PLENARY REPORT

COMPOSITE MODELS OF HADRONS AND RELATIVISTIC BOUND STATES

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1. Introduction

The main problems to be considered in this report are the following: 1) What are the constituents of the hadrons, what are their quantum numbers, and what are their broken and unbroken symmetries? 2) What is the dynamics of the constituents (equations, binding forces and the origin of symmetry violations)? The most puzzling question is: why the constituents "escape from freedom" and are confined inside the hadrons? 3) What experimentalists can tell us about the hadron constituents and their dynamics, if not finding them?

There are no final answers to all these questions. Today we can only give more or less plausible answers demonstrating that the questions are sensible.

Due to a great complexity of the matter, this review is by no means impartial. Never - theless, the attempt is made to present also alternative views on the same problems. Many people have been thinking of these problems for years and it would be proper to remind their results having to do with the present-day concerns. "Those who do not remember the Past are condemned to repeat it" (Jaspers).

2. Constituents

2.1. A way to "Colour-ado"

The first model of composite hadrons was constructed by Fermi and Yang (1949) (F-Y). The constituents were the proton and the neutron, the pions being composed of them: $\pi^+ = p\bar{n}$, $\pi^0 = \frac{p\bar{p} - n\bar{n}}{\sqrt{2}}$. The main difficulty of the model was the absence of $\pi^0_0 = \frac{p\bar{p}+n\bar{n}}{\sqrt{2}}$ with mass roughly equal to that of the pion (Baldin et al.). The F-Y model was treated by physics community without any enthusiasm, but, as strange particles were being discovered, the attitude to composite models was becoming more friendly. In 1953 Goldhaber tried to compose all particles from p, n, K^0, K^- (more symmetric form of this model was given by Frisch (1960), with constituents $\Lambda, K^0, K^+, \bar{K}^0, K^-$. A more natural extension of the F-Y model was proposed by Sakata (1956). All particles were assumed to be made of p,n and Λ (F-Y-S). This model successfully explained many facts of strong, electromagnetic and weak interactions, and in terms of it the SU(3) symmetry was first formulated (Markov, Okun, Ikeda, Ogawa, Ohnuki, Yamaguchi, Zeldovich et al.). With proper modifications, it still can be applied to the mesonic states. However, it could not naturally explain the baryon spectrum, and after the success of the eight fold-way approach to the SU(3)-symmetry (Gell-Mann and Ne'eman) it made the way to the quark model (QM) (Gell-Mann and Zweig). The essential difference of QM from F-Y-S model lies in the three-fermion structure of the baryons. The elementary particles in QM are the three spin -1/2 fermions(quarks) q: $u\ (Q=2/3, I=1/2, I_3=1/2)$, $d\ (Q=-1/3, I=1/2, I_3=-1/2)$ and $s\ (Q=-1/3, I=0)$, mesons being $\bar{q}q$ and baryons — qqq.

This model enables us to formulate the SU(6) spin-unitary spin symmetry (Gursey, Radicati, Sakita, Pais, B.Lee, Beg et al.) which, if properly formulated and used, gives qualitative understanding of hadron spectroscopy as well as collinear decays and scattering processes. The successes of this approach to hadron dynamics were summarized at the London conference /1/ and they are really impressive. Most impressive is the remarkable simplicity of the quark dynamics to be discussed below.

But let us turn to the difficulties. The main difficulty lies in understanding the baryon spectrum. In terms of the $SU(6) \times O(3)_L$ symmetry (L is the total orbital momentum of the quarks) the observed baryons are classified (see e.g. /2/) into the (mult. SU(6), L^P)

multiplets $(\underline{5}^6, 0^+), (\underline{70}, 1^-), (\underline{56}, 2^+), (\underline{56}, 0^+)_R$ ... with no candidates for $(\underline{20}, L^P), (\underline{56}, 1^-), (\underline{70}, 0^+)$ etc. Here eight $J^P = 1/2^+$ baryons and ten $3/2^+$ resonances nicely complete $(\underline{56}, 0^+) = (\underline{8}, 1/2^+) + (\underline{10}, 3/2^+)$ But why do they belong to $(\underline{56}, 0^+)$?

With Fermi statistics for quarks, it is extremely difficult, if not impossible, to construct a potential which gives a large mass for antisymmetric in spin and unitary spin $\underline{20}$ and the mentioned peculiar correlation of the SU(6)-structure to L^P (Instead , it is easily understandable with Bose-quarks). Even worse, the nonexistence of exotic states $qq\bar{q}$, $qqq\bar{q}$ (the triality puzzle) and qq (the diquark puzzle) says us that something essentially new must be added to the Gell-Mann -Zweig model. In view of these difficulties a new degree of freedom was introduced in 1965 by Bogolubov, Struminsky, Tavkhelidze et al. /3/ and by Han and Nambu /4/ . Each quark was supplied by a new quantum number (now called the "colour"), and it was postulated that the lowest baryon multiplets are made of quarks of different colours(i.e., "colourless"). The difficulty with the statistics was immediately resolved while the triality and diquark puzzles were reformulated in terms of the problem that coloured states must be much heavier than colourless ones.

Now we have three families of quarks (u_R, d_R, s_R), (u_B, d_B, s_B) and (u_G, d_G, s_G) (red, blue and green). There are two essentially different possibilities for prescribing to these quarks the usual quantum numbers (now called the "flavours"). It is most natural to introduce a symmetry in the colour space, say $SU(3)^c$ (the other possibilities were also discussed). Then, if the electric charge Q is the colour--singlet, the quarks of all colours have identical flavour quantum numbers (fractional charges) and the SU(3) - symmetry must be unbroken. Alternatively, if we assume Q to be colour non-singlet, the quarks may have integer charges and the $SU(3)^c$ need not be the exact symmetry.

Other approaches to the statistics paradox (e.g. Greenberg /5/) are generally connected with violations of the spin-statistics theorem for quarks and, hence, require a revision of the foundations of the quantum field theory.

2.2. Charmed Colour-ado. How many flavours?

The primary motivation for introducing new flavours (charmed quarks) was simple: why not have a higher flavour symmetry, say SU(4) (Tarjanne, Teplitz; Hara; Maki, Ohnuki; Vladimirsky; Gerstein, Whippman; Bjorken, Glashow, Amati et al.; Okun et al., for refs. see eg./6/).

Later it became clear that the fourth quark (c) is indeed useful for making weak interactions internally consistent (Glashow, Illiopoulos, Miani). Without extra quarks and/or leptons all usual formulations of weak interactions badly violate universality and cannot survive. With the charmed quark we can also restore the lepton-quark universality. Finally, the existence of strangeness - conserving neutral weak current and the absence of the strangeness-changing ones is easily explained only in the theories with extra quarks. Thus we have very good reasons to believe in the fourth quark.

The discovery of J/ψ family (Ψ) tells us that we are on a right track interpreting them as $c\bar{c}$ /7/ . Of course, there are more exotic explanations of the Ψ- particles but we can say that in general the charmonium $(c\bar{c})$ spectroscopy is today in a good shape (see esp. the talk given by De Rujula at this conference). As we have heard at this conference there is a reason to believe that a particle $D = u\bar{c}$ (c is the charmed quark), was discovered at SLAC. If confirmed, this would give the c- quark the same status as the older, uncharmed quarks (being necessary though invisible). But are we in need of more quarks?

The experiments, much discussed at this conference, seem to tell us that the introduction of one or two extra quarks ("t" and "b", e.g.) would be harmless and even agreable. The large R in e^+e^- annihilation into hadrons, dilepton events, anomalies in ν - reactions, etc., are most naturally understood with five or six flavours of quarks and some new leptons /8/ . Such new flavours were discussed as soon as Ψ was born (Barnett; Harari; P.Bogolubov, Matveev, Kuz'min, Tavkhelidze et al.; Mohapatra, Pati; Fritzsch, Minkowski; Wilczek et al. /9,10/). Later it was suggested that more than three quark flavours are required in unified theories of the weak, electromagnetic, and strong interactions based on a single gauge group, if one wants to avoid the introduction of extremely heavy gauge bosons ($\gtrsim 10^{15}$GeV). The minimal number of quark flavours in this approach is six (for a detailed explanation and references see: Fritzsch /10/). Introduction of new quarks and leptons allows one to construct beautiful vector- like theories of weak interaction (for refs. see /10/) ; unfortunately they are in a bad shape at this conference. However, there exist other (less symmetric) theories of weak interactions with extra quarks and leptons /11/,/12/ which do not contradict present experiments /8/ .

There exists an entirely different approach to the problem of the flavour degrees of freedom which attempts to give the internal symmetry space a geometric meaning. As has been argued by Arbuzov and Filippov /13/ for weak interactions, and by Ne'eman /14/ for strong interactions, giving a geometric meaning to the elementary particle interactions (i.e., connecting them with a curvature of the space-time) requires an embedding of the space-time into some N-dimensional space. The minimum value of N is 10, as we can locally embed any curved 4-dimensional space into some 10-dimensional space ($g_{\mu\nu}$ has 10 independent components). Hence, the dimension of the compact internal space must be not less than 6. The corresponding symmetry group according to /13,14/ must be SO(6) which is locally isomorphic to SU(4). It should be stressed that this is not the final answer. In fact, in a geometric theory of this type the dimension and the structure of the internal space is dependent of the interaction which, in turn, has to be deduced from geometric constraints (see /13/). The solution of all such constraints is a challenge, very little known as yet even how to approach this problem.

The geometric approach was recently revived in an attempt to construct a consistent dual theory of hadrons in the four-dimensional space-time. As has been earlier suggested by Fubini and Veneziano /15/ , the extra dimensions (26-4 or 10-4) required by consistency of dual models can be ascribed to the "internal" space ("flavour space"). Then for a large enough dimension of the flavour space the dual theory could be realized in the 4-dimensional space-time . This idea combined with the geometric approach to the origin of symmetries is being investigated by Scherk, Schwartz et al. /16/ . As the most difficult problem of solving geometric constraints is not yet clarified, it is premature to deduce from their results any predictions concerning flavour symmetries. The only firm prediction is that the flavour space must be large enough; the larger is the dimension, the easier is the solution of the geometric constraints.

To my mind, the most unsatisfactory feature of all above mentioned approaches to flavours is their inability to explain colours. Usually one asks the questions: how many quarks are there? and what is the flavour symmetry group? Probably these questions are unfair to Nature as stressed by Salam at this conference. Our primary concern must better be not the number of quarks and leptons but the number of conserved charges and the nature of fundamental laws, controlling the "elementary" particle phenomena. We think that a "colour-blind" person cannot find such laws.

2.3. Confining colours give birth to flavours

What is the most fundamental thing in hadrodynamics? We cannot answer this question, but we know the most puzzling thing: coloured particles (quarks, diquarks e.a.) do not occur in the physical spectrum. If this is not merely the low energy phenomenon, then the colour, being exactly conserved, has a chance to be the most fundamental property of quarks, and the colour conservation has to be considered as one of the most fundamental laws of Nature. With permanently frozen (confined) colour degrees of freedom, we face the novel feature of particle dynamics, probably requiring a revision of some basic ideas. Some people proposed that the phenomenon of confinement is "simply" reduced to disappearing the quark pole from the quark propagator in nonperturbative solutions of quark field theory. It is possible, but our task is much more ambitious: to construct colourless bound states and to prove that coloured states never appear in the physical world. There might be some analogy with quantum electrodynamics, where the longitudinal and "time-like" photons play a significant role in intermediate states but completely decouple from all asymptotic physical states (in modern usage they are "confined gluons" . For quantum description of this decoupling it is necessary to supply the Hilbert space with indefinite metrics. This simple trick does not help to confine much more complicated colour degrees of freedom, and it is quite possible that a more radical modification of the physical laws is necessary for describing coloured quarks.

