untested by fixed target experiments. As already mentioned above, despite our ignorance of these distributions at low x and Q^2, the AP equations are expected to give a reliable estimate at low x and higher Q^2. This is because the growth at low x, due to parton cascade from higher x, is so much larger than the presumed starting value at low x. The issue is whether the AP equations with first order kernels are an accurate representation of the behaviour of the theory in this region. The one loop evolution equations at low x are dominated by the poles at x = 0 which appear in the splitting functions. In the limit $x \to 0, (C_A=3,\ C_F=4/3)$,

$$P_{GG}^{(1)}(x) \simeq \frac{\alpha_s}{2\pi}\frac{2C_A}{x} \quad ; \quad P_{Gq}^{(1)} \simeq \frac{\alpha_s}{2\pi}\frac{2C_F}{x} \tag{4}$$

In this approximation, the gluon distribution function evolves according to

$$\frac{dG(x,Q^2)}{d\ln Q^2} = \frac{C_A}{x}\frac{\alpha_s(Q^2)}{\pi}\int_x^1 dz\ G(z,Q^2) \tag{5}$$

The solution to this equation in the limit in which $\ln(1/x)\ln(\ln Q^2) >> 1$ is,[8]

$$G(x,Q^2) = \frac{1}{x}\exp\sqrt{\frac{4C_A}{\pi b}\ln\left(\ln\frac{Q^2}{\Lambda^2}\right)\ln\frac{1}{x}} \tag{6}$$

where $[b\alpha_s(Q^2)]^{-1} = \ln Q^2/\Lambda^2$. The second order splitting function does not lead to a large modification of this behaviour; at small x the matrix of evolution kernels is given by,[5]

$$xP^{(2)} = \left(\frac{\alpha_s}{2\pi}\right)^2 \begin{pmatrix} \frac{40}{9}C_F T_R n_f & \frac{40}{9}C_A T_R n_f \\ & \\ C_F C_A - \frac{40}{9}C_F T_R n_f & \frac{4}{3}(C_F - \frac{C_A}{2}) - \frac{4C}{,} C_A T_R n_f \end{pmatrix} \tag{7}$$

This equation should be compared with the corresponding results for the timelike case. For example, the function which controls the fragmentation of a gluon is given by,

$$P_{GG}^T \underset{x\to o}{=} \frac{\alpha_s}{2\pi}\left(\frac{2C_A}{x}\right) - \left(\frac{\alpha_s}{2\pi}\right)^2 \frac{4C_A^2\ln^2 x}{x} + \ldots \tag{8}$$

and after resummation to all orders the moments of this function are known to be given by,[9]

$$\gamma_{GG}^T(n) = \frac{1}{4}\left[-(n-1) + \sqrt{(n-1)^2 + \frac{8\alpha_s C_A}{\pi}}\ \right] \tag{9}$$

Returning to the spacelike case we see from Eq. (7) that terms of order $\alpha_s^2\ln(1/x)^m/x$ for m=1,2 are absent. Indeed it is known that to all orders the most singular terms in the perturbation series for the splitting function are of the form,[10]

$$P_{GG}(x) = \sum_{j=0}^{\infty} a_j\frac{\kappa}{x}\left(\kappa\ln\frac{1}{x}\right)^j \quad ; \quad \kappa = \left(\frac{\alpha_s C_A}{\pi}\right) \tag{10}$$

The values of the coefficients a_j are known.[6] Note that $a_1=a_2=0$. Since the correction terms are of order $\kappa\ln(1/x)$, we should not envisage any problems with perturbation theory until $\kappa\ln(1/x) \simeq 1$. Thus the first order equations provide an adequate description at least down to values,

$$x > 10^{-3} \text{ at } Q^2 = 10^4\text{GeV} \tag{11}$$

In ref. (2) it is argued that lowest order perturbation theory should be valid to even smaller values of x, because of the steepness near x = 0, with which the splitting function is convoluted. However, Eq. (11) is sufficient for most purposes at energies E < 40 TeV.

In order to make numerical estimates of the cross-sections we will use the results of numerical integration the of Altarelli-Parisi equation given in the literature. The parameterisations which we consider are those of Duke and Owens (DO),[7] Gluck Hoffmann and Reya (GHR)[11] and Eichten, Hinchliffe, Lane and Quigg (EHLQ).[12] None of the parameterisations is entirely satisfactory throughout the range $\sqrt{S}$=0.5-40 TeV. A satisfactory parameterisation must
a) be compatible with the data at fixed target energies,
b) give a satisfactory fit to the result of numerical evolution of the low energy distributions throughout the range of collider and super-collider energies (cf.eq.(11)). The stated range of accuracy of the three sets is

DO:	$5.10^{-3} < x < 1$	$2 < Q < 10^3$ GeV	(few %)	
GHR:	$10^{-2} < x < 1$	$2 < Q < 200$ GeV		
EHLQ:	$10^{-4} < x < 1$	$2.3 < Q < 10^4$ GeV	(5%)	(12)

where the percentage is the estimated maximum deviation of the parameterisation from the result of the numerical evolution of the starting distributions. Thus we see that the first two sets have an x range somewhat less than desired for super-collider energies.

Not all features of fixed target data are reproduced by the parameterisations, although there is some degree of choice in the data sets which are used. The DO parameterisations have an SU(3) symmetric sea which appears to be excluded by the data.[3] Since sea distributions are important at super-collider energies, this deficiency can lead to noticeable differences. The ratio of valence down and up quarks is measured to be approximately given by

$$d_v(x)/u_v(x) = 0.57(1-x) \tag{13}$$

The EHLQ structure functions fit this ratio rather poorly (see ref.(12)) and hence somewhat underestimate W production cross-sections at CERN collider energies. Different theoretical treatments of the charm quark threshold can be lead to appreciable differences at small values of x. Generally speaking these incompatibilities of the parton distribution functions with data lead to less than 20% effects in the final cross-sections, nevertheless they introduce an avoidable source of error.

The total cross-sections for vector boson production in $p\bar{p}$ collisions at CERN collider energies including the $O(\alpha_s)$ corrections have been presented in ref. (13). The gluonic radiative corrections were implemented following the basic strategy of ref. (14). Inclusion of the $O(\alpha_s)$ corrections increases the zero order cross-section - the so-called K factor - by about 30%. This is to be compared with the $O(\alpha_s)$ correction in Drell-Yan production at fixed target energies which is about 80%.[14] This decrease in the size of the radiative correction is mainly due to the decrease in the size of the running coupling α_s. The contribution of the initial gluons after factorisation is negative and less than a 5% correction.

