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IMPLICATIONS OF DYNAMICAL SYMMETRY BREAKING FOR HIGH ENERGY EXPERIMENTS

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IMPLICATIONS OF DYNAMICAL SYMMETRY BREAKING FOR HIGH ENERGY EXPERIMENTS *

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Abstract

A scenario of dynamical symmetry breaking as an alternative to the canonical Higgs mechanism with elementary spin-0 fields is described, and its implications for high energy experiments contrasted with those of the canonical theory. The potential role of e^+e^- annihilation physics in unravelling the nature of spontaneous symmetry breaking is emphasized.

Introduction

The problem of generating fermion and gauge boson masses in the framework of Quantum Flavor Dynamics, QFD, is one of the least understood aspects of these theories. Hardly was the jubilation over the phenomenological successes of the unified electroweak $SU(2)_L \otimes U(1)$ theory over, then it was noticed that the mechanism for spontaneous symmetry breaking in the form of explicit spin 0 Higgs fields is one of the less attractive features of the Weinberg-Salam theory [1]. Ever since there has been a growing suspicion that the Higgs mechanism is more of a phenomenological prop rather than a fundamental mechanism in nature [2].

The reason of distrusting the Higgs mechanism as a fundamental phenomenon stems from the large mass scales involved in the unified theories of electroweak and strong interactions, the so-called Grand Unified Theories, GUTS. The different scales involved in such theories can be exemplified through the decay rate of the neutron (the ordinary β-decay) and of the proton. In GUTS, the former is mediated by the $W^{\pm}$ exchange and the latter by the very massive X bosons. Thus, there are two scales involved in these theories and there are two Higgs multiplets having vacuum expectation values, v.e.v's, typically of order $\langle \varphi_W^0 \rangle \sim (G_F^{-1})^{1/2} \simeq$ 250 GeV and a much higher one $\langle \varphi_x^0 \rangle \sim (G_x^{-1})^{1/2} \simeq 10^{14} - 10^{15}$ GeV. Thus, a ratio of $\sim 10^{13}$ has to be accounted for. An obvious proposal would be to have a potential V with minimum at

$$\langle \varphi_x^0 \rangle \sim 10^{13} \langle \varphi_W^0 \rangle \qquad (1.1)$$

and

$$V(\varphi_W, \varphi_x) = \lambda_1 \left(\phi_x^2 - \langle \varphi_x^0 \rangle^2 \right) + \lambda_2 \left(\varphi_W^2 - \langle \varphi_W^0 \rangle^2 \right) \qquad (1.2)$$

However, there are quantum radiative corrections to V which introduce couplings between the Higgs fields φ_x and φ_W through the gauge bosons to which both the Higgs bosons couple. The additional contribution to V is $\sim \frac{g^2}{8\pi^2} \varphi_W^2 \varphi_x^2$. If one now minimizes the potential it leads to a ratio

$$\langle \varphi_W^0 \rangle / \langle \varphi_x^0 \rangle \sim g^2/4\pi \lambda_2^{1/2} \sim \alpha \qquad (1.3)$$

instead of 10^{-13}, which is unacceptable. Thus, in each order one has to go back and retune the parameters in V so that (1.1) is maintained. This retuning, which for the present example is 1 part in 10^{26}, is unnatural and hence is the basis of doubting the philosophy of introducting explicit Higgs multiplets with vev's, differing by so many orders of magnitude [3].

Let us recall that the Higgs mechanism was introduced to give masses to the gauge bosons, which if the underlying theory is renormalizable, can only get masses through the spontaneous symmetry breaking mechanism. This is a general consequence. What is not in general true but holds for the $SU(2)_L \otimes U(1)$ theory [1] is that the fermion masses must also be generated spontaneously.

There is another example of spontaneous symmetry breaking, namely the spontaneous breaking of chiral symmetry in Quantum Chromodynamics, QCD [4]. For example, in massless QCD with two flavors u,d, there is a global chiral $SU(2)_L \otimes SU(2)_R$ symmetry. At some scale typical of strong interactions, Λ_{QCD}, the QCD coupling constant $\alpha_S = g^2/4\pi$ becomes strong, i.e. of order 1. It is then generally assumed, though nobody has proven it, that the $SU(2)_L \otimes SU(2)_R$ symmetry is broken down to $SU(2)_{L+R}$ due to the non-perturbative mechanism of chiral symmetry breaking

$$\langle \bar{u}u + \bar{d}d \rangle_0 \neq 0 \tag{1.4}$$

leading to a triplet of Goldstone bosons, the ordinary pions of the hadronic world. The condensate $\langle \bar{u}u + \bar{d}d \rangle_0$ has dimensions $(\text{mass})^3$ and in QCD $\langle \bar{u}u + \bar{d}d \rangle_0 = (250\ \text{MeV})^3$.

Thus, there is a dimensional quantity in QCD, namely $\langle \bar{q}q \rangle_0$, which is due to the spontaneous breakdown of chiral symmetry, and perhaps the mechanism which generates this dimensional quantity could be utilized to generate masses for the gauge bosons and fermions.

Weinberg and Susskind independently have demonstrated how to generate the $W^{\pm}$ and Z^0 boson masses using the spontaneous chiral symmetry breaking mechanism (5) Following the example of QCD with 2 flavors, they introduced a doublet of heavy quarks

$$\psi = \begin{pmatrix} U \\ D \end{pmatrix} \tag{1.5}$$

With $m_U = m_D = 0$, the global chiral symmetry group is $G = SU(2)_L \otimes SU(2)_R$. At some scale typical of the new strong interactions, which we call Hypercolor, HC, the Hypercolor gauge coupling constant becomes strong, i.e. $\alpha_{HC} \simeq 1$ at $Q \sim \Lambda_{HC}$. The non-perturbative HC strong interactions come into play and G is broken spontaneously by the HC condensate.

$$\langle \bar{U}U + \bar{D}D \rangle_0 \neq 0 \tag{1.6}$$

to $SU(2)_{L+R}$. This symmetry breaking is accompanied by a triplet of exact Goldstone bosons, π'^+, π'^-, π'^0, which are absorbed by the weak interaction quanta W^+, W^- and Z^0, respectively, which in turn become massive.

To be precise, the Goldstone bosons $\pi'^{\pm}, \pi'^0$ couple to the weak currents via the coupling strengths

$$\begin{aligned} \langle 0 | J_\mu^{\pm,0} | \pi'^{\pm,0}(q) \rangle &= F_{\pi'}\, q_\mu \\ \langle 0 | J_\mu^{Y} | \pi'^{0}(q) \rangle &= F_{\pi'}\, q_\mu \end{aligned} \tag{1.7}$$

with $J_\mu^{\pm,0}$ being the $SU(2)_{\text{left}}$ and J_μ^Y the hypercharge current. The $W^{\pm,3}$ couple to the current $J_\mu^{\pm,0}$ with the gauge coupling $g_2/2$ and the current J_μ^Y couples to B with the gauge coupling $g_1/2$.

Since the chiral group G is broken spontaneously in such a way that an $SU(2)_{L+R}$ symmetry is still preserved, characterized by the equality of the coupling constant $F^1_{\pi'} = F^2_{\pi'} = F^3_{\pi'} \equiv F_{\pi'}$, the Goldstone boson-weak current coupling leads to the following mass matrix for the gauge bosons $W^{\pm,3}$ and B

$$m^2 = \begin{pmatrix} g_2^2 & g_1 g_2 \\ g_1 g_2 & g_1^2 \end{pmatrix} \frac{F_{\pi'}^2}{4} \tag{1.8}$$

which has the eigenvalues

$$\begin{aligned} m_\gamma^2 &= 0 \\ m_Z^2 &= (g_1^2 + g_2^2)\, F_{\pi'}^2/4 \end{aligned} \tag{1.9}$$

and one retains the ratio obtaining in the standard model:

$$\frac{m_W^2}{m_Z^2} = \frac{g_2^2}{(g_1^2 + g_2^2)} = \cos^2\theta_W \tag{1.10}$$

where θ_W is the Glashow-Salam-Weinberg weak angle. Thus the equality of the coupling constants $F^i_{\pi'}$ (a consequence of the residual SU(2) symmetry) ensures the weak $\Delta I = 1/2$ rule. In the usual Higgs mechanism, the Higgs potential has an O(4) invariance broken down to O(3) as a result of spontaneous symmetry breaking, which ensures Eq.(1.9). In both the canonical Higgs mechanism [1] and the Dynamical Scenario [5] Eq.(1.10) receives $O(\alpha)$ radiative corrections.

Let us go back to the question of gauge herarchy, i.e. the problem of generating vastly different mass scales in a natural way. In the usual GUTS, these scales come about by the vastly differing v-e-v's of the Higgs fields. However, it is instructive to see how the scale of the global chiral symmetry breakdown of QCD, Λ_{QCD}, is set in GUTS. To be specific let us take the SU(5) group which is broken at a mass scale M to the observed group $SU(3)_c \otimes SU(2)_L \otimes U(1)$. Now SU(5) is an asymptotically free theory and the coupling constant, $\alpha_s(Q^2)$, has a logarithmic dependence on Q^2. In the case of SU(5) one has

$$\partial \alpha_s(Q^2)/\partial \ln(Q^2) = \beta(\alpha_s) = -11/4\pi \, \alpha_s^2 + \cdots \qquad (1.11)$$

At the scale M, the QCD sector has a small coupling constant. However, when Q^2 decreases $\alpha_s(Q^2)$ increases and at some Q_o^2 the QCD coupling constant becomes large, $\alpha_s(Q_o^2) \sim 1$; QCD non-perturbative effects set in and chiral symmetry breaking takes place. The scale Λ_{QCD}, at which this symmetry breaking occurs can be obtained by integrating Eq.(1.11) and finding the value of Q^2 at which $\alpha_s \simeq 1$. In this way one obtains

$$\Lambda_{QCD} \sim M \exp(-11/8\pi\alpha_o) \qquad (1.12)$$

Thus, for a reasonable value of α_o, $\frac{\Lambda_{QCD}}{M}$ can naturally be of order 10^{-15}. This is the model which the Hypercolor Scenario would like to make use and generate the scale of the weak interactions which is of order ~ 100 GeV. It is in this spirit that Weinberg and Susskind [5] postulated the existence of a new strong interaction with a chiral symmetry breaking scale Λ_{QHD} $(G_F^{1/2})^{-1} \sim 250$ GeV.

The idea of describing weak interactions in the spirit of a QCD-like theory, with no fundamental Higgs fields and hence no unknown Yukawa couplings, is certainly a very attractive one. However, it is clear that the Weinberg-Susskind model [5] is only the first step in realizing this goal. First of all there is the all important question of generating fermion masses. Then, one would like to unify the hypercolor gauge interaction together with the usual $SU(3)_c \otimes SU(2)_L \otimes U(1)$ interactions at some finite scale. This circumstance by itself suggests that the HC group is much larger than the minimal Weinberg-Susskind model. To see this let us assume that there is a unifying scale m at which QCD and QHD coupling constants are equal

$$\bar{g}_H(m) = \bar{g}_c(m) \qquad (1.13)$$

This equality could also be expressed as

$$(11 N_{c'} - 2 N_{f'}) \ln\left(\frac{m}{\Lambda_{HC}}\right) = (11 N_c - 2 N_f) \ln\left(\frac{m}{\Lambda_c}\right) \qquad (1.14)$$

where $N_{c'}$ and $N_{f'}$ are, respectively, the number of hypercolors and hyperflavors. Assuming $N_c = 3$, $N_F = 6$ for QCD, and noting that $\Lambda_{HC} \gg \Lambda_c$, Eq. (1.14) can be satisfied only if

$$11 N_{c'} - 2 N_{f'} > 21 \qquad (1.15)$$

So, $N_c' \geq 3$ if $N_f' \leq 4$ etc. Thus, if QCD and QHD are to be unified at some finite scale, then the hypercolor group is certainly bigger than the minimal Weinberg-Susskind model.

In this talk I shall describe the implications of an extended Weinberg-Susskind model for high energy experiments. Though, for purely pedagogical reasons I will stick to a particular extension due to Farhi and Susskind and Dimopolous [6], the existence of pseudo Goldstone bosons, PGB's, which are Goldstone bosons of the HC scenario and get masses through the various interactions, is a matter of general consequence. These PGB's are inescapable in any realistic DSB realization, being an artifact of the large underlying chiral symmetry of the HC group.