The most radical approach to colour was proposed by Gursey /17/ et al. He suggested that the matrix elements of the quark field operators are octonions (Cayley numbers) instead of being complex numbers. The octonion can be written as $\tau_o + \sum_{A=1}^{7} \tau_A e_A$ where τ are real (or complex) numbers and e_A are "imaginary units" : $e_A^2 = -1$. The multiplication law of e_A is non-commutative and non-associative (generally $(ab)c \neq$ $\neq a(bc)$). Like real and complex numbers, and quaternions, which can be used for describing the spin, the octonions form a normed algebra $\mathcal{O}$ having a unit element (in fact $\mathcal{O}$ is the highest dimensional algebra having such properties). Replacing complex numbers by octonions we are forced to replace the i $(i^2=-1)$ in the translation operator $e^{-ip_\mu x_\mu}$ by one of the new imaginary units, say e_7 . It appears that the subalgebra of $\mathcal{O}$ which commutes with translations $e^{-e_7 p_\mu x_\mu}$ is isomorphic to the algebra of SU(3). This is identified with $SU(3)^c$ and the quark fields are represented by $q_\alpha^n u_n$ where $u_n \in SU(3)^c$ and α is the flavour index. Due to the non-associativity of octonions, only colourless operators can be observable quantities.

With octonionic quarks the nightmarish quark-parton paradox can be resolved. It can be formulated as follows /18/ : the m.e. $\langle p|\bar{\Psi}_1(x)\Psi_2(0)|p\rangle$ has vanishing Fourier components if p_0 is below the colour-production threshold (∞ for permanently confined colour) and this is inconsistent with the early scaling in deep inelastic processes. In other words, how can one explain scaling with zero imaginary part of the quark propagator? The proof is based on inserting $I = \sum |n\rangle\langle n|$ between quark fields Ψ_1 and Ψ_2 and on using the translations. For octonions $\langle\alpha|[(\sum_n |n\rangle\langle n|)|\beta\rangle] \neq$ $\neq \sum_n (\langle\alpha|n\rangle)(\langle n|\beta\rangle)$ and the proof is no longer valid. This simple example is given to dramatize the difference between the usual and octonionic quark theory, the novel aspects of which deserve a careful investigation.

The octonionic approach also brings a new light on the choice of flavour groups. Octonions are associated with representations of exceptional Lie algebras (e.g. G_2 is the automorphism group of the octonion algebra, for other groups the relation is more complex). For all exceptional simple Lie algebras the representations cannot be constructed in terms of usual matrices. To construct the representations of the exceptional groups G_2, F_4, E_6, E_7, E_8 the octonionic matrices (Jordan algebras) or their direct products (B.A.Rosenfeld algebras) must be used. For example the fundamental 27 dimensional representaion of E_6 can be realized in terms of the Jordam matrices which are 3x3 hermitian matrices with octonionic entries, etc. All exceptional groups contain G_2 as a subgroup which, as we have seen above, contain $SU(3)^c$. The group structure is as follows:

$F_4 \supset SU_3^c \times SU(3);$

$E_6 \supset SU_3^c \times SU(3) \times SU(3);$

$E_7 \supset SU_3^c \times SU(6);$

$E_8 \supset SU_3^c \times SU(3) \times SU(3) \times SU(3).$

Remark that all $G_{(\text{flavour})}$ have no abelian part. The fundamental representations of the most interesting group E_7 are:

$$\underline{56} = \underbrace{(\underline{20},\underline{1})}_{\text{leptons}} + \underbrace{(\underline{6},\underline{3}) + (\underline{6},\underline{3})}_{\text{quarks}} \;/(SU_6, SU_3^c)\text{-classification}/$$

As Q is the colour singlet, $Q \subset SU_6$ and we have $\sum_{\text{quarks}} Q = \sum_{\text{lept.}} Q = 0$. The charges of quarks are naturally fractional. If we suppose that there is the usual SU(3)-triplet of quarks u, d, s with $Q_u - Q_d = +1, Q_d = Q_s$, then the charges of all quarks are uniquely determined $Q_q = (\frac{2}{3}, \frac{1}{3}, \frac{1}{3}, \frac{2}{3}, \frac{1}{3}, \frac{1}{3})$ and similarly the charges of leptons $\left((\text{lept.}) \underset{SU_6}{\sim} (6 \times 6 \times 6)_{(\text{antisymm.})}\right)$ must be $(\pm 1 \pm 1 \pm 1 \pm 1$; 12 neutrals). If we introduce gauge bosons (which is by no means necessary) we have the adjoint representation

$$\underline{133} = \underbrace{(\underline{35},\underline{1})}_{\substack{\text{gauge}\\ \text{bosons}}} + \underbrace{(\underline{1},\underline{8})}_{\text{gluons}} + \underbrace{(\underline{15},\underline{3}) + (\underline{15},\underline{3})}_{\substack{\text{leptoquark}\\ \text{bosons}}}$$

It can be shown that theories based on exceptional droups /12/ are welcomed by recent experimental discoveries /8/ , but we will not dwell upon this. We only stress that among "simple" (mathematical term!) exceptional groups only E_7 is a plausible candidate (giving more than 3 quarks and large enough number of leptons, but not containing colour octet quarks). Hence, there must be six and only six quakrs -NO MORE QUARKS!

2.4. Alternatives

There are other approaches to the flavour symmetries not considering the colour conservation as a fundamental law. Pati and Salam /19/ introduce the following fundamental fermions: $\ell = (\nu_e, e^-, \mu^-, \nu_\mu \,;\; L=1,\ B=0)$ $\mathscr{B} = C_R^0, C_B^-, C_G^-$ ($L=1$, $B=1$, coloured) and S^0 ($B=1$, $L=0$, colourless). The quarks are made of leptons ℓ and prequarks $\mathscr{B}$, S^0: $u_R = \nu_e \bar{C}_R^0 S^0$, $d_R = e^- \bar{C}_R^0 S^0$, $S_R = \mu^- \bar{C}_R^0 S^0$, $C_R = \nu_\mu \bar{C}_R^0 S^0$ etc. Their integer-charge quarks are unconfined and unstable, with the lifetime $\sim 10^{-11}$-10^{-12} sec and hence can be found in emulsion experiments. This scheme is rich enough and flexible enough to be compatible with present experiments.

There exist suggestions to revive the Goldhaber model (Lipkin /20/) and F-Y-S model (Tyapkin /21/) which are not yet elaborated to be confronted with experiment. Finally, we note that some dualists propose the infinity of quarks and the corresponding additive quantum numbers (D.V.Volkov [22]). Others, however, insist on finite number of quakrs [16] . We cannot go into discussion of these ideas.

3. Dynamics of quakrs

3.1. Independent quarks

As has been mentioned above, the simplest versions of the quark dynamics in which quarks are supposed to be almost freely moving inside a sphere gives very good quantitative results. Essential ingredients of this dynamics ("quasi-independent quark" model, which we now call the "Dubna bag") proposed by N.N.Bogolubov e.a. (see [3]) and further developed and improved by P.N.Bogolubov [23] are as follows. The quark mass outside the sphere (bag) is very large (or infinite) and it is very small inside the bag ($\sim \frac{m_M}{2}$ for mesons M and $\sim \frac{m_B}{3}$ for baryons B). The magnetic moments of baryons are explained by small effective masses of quarks inside the bag. The important new results which could not be obtained in the phenomenological SU(6) approach are: the magnitude of μ_p , the correct results for G_A/G_V , and $\langle r_p^2\rangle_{e.m.}$ [23] . Of special importance is the good result for G_A/G_V (instead of 5/3 from SU(6)). At this point the relativistic corrections of the order of $\frac{\langle \vec{p}_q^{\,2}\rangle}{m_q^2}$ or the contribution of the orbital motion of quarks inside the bag are crucial. The explicit expressions are roughly the following $G_A/G_V = \frac{5}{3}\langle\sigma_z\rangle$, $\mu_p = 3\left(\frac{e}{2m_p}\right)\langle\sigma_z + L_z\rangle$. As $\langle J_z\rangle = \langle\frac{1}{2}\sigma_z + L_z\rangle = \frac{1}{2}$ we can express G_A/G_V and μ_p in terms of $\langle L_z\rangle$. Describing the bag by a scalar spherical well potential (the cavity) and finding the Dirac wave functions of quarks moving inside the bag, it can be found that $G_A/G_V \simeq 1.1$, $\mu_p \simeq 2.5$, $\langle r_p^2\rangle \simeq 0.43/m_\pi^2$ [23] . We do not discuss the applications of this approach to mass formulae where the results are similar to other models (e.g., the nonrelativistic models with the oscillator potential [24]). For further discussion see papers by P.Bogolubov and B.Struminsky (these proceedings).

Relativistic corrections are also of importance in weak and electromagnetic decays of hadrons involving the annihilation of quarks ($p \to \ell\nu$, $\nu \to e\bar{e}$). The naive non-relativistic treatment of such processes has led to a rather paradoxical conclusion that e.g. $\left|\frac{\Psi_{q\bar{q},\pi}(0)}{\Psi_{q\bar{q},K}(0)}\right| \sim \frac{m_\pi}{m_K}$ in contrast to the supposed SU(3) symmetry of the $q\bar{q}$ potential. These processes were first treated by Matveev, Struminsky and Tavkhelidze [25] , very clear and comprehensive representation of the quark model results for different decays was given by Van Royen and Weisskopf [26] . The resolution of this paradox lies also in relativistic corrections (Struminsky, Llewellyn-Smith [27]). In relativistic theory, based e.g. on the Bethe-Salpeter equation, there are no apparent paradoxes with SU(3)-symmetry. The main idea of calculating the processes

$$h \to h'(e^+e^-),\ h \to h'(\ell\nu),\ h \to h'\pi(K)$$

consists in taking into account only one-quark transitions. This assumption is a generalization of the well-known Okubo-Zweig-Iizuka (OZI) rule and it was successfully applied for the mentioned hadronic decays. The calculations of L.Soloviev, Anisovitch et al.; Thirring, Becchi Morpurgo et al. used non-relativistic approximation. For a very clear and comprehensive presentation of these and other results of N.R.Q.M. see [28] .

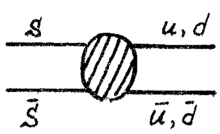
The consistent relativistic approach based on the Logunov-Tavkhelidze quasipotential equations for relativistic bound states was developed by N.Bogolubov, Matveev and Tavkhelidze [29] . The essence of their approach is the calculating of the moments of the currents V_μ^i, A_μ^i between bound states. With this aim they introduce external field v_μ^i, a_μ^i

interacting with corresponding currents and consider the case of small, slowly varying external fields. Then the variations of the energy of the bound state with respect to external field gives the matrix elements of the currents. This method enables one to reproduce all the nice results of the model of quasiindependent quarks and to obtain more general results. For example,

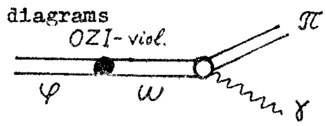
$$\frac{G_A}{G_V}=\frac{5}{3}\left(1-\frac{\langle \vec{p}_q^{\,2}\rangle}{3\,m_q^2}+\dots\right),\quad \mu_P=\frac{e}{2m_q}\left(1-\frac{\langle \vec{p}_q^{\,2}\rangle}{6\,m_q^2}+\dots\right).$$

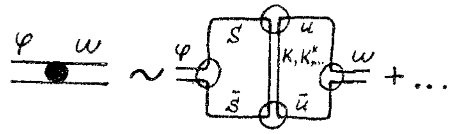
Similar results were derived by Gell-Mann, using the algebra of "good" components of currents in the $p_z\to\infty$ frame /30/ and by Shelest et al. /31/ in the framework of relativistic bound-state equations supplemented by the Markov-Yukawa condition /55/. Using PCAC one can calculate the processes $h\to h'\pi$ etc. With different binding "potentials" (e.g., square well, oscillator) the more detailed predictions can be obtained. However, it is difficult (if not impossible) to describe all existing data by a single potential, and the introduction of some phenomenological parameters is necessary. This is the essence of the so-called current -constituent quarks approach to hadronic decays summarized at the London conference /1/.