The theoretical calculations of the cross-sections for $p\bar{p}$ collisions at $\sqrt{S}$ = 0.54 TeV are,[13]

$$\sigma^{W^+ + W^-} = \left(4.2 \begin{array}{l} +1.3 \\ -0.6 \end{array}\right)\text{nb} \quad ; \quad \sigma^{Z^0} = \left(1.3 \begin{array}{l} +0.4 \\ -0.2 \end{array}\right)\text{nb} \qquad (14)$$

The theoretical uncertainties in these cross-sections have been estimated by using different sets of parton distributions and different arguments for the running coupling. The value for the W cross-section found using the EHLQ structure function is somewhat low but lies within the range given in Eq. (14). The ratio of the two cross-sections, important for counting neutrinos is less subject to theoretical error,

$$\frac{\sigma^{W^+ + W^-}}{\sigma^{Z}} = 3.3 \pm 0.2 \qquad (15)$$

Multiplying Eqs. (14) by the branching ratio into electrons,

$$B(W \to e\nu) = 0.089 \qquad B(Z^0 \to e^+e^-) = 0.032 \qquad (16)$$

which are the values obtained for a top quark mass m_t = 40 GeV and α_s/π = 0.04, we find that the product of the cross-section and decay branching ratio is,

$$(\sigma B)^{W^\pm \to e^\pm} = \left(370 \begin{array}{l} +110 \\ -60 \end{array}\right)\text{pb}$$

$$(\sigma B)^{Z^0 \to e^+e^-} = \left(42 \begin{array}{l} +12 \\ -6 \end{array}\right)\text{pb} \qquad (17)$$

The corresponding experimental results are[15,16]

$$\text{UA1} : (\sigma B)^{W^\pm} = 530 \pm 80 \pm 90\text{pb} \qquad (\sigma B)^{Z^0} = 71 \pm 24 \pm 13\text{pb} \qquad (18)$$

$$\text{UA2} : (\sigma B)^{W^\pm} = 530 \pm 100 \pm 100\text{pb} \qquad (\sigma B)^{Z^0} = 110 \pm 40 \pm 20\text{pb} \qquad (19)$$

Theoretical predictions for higher energies are given in Table (1). These results are also subject to theoretical error. Fig. (2) displays these results for a fixed set of parton distribution functions (Duke and Owens[7], Set 1) and a given choice of scale for $\alpha_s(Q=M_W)$. The solid curve is for proton-antiproton and the dotted curve is for proton-proton collisions. Above $\sqrt{S}$ = 10 TeV the two curves are essentially identical because of the dominance of sea quarks. Also shown plotted are the cross-sections for the production of hypothetical bosons of mass 0.2, 0.5 and 1 TeV which couple to quarks exactly in the same way as the normal W boson. These curves are also subject to theoretical uncertainties similar to those in Table 1. Although the cross-section for the production of W bosons at 40 TeV is large, it should be borne in mind that only about 30% of them occur at observable rapidities y < 2. A W produced at rapidity greater than 2 lies within 15° of the beam pipe.

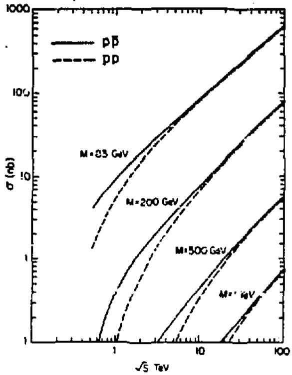

Fig. 2

The total cross-section for the production of $W^+ + W^-$ bosons, M = 83 GeV in proton antiproton collisions (solid line) and proton proton collisions (dashed line). The other curves refer to heavier charged bosons with the same couplings to quarks as the W of the standard model.

$\sqrt{S}$(TeV)	$\sigma^{W^+ + W^-}$ (M_W=83 GeV)(nb)	σ^{Z^0} (M_{Z^0}=94 GeV)(nb)
0.54	$4.2^{+1.3}_{-0.6}$	$1.3^{+0.4}_{-0.2}$
0.63	$5.3^{+1.6}_{-0.9}$	$1.6^{+0.5}_{-0.3}$
1.6	$16.0^{+4.0}_{-2.5}$	$4.9^{+1.2}_{-0.8}$
2.	$20.^{+6.}_{-4.}$	$6.2^{+1.9}_{-1.2}$
10.	$75.^{+35.}_{-25.}$	$27.^{+12.}_{-9.}$
20.	$130.^{+70.}_{-55.}$	$46.^{+24.}_{-20.}$
40.	190.±100.	70. ±30.

Table 1

Theoretical results for the W and Z total cross-sections in $p\bar{p}$ interactions at various energies. Estimates of the theoretical error are also given.

We now consider the transverse momentum of the produced vector bosons in more detail. This is a subject of both theoretical and practical importance. They are theoretically important because it has been shown that essentially the whole q_T distribution (including the low q_T region) can be predicted. The procedure for the resummation of multiple gluon emission including transverse momentum conservation was introduced in ref. (17) and further developed in refs. (18,13,19). The comparison with W boson production data at $\sqrt{S}$ = 0.54 TeV is shown in Fig. 3.

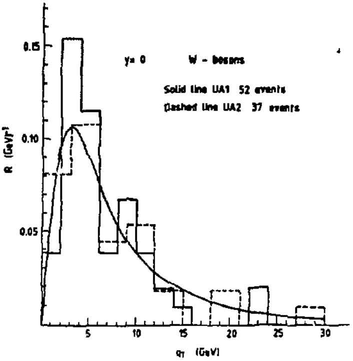

Fig. 3
The normalised differential cross-section R for the production of (W^++W^-) bosons as a function of q_T at $\sqrt{S}$ = 0.54 TeV. The dotted and dashed histograms are the suitably normalised data of the UA1 and UA2 collaborations respectively. The solid line is the theoretical prediction for

$$R = \frac{d\sigma(y=0)}{dq_T\, dy} \Big/ \frac{d\sigma(y=0)}{dy} \qquad (20)$$

based on the parton distributions of Gluck et al.[11] A full analysis[13] of the uncertainty in the theoretical prediction due to the form of the parton distribution functions, the size of Λ, and the uncalculated higher order corrections shows that it is about 25%. Within the limited statistics the agreement between theory and data is acceptable. The change of the ratio R with increasing centre-of-mass energy is illustrated in Fig. 4.

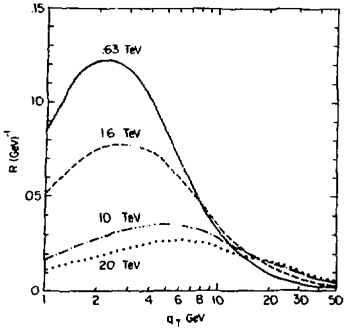

Fig. 4
The normalised differential cross-section R for the production of W^++W^- bosons in proton anti-proton collisions at various centre of mass energies.

With increasing energy a larger fraction of the events lie above q_T = 30 GeV. It is therefore to this large transverse momentum tail, which is well described by the simple perturbative formula, that we turn our attention.

At super colliders the large transverse momentum region is of most interest because it is in this region that the search for physics beyond the standard model will take place. W and/or Z production at large q_T could cause "monojets" or "lepton + jet" events with missing transverse energy. Both of these types of events are typical triggers in the search for new phenomena. In order to estimate the probability of such events from conventional QCD sources, we define the quantity

$$\pi(q_T) = \int_{q_T}^{A_T} \frac{d\sigma(y=0)}{dp_T\, dy}\, dp_T \Big/ \int_0^{A_T} \frac{d\sigma(y=0)}{dp_T\, dy}\, dp_T \qquad (21)$$

where A_T is the kinematic limit of the transverse momentum.

q_T GeV	$\pi^W(q_T)$%		
	$\sqrt{S}$=0.54 TeV $p\bar{p}$	$\sqrt{S}$=10 TeV $p\bar{p}$	$\sqrt{S}$=40 TeV pp
25	3.4 ±0.4	----	---
30	2.0 ±0.2	26.0	---
40	0.8 ±0.1	16.9	---
50	0.40±0.05	11.7	15.
60	0.10±0.02	8.3	11.
70	---	6.0	8.
80	---	4.5	6.
90	---	3.4	5.
100	---	2.6	4.
110	---	2.1	3.
120	---	1.7	2.
130	---	1.3	2.
140	---	1.1	1.
150	---	0.9	1.

Table 2
The probability $\pi(q_T)$ of finding a W boson above a certain q_T at various centre-of-mass energies.[21]

In Table (2) the values of π at $\sqrt{S}$ = 0.54 TeV and 10 TeV for $p\bar{p}$ collisions and at $\sqrt{S}$ = 40 TeV for pp collisions are given. The results at $\sqrt{S}$ = 0.54 have been calculated using the $O(\alpha_s^2)$ contribution[20] coming from quark-antiquark annihilation. The difference between π calculated in order α_s and calculated in order α_s^2 is small, but inclusion of the $O(\alpha_s^2)$ term leads to a substantial decrease in the error which is mainly due to the scale ambiguity in the running coupling constant. At the other two energies the percentage errors on π are of the same order as the percentage errors given in Table 1 at the corresponding energies. The figures are therefore for illustration only. Table 2 indicates that it is most unlikely to find more than 3% of the W's (or Z's, for which a similar result holds) with an associated jet of $q_T \geq 35$ GeV. Taking into account the factor 6 between $\Gamma(Z \to \nu\bar{\nu})$ and $\Gamma(Z \to e^+e^-)$ it follows that at $\sqrt{S}$ = 0.54 TeV we should expect about five times fewer monojets with $q_T \geq 35$ GeV, than regular Z decays to electron pairs at $\sqrt{S}$ = 0.54 TeV.

Acknowledgements

The work described in this report was performed in collaboration with G. Altarelli and G. Martinelli.

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CAN THE PREON SCALE BE SMALL?

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SUMMARY

The preon scale Λ_p is bounded from below by rare or unobserved processes and from above by the cosmological abundance of stable heavy composites. On the other hand composite models can be tested by the Superconducting Super Collider (SSC) or by low energy precision experiments only if Λ_p is allowed to be at most 5-10 TeV. In search of such models we re-examine some conditions that must be fulfilled if Λ_p is small, and point out the possibility of certain mechanisms that could avoid the dangerous rare processes. In addition, certain properties of exotic composite particles, their possible role in breaking the electroweak symmetry and in producing observable signals beyond the standard model are also discussed.

1. LOW ENERGY CONSEQUENCES OF PREON SYMMETRIES

The structure of a preon theory is similar to QCD in many ways. Quarks are confined by color forces at a scale Λ_{QCD} to form hadrons; preons are confined by precolor forces at the scale Λ_p to form composite quarks and leptons (and maybe some exotics). Like the quarks, preons come in several (pre) flavors that define the preonic symmetries. The major difference from QCD is that the preonic chiral symmetries must remain unbroken in the vacuum[1]. They are slightly broken when perturbed by another force which is small compared to precolor. This generates the small masses, $m \ll \Lambda_p$, of quarks and leptons.

At low energies ($E \ll \Lambda_p$), in analogy to the sigma model that follows from QCD, we may write an effective theory (see e.g. ref. 2) that describes the low lying composite states of the preon theory. This must have the form

$$L_{eff} = L(\text{standard}) + L(\text{non-renormalizable}).$$

The symmetry structure of L_{eff} is dictated at the scale Λ_p where the bound states form. At Λ_p all known forces (including QCD) are small compared to the confining precolor interactions. It is therefore useful to consider the limit in which all forces except for precolor is turned off. The fully conserved preonic flavor symmetries G_F that show up in this limit govern the classification of all composite states. These may include

(i) 3 or more generations of massless quarks and leptons

(ii) Massless exotics (color, weak isospin, charge)

(iii) Heavy composites $m \gtrsim \Lambda_p$ classified in irreducible representations $\{r\}$ of G_F.

Only the states (i) and (ii) are included in L_{eff}. At energies $E \gtrsim \Lambda_p$ the states (iii) are also considered.

The symmetries G_F also govern the structure of the 4-fermi and other non-renormalizable interactions that appear in the effective low energy Lagrangian. $SU(3) \times SU(2) \times U(1)$ must be a subgroup of G_F. It is gauged. The classification and structure of interactions provided by G_F are only slightly changed when QCD, electroweak or other mass generating interactions are turned on (however, the model should have the property that these symmetry breaking interactions must not mediate undesirable levels of neutral $\Delta S = 1,2$ or other reactions that may be introduced via mass generation and "Cabibbo" mixing).

The important role of the 4-fermi interactions for testing compositeness at low energies was first discussed in Ref. 2 and later in the 82 workshop[3] and other articles[4]. In the effective theory the 4-fermi interactions are assumed to have the strength $\lambda^2/2\Lambda_p^2$. If they mediate a rare or unobserved process then Λ_p may be required to be large. Here are some of the bounds on Λ_p taken from Ref. 2.

Process	Limit on Λ_p
Proton decay	$\Lambda_p \gtrsim \lambda \times 10^{15}$ TeV
$K^0-\bar{K}^0$ mixing	$\Lambda_p \gtrsim \lambda \times 400$ TeV
$D^0-\bar{D}^0$ mixing	$\Lambda_p \gtrsim \lambda \times 50$ TeV
$K^+ \to \pi^+ \mu e$	$\Lambda_p \gtrsim \lambda \times 30$ TeV
$K_L \to \mu e$	$\Lambda_p \gtrsim \lambda \times 25$ TeV

Naively the magnitude of λ (unless $\lambda=0$ because of symmetry) is estimated to be of order 1 by analogy[2,6] to QCD. [Note different definitions of the scale Λ used by others authors[3,4].] We see that from the point of view of the SSC the most interesting models are those with enough symmetries that require $\lambda=0$ to supress each one of the above (and similar rare) processes.

It is remarkable that many of the proposed preon models can be banned from the TeV regime (i.e. $\Lambda_p \gg$few TeV) thanks to the existence of the few precision measurements listed above. There are proposed experiments to improve the limits of K-decays. The impact of future experiments on Λ_p can be estimated by noting that the dependence of the decay rates on Λ_p is quartic[2]: Γ (K-decay)$\sim(1/\Lambda_p)^4$.

It is not difficult to find models[2] with symmetries that suppress the 4-fermi and higher dimension interactions (i.e. $\lambda=0$ indentically) that mediate (1) proton decay, (2) $K^0-\bar{K}^0$ mixing and (3) $D^0-\bar{D}^0$ mixing. The criteria to eliminate these are as follows[2]: (1) Baryon number must be one of the conserved quantum numbers in the form of a U(1) embedded in G_F. (2) There must be no symmetry embedded in G_F that can transform the left-right components of the composite strange quark when written in the form (s_L, s^c_L), where s^c_L is the charge conjugate of s_R. This may be assured by requiring s_L, s_R to belong to distinct representations of the (sub)group(s) of G_F. (3) There must be no symmetry in G_F that can mix the left-right components of the composite charmed quark in the form (c_L, c^c_L) where c^c_L is the charge conjugate of c_R. Again, this may be assured by requiring c_L, c_R to belong to distinct representations of the (sub)group(s) of G_F. [The following provides an undesirable example: if the Georgi-Glashow SU(5) is embedded in G_F then the 10 contains (c_L, c^c_L) and they can mix via a generator of $SU(5) \subset G_F$. If this happens then $D^0-\bar{D}^0$ mixing will occur via the 4-fermi interactions, and will require $\Lambda_p > 50$ TeV]. These criteria are compatible with the symmetry structure of the standard model based on $SU(3) \times SU(2) \times U(1)$ which is expected to emerge as the low energy limit of the preon theory.

However, as pointed out in ref. 2, the case of K-decays is more delicate because, unlike the other proceses, $\lambda=0$ may not be so easy to achieve by symmetries which classify the quarks and leptons together in repetive families. $K^+ \to \pi^+ \mu e$ or $K_L \to \mu e$ can be eliminated by symmetries only by deviating from the intuitive classification of families suggested by the standard model as described below.

The mass spectrum of quarks and leptons together with SU(3) x SU(2) x U(1) anomaly cancellation arguments within the standard model have led to the notion that a single family contains both quarks and leptons and that there exists at least 3 families of increasing masses. A complete family contains 16 or 15 fermion degrees of freedom. [The structure of Grand Unified Theories reinforces the notion that quarks and leptons belong together in one family.] The repetition as replicas of the first one is not explained in theories of elementary quarks and leptons. In composite models it has been suggested[3] that the repetition is required at least in certain classes of models, due to anomaly cancellation of precolor in the underlying preon theory, thus connecting the existance of families to underlying dynamics.

In the limit of zero gauge couplings for SU(3) x SU(2) x U(1), and absence of a Higgs, the standard model shows a big symmetry: SU(48) (or SU(45)→15 per family) corresponding to 48 (or 45) left handed free fermions. Thus, in the absence of the gauge couplings and masses in the standard model the family structure is completely washed out. This is an accident simply because L (standard) is quadratic in the fermions. However, in a composite model, if there is a family structure, it will show up in the structure of the *4-fermi* and *other non-renormalizable interations*. Thus the preonic symmetry G_F that provides a family structure must be a subgroup of SU(48) or a larger group if there are more families. There are, of course, many possibilities, but the one that suggests itself most intuitevely (when the masses and gauge couplings are turned on) is a cross product of the form

$$SU(48) > (G_V \times G_H) = G_F \tag{1.1}$$

where G_V (V for vertical) acts on the 16 (or 15) members of a family, and is the same for all families,

$$SU(16) > G_V, \tag{1.2}$$

While G_H (H for horizontal) acts on the 3 families. In the limit of zero mass G_H might satisfy $U(3) > G_H$ or $U(3) \times U(3) > G_H$, etc, depending on the number of irreducible representations in which G_V classifies the 16 fermions. [Examples of such structures occur also in grand unified theories; e.g. for SO(10) grand unification G_V= SO(10), G_H=U(3); for SU(5) grand unification G_V SU(5), G_H=$U(3)_5 \times U(3)_{10}$; for Pati-Salam unification G_V = $SU(4) \times SU(2)_L \times SU(2)_R$, $G_H = U(3)_L \times U(3)_R$ etc]. The main thing to notice is not the particular group, but the vertical x horizontal structure that one might expect if families are to be explained by compositeness, and that such an explanation is likely to lump together quarks and the leptons of 1 family within representations of G_V. This type of structure includes the possibilities that

a) Family quantum numbers are carried by a set of family preons while the rest of the usual quantum numbers are carried by other preons.

b) Family quantum numbers come from scalars or pairs of fermions that occur different number of times in different families.

c) Family quantum numbers come from radial quantum numbers.

Thus, under the assumption $G_F = G_V \times G_H$, where G_V lumps quarks and leptons in one family, and G_H distinguishes families, we may analyze the kinds of 4-fermi interactions that must occur with a coupling $\lambda^2/2\Lambda_p^2$, where λ is of order 1. Here we find that there is always a term that mediates $K^+ \to \pi^+ \mu e$ and/or $K_L \to \mu e$, namely

$$\frac{\lambda^2}{2\Lambda^2}\,[\bar{s}_0\,\frac{(1\pm\gamma_5)}{2}d_0][\bar{e}_0\,\frac{(1\pm\gamma_5)}{2}\mu_0] + G_F\text{symmetric terms} \tag{1.3}$$

where the o-index implies that these are the states before mass generation or Cabibbo mixing is taken into account. Assuming that these mixing angles are not large we see that the symmetry $G_F = G_V \times G_H$ can never eliminate this term and thus we must require

$$\Lambda_p \geq (20\text{-}30)\ \text{TeV}. \tag{1.4}$$

[Note that the decays occur for zero Cabblbo angles.] Models satisfying the reasonable assumptions above are therefore just beyond the reach of the SSC (E (max)= 10 TeV in parton + parton center of mass with any appreciable luminosity).