As we shall see later, there are always some composite particles in HC models in general, which are parametrically light on the scale of the HC interactions, characterized by $\Lambda_{HC} \sim (G_F^{-1})^{1/2} \sim 250$ GeV. Some of them receive masses only through higher order electroweak interactions and are expected to be lighter than the others [7,8]. They are to the HC interactions what π, K, η are to the color interactions, QCD. Though the exact mass of these PGB's is a matter of controversy, most estimates [7,9,10,11] put them in the range of 2-40 GeV. Thus when produced in reactions involving $Q^2 \ll \Lambda^2_{HC}$ they would appear point-like particles -- very much like the pointlike Higgs of the standard Weinberg-Salam theory.

The discovery of a pointlike scalar particle in an experiment, though in its own right a matter of tremendous joy for our experimental colleagues, has to be carefully analysed to check whether it is a Higgs or a pseudo Goldstone boson. It is gratifying to note that the <u>scalar 0^+ nature of the fundamental Higgs vs. the pseudoscalar 0^- nature of the PGB's</u>, which is a feature common to all <u>natural</u> realizations of Dynamical Symmetry breaking scenario, <u>introduces a fundamental difference in a large class of production processes</u> which are of interest in the next generation of e^+e^- experiments [12,13].

The topics which I would like to discuss are as follows:

(i) spectroscopy [6] and estimates of the masses of the HC 0^- bosons [7,9,10,11].
(ii) couplings of the PGB's to gauge bosons and the production of neutral PGB's vis à vis the standard Higgs [12,13,15].
(iii) couplings of the PGB's to ordinary fermions [7,8,12,13] and their phenomenological implications in the decays of the K,B,D and T mesons and in top onia, J_T [12,13,16,17].
(iv) the problem of the Flavor Changing Neutral Current, FCNC, transitions in Extended Hypercolor theories [8,14] and
(v) the production and decays of the charged PGB's in e^+e^- annihilation [7,12,13,17 and the Z decays [15] including the signatures that experimentalists at PETRA/PEP and LEP should be made aware of.

Some of these points have been discussed in review articles [18]. I have left out production processes other than e^+e^- annihilation since they have only marginal chance of becoming experimentally relevant. High energy e^+e^- annihilation is expected to be the main theatre where the activity concerning the nature of spontaneous symmetry breaking would be concentrated.

II. Spectroscopy and Mass Estimates of Pseudoscalar Bosons in the HC Scenario

In writing a realistic model for the Hypercolor interactions, it is necessary to take care of the breaking of the electroweak $SU(2)_L \otimes U(1)$ symmetry down to the U(1) symmetry, since phenomenology demands that the relations $m_W = m_Z \cos\theta$ and $m_\gamma = 0$ must emerge; aesthetics demands that they should emerge naturally. Thus one possible scheme is to follow the Weinberg-Susskind prescription [5] by setting the HC interactions to be <u>vectorlike</u> and assign to a pair of Hyperfermions the $SU(2)_L \otimes U(1)$ quantum numbers of a quark or a lepton. In this way the chiral symmetry breaking of the HC interaction will induce the correct pattern of $SU(2) \otimes U(1)$ breaking. Though to start with one could remove the requirement of vectorlike HC interactions and argue that the chiral symmetry breaking is a dynamical phenomenon and its realization a priori not obvious, we find it hard to understand that the weak $\Delta I = 1/2$ rule would emerge naturally if the HC interaction were <u>not</u> vectorlike. So, whereas more general scenario exist [19], we would not pursue them here.

The simplest extension of the one-doublet Weinberg-Susskind model is the one-family model of ref. (6). Following the familiar pattern of the ordinary fermions, it is assumed in this model that the Hyperfermions also come in the form of a standard family. This would guarantee the cancellation of the $SU(3)_{Colour} \otimes SU(2)_{Left} \otimes U(1)_{HYP}$ anomalies. Thus one has four weak doublets:

$$\text{Hyper quarks:} \quad \begin{pmatrix} U_i^\alpha \\ D_i^\alpha \end{pmatrix}_L , \quad (U_i^\alpha)_R , \quad (D_i^\alpha)_R$$

$$\text{Hyper leptons:} \quad \begin{pmatrix} N^\alpha \\ E^\alpha \end{pmatrix}_L , \quad (N^\alpha)_R , \quad (E^\alpha)_R \tag{2.1}$$

where $i = 1,2,3$ is the usual color index and α is a Hypercolor index. Notice that there is a right handed hyperneutrino whose presence is necessary to generate the weak $\Delta I = 1/2$ rule. The weak interaction properties of this model are the ones that would obtain for four-doublet models, where the doublets are distinguished by the usual color labels $i = 1,2,3$, with the fourth label being the lepton number.

Let us look at the spectroscopy of the one-family HC model [6]. Since there are eight species of Hyperfermions then in the absence of any bare mass terms and interactions the global HC chiral symmetry could be realized in one of the two forms:

case (i) : $G = SU(8)_L \otimes SU(8)_R \otimes U(1)$ if HC fermions belong to a complex representation of the HC group.

case (ii): $G = SU(16)$ if the HC fermions belong to a real representation of the HC group.

When the HC interactions become strong the symmetry group G is broken spontaneously by the condensates $\langle \bar{U}_i U_i \rangle \neq 0$, $\langle \bar{D}_i D_i \rangle \neq 0$ etc. However, there is still a symmetry group, H, that survives the spontaneous symmetry breaking. This could have one of the following forms:

case (i) $H = SU(8) \otimes U(1)$

case (ii) $H = O(16)$ if the symmetric product of the fermion representation is invariant

or $H = SP(16)$ if the antisymmetric product of the fermion representation is invariant.

The case when SU(16) symmetry is broken spontaneously to SP(16) has inherent phenomenological troubles as studied by Peskin [9] and we shall neglect it. Since the purpose of this talk is illustrative, I shall concentrate only on case (i) which is realized through an SU(8)-symmetric condensate

$$\langle \bar{U}_{iR} U_{iL} \rangle_o = \langle \bar{D}_{iR} D_{iL} \rangle_o = \langle \bar{E}_R E_L \rangle_o = \langle \bar{N}_R N_L \rangle_o = \mu^3 \neq 0 \qquad i = 1,2,3 \tag{2.2}$$

For this case the classification group $H = SU(8)_{L+R}$ and there are $8^2 - 1 = 63$ Goldstone bosons, which are exactly massless at this stage. These Goldstone bosons, which are also present in the O(16) realization, can all be characterized through the transformation properties

$$P^{\delta, r} \equiv i \bar{F} \wedge \tau^r F \qquad \delta = 0,1,\dots;15 \quad r = 0,1,2,3 \quad (\delta, r) \neq (0,0) \tag{2.3}$$

where the HC (weak) fermion multiplet F has the form

$$F = \begin{pmatrix} U_1 & U_2 & U_3 & N \\ D_1 & D_2 & D_3 & E \end{pmatrix} \tag{2.4}$$

The HC scenario is complete when the electroweak and strong interactions are switched on. The residual $SU(8)_{L+R}$ symmetry is broken by the $SU(3)_c \otimes SU(2)_L \otimes U(1)$ forces plus any $F \rightarrow f$ forces which, for example, are present in the Extended Hypercolor, EHC models [8]. The Goldstone bosons get their masses through these interactions. Depending on the transformation properties under $SU(3)_c$ there are three distinct classes of pseudoscalar HC bosons in the one-family model [6].

Color-singlets

There are seven of them in all. The combination

$$\begin{gathered} \bar{U}\gamma_5 U - \bar{D}\gamma_5 D + \bar{N}\gamma_5 N - \bar{E}\gamma_5 E \\ \bar{U}\gamma_5 D + \bar{N}\gamma_5 E \\ \bar{D}\gamma_5 U + \bar{E}\gamma_5 N \end{gathered} \tag{2.5}$$

are eaten by the Z, W^+ and W^-, respectively, which in turn become massive. The weak gauge bosons $W^\pm$ and Z in this model are hybrid objects with the longitudinal degrees of freedom being the (HC)-composite Goldstone bosons. Thus, the pair production of longitudinally polarized $W^\pm$ bosons is expected to show a resonating structure in contrast to the pair production of transversely polarized $W^\pm$ bosons, which, as in the Weinberg-Salam model, would have a point-like cross-section. This is true of all models in which the $W^\pm$ and Z bosons exist as fundamental entities but receive masses through a dynamical symmetry breaking mechanism; it is to be contrasted with models in which the $W^\pm$ and Z are purely composite, where both the transverse and the longitudinal cross-sections would show resonating structures.

In addition to the fictitious bosons (2.5) there are four genuine 0^- color-singlet bosons, which we call π''s. In the one-family SU(8) model they have the following representations

$$\begin{aligned} \pi'^+ &= \bar{U}\gamma_5 D - \bar{N}\gamma_5 E \\ \pi'^- &= (\pi'^+)^\dagger \\ \pi'^0 &= \bar{U}\gamma_5 U - \bar{D}\gamma_5 D + \bar{E}\gamma_5 E - \bar{N}\gamma_5 N \\ \pi'^0_\eta &= \bar{U}\gamma_5 U + \bar{D}\gamma_5 D - 3(\bar{N}\gamma_5 N + \bar{E}\gamma_5 E) \end{aligned} \tag{2.6}$$

These Goldstone bosons remain massless in the absence of interactions but become massive through the electroweak and any $F \rightarrow f$ EHC forces. We shall call them pseudo Goldstone bosons, PGB's, to distinguish them from other HC 0^- bosons.

The second neutral π'^{0} is being called π'^{0}_{η} due to its (HC) 0^{-} isoscalar nature.

Color-Octets

There are thirtytwo of such 0^{-} objects in the one-family model under consideration [6], having the composition

	charge	
$\bar{U}\gamma_5 \frac{\lambda^a}{2} U \pm \bar{D}\gamma_5 \frac{\lambda^a}{2} D$	0	
$\bar{U}\gamma_5 \frac{\lambda^a}{2} D$	+1	(2.7)
$\bar{D}\gamma_5 \frac{\lambda^a}{2} U$	-1	

$a = 1,\ldots 8$

Color-Triplets

There are twentyfour of such objects, transforming like a triplet under $SU(3)_c$. They can also be termed as the HC lepto-quarks and have the composition

	charge	
$\bar{E}\gamma_5 U_i$	5/3	
$\frac{1}{\sqrt{2}}\left(\bar{E}\gamma_5 D_i - \bar{N}\gamma_5 U_i\right)$	2/3	(2.8)
$\bar{N}\gamma_5 D_i$	- 1/3	
$\frac{1}{\sqrt{2}}\left(\bar{N}\gamma_5 U_i + \bar{E}\gamma_5 D_i\right)$	2/3	

with i = 1,2,3, and the conjugates are not shown here separately. Both the color-octets and the color-triplets will be produced in the high energy scattering and decay processes in the sense of ordinary gluons and quarks, i.e. jets instead

of sharply defined mass states.

Mass Estimates

The starting point of all such mass estimates is the Dashen's formula [20]

$$
\begin{aligned}
M^2_{ab} &\equiv \langle \pi'^a | \delta\mathcal{L} | \pi'^b \rangle \\
&= \frac{1}{F^2_{\pi'}} \langle 0 | [Q^a_5 , [Q^b_5 , \delta\mathcal{L}]] | 0 \rangle
\end{aligned} \tag{2.9}
$$

$$\delta\mathcal{L} = \delta\mathcal{L}^{EW} + \delta\mathcal{L}^{QCD} + ----$$

where --- indicates any possible EHC contribution arising from the exchange of the massive E bosons [6], inducing an $F \rightarrow f$ transition.

$$\delta\mathcal{L}^{EW} = \delta\mathcal{L}^{em} + \delta\mathcal{L}^{Weak}$$

$$\delta\mathcal{L}^{em} \propto \int d^4q/q^4 \text{ F.T. } T\left(J^{em}_{\mu}(x)\, J^{em}_{\mu}(0)\right) \tag{2.10}$$

and so on. (F.T. stands for Fourier transform). Perhaps it should be remarked that the Dashen's formula (2.9) with $\delta\mathcal{L} = \delta\mathcal{L}^{em}$ leads to the successful electromagnetic mass difference for the ordinary pions [21].

The diagrams that could give masses to the π''s are shown in Fig. (1). It turns out that all the three neutral Goldstone bosons (π'^{0}, π'^{0}_{η} and the one eaten by the Z) remain <u>massless</u> to all orders in the $SU(2)_L \otimes U(1)$ interactions ($SU(3)_c$ does not contribute since these bosons are color-singlets). The reason of this could be traced to the circumstance that these neutral bosons are coupled to the conserved currents. Thus,

$$\left[J^{k}_{\pi'},\ \delta\mathcal{L}^{EW} \right] = 0 \qquad \text{for } k = 1,2,3 \tag{2.11}$$

leading to

$$m^{k} = 0 \tag{2.12}$$

for the color-singlet neutral 0^- bosons.

The mass of the charged PGB's, $m^{EW}_{\pi'^{\pm}}$, estimated from Fig. (1) has the value

$$\left(m^{EW}_{\pi'^{\pm}}\right)^2 \approx (3\alpha/4\pi)\, m_Z^2 \ln\left(m^2_{HC}/m^2_Z\right) \tag{2.13}$$

where m_Z is the mass of the Z boson and m_{HC} is the constituent Hyperquark mass, probably $\sim$ 1 TeV. Thus

$$\left(m^{EW}_{\pi'^{\pm}}\right)^2 \simeq (5 - 10\ \text{GeV})^2 \tag{2.14}$$

If the estimates (2.12) and (2.13) were the only contributions to the masses of the GB's and PGB's, then this particular HC alternative would have been dead already at present energies! This is so because there is no evidence for a massless, spinless chargeless particle anywhere; neither is there any indication of a charged PGB pair production at PETRA, which is good enough to detect $\pi'^{+}\pi'^{-}$ with the mass in the range (2.14).

The PGB's $\pi'^{\pm}$, π'^{0}, π'^{0}_{η} receive additional masses from the four-fermion interactions amongst the Hyperfermions. These four-fermi interactions are generated in an effective Lagrangian sense through the exchange of Extended Hyperbosons, E, between a Hyperfermion and an ordinary fermion as shown in Fig. (2). Apart from giving masses to the ordinary fermions this interaction

also breaks chiral symmetries thus giving masses to the otherwise exact massless Goldstone bosons, π'^{0}, π'^{0}_{η} [6]. The estimate of the EHC contribution to the PGB masses is model dependent and not very reliable. Farhi and Susskind [11] estimate that the EHC forces contribute very substantially to the PGB masses, giving

$$m_{\pi'^0} \simeq m_{\pi'^0_\eta} \simeq R \times 60\ (\text{GeV}) \tag{2.15}$$

where R is a model dependent factor of order (1). Other authors estimate much smaller contributions due to the EHC forces [10]; the difference essentially lying in the determination of the EHC mass scale. Thus, there is considerable uncertainty in the mass estimates of the PGB's. As a ball park estimate we shall take

$$m_{\pi'} \simeq (10 - 40)\ \text{GeV} \tag{2.16}$$

though neutral PGB's may be much lighter.

The contribution to the color-octet 0^- bosons comes from a single gluon exchange. This can be compared to the single photon exchange which gives the mass difference $m^2_{\pi^+} - m^2_{\pi^0}$ [21]. Following Farhi and Susskind [11], one has (O is for color octets)

$$\frac{m^2_{O}}{m^2_{\pi^+} - m^2_{\pi^0}} \simeq \frac{\Lambda^2_{HC}}{\Lambda^2_{QCD}} \cdot \frac{\alpha_{strong}(\Lambda_{HC})}{\alpha_{em}} \cdot \frac{3}{1} \tag{2.17}$$

Here 3 is the color Casimir for the octet. The factor $\left(\Lambda_{HC}/\Lambda_{QCD}\right)^2$ which characterizes the ratio of the bound state structure due to QHD and QCD can be estimated using the scaling argument [11]. If the Hypercolor group is $SU(N_{c'})$ then

$$\sqrt{\frac{N_{c'}}{3}}\ \frac{\Lambda_{HC}}{\Lambda_{QCD}} \sim \frac{F_{\pi'}}{f_{\pi}}$$

and

$$F_{\pi'} = \frac{2m_W}{\sqrt{n_{f'}}\, g} \simeq \frac{246}{\sqrt{n_{f'}}} \text{ GeV} \tag{2.18}$$

where $n_{f'}$ is the number of HC doublets (= 4 in the one-family model). Using $\alpha_{strong}(\Lambda_{HC}) \simeq 0.1$ gives

$$m_\theta \simeq \sqrt{\frac{4}{N_c'}} \times 260 \text{ GeV} \tag{2.19}$$

The estimate (2.19) is a lower limit, since like the color singlets π''s, there will be additional contributions to the mass of the octets from the EHC forces but they should not substantially change the value obtained in (2.19).

The masses for the color-triplet bosons, called T_c and T_c^i, are likewise estimated by a single gluon exchange and can be obtained by replacing the color Casimir, 3 for the octets, by the corresponding factor 4/3 for the triplet. This gives

$$m_{T_c}/m_\theta \simeq \left(\frac{4/3}{3}\right)^{1/2} = \frac{2}{3}$$

yielding a mass

$$m_{T_c} \simeq \sqrt{\frac{4}{N_c'}} \times 170 \text{ GeV} \tag{2.20}$$

Thus, the three classes of the HC 0^- bosons are fairly split in masses as indicated through the estimates (2.16), (2.19) and (2.20), with (2.16) the lowest and uncertain. However, since the mass of the charged PGB's $\pi'^{\pm}$ arises

at the one loop level, it should be of the order of magnitude

$$m_{\pi'}^2 \sim \alpha F_{\pi'}^2 \Longrightarrow m_{\pi'} < 20 \text{ GeV} \tag{2.21}$$

In estimating these masses, I have neglected the contribution from the chiral perturbation theory self interactions. These have been very carefully estimated by Chada and Peskin [9], but considering the uncertainty in the value of $m_{\pi'}$ and the model dependence of the chiral perturbation estimates, the numbers quoted above are still correct as to the order of magnitude.

The picture that emerges from these estimates is that the typical scale of the Hypercolor interaction, characterized by Λ_{HC} is of order $\sim$ 1 TeV. One expects Hyperbaryons in that region (fermions or bosons depending on N_c, being odd or even) which may or may not be stable and a whole new attendant spectroscopy. This, however, will be accessible only at Trans-TeV accelerators. Somewhat below lie the color octets (around 270 GeV) and the color triplets (around 170 GeV) which may be detected using Isabelle and SUPERLEP. However, the lowest lying of these bosons, the PGB's π''s may act as the outposts of the new hadronic land and may be accessible at the maximum PETRA/PEP energies and almost certainly at LEP energies. In the rest of this talk I shall only discuss the production and decay mechanisms of the π''s for obvious reasons.

III. Couplings of PGB's to Gauge Bosons and their Implications

We turn first to the couplings of the PGB's among themselves and with the gauge bosons. These couplings can be induced using an effective Lagrangian approach, as advocated by Bèg, Politzer and Ramond [7]. This effective Lagrangian, which is useful so long as π''s are soft on the scale of $F_{\pi'}$, in general must contain all terms describing (the global chiral) group invariant interactions, each appearing with an arbitrary dimensionful coefficient. Since the field is dimensionless, a term in the Lagrangian with n derivatives has a coefficient of dimension $(\text{mass})^{4-n}$. The scales of these factors are set by the natural scale of the theory, Λ_{HC}. For energies $E \ll \Lambda_{HC}$ each extra derivative gives a suppression factor E/Λ_{HC}. Let us analyse first the effective Lagrangian

$$\mathcal{L} = \frac{1}{2} F_{\pi'}^2 \left(D_\mu \pi' \right)^2 \tag{3.1}$$

where D_μ is the appropriate gauge-covariant derivative. The relevant vertices and Feynman rules based on (3.1) are shown in Fig. (3), which involve the couplings of a pair of PGB's with either a single boson or a pair of electroweak gauge bosons (γ, $W^\pm$ and Z^0); the couplings of π''s to the gluons is zero because of the color singlet nature of the π''s.

There are two immediate consequences of the couplings shown in Fig. (3), namely they lead to the following rates for the pair production of PGB's in e^+e^- annihilation [7,15]:

$$\begin{aligned} \Delta R(\pi'^+\pi'^-) &= \sigma(e^+e^- \xrightarrow{\gamma} \pi'^+\pi'^-)/\sigma(e^+e^- \to \mu^+\mu^-) \\ &= 1/4\,(1 - 4m_{\pi'^\pm}^2/s)^{3/2} \\ \sigma(e^+e^- \xrightarrow{\gamma} 2\pi'^0,\ 2\pi'^0_\eta,\ \pi'^0\pi'^0_\eta) &= 0 \end{aligned} \tag{3.2}$$

Likewise, the Z decays would also lead to pair production of PGB's with the general result

$$\Gamma(Z^0 \to \pi'\pi') = \frac{\alpha m_Z}{12(\sin 2\theta_W)} \left(I_3 - 2\sin^2\theta_W Q \right)^2 \times \left(1 - 4m_{\pi'}^2/m_Z^2\right)^{3/2} \tag{3.3}$$

where I_3 and Q are, respectively the third component of isospin and charge of the PGB's, π' . Thus, one would have [15]

$$\begin{aligned} \Gamma(Z \to \pi'^+\pi'^-)/\Gamma(Z^0 \to \mu^+\mu^-) &= \frac{1}{2}(1 - 2\sin^2\theta_W)^2 (1 - 4m_{\pi'}^2/m_Z^2)^{3/2} / (1 - 4\sin^2\theta_W + 8\sin^4\theta_W) \\ &\simeq 0.12\,(1 - 4m_{\pi'}^2/m_Z^2)^{3/2} \\ \Gamma(Z \to \pi'^0\pi'^0) = \Gamma(Z \to \pi'^0_\eta \pi'^0_\eta) &= \Gamma(Z \to \pi'^0\pi'^0_\eta) = 0 \end{aligned} \tag{3.4}$$

The equations in the second line in (3.4) follow from $Q(\pi'^0, \pi'^0_\eta) = I_3(\pi'^0, \pi'^0_\eta) = 0$ and the numerical estimate in the first of Eq. (3.4) is obtained for the value $\sin^2\theta_W = 0.23$. Thus it is clear that given the phase space, the production of a pair of $\pi'^+\pi'^-$ in both the e^+e^- annihilation and the Z decays is not expected to be small.

Next, we turn to the couplings of a single PGB to a pair of electroweak gauge bosons. Note that being higher order they are not contained in (3.1). In gauge theories the vertex functions $\langle VVA \rangle_0$ and $\langle AAA \rangle_0$ are controlled by the Adler-Bell-Jackiw anomalies at momenta less than the typical scale of the interaction. Thus, the decay rate for $\pi^0 \to 2\gamma$, calculated via the anomaly, agrees with the experiment to within 10 %. Even the decay rates for $\eta \to 2\gamma$ and $\eta' \to 2\gamma$ calculated using the anomaly, which involve momenta $\approx O(m_\eta, m_{\eta'}) > \Lambda_{QCD} \simeq O(100\text{ MeV})$, are in agreement with the experiments within a factor 2. Since the scale of the HC interactions, $\Lambda_{HC} \simeq 1$ TeV and we have estimated the masses of the PGB's to be of order (10-50 GeV), their couplings to a set of electroweak gauge bosons can be reliably calculated using the triangle diagram.

The transitions of interest are

$$V^i(k_1) \longrightarrow V^j(k_2) + \pi'^a(q) \tag{3.5}$$

where $i,j = \gamma, Z^0, W^\pm$ and the hyperpion index 'a' takes the values relevant for the processes involving π'^+, π'^-, π'^0 and π'^0_η. Recalling that the π''s are pseudoscalar one could express the invariant amplitudes for (3.5) as

$$\mathcal{M}\left(V^i(k_1) \rightarrow V^j(k_2) + \pi'^a(q)\right)$$
$$= f\left(k_1^2/M^2, k_2^2/M^2, q^2/M^2\right)^{aij} \epsilon_{\mu\nu\lambda\sigma}\, k_1^\mu k_2^\nu \epsilon_1^\lambda \epsilon_2^\sigma$$

where M is the HC constituent mass $\sim$ O(1 TeV). The transition form factors $f(\frac{k_1^2}{M^2}, \frac{k_2^2}{M^2}, \frac{q^2}{M^2})$ are evaluated using the triangle diagram shown in Fig. (4).

For PGB's and available Q^2 in the foreseeable future one could approximate

$$f\left(k_1^2/M^2, k_2^2/M^2, q^2/M^2\right)^{aij} \simeq f(0,0,0)^{aij}$$

where

$$f(0,0,0)^{aij} = 1/(4\pi^2 F_{\pi'}) \cdot n_{f'}\, \mathrm{Tr}\left[\pi'^a \{\Gamma^i, \Gamma^j\}\right] \tag{3.6}$$

The matrices Γ^i, Γ^j for the processes (3.5) are given by the following expressions (assuming the usual SU(2) $\otimes$ U(1) properties for the hyperfermions).

$$\Gamma_V^{W^\pm} = g\left(\frac{\tau^1 \pm i\tau^2}{2}\right) = -\Gamma_A^{W^\pm} \tag{3.7}$$
$$\Gamma_V^Z = g/\cos\theta_W\, \left(\tau^3/4 - \sin^2\theta_W Q\right)$$
$$\Gamma_A^Z = -g/\cos\theta_W\, \left(\tau^3/4\right)$$
$$\Gamma_V^\gamma = eQ$$
$$\Gamma_A^\gamma = 0$$

The transition charges $f(0,0,0)^{aij}$ can then be calculated by specifying the HC model and the representation for the PGB fields π''s. In general note that the triangle diagram in (4) depends both on the number of Hypercolors and Hyperflavors. For example the transition charge $f_{ZZ\pi'^0}$ is given by

$$f_{ZZ\pi'^0} = -\, 2\alpha/(3\pi F_{\pi'})\, N_c' n_f' \times \sin^2\theta_W (1 - 2\sin^2\theta_W)/(\sin 2\theta_W)^2 \tag{3.8}$$

This expression could be used to set the order of magnitude for the transition charges relevant for the processes (3.5). It is instructive to compare the coupling, Eq. (3.8), with the Higgs coupling in the standard model [1] which is given by

$$g_{ZZ\varphi^0} = 2\left(\sqrt{2}\, G_F\right)^{1/2} m_Z^2 \tag{3.9}$$

This shows that the PGB coupling, (3.8), is several orders of magnitude smaller than the corresponding Higgs coupling. Perhaps it is worth reiterating the assumptions that have been made in the derivation of the π' couplings (3.6). We have made use of the following two inputs:

(i) the 0^- nature of the PGB's (a feature common to all natural DSB realizations) and

(ii) the hypercolor PCAC relation

$$q^\mu J_\mu^5 \sim 1/\sqrt{2}\; F_{\pi'}^2\, q^2 \pi' \tag{3.10}$$

which should be a better approximation for π''s in QHD, than the ordinary PCAC relation for the pions in QCD.

The vertices (3.5) appear in a large class of high energy reactions. In particular, they would be responsible for the processes

$$Z \longrightarrow \pi'^0 + (\ell^+\ell^-) \qquad (\ell = e, \mu, \tau) \tag{3.11}$$

$$Z \longrightarrow \pi'^0 + \gamma \tag{3.12}$$

$$e^+e^- \longrightarrow \pi'^0 + Z \tag{3.13}$$

as well as the reaction

$$e^+e^- \longrightarrow \pi'^0 + \gamma \tag{3.14}$$

The processes (3.11) - (3.14) are to be compared with the corresponding processes in the standard Weinberg-Salam model, where π'^0 is replaced by the Higgs scalar, φ^0. Recall that the processes analogous to (3.11) and (3.13) involving a Higgs namely [22]

$$Z \longrightarrow \varphi^0 + \ell^+\ell^- \tag{3.11'}$$

and the production process [23]

$$e^+e^- \longrightarrow \varphi^0 + Z \tag{3.13'}$$

are tree level processes in the standard model [1] and hence à priori expected to be much bigger than the corresponding one-loop (QHD) processes (3.11) and (3.13). The reactions analogous to (3.12) and (3.14) involving a Higgs, namely the radiative decay of Z [24]

$$Z \longrightarrow \varphi^0 + \gamma \tag{3.12'}$$

and [25]

$$e^+e^- \longrightarrow \varphi^0 + \gamma \tag{3.14'}$$

are on the other hand a one-loop (QFD) processes. Hence à priori we expect them to be comparable to the corresponding PGB processes, though the details would depend upon the transformation properties of the hyperquarks under the electroweak group, $G_W = SU(2) \otimes U(1)_{Hyp}$. The diagrams that contribute to the PGB processes (3.11) - (3.14) and the corresponding processes (3.11') - (3.14') involving a Higgs are shown in Fig. (5).

The ratio $\Gamma(Z \to \pi'^0\mu^+\mu^-)/\Gamma(Z \to \varphi^0\mu^+\mu^-)$ is shown in Fig. (6) for equal values of π'^0 and φ^0 mass. Note that in the entire range accessible in the Z decay, one has

$$\Gamma(Z \to (\pi'^0, \pi'^0_\eta) + \mu^+\mu^-)/\Gamma(Z \to \varphi^0 + \mu^+\mu^-) \simeq 10^{-7}\, n_{f'}^3\, N_{c'}^2 \leqslant 10^{-4} \tag{3.15}$$

Thus unless one is willing to consider absurdly large values of $N_{c'}$ and $n_{f'}$, a <u>positive signal in $Z \to \mu^+\mu^-$ + (spin 0) would be indicative of a Higgs φ^0.</u> In addition to the rates, one could check the dimuon invariant mass distribution to see which particle is being produced; the two distributions are shown in Fig. (7) and are quite different due to the contribution of the photon propagator in the PGB process (3.11) (see Fig. 5), which is absent in the tree approximation for the Higgs process (3.11').

The ratio $\sigma(e^+e^- \to \pi'^0 Z)/\sigma(e^+e^- \to \varphi^0 Z)$ is likewise very small. It is shown in Fig. (8) as a function of $m_{\pi'^0} (= m_{\varphi^0})$ for some representative values of the LEP centre of mass energies. The order of magnitude estimate is [12,13]

$$\sigma\left(e^+e^- \to (\pi'^0, \pi'^0_\eta) + Z\right) / \sigma\left(e^+e^- \to \varphi^0 + Z\right) < 10^{-8}\, n_{f'}^3 N_{c'}^2 \leqslant 10^{-5} \tag{3.16}$$

Again the production of <u>a spin 0 object in e^+e^- annihilation recoiling against the Z boson would be a clear vote for the Higgs mechanism as opposed to the Dynamical Symmetry breaking scenario.</u>

The processes (3.12) and (3.14) are likewise calculated using the amplitudes calculated through Eqs. (3.6) and (3.7). The corresponding processes involving a Higgs are calculated using the (heavy) fermion and $W^\pm$ boson loops, as shown in Fig. (5) and we quote the results [12]

$$\begin{aligned} \Gamma(Z \to \pi'^0 + \gamma) / \Gamma(Z \to \varphi^0 + \gamma) &\simeq 10^{-4} N_{c'}^2 n_{f'}^3 \quad \text{for } \sin^2\theta_W = 0.2 \\ &= 0 \quad \text{for } \sin^2\theta_W = 0.25 \end{aligned} \tag{3.17}$$

$$\sigma\left(e^+e^- \to \pi'^0 + \gamma\right) / \sigma\left(e^+e^- \to \varphi^0 + \gamma\right) \simeq \frac{1}{36} N_{c'}^2 n_{f'}^3 \tag{3.18}$$

The smallness of the ratio in (3.17) is due to the fact that in the process (3.12) only the vector current contributes, which in the model under consideration gives a contribution proportional to $(1-4\sin^2\theta_W)^2$. Among the reactions considered so far only the process $e^+e^- \to \pi'^0 + \gamma$ could be potentially large (as compared to $e^+e^- \to \varphi^0 + \gamma$) but considering the estimates $\Delta R(\varphi^0\gamma) \leqslant 10^{-6}$ and (3.18), it would be hard to detect it.

In principle, if $m_Z > m_{W^\mp} + m_{\pi'^\pm}$, then a charged PGB could be produced in the decay of Z via

$$Z \to W^\pm + \pi'^\mp \tag{3.19}$$

The decay width for (3.19) can be estimated using the transition charge

$$\left(f_{ZW^\mp\pi'^\pm}\right) \simeq -e^2 / (24\pi^2 F_{\pi'} \cos\theta_W) \times N_{c'} n_{f'} \tag{3.20}$$

which gives

$$\begin{aligned} \Gamma(Z \to W^\pm \pi'^\mp) / \Gamma(Z \to \mu^+\mu^-) \\ \simeq \alpha^2 N_{c'}^2 n_{f'}^3 K / (72\,\pi^2 \cos^4\theta_W) \end{aligned} \tag{3.21}$$

where the factor K is given by

$$\begin{aligned} K &= 1/m_Z^6 \left[\left\{ (m_Z^2 + m_W^2 - m_{\pi'}^2)^2 - 4 m_Z^2 m_W^2 \right\} \times \left\{ (m_Z^2 + m_{\pi'}^2 - m_W^2)^2 - 4 m_Z^2 m_{\pi'}^2 \right\}^{1/2} \right] \\ &\xrightarrow[m_{\pi'} \approx 0]{} (1 - m_W^2/m_Z^2)^3 \end{aligned} \tag{3.22}$$

which for $n_{f'} = N_{c'} = 4$ gives

$$\Gamma(Z \to W^\pm \pi'^\mp) / \Gamma(Z \to \mu^+\mu^-) \leqslant 10^{-6} \tag{3.23}$$

and it looks hopeless!

The process

$$\begin{aligned} Z \to W^\pm_{\text{vir.}} &+ \pi'^\mp \\ \hookrightarrow (\ell^\pm \nu_\ell) \end{aligned} \tag{3.24}$$

has a branching ratio similar to the process $Z \to \pi'^0 + \mu^+\mu^-$ and we expect

$$BR\left(Z \to \ell^\pm \nu_\ell + \pi'^\mp\right) \leqslant 10^{-8} \tag{3.25}$$

The decays $Z \to W^{\pm} + \varphi^{\mp}$ and $Z \to \underset{\longrightarrow}{W^{\pm}_{\nu u}} + \underset{(\ell^{\pm}\nu_\ell)}{\varphi^{\mp}}$ in an extended Higgs Weinberg-Salam model are also one-loop (QFD) processes like the radiative decay $Z \to \varphi^0 + \gamma$. Hence we expect that the branching ratios are comparable to the analogous decays involving a $\pi'^{\pm}$. Thus, the PGB processes (3.19) and (3.23) and the corresponding ones involving a charged Higgs are not of any great experimental interest.

The multiple production of PGB's in the decays of Z, e.g.,

$$Z \longrightarrow \pi'^0 \pi'^+ \pi'^- \tag{3.26}$$

has been studied in ref. (15). One expects

$$\Gamma(Z \to \pi'^0 \pi'^+ \pi'^-) / \Gamma(Z \to \mu^+\mu^-) \leqslant 10^{-4} \tag{3.27}$$

Considering that in the canonical framework $BR(Z \to \mu^+\mu^-) \simeq 3\ \%$ and the difficulty in reconstructing the manybody final states, the process (3.23) is at best of remote experimental interest.

IV. Couplings of PGB's to light fermions and their implications

In the Weinberg-Salam theory [(1)], the mechanism which breaks the chiral symmetry (i.e. gives masses to the fermions) also sets the scale for the Higgs-fermion-fermion couplings. In particular, one has the $(SU(2)_L \otimes U(1))$ invariant Yukawa couplings of the Higgs doublet with the fermions

$$\mathcal{L}'_{WS} = \bar{L} \phi h_f R + h.c. \tag{4.1}$$

where h_f is a coupling matrix. After spontaneous symmetry breakdown

$$\varphi \to \varphi + \langle \varphi^0 \rangle \Longrightarrow m_f = h_f \langle \varphi^0 \rangle_c \tag{4.2}$$

Thus, the Higgs-fermion-fermion coupling is given by

$$h_f = m_f / \langle \varphi^0 \rangle_0 \tag{4.3}$$

However, the fact that the fermion masses are so widely split or equivalently why $h_e/h_b = 10^{-4}$ is not explained. In theories with dynamical symmetry breaking, the mechanism which breaks the chiral symmetries so as to generate the right (order of magnitude) fermion masses is not understood. In such scalarless theories, Yukawa couplings like (4.1) are absent; moreover the bare mass terms for the ordinary fermions are not allowed due to non-invariance under the $SU(2)_L \otimes U(1)$ transformations. It is almost certain in these circumstances that the mechanism of generating fermion masses has to have a gauge origin.

In one such mechanism the ordinary fermions are put together with the hyperfermions in a single multiplet of a large symmetry group and this symmetry is then gauged. The resulting theory is called the Extended Hypercolor, EHC [(6)]. Thus, for example one has an EHC multiplet

$$\Psi_{EHC} \equiv \begin{pmatrix} F \\ F \\ F \\ \vdots \\ f \\ f \\ f \end{pmatrix} \tag{4.4}$$

and for the sake of definiteness let the EHC group be SU(N+3). At some scale Λ_E, the SU(N+3) chiral symmetry is broken spontaneously

$$SU(N+3) \supset SU(N)_{HC} \otimes SU(3)_c$$

still preserving a hypercolor SU(N) and the ordinary color SU(3) symmetry.

The gauge bosons mediating the $F \to f$ transitions become massive. These transitions break chiral symmetries of the ordinary fermions as well as of the Hyperfermions thus generating masses for them. The symmetry breaking chain could go through several stages involving different HC mass scales, providing a mechanism to understand the vast differences in the observed fermion masses.

To be somewhat more precise let us consider the EHC-boson-induced currents

$$J_\mu^E \sim \bar{F}\gamma_\mu f \qquad (4.5)$$

and the resulting effective Lagrangian

$$\mathcal{L}^E = 1/2\,(g_E/m_E)^2(\bar{F}\gamma_\mu f)(\bar{f}\gamma_\mu F) \qquad (4.6)$$

Fierz transformation leads to

$$\mathcal{L}^E = -1/2\mu^2\left[(\bar{F}F)(\bar{f}f) - (\bar{F}\gamma_5 f)(\bar{f}\gamma_5 F) + \cdots\right] \qquad (4.7)$$

where we have set $\mu^2 \equiv m_E^2/g_E^2$, and g_E is the EHC gauge coupling constant. Now at some scale, the Hypercolor interactions become strong and the HC chiral symmetry is broken spontaneously due to the non-zero v.e.v. of the condensate $\langle\bar{F}F\rangle_0$. This gives constituent masses to the Hyperquarks, of order $\Lambda_{HC} \simeq 1$ TeV, the constituent mass of the hyperquarks is fed down via the diagram shown in Fig. (2) to generate the ordinary quark (current) mass. Thus for ordinary fermions one has [6] (μ depends on the type of fermion)

$$m_f = \langle\bar{F}F\rangle_0/2\mu^2 \qquad (4.8)$$

This then also sets the scale for the pseudo Goldstone boson couplings to a

pair of fermions

$$g_{f\bar{f}'\pi'} \sim \left(m_f/\langle\bar{F}F\rangle_0\right)F_{\pi'}^2\,\Gamma_{ff'} \sim m_f/F_{\pi'}\,\Gamma_{ff'} \qquad (4.9)$$

where $\Gamma_{ff'}$ are possible mixing angles and we have set $\langle\bar{F}F\rangle_0 \simeq F_{\pi'}^3$. Now, if $\Gamma_{ff} \sim O(1)$, then the π'^0 $f\bar{f}$ couplings (4.9) are of the same order of magnitude as the $\varphi^0 f\bar{f}$ couplings (4.3) in the Weinberg-Salam model. Let us consider the implications of the $\pi' f\bar{f}'$ couplings, Eq. (4.9), for the K_L-K_S mass difference. The diagram shown in Fig.(9) involving the exchange of a single pseudo Goldstone boson π' generates an effective four-fermion $\Delta S = 2$ interaction given by

$$\begin{aligned}\mathcal{L}_{\pi'}(\Delta S=2) &= G_F^2 F_{\pi'}^2 (m_s+m_d)^2/2m_{\pi'}^2 \\ &\quad \times (\bar{d}\gamma_5 s\,\bar{s}\gamma_5 d) \\ &\simeq \frac{G_F}{2\sqrt{2}\,n_{f'}}\left((m_s+m_d)/m_{\pi'}\right)^2(\bar{d}\gamma_5 s\,\bar{s}\gamma_5 d)\end{aligned} \qquad (4.10)$$

which gives an effective Fermi type coupling with the ratio

$$G_{\pi'^0}/G_F \sim \frac{1}{2\sqrt{2}\,n_{f'}}\left((m_s+m_d)/m_{\pi'}\right)^2 \qquad (4.11)$$

which for $m_{\pi'^0} \simeq 10$ GeV gives $G_{\pi'^0}/G_F \simeq 10^{-3}$ --- in conflict with the limit from the observed K_L-K_S mass difference. The observed Δm_k implies an upper bound on $G_{\pi'^0}$:

$$\Delta m_K/m_K \simeq 0.7\times 10^{-14} \Longrightarrow G_{\pi'^0}/G_F < 10^{-5} \qquad (4.12)$$

Thus, if $m_{\pi'^0} \simeq O(10\ \text{GeV})$, there is a severe conflict between experiment and the single π'^0 exchange $\Delta S = 2$ amplitude. This problem can be cured if no FCNC transitions are allowed involving a single PGB exchange. John Ellis and co-workers [13] have invented such a cure which they call the principle of "monophagy". It is an extension of the Glashow-Weinberg criterion [26] for the case of a fundamental Higgs, which ensures natural diagonal Higgs-fermion-fermion couplings. In essence, "monophagy" states that fermions of a given charge get their mass from a single hyperfermion condensate. Thus, for example, only transitions of the following kinds are allowed

$$(u,c,\ \dots) \longleftrightarrow U \qquad (4.13)$$

$$(d,s,\ \dots) \longleftrightarrow D$$

but not the ones coupling (u,c,..) to both U and D.

It is probable that if $m_{\pi'^0} \simeq O(10\ \text{GeV})$, then a condition analogous to (4.13) will have to be satisfied by the EHC models.

Let us assume that the EHC/HC generators follow (4.13). Then the EHC currents are given by [6,13]

$$J_\mu^E = \sum_{f,n} \left(\bar{f}_{Ln}\, g^E_{Ln}(f) \gamma_\mu F_L(f) + \bar{f}_{Rn}\, g^E_{Rn}(f) \gamma_\mu F_R(f) \right) + h.c. \qquad (4.14)$$

which are coupled to the EHC/HC vector boson mass eigenstates. Here $n \equiv$ generation = 1,2,3 and $f \equiv$ type of fermions = (u,...), (d,...), (e,...), (ν,...) and the coefficients $g^E_{Ln}(f)$ are constrained by the $SU(2)_L \otimes U(1) \otimes SU(3)_c$ symmetry. The exchange of the EHC gauge bosons leads to the effective Lagrangian

$$\mathcal{L} = -\sum_E g_E^2/2\mu_E^2 \; J_\mu^E J_\mu^E$$

$$\begin{aligned} \sim -\sum_{f,f'} \sum_{n,m} \sum_E g_E^2/\mu_E^2 \times \Big\{ & \left(\bar{f}_{Ln}\, g^E_{Ln}(f)\gamma_\mu F_L\right)\left(\bar{F}'_R\, g^{E*}_{Rm}(f')\gamma_\mu f'_{Rm}\right) \\ & + \left(\bar{f}_{Ln}\, g^E_{Ln}(f)\gamma_\mu F_L\right)\left(F'_L\, g^{E*}_{Lm}(f')\gamma_\mu f'_{Lm}\right) + (L\leftrightarrow R) \\ & + h.c. + \cdots \Big\} \end{aligned} \qquad (4.15)$$

where the dots indicate possible contributions for the case when the fermion HC representation is not complex. The first term can be related, after Fierz-transformation and some manipulations to the fermion mass matrix. Thus, it contributes an interaction term of the form [13]

$$-2 \sum_{f,f'} \left(\bar{F}'_R F_L\right)\left(\sum_{n,m} \bar{f}_{Ln} \left(\Gamma^{ff'}_{LR}\right) f'_{Rm}\right) + h.c. \qquad (4.16)$$

where

$$\left(\Gamma^{ff'}_{LR}\right)_{nm} \equiv \sum_E g_E^2/\mu_E^2 \; g^E_{Ln}(f)\, g^{*E}_{Rm}(f') \qquad (4.17)$$

and one could show that

$$m_f = \langle \bar{F}F \rangle_0 \; \Gamma^{ff}_{LR}$$

Thus, in the basis in which the mass matrix is diagonal, so is the coupling matrix $\Gamma^{ff'}_{LR}$. Eq. (4.16) leads to the following interaction strengths

$$\mathcal{L}(\pi'^{0} f\bar{f}) \sim m_f/F_{\pi'} \quad \pi'^{0}(\bar{f}\gamma_5 f)$$

$$\mathcal{L}(\pi'^{+} f\bar{f}') \sim 1/F_{\pi'} \quad \pi'^{+}\bar{f}'\left[\mathcal{U}_{KM}\, m_f (1+\gamma_5)/2 - \mathcal{U}_{KM}\, m_{f'}(1-\gamma_5)/2\right] f$$

$$\mathcal{L}(\pi'^{+} \ell^{+}\nu_\ell) \sim m_\ell/F_{\pi'} \quad (\pi'^{+}\bar{\nu}_\ell\, m_\ell (1+\gamma_5)/2\, \ell) \tag{4.18}$$

where U_{KM} is the Cabibbo-Kobayashi Maskawa mixing matrix. Note that the flavor diagonal couplings of π'^{0} are pure γ_5. This has to be contrasted with the parity conserving fermion couplings of the Higgs in the Weinberg-Salam theory.

Presumably the PGB-fermion-fermion couplings (4.18), which are here obtained in the so-called "monophagic" models are much more general. For example, arguments based on a sigma model type realization of the Hypercolor chiral symmetry [7] will give rise to the generalized Goldberger-Treiman relations among the PGB's and the fermions, also yielding the same order of magnitude as in (4.18).

Perhaps, the reader should be warned that the requirement of monophagic $f \rightarrow F$ couplings is natural only if $m_{\pi'^0} < 100$ GeV. More generally the couplings $\pi'^{0} f\bar{f}'$ could be non-diagonal in which case the $\pi'^{0} f\bar{f}'$ couplings could be both parity violating and parity conserving. Such non-monophagic realizations of EHC symmetry are also present in the literature [27], but apart from pointing out this possibility I shall not follow their consequences.

Assuming the order of magnitude validity of the couplings (4.18) leads to immediate experimental implications. I shall point out the two most obvious reactions, which in forthcoming experiments at e^+e^- accelerators may provide the best chance of observing the neutral as well as the charged pseudo Goldstone bosons, π''s.

First, note that if $m_t > m_{\pi'^{\pm}} + m_b$, then the dominant decays of the t-quark would be

$$t \rightarrow b\,\pi'^{+} \tag{4.19}$$

Not only will the semiweak decay (4.19) dominate the normal(W^+ exchange) weak decays $t \rightarrow b q\bar{q}$, $b\,\ell^{+}\nu_\ell$ of the top quark, but it will also dominate the decays of the toponia, J_T, $J_{T'}$, ... resulting in the final states [12]

$$(J_T,\ J_T',\ \dots) \longrightarrow \pi'^{+} b\bar{t} \qquad \longrightarrow \pi'^{+}\pi'^{-}\, b\bar{b} \tag{4.20}$$

suppressing the canonical decays of J_T

$$J_T \longrightarrow \ell\bar{\ell},\ q\bar{q},\ ggg,\ gg\gamma \quad (\ell = e, \mu, \tau) \tag{4.21}$$

In Fig. (10), a plot for the decay rate $\Gamma(J_T \rightarrow \pi'^{+} b\bar{t})$ is presented for various values of m_t and $m_{\pi'^{\pm}}$. Estimates for the standard decays (4.21) for J_T are O(50 KeV). It is clear that the $\pi'^{\pm}$ mode (4.20) would dominate the J_T decays upto almost the threshold $m_{\pi'^{\pm}} \approx m_T - m_B$. The signatures of (4.20) are: (almost) isotropic mixed lepton-hadron events --- very different from the pure leptonic and hadronic 2- and 3-jet decays from (4.21).

The second, and perhaps the most promising reaction to detect a neutral PGB is the radiative transition

$$J_T \rightarrow (\pi'^0, \pi'^0_\eta) + \gamma \tag{4.22}$$

This is the DSB analog of the Wilczek mechanism [28] involving a fundamental Higgs

$$J_T \rightarrow \varphi^0 + \gamma \tag{4.23}$$

Note that if the π'-f-$\bar{f}$ couplings (4.18) are valid then we expect [12]

$$\frac{\Gamma(J_T \rightarrow \pi'^0 + \gamma)}{\Gamma(J_T \rightarrow \varphi^0 + \gamma)} \simeq \sqrt{2}\, G_F F_{\pi'}^2 = (n_{f'})^{-1}$$

whereas

$$\frac{\Gamma(J_T \rightarrow \varphi^0 + \gamma)}{\Gamma(J_T \rightarrow \mu^+\mu^-)} \simeq \frac{G_F m_{J_T}^2}{4\sqrt{2}\,\pi\alpha}\left(1 - \frac{m_{\varphi^0}^2}{m_{J_T}^2}\right) \simeq 12\% \tag{4.24}$$

where the numbers correspond to m_{J_T} = 40 GeV, $m_{\varphi^0} \simeq$ 15 GeV. Thus, a branching ratio of $\sim$ 1 % is expected for <u>either</u> of the radiative decays (4.22) and (4.23). <u>In fact the radiative decay of the Toponium holds the greatest promise of observing a π'^0 or a Higgs.</u>

The monophagic estimates for the couplings $\pi' f\bar{f}$ also set the decay pattern of the π''s, namely the π''s will decay to the heaviest pair of fermions allowed due to phase space and the Cabibbo angles. Thus, the dominant decays of π'^0 would be

$$\begin{aligned}
(\pi'^0, \pi'^0_\eta) &\rightarrow \tau^+\tau^-, c\bar{c} && \text{for } 2m_b > m_{\pi'^0} > 2m_c \\
(\pi'^0, \pi'^0_\eta) &\rightarrow b\bar{b} && 2m_t > m_{\pi'^0} > 2m_b \\
(\pi'^0, \pi'^0_\eta) &\rightarrow t\bar{t} && m_{\pi'^0} > 2m_t
\end{aligned} \tag{4.25}$$

Recalling the result from the last section that $\Gamma(Z \rightarrow \pi'^0\mu^+\mu^-)/\Gamma(Z \rightarrow \mu^+\mu^-)$ we note that the observation of a positive signal in the radiative transition $J_T \rightarrow$ scalar + γ and the lack of evidence for the process Z $\rightarrow$ scalar + $\mu^+\mu^-$ would be a vote in favour of Dynamical Symmetry breaking, observation of positive signals in both the J_T and Z decays a vote for the canonical Higgs; non-observation of any signal in these processes presumably means that the nature of spontaneous symmetry breaking would remain a puzzle for a long time to come!

So far we have concentrated on the dominant $\pi' f\bar{f}$ couplings that arise from the first term of the effective Lagrangian (4.15). Let us now consider the second term in the Lagrangian. The second part of the Lagrangian (4.15) gives smaller contribution $\sim O(g^2 m_f^2/m_W^2)$ and could be cast as [16]

$$\frac{1}{4\sqrt{2}} F_{\pi'}\, \pi' \bar{f} \left\{ m_f \left[\Gamma_{RR}^{ff'}(1+\gamma_5) - \Gamma_{LL}^{ff'}(1-\gamma_5) \right] - \left[\Gamma_{RR}^{ff'}(1-\gamma_5) - \Gamma_{LL}^{ff'}(1+\gamma_5) \right] m_{f'} \right\} f' \tag{4.26}$$

where

$$\left(\Gamma^{ff'}_{\substack{LL\\RR}}\right)_{nm} \equiv \sum_E \frac{1}{\mu_E^2}\, g^E_{\substack{Ln\\R}}(f)\, g^{E*}_{\substack{Lm\\R}}(f') \tag{4.27}$$

Now in the basis in which the fermion mass matrix is diagonal, Γ_{LL}^{ff} and Γ_{RR}^{ff} are in general not simultaneously diagonalizable. Thus, (4.26) leads to flavor-changing couplings of π'^0 and π'^0_η <u>either</u> for charge 2/3 quarks <u>and/or</u> for charge -1/3 quarks. These non-diagonal couplings are of the order of magnitude

$$\frac{G_F}{\sqrt{2}}\; m_{2/3}\; m_{(-1/3)}\; \sin\theta_c \qquad (4.28)$$

Even if the tree-level couplings (4.28) are absent for the charge 2/3 or charge -1/3 quarks, there would be radiative corrections [29] inducing off-diagonal transitions at the same rate as for the Weinberg-Salam fundamental Higgs case, where $(1-\gamma_5)$ couplings are induced and are of order

$$\left(\sqrt{2}\, G_F\right)^{3/2}/8\pi^2 \left(m_{(-1/3)}\, m_{2/3}^2 \text{ or } m_{2/3}\, m_{(-1/3)}^2\right) \sin\theta_c \times \ln\left(m_W^2/m_f^2\right) \qquad (4.29)$$

The couplings (4.29) are rather small. However, both (4.28) and (4.29) have interesting consequences for the $K^{\pm}$ decays, namely if the decay $K^+ \to \pi^+ \pi'^0$ were kinematically accessible then (4.28) would give [16]

$$\frac{\Gamma(K^+ \to \pi^+ \pi'^0)}{\Gamma(K^+ \to \pi^+ \pi^0)} \geqslant O(1) \qquad (4.30)$$

whereas (4.29) would give $\sim O(10^{-5})$ for the same ratio. Since even this branching ratio is ruled out by data, one gets a bound for the mass of π'^0 and π'^0_η

$$\left(m_{\pi'^0}, m_{\pi'^0_\eta}\right) > 350 \text{ MeV} \qquad (4.31)$$

The existence of the π'^0 below the D and B meson masses having couplings of order (4.28) can be tested in the neutral current decays of the b and c quarks [16]

$$c \to u + (\pi'^0, \pi'^0_\eta) \qquad (4.32)$$

and

$$b \to (d, s) + (\pi'^0, \pi'^0_\eta) \qquad (4.33)$$

Presumably (4.33) as the <u>dominant</u> transition for the b quark decay is already ruled out on the basis of the CESR result $Y^B_{\mu\mu} \equiv BR(B \to \mu^+\mu^- X) < 1.3\ \%$ [30]. The coupling (4.28) for the charge -1/3 quark decays would give dominance of the decay mode (4.33) and a concomittant branching ratio $Y^B_{\mu\mu} \geqslant 2\ \%$. Non-observation of (4.33) would mean that either the estimates (4.28) do not apply to the charge = -1/3 quarks and/or else $m_{\pi'^0} > m_B - m_K$. It would be interesting to check (4.32) through the decays

$$D \to (\pi'^0, \pi'^0_\eta) + \text{pions} \qquad (4.34)$$

which could have as large a branching ratio as 1 %. The non-existence of both (4.32) and (4.33) would mean that $m_{\pi'^0} \geqslant 1.7$ GeV, or else the couplings (4.28) are gross overestimates. It is quite conceivable that $m_{\pi'^0} > m_D$ but the coupling (4.28) is applicable <u>only</u> in the decays of charge = + 2/3 quarks, in which case none of the processes (4.32) and (4.33) are permitted. It has been argued [16] that a value of $m_{\pi'^0} \simeq O(2\text{ GeV})$ would have interesting consequences for the decay $K_L \to \mu e$ in monophagic models, though the argument is somewhat circular. I shall repeat the argument for the sake of completeness.

In "monophagic" models the neutral Goldstone bosons π'^0, π'^0_η <u>do not</u> receive any mass from the $U(1) \otimes SU(2)_L \otimes SU(3)_c \otimes HC$ forces. Hence, there must be Pati-Salam type leptoquark forces to prevent the π'^0, π'^0_η from remaining exactly massless. The contribution of the Pati-Salam gauge bosons to $m_{\pi'}$ has been estimated in ref. (10). They find

$$m_{\pi'^0} \sim \left(\frac{310\text{ TeV}}{m_{PS}}\right) \times 1.7\text{ GeV} \qquad (4.35)$$
$$m_{\pi'^0_\eta} \sim \left(\frac{310\text{ TeV}}{m_{PS}}\right) \times 1.5\text{ GeV}$$

where m_{PS} is the mass of the Pati-Salam boson in TeV. This gives an upper bound on m_{PS}, which would lead to [10]

$$R_{\mu e} \equiv \Gamma(K_L^0 \to \mu e)/\Gamma(K^+ \to \mu^+ \nu_\mu) \geqslant 6 \times 10^{-10} \,(310\,\mathrm{GeV}/m_{PS})^4 \qquad (4.36)$$

The value (4.36) grazes the present experimental limit. Of course, there are uncertainties of at least an order of magnitude in (4.35) and (4.36) but it is interesting that the EHC theories predict rare decay modes which would soon be subject to experimental tests in the decays of B, D and K mesons. Perhaps it should be stressed that the bounds (4.35) and the related $K_L{}^0 \longrightarrow \mu e$ rate are valid only in the monophagic EHC models. A much bigger estimate [11] for $m_{\chi'^0}$ in non-monophagic models also exists in the literature and has already been presented in the previous section. In such models the Pati-Salam leptoquark bosons which could be responsible for the $K_L \to \mu e$ decay do not set the scale of the PGB masses, $m_{\chi'^0}, m_{\chi'^0_\eta}$, and the non-existence of the process (4.32), (4.33) and the decay $K_L \longrightarrow \mu e$ are all compatible with each other and with the EHC theories.

V. Problems with Flavor Changing Neutral Currents in Extended Hypercolor Theories

The generic structure of the Extended Hypercolor Theories is shown in Fig. (11), which suggests that, in general, there are EHC bosons inducing horizontal transitions $d \longleftrightarrow s$, $s \longleftrightarrow b$ etc. among ordinary quarks [8]. The point is that if the EHC generators commute with the $SU(2)_L$ group, which I guess is natural to expect, then the eigenstates of the EHC are the same Cabibbo rotated eigenstates of $SU(2)_L$. Thus the EHC current has the form

$$J_\mu^{q,E} \sim \bar{Q}_L \gamma_\mu q_{L\theta} \qquad (5.1)$$

and the horizontal transition is generated by the commutator

$$\left[J_0^{s,E},\ J_\mu^{\dagger\, d,E}\right] \sim \bar{s}_L \gamma_\mu d_L \sin\theta \cos\theta \qquad (5.2)$$

$$\Rightarrow \quad m(\Delta S = 2) \sim \left(g_E^2/m_E^2\right) \sin\theta\cos\theta \qquad (5.3)$$

It follows that unless $g_E^2/m_E^2 < 2 \times 10^{-5}\ \mathrm{TeV}^{-2}$, there is a severe problem of generating unacceptably large flavor changing neutral currents FCNC in the EHC theories. Recall that structurally, there was a similar problem with the Cabibbo current with the u,d and s quarks before the GIM mechanism [31] was discovered. The strength of the EHC induced effective Fermi type coupling g_E^2/m_E^2 is more or less fixed by the diagram shown in Fig. (2) and the requirement that it generates the correct fermion masses through the expression

$$m_f = g_E^2/2m_E^2 \,\langle \bar{F} F \rangle_0 \qquad (5.4)$$

Dimopolous and Ellis [8] have evaluated $\langle \bar{F}F \rangle_0$ using scaling arguments and PCAC for ordinary quarks, obtaining

$$\langle \bar{F}F \rangle_0 = \frac{m_{\pi'}^2}{(m_d + m_u) f_\pi} \left(\frac{1}{\sqrt{2} G_F n_{f'}}\right)^{3/2} \left(\frac{N_c'}{3}\right)^{1/2} \qquad (5.5)$$

which for $n_{f'} = 4$, $m_d/m_u = 1.8$, $m_d = 10$ MeV and $f_\pi = 95$ MeV gives

$$m_f = 0.0035/(v_E^2)_f \left(\frac{N_c'}{4}\right)^{1/2} \mathrm{TeV}$$
$$(v_E^2)_f = \left(m_{E_f}^2/g_E^2\right) \qquad (5.6)$$

which states that in order to generate m_d = 10 MeV, $(v_F^2)_f$ has to be $\sim 650\ (\mathrm{TeV})^2$ which is in conflict with the requirement from $\Delta S = 2$ transition $(m^2_{E_d}/g_E^{\,2}) \gtrsim 5 \times 10^4\ \mathrm{TeV}^2$ by at least two orders of magnitude. The same problem exists with the $\Delta c = 2$ transitions. The resolution of this problem lies in preventing the FCNC transitions at the tree diagram level shown in Fig. (12) which for example generates $\Delta S = 2$ transitions at $\sim G_E \sin\theta \cos\theta$. How to go about it without generating unwanted Goldstone bosons or rotating away the Cabibbo angles is not yet clear. Once such a mechanism is discovered it would be interesting to investigate its consequences in order $G_E\,\alpha$ and $G_E^{\,2}$.

It is conceivable that there are some unsuspected symmetries that prevent (5.2) and arrange cancellation of FCNC transitions at the tree level but that they leave the flavor diagonal couplings $\pi'^{\,0} f\bar{f}$ and the charged PGB's couplings $\pi'^{\,\pm} f\bar{f}'$ very much intact as expressed through Eqs. (4.18). In that case all that has been said about the couplings of the PGB's with fermions and their phenomenological implications will hold in the forthcoming complete theory. However, it is also conceivable that there is some other (as yet undiscovered) mechanism to generate fermion masses and the EHC forces are indeed of the Pati-Salam type, thereby inducing very feeble $\pi' f\bar{f}$ couplings. The search of PGB's has to be carried out in a more general way than that of the fundamental Higgs. This is the next topic that I want to discuss.

VI. Production of charged PGB's ($\pi'^{\pm}$) in e^+e^- annihilation

I shall concentrate on e^+e^- annihilation, not because of any institutional affinity, but simply because high energy e^+e^- machines provide the best chance of detecting the PGB's (both charged and neutral). I have already noted that

the best chance of detecting a neutral PGB, $\pi'^{\,0}, \pi'^{\,0}_{\eta}$ lies in the radiative decays of the toponium, $J_T \to (\pi'^{\,0}, \pi'^{\,0}_{\eta}) + \gamma$, provided $(m_{\pi'^0}, m_{\pi'^0_\eta}) < m_{J_T}$. One has to look for monoenergetic photons recoiling against a mixed lepton-hadron jet, almost back to back.

The processes which have the best chance of producing the charged PGB's have all been enumerated earlier, and I reiterate them here at the risk of repetition:

$$e^+e^- \longrightarrow \pi'^{+}\pi'^{-} \tag{3.2}$$

$$Z^0 \longrightarrow \pi'^{+}\pi'^{-} \tag{3.3}$$

and if the estimates of the π'-f-$\bar{f}$ couplings of section 4 are correct, then

$$t \longrightarrow b\,\pi'^{+} \tag{4.19}$$

and

$$J_T \longrightarrow b\bar{t}\,\pi'^{+},\quad b\bar{b}\,\pi'^{+}\pi'^{-} \tag{4.20}$$

with (3.2) being theoretically the cleanest. Both the processes (3.2) and (3.3) have the all important feature that they have an angular dependence

$$d\sigma/d\Omega \sim \sin^2\theta \tag{6.1}$$

which is to be contrasted with the usual e^+e^- background which follows

$$d\sigma/d\Omega \sim 1 + \cos^2\theta + 1/\gamma\,\sin^2\theta \tag{6.2}$$

$$\gamma = E/M$$

This for example, is the angular dependence of the process $e^+e^- \to 2$ jets in QCD and of a charged heavy lepton pair production $e^+e^- \to L^+L^-$. M in (6.2) is then m_Q or m_L.

What are the signatures of (3.2) and (3.3)? It all depends on how the $\pi'^{\pm}$ decay. Of course, this is intimately related to the couplings of $\pi'^{\pm}$ with a fermion pair. In any case there are three possible decay modes of $\pi'^{\pm}$:

(i) the purely fermionic mode $\pi'^{\pm} \to f\bar{f}'$ (7) $\qquad$ (6.3)

(ii) the semifermionic modes (17) $\quad \pi'^{\pm} \to \pi'^{0} + (q\bar{q}', \ell^{\pm}\nu_\ell)$ $\qquad$ (6.4)

and

(iii) the semifermionic mode involving a photon (17,32)

$$\pi'^{\pm} \to \gamma + (\ell^{\pm}\nu_\ell, q\bar{q}') \tag{6.5}$$

The relevant Feynman (51) diagrams are shown in Fig. (13). If the estimates of section 4 are of the right order of magnitude, then the decays (6.3) are semiweak $\Gamma \sim G_F$ whereas (6.4) are weak ($\Gamma \sim G_F^2$) and (6.5) are higher order processes $\quad \Gamma \sim G_F \alpha^3$. Presumably these orders of magnitude should be good guesses of the relative importance of (6.3)-(6.5). Anyway, it is straight-forward to calculate their signatures in e^+e^- annihilation and they are listed below.

(i) Fermionic decays of $\pi'^{\pm}$

The general form of the coupling $\pi'^{\pm} f\bar{f}'$ can be expressed as

$$g_{\pi'^{\pm} f\bar{f}'} \sim m_f/F_{\pi'} \; U_{f\bar{f}'} \; K_f \tag{6.6}$$

where $U_{f\bar{f}'}$ are mixing angles and in monophagic models one could show that they

are just the Cabibbo-Kobayashi-Maskawa angles. K is a factor which depends upon the details of the EHC/HC representation, i.e. upon the $SU(3)_c \otimes SU(2)_L \otimes U(1)$ properties of the massive bosons coupled to the generators of the EHC/HC quotient group. In monophagic models it is normally given by $K_f \sim (\sqrt{3})^{\pm 1}$ with the sign of the exponent depending on the details of the EHC/HC generators.

The estimate (6.6) states that the π''s couple to the pair of heavy fermions allowed by kinematics and Cabibbo angles. Thus, for $m_{\pi'^{\pm}} < m_t + m_b$ the dominant decays are

$$\pi'^{+} \to c\bar{s} \tag{6.7}$$

$$\pi'^{+} \to \tau^{+}\nu_\tau \tag{6.8}$$

The decay $\pi'^{+} \to c\bar{b}$ presumably is smaller than (6.7) being Cabibbo suppressed. Above $m_t + m_b$, the dominant mode would be

$$\pi'^{+} \to t\bar{b} \tag{6.9}$$

Since the present PETRA limit on the quantity $m_t + m_b = 23$ GeV (33) and the uppermost PETRA/PEP energy is ~ 46 GeV, let me just concentrate on the decays (6.7) and (6.8). The model dependence through K would reflect itself as

$$\Gamma(\pi'^{+} \to c\bar{s}) / \Gamma(\pi'^{+} \to \tau^{+}\nu_\tau) \simeq \chi \frac{m_c^2}{m_\tau^2}$$

where $\chi \equiv (K_{c\bar{s}}/K_{\tau\nu_\tau})^2$. Now, there are three possibilities

(a) $\chi \sim 1$: the favorable decay chain in this case would be

$$
\begin{array}{l}
e^+e^- \longrightarrow \chi'^+\chi'^- \\
\quad \hookrightarrow \bar{c}s \\
\qquad \hookrightarrow \text{hadrons} \\
\quad \hookrightarrow \tau^+\nu_\tau \\
\qquad \hookrightarrow (e^+, \mu^+, \cdots) + \nu_\tau
\end{array}
\tag{6.10}
$$

This will lead to the signature

$$e^+e^- \rightarrow \chi'^+\chi'^- \longrightarrow (e,\mu)^\pm + \text{hadrons} \tag{6.11}$$

The background to (6.11) from the usual $\tau^+\tau^-$ production $e^+e^- \rightarrow \tau^+\tau^- \longrightarrow (e,\mu)^\pm$ + hadron is well separable because of kinematics and topology. This is shown in Fig. (14) for some representative values of $m_{\chi', t}$.

(b) $\chi \ll 1$: The dominant decay chain in this case would be

$$
\begin{array}{l}
e^+e^- \longrightarrow \chi'^+\chi'^- \\
\quad \hookrightarrow \tau^-\nu_\tau \\
\qquad \hookrightarrow (e^-, \mu^-, \cdots) + \nu's \\
\quad \hookrightarrow \tau^+\nu_\tau \\
\qquad \hookrightarrow (e^+, \mu^+, \cdots) + \nu's
\end{array}
\tag{6.12}
$$

This will give rise to very low multiplicity two-jet events including "anomalous dilepton events" -- very much reminiscent of the $\tau^+\tau^-$. However, the process (6.12) will differ from the $\tau^+\tau^-$ induced background in several important details: the e^+e^-, $e^\pm\mu^\mp$ and $\mu^+\mu^-$ events arising from (6.12) would be highly non-collinear and <u>non-coplanar</u> with a large missing momentum (17). Since the e (or μ) from (6.11) and (6.12) are expected to be very soft with $\langle E_{e,\mu}\rangle \simeq \frac{1}{6}E_{beam}$, <u>a good low energy lepton detection is of utmost importance.</u>

The distinction between the processes (6.11) and (6.12) and the similar final states -- $\ell^+\ell^-$ and $\ell^\pm$ hadrons -- which can also be produced from the decays of a pair of heavy leptons $e^+e^- \rightarrow L^+L^-$ can be made through the threshold factor $\Delta R\,(\chi'^+\chi'^-) \sim (1/4)\beta^3$ vs. $\Delta R(L^+L^-) \sim \beta$ and the angular dependence (6.1) for $\chi'^+\chi'^-$ vs. (6.2) for L^+L^- production. The distributions $\frac{1}{\sigma}\frac{d\sigma}{d\cos\theta_\ell}(e^+e^- \rightarrow \chi'^+\chi'^- \rightarrow \ell^\pm_{+\cdots})$ and $\frac{1}{\sigma}\frac{d}{d\cos\theta_\ell}(e^+e^- \rightarrow L^+L^- \rightarrow \ell^\pm_{+\cdots})$ are shown in Fig. (15) for some representative values of $m_{\chi'}$ and m_{L} t .

(c) $\chi \gg 1$: This would lead to events of the type

$$e^+e^- \rightarrow \chi'^+\chi'^- \longrightarrow \text{hadrons} \tag{6.13}$$

Near the threshold (6.13) will give rise to almost isotropic events -- large sphericity and acoplanarity -- but with $\Delta R \sim (1/4)\beta^3$ and hence seemingly imperceptible. Away from the threshold the hadronic events tend to be 2-jet-like, however, with an angular distribution given by (6.1). This could be used to establish a signal though on an event by event basis, it would be hard to tell (6.13) apart from the usual 2-jet background --- so abundant in e^+e^- annihilation. In high statistic experiments (6.13) could be searched through a step in R. This is going to be painstaking!

<u>(ii) Lowest order Weak Semifermionic modes</u>

The decay

$$
\begin{array}{l}
\chi'^\pm \longrightarrow \chi'^0 + W^\pm_{\mathrm{vir}} \\
\qquad\qquad\qquad \hookrightarrow \ell^\pm\nu_\ell,\ q\bar{q}'
\end{array}
$$

is analogous to the well known χ^0_{e3} decay $\chi^0 \rightarrow \chi^+ e^- \bar{\nu}_e$
A straightforward calculation (neglecting q^2 dependence) gives

$$\Gamma(\pi'^{\pm} \to \pi'^{0} \ell^{\pm} \nu_\ell) = \frac{G_F^2 m_{\pi'^\pm}^5}{384\pi^3} f\left(m_{\pi'^0}^2 / m_{\pi'^\pm}^2\right) \qquad (6.14)$$

where $f(x) = 1 - 8x + 8x^3 - x^4 - 12x^2 \ln x$

The process (6.4) would lead to the final state

$$\begin{aligned} e^+e^- \to\ & \pi'^+ \pi'^- \\ & \hookrightarrow \pi'^0 + (\ell^- \nu_\ell,\ q\bar{q}') \\ & \qquad \hookrightarrow \text{hadrons} \\ & \hookrightarrow (\text{hadrons} + (\ell^+ \nu_\ell,\ \bar{q} q')) \end{aligned} \qquad (6.15)$$

leading to mixed lepton-hadron events which could be separated from the $c\bar{c}$,$b\bar{b}$ decay products due to topology and angular dependence. However, in the scenario where $\pi'^{\pm} f \bar{f}'$ couplings are suppressed, it may be that the couplings $\pi'^0 f\bar{f}$ are also suppressed. In that case the decay mode

$$\pi'^0 \to 2\gamma$$

may become important. This would lead to spectacular signatures in e^+e^- annihilation through

$$\begin{aligned} e^+e^- \to\ & \pi'^+ \pi'^- \to 4\gamma + \ell^+\ell^- \\ \to\ & 4\gamma + \ell^{\pm} + \text{hadrons} \\ \to\ & 4\gamma + \text{hadrons} \end{aligned} \qquad (6.16)$$

(iii) Higher order semifermionic modes

The decays

$$\pi'^{\pm} \to \gamma + W^{\pm}_{\text{vir}} \hookrightarrow \ell^{\pm} \nu_\ell,\ q\bar{q}'$$

is reminiscent of the second order electromagnetic decay $\pi^0 \to 2\gamma$, and likewise goes via the triangle diagram. The decay width has been calculated in ref. (17)

$$\begin{aligned} \Gamma(\pi'^{\pm} \to \gamma \ell^{\pm} \nu_\ell) = \ & \frac{G_F (\alpha/\pi)^3 m_{\pi'}}{192\sqrt{2}\,\pi \sin^4\theta_W} \left(m_{f'}^3 N_{c'}^2/9\right) \\ & \times \Big[\frac{1}{9 m_{\pi'}^2} \left\{ m_{\pi'}^4 - 6(m_W^2 - m_{\pi'}^2)(4m_W^2 - 3m_{\pi'}^2) \right\} \\ & + \frac{2}{3 m_{\pi'}^4} \left\{ m_{\pi'}^6 - 6 m_{\pi'}^4 m_W^2 + 9 m_{\pi'}^2 m_W^4 - 4 m_W^6 \right\} \\ & \times \ln\left((m_W^2 - m_{\pi'}^2)/m_W^2\right) \Big] \end{aligned} \qquad (6.17)$$

Rates for the semileptonic processes (6.4) and (6.5) are given in table (1), which corroborates the remarks made earlier about the relative importance of the $\pi'^{\pm}$ decays. In any case the process (6.5) has beautiful signatures in e^+e^- annihilation giving rise to events of the type

$$\begin{aligned} e^+e^- \to\ & 2\gamma + \text{hadrons} \\ \to\ & 2\gamma + (\ell^+, \ell^-) + \text{hadrons} \\ \to\ & 2\gamma + (\ell^+\ell^-) + \nu\text{'s} \end{aligned} \qquad (6.18)$$

with energetic photons and leptons.

The signatures (6.11), (6.12), (6.13), (6.15), (6.16) and (6.18) also hold for the decays $Z^0 \to \pi'^+\pi'^-$ and could be used to search for the charged PGB's $\pi'^\pm$ up to a mass $m_{\pi'^\pm} \sim (m_{Z/2}$ - a few GeV).

VII. Conclusions

For many the Weinberg-Salam theory is a provisional framework whose phenomenological success as an effective Lagrangian is certainly impressive but which has to be replaced by a forthcoming complete theory, following the example of QCD. Dynamical symmetry breaking is an attempt in that direction where the order parameter $\langle\varphi^0\rangle_0$ is replaced by a condensate $\langle\bar{F}F\rangle_0$ whose value is related to the scale of a new strong interaction --- very much like $\langle\bar{q}q\rangle_0$ in QCD signals the presence of a hadronic scale characterized by $\Lambda_{QCD} \sim 0(100$ MeV).

The example of the Weinberg-Susskind [5] model of DSB shows very elegantly how to generate the weak gauge boson masses and yet keep the weak $\Delta I = 1/2$ rule intact naturally. This part of DSB is certainly impressive. However, it is equally likely that the Weinberg-Susskind model is not complete; a realistic realization of the Weinberg-Susskind idea will almost certainly entail larger groups and representations. Almost all realizations involve pseudo Goldstone bosons with masses that are parametrically small compared to the scale of the hypercolor interactions, $\Lambda_{HC} \sim 1$ TeV. Some of them are protected from getting masses in the lowest order and could therefore be very light. Most estimates predict the masses of these colorless PGB's to be of order 10 GeV though they could be as large as 50 GeV [9-11].

These PGB's, particularly the neutral ones, could be confused with the standard fundamental Higgs, apart from their parity assignment. It turns out

that the 0^- nature of the PGB's introduces a fundamental difference in their production rate vis à vis the 0^+ Higgs particle giving rise to the ratios [12,13]

$$\frac{\Gamma(Z \to (\pi'^0, \pi'^0_\eta) + \mu^+\mu^-)}{\Gamma(Z \to \varphi^0 + \mu^+\mu^-)} < 10^{-4}$$

and

$$\frac{\sigma(e^+e^- \to (\pi'^0, \pi'^0_\eta)\, Z)}{\sigma(e^+e^- \to \varphi^0 Z)} < 10^{-5}$$

The hypercolor approach to generate the gauge boson masses is impressive; attempts to generate fermion masses in an Extended Hypercolor framework [6] are equally unimpressive. More importantly the EHC theories in their present form are ruled out by experiments on Flavor Changing Neutral Currents. The difficulty with the FCNC transitions has a bearing on the couplings of the π''s with a fermion pair. It is conceivable that a generalized Goldberger-Treiman type estimate $g_{\pi' f\bar{f}} \sim m_f/F_{\pi'}$ would survive the act of exorcising the EHC theories of the FCNC ghost. These estimates could be used to guess the production and decays of π''s involving ordinary fermions. Based on these estimates it is expected that the <u>best place of producing neutral PGB's, π'^0, π'^0_η is in the radiative decays of toponium, $(J_T \to \pi'^0, \pi'^0_\eta) + \gamma$</u> [12] for which a branching ratio $\sim$ (1 %) is expected.

The production of charged PGB's, $\pi'^\pm$, on the other hand would be copious in e^+e^- annihilation, in Z decays and if kinematics permits in the decays of the top quark as well as in toponia. The signatures of $\pi'^{\pm}$ in all these processes are very healthy and our experimental colleagues at PETRA and PEP are well advised to search for these signals.

There are other consequences of hypercolor theories; some of which in the form of rare decays $K_L \to \mu e$ and $\mu \to e\gamma$ are well documented in the literature [10,16,18] and we discussed them in the context of monophagic models. Production mechanisms of the PGB's other than the ones discussed here also exist in literature, in particular the processes [34] $e^+e^- \to e^+e^- + \pi'^0$, $e^+e^- \to e^+e^- + \pi'^+\pi'^-$, $ep \to ep + \pi'^0$ [35], $pp(\bar{p}) \to (\pi'^0, \pi'^0_\eta) + X$ [10] and the production of the colored 0^- bosons θ_η [10] in the process $pp(\bar{p}) \to \theta_\eta + X$ have been worked out. However, none of them appear very encouraging either due to the miniscule production cross-sections and/or due to the problem of finding a scalar particle inside a hadronic junk.

In conclusion, e^+e^- machines hold the greatest promise of unravelling the nature of spontaneous symmetry breaking.

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References

1) S. Weinberg, Phys. Rev. Lett. 19, 1264 (1967);
A. Salam, in Elementary Particle Physics (ed. N. Svartholm) (Almqvist and Wiksells, 1968).

2) The point of view of generating masses dynamically was first expressed by M.A.B. Bèg and A. Sirlin in the context of electroweak gauge theories; M.A.B. Bèg and A. Sirlin, Ann. Rev. Nucl. Sci. 24, 379 (1974) and given a concrete form by S. Weinberg, Phys. Rev. D13, 974 (1976). Of course, the idea of dynamical symmetry breaking is very old and can be traced to the seminal paper of G. Jona Lasinio and Y. Nambu, Phys. Rev. 122, 345 (1961). For subsequent development of dynamical symmetry breaking see also
J. Schwinger, Phys. Rev. 125, 397 (1962), 128, 2425 (1969);
F. Englert and R. Brout, Phys. Rev. Lett. 13, 321 (1964);
R. Jackiw and K. Johnson, Phys. Rev. D8, 2386 (1973);
J.M. Cornwall and R.E. Norton, Phys. Rev. D8, 3338 (1973);
S. Coleman, Comm. in Math. Phys. 31, 259 (1973);
L. Dolan and R. Jackiw, Phys. Rev. D9, 3320 (1974).

3) It is conceivable that one could define a renormalization group improved perturbation theory where this problem could be circumvented. Nevertheless, the problem of understanding the mass scales still remains very much intractable.

4) See, for example, the review article by H. Pagels, Phys. Rep. 16C, 219 (1975) and references quoted therein.

5) Most of what is being discussed in this talk is in the spirit of
S. Weinberg, Phys. Rev. D19, 1277 (1979) and
L. Susskind, Phys. Rev. D20, 2619 (1979).

6) For one-Hypercolor-family model see
E. Farhi and L. Susskind, Phys. Rev. D20, 3404 (1979) and
S. Dimopolous, Nucl. Phys. B168, 69 (1980)

7) M.A.B. Bèg, H.D. Politzer and P. Ramond, Phys. Rev. Lett. 43, 1701 (1979).

8) S. Dimopolous and L. Susskind, Nucl. Phys. B155, 237 (1979) and
E. Eichten and K. Lane, Phys. Lett. 90B, 125 (1980).

9) J.D. Bjorken, unpublished;
S. Dimopolous in ref. (6);
M.E. Peskin, C.E.N. Saclay Preprint DPh-T-80-46 (1980);
J. Preskill, Harvard Univ. Preprint HUTP-80-A033 (1980);
S. Chada and M.E. Peskin,CERN Report TH-3038 (1981).

10) The smallest PGB mass estimate is due to
S. Dimopolous, S. Raby and G.L. Kane, Univ. of Michigan Preprint
UM HE 80-22 (1980).

11) The largest (published) PGB mass estimate is due to
E. Farhi and L. Susskind, CERN Report TH-2975 (1980).

12) A. Ali and M.A.B. Bèg, DESY Report 80/98 (1980), to appear in Phys. Lett. B.

13) J. Ellis, M.K. Gaillard, D.V. Nanopolous and P. Sikivie,
Nucl. Phys. B182, 529 (1981).

14) M. Dine, E. Farhi and L. Susskind, unpublished;
S. Dimopolous and J. Ellis, Nucl. Phys. B182, 505 (1981).

15) S. Chada and M.E. Peskin, CERN Report TH-3023 (1981).

16) J. Ellis, D.V. Nanopoulos and P. Sikivie, CERN Report TH-3030 (1981).

17) A. Ali, H.B. Newman and R.Y. Zhu, DESY Report 80/110 (1980), to be published
in Nucl. Phys. B.

18) For earlier reviews on the implications of dynamical symmetry breaking see
E. Farhi and L. Susskind (ref. 11);
K.D. Lane and M.E. Peskin, Nordita Report 80-33 (1980);
P. Sikivie, CERN Report TH-2951 (1980);
M.A.B. Bèg, Recent developments in High Energy Physics (ed. B. Kursunoglu,
A. Perlmutter and L.F. Scott) (Plenum Publishing Corporation, 1980).

19) P. Sikivie, L. Susskind, M. Voloshin and V. Zakharov, Nucl. Phys. B173,
189 (1980).

20) R. Dashen, Phys. Rev. 183, 1245 (1969).

21) T. Das, G.S. Guralnik, V.S. Mathur, F.E. Low and J.E. Young,
Phys. Rev. Lett. 18, 759 (1967).

22) J.D. Bjorken, SLAC Report No. SLAC-PUB-1866 (1977).

23) J. Ellis, M.K. Gaillard and D.V. Nanopoulos, Nucl. Phys. B106, 292 (1976);
B.W. Lee, C. Quigg and H.B. Thacker, Phys. Rev. D16, 1519 (1977).

24) R.N. Cahn, M.S. Chanowitz and N. Fleishon, LBL Report No. LBL-849 (1978).

25) J.P. Le.veille, Wisconsin Report COO-881-86 (1979).

26) S.L. Glashow and S. Weinberg, Phys. Rev. D15, 1958 (1977).

27) See for example K. Lane, Ohio State University Report DOE/ER/01545-306 (1981)

28) F. Wilczek, Phys. Rev. Lett. 39, 1304 (1977).

29) M. Wise, Harvard University Preprint HUTP-80/A086 (1980).

30) K. Chadwick et al., Phys. Rev. Lett. 46, 88 (1981).

31) S.L. Glashow, J. Iliopolous and L. Maiani, Phys. Rev. D2, 1285 (1970).

32) G.L. Kane private communication to Lane and Peskin (ref. 18).

33) See for example B. Wiik, in Proceedings of the XX International Conference on High Energy Physics, Madison, Wisconsin (1980), (Editors: L. Durand and L.G. Poudron)

34) J.A. Grifols, Universitat Autónoma de Barcelona, Report UAB-FT-69 (1981).

35) S. Rudaz and J.A.M. Vermