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AN INTRODUCTION TO WEAK-INTERACTION THEORIES WITH DYNAMICAL SYMMETRY BREAKING

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ABSTRACT

We present a straightforward introduction to theories of the weak interactions with dynamical symmetry breaking-theories of "technicolor" or "hypercolor". Our intent is to inform experimentalists, but also to goad theorists. We first describe the motivation for considering theories of this type. We then outline the structure that such a theory must possess, including new gauge interactions at mass scales of 1-100 TeV. We argue that, despite their reliance on phenomena at such enormous energies, these theories contain new phenomena observable at currently accessible energies. We describe three such effects which are especially likely to be observed.

RÉSUMÉ

Nous donnons une présentation simple des theories des interactions faibles avec brisure dynamique de la symétrie, les theories dites de "technicolor" ou "hypercolor". Notre intention est d'informer les expérimentateurs, mais aussi d'aiguillonner les théoriciens. D'abord, nous décrivons nos motivations pour considérer les théories de ce genre. Ensuite, nous esquissons la structure qu'une telle theorie doit posséder, et qui comprends des interactions de jauge nouvelles á des energies de 1-100 TeV. Nous montrons que ces théories, quoiqu'elles dépendent de ce qui se passe á des énergies trés élevées, contiennent en fait des phénomènes nouveaux observables aux énergies accessibles actuellement. Nous décrivons trois effets de ce type pour les quels la probabilité de les observer est la plus grande.

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1. INTRODUCTION

During the past few years, the experimental study of weak interaction processes has consolidated our understanding of these interactions. It is no longer a contested proposition but, now, a commonplace to claim that the weak interactions are mediated by vector mesons which are, in fact, gauge mesons of a gauge theory built on the group SU(2)xU(1). The coupling of fermions to these gauge mesons are well-measured and consistent among different reactions. But this illumination, impressive as it is, is only a piece of a complete picture. The gauge symmetry of the weak interactions must be a spontaneously broken symmetry; otherwise the W-bosons could not be massive. Whatever it is that causes this symmetry-breaking, whatever it is that gives the W bosons masses, is something that stands outside the SU(2)xU(1) gauge theory. It is a new fundamental interaction - new in the sense that it has not yet been observed directly and, in fact, has not yet become seriously constrained by experiment.

What is the nature of this new interaction? In the original theories of Weinberg and Salam[1)] this interaction was constructed as a theory of an elementary scalar particle - the Higgs meson - with a self-interaction and a coupling to fermions as well as its coupling to the SU(2)xU(1) gauge bosons. As the Weinberg--Salam gauge theory became accepted as the correct theory of the weak interactions, the notion that its gauge symmetry is broken by an elementary scalar field also became part of the standard dogma. We regard this as unfortunate, for two reasons: First, we consider it unlikely that the Higgs meson actually is a fundamental particle, in the sense that quarks and leptons are fundamental, for reasons that we will discuss at length in section 2. Secondly, the theory of elementary Higgs scalars has consumed much effort which might have been better spent trying to imagine more interesting alternatives.

In these lectures, we will review one such alternative, more physical picture of what it is that breaks SU(2)xU(1), a picture in which this breaking is accomplished by a new strong interaction theory of fermions and gauge mesons. The Higgs mesons of the Weinberg-Salam model become, in this picture, composites

- bound states of the new fermions. We will refer to this picture as "dynamical" symmetry breaking. Because this picture contains strong interactions, it is much richer than the conventional model, both in its theoretical structure and in its observable implications. Our aim in these lectures is to discuss the structure of this new picture in as straightforward a manner as we can, and then to detail its implications which can and should be tested in the near future. Little that we will say here is novel; our intention is to assemble various pieces of the theoretical understanding of dynamical symmetry breaking into a coherent and compelling argument. We warn the reader that a few arguments have been oversimplified for the sake of clarity; in each case, the precise argument may be found in the original papers cited[2)].

The plan of these lectures is the following: Sections 2 and 3 discuss the motivation for postulating a new strong interaction just to break SU(2)xU(1). We apologize in advance that our arguments here will be of a theoretical (as opposed to a phenomenological) nature. We hope, however, that these sections will make clear why we feel it important to study models of this type. Sections 4 through 6 review more specifically the structure of models of dynamical symmetry-breaking. Section 4 also completes our explication of a theorist's appreciation of these models, by showing that the simplest such model gives rather directly the symmetry-breaking pattern for SU(2)xU(1) required by weak interaction phenomenology[3,4)].

Sections 7 through 9 review some specific experimental implications of these models. In each case, the effects we discuss are not specific to particular realizations of dynamical symmetry breaking but, rather, should appear in almost any realistic model. In section 7, we note that such models contain weak flavor-changing neutral current interactions which can mediate rare processes such as $K_L \to \mu e$ and $\mu \to e\gamma$ [5)]. Such processes should occur with branching ratios between 10^{-10} and 10^{-14}. In section 8, we demonstrate a natural mechanism for the appearance of CP violation in theories with dynamical symmetry breaking[6)]. This mechanism yields a neutron electric dipole moment of 10^{-24} - 10^{-25} e-cm, just below the present upper limit[7)]. In section 9, we argue that several pseudoscalar mesons built of the new fermions - including, at least, an electrically charged pair and two neutral particles - should be quite light, with masses of at most 30 GeV, but possibly as small as 5 GeV[5]. We will discuss the signatures of these particles in e^+e^- annihilation experiments, where they might be most readily detected.

2. WHAT'S WRONG WITH ELEMENTARY SCALARS ?

Why should one introduce a new theory of strong interactions just to break SU(2)xU(1) ? In the previous section, we had put this idea forward as an interesting possibility, one which asks to be explored. In this section, we will explain why it is a necessity. We will argue that serious problems arise in viewing the Higgs meson as an elementary scalar field. These problems do not impeach the consistency of the theory as a phenomenological description of the Weinberg-Salam symmetry-breaking, but they do indicate that this description is no more than a phenomenology. Arguments similar to those given here have been offered by Weinberg[8], 't Hooft[9], and Susskind[4].

To understand these problems of the standard Weinberg-Salam theory, it is worth comparing this theory to another familiar and fashionable theory, Quantum Chromodynamics. Which of these theories, we should ask, is the more predictive ? In practice, given our current computational skills, the Weinberg-Salam theory clearly gives the more detailed predictions, since, in this theory, all coupling constants are weak, and, therefore, the rate of any process may be determined by straightforward computation of Feynman diagrams. In QCD, only a handful of processes of a very special type involve small coupling constants; for the most important questions - the ratio of the ρ-meson and proton masses, for example - we have, at the moment, no adequate scheme of computation. This fact makes the Weinberg-Salam theory much easier to understand in detail and to test.

If one speaks of matters of principle, however, the situation is entirely reversed. If we make the approximation of ignoring the bare masses of the u and d quarks and the influence of heavier flavors, QCD has only one adjustable parameter, its gauge-field coupling constant g_C. But, further, the value of g_C varies with the momentum transfer involved in a particular process, growing small at large momenta and vice versa. This means that what we actually have the freedom to adjust is the value of g_S at a certain momentum, or, alternatively, the momentum Λ at which g_S reaches a certain fixed (large) value[10]. Since this momentum Λ is the only adjustable parameter in the theory, all other quantities appearing in the theory are, in principle, computable in terms of it. But since Λ carries the dimensions of mass, it can do no more than set the scale of masses. All dimensionless quantities appearing in the strong interactions must be simply pure numbers, computable without adjustable input. At the level of principle, this is all that one could ask from a physical theory. And, even if we cannot yet compute these numbers precisely, we can at least estimate such quantities as m_ρ/m_N or m_Δ/m_N from little more than our intuition about how QCD behaves.

In the standard Weinberg-Salam theory, the situation is quite different. In order to define the theory, it is necessary to specify the values of a relatively large number of parameters. Each of these parameters may be freely adjusted without affecting the consistency of the theory. These parameters include the SU(2)xU(1) gauge meson couplings g and g', but also the masses and self-couplings of the elementary Higgs mesons. This means that the theory cannot predict the masses of the Higgs mesons[11), indeed, it cannot even predict the number of these particles which should exist. This set of parameters also includes the couplings λ of Higgs mesons to quarks and leptons. Even in the simplest case of one Higgs field, these couplings form a matrix λ_{ij}: the fermion of flavor i converts to one of flavor j with the emission of a Higgs meson. The eigenvalues λ_a of this matrix give the quark and lepton masses through relations of the form

$$m_a = \lambda_a \langle\phi\rangle, \tag{1}$$

where $\langle\phi\rangle$ is the Higgs field vacuum expectation value. One of the angles in the rotation which diagonalizes the matrix λ can be identified with the Cabibbo angle θ_C. None of these quantities can be predicted by the theory.

A possible way of evading this problem of the adjustability of the λ_{ij} might be to impose symmetry relations among these quantities. An example of such a relation is that present in the SU(5) grand unified theory[12)] which gives a successful prediction of the ratio of the b-quark and τ-lepton masses[13)]. Unfortunately, that one successful result of this program remains unique. We know of no other substantial progress toward calculating the quark and lepton masses since 1972, when Georgi and Glashow failed to compute the mass of the electron[14)].

To find a large number of adjustable parameters in a theory is itself a sign that some ingredient is missing. But a more certain sign of this is to find that some parameters must be adjusted to excessively small values. In the standard Weinberg-Salam theory, one finds several parameters of this type: A first example reflects the small size of the u quark and electron masses. Since the quantity $\langle\phi\rangle$ which appears in (1) also sets the scale of the W-boson mass μ_W, one can determine that $\langle\phi\rangle \sim 250$ GeV. Then the observed values, of MeV order, for the u and e masses require

$$\lambda_u,\ \lambda_e \sim 10^{-5} \tag{2}$$

A second example concerns the possibility of an overall complex phase θ for the matrix λ_{ij}, which would become an overall, CP-violating phase in the quark mass matrix. The matrix λ_{ij} contains another CP-violating angle δ which, in the picture of CP-violation (now considered standard) due to Kobayaski and Maskawa[15)], must be of order 1. However, θ contributes to the neutron electric dipole moment d_N a term of magnitude[16)]

$$d_N \sim (10^{-16}\ \text{e-cm}) \times \theta \tag{3}$$

model contains new vector mesons of large mass μ_X, such that $\mu_W/\mu_X \sim 10^{-13}$. To obtain this large ratio in a theory with elementary Higgs mesons, one must adjust a Higgs meson $(\text{mass})^2$ in the grand unified theory to a tolerance of one part in 10^{26}.

How can one obtain these small numbers in a natural way? We have stressed that, in the Weinberg-Salam theory with elementary Higgs scalars, this question is not merely technically difficult but unanswerable as a matter of principle. We must, then, seek out alternatives to that theory.