Any model that manages to avoid the conditions of the theorem above is likely to do it in one of the following ways: either

(i) Quarks and leptons are not linked within a family.

or (ii) There is a set of one or more preonic U(1)'s that assign different quantum numbers to quarks than leptons and simultaneously distinguish families.

or (iii) The mixing angles are large so that the mass eigenstate τ, μ, s, d correspond to $e_0 = \tau$, $\mu_0 = \mu$, $s_0 = s$, $d_0 = d$. Instead of τ, e. may correspond to an even heavier lepton.

To these one could add less attractive possitilities that destroy the repetitive family structure, but we will not consider them here, since understanding family repetitions is one of the goals of compositeness.

In the first case it is evident we must give up a simultaneous explanation of quarks and leptons belonging to the same family. In such models it may turn out that leptons could artificially be added to the models by throwing in preonic degrees of freedom that are not required by the precolor dynamics. That is the model could be constructed for only the quarks[7]. We recall that the $U(1)_Y$ gauge anomaly in the standard model is the only evidence of a link between quarks and leptons of the same family. this guage coupling has nothing to do with the precolor dynamics that yield composite quarks and leptons. A model which does not provide a dynamical link between quarks and leptons (in the absence of negligible couplings) may be possible, but we have to ask how palatable it is, since it breaks one of our intuitive expectations.

In the second case I suggest that it is attractive to associate the desired global U(1)'s with the hypercharge Y of the standard model, since this is the only apparent link between quarks and leptons in each family. For example, consider 3 conserved preonic U(1)'s that assign separately the hypercharges in each family. [The gauge $U(1)_Y$ is the "diagonal" U(1)]. These U(1)'s or an appropriate discrete subgroup embedded in them are sufficient to eliminate the dangerous terms of type (1.3). While this sounds attractive a model of this type has not yet been constructed.

The third case* of large mixing angles is also counter intuitive. However, here there may be room for much further investigation since an attractive mass generating mechanism does not yet exit. Note that even though mixing angles may completely be rotated away in the lepton sector in L (standard) (certainly so, if ν_R do not exist), this is not *necessarily the case in L(4-fermi), Since L(4-fermi)* is not quadratic in the fermions. Thus, in this mechanism the burden of suppressing K_L, K+ rare decays rests with the mass generating mechanism without compromising the suspected linkage between quarks and leptons. The classification scheme for mass eigenstates is then expected to look as follows

$$\begin{array}{ll} \text{1st family} & \binom{\tilde{u}}{d}_L \; u_R \; d_R \; \binom{\nu_\tau}{\tau}_L \; \tau_R \; \nu_{\tau R} \\ \text{2nd family} & \binom{\tilde{c}}{s}_L \; c_R \; s_R \; \binom{\nu_\mu}{\mu}_L \; \mu_R \; \nu_{\mu R} \\ \text{3rd family} & \binom{\tilde{t}}{b}_L \; t_R \; b_R \; \binom{\nu_e}{e}_L \; e_R \; \nu_{eR} \end{array} \qquad (1.5)$$

where $\tilde{u}$, $\tilde{c}$, $\tilde{t}$ are the (u,c,t) mass eigenstates rotated by the Cabibbo-Kobayashi-Maskawa mixing angle. With such a mass scheme, e.g. some of the models discussed in ref. 2,6 would completely avoid all the bounds discussed above.

Furthermore, by mixing the (u,c,t) quarks rather than the (d,s,b) quarks, $\Delta S=1$, neutral current 4-fermi interactions do not occur. the family changing interactions that are generated by this mixing scheme are not restricted by known phenomenology. In L(standard) it does not matter whether the ups or the downs mix, however, in L(4-fermi) it makes an important phenomenological difference. Of course, the mass generating mechanism holds the secret for why the ups rather than the downs (or both?) should mix.

An example of trouble free 4-fermi interactions *that illustrate the points above is explicity* exhibited in section 3.

2. COSMOLOGICAL UPPER BOUND ON Λ_p

In the previous section we discussed bounds coming from low energy physics. However, cosmological consideration can help probe the heavy sector $M \sim \Lambda_p$ of a preon model if there are long lived states. This idea was first implemented in ref. 6, as outlined below.

A preon model often has some (naively) conserved U(1) quantum numbers. The low mass quarks and leptons can be taken neutral under some U(1) but some heavy states are charged. Then, in the same way that the proton is stable, such states are also (naively) stable.

Note that I emphasized naively conserved U(1). This is because after stronger precolor instanton effects this U(1) may be broken (it is broken in ref. 6). However, one must still analyze the effective instanton interaction and estimate the rate at which the heavy state is allowed to decay. Then, an interesting huge suppression may be found if the only allowed decays are to a large number of particles, despite a strong effective coupling constant. For example, the lifetime of a heavy scalar particle, $M \sim \Lambda_p$, that decays to N massless particles in the final state must be larger than

$$\tau \geq \frac{1}{(G^2 \Lambda_p)} \frac{(16\pi^2)^{N-1}}{\pi} \frac{(3N-4)!}{(4N-4)!} (2N-1)! \; (2N-2)! \qquad (2.1)$$

Here G is a dimensionless effective coupling that measures the strength of the (instanton) interaction. A realistic model may require N of order 16, corresponding to the 16 members of a family, as in the example considered in ref. 6. Then

$$\tau \geq \frac{(100 \text{ TeV})}{G^2 \Lambda_p} \quad (4 \times 10^{34}) \quad \text{years.} \qquad (2.2)$$

Thus, even for a large value of Λ_p, the lifetime of such a particle is larger than the lifetime of the universe. This illustrates that U(1)'s that are broken by instanton effects should not be dismissed, as they may still lead to almost stable particles.

In the event that a preon model has long lived particles (even for lifetimes than several minutes), cosmological considerations can put limits on its Λ_p. In ref. 6, mainly the case of $\tau \geq \tau$ (universe) was discussed. It is estimated that the abundance of such stable particles in today's universe is

$$\left(\frac{N}{N_Y}\right)_{\text{today}} \sim \left(\frac{\Lambda p}{M_{\text{planck}}}\right) \; \ln \frac{(M\text{planck})}{\Lambda_p} \qquad (2.3)$$

For these not to dominate today's matter (baryons) dominated universe, we must require

$$\Lambda_p \leq 250 \text{ TeV.} \qquad (2.4)$$

It may be possible to improve this bound by taking into account clustering of such particles in the form of galaxies. In any event, the fact that there is an upper bound in certain potentially realistic models and that the bound is fairly low is rather interesting from the point of view of the SSC.