The essential ingredient of all these calculations is the OZI-rule, used to describe the different construction of the PS-multiplet and V (or T) multiplets. For the vector (tensor) mesons it is supposed that the process

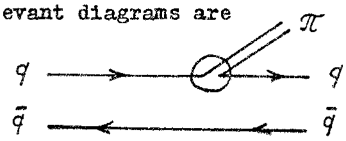


is very small while for PS-mesons it is appreciable. Then φ, f' are almost pure $s\bar{s}$-states, ω, f almost pure $\frac{u\bar{u}+d\bar{d}}{\sqrt{2}}$ states, while η and η' are certain mixtures of $s\bar{s}$ and $\frac{u\bar{u}+d\bar{d}}{\sqrt{2}}$. The effect of mixing is easily taken into account by considering e.g., the diagrams

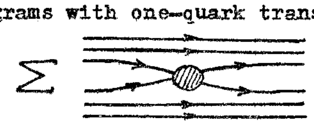




For pure hadronic decays the only relevant diagrams are



If we consider the scattering process then, as first suggested by Levin and Frankfurt /32/ the process is described by the sum of the diagrams with one-quark transition.

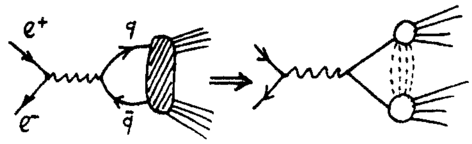


This is a generalization of OZI-rule for scattering processes. The predictions of the model are in reasonable agreement with experiment but show systematic deviations which can be naively interpreted as the result of somewhat smaller radius of the strange mesons /33/, we will discuss some related considerations in meson spectroscopy in what follows.

The simple-minded approach of independent quarks moving inside a cavity is supported by ideas of quark-parton models of Feynman, Bjorken, Weisskopf, Kuti et al. /34/ giving very clear and good description of all existing data on deep inelastic scattering of leptons on nucleons. The only cloud in this clear picture is the fact that the quarks are carrying only a half of the total momentum of the nucleon. The missing momentum is usually attributed to fashionable "gluons", but this is only another way of stating our essentially incomplete understanding of the quark-parton structure of the hadrons. The results of comparison of $\nu N\to\dots$ and $eN\to\dots$ are in good agreement with the fractionally charged quarks. However, as was argued by Salam, Pati, Roy and Rajakesran (see [35]) in theories with unconfined colour there exists a possibility that the quark charge sum rule is also consistent with integrally charge quarks.

Still other very impressive prediction of simple quark model is the quark counting rule for exclusive processes with high transverse momentum (Matveev, Muradyan, Tavkhelidze; Brodsky, Farrar [36]). Not going into discussion of its nature we only mention that it rests on the assumption that the wave function of constituent quarks is finite when all quarks are at the same point, i.e., $\Psi(x=0) = \int d^4k\, \tilde{\Psi}(k) < \infty$. Hence, in such processes the short-distance behaviour of the bound-state wave function is directly proved. An interesting problem is to extract from the scattering data some information concerning the wave functions.

Remind finally the application of the quark-parton model to the reaction $e^+e^- \to$ hadrons

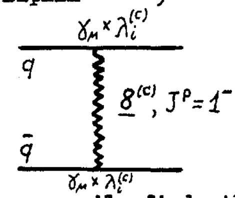


with clearly visible in SLAC experiment " jet structure", corresponding to $q\bar{q}$ pair. This dramatizes the mechanism of hadron production through the intermediate state of two quarks.

3.2. Dynamical role of colour: confining independent quarks.

Here we discuss some other ideas about the dynamical role of colour. The radical octonion approach is attractive, but, even if it is correct, it only gives a new frame for dynamics. It is also probable that there are different ways leading colour to confinement, and we are free to choose one that provides us with the simplest understanding of the hadron phenomenology.

The main facts of baryon spectroscopy can be explained in remarkably simple terms. Considering the SU(3)-invariant potential, corresponding to the exchange of colour-octet vector mesons (Nambu; O.Greenberg and Zwanziger; Lipkin [37])



$$\mathcal{L}_{int} = g\bar{q}\gamma_\mu \lambda_i^{(c)} q V_\mu^i$$

one easily finds that the effective coupling constants in different channels are

channel	qq	qq	$q\bar{q}$	$q\bar{q}$	qqq	...
SU(3)	$\underline{3}^c$	$\underline{6}^c$	$\underline{1}^c$	$\underline{8}^c$	$\underline{1}^c$	
$g^2_{eff.}$ (arb. units)	−4	+2	−8	+1	−12	

If the free quarks are heavy ($m_q \to \infty$) then the colourless states are, on this scale, massless, i.e., "confined". From the above table one can infer that the mass of the colour-triplet diquark state is of the order of the quark mass. This makes the successful quark-diquark picture of the baryon [125] quite naturally emerging in this approach.

The colour-exchange vector potential was originally used for integrally charged quarks. If we take the colour-singlet charge operator (hence, fractionally charged quarks) and suppose that coloured vector mesons are massless Yang-Mills mesons, corresponding to the exactly conserved SU(3) gauge group, we arrive at "Quantum Chromodynamics" QCD (for refs. see e.g., [10]) (Gell-Mann, Weinberg et al.). As is by now well-known, such theories enjoy asymptotic (ultraviolet, U.V.) freedom (A.F.), i.e., the effective coupling constant $g(p^2)$ is vanishing for $p^2 \to \infty$.

It is generally believed that in asymptotically free theories there is a good reason to rely upon perturbation theory. Even if this is true for the estimate of the small-distance behaviour of the coupling constant, the perturbative results for quark scattering amplitudes and bound states cannot be trusted. For example, consider the Logunov-Tavkhelidze or B.-S. equation with asymptotically-free potential. It is not difficult to demonstrate that for small r it can be reduced to a Schrödinger-type equation with the potential of the form $V(r) \underset{r\to 0}{\approx} g^2 r^{-2}\left(1+g \ln \frac{r_0}{r}\right)^{-\alpha}$, $\alpha > 0$ /38/. The solutions of this equation have an essential singularity in the g-plane for which cannot be traced in perturbation theory. In addition, the scattering amplitude has an essential singularity in the ℓ-plane ($\ell = 0$ for spinless particles). In the non-relativistic theory the singularity in the ℓ-plane has later been investigated by Oehme et al./39), who treated the simpler potentials: $V(r) = g^2 r^{-2}\left(\ln \frac{r_0}{r}\right)^{-\alpha}$, $\alpha = 1, 2$. Some general arguments in favour of the existence of the g-plane essential singularity were recently given by Shirkov /40/ who considered the theories with the Landau-Pomeranchuk null-charge phenomenon, using the renormalization group equations and a spectral representation for the invariant charge. It would be of great interest to extend his analysis to A.F. theoreis. The moral of this discussion is: using perturbation theory arguments for investigating U.V. behaviour in A.F. theories requires some caution.

The attitude of the QCD proponents to the infra-red (I.R.) behaviour of the theory with massless gluons is strikingly different. In fact, it is proposed to entirely disbelieve perturbation theory and to search for some peculiar non-perturbative solutions. The reason for this lies in the fact that the QCD can make sense only if coloured states are confined. So it is supposed that the effective interaction constant $g(p^2)$ which tends to 0 for $p^2 \to \infty$ (small distance), infinitely grows for $p^2 \to 0$ (large distance) providing the desirable confinement. Hence, the slogan of chromodynamists (Wilson, Susskind, Kogut; A.A.Migdal, Polyakov et al., for review and refs. see e.g. /41/): "Ultraviolet Freedom is Infrared Slavery".

The existence of the confined phase was demonstrated for some lattice theories as long as the lattice constant R was kept finite. The principal questions are; 1) is confinement preserved when $R \to 0$; 2) is the confined phase stable against external perturbation? 3) does the confinement work for all sectors of the Hilbert space and for all energies? (remind that the structure of the Hilbert space and observables in such a theory would be very unusual). Up to now there are no convincing answers to these questions. The recent progress in understanding the confinement of quarks in lattice gauge theories has been reviewed at this conference by Wilson /41/ . An attempt to construct a theory of composite hadrons in such a theory is presented at this conference by Bardeen /42/ . The light-front formulation (to be discussed below) of the QCD is used with a transition to the lattice variables in the transverse direction. Supposing the existence of a new phase of the theory (not realized in the usual perturbative solution), in which the transverse gauge invariance is exact, an attractive theory of composite hadrons can be formulated. However, the proof of the existence of such a phase is still lacking, the proof probably could be given for the finite lattice constant but there is no idea how to pass to the continuum limit.

The main difficulty of the confinement theories lies in that the perturbatien theory exhibits no hints for finding the confinement mechanism (Appelquist et al. /43/). The I.R.behaviour in nonabelian gauge theories seems to be very similar to that of QED. Alternative calculations were performed by Cornwall and

Tiktopoulos. They claim that the summation of leading logarithms gives confinement, i.e., $\lim_{\mu\to 0} \sigma(\mu) = \lim_{\mu\to 0} \sum_{leading} \sigma^{(n)}(\mu)$ is zero for some matrix elements $\sigma(\mu)$ with colour creation. Even if these (very difficult) calculations are technically correct, we cannot rely upon this result. Some time ago it was shown (Arbuzov et al. /44/) that when you sum up a logarithmic series which in this case is almost certainly not convergent (most probably it is the asymptotic series) the summation of the leading logarithms usually gives a result which has very little in common with the exact sum. Therefore, unlike the U.V.freedom, the I.R.slavery is not in good position.

Concluding this discussion we stress that the QCD is a very promising theory of hadrons composed of confined coloured quarks, even if we forget the most ambitious attempts to unify all the interactions /10/ and some interesting phenomenological applications to be discussed below. However, it can be regarded as a real theory (not merely a new "religion" of theorists) only after having answered the fundamental questions discussed above.

3.3. Relativistic bound states

A) Now we turn to equations describing relativistic bound states. The systematic approach to this problem has been developed by Fock and Podolsky in 1932 (F.-P). It is based on a three-dimensional one-time formulation of the bound-state equations. In the alternative approach (Dirac, Fock, Podolsky, 1932, D-F-P) an individual time variable is assigned to each particle (for further details and refs. see /45/). The first is not covariant while the second is obviously covariant. Both approaches have given many fruits. The F-P approach in the formulation of Tamm and Dancoff was applied to meson theory (Dyson, Low) and for nucleon-nucleon interaction (Klein, Levy, Macke). The development of the D-F-P approach resulted in covariant equations for bound states (Nambu; Salpeter, Bethe, Schwinger, Gell-Mann, Low) /47/ (Some relations between both approaches have been investigated by Zimmermann /46/ .

However, in both approaches serious difficulties were found. After having found the rules for calculating matrix elements of currents between bound states and the normalization conditions (Nishijima; Mandelstam) and with some exact solutions in the ladder approximation (Wick, Cutkosky; Okubo, Feldman; Goldstein; Nakanishi; Kummer, et al.) it became clear that the price for covariance was quite high-physical interpretation of the bound state solution is unclear /47/ . There exist solutions of the Bethe-Salpeter equation with negative norm (violating unitarity) and solutions with exotic J^{PC} not occuring in the nonrelativistic limit, e.g. for $N\bar{N}$ or $q\bar{q}$ system $J^{PC} = 0^{--}, 0^{+-}, 1^{-+}, 2^{+-}, \ldots$ are exotic and we never saw such states. Both phenomena have a common source - the presence of the relative time in the bound state equations. The states with negative parity with respect to the relative-time reflection may have negative norm and/or exotic J^{PC} /47/ . For example, the solutions of the B-S equation for the pion are of the two types: the normal solution $F_\pi \cong \gamma_5 f(p^2, (pK)^2)$ and the anomalous solution $F_\pi \cong \gamma_5 (pK) f(p^2, (pK)^2)$.
Due to a factor $(pK) = p_0 K_0$ for $\vec{K} = 0$ the last solution has $J^{PC} = 0^{--}$. All such anomalous states disappear in the nonrelativistic limit or in the equal time limit. However, as we have argued earlier, any reasonable theory of $q\bar{q}$ bound states must be relativistic. Besides this, the old equal-time formulations were extremely complicated in contrast to the relatively simple B-S equations.

B) The escape from these difficulties was found by Logunov and Tavkhelidze /48,45/ who proposed the quasipotential approach to quantum field theory which unifies the physical simplicity of the equal time formulation and the mathematical simplicity of the covariant formulation. The main idea is that only the on-mass-shell scattering amplitudes are relevant for calculating bound states, and the potentials are expressed in terms of these amplitudes. The simplest example of such an equation for equal-mass spinless constituents is, in coordinate space,

$$(\Delta+K^2)\,\Psi_{\bar{K}}(\bar{\tau})=(m^2-\Delta)^{-1/2}\, g\,\frac{V(\tau,K^2)}{\tau}\,\Psi_{\bar{K}}(\bar{\tau})$$

where $\bar{K}^2=\frac{1}{4}M^2-m^2$, m – the constituent mass, M — the bound state mass $(S=M^2)$ $\bar{\tau}$ — the three dimensional relative coordinate, $\bar{\tau}=(\bar{\tau}_q-\bar{\tau}_{\bar{q}})$ in the CMS. This is the Fourier transform of the momentum-space equation of the form

$$\varphi_{\bar{K}}(\bar{p})=\int\frac{d^3q}{\sqrt{m^2+q^2}}\;\frac{g\,\mathcal{V}((\bar{p}-\bar{q})^2,K^2)\,\varphi_{\bar{K}}(q)}{\bar{K}^2-\bar{q}^2+i0}$$

Formally it can be derived from Lippman-Schwinger (L-S) equation by substitution $d^3q\to\frac{d^3q}{\sqrt{m^2+q^2}}$

This equation was generalized for unequal-mass case and for bound states of particles with spin 1/2, and was successfully applied to numerous problems in elementary particle /45/ and nuclear physics /49/ . Mathematically, the Logunov-Tavkhelidze (L-T) equation is somewhat more complicated than the non-relativistic Schrödinger equation. Nevertheless for large classes of the potentials different mathematical methods can be successfully employed for solving bound-state problems /50,51/ . The L-T equation has the correct small distance behaviour in the sense that it correctly reproduces the small distance behaviour of the wave function and the singularities of the scattering amplitude in the l-plane (poles and cuts) obtained in the corresponding field theories. For example, if we calculate the quasipotential in any given quantum field theory by using the perturbative expansion, the resulting L-T equation for the scattering amplitude $M(S,t)$ correctly reproduces each order of perturbation theory, and the asymptotic behaviour of $M(S,t)$ for $S<0,\ t\to\infty$ coincides with that of the sum of the corresponding ladder diagrams /52/ . Moreover, the differential formulation of the L-T equation /50,51/ can be used for finding scattering amplitudes and bound states in non-renormalizable theories when perturbation theory is inapplicable /51,53/.

A large class of other quasi-potential equations can be obtained by substitution

$$(m^2-\Delta)^{-1/2}\to f\left((m^2-\Delta)^{1/2},(m^2+K^2)^{1/2}\right)(m^2+K^2)^{-1/2},$$

where $f(x,x)\equiv 1$. All these equations, like L-T equation, satisfy two-particle unitarity and, with energy-dependent potential $V(\tau,K^2)$ can incorporate many-particle unitarity. A rather simple equation useful for the description of tightly bound states can be obtained with $f=\frac{m^2+K^2}{m^2-\Delta}$ (Filippov, see also /50/). Then the bound state equation

$$(m^2-\Delta)(\Delta+K^2)\Psi_K(\tau)=\sqrt{m^2+K^2}\,g\,\frac{V(\tau,K^2)}{\tau}\Psi_K(\tau),$$
$$\sqrt{m^2+K^2}\equiv M/2$$

is of the fourth order (with four boundary conditions) and in it the effective potential automatically vanishes for $M\to 0$ (Correspondingly, for not very singular potentials there are no $M=0$ bound states). Most of other formulations of quasipotential equations as well as of the B-S equation have the difficulty that with any deep enough attractive potential the mass of the bound state is imaginary.

Another variation on this theme is the equation proposed by Todorov /45/

$$(\Delta+K^2)\Psi_K(\tau)=(m^2+K^2)^{-1/2}\,g\,\frac{V(\tau)}{\tau}\Psi_K(\tau).$$

This equation is useful in calculating high-energy ($K^2 \to \infty$) behaviour of the elastic scattering in $\lambda\varphi^3$- type theories but it is singular for $K^2 \to -m^2\ (M \to 0)$ and the potential $V(\tau)$ is more singular for $\tau \to 0$ than in the L-T equation (in the above case of the fourth-order equation the potential is effectively less singular). These modifications are mathematically simpler, than the L-T equation, and there are rather general methods allowing one to obtain some analytical solutions for more or less simple potentials /50,51/ .

There have been proposed many other modifications of the original L-T equations (Kadyshevsky, Gross, Thompson, Fronsdal, Todorov, Yaes, Klein et al.) /45,49/ . They all differ either in the choice of the f - function or in the choice of the propagator: $(\Delta + K^2)^{-1} \to \ldots$ A freedom in the choice of the quasipotential equation (QPE) corresponds to a freedom in extrapolating the scattering amplitude off the mass shell.

C) The quasipotential equation can be also rewritten in an explicitely covariant form (Matveev, et al. /54/ , see also /29/):

$$\varphi_K(p) = \int d^4q\ \delta(n \cdot q)\ \frac{g\,\mathcal{V}(p,q;K^2)}{K^2/4 + p^2 - m^2}\ \varphi_K(q)\,;$$

$$K \cdot p = 0\,, \quad n_\mu \equiv K_\mu/\sqrt{K^2}.$$

The condition $k \cdot q = 0$ was earlier used by Markov and later by Yukawa (M-Y) /55/ for exluding from the theory the dependence on relative time. (It can also be used for choosing the solutions of B-S equations having a finite nonrelativistic limit). The M-Y condition was originally invented as a mathematical device for a consistent treatment of a bilocal theory of composite particles. These ideas were recently revived by several authors (see e.g. /56/). The physical consequences of a bilocal quark theory of hadrons, based on equations somewhat intermediate between quasipotential and bag equations, are most detailly elaborated by Preparata /57/ . The main idea of the Preparata approach is to entirely exclude the quark variables from the physical quantities, the only dynamical trace of the quark structure being supposed the bilocal nature of the hadron fields. This is in parallel with attempts to use the bilocal current algebra instead of more specific assumptions of the parton model (see e.g., /34/). It is well known, that the predictive power of the bilocal current algebra is somewhat weaker than that of the parton model /34/ . Similarly, the bilocal quark theory of hadrons, reproducing many nice results of the constituent quark models, fails to give any definite prediction in several important points (e.g., for $R = \frac{\sigma(e^+e^- \to hadrons)}{\sigma(e^+e^- \to \mu^+\mu^-)}$ /58/, it is also difficult to imagine a simple explanation of the jet structure in $e^+e^- \to$ hadrons). The main advantage of the bilocal theories over constituent theories lies in avoiding the quark-parton paradox (also in parallel with the bilocal current algebra). Our point of view is that a resolution of the paradox can only be found on a more fundamental level (octonions?) and that the quarks, while not existing as free particles, can otherwise be regarded as real particles of which the hadrons are composed. For these reasons we concentrate in what follows on the quasipotential quark models and on bags.

Return now to the covariant quasipotential equation. Generalizing the M-Y condition we can replace K_μ by some vector λ_μ . If this vector is light-like $(\lambda^2 = 0,\ \lambda = (1, 0, 0, \pm 1))$ we arrive at the simplest light-front (L-F) formulation of the QP equation. Several forms of such an equation are presented to this conference /59/ (see also /60,61/). There also exists an extensive literature on closely related approach of infinite-momentum-limit bound-state equation (see e.g. /63/). Here the people depart from Weinberg's formulation of the quasipotential equation in the infinite momentum system [63]. The common feature of all these approaches is to describe the bound states in the light-front system (or in the infinite-momentum system).

D) Why the LIGHT-FRONT?

The L-F dynamics has been discovered by Dirac (1949) but until recently it was practically unknown to physics community. Later Fubini and Furlan [64] realized that the current algebraist's life is much more comfortable in the infinite momentum "frame" which essentially coincides with the L-F "frame". As the partons can live only in such a system, it is now most popular among theorists. The experimentalists gradually approach this system with growing available energies in CMS.

If we boost any Lorentz system in z - direction then for $p_z \to \infty$ the most natural variables are the light-front variables $t+z, t-z, x, y$ Denote them as

$$x_+ = \frac{t+z}{2}, \; x_- = \frac{t-z}{2}, \; x_\perp^1 = x, \; x_\perp^2 = y$$

and consider the variable x_+ as a substitute for the time variable t. The classical dynamics in such variables is not simple.

E.g. $(\Box + m^2)\varphi = 0 \Rightarrow (2\partial_+\partial_- - \partial_\perp^2 + m^2)\varphi = 0$ and the initial value problem on the surface $x_+ = 0$ is known by mathematicians as being "incorrect" (infinitely many solutions). However, if we require that these solutions correspond to the finite energy, there would be no ambiguity in finding such solutions [65].

The L-F quantum theory for finite degrees of freedom is not much different from usual (as realized by Dirac). However, for infinite degrees of freedom (QFT) the new theory is radically different. Due to the existence of the positively definite conserved operator $p_0 + p_z$ (for particles with $m \neq 0$, $p_0 = \sum\sqrt{\bar{p}^2+m^2} > \sum p_z$ the interactions do not produce particle-antiparticle pairs, and one can hope to avoid such frightening theorems as Haag's theorem and Coleman's theorem [65,61]. Let us explain this point in some detail (for details and refs. see [66]):

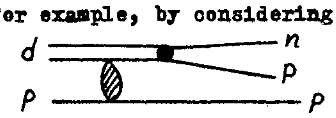
1) For systems with finite degrees of freedom we have very nice von Neumann's theorem that all irreducible representations of the commutation relations $[a_i(\bar{p}), a_j^\dagger(\bar{p}')] = \delta_{ij}\,\delta(\bar{p}-\bar{p}')$ are unitary equivalent, and so we can define the physical vacuum such that $a_i(\bar{p})|0\rangle_{ph} = a_i(p)|0\rangle_{bare} = 0$.

2) For systems with infinite degrees of freedom, according to Haag's theorem $|0\rangle_{ph.} = |0\rangle_{bare}$, and even free fields with different masses are unitary inequivalent! Stated in other way, this means that $S(t,-\infty)|0\rangle_{bare}$ does not exist, due to the pair production from the bare vacuum. In contrast, $S(t+z,-\infty)|0\rangle_{bare}$ probably may exist, due to the conservation of $p_0 + p_z$ forbidding the pair production.

Similarly, Coleman's theorem ("the invariance of the vacuum is the invariance of the world") is not true in the L-F variables, as the vacuum is stable under the action of any charge operator: $Q_i|0\rangle = \lambda_i|0\rangle$

If the charge is not conserved, then $\lambda_i = 0$ and Q_i annihilates the vacuum. We see that the vacuum is always automatically invariant. This fact is especially useful in considering dynamical realization of chiral symmetries [67]. These and other conceptual advantages of the L-F variables give us all the reasons for using the L-F dynamics in relativistic bound state theory. There are also several practical advantages of these variables: 1) The stability group (the little group) of the light-like system, $(E_2 \times \mathcal{D}) \times T_3$, is larger than the corresponding stability group of "space-like" systems, $SO_3 \times T_3$, and the first has the Galilei group as a subgroup. For this reason the dynamics in L-F system is (somewhat paradoxically!) very similar to nonrelativistic dynamics /68/ . 2) It follows that the bound state equations in this system must be of the three-dimensional quasipotential nature /63/ . 3) The concept of the L-F variables proved to be very useful in parten model, in the light-cone current algebra, and it revealed its practical advantages in treating deep-inelastic processes. To unify these semiphenomenological approaches with quark-bound-state models is hardly possible without using the L-F formalism.