3. AN OVERVIEW OF DYNAMICAL SYMMETRY BREAKING

In the previous section, we attacked the standard model of weak-interaction symmetry-breaking by an elementary Higgs field. We wish that we had a complete, consistent theory which remedies the defects of the standard theory which we have displayed. Unfortunately, we do not. Thus far, only pieces of such a theory have actually been constructed; these pieces will be reviewed in later sections of these lectures. First, however, we wish to explain the perspective from which we attempt the resolution of these problems, the program which we will use in constructing specific theories. It is this program, rather than any specific realization of it, which we mean when we speak of dynamical symmetry breaking. And, to a great extent, it is this program, rather than specific models, which the experiments described later in these lectures test.

The program is, simply, to build up the Higgs mesons in just the way that the hadrons are built in QCD, from a theory of fermions, of zero bare mass, bound by a strongly-coupled gauge field. Such a theory will have a rich structure, but it will introduce (in addition to the Weinberg-Salam gauge couplings g and g') only a single parameter - a mass scale Λ - for each new strong-interaction gauge group. In the best case, there should be only one such parameter Λ. This scale could then be determined from any other dimensionful weak-interaction quantity, for example, the Fermi constant G_F; all other weak-interaction quantities, including the quark and lepton masses, could then, in principle, be computed. Actually, though, the models which have been presented in the literature [5,17] are rather less elegant. They contain two scale parameters Λ : one sets the scale of μ_W ; the other enters in determining the scale of quark and lepton masses.

Before describing what we know of the realization of this program, let us answer three questions of a more general nature. The first question is the following : Our computational ability, in theories of strongly interacting gauge fields, is extremely limited. How, then, can we expect to make any useful predictions concerning these new strong interactions? To understand the answer to this question, recall that, in the ordinary strong interactions, many aspects of low-energy dynamics can be computed unambiguously using symmetries of the theory, despite the presence of strong interactions. For example, Weinberg showed, a decade ago, how to compute the low-energy scattering of pions from nucleons and from each other by using the methods of current algebra[18]. The theoretical apparatus of current algebra has a direct analogue in any strong-interaction theory of fermions with small bare masses; hence, it may be used to discuss some aspects of the new strong interactions which we introduce. It will become clear as we proceed that these aspects include most of the important effects of these interactions on physics at presently accessible energies.

The second question is that of clarifying to what extent this program has actually been realized concretely. The most basic question, how to break the SU(2) x U(1) symmetry of the weak interactions, has been given an elegant answer, which we will describe in the next section. The theory of the spectrum of observable Higgs mesons is also well-understood; in section 5, we will review this theory and illustrate it in a simple model. Two other aspects of the weak interactions have been recovered from this viewpoint, although only at a semi-quantitative level: In section 6, we discuss our understanding of the origin of quark and lepton masses; in section 8, we discuss the origin of CP violation.

The third question is, to any theorist, the most crucial: Having introduced a new gauge interaction, what should one name it ? Dimopoulos and Susskind[17] have labelled this new interaction "technicolor". Other suggestions have been put forward by a variety of notables, including the editor of the Physical Review.

Throughout these lectures we will, against the better judgement of one of us, use the name "hypercolor" suggested by the other[5].

4. A BALL-AND-STRING MODEL WHICH BREAKS SU(2) x U(1)

In this section, we will review a mechanism by which a new strong-interaction can break the Weinberg-Salam SU(2) x U(1). Actually, we will begin by posing the following more innocent-looking question: Given a pair of massless fermions, coupled strongly to hypercolor and weakly, in the fashion standard for a quark or lepton doublet, to the Weinberg-Salam gauge bosons, what happens? The answer is a theory of SU(2) x U(1) breaking in which the input is minimal and the discovery maximal. The discovery in question was made, independently, by Weinberg[3] and Susskind[4].

To analyze the behavior of this theory, we proceed in two stages. First, we turn off the Weinberg-Salam coupling and ask what structure develops from the pure strong-interaction theory. Then we will add the electroweak interactions as a perturbation and ask how they affect that structure.

What physics do we expect, then, from a set of two massless fermions (call them U , D) in a gauge theory of strong interactions? The familiar strong interactions, involving the almost massless quarks u , d coupled to QCD, provide an example of such a theory. In that example, we observe that the quarks are confined into color-singlet bound states. The isospin symmetry linking u , d remains a good symmetry, but another symmetry of the theory, chiral SU(2), is spontaneously broken. The simplest expectation for our hypercolor theory is that it manifests these same three features. This is all we need to assume about hypercolor dynamics, so at this point we could simply go on to the second stage of our analysis. However, the most crucial of our three assumptions is also the least appreciated: Though it is well-known that chiral symmetry breaking is a feature of QCD, it is less broadly recognized that this is an expected or natural feature of this theory. Let us now digress to explain why, indeed, it is. In the course of this digression we will introduce notations which will be useful in the next several sections.

In a gauge theory, the coupling of a gauge boson to fermions preserves the fermion helicity. If the fermions are massless, there are no other terms in their equations of motion which mix different helicities. Thus, any symmetry interchanging fermion flavors is an equally good symmetry if performed on fermions of one helicity only. In QCD, if the u and d quarks had precisely zero bare masses, one would have two separate SU(2) isospin symmetries, one rotating only the left-handed components of u, d, the other only the right-handed components. We will use the symbols L, R to label helicities. To these two symmetries would be associated conserved currents $J_L^{\mu a}$, $J_R^{\mu a}$:

$$J_L^{\mu a} = \bar{q}\gamma^\mu\left(\frac{1-\gamma^5}{2}\right)\tau^a q \quad ; \quad J_R^{\mu a} = \bar{q}\gamma^\mu\left(\frac{1+\gamma^5}{2}\right)\tau^a q \qquad (5)$$

where $q = (u,d)$ and $\tau^a = \sigma^a/2$ is an isospin matrix, and the isospin index $a = 1,2,3$. We will refer to the group of these handed flavor symmetries, SU(2) x SU(2) in this case, as the chiral group.

Running this argument in the other direction yields some useful nomenclature and a distinction which we will apply repeatedly in these lectures. The separate symmetries of left-and right-handed isospin, if exact, would forbid the appearance in equations of motion of quark mass terms. We may, then, say that the non-conservation of the currents (5) is a measure of the u and d quark masses. The quark bare masses, defined in this way, enter into and may be determined from various predictions of current algebra[19). It is these quantities, often called current-algebra quark masses, that we mean when, in these lectures, we use the term "quark masses". With the normalization used in ref. 19, the current-algebra masses of the u,d and s quarks are approximately 5,8 and 165 MeV. In addition to these masses, quarks also have a so-called dynamical mass, which is about 1/3 of the proton mass. We will see in a moment that the dynamical masses of u, d, s, etc. need not vanish even if their current-algebra masses are zero. Finally, the term "constituent quark mass" refers, roughly speaking, to the sum of their current-algebra and dynamical masses. Thus, the constituent masses of u, d and s are 1/3, 1/3 and

1/2 GeV, respectively; if we take the constituent mass of the charmed quark to be about 1.6 GeV, we infer that its current-algebra mass is roughly 1.3 GeV.

Having introduced the handed symmetries (5) of a theory of massless quarks, we wish to argue that they are spontaneously broken. By this we mean that these symmetries, though they are invariances of the Hamiltonian H of (massless-quark) QCD, are not respected by the vacuum state. To see this, we can attempt to construct the vacuum, the eigenstate of H of lowest energy[20]. Split H into two pieces:

$$H = H_d + H_c , \qquad (6)$$

in which H_d contains terms which preserve the number of quarks and of antiquarks, including the kinetic energy terms and interactions such as those of Fig. 1a, and H_c contains terms which create and annihilate pairs, such as that shown in Fig. 1b. H_c has expectation value zero in the vacuum of perturbation theory, the state containing no quarks or antiquarks. However, it has large off-diagonal matrix elements. We can clearly form a lower-energy state by taking advantage of these terms, by mixing the

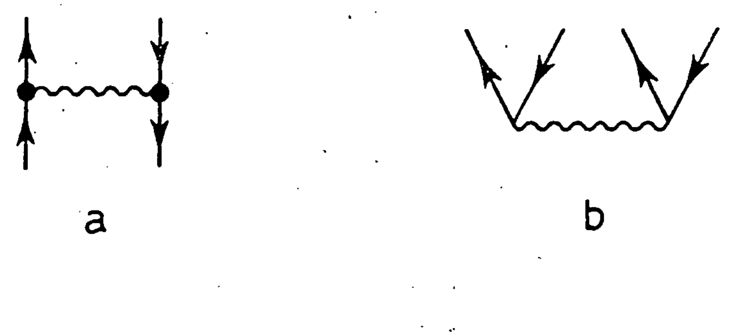

Figure 1. Interaction terms in the QCD Hamiltonian.

Figure 2. A $q\bar{q}$ pair with vacuum quantum numbers.

perturbative vacuum state with states containing a certain number of quark-antiquark pairs. If the interaction terms of Fig. 1 are weak, this mixing will be a small effect. However, if the interaction is made strong, two factors alter this balance: The attractive interaction in color-singlet channels of Fig. 1a lowers the energy H_d of a $q\bar{q}$ pair, lowering the cost of producing an additional pair, and the strength of Fig. 1b increases, increasing the gain from diagonalizing H_c. For a sufficiently strong coupling, H_c becomes the dominant effect; then the lowest-energy state of H_c becomes a state containing a large and indefinite number of $q\bar{q}$ pairs, similar to the pair condensate of a superconductor.

To connect this dynamics to chiral symmetry, we need only look at one of these pairs. Such a pair must have vacuum quantum numbers, in particular, zero momentum and zero angular momentum. We have shown in Fig. 2 a pair satisfying these constraints. Such a pair cannot, obviously, have zero helicity; thus, helicity conservation is not respected by the vacuum state. Another way to see this is to observe that, in Fig. 2, the q_R has been forced to pair with an anti-q_L. This means that the two isospin symmetries (5) can no longer be separate; they are linked by the pairing. The ordinary isospin symmetry, generated by

$$J^{\mu a} = J_R^{\mu a} + J_L^{\mu a} = \bar{q}\gamma^\mu \tau^a q \qquad (7)$$

is still preserved, but the symmetries generated by

$$J^{\mu 5 a} = J_R^{\mu a} - J_L^{\mu a} = \bar{q}\gamma^\mu \gamma^5 \tau^a q \tag{8}$$

are spontaneously broken. Since the vacuum does not respect the separate symmetries (5), the u, d quarks can acquire dynamical masses through their interaction with the pair condensate.

There are two mathematical signals of chiral symmetry breaking which we will make use of in these lectures. The first is the appearance of a vacuum expectation value of the mass operator $\bar{q}q$: If i,j are isospin indices u,d,

$$< o|\, \bar{q}_i\, q_j |o> = \delta_{ij}\, \Delta \tag{9}$$

where Δ has the dimensions of $(\text{mass})^3$. The second arises from the following considerations: We are assured that to every spontaneously broken symmetry current there corresponds a massless scalar particle, a Goldstone boson[21]. The spontaneous breaking of (8) yields, then, an isospin triplet of pseudoscalar particles — pions — which would be massless if the u,d masses were zero. An axial SU(2) current (8) may create from the vacuum a single pion of momentum p. We may write this as $(a,b = 1,2,3)$

$$< o|\, J^{\mu 5 a}\, |\pi^b(p)> = i\delta^{ab}\, p^\mu\, f_\pi\ . \tag{10}$$

The pion decay constant f_π has dimensions of mass. In QCD, if we write $\Lambda = 300$ MeV, we can determine that[19]

$$\Delta \cong (0.4)\,\Lambda^3 \quad , \quad f_\pi = (0.3)\,\Lambda\ . \tag{11}$$

Having now seen that one should expect chiral SU(2) to be a spontaneously broken symmetry in a strong interaction theory of massless fermions U, D, let us now return to the main course of our argument. What happens when we couple this theory, with its broken-symmetry state, to the Weinberg-Salam gauge bosons? The crucial feature of this coupling is that the weak interactions involve the handed currents (5) and not simply the vector currents $J^{\mu a}$. This means that the gauge mesons couple to the pair condensate, which can communicate to them its symmetry breaking and, as a result, give the gauge mesons masses. This mass generation is easy to understand qualitatively by comparing the pair condensate to a plasma, in which massless modes of oscillation, sound waves, coupled to a massless photon, produce a plasma oscillation, a mode with finite frequency at zero momentum. At a deeper level, the three examples of a plasma, a gauge theory coupled to an elementary Higgs field, and a gauge theory coupled to a pair condensate, generate gauge boson masses through the same mathematical mechanism, the Higgs mechanism[22].

We would like to actually compute the masses generated for weak gauge bosons in our hypercolor model. This is difficult to do by straightforward examination of the condensate, but it can be done easily using a more indirect method[23]: If a vector boson is to acquire a mass μ^2, its self-energy $\Pi^{\mu\nu}(p)$ must tend to μ^2 as $p \to 0$. But since $\Pi^{\mu\nu}$ is also the vacuum polarization tensor (see Fig. 3), this object must be transverse: $p_\mu \Pi^{\mu\nu}(p) = 0$. Thus, it must have, as $p \to 0$, the form

$$\Pi^{\mu\nu}(p) \underset{p \to 0}{\sim} \left(g^{\mu\nu} - \frac{p^\mu p^\nu}{p^2}\right) \mu^2 \tag{12}$$

The term in (12) of the form $g^{\mu\nu}$ is difficult to compute. But the term with a $1/p^2$ is easy to isolate, since such a term can arise from Fig. 3 only if one current creates and the other annihilates a massless particle. The matrix element in Fig. 3 is to be computed in the hypercolor theory, and the only massless particles in this theory are the three pions. Hence, the only processes which contribute to this term are those of the form of Fig. 4.

To evaluate these terms for the case at hand, we should write the coupling of the doublet $Q = (U, D)$ to the SU(2) x U(1) gauge bosons:

$$\delta\mathcal{L} = g A^a_\mu \bar{Q} \gamma^\mu \left(\frac{1-\gamma^5}{2}\right) \tau^a Q \tag{13}$$
$$+ g' B_\mu \left[\bar{Q} \gamma^\mu \left(\frac{1+\gamma^5}{2}\right) \tau^3 Q + (\text{vector current})\right]$$

Both currents can produce hypercolor pions; to find the amplitudes for this production, note that these currents contain the form (8) and use (10), replacing f_π by a new hypercolor constant F_π. We may then evaluate the contribution of Fig. 4 explicitly:

$$\text{(Fig. 4)} = i\left(i\,\frac{gF_\pi}{2}\,p^\mu\right)\frac{i}{p^2}\left(-i\,\frac{gF_\pi}{2}\,p^\nu\right) \tag{14}$$

This is of the form of the second term in (12) and would yield

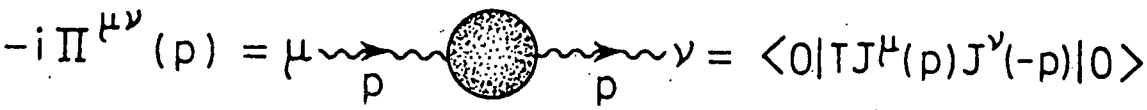

Figure 3. Vector meson self-energy.

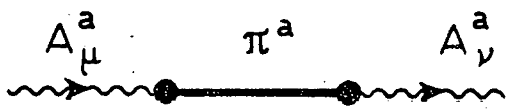

Figure 4. A contribution to $\Pi_{\mu\nu}(p)$ which contains a $1/p^2$.

$\mu_A^2 = (\frac{1}{2}gF_\pi)^2$. However, for the $a = 3$ component of A_μ^a there is a complication: Both A_μ^3 and B_μ couple to the neutral pion, with couplings given in Fig. 5. Thus, the linear combination of these bosons indicated in Fig. 5,

$$\frac{1}{(g^2+g'^2)^{\frac{1}{2}}}\left(-g\,A_\mu^3 + g'B_\mu\right), \tag{15}$$

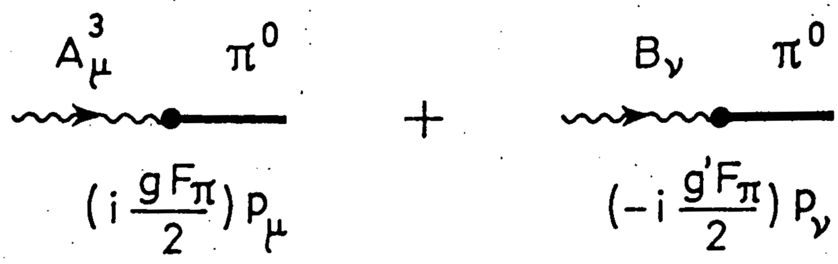

Figure 5. Mixing of Weinberg-Salam bosons with the hypercolor π^0 .

will get a mass equal to $\frac{1}{2}F_\pi\,(g^2+g'^2)^{\frac{1}{2}}$. The orthogonal combination will be left massless; it is the photon.

To see how these results accord with the standard picture of the weak interactions, let us call the massive doublet of vectors $W_\mu^\pm$, the vector (15) Z_μ , and the orthogonal combination A_μ , and label $g'/g = \tan\theta_W$. The results of the above paragraph may be written:

$$\mu_W = \frac{1}{2}\,g\,F_\pi$$

$$Z_\mu = -\cos\theta_W\ A_\mu^3 + \sin\theta_W\ B_\mu \qquad (16)$$

$$\mu_Z = \mu_W/\cos\theta_W \quad , \ \mu_A = 0$$

These results are exact to all orders in the hypercolor interactions. The last two lines of (16) give precisely the relations among the Z , W and photon which are the starting point for deriving the conventional phenomenology of the Weinberg-Salam model. We have obtained these relations, not by adjusting the theory to give them[24)], but by writing the very simplest model of dynamical symmetry breaking, and following our noses.

5. HYPERCOLOR'S PSEUDOSCALAR MESONS

The relations (16) do contain one not-so-innocent feature. To make the first line of (16) agree with the W-boson mass required by phenomenology, we must set

$$F_\pi = 2^{-1/4} G_F^{-1/2} = 250 \text{ GeV} \qquad (17)$$

Thus, the new strong interactions that we postulate must have an extremely large characteristic energy. Scaling from (17) using (11), we would estimate that, for hypercolor,

$$\Lambda_H \sim 800 \text{ GeV} . \qquad (18)$$

The dynamical mass of a hyperfermion is of order Λ_H . The new strong interactions would be expected to produce a new family of hadrons, but the estimate (18) indicates that these hadrons will have masses of order several TeV. The lightest scalar meson, the analogue in this theory of the neutral Higgs meson of the standard Weinberg-Salam model, would be a broad feature like the ϵ of the usual strong interactions, but centered on a mass of 2 TeV. Such particles will remain inaccessible to experiment in the foreseeable future. If we wish to test this theory of SU(2) x U(1) breaking, we need to ask what its manifestations are at energies very small compared to the enormous scale (18). In this section we will examine the question of whether, despite the forbidding size of (18), some mesons built of hypercolored fermions might, nevertheless, be accessibly light.

The Weinberg-Susskind model discussed in the previous section has, unfortunately, no such particles. The breaking of chiral symmetry in this model produced an isospin triplet of massless Goldstone bosons. However, these were seen to combine with the $W^\pm$ and Z through the Higgs mechanism. As a result, these states would not be directly observable; their presence would be felt only through some rather subtle effects on the interactions of W's and Z's .

This situation changes, however, if the Weinberg-Susskind model is at all enlarged. One might imagine writing a model with N , rather than 1 , hyperfermion doublets, that is, with 2N , rather than 2 , hyperfermions. In such a model, the chiral symmetry is no longer (5), the handed isospin SU(2) x SU(2), but now a group SU(2N) x SU(2N). The arguments of the previous section would indicate that this large group of chiral symmetries should be spontaneously broken, with only the vector symmetries preserved. Since there are more broken symmetry currents, one must find more Goldstone bosons, no longer 3 but now $(2N)^2-1$. Of these, only 3 will combine with the $W^{\pm}$ and Z ; the rest will be left as physical pseudoscalar mesons. Before we couple the hypercolor theory to the weak and electromagnetic interactions, these particles are massless; even after we account for this coupling, they will be much lighter than typical hypercolor hadrons. Relatively light pseudoscalars of this type, arising in theories of dynamical symmetry breaking, have been labeled technions[25)]. The very lightest of these were first discussed, under another name, in ref. 5 .

We have introduced the idea of enlarging the simplest hypercolor model as a theoretical possibility. However, we actually expect such an extension to be required in realistic hypercolor models. The reason for this expectation stems from the fact that the mechanism of SU(2) x U(1) symmetry-breaking must do more than give masses to the $W^{\pm}$ and Z ; it must also give current-algebra masses to the quarks and leptons. The mechanism which accomplishes this is described in the next section. For now, we mention that it works by enlarging the gauge-group structure to include a weak coupling between the hyperfermions, which have acquired dynamical masses through breaking of their chiral symmetries, and the ordinary fermions. The large multiplets that necessarily arise, involving at least three doublets each of quarks and leptons together with the hyperfermions, invariably lead to two or more doublets of hyperfermions.

Such enlarged models are complex and, since they normally contain the simultaneous action of several different gauge groups, a bit difficult to understand. To ease the discussion of these models, in this and the next section, we will refer to these various groups using a standard notation, given in Table 1. These groups are the essential components of a hypercolor model, and readers should keep them in mind as they proceed. In an elegantly constructed model, these groups will not all be distinct, but it is useful to separate them for the purpose of explanation. The groups $\mathcal{G}_S$, $\mathcal{G}_A$ will be introduced and explained in the next section.

Table 1: Groups arising in hypercolor models

Gauge Groups

$\mathcal{G}_W$	Weinberg-Salam SU(2) x U(1)
$\mathcal{G}_C$	QCD color
$\mathcal{G}_H$	Hypercolor
$\mathcal{G}_S$	Sideways interaction
$\mathcal{G}_A$	Higher-level strong interaction

Global Symmetries

G_H	Chiral symmetry of hyperfermions
G_{QL}	Chiral symmetry of quarks and leptons

Let us now examine some specific toy models which enlarge the Weinberg-Susskind model, due to Farhi and Susskind[26)] and Dimopoulos[27)]. These authors allow some of the hyperfermions to carry the ordinary color of QCD ($\mathcal{G}_C$) as well as hypercolor ($\mathcal{G}_H$). They utilize a set of 3 hyperfermion doublets, transforming under g_C as a color triplet, to give mass to s, c, μ and ν_μ, and an additional hyperfermion doublet, not coupled to QCD, to give mass to u and d. In all, one has 4 doublets and, therefore, an SU(8) x SU(8) chiral symmetry. The spontaneous breaking of this symmetry to SU(8) produces 63 Goldstone bosons. Three of these combine with the $W^\pm$ and the Z; the rest are physical pseudoscalars, technions.

What are the masses of these particles? In the absence of their coupling to g_C and g_W, they are massless. But the QCD and Weinberg-Salam gauge boson exchanges can contribute masses to these particles; to lowest order[28)], one must compute diagrams of the form of Fig. 6. The magnitude of the mass generated will be of order

$$m_t^2 \sim \alpha_C \Lambda_H^2 , \alpha \Lambda_H^2 \sim (100\ \mathrm{GeV})^2 \qquad (19)$$

Hence, these particles are expected to be an order of magnitude lighter than the hypercolor ρ meson, though still extremely heavy by ordinary standards.

Despite the size of the estimate (19), however, it is worth investigating these particles a bit further. They are Goldstone bosons, related intimately to the currents of broken symmetries, after the fashion of the familiar π and K mesons. This implies that, just as for π and K, many of their properties may be determined from current algebra. It is, in particular, known that the purely electromagnetic contribution to the pion mass, which is of the form of Fig. 6 with an exchanged photon and an external π, may be computed

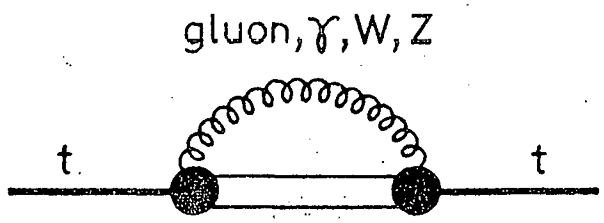

Figure 6. Leading-order contribution to technion masses.

using current-algebraic methods[29]. The contribution of Fig. 6 to technion masses may be computed in the same way [5,8,25,30]. The result of this computation, for this SU(8) x SU(8) model, is shown in Fig. 7. (The overall scale of masses shown here assumes that $\mathcal{G}_H$ = SU(4); for a larger group, the masses will be lower, but in the same ratios.) Dimopoulos points out that this structure, like that of other speculative theories, can be made much richer by simple changes in one's assumptions. In this case, the technical modification of assigning the hyperfermions to a real, rather than a complex, representation of $\mathcal{G}_H$ converts the spectrum of Fig. 7 to that of Fig. 8.

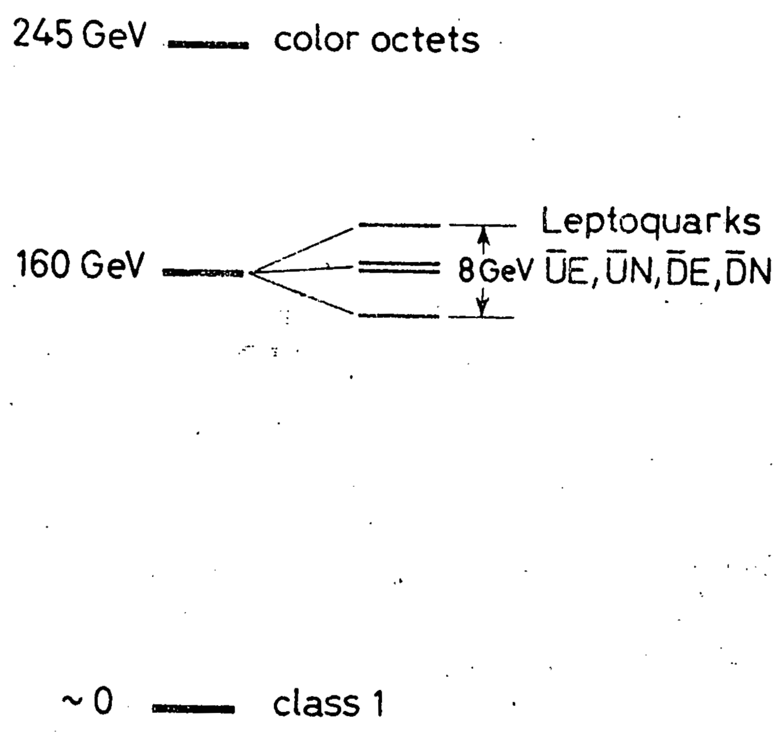

Figure 7. Spectrum of technions in a model of Dimopoulos, from ref. 25. U,D,N,E are hyperfermions with the quantum numbers under color and the electroweak interactions of u,d,ν,e, respectively; in the figure they label the quantum numbers of bound states. Class 1 contains 4 non-exotic mesons, 2 electrically charged, 2 neutral.

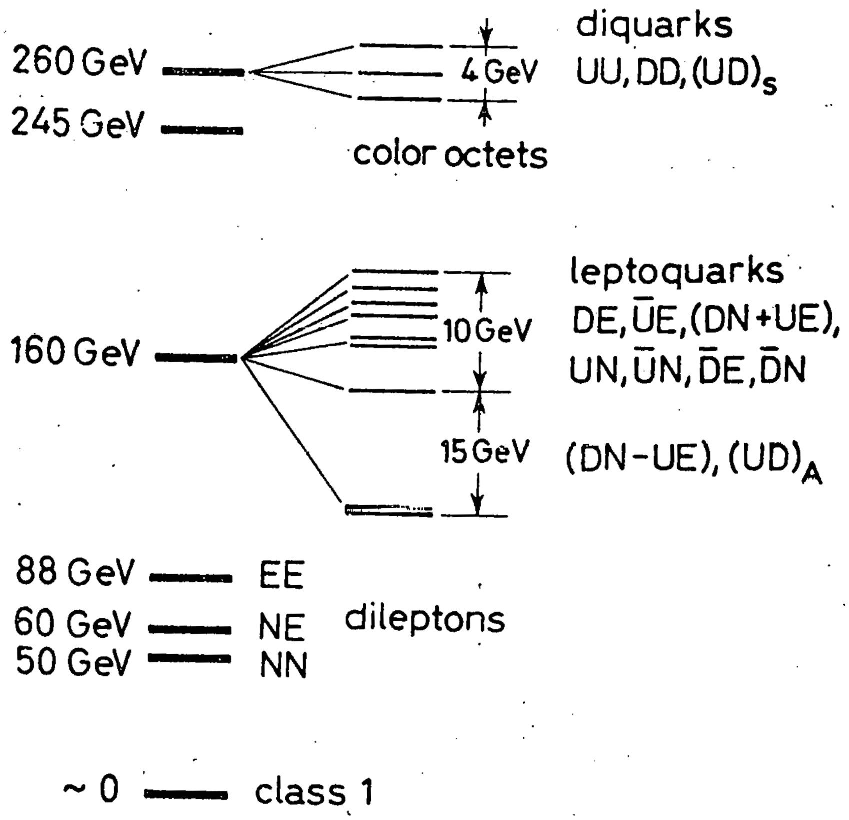

Figure 8. Spectrum of technions in a second model of Dimopoulos. The notation is that of Fig.7.

The masses of most of the particles displayed in Figs. 7 and 8 are large. Still, the qualitative features of the spectra are very weird and this weirdness, following as it does from rather plausible assumptions, merits some comment. Both models contain pseudoscalar mesons which are octets under QCD color; these might be produced at observable rates in gluon-gluon collisions at the TeV-energy $p\bar{p}$ collider envisioned at Fermilab[31)]. Both models contain many pseudoscalars with quark-lepton quantum numbers; in the second model these states alone produce, in e^+e^- annihilation, a step of 5 units of R at a center-of-mass energy of about 300 GeV. The second model also contains a pseudoscalar (EE) of electric charge (−2), which decays into two leptons; it is light enough to be pair-produced at the highest LEP energies.

The most interesting feature of these two figures, however, lies at the bottom of the spectrum. In both variants of this model, there are two neutral technions and one electrically charged pair, all color singlets, for which the $O(\alpha)$ contributions shown in Fig. 6 precisely cancel. This cancellation is due to the group-theoretic and chiral structure of the weak SU(2) x U(1) currents. For example, the contribution of a charged W-boson has two terms, one from the vector and one from the axial vector piece of its coupling, equal in magnitude and opposite in sign. Also, to $O(\alpha)$, we should ignore mixing of A_μ^3 and B_μ (see (13)-(15) and Fig. 4), so that the chiral-cancellation argument still holds. More generally, it has been shown that, for the color-singlet technions, this cancellation occurs in any hypercolor theory constructed along the lines set out here [5,25,30]. In any such theory, there will be at least two charged and two neutral pseudoscalars, and the electroweak contribution to their mass will be at most $O(\alpha^2 \ln \alpha^{-1})$, much less than the estimate (19).

Remarkably, then, the new strong interaction theory we have postulated, whose natural scale of masses is 1 TeV, possesses a few mesons which are accessibly light. We will discuss the masses and phenomenology of these particles in some detail in section 9 .

6. HOW TO GIVE MASSES TO QUARKS AND LEPTONS

In the previous few sections, we have outlined the basic structure of a weak interaction theory with dynamical symmetry breaking. We have introduced a set of new fermions, coupled to a new strong interaction gauge group hypercolor ($\mathcal{G}_H$) as well as the more familiar Weinberg-Salam SU(2) x U(1) ($\mathcal{G}_W$) and, possibly, QCD color ($\mathcal{G}_C$) . The hypercolor strong interactions break the chiral symmetries of these fermions; this in turn breaks the gauge symmetry $\mathcal{G}_W$. We have shown how masses are generated for the W and Z bosons and for various pseudoscalar bound states of the hyperfermions. But so far we have not included the coupling of the hypercolor sector to the ordinary quarks and leptons. In this section we will complete our explication of the structure of a dynamically broken weak-interaction theory by describing this coupling and its effects.

In deciding how to introduce quarks and leptons, there is a special problem that we must confront. The rules for model-building which we set ourselves in section 3 insist that all fundamental couplings of fermions in a dynamically broken theory must be couplings to gauge fields. These couplings conserve fermion helicity and lead to separate handed fermion flavor symmetries of the form of (5). These symmetries, considered for hyperfermions alone and for ordinary fermions alone, were labeled G_H and G_{QL} , respectively, in Table 1 . Some component of these symmetries must be exact, since we have coupled the SU(2) x U(1) gauge bosons to chiral currents. But G_{QL} alone cannot be an exact symmetry, since it prohibits the appearance of quark and lepton masses. In such a world, π^0, K^0, η and D^0 are massless, while the mass of $\pi^{\pm}$ is 35 MeV, due to electromagnetism alone. In the standard Weinberg-Salam theory, these unwanted symmetries are broken by couplings to the elementary Higgs mesons, at a price that we have indicated in section 2. We must be sure that the couplings of quarks and leptons to the hypercolor theory include some alternative mechanism for breaking G_{QL} .

The only way to break G_{QL} in a gauge theory with dynamical symmetry breaking, while maintaining invariance under $\mathcal{G}_W \times \mathcal{G}_C \times \mathcal{G}_H$, is to allow gauge couplings between the light fermions and the hyperfermions. If this new interaction implies that quarks have strong self-interactions at a mass scale of 1 TeV, the arguments of section 4 would indicate that their chiral symmetries are spontaneously broken and that they acquire dynamical masses as large as those of hyperfermions[32]. For this and other good reasons, the new gauge couplings must induce only weak transitions between hyperfermions and ordinary fermions (Fig. 9). One way to implement

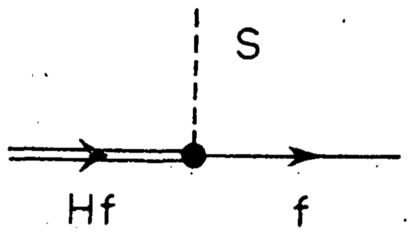

Figure 9. A transition from a hyperfermion to an ordinary fermion mediated by a sideways boson.

this is to introduce two additional new gauge interactions. The first mediates these new transitions. This interaction has been referred to by various authors as the sideways interaction[5] or extended technicolor[17]. We will use the former name here and label its gauge group $\mathcal{G}_S$. Transitions between ordinary fermions and hyperfermions will be weak if $\mathcal{G}_S$ describes a very strongly broken gauge interaction. The second new gauge interaction is strong, and designed to give very large masses to $\mathcal{G}_S$ bosons via dynamical symmetry breaking. This we will label $\mathcal{G}_A$. We remarked in section 4 that not all the new gauge interactions in Table 1 need be distinct. In fact, as we will see shortly, they must not be.

Let us, then, suppose a mechanism similar to that discussed in section 4 by which the $\mathcal{G}_A$ strong interactions break $\mathcal{G}_S$ and give mass to a set of $\mathcal{G}_S$ bosons. By analogy with (16), these masses will be of order

$$\mu_S^2 \sim (g_S F_S)^2 , \tag{20}$$

where g_S is the $\mathcal{G}_S$ coupling constant and F_S is the decay constant (10) associated with $\mathcal{G}_S$ breaking. We will shortly determine that $F_S \gtrsim 30$ TeV. These sideways gauge bosons appear to the hyperfermions as W bosons do to the ordinary fermions: They are heavy on the hyperfermion mass scale and mediate a weak current interaction

$$H' = \sum_{\mathcal{G}_S \text{ bosons } a} g_S^2 \frac{1}{\mu_S^2(a)} J_\mu^{a\dagger} J^{\mu a} \tag{21}$$

The various sideways currents J_μ^a produce exotic transitions of the form of Fig. 9 ; they may also connect hyperfermions to one another.

The existence of transitions between hyperfermions and ordinary fermions tells us that G_H and G_{QL} are not separate invariances of the theory which includes $\mathcal{G}_S$: A rotation of the hyperfermions (in G_H) is not a symmetry of the full theory unless accompanied by a rotation of the quarks and leptons (in G_{QL}). But we have seen already that G_H must be spontaneously broken; the sideways interactions will propagate this into a breaking of G_{QL}. This amounts to an explicit breaking of G_{QL} symmetries, and it manifests itself in the appearance of quark and lepton masses.

The specific diagrams which generate masses are those of the form of Fig. 10 [5,17]. The blob which actually creates the mixing of left-and right-handed fermions is the dynamical mass term acquired

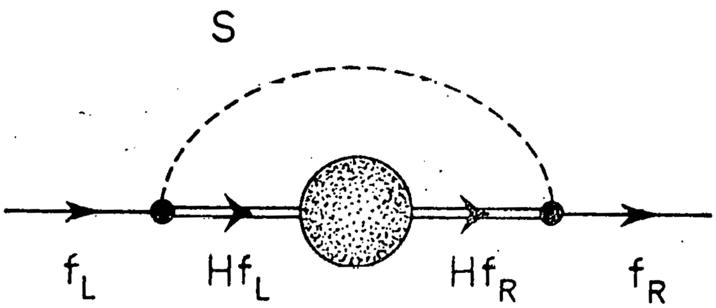

Figure 10. Typical graph producing a mass for a quark or lepton. The shaded blob is the dynamical mass of the hyperfermion.

by the hyperfermions through their chiral symmetry breaking. If this were an ordinary fermion mass, the diagram shown in Fig. 10 would be infinite. Presumably, it could be renormalized, but the renormalization condition would introduce into the theory a new adjustable parameter for each such generated mass. This is precisely the problem in the standard Weinberg-Salam theory which leads to the difficulties described in section 2. But it is known that the dynamically-induced dynamical mass is momentum-dependent and falls off rapidly for large momenta [33]. Thus, in a theory of dynamical symmetry breaking, Fig. 10 is finite and gives a light fermion mass of the order of magnitude

$$m \sim g_S^2 \frac{1}{\mu_S^2} \Delta_H \sim \frac{1}{F_S^2} \Lambda_H^3 \tag{22}$$

where Δ_H is the measure of the hyperfermion condensate defined in (9). In a detailed theory of sideways interactions, the quark and lepton spectra would be computable in terms of their couplings to hyperfermions (group-theoretic factors of O(1)) and the parameter F_S .

We have thus sketched a mechanism by which the addition to a hypercolor model of a sideways gauge coupling and an additional strong interaction can generate quark and lepton masses. To turn this mechanism into a concrete model, it is obviously important to specify what the two new gauge groups are. Unfortunately, the problem of setting up the $\mathcal{G}_S$ breaking turns out to be more difficult than the corresponding problem for $\mathcal{G}_W$, and we have not yet found a reasonable realistic solution. Since the fermion mass spectrum results from a delicate interplay of the $\mathcal{G}_H$ dynamics, the $\mathcal{G}_S$ couplings, and the $\mathcal{G}_S$ breaking, the search for the correct model will be a difficult but intriguing problem for theorists.

Despite our inability to specify $\mathcal{G}_S$ and $\mathcal{G}_A$, we can extract some general constraints on these groups. The mass scale associated with the $\mathcal{G}_A$ interactions is set by its influence on the scale of ordinary fermion masses: If we assume that a typical quark or lepton mass is of order 1 GeV, we find from (22):

$$F_S \sim 30 \text{ TeV} . \tag{23}$$

The associated Λ would be a few times larger, of order 100 TeV . To explain the small current-algebraic masses of u, d and e , it has been proposed[34] that there is a hierarchy of $\mathcal{G}_S$ - breaking scales, $\Lambda_1 \sim 1000$ TeV to produce the lightest fermion masses, and $\Lambda_2 \sim 100$ TeV to produce the O(1 GeV) masses. We believe such hierarchies are also needed to suppress $|\Delta S| = 2$ processes to acceptable levels.

$\mathcal{G}_S$ is actually quite strongly constrained[5]: Since sideways bosons change ($\mathcal{G}_H$-singlet) quarks and leptons into hyperfermions to break G_{QL} , $\mathcal{G}_S$ cannot commute with $\mathcal{G}_H$. But it must also spoil symmetries in G_H which, when spontaneously broken, would have associated with them either exactly massless Goldstone bosons or very light bosons resembling the axion[35]. Such bosons are easily generated in a theory of hypercolor alone; they are certainly not observed. This requirement may be seen to imply that $\mathcal{G}_S$ commutes neither with ordinary color nor electric charge. Thus, $\mathcal{G}_S$ must contain a subgroup which remains unbroken and contains $\mathcal{G}_H \times \mathcal{G}_C$. (As far as we know, $\mathcal{G}_S$ may commute with the weak-interaction SU(2) and, for simplicity, we will assume that it does.) Further, all quarks, leptons and hyperfermions must be contained in at most four multiplets of $\mathcal{G}_S$: a pair, transforming under equivalent representations, forming a doublet under weak SU(2), and at most two multiplets for the weak isosinglet fermions. To ensure that the up and down quarks have different masses, these last two multiplets must transform as inequivalent representations of $\mathcal{G}_S$. The required structure is, then, remarkably compact. $\mathcal{G}_S$ appears as a large group, broken at a mass scale of 100 TeV. Some of its gauge bosons get masses and mediate the sideways interactions. Others remain massless, and at smaller momenta, become strongly coupled to form the hypercolor and color interactions. Quarks, leptons and hyperfermions are unified into multiplets of $\mathcal{G}_S$ [36].

This picture of new interactions is rich with new phenomena. Although the characteristic mass scales are very large, 1 - 100 TeV, several predictions of this general picture apply to lower-energy processes which are currently accessible to study. While masses, cross sections and branching ratios may be calculated with precision only after making specific realistic choices of $\mathcal{G}_H$ and $\mathcal{G}_S$, the order of magnitude of the effects we will describe follow directly from the broad outlines of the theory set out here.

7. TEST # 1: FLAVOR-CHANGING NEUTRAL CURRENTS

In the standard SU(2) x U(1) model, the presence of neutral currents is a consequence of group theory: commuting two oppositely charged weak currents, one finds a neutral current. In the same way, commuting two sideways currents linking quarks or leptons to hyperfermions reveals the presence of additional sideways currents of the form

$$\bar{u}_L\gamma^\mu u'_L + \bar{d}_L\gamma^\mu d'_L \; ; \; \bar{u}_R\gamma^\mu u'_R \; ; \; \bar{d}_R\gamma^\mu d'_R \; ;$$

$$\bar{\nu}_L\gamma^\mu \nu'_L + \bar{e}_L\gamma^\mu e'_L \; ; \; \bar{e}_R\gamma^\mu e'_R \tag{24}$$

In (24), we have used u,d,e,ν as generic labels: u and u′, for example, denote any two charge +2/3 quarks. The usual neutral current is flavor-conserving. But these new currents, which couple to $\mathcal{G}_S$ rather than $\mathcal{G}_W$ bosons, cannot be, if all the various flavors reside in only four $\mathcal{G}_S$ multiplets. (Recall that, in section 6, this requirement was forced on us by the need to avoid axion-like mesons.) $\mathcal{G}_S$ interactions, then, mediate flavor-changing neutral current processes[5]. Rare K^0 decays such as

$$K_L \to \mu^\pm e^\mp \qquad K^+ \to \pi^+\mu^\pm e^\mp \, , \; K^+ \to \pi^+ \nu\bar{\nu} \tag{25}$$

and rare μ decays such as

$$\mu^- \to e^-\gamma \;\; , \;\; \mu^- \to e^- e^+ e^- \tag{26}$$

provide the best tests for the presence of these unusual neutral currents.

These processes are mediated by $\mathcal{G}_S$ gauge bosons whose mass is comparable to that (eq.(20)) encountered in quark and lepton mass generation. This fact permits the following rough estimates of the effective Fermi constants, G_{eff}, for rare K decays and for $\mu \to eee$. ($\mu \to e\gamma$ will be discussed separately.) The most naive guess[5] for G_{eff} is

$$G_{eff} \cong \frac{1}{F_S^{\,2}}\mathcal{R} \cong 10^{-5}\, G_F \mathcal{R} \, , \tag{27}$$

where $F_S \cong$ 30 TeV and $\mathcal{R}$ is a factor containing the sines and cosines of mixing angles which inevitably appear. The small masses of u, d, s and e, μ suggest an even smaller Fermi constant[34,37), roughly

$$G_{eff} \sim (10^{-6} - 10^{-7})\, G_F \mathcal{R} \tag{28}$$

Thus, we expect branching ratios

$$10^{-14} \lesssim B(\text{rare K decay}),\ B(\mu\to eee) \lesssim 10^{-10} \tag{29}$$

The transition dipole moment d for $\mu \to e\gamma$ may be estimated from the graphs of Fig. 11. The blob is, as in Fig.10, the hyperfermion's dynamical mass. As in that case, the momentum dependence of

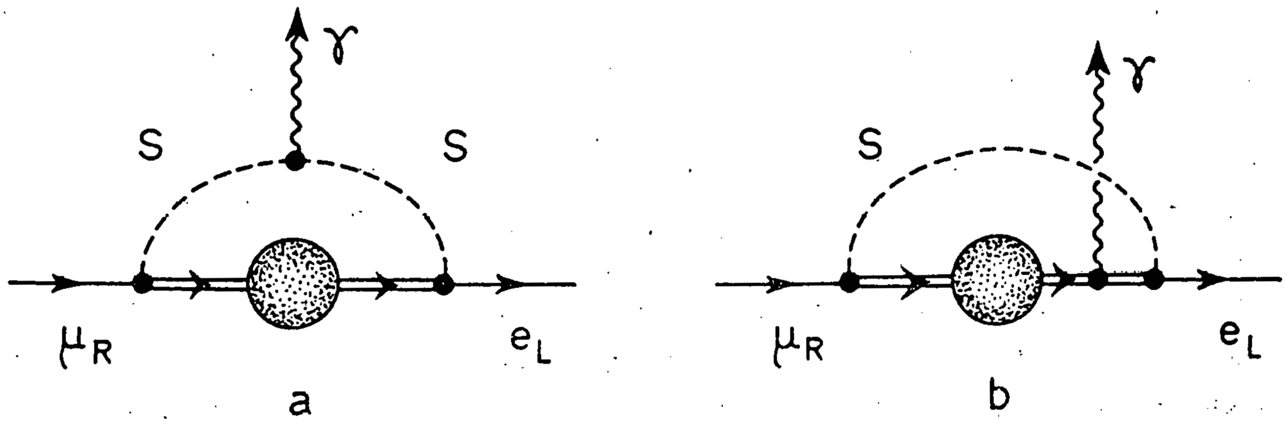

Figure 11. Contributions to $\mu \to e\gamma$. The notation is as in Fig. 10. These graphs assume assume that the relevant $\mathcal{G}_S$ bosons and hyperfermions are electrically charged.

this mass suppresses high momenta in the loop integral. The contribution of Fig. 11a is then roughly $d_a \cong eg_S^2\, \Lambda_H^3/\mu_S^4$ × mixing angle factors. Fig. 11b seems to give a contribution of order $eg_S^2\, \Lambda_H/\mu_S^2$, which is far too large, but a careful analysis shows that this graph's contribution to d is actually of the order of d_a. Thus, we estimate,

$$d \cong e g_S^2 \frac{\Lambda_H^3}{\mu_S^4} \mathcal{R} \cong e \frac{m_\mu}{\mu_S^2} \mathcal{R} \cong (10^{-11}\, e\text{-GeV}^{-1})\, \mathcal{R} \tag{30}$$

We have used (22), noting that the relevant $\mathcal{G}_S$ bosons also determine the μ and e masses. The branching ratio for $\mu \to e\gamma$ is

$$B(\mu \to e\gamma) = \frac{3\pi d^2}{G_F^2 m_\mu^2} \cong 5\mathcal{R}^2 \times 10^{-11} \quad , \tag{31}$$

which is to be compared with the present experimental upper limit[38] of 1×10^{-10}.

All these estimates are crude, but it is tempting that they lie not far below present upper limits. It is worthwhile to push these limits down another order of magnitude or so.

8. TEST # 2: CP-VIOLATION AND THE NEUTRON ELECTRIC DIPOLE MOMENT

Since the discovery of CP-violation in the K^0-$\bar{K}^0$ system[39], there have been many theoretical attempts to explain and understand it[40]. Many of these proposals are attractive and, indeed, remain viable. However, none of them really does explain the origin of CP-violation: To some extent, it is always put in by hand, as phases of scalar-field couplings, for example, in the Kobayashi-Maskawa (KM) model, or by an arbitrary choice of parameters that leads to spontaneous CP-violation in the Higgs models of Lee and Weinberg. Moreover, none of these models seriously addresses the issue of strong CP-violation[35]: As we recalled in section 2, a CP-violating imaginary part of the quark mass matrix contributes directly (eq.(3)) to the neutron electric dipole moment, for which there exist stringent bounds[7]. Why should this particular parameter be especially small?

Dynamically-broken weak interaction theories offer a natural explanation for the origin of weak CP-violation and, possibly, for the absence of strong CP-violation[6]. Because the fermion Hamiltonian is generated solely by gauge interactions, without elementary scalar couplings or fermion bare mass terms, it is natural that it be CP- and T-invariant. The mechanisms of section 4 do not introduce CP-violation in the breaking of $\mathcal{G}_S$ or in the spontaneous breaking of G_H. By this we mean that the pair condensate vacuum of the hyperfermions and quarks, which we label $|\Omega\rangle$, and the $\mathcal{G}_S$ perturbation H', eq.(21), are each, considered separately, invariant under CP transformations. One possibility, however, remains for CP-violation: These two CP transformations can be different and, in fact, conflict, so that the full theory has no CP-invariance. In this section, we will explain this mechanism of CP-violation and its implications. We apologize that our discussion here will be slightly more technical than that of the other sections of these lectures.

Let us begin by discussing the Hamiltonian H_o of hyperfermions, quarks and leptons, in the absence of H'. G_H and G_{QL} are symmetries of H_o but, because of spontaneous symmetry breaking, they are not symmetries of its ground state $|\Omega\rangle$. In fact, $G_H \times G_{QL}$ transformations will rotate this state into other, distinct states of the same energy. More precisely, if W is a unitary transformation in $G_H \times G_{QL}$, $W|\Omega\rangle$ is an equally good ground state of H_o. (The observant reader will note that, since lepton chiral symmetries are explicitly, but not dynamically, broken, leptons are spectators in this discussion and may be ignored.)

Our unperturbed Hamiltonian H_o has, then, a family of degenerate ground states. The perturbation H', which explicitly breaks the symmetry $G_H \times G_{QL}$, will lift this degeneracy and pick out the true ground state of the theory. As Dashen[41] has emphasized, this true ground state is the one for which the vacuum expectation value of H' is a minimum. That is, if

$$E(W) = \langle\Omega| W^{-1} H' W |\Omega\rangle \qquad (32)$$

is minimized by $W = U$, the true vacuum is $U|\Omega\rangle$. Equivalently, we can take the true vacuum to be simply $|\Omega\rangle$ and regard $H'_U = U^{-1} H' U$ as the true perturbation.

If $|\Omega>$ and H' have the same CP-transformation, H'_U will be CP-invariant along with $|\Omega>$ if and only if U is a real unitary transformation, $U = U^*$. This is a possibility, because the CP-invariances of $|\Omega>$ and H' imply that $E(W) = E(W^*)$. Hence, E has T-invariant extrema $E(W_o)$ with $W_o = W_o^*$. But it is also possible that E is minimized away from such a point, at U and U^* such that $U \neq U^*$. This corresponds to a spontaneous CP- and T-violation; indeed, this violation disappears for energies well above 100 TeV where $\mathcal{G}_S$ - symmetry is restored. We see now that the origin of CP-violation lies in the pattern of $\mathcal{G}_S$-breaking: This pattern is, in principle, unique, given $\mathcal{G}_A$; it determines H_o and H', and therefore it ultimately determines whether or not spontaneous CP-violation occurs.

It is useful to parametrize W more explicitly. A set of fermions with the same quantum numbers under $\mathcal{G}_H \times \mathcal{G}_C$ will be mixed by such a chiral symmetry. Let us label each such set of fermions by an index ρ. All the quarks belong to a single set; the hyper-fermions must belong to at least one different set. We may then label the quark and hyperfermion fields by ψ^{ρ}_{Lr}, ψ^{ρ}_{Rr}, where r is a flavor index. To guarantee that $\mu_W / \mu_Z = \cos\theta_W$ and that electric charge is conserved after $\mathcal{G}_W$ is spontaneously broken, it is assumed that the flavor symmetry $G_H \times G_{QL}$ does not transform weak doublets ψ^{ρ}_{Lr} into weak singlets ψ^{ρ}_{Rr}. A transformation W in $G_H \times G_{QL}$ then transforms

$$\psi^{\rho}_{Lr} \rightarrow \sum_{r'} W^{L}_{\rho rr'} \psi^{\rho}_{Lr'} \quad ; \quad \psi^{\rho}_{Rr} \rightarrow \sum_{r'} W^{R}_{\rho rr'} \psi^{\rho}_{Rr'} \tag{33}$$

W^L_ρ, W^R_ρ are unitary matrices. The overall phases of these matrices are constrained by the requirement that the currents generating these symmetries are free of axial vector anomalies arising from either the $\mathcal{G}_H$ or $\mathcal{G}_C$ gauge interactions. If the vector flavor symmetries (those with currents of the form (7)) are respected by the condensation, $E(W)$ depends only on the combinations $W_\rho \equiv W^R_\rho W^{L\dagger}_\rho$ and $W^\dagger_\rho$.

The condition that $U = W$ is an extremum of $E(W)$ is the following[6]:

$$(M_\rho - M_\rho^\dagger)\,\Delta_\rho = i\,\nu\,T_\rho\,\mathbb{1}_\rho \tag{34}$$

Here ν is a real quantity of dimension $(\text{mass})^4$, the same for all sets ρ. T_ρ is a dimensionless group-theory factor, $\mathbb{1}_\rho$ is the unit matrix in the flavor space. Δ_ρ is the condensate defined in (9); this will be different for each fermion species ρ : $\Delta_\rho = O\big((1\text{GeV})^3\big)$ for quarks and $O\big((1\ \text{TeV})^3\big)$ for hyperfermions. The "mass" M_ρ is defined by

$$M_{\rho r r'}\,\Delta_\rho = \sum_{s,s'} \left[W^{L\dagger}_{\rho r s} \frac{\partial E(W)}{\partial W_{\rho s' s}} W^{R}_{\rho s' r'} \right] \Bigg|_{W=U} \tag{35}$$

The right-hand side of (34) arises from the phase constraint on the W_ρ's ; the quantity ν is the same for all ρ because $\mathcal{G}_H$ and $\mathcal{G}_C$ are both contained in a single larger group $\mathcal{G}_S$. With $U = \{U_\rho\}$ chosen to minimize E, the remaining arbitrariness in U^L_ρ and U^R_ρ is removed by requiring that M_ρ be diagonal. Then $(\nu T_\rho / 2\Delta_\rho)\,\mathbb{1}_\rho$ is the imaginary part of M_ρ.

For the set ρ of quarks $(\rho = q)$, M_q is almost the quark mass matrix m_q. It is possible to show[6)] that M_q differs from m_q by terms of relative order $(\Lambda/\Lambda_H)^3$ or $(\alpha_S(\Lambda_H)/2\pi)\cdot(\Lambda/\Lambda_H)^2$, where Λ is the QCD scale; both terms are $\sim 10^{-9}$. This has important consequences for CP-violation. There are three possibilities for the solution of (34):

(1) <u>$U = U^*$</u>. In this case, CP is not spontaneously broken; in particular, $\nu = 0$.

(2) <u>$U \neq U^*$, $\nu \neq 0$</u>. In this case, strong CP-violation occurs. The natural scale of ν/Δ_q is m_u, so we expect a neutron electric dipole moment $\cong 10^{-15}$ e-cm [16)], nine orders of magnitude larger than the experimental upper bound.

(3) <u>$U \neq U^*$, $\nu = 0$</u>. This case occurs if E is minimized by a complex U in $G_H \times G_{QL}$ even if the phase constraint is removed. In this case, there is no strong CP problem.

The last of these cases is the only realistic one. Let us ask first its implication for the size of the neutron electric dipole moment d_N. $Im(m_q)$ is at most $\sim 10^{-9} m_u$ and contributes $\sim 10^{-24}$ e-cm to d_N. CP-violating phases of order 1 in broken $\mathcal{G}_S$ interactions involving four quarks will contribute $\sim eM_N\, m_u/\Lambda_H^3 \sim 10^{-25}$ e-cm to d_N. Rough estimates of other contributions are comparable to these two. Together, they imply

$$d_N \sim (10^{-24} - 10^{-25})\,e-cm\,, \tag{36}$$

just below the present upper limit of 2×10^{-24} e-cm [7]. As a comparison, the standard KM model predicts a value of $d_N \sim 10^5$ times smaller than this [42].

What about CP-violation in the electroweak interactions? Since the quark matrices $U_q^{L,R}$ commute with electric charge, they will be block-diagonal into blocks for up-and down-quarks, $U_u^{L,R}$ and $U_d^{L,R}$. The U^L appear in the charged weak currents, and the Cabibbo-Kobayashi-Maskawa mixing matrix is $(U_u^{L\dagger} U_d^L)$. If this matrix contains phases which cannot be removed by an allowed redefinition of the fermion fields, there is CP-violation in the electroweak interactions, and it appears exactly as in the KM model. In particular, the electroweak contribution to d_N is $\sim 10^{-29}$ e-cm.

In the same way, CP-violating phases appear in the effective 4-quark terms generated by the sideways interactions. Apart from d_N, their most notable effect is in the $K^0-\bar{K}^0$ system. If $|\Delta S| = 2$ terms, such as $\bar{s}\gamma_\mu d\,\bar{s}\gamma^\mu d$, appear, then a naive estimate of their strength would be as in (28): $G_{eff} \sim (10^{-6} - 10^{-7})G_F\,\mathcal{R}$, with the phase of $\mathcal{R}$ of order one. A PCAC calculation of the contribution of such terms to the CP-violation parameter ε gives $\varepsilon \sim 1$, 1000 times too large. It is not clear what mechanism might give the required suppression. The naive $\mathcal{G}_S$ contribution to the parameter ε' is $\lesssim 10^{-6}$. As with most of the rate estimates presented here, an explicit realistic model is needed to pin down the details.

9. TEST # 3. ACCESSIBLE HIGGS MESONS

We saw in section 5 that at least two charged and two neutral pseudoscalar bound states of hyperfermions escape acquiring the large masses which characterize the hypercolor theory. They are, in fact, sufficiently light that they will be produced at currently planned accelerators; it is likely, in fact, that they are already being produced at PETRA and PEP energies. The appearance of these particles is the most obvious and most characteristic manifestation of dynamical symmetry breaking. We devote this section to a discussion of their phenomenology. We will refer to these particles as the $P^{\pm}$, P^{0}, P'^{0} [5].

Our first task is to estimate the masses of these particles. We explained in section 5 that the order - α effects which give mass to these mesons (Fig. 6) cancel. Thus, these mesons will receive mass only from order - α^2 electroweak effects and from sideways interactions. Let us consider first the higher-order electroweak effects: If one includes in Fig. 6 the effect of the Z^0 mass, one finds a contribution to the mass $m_{\pm}$ of the $P^{\pm}$ [5,25,30]:

$$(m_{\pm}^2)_W = \frac{3\alpha}{4\pi}\,\mu_Z^2\,\ln\frac{\Lambda_H^2}{\mu_Z^2} = (5-8\ \text{GeV})^2 \qquad (37)$$

This effect is of order α^2, since $\mu_Z^2 \sim \alpha F_\pi^2$. It can be shown that, if the hypercolor chiral symmetry G_H is simply a handed SU(N) x SU(N) symmetry, this is the only contribution of this order[25]. (In more complex models, one might find additional terms.) The uncertainty in (37) arises only from the ambiguity of Λ_H. The electroweak interactions may be seen to contribute nothing at all to the masses of P^0 and P'^0.

Since we lack a detailed theory of the sideways interactions, we cannot compute the $\mathcal{g}_S$ contribution to these masses in the same precise manner. However, it is possible to make a reliable estimate of this contribution. The effect of the interaction H', eq.(21), in producing P masses may be related, using current algebra, to a vacuum expectation value of another operator of the form of H', actually a commutator of H' with chiral charges. This allows us to estimate any one m^2 of the four masses by

$$(m^2)_S \approx F_\pi^{-2} \langle\Omega|\, g_S^2\,\frac{1}{\mu_S^2}\,(\bar{\psi}_L\psi_L)(\bar{\psi}_R\psi_R)\,|\Omega\rangle \qquad (38)$$

The largest contribution comes when all four fermion fields correspond to hyperfermions; then we may estimate (38) using (9) and find

$$(m^2)_S \sim \frac{1}{F_\pi{}^2} \frac{1}{F_S{}^2} \Delta_H^2 \sim \left(\frac{1}{F_\pi F_S} \Lambda_H^3\right)^2$$

$$\sim (10 - 30\ \text{GeV})^2 \tag{39}$$

A model which accounts for the small masses of u, d, e might introduce into (39) a large value of F_S and so produce a smaller $(m^2)_S$.

Using $m = [(m^2)_W + (m^2)_S]^{\frac{1}{2}}$, we conclude that the P-meson masses are

$$m_\pm \cong m_o \cong m_o' \cong 10 - 30\ \text{GeV}, \quad \text{if } (m_\pm^2)_S \gg (m_\pm^2)_W$$

$$m_\pm \cong 5 - 8\ \text{GeV}; \quad m_o \cong m_o' \cong 2 - 5\ \text{GeV}, \quad \text{if } (m_\pm^2)_S \ll (m_\pm^2)_W \tag{40}$$

As far as we know, no model with elementary Higgs bosons requires charged Higgs mesons in this mass range. If such mesons are found, the case for dynamical symmetry breaking will be particularly compelling. Moreover, the relative masses of $P^\pm$ and P^0, P'^0 will give important information on the scale F_S appearing in (39).

The $P^\pm$ is produced copiously in e^+e^- annihilation. Since this particle is composite, one might expect its coupling to the photon to contain a form factor; however, effects of this form factor become noticeable only for photon $Q^2 \gtrsim (1\ \text{TeV})^2$. At ordinary energies, the $P^\pm$ couples to the photon like an elementary spinless meson. Its contribution to R in e^+e^- annihilation is

$$\Delta R = \frac{1}{4}\beta^3 \quad , \quad \beta = \left(1 - \frac{4m_\pm^2}{Q^2}\right)^{\frac{1}{2}} \tag{41}$$

It is produced with an angular distribution

$$\frac{d\sigma}{d\Omega} \propto \sin^2\theta \tag{42}$$

Since this distribution peaks at $\theta \sim 90^0$, the fraction of reconstructable events may be greatly increased [43].

The $P^\pm$ decays primarily into pairs of quarks or leptons. Since the $P^\pm$ contains a combination of hyperfermions orthogonal to those of the hypercolor pion, the particle which appears in Fig. 4 and combines with the $W^\pm$, this particle cannot decay through a single $W^\pm$ (Fig. 12a). It can decay, however, by exchanging a sideways boson between its hyperfermion constituents, allowing them to turn into ordinary fermions (Fig. 12b). (Theorists should note that Fig. 12b is intimately related to Fig. 10 through current algebra.) The process of Fig. 12b yields a decay rate of the form

$$\Gamma(P^+ \to f\bar{f}') = \frac{1}{4\pi}\,\lambda^2\left[1 - \left(\frac{m_f - m_{f'}}{m_\pm}\right)^2\right] p \tag{43}$$

where p is the momentum of a final-state fermion and λ is a dimensionless amplitude, computed from figure 12b, of order

$$\lambda \sim \frac{g_S^2}{\mu_S^2}\,\Lambda_H^2 \sim \left(\frac{m_{ff'}}{\Lambda_H}\right) \tag{44}$$

where we have used (22). We expect $m_{ff'}$ here to be the larger of

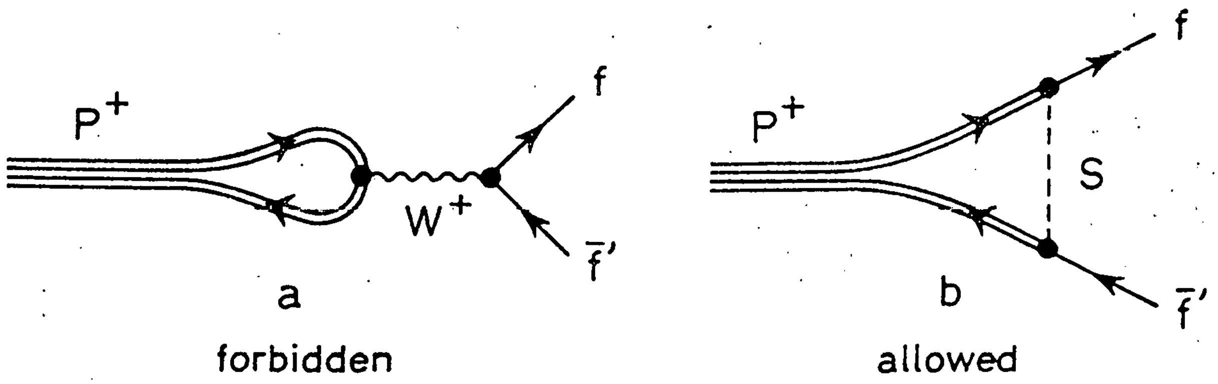

Figure 12. Possible contributions to $P^+ \to f\bar{f}'$.

the masses of the product fermions, but this estimate is sensitive to effects of mixing angles. The amplitude (44) depends on the hypercolor scale as $(\mu_W^2)^{-\frac{1}{2}}$; this process is, then, semi-weak, like the decay of an elementary Higgs scalar in the standard Weinberg-Salam model.

A specific theory of the sideways interactions would predict the amplitudes λ for various final states. Lacking such a theory, we can still offer some general comments on the decay modes of the $P^{\pm}$. These general comments also apply to decay modes of elementary charged Higgs scalars, which, in the standard model, might, but need not, exist. Conversely, signatures for elementary charged Higgs scalars[44)] are, equally well, signatures for the $P^{\pm}$.

The amplitude (44) favors decay into high-mass fermions. Thus, we expect the major decay modes of a P^+ which might be found at PETRA or PEP to be

$$\begin{aligned} &P^+ \rightarrow \tau^+ \nu && \text{(branching fraction } = x) \\ &P^+ \rightarrow \bar{b}c\,,\ \bar{b}u && \text{(branching fraction } = (1-x)) \end{aligned} \qquad (45)$$

Naively, one would expect $x \sim (m_\tau/m_b)^2 \sim (0.1)$, though mixing angles in the amplitude (44) might alter this result drastically. Fortunately, the direct production of $P^{\pm}$ in e^+e^- annihilation leads to three different types of distinctive events, depending on the decay channels:

(1) The process $e^+e^- \rightarrow P^+P^- \rightarrow (\bar{b}c)(b\bar{c})$ produces nonplanar, high sphericity events, a class of events to which experiments at PETRA are already very sensitive.

(2) The process $e^+e^- \rightarrow P^+P^- \rightarrow (\tau^+\nu)(\tau^-\bar{\nu})$ produces an excess of μe events, enhancing the signal from direct τ production, concentrated at large missing energy.

(3) The process $e^+e^- \rightarrow P^+P^- \rightarrow (\bar{b}c)(\tau^-\bar{\nu})$ produces a distinctive signal of a high multiplicity jet recoiling against a lepton.

Experiments collecting data at 30 GeV might, if we are unlucky, produce a plot similar to Fig. 13, enabling them to rule out $P^{\pm}$, within a fixed mass range, independently of the value of x.

We should note, however, that while the decay pattern (45) is the most likely situation, unusual circumstances could lead to stranger possibilities. Here are two such possibilities[45]: First, if

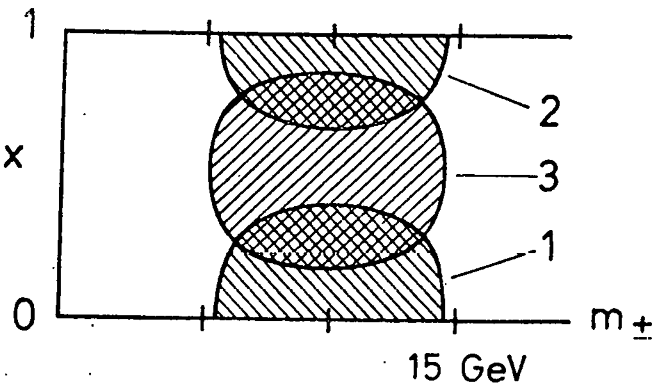

Figure 13. Result of a hypothetical unsuccessful search for $P^{\pm}$. Each shaded region denotes the region excluded by the non-observation of a particular type of unusual event.

$m_{\pm} < m_b$, the second decay path of (45) will be forbidden. It will be replaced by decays to charmed final states, which also yield nonplanar events. In this case, however, $\bar{b} \to P^- u$ should be a major decay mode of the b quark. Secondly, if, for some reason, both decays (45) of $P^{\pm}$ were suppressed, one might see the second-order-electroweak decay

$$P^+ \to \gamma W^+ \to \gamma \ell^+ \nu \tag{46}$$

where ℓ^+ is a lepton. This leads to $e\mu$ events containing 2 hard photons, a most unusual signature.

The P^0 and $P^{\prime 0}$ are less conspicuous. Here we will say only that they appear qualitatively like the neutral Higgs meson of the standard Weinberg-Salam theory; the most important differences are that P^0, $P^{\prime 0}$ do not couple to $f\bar{f}'$ exactly as $G_F^{1/2}(m_f+m_{f'})$ and that the P^0's are pseudoscalar and not scalar particles. Signatures of the neutral Higgs meson[46] apply also to the P^0's. The most promising way to find P^0 and $P^{\prime 0}$, proposed by

Wilczek[47] for the standard Higgs, is to search for the decay

$$V \rightarrow \gamma P^0 \tag{47}$$

where V is the Υ resonance or the analogous $(t\bar{t})$ state (once it is found). The signature is the monochromatic photon against which the P^0 recoils. The branching ratio for (47) is

$$\frac{\Gamma(V \rightarrow \gamma P^0)}{\Gamma(V \rightarrow \mu^+\mu^-)} = \frac{G_F M^2}{4\sqrt{2}\,\pi\alpha}\,\mathcal{R}\,\left(1 - \frac{m_o^2}{M^2}\right) \tag{48}$$

where M is the mass of the V and $\mathcal{R}$ is a factor containing mixing angles. Some values of (48) are given in Table 2 .

Table 2: Branching ratios for $V \rightarrow \gamma P^0$

V	$\frac{\Gamma(V \rightarrow \gamma P^0)}{\Gamma(V \rightarrow \mu^+\mu^-)}$	$B(V \rightarrow \gamma P^0)$
$\psi(3.1)$	8×10^{-4}	7×10^{-5}
$\Upsilon(9.5)$	8×10^{-3}	3×10^{-4}
$\Xi(36.)$	1×10^{-1}	1×10^{-2}

This table ignores phase space suppression and sets $\mathcal{R} = 1$. Ξ is a hypothetical $(t\bar{t})$ resonance at 36 GeV.

Should the P mesons be too heavy to be produced at PETRA and PEP, they can be produced at the Z^0 resonance. The pattern of P couplings to the Z^0 is simple and independent of assumptions about the sideways interactions[48]. To discuss Z^0 couplings, it is useful to normalize to $\Gamma(Z^0 \rightarrow \nu\bar{\nu})$; though this is not directly observable, it is relatively independent of the value of $\sin^2\theta_W$. For comparison,

$$\frac{\Gamma(Z^0 \rightarrow \mu^+\mu^-)}{\Gamma(Z^0 \rightarrow \nu\bar{\nu})} = (1 - 4\sin^2\theta_W + 8\sin^4\theta_W) \approx (0.5) \tag{49}$$

For the $P^{\pm}$,

$$\frac{\Gamma(Z^0 \to P^+P^-)}{\Gamma(Z^0 \to \nu\bar{\nu})} = \frac{1}{2}(1-2\sin^2\theta_W)^2\beta^3 \approx (0.15)\beta^3 \qquad (50)$$

where β is as in (41). The branching ratio to P^+P^-, if there are only 6 quarks and no new leptons, should be

$$BR(Z^0 \to P^+P^-) \cong 0.8 \times 10^{-2} \qquad (51)$$

The angular distribution is, again, (42). In general, for any of the technions of section 5 whose pair-production thresholds lie below the Z^0,

$$\frac{\Gamma(Z^0 \to T\bar{T})}{\Gamma(Z^0 \to \nu\bar{\nu})} = \frac{1}{2}(I^3 - 2Q\sin^2\theta_W)^2\,\beta^3 \qquad (52)$$

where I^3 and Q denote the isospin and electric charge. Two processes expected for a theory with elementary charged Higgs mesons:

$$\begin{aligned} &Z^0 \to P^+W^-,\ P^0 Z^0 \\ &Z^0 \to P^0P^0,\ P'^0P'^0 \end{aligned} \qquad (53)$$

do not occur in a dynamically broken theory; hence Z^0 physics allows one to discriminate a more standard picture of charged Higgs mesons from the one given by dynamical symmetry breaking.

10. CONCLUSIONS

We have presented a physical mechanism for understanding the pattern of weak-interaction gauge symmetry breaking. It also offers a natural explanation for CP-violation, and for the absence of strong CP-violation. The price we must pay is a hierarchy of new gauge interactions whose characteristic scales are 1 TeV and 100 TeV. Remembering that new physics at higher energy scales has always been a fact of life, we feel this is not too great nor surprising a price.

Although a realistic model incorporating dynamical symmetry breaking is yet to be constructed, enough is known about its general structure and consequences to test these ideas now. We have predicted: Flavor-changing neutral currents mediating rare K and μ decays at rates which may be measured with a modest improvement in sensitivity; a neutron electric dipole moment $\sim (10^{-24} - 10^{-25})$ e-cm; and novel pseudoscalar mesons, with masses between 5 and 30 GeV, readily detectable in e^+e^- annihilation. By themselves, positive results in any of these experiments would be very exciting. More exciting still, they would give us a good look into the world above 1 TeV.

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