3. A MODEL WITH EXOTICS

A preon model can be tested at low energies if it has exotic bound states that are G_F- partners of the (massless) quarks and leptons. The mass of such states is likely to be in the range

$$m_{top} < m < \Lambda_p, \qquad (3.1)$$

thus requiring energies lower than Λ_p for discovering them. The recent jet activity around m~150 seen at the UA1 and UA2 detectors at CERN may be attributed to exotics, as discussed in the Compositeness Subgroup at the SSC Workshop[9]. The model presented here is an example which has a minimal number of exotics [1 color nonet (8+1)], and can provide signals of the type seen at CERN.

The precolor group is taken as G_p=SU(4)xSU(4) and the preons are placed in the three representations $R_1=(\underline{4},\underline{4})$, $R_2=(4,1)$, $R_3=(1,4)$. The numbers and helicities of the preons are

$$1_L \; R_1 + 4_L \; R_2 + (10_L + 6_R) R_3 \qquad (3.2)$$

Thus, the preflavor symmetry G_F which classifies the preons and composites is (after instanton effects)

$$G_F = SU(4) \times SU(10) \times SU(6) \times [U(1)]^2 \times Z^2 \qquad (3.4)$$

The massles composites which satisfy anomaly, decoupling and certain other conditions for the entire conserved G_F are:

$$(4,\ 10,\ 1)_L^{(1,0)}\ (4,\ 1,\ 6)_R^{(0,1)} \qquad (3.4)$$

This solution was used before in refs. (2,6) (without exotics) with a different interpration of the "flavor" quantum numbers than the one suggested below.

We embed $SU(3) \times SU(2) \times U(1)$ in G_F so that the preons are classified as follows:

$$R_2:\ 4_L \to (3,1)_{1/6} + (1,1)_{-1/2} \qquad (3.5)$$

$$R_3:\ \begin{array}{l} 10_L \to (1,2)_0 + (1,2)_0 + (1,2)_0 + (\bar{3},1)_{-1/6} + (1,1)_{1/2} \\ 6_R \to (1,1)_{1/2} + (1,1)_{-1/2} + (1,1)_{1/2} + (1,1)_{-1/2} + (1,1)_{1/2} + \\ \qquad + (1,1)_{-1/2} \end{array}$$

The subscripts are the U(1) quantum numbers. Note that this embedding is anomaly free for gauged $SU(3) \times SU(2) \times U(1)$, as it should be. QCD is embedded in SU(4) a la Pati-Salam. Therefore, the composites are classified as $(4 \to 3 + 1)$

$$(4,\ 10,\ 1)_L \begin{cases} \nearrow 3:\ \ 3\times(3,2)_L^{1/6} + (3,1)_L^{2/3} + (1,1,)_L^0 + (8,1)_L^0 \\ \searrow 1:\ \ 3\times(1,2)_L^{-1/2} + (\bar{3},1)_L^{-2/3} + (1,1)^0_L \end{cases}$$

$$(4,\ 1,\ 6)_R \begin{cases} \nearrow 3:\ \ 3\times(3,1)_R^{2/3} + 3\times(3,1)_R^{-1/3} \\ \searrow 1:\ \ 3\times(1,1)_R^0 + 3\times(1,1)_R^{-1} \end{cases} \qquad (3.6)$$

This corresponds to 3 usual families of quarks and leptons plus a fourth up quark, plus a color nonet $(3\times3^*,1) = (1,1)_L^0 + (8,1)^0_L$ and a singlet $(1,1)^0_L$. The quarks and leptons may be indentified as in (1.5) so that Λ_P is not restricted by the rare processes discussed in section 1.

The point of this model is the presence of the nonet so that the singlet and octet have the same global quantum numbers, corresponding to a conserved U(1) embedded in G_F. Suppose the octet is heavy. If produced in $p\bar{p}$ reactions at CERN it can decay to a pair of quark + antiquark plus the neutral singlet that carries the same global quantum number as the octet. Thus in the final state one would see a pair of highly energetic jets plus missing energy. Since one of the quarks may sometimes be slow, the event (after the cuts) can also look as 1 energetic jet plus missing energy. The cross section for production + decay is quite large and can explain the rates seen at CERN, as discussed in the compositeness group in this workshop.[9] Note that the octet of this model has some properties similar to the gluino in supersymmetric theories, if the gluino its taken at around the same mass, and may be confused with it.

More model independent properties of exotics, are discussed in ref. 9.

I wish to propose another important role for exotics in a composite model. Marciano[11] suggested that high color states (6, 8, 10 etc.) may condense at the electroweak scale $F\pi \sim 250$ GeV, thus providing a mechanism of mass generation analogous to technicolor but only with QCD forces. In the context of composite models this idea is quite attractive because

(i) Exotics occur naturally

(ii) The 4-fermi interactions provide masses for quarks and leptons after condensation.

In the models of elementary quarks and leptons discussed in ref. 10, it was difficult or unattractive to implement a substitute for (ii).

To use this mechanism one must address questions[12] about the asymptotic freedom of QCD because, if QCD looses its asymptotically free behaviour due to many exotics, condensation would take place at the highest values of α_{QCD}, thus at the highest scales. This is not desirable. For this I emphasize that in a composite model we must separately consider the calculation of β_{QCD} in the regimes below Λ_P and above Λ_P. Below Λ_P there are few and non-exotic preons. In terms of preons β_{QCD} must and can easily be negative for asymptotic freedom to be correct. Below Λ_P the behaviour of α_{QCD} or β_{QCD} may be smooth or complicated depending on the number of exotics and their thresholds In the range $\Lambda_{QCD} < \mu < \Lambda_P$ condensation will occur if $\alpha_{QCD}(\mu)$ attains the critical value at $\mu = F\pi = 250$ GeV

$$\alpha_{critical} = \alpha_{QCD}\ (F_\pi) \qquad (3.7)$$

[$\alpha_{critical}$ may approximately be estimated [10,11] via the quadratic casimir for the exotic representation R, $C_2(R)\alpha(F_\pi) = 1$]

For $\mu > F\pi$, $\alpha(\mu)$ must never exceed αcritical, otherwise the scheme will not have any meaning. Two possibile situations are shown in Figs. 1 and 2

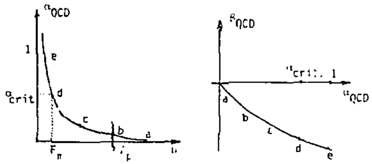

Fig. 1. Few exotics. $\beta < 0$ for all scales.

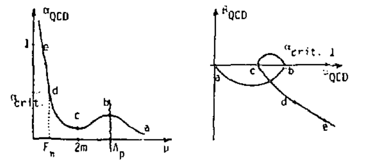

Fig. 2. Many exotics. $\beta > 0$ above 2m threshold; $\beta < 0$ above Λ_p.

In Fig. 1, even below Λ_p, there are few exotics so that the β function (slope of $(\alpha(\mu))$ always remains negative. In Fig. 2, there are too many exotics below Λ_p. The threshold for producing the exceeding exotics is $\mu=2m$, above which β_{QCD} is positive. However beyond Λ_p β_{QCD} is again negative since the computation is done in terms of preons. Note the interesting multivalued plot of β versus α for this case which, as explained, can happen quite naturally in a composite model. Each branch of this curve is computed pertubatively since $\alpha_{QCD}(\mu)$ is small. The non-perturbative phenomena occuring via the underlying precolor forces is what gives rise to such a non-pertubative looking curve.

For these mechanisms to be useful for electroweak symmetry breaking there should be some exotics carrying electroweak quantum numbers, such that $\Delta I^w=1/2$. These could be of the form $(r,2)_L+(r,1)_R$ where r is a complex representation of SU(3), such as r=6, 10, etc., and 2 is a doublet, 1 is singlet of $SU(2)_w$. The numbers of doublets and singlets sould be such that the symmetry breaking preserves a custodial SU(2) (approximately). We cannot allow r = real (e.g.(8,2)) since this would lead to $\Delta I_w=1$ via $(r,2)_L \times (r,2)_L \sim (1,3)$. Any real exotic representation should not simultaneously be a doublet of $SU(2)_w$ e.g.(8,1) is o.k.). As Marciano estimates, 2 sextets together with the usual 3 families just about saturate asymptotic freedom for QCD. Thus, although there is the possibility of a composite model described by Fig. 1, most models with exotics are likely to be described by Fig 2, if they play any role in electroweak symmetry breaking.

Models with exotics now being investigated will be described in future publications.

ACKNOWLEDGEMENT

This research was supported by a grant from the Department of Energy, No. PR-0384 ER40168.000, and by the USC Research and Innovation Fund.

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(5) The reader should be aware that the different definitions of compositeness scales imply bounds that may seem different by a factor of 5, but can be translated to be in agreement with each other. In QCD the mass of the rho satisfies $m\rho=(3-5)$ ΛQCD where ΛQCD is thought to be in the range 150-250 MeV. Λ_p is defined in analogy to ΛQCD. Other authors (e.g. ref 3,4) have used the mass of the expected composite vector meson Λ^* as a convenient scale. By analogy to QCD we might expect $\Lambda^* = (3-5)\ \Lambda_p$. Thus the strength of the 4-fermi interaction parametrized in ref 2 as $\lambda^2/2\Lambda_p{}^2$, with λ of order 1, is consistent with $4\pi/2\Lambda^{*2}$ used in ref. 3,4.

(6) I. Bars, M. Bowick, K. Freese, Phys. Lett 138B, 159 (1984).

(7) A model which appears to fall in this category was described to me by O. W. Greenberg and S. Nussinov after this workshop: The composites contain a fermion and a boson $(\psi\phi)$; $\psi_{L,R}$ carry $SU(2)_L \times SU(2)_R$, while ϕ^{ai}, ϕ^{4i} carry 3=color, 4=lepton, i=quark family, I-lepton family (different than quark family).

(8) I thank H. Harari, M. Bowick and M. Dine for discussions.

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THE CASE FOR LIGHT GAUGINOS

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Summary

I consider the possibility that gauginos be substantially lighter than quark and lepton superpartners. *After a brief review of the present limits, I go* through several possible spectra for light gauginos, photinos and w-inos, each leading to quite different experimental signatures.

1. *At this meeting, in most of the discussions* on supersymmetry signatures, we have taken gaugino masses ranging from a few ten of GeV up to the Tev region.*

Here I consider the possibility that gaugino masses -in many instances not simply related to the scale of supersymmetry breaking- be considerably lighter than the scalar superpartners of quarks and leptons, thus making some gauginos more easily amenable to experimental search.

As a mere way to produce sensible mass spectra for light gauginos and discuss the corresponding signatures, I refer to N = 1 supergravity models giving rise at low energy to a softly broken, globally supersymmetric SU(3) x SU(2) x U(1) theory[1]. Furthermore, I invoke an approximate symmetry (U(N), U(N,1)) of the basic Lagrangian enforcing the vanishing of the tree level gaugino Majorana masses[2]. I mimic in this way the situation of a renormalizable supersymmetric field *theory which is known to produce gaugino Majorana* masses only via loops. For example, in the gluino case, I believe that it makes physical sense to start with a tree level vanishing mass, $m_{\tilde{g}}^{(o)} = 0$, to avoid unwanted CP violation effects possibly related to an imaginary part of $m_{\tilde{g}}^{(o)}$ in the neutron electric dipole moment or in the QCD θ-parameter.

2. The laboratory limits for a light gluino ($m_{\tilde{g}} \lesssim 5$ GeV) coming from beam dump experiments and stable particle searches, are summarized in Figure 1, taken from Dawson, Eichten and Quigg[3]. These limits are probably conservative, since they do not include the implications of the decay $J/\psi \to \tilde{g}\tilde{g}$ (but who knows the threshold effects due to the physical mass of the gluino containing R-hadrons ?) or the limits on neutral stable R-hadrons produced in relatively low energy experiments (but there the relevant production cross sections are uncertain). Nevertheless a gluino lighter than 5 GeV is not excluded. On the other hand, for a heavier gluino which decays into hadrons and an unseen photino, its copious pair production in $p\bar{p}$ collisions allows to exclude a mass window 5 GeV $\lesssim m_{\tilde{g}} \lesssim$ 30 GeV already at present*.

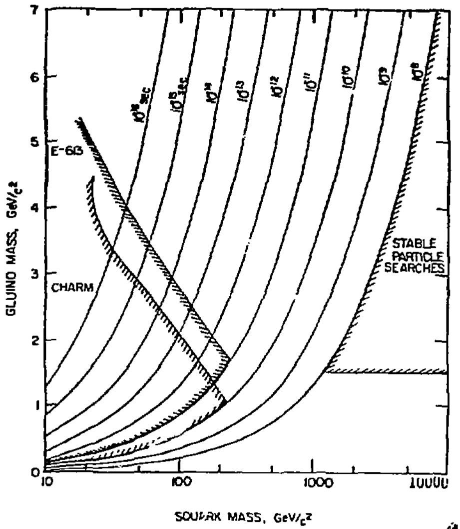

Fig. 1. : Limits on the gaugino mass as function of the lightest squark mass. The gluino is assumed to decay into a $q\bar{q}$ pair and a massless photino. The corresponding lifetimes are also shown.

The photino is more elusive than the gluino. For a scalar electron heavier than about 100 GeV -the situation I will mostly consider- there is no limit at present from laboratory experiments on stable or quasi-stable photinos.

As is well known, assuming the conservation of a discrete R-parity, the consequent stability of the lightest supersymmetric particle allows to limit its mass due to cosmological considerations. The photino can be stable only if its pair annihilation into quarks or leptons via scalar exchange is efficient enough to reduce its cosmological mass density below 2×10^{-29} gr/cm^3. In turn, this sets a lower limit on its mass, $m_{\tilde{\gamma}}$, which

* see the talk by John Ellis

* see the talk by John Ellis

rapidly increases with the mass M_S of the scalar quark or lepton exchanged : typically $m_{\tilde{\gamma}} \gtrsim 1 \div 2$ GeV for M_S = 40 GeV, $m_{\tilde{\gamma}} \gtrsim 10 \div 15$ GeV for M_S = 100 GeV [4]. On the other hand, odd as it might be, a stable gluino would always annihilate sufficiently fast into a virtual gluon not to give any appreciable contribution to the cosmic mass. Still there would be residual gluinos with a number density relative to hydrogen of about 10^{-10}, forming neutral stable hadrons, like ($\tilde{g}g$) bound states. In turn these states would bind to nuclei, giving rise to anomalous isotopes. From studies[5] of O_{18} this is excluded for masses between about 5 to 35 GeV. On the other hand, the comparison of chemically versus physically determined masses is not accurate enough to give any other constraint.

3. With reference to the theoretical framework I have already mentioned, one can discuss[6] four types of light gaugino spectra with different experimental implications and different degrees of plausibility, depending on the value of the gravitino (m) and the top quark mass (m_t)*

i) (1TeV $\lesssim m \lesssim$ 10 TeV, $m_t \approx$ 50 GeV)

$m_{\tilde{W}} = m_{\tilde{Z}} = 20 \div 50$ GeV

$m_{\tilde{g}} = 20 \div 50$ GeV

$m_{\tilde{\gamma}}$ $7 \div 15$ GeV

One is led here to a rather classical phenomenology for supersymmetry searches, the present $p\bar{p}$ collider being quite suited to perform this search. The reactions of interest are

$p\bar{p} \to \tilde{g}\tilde{g} \to$ Jets + missing p_T,

as already mentioned, and

$$p\bar{p} \to W + \text{anything}, \quad W \to \tilde{W}\tilde{\gamma}$$

$$p\bar{p} \to Z + \text{anything}, \quad Z \to \tilde{W}\tilde{W}$$

with the w-ino decaying, most of the time, into a photino and a virtual W. The signature for w-inos is the missing energy carried away by the undetected photinos.

ii) (1TeV $\lesssim m \lesssim$ 10 TeV, $m_t \lesssim$ 35 GeV)

Comparing with the spectrum in i) the light top could give rise to a gluino lighter than the photino. Odd and unlikely as it appears, this situation is, however, not excluded. Now it is the photino that decays into a gluino plus a $\bar{q}q$ pair and not vice-versa, with typical lifetimes of $10^{-10} \div 10^{-11}$ secs. The most striking consequence would be the absence of the missing p_T signature associated with supersymmetric particle production. Beam dump experiments are obviously irrelevant. This seems to be the worst case for an experimental search, although probably a careful analysis of W decays could still reveal the $\tilde{w}$-ino decay mode

$$\tilde{W} \to \tilde{\gamma} + \text{virtual } W, \quad W \to \bar{q}q,\ \bar{l}l, \quad \tilde{\gamma} \to \tilde{g} + \bar{q}q$$

iii) (150 GeV $\lesssim m \lesssim$ 400 GeV, $m_t \approx$ 50 GeV)

$m_{\tilde{Z}} \approx m_{\tilde{W}} \approx 20 \div 35$ GeV

$m_{\tilde{g}} \approx 1 \div 2$ GeV

$m_{\tilde{\gamma}} \approx 400 \div 700$ MeV

This is the classical case already discussed by Farrar and Fayet[7] a long time ago. Although the gluinos are copiously produced in hadronic collisions, their detection could be obscured by nonperturbative effects. On the other hand, the relatively heavy scalar quarks ($m_S \approx m \gtrsim 150$ GeV) can make the gluino lifetime long enough ($\gtrsim 10^{-10} \div 10^{-9}$ secs.) so as to make the beam dump experiment uneffective and rather suggest searches for anomalous tracks or decay paths in hadronic collisions. In this situation the most promising way to look for supersymmetric signals is the search for light w-inos and z-inos in W and Z decays. A non negligible fraction of light gluinos in the proton could give rise to single squark production in high energy hadronic collisions[8].

iv) ($m \lesssim$ 400 GeV, $m_t \lesssim$ 35 GeV)

Again, one may consider, as in ii) an inversion of the spectrum between gluinos and photinos. Suppose that the only stable gluino containing R-hadron, with a mass near 1 GeV, be electrically neutral. It could possibly have escaped detection so far. The remarks made in ii) on the missing p_T signal apply here as well.

Looking back at the cosmological constraints, in all cases, i), ii), iii), except the last one, iv), there is a conflict of cosmology with the spectra that we find : another particle is required to play the role of the lightest superpartner. As a possibility, one can think of a gauge singlet fermion, z, maybe needed for independent reasons, with a mass $\lesssim$ 1 KeV[6]. In this case, one would have photino or gluino decays $\tilde{\gamma} \to \gamma + z$ ($\tilde{g} \to g + z$) with typically long lifetimes, $1 \div 10^3$ secs.

4. The present status of supersymmetric models does not allow any firm prediction of the various gaugino masses. For this reason, it is useful to keep in mind

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* The top quark mass enters into the determination of the gluino mass, scaling approximately as m_t^2, since I assume, in order to make a prediction, that the top is the heaviest coloured particle exchanged in the loop for the radiative gluino mass. For more details see ref. 6.

that a light gaugino spectrum is far from being excluded on an experimental basis. One has in fact, various options which, even with different degrees of likelihood give rise to different alternative phenomenologies not to be discarded a priori.

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