Some preliminary attempts in this direction were presented to this conference. For example, the quark-counting rules are naturally emerging from the quasipotential equations in LFV (Garsevanishvili et al. /69/ , Brodsky et al. /70/ , Khelashvili /71/ , Kvinikhidze /72/). A connection of large and small momentum behaviour of meson form factors is discussed by Terent'ev /73/ . An interesting field of application of light-front formalism is in high-energy hadron-nucleus reactions where the constituents (nucleons) are unconfined. For example, by considering $pd \to ppn$



one can directly measure the deuteron wave function $\varphi_d(x, \bar{p}_\perp^{sp})$, as the differential cross section is of the form /74/

$$\frac{d^2\sigma}{d p_\perp^{sp} dx} \sim \sigma_{el}(pp) \left| \varphi_d(x, \bar{p}_\perp^{sp}) \right|^2$$

In such a way the equations describing bound states of nucleons in extremely nonrelativistic situations can be confronted with experiment. Similar equations can be used for other relativistic bound states (e.g., for positronium, Faustov et al. /75/). Relativistic nuclear physics and atomic physics provide us a very promising field for applications of the described formalism. Here we know the constituents and have a very good knowledge of the binding forces. Confronting theoretical results with experiment we can probe our ideas on relativistic bound states.

E) Some other approaches to relativistic bound states are based on the Kadyshevsky formulation of the relativistic Hamiltonian quantum field theory which elegantly generalizes the non-covariant perturbation theory /76/. Starting from this formulation a class of quasipotential equations was derived. The simplest one can be obtained (formally!) from the Lippman-Schwinger equation by substitutions

$$E_p = m + \frac{\bar{p}^2}{2} \to E_p = \sqrt{\bar{p}^2+m^2}; \quad d^3q \to \frac{d^3q}{\sqrt{m^2+q^2}}$$

and reads

$$\varphi_{\bar{k}}(\bar{p}) = \frac{1}{2} \int \frac{V(\bar{p}-\bar{q}, k^2)\, \varphi_{\bar{k}}(\bar{q})}{E_k - E_q + i0} \cdot \frac{d^3q}{\sqrt{m^2+q^2}}$$

By means of a transition to the relativistic coordinate space it can be transformed into a differential-difference equation. The relativistic coordinate space /76/ is related to the momentum space through the Shapiro /77/ transformation

$$\varphi_{\bar{k}}(\bar{p}) = \int d^3\tau \left(\frac{p_0 - \bar{p}\bar{n}}{m} \right)^{-1-i\tau M} \Psi_{\bar{k}}(\bar{\tau}), \quad \bar{n} \equiv \frac{\bar{\tau}}{\tau}$$

This is the natural generalization of the usual Fourier transformation used in nonrelativistic quantum mechanics. In fact, the function entering this expression is the relativistic form of $e^{i\bar{p}\bar{\tau}}$ and its limit for $m \to \infty$ is exactly $e^{i(\bar{p}\bar{n})\tau} = e^{i\bar{p}\bar{\tau}}$.

An interesting variation of this theme is presented at this conference by Mir-Kasimov et al. /78/ . It is suggested to use the rapidity variable instead of E_q: $\chi_q = \ln(E_q + \sqrt{E_q^2 - m^2})$. Then $(E_k - E_q)^{-1} \to \frac{\chi_q}{sh\,\chi_q}(\chi_q^2 - \chi_k^2)^{-1}$ and the resulting equation in the coordinate space is a second order differential equation similar to the Schrödinger equation. For the S-wave it reads

$$\frac{d^2\Psi_k(\tau)}{d\tau^2} + \chi_k^2 \Psi_k(\tau) = \frac{\chi_k}{sh\,\chi_k} V(\tau) \Psi_k(\tau)$$

This equation is as simple as the Schrödinger equation. The extension of this approach to spinor particles and to the unequal-mass problems is very desirable .

Starting from the Kadyshevsky Hamiltonian formulation of quantum field theories a somewhat different approach to relativistic equations for bound states can be developed. In this formulation momenta of all particles belong to the mass shell (as in nonrelativistic theory). The formalism is particularly convenient for constructing the Fock space in the light-front variables. The corresponding equations for n-particle bound states were considered by Karmanov /79/ . It would be interesting to investigate such equations in detail and to extend them to spinor particles. A new feature of this equation is a dependence of the Fock amplitudes on a unit vector which is somehow related to $\bar{P}/|\bar{P}|$, $|\bar{P}| \to \infty$. However, the necessity of these new parameters and their meaning is not completely clarified. The consideration of some physical problems would be most instructive.

A simpler approach to the relativistic two-particle bound states without extra variables is applied by Terent'ev /73/ to different problems of the relativistic quark model, especially to radiative decays of mesons. Equations used by him are similar to quasipotential equations, constructed earlier by Sokolov /60/ , who starts from Dirac's formulation of the relativistic Hamiltonian theory. This approach avoids using the quantum field theory and only deals with the generators of the Poincare group. The quasipotential is introduced phenomenologically, the theory is only giving us a prescription for doing this in a covariant way. Sokolov's method can also be applied to many-particle bound-state problem /60/ , however, the practical realization of this possibility is not yet elaborated.

The methods described above do not allow one to specify the binding potential, and it should be extracted from some field theory or be somehow guessed. We discuss several popular potentials in the next section. Here we mention a possibility of constructing the two-particle relativistic bound-state theory, in which the "potential" is completely defined by the physical scattering matrix of the constituents. As proposed by Logunov, Khrustalev et al. /80/ the relativistic generalization of the Heitler-Sokolov-Wilson equation /81/ can be obtained in the framework of the equal time formulation of the relativistic two-body problem in QFT. The corresponding "quasipotential" is expressed in terms of the elastic and inelastic cross sections of the two constituent particles. This method has been successfully applied to the description of the high-energy two-particle scattering. It is potentially useful for describing the two-particle bound states. In relativistic atomic and nuclear physics and in hadron phenomenology it allows one to take phenomenologically into account many-particle contributions to scattering and to bound-state energy. In the quark theory of hadrons its application is less justified as some knowledge of the quark scattering amplitudes is required.

F) We have mostly reviewed above the quasipotential type formulations of the relativistic bound-state problem. There are presented to this conference a few more conventional treatment of the problem. Cung et al. /82/ summarize the results of their investigations of the two-fermion B-S equation with the kernel restricted to the zero relative time (static interaction). The approach is essentially equivalent to the quasipotential approach of Faustov and Todorov /45/ . In the paper presented by Ladanyi /83/ the small distance behaviour of the B-S equation for bound states of a fermion with a massive vector meson is investigated (see also Ciafaloni and Ferrara /83/). For similar (but more relevant to the quark model) investigations see e.g./84,85/.

Note that the authors of ref. /85/ start from the B-S equation but subsequently reduce it to the Logunov-Tavkhelidze equation to simplify the calculation of the asymptotic behaviour of the bound-state form factors. (Compare to /68-72/). Some other calculations of this behaviour appear to a related method of summing the "leading" contributions of Feynman diagrams /86,87/ . In this case the bound state is conveniently defined in terms of a pole in the angular-momentum plane. (See especially Efremov et al. /87/ , where the method is consistently used for investigating the asymptotic behaviour of form factors, and the validity of the quark counting rules. Finally we mention some diverse results in the theory of the B-S equation which are related to the problems discussed above. New exact solutions of spinor-spinor B-S equations are obtained in references /88/ . In ref. /89/ exact upper and lower bounds for the sum of scalar ladder diagrams are found. Glimm and Jaffe /90/ have given a rigorous proof of existence of two-particle and three-particle B-S kernels in the Euclidean region for a wide class of two-dimensional scalar theories. The structure of the three-quark B-S equation is poorly known. An investigation of the general spinor structure of the bound-state wave functions is attempted in /91/ (see also /61/).

Concluding this rather lenghthy and by no means complete discussion of present trends in the relativistic theory of bound states it is to be emphasized that up to now there is no formulation of the theory which is adequate for solving all problems, occuring in physical applications. For different problems we have to use different methods. In general, the L-F quasipotential equations seem to be most appropriate for describing relativistic bound system. However, the B-S equation is better suited for extremely tightly bound states (e.g., for zero mass bound states to be discussed in section 3.5).

3.4. Interquark forces

Once the equation is chosen the next question is: What is the (quasi)potential acting between quarks. The symmetry properties with respect to the colour and flavour groups have been discussed earlier: 1) The potential corresponds to the exchange of the colour-gauge bosons and most probably is colour-conserving. 2) It either is flavour-conserving or has a small symmetry violating term. The main flavour-symmetry violation is assumed to be attributed to the different masses of quarks. 3) As is argued in the next section it probably involves a piece corresponding to the exchange of flavour-gauge bosons. Such terms are desirable for spontaneously generating quark mass differences. As to the spatial dependence of the potential, the choice between different possibilities is much more difficult. We summarize here the most popular potentials together with new ones presented at this conference.

1) The "good old" oscillator potential (see e.g. /2,5,24,31/ $V(r) = \alpha r^2$ or the "bag-like" oscillator potential $V_{osc} = -V_0 + \alpha r^2$ (e.g. /73,92/). These are most popular due to availability of the exact analytic solutions for some of the bound-state equations mentioned above. The Regge-trajectories for these potentials are linearly growing with M^2 (or S). However, the form factors $F(q^2)$ of bound states have a pathological dependence on the momentum transfer q^2 and the predictions for excited states are unrealistic. Thus the oscillator potential can be considered only as an approximation to the "realistic" potential, which is only adequate for describing some properties of low-lying states of composite hadrons.

2) The QCD-potential $V_{QCD}(r)$ /10/

$$V_{QCD} \underset{r\to 0}{\sim} \frac{\beta(r)}{r}, \quad V_{QCD} \underset{r\to\infty}{\sim} \alpha r$$

Here $\beta(r) \underset{r\to 0}{\sim} [\ln(r/r_0)]^{-\alpha}, \alpha > 0$. This potential has a singularity at infinity. For $r \to 0$ it is singular if $\alpha < 1$ and regular if $\alpha \geqslant 1$ /38/. Such potentials have been applied to a description of the $q\bar{q}$ states and especially to the charmonium spectroscopy. For a summary of the corresponding calculations see the invited paper of Mir-Kasimov /93/ (see also /94/). In all these calculations either Schrödinger or some quasipotential equations have been used. As many questions to theory and experiment are yet to be answered, it would be premature to draw from these calculations any definite conclusion.

3) Some other confining potentials are discussed at this conference. Skachkov /95/, generalizing Kadyshevsky approach /76/, obtaines a quasipotential equation for the $q\bar{q}$-system with the potential $V_S(r)=(4\pi r)^{-1} ctg(\pi m r)$ where m is the mass of the quark. m is the only free parameter in the equation, and fixing it, say, by the requirement that the lowest state is the ρ-meson one can predict a sequence of the excited states ($M_{\rho'} = 1100$ MeV, $M_{\rho''} = 1465$ MeV etc.). Unfortunately, the quarks are supposed to be scalars and so the spin effects have not been discussed. Another attempt to confine quarks is presented by Guenin /96/, who "simply" changes the sign of the mass of the gluon ($\mu \to -\mu$) in the space-like part of the gluon propagator, thus arriving at the potential $V_G(r) \underset{r\to\infty}{\sim} r^{-1} e^{\mu r}$ (the quark propagator is not modyfied). It is not clear at the moment whether the corresponding theory remains causal. The phenomenological applications are not discussed.

4) Dolgov /97/ gives some arguments in favour of a double-well structure of the $q\bar{q}$-potential. He starts from the Blokhintsev et al. /98/ quasipotential equation and observes that the structure of the equation itself dictates a double-well form of the effective potential in the radial equation. It is possibly true for other quasipotential equations for spinor particles. This idea is attempted to be applied to explaining Ψ-particles without new quarks.

5) It is known for long that a spherically symmetric well potential gives nice phenomenological results in the quark model /23/. Such potentials naturally arise in an approximation to MIT-bag model (see sect. 3.1 and 3.6). Another source of similar potentials is the exchange of infinite number of resonances with an exponentially growing mass spectrum $\sigma(\mu) \sim e^{\mu a}$ /99/. As the exchange of one particle results in the Yukawa potential $r^{-1} e^{-\mu r}$, the exchange of $\int^{\infty} d\mu\, \sigma(\mu)$ particles gives rise to the potential

$$V(r) \sim \int d\mu\, \sigma(\mu) \frac{e^{-\mu r}}{r} \underset{r\to a}{\sim} (r-a)^{-1}$$

having a singularity at the finite distance from the origin (FDS-potential). Such potentials can be obtained in a non-polynomial field theory /99/, or in theories with infinite-component fields /100/. Properties of the pion have been investigated in the model, supposing the quark motion can be described by the Euclidean B-S equation with the kernel (potential) $V(r) = g^2(r^2-a^2)^{-1}$. The parameter a is fixed ($a = 3 \div 4$ GeV) by considering the empirical mass spectrum which in fact is exponentially growing up to 2 GeV. The remaining two parameters g and m_q are determined from the eigenvalue condition for the pion and from $\Gamma(\pi \to \mu\nu)$. The predictions for $\Gamma(\pi\to\gamma\gamma), \langle r^2\rangle_\pi$ and for the slope of the pion Regge-trajectory α'_π are in good agreement with experiment. The most remarkable prediction is the presence of oscillating terms $\sim t^{-6} \cos(\frac{a\sqrt{t}}{2})$ (in form factors $F_\pi(t)$ and elastic cross sections $\frac{d\sigma}{dt}(pp)$

for large space-like t . The period of the oscillations is predicted to be $\Delta\sqrt{t} \simeq \frac{4\pi}{a} \sim 3 \div 4\,GeV^{-1}$ in a striking agreement with the observation of Schrempp and Schrempp [101] . We are not aware of any other natural explanation of the oscillations in $\frac{d\sigma}{dt}(pp)_{el}$ found in [101] . Note that the Regge trajectories for the FDS--potential are approximately linearly growing with mass M (not M^2!). It is assumed that a faster (linear in M^2) behaviour will result from the contributions of inelastic channels opening for large M . The theory with energy independent potentials is supposed to be applicable only to low-mass hadron states. Finally, the FDS potential strongly confines quarks but the confinement is only partial[99] .

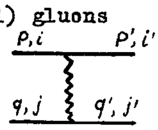
We discussed deverse coordinate dependences of the interquark potential for $q\bar{q}$-system. For choosing the most realistic one it is, first of all, necessary to consider the corresponding three quark potentials and to investigate the radially excited bound states of the three quarks. Very little has been done along this line (except for non-relativistic and simplest relativistic equations with $V_{osc}(r)$). To probe the radial dependence of the $q\bar{q}$- potential the decays and the radially excited states of mesons should be carefully investigated. Due to opening inelastic channels ($(q\bar{q}) \to (q\bar{q})(q\bar{q})$ etc.) this is (at least!) a many channel problem which has not been discussed in detail. In addition, the experimental status of excited mesons is rather unclear. We discuss a possibility to by-pass these difficulties in sect. 4.

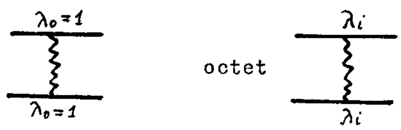
3.5. Chiral symmetry and quark masses

There are other difficult problems of the quark dynamics which have not been discussed above. 1) What is the origin of the flavour-symmetry violation (assuming the fundamental interaction is symmetry preserving)? 2) What is the origin of the approximate chiral symmetries (e.g. $SU(2)_R \times SU(2)_L$, $SU(3)_R \times SU(3)_L$ ...)? 3) What is the origin of the quark masses and of their differences? All these questions are obviously interdependent.

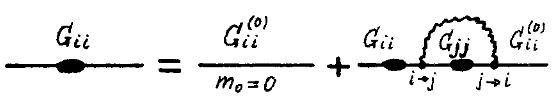
In semi-phenomenological theories it is usually assumed that the strong interaction of quarks is $SU(3)$ and $SU(3)_R \times SU(3)_L$ symmetric, and the observed symmetry-breaking effects are ascribed to the quark-masses. At a more fundamental level we have to investigate seriously the third question. One promising approach to this problem is based on the unified gauge theories of all interactions (see e.g. [10] and Slavnov's talk at this conference). Another, less ambitious one, is formulated within a semiphenomenological scheme of quark-quark interaction which simultaneously gives two apparently different effects: binding quarks and providing them with masses and mass splittings. This approach uses the mechanism of dynamical realization of symmetries which first has emerged in Bogolubov's theory of superfluidity [102] and subsequently has been applied to ferromagnetism, superconductivity, etc. The main idea is that the invariance of the Hamiltonian needs not to be the invariance of the ground state. To obtain such a solution we have first to remove the degeneracy of the Hamiltonian by adding some symmetry-breaking term. This symmetry breaking is switched off only after finding the desired solution. If there exists such a symmetry-breaking solution then, generally, there appear some zero-mass excitations (quasiparticles) which, in a sense, restore the original symmetry. The ground state contains an infinity of such quasiparticles (magnons, Cooper pair, etc.). These ideas in the statistical physics were first formulated by Bogolubov [103]. Their relevance to problems of elementary particle physics was discovered by Nambu and Goldstone [104] . We will call this approach the Bogolubov-Nambu-Goldstone realization of symmetry (BNG). Nambu also suggested to treat the pion as the massless particle corresponding to BNG-realization of the chiral symmetry $SU(2)_R \times SU(2)_L$. Examples of the quantum field theories with BNG-realization of the chiral symmetry $U(1)_R \times U(1)_L$ were first treated by Arbuzov et al. /105/ (two-dimensional) and by Nambu and Jona - Lasinio /106/ (four-dimensional).

Following this line of thinking consider /107/ the $U(n)_R \times U(n)_L$ symmetric theory (for definiteness consider n=3) of n massless quarks interacting through exchange of vector (or axial) gluons

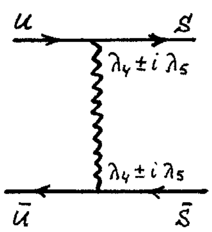
$V_{ij,i'j'}(p-p') =$ 



We treat this interaction as an effective potential (propagator). As we are not talking about three quark states we can be temporally "colour-blind". Then the equations for the propagator of quarks are of the form



For different potentials (e.g., for FDS-potential) these equations have solutions corresponding to $m_i \equiv m \neq 0$. If there is only the SU(3)-singlet interaction, then there are 9 massless pseudoscalar bosons. If there is also the SU(3)-octet interaction, then different possibilities arise due to strong mixing of the quark configurations



A very preliminary statement is that in this case only one pseudoscalar state remains massless, others can acquire a mass. One can also hope to arrange the relative singlet-octet coupling strengths so as to split the masses of quarks. This has been done in some simple models with factorizable partial wave potentials $V_\ell(p,q)$ (V_ℓ is the angular-momentum projection of $V(p-q)$) /107/ . This probably opens new way for solving the three distinguished problems. Unfortunately, "Things Take Time" . To demonstrate the consistency of this approach we have to do a lot of job: 1) to find a non-trivial symmetry-breaking quark propagators by solving the system of nonlinear equations with a realistic potential; 2) to find the solutions of the corresponding linear equations for pseudoscalar meson bound states (the B-S equations with the "exact" non-symmetric quark propagators); 3) to demonstrate that the $q\bar{q} \to q\bar{q}$ Green's function has corresponding poles and/or to incorporate in this scheme a confinement mechanism.

These problems are essentially unsolved even in the technically simpler "finite quantum electrodynamics" of Johnson et al. /108/ (for new results and refs. see /109/). In the paper presented at this conference Fukuda and Kugo /110/ attempt to solve the non-linear equation for the electron propagator introduced in refs. /108,111/ . They claim the propagator to have in the time-like region neither poles nor cuts for arbitrary e^2 and interprete this as a "confinement". The absence of the pole can be proved quite convincingly, but they give no proof of the absence of a branch-point singularity. As a matter of fact, expanding the self-energy part $\sigma(p^2)$ of their electron propagator in a series of powers of $\alpha = e^2/4\pi$ (which is convergent for small enough values of α) one can easily demonstrate that any approximation to $\sigma(p^2)$ has a branch point at $p^2 = -\sigma(0)$ It is rather difficult to understand why this singularity could completely disappear in the sum of the series. This sum most probably has a branch point either in the time-like region or in the complex p^2- plane (in a vicinity of $p^2 = -\sigma(0)$ for small α). Remark in passing that there exist suggestions (Dubničkova, Efimov /112/) to describe confined particles by "propagators" having no singularities in the complex p^2- plane except infinity (an integer function of p^2). An interesting question is: can such "integer" propagator naturally emerge in any quantum field theory? We think this problem has something to do with colour-confinement mechanism but a more serious discussion of this point is impossible at this moment. The propagator of ref. /110/ is almost certainly not an integer function.

Some other aspects of the BNG-realization of chiral symmetries are discussed at this conference. N.N.Bogolubov (Jr.) et al. /113/ investigate in detail the structure of the vacuum in the four-fermion theory of ref. /106/ by using Bogolubov's transformation. Kleinert /114/ and Pervushin and Ebert /116/ try to avoid the detailed discussion of the quark dynamics and to construct (without really solving the dynamical equations) a semiphenomenological theory which can be confronted with the usual $SU(3)_R \times SU(3)_L$ algebra of fields. This is achieved by "hadronizing" the quark interactions, i.e., by excluding the quark fields from the dynamics. This approach looks interesting but the important things must be clarified before we can reach some definite conclusions. Without solving dynamical equations the meaning of such approaches is not clear. In addition, some intriguing problems of the chiral quark theory - the η-η' mixing, the problem of the BNG-nature of pseudoscalar mesons (the so-called U(1) problem)are not touched upon in this approach. The U(1) problem can be formulated as follows. In any quark-gluon theory the chiral symmetry is $U(3)_R \times U(3)_L$, instead of the phenomenological $SU(3)_R \times SU(3)_L$. In the simplest models this results in obtaining 9 pseudoscalar massless mesons instead of desired 8 ones. This is reflected in some unpleasant features of the corresponding current algebra which can not be discussed here. The approach based on the nonlinear equations for the quark propagator probably offer a new possibility for the solution of this problem. An alternative approach based on the unified field theories is developed by Weinberg /116/. The present status of chiral phenomenology has been recently summarized by Pagels (see /117/ where further references can be found).

3.6. Attempt of synthesis in bags

The modern fashionable bags contain the quarks and gluons and pretend to simultaneously incorporate the equations of motion as well as the forces keeping the quarks inside hadrons. There are different sorts of bags which I will not try to describe here. On the parallel session they were discussed in some detail by Weisskopf, Kuti, P.Bogolubov, Struminsky and Mat'ev and here I only summarize several important points (for further references and details see these Proceedings and /118-121/ .

The M.I.T. bag is the most natural relativistic generalization of the Dubna bag. The new features are the following: 1) The external pressure B is introduced to balance the internal pressure of quarks and gluons moving inside a sphere of the radius. 2) The radius R is not fixed and is determined by the condition of the minimum energy of the system. This energy is the sum of the three terms:

$$M = \left[n_u\left(m_u^2 + \frac{X_u^2}{R^2}\right)^{1/2} + n_s\left(m_s^2 + \frac{X_s^2}{R^2}\right)^{1/2}\right] + \left[\frac{4}{3}\pi R^3 B - \frac{Z_0}{R}\right] + \Delta E_c .$$

Here n_u, n_s are the numbers of non-strange and strange quarks resp.; x/R is the momentum of the quark, which is derived by solving the Dirac equation in the infinitely deep spherical well. The second term represents a "renormalized zero-point fluctuation" energy, and the last term is the colour interaction energy which is responsible for spin-spin (hyperfine) splitting of hadron masses (this effect was first observed in the frame of QCD by De Rujula et al. /122/). 3) This expression was derived by using an analogy between massless colour gluons and photons, the colour playing the role of the electric charge. The colour gluons were confined by brute force inside the bag and the result of such a brutality is nice, only the colourless states can be stable. 4) The spectrum of excited states is exponentially growing in this model ($\sigma(M) \sim e^{MR}$). The Regge trajectories $\alpha(M^2)$ are also infinitely rising but for the spherically symmetric bag the dependence on M^2 is nonlinear.

We have just described a somewhat modified version of the M.I.T. bag. The main modification concerns the introduction of the quarks as point-like massive objects interacting with the coloured gluons. This modification of the original M.I.T. bag has been suggested by Kuti et al. /119/ and by De Grand et al./118/. When confined to a fixed sphere, the modified M.I.T. bag reproduces phenomenological results of the Dubna bag and, in addition, incorporates all good features of the QCD-approach/107,122/ to composite hadrons. Note that the confinement of the colourless bound states is in this approach an immediate consequence of the confinement of the gluons inside the bag. In general this bag picture is successful in qualitatively describing the lowest-lying states of baryons and mesons. However, further improvements are required if we wish to account for excited states and scattering processes.

First, the shape of the bag should be not fixed if we are to consider the processes of the fusion and fission of bags. As shown by Low /123/ the high-energy scattering of two bags can successfully reproduce the main features of elastic and inelastic processes of hadrons, provided that the bags are allowed to assume highly non-spherical shapes. With strongly deformed bags, we can also obtain a good description of hadrons with high values of the angular momentum of quarks and explain the linear growth of the Regge trajectories with M^2 /124/ . The diquark structure of the baryons /125/ is also naturally included in this picture /124/ . A variational approach to treating the static properties of deformed bags is presented at this conference by De Tar /126/ . A more radical modification of the bag model is suggested by the Budapest group /119/ . They supply the bag with an elastic skin (or "membrane") which enters into dynamical equations as a new variable, thus allowing for the canonical quantization of the whole system. The phenomenological motivation of this step is mainly in the fact that with the soft gluon-quark interaction (this hypothesis lies at the heart of the bag-phenomenology) it is difficult to explain the momentum sum rule in the deep-inelastic scattering (the missing momentum is ascribed to gluons, and yet the interaction of quarks with gluons is presumably weak). In the Budapest bag the missing momentum is possibly carried by the membrane. But now the question is: why the interaction of the membrane with quarks does not produce a large number of $q\bar{q}$? Being conceptually transparent the Budapest model is technically more complicated than the MIT model and phenomenological applications are still to be worked out. We hope that the relevant questions will be answered next year at the Budapest conference.

An interesting question to the bag-theory is: how to explain the nuclear structure? Without colour gluon exchange the lowest energy state of the 6 quarks would be a six quark bag and not a two-bag system representing a deuteron. With colour gluon exchange the six quark bag can be viewed as a system of two three-quark bags thus really representing the deuteron system /127/ . For many-nucleon systems an interesting phenomenon is predicted /128/ . At some quark density, higher than in nuclear matter, the "bag" will become the lowest state again, and a phase transition from nuclear matter to "quark matter" is possible. A simplified treatment of the deuteron as a six quark system is presented to this conference by Babutsidze and Machabeli /129/ . They put all six coloured quarks in an effective potential well, described by an oscillator potential, and classify the colourless states by using the methods of the nuclear shell model. The phenomenological results seem to be satisfactory yet the physical motivation of the calculations is not convincing. There is no two-bag structure of the deuteron, and it is not clear why the energy, say, of twelve-quark systems is not lower than that of the "deuteron". In general, the bag approach to nuclear physics opens new ways for investigating the nuclear structure, but before a quantative approach is possible, many important points have to be clarified.

Recently, it has been realized that the bag-like models predict an essentially richer spectrum of hadrons than non-relativistic potential models (including the Dubna bag). In fact, all kinds of exotics are predicted to exist with masses comparable to masses of the usual hadrons: ($qq\bar{q}\bar{q}$) bound states, the mesons with exotic J^{PC} , excitations corresponding to center of mass motion, etc. /120/ . These predictions are not in agreement with present experiment, as the empirical mass spectrum is rather sparse. In the contribution by Jaffe /120/ an attempt is made to identify the predicted $(qq)(\bar{q}\bar{q})$ states ("cryptoexotic" mesons) with some more or less established resonances. However, this seriously aggravates the well-known difficulty of missing $q\bar{q}$ - states. We consider the whole problem as essentially unsettled both from the experimental and theoretical sides.

The most interesting alternative to the MIT bag is the Vinciarelli-SLAC bag /121/ Unlike the MIT crew, the SLAC-crew starts from a field theory with a spontaneous symmetry breaking of the vacuum. Hence, the fundamental role of scalar fields in this approach. However, the surprising feature of the SLAC-bag is that the quarks concentrate near the surface of the bag which results in some not pleasant phenomenological predictions. The modern development of SLAC-bag is connected with solitons and is outside the scope of the present review. Some interrelations between SLAC and MIT-bags are discussed by Huang and Stump /130/. Using the variational approach to a model of quarks interacting with a scalar field, they obtain two solutions. One is similar to the MIT bag, either to the Vinciarelli-SLAC bag.

We have forgotten to mention two more problems of the bag theories. In the MIT-calculations it is supposed that the quark-gluon coupling is rather weak so as perturbation theory with respect to this interaction be sensible. In fact, the phenomenological applications require rather a large value for the coupling constant α_c ($\alpha_c \approx 2.2$) similar to the Sommerfeld constant $\alpha = 1/137$ (the authors of refs. /118,120/ erroneously quote the value $\alpha_c \simeq 0.55$, see /119/); Kobzarev and Mat'ev /131/ suggest a remedy to cure this desease at the price of the introducing new parameters in the theory (see these proceedings). We also have to note a difficulty of the Vinciarelly-SLAC model in explaining the observed scaling in deep-inelastic scattering processes. As suggested by Giles /121/, this difficulty can be resolved at the expense of supposing the surface of the bag to be extremely soft to deformations. Then the surface is considered as a dynamical object (like the Budapest membrane) and the theory becomes much more complicated than the original one. Only semiclassical solutions have been investigated up to now.

Concluding this rather sketchy discussion of bags we may generally state that the bag theories are successful phenomenological theories of hadrons made of coloured quarks and coloured gluons but they certainly do not constitute a fundamental theory of matter. The origin of the volume or surface tension, of symmetries and of their breaking and of quark masses is not explained. For example, bags are well suited for a description of the broken SU(6)-symmetry but not for the more fundamental $SU(3)_R \times SU(3)_L$ or at least $SU(2)_R \times SU(2)_L$ chiral symmetries. There are some attempts to incorporate PCAC in a bag-theory at a purely phenomenological level (see e.g. papers /132/). In these papers the pion is treated as an unconfined field interacting with a bag surface.

4. Quarks and Experiment. Conclusions.

Now we briefly consider some problems concerning the comparison of the quark model with experiment. The status of the baryon spectroscopy has not been significantly changed after the London conference /1/ (see also /133/), and we will not discuss it here. As to the meson spectroscopy, there is a dramatic change due to discovery of the new heavy resonances which we identify with charmed particles. Here we will not touch upon the details of the charmonium spectroscopy as well as the new data on the "old" particles. Instead, we concentrate on some of long standing contradictions between the quark model and experiment 1) The masses of all well-established mesons (except pseudoscalars which require a special treatment) can be described by a remarkably simple formula /107/. The formula is obtained as follows. Consider some equation for the $q\bar{q}$ bound-state wave function Ψ_{ij} of the i-th and j-th quarks (i and j are the flavours of the quarks) which we write in a rather general form

$$[\hat{R}_{ij} - K^2(M,m_i,m_j)]\Psi_{ij} = 0$$

$$K^2(M,m_i,m_j) = \frac{M^2}{4} - \frac{m_i^2+m_j^2}{2} + \frac{(m_i^2-m_j^2)^2}{4M^2}$$

Here $\hat{R}_{ij}$ is assumed to be some operator which does not depend on the quark masses m_i and m_j. We suppose that R_{ij} has the eigenvalues τ_{ij} depending on the orbital angular momentum $\vec{L}$ and on the total spin of the quarks as follows:

$$\tau_{ij} = \frac{m_0^2}{4} + \frac{\beta_{ij}}{4}L - \frac{m_1^2}{8}S(S+1) + $$
$$+ \frac{m_2^2}{8}[J(J+1) - L(L+1) - S(S+1)]$$

where J is the total spin of the bound state. We introduce here the spin-spin and spin-orbital splittings and a linear dependence of the eigenvalue on L (this corresponds to linear Regge trajectories of mesons). The equation of such an abstract form can be obtained in different quasipotential formulations of the bound-state problem; the B-S equation for $M \gg m_i, m_j$ can also be approximately reduced to a similar equation /107/. However, our specific Ansatz for τ_{ij} is of non-relativistic origin. We simply try to dramatize some problems concerning the meson mass spectrum. (Without using the above expression for τ_{ij}, the mass relations for states with equal J, L, S can be obtained supposing τ_{ij} is independent of i, j).

The expression for the meson masses is now obtained by setting $K^2(M,m_i,m_j) = \tau_{ij}$

The resulting mass formulae neglect the mixing of different quarks (say $u\bar{u} \leftrightarrow s\bar{s}$) in the isospin-zero $q\bar{q}$- states. The mixing can be considered by writing the equations for these states

$$[\hat{R}_{ii} - K^2(M, m_i, m_i)]\Psi_{ii} = \varepsilon^2 \sum_j \Psi_{ji} \equiv$$
$$\equiv \varepsilon^2 \sum_j \theta_{ij} \Psi_{ji} ; \qquad \theta_{ij} \equiv 1$$

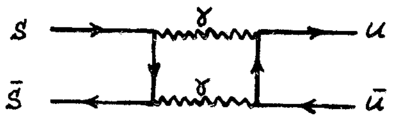
This mixing follows from the specific flavour--exchange mechanism discussed in Sect. 3.5. A similar form of the mixing matrix has been proposed by De Rujula et al. /122/ (mixing in the mass matrix) and by Fritzsch and Minkowsky (mixing in the mass-squared matrix, see /134/, where further references can be found). Our equations generalize the previous approaches, the quark dynamics is implicitly included in the dependence of K^2 on masses and in the eigenvalues τ_{ij} . Note that our mass formulae in general are neither linear nor quadratic in masses. For $L=0$, with no quark-mixing $(\varepsilon=0)$ and with $m_u = m_d$ we obtain the linear mass formulae $m_\rho = m_\omega$, $M_{K^*} = \frac{1}{2}(M_\rho + M_\varphi)$ which are satisfied within 1%. To account for ω-ρ mixing, corresponding to $s\bar{s} \leftrightarrow u\bar{u}, d\bar{d}$ mixing in $I=0, L=0, S=1, J=1$ state consider the equations for Ψ_{uu}, Ψ_{dd} and Ψ_{ss} with some mixing parameter $\varepsilon^2_{\omega\varphi}$ By applying the Schrödinger method of factorization (see e.g., /135/) one easily obtains the expressions for m^2_ω and m^2_φ in terms of one unknown parameter $\varepsilon^2_{\omega\varphi}$ (other parameters in this case are determined by the masses of K^* and ρ). The predicted mass of the φ- meson is in good agreement with the experimental value /139/ . The treatment of mixing the $I=0, L=1, S=1, J=2$ state requires some additional information on the coefficients β_{ij} .

These can be determined by fitting the general mass formulae (with mixing) to the masses of the well-established mesons. The result is rather interesting - the parameters $\beta_{uu}, \beta_{us}, \beta_{ss}$ ($\beta_{dd} = \beta_{uu}$, $\beta_{ds} = \beta_{us}$ - by isospin invariance) satisfy the relation $\beta_{us} = \frac{1}{2}(\beta_{uu} + \beta_{ss})$ As the differences between these parameters are in fact not large ($\beta_{uu} = 0.872$, $\beta_{us} = 0.942$, $\beta_{ss} = 1.015$) the multiplicative relation $\beta^2_{us} = \beta_{uu}\beta_{ss}$ is also very well satisfied. To our knowledge there are no arguments in favour of the additive relation but in the contribution to this conference by Pasupathy /136/ it was demonstrated that the multiplicative one probably follows from duality /137/ and from factorization property of the Regge-pole residues /138/ .

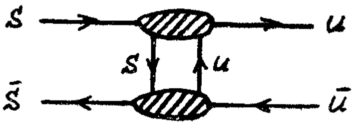
An interesting property of the Regge trajectories is that they seem to intersect in the same point of the L, m^2 plane. This fact for the J - trajectories was also observed by Becher and Böhm /147/ . It can be qualitatively explained by a somewhat smaller radius of the particles containing heavier quarks. Azimov, Frankfurt and Khoze also proposed that the radius of charmed particles is dramatically smaller than that of "usual" particles .

2) As was emphasized above the pseudoscalar mesons require a special treatment. Here we mention the most mysterious η-η' problem and the pion mass provlem. It is now generally believed that any solution of both problems is possible only in a theory explaining the broken chiral symmetry. As discussed above there exist two approaches to the η-η' problem. Both relate the large mass difference between η and π and the violation of both quadratic and linear relations $\eta' \neq 2K - \pi$ to a strong mixing of $s\bar{s}$ and $u\bar{u}$ (or $d\bar{d}$) in $I=0, L=0, S=0$ channel.

In QCD this mixing is due to the diagram



where coloured gluons γ are exchanged in the S-channel.. A more phenomenological explanation is presented diagrammatically as



where strange (flavoured) bound states and resonances are exchanged in the t-channel (see Sect. 3.5).

For both mechanisms the mixing matrix can be written in the form $\varepsilon_{ij} = \varepsilon^2$ for all i and j. This mixing matrix was introduced above in the equation for Ψ_{ii}. The resulting expressions for the masses of η and η' are of the form

$$\eta^2 = m_\rho^2 - 6\varepsilon^2 - 2\sqrt{9\varepsilon^4 + \Delta^2(\Delta^2 + 2\varepsilon^2)}$$

$$\eta'^2 = m_\rho^2 - 6\varepsilon^2 + 2\sqrt{9\varepsilon^4 + \Delta^2(\Delta^2 + 2\varepsilon^2)}$$

Here $\Delta^2 = m_s^2 - m_u^2 = K^*(K^* - \rho) = 0.1064$, $m_\rho^2 = K^2 + \Delta^2/K^2 = 0.2916$ the only unknown parameter being ε^2. As $m_\rho^2 \simeq \eta^2 = 0.3012$, the approximate value of ε^2 is $\varepsilon^2 \approx \frac{\Delta^2}{2} = -0.053$. With this value of ε^2 the prediction for η' is $\eta' \simeq 0.963$ which is in very good agreement with the hypothesis that η' – meson is $X(958)$. However, in this approach the pion mass is defined by the relation $\pi^2 = m_P^2 - 2\Delta^2$ ($= \eta^2$ for $\varepsilon = 0$) and $\pi \simeq 0.280$ is two times as large than the experimental value. We conclude that the pion wave function cannot be described by this simple equation. The ideas described in Sect. 3.5 might be relevant to this problem but no successful model is available at this moment.

Fritzsch and Minkowsky /134/ used the same mixing matrix for the mass squared matrix. Their results can be obtained from our formulae if we write $m_\rho^2 = K^2 = 0.2457$, $\Delta^2 = \frac{1}{2}(K^2 - \pi^2) = 0.1134$. Note that this value of Δ^2 disagrees with that obtained from the vector and tensor meson masses: $\Delta_V^2 = (K^{*2} - \rho^2)/2 = 0.099$, $\Delta_T^2 = (K^{*2} - A_2^2)/2 = 0.150$. If we nevertheless, try to describe η and η' by their formulae with ε^2 defined by the η - mass, $\varepsilon_\eta^2 = -0.200$, the prediction for η' is $\eta' \simeq 1.61$. Alternatively, defining ε^2 from X-mass, we find $\varepsilon_X^2 \simeq -0.056$, $\eta \simeq 0.50$. This clearly shows that η and X do not satisfy the equations. A much better fit can be obtained with $\eta' = E(1.42)$. There are other schemes in which the mass of the η' is predicted to be close to the mass of the E-meson. For example, Caser and Testa /140/ came to this conclusion by using a variant of infinite momentum frame current algebra for describing the chiral symmetry breaking. They also suggest identifying the X(.958) with an almost pure glue state.

Attempts to preserve the identification $\eta' = X(.958)$ are based either on introducing some admixture of glue states in the η and η' /141/ or on using different mixing angles θ_η, $\theta_{\eta'}$ /142,122/. To solve this long standing problem it is badly desirable to establish the J^P – quantum numbers of the X(.958) and E(1.42) and to obtain a more detailed and credible experimental information on radiative decays in which these mesons participate. In contrast with the statement of PDG /139/, the present status of the quantum numbers of X(.958) is very controversial. This was clearly demonstrated at the Conference by Ogievetsky and Lednicky (see these Proceedings). Unfortunately, the state of the art in the meson radiative decays is also far from satisfactory (see e.g. /143/ and the invited paper by Gerasimov). In addition to defining the J^P – quantum numbers of the X and E the most important experimental problems are the measurements of $\Gamma(X \to \rho^0\gamma)$, $\Gamma(X \to \gamma\gamma)$, $\Gamma(\rho \to \pi\gamma)$

$\Gamma(K^{*0}\to K^0_{\ }\gamma)$. New measurements of $\Gamma(\eta\to\gamma\gamma)$ and $\Gamma(\pi^0\to\gamma\gamma)$ not using the Primacoff effect would be also welcomed, in view of their utter importance for theory (esp. for quark models).

In recent paper /144/ Greco and Etim-Etim have constructed a model successfully describing all the known meson radiative widths expect $\Gamma(\rho\to\pi\gamma)$. Not judging their general reasoning, we only remark that the naive quadratice $\eta-\eta'$ mixing is used for calculating decays with η and η' . As must be clear from the above discussion, this unavoidably results in severe difficulties with mass formulae, which are not discussed in the paper.

Finally, consider the new particles. If we suppose that Ψ is a pure $c\bar{c}$ state then our formulae immediately give the predictions

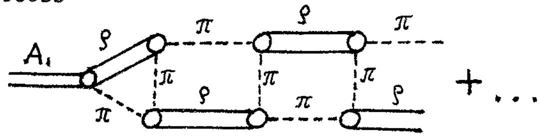
$$\mathcal{D}^*=\tfrac{1}{2}(\Psi+\rho)\approx 1.93,\quad F^*=\tfrac{1}{2}(\Psi+\varphi)\approx 2.06.$$

For the pseudoscalar mesons if we use the same m_p^2 as above, we will find $\mathcal{D}\approx 1.64$, $F\approx 1.87$, $X_{c\bar{c}}\approx 3.01$. If we use as the input $\mathcal{D}=1.87$, we find $X_{c\bar{c}}\approx 3.1$. Trying all possible modifications of our equations we never obtain the mass of the $X_{c\bar{c}}$ as low as 2.8 GeV. We think that the most plausible explanation of these discrepancies is the possibility of mixing the $c\bar{c}$ states with $t\bar{t}$ or $b\bar{b}$ states (the admixture of u,d quarks does not help). There exist good candidates for $L=1$ $c\bar{c}$ mesons (see Wiik's talk, these Proceedings). If we draw the straight-like L-trajectory for $c\bar{c}$ through the point in the L,m^2 plane in which $s\bar{s}$, $u\bar{u}$ and $u\bar{s}$ trajectories intersect, we find that the orbitally excited states of $c\bar{c}$ must lie near 3.6 GeV. More explicitly, we have obtained that

$$\frac{\beta_{ss}-\beta_{su}}{\Delta_{su}}\simeq\frac{\beta_{su}-\beta_{uu}}{\Delta_{su}}\approx 0.67\,.$$

Let us suppose that the same is true if we replace the S-quark by the c-quark. Then we obtain $\beta_{cc}\approx 3.9$ which allows us to estimate the masses of $^3P_1$ and $^3P_2$ $c\bar{c}$ mesons. However, the SS and LS splittings for present candidates are difficult to explain in usual terms and this possibly tells us that we have no simple $c\bar{c}$ states but some more complicated mixtures of $c\bar{c}$ quarks with other new quarks.

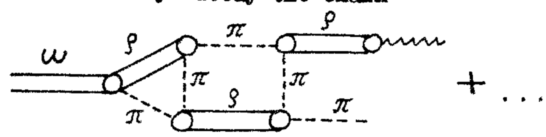
3) In conclusion we briefly discuss the problem of missing particles. A more detailed discussion of this problem can be found in Ref. /133/ . We mention here only the most notorious A_1-problem. Practically all variants of the quark model predict $J^{PC}=1^{++}$ particle with mass $\sim$ 1.1 GeV. However, the latest very good experiments fail to confirm that the $A_1(1.1)$ bump can be interpreted as a resonance, and there are no other candidates for such a particle. A possible explanation of this phenomenon may be searched in the influence of the $(q\bar{q})(q\bar{q})$ channels. For example, in the decay the contribution of the rescattering process



is rather large (due to the large $\rho\pi\pi$ coupling and the large radius of the π-exchange interaction). In addition, there are other two-meson channels strongly coupled to A_1 and to each other. It is possible that the interaction of all these channels can spoil the simple quark model picture in which A_1 is regarded as the pure $q\bar{q}$- state. The detailed investigation of this problem would be very desirable. A preliminary discussion of some related ideas was attempted by Dashed and and Kane and by Badalyan, and Simonov /145/, see also /133/ .

A similar mechanism can spoil the quark model prediction for $\Gamma(\rho\to\pi\gamma)$ and $\Gamma(\omega\to\pi\gamma)$.

In the $\omega\to\pi\gamma$ decay the chain



may give a large contribution to the decay rate, and there is no similar contribution to the $\Gamma(\rho\to\pi\gamma)$. Such mechanisms could be relevant to $\mathcal{D},\mathcal{D}^*$ interactions, as recently observed by Okuň and Voloshin. They proposed the "hadronic molecules" made of $\mathcal{D}$ and $\mathcal{D}^*$ which are bound by the pion exchange.

The first discussion of the interaction in the exotic channels was given by Shapiro et al. /146/ who investigated the interactions in $N\bar{N}$ channels and demonstrated a possibility of exsistence of rather narrow $N\bar{N}$ resonances. The present state of arts in this field was summarized at the conference by Shapiro (these Proceedings). Additional information can be found in Rosner's review /133/.

The moral of this sketchy discussion is as follows. The naive two particle $(q\bar{q})$ model of massive meson resonances is certainly too naive. The exotic $(q\bar{q})(q\bar{q})$ channel cannot be neglected for large masses when many channels are open or almost open, and we face an unpleasant situation: with growing mass of the $q\bar{q}$ bound state, the $q\bar{q}$ interaction is becoming simpler (the exchange forces are dying away, the OZI-rule is becoming exact), but the influence of exotic $(q\bar{q})(q\bar{q})$ channels can spoil the usual quark model predictions. Fortunately, the existence of the new (charmed) particles provides us with the unique possibility of the pure $q\bar{q}$ high mass resonances which are not spoiled by $(q\bar{q})(q\bar{q})$ admixture.

In this brief discussion of the experimental status of the quark model we concentrated on some unsolved problems, leaving its numerous successful predictions aside. It must be stressed that there is no substitute today for the quark model in explaining diverse experimental facts in strong, weak, and electromagnetic interactions of hadrons. Despite the existence of some unsolved theoretical and experimental problems we may conclude that the quark model is in a very good shape in Tbilisi!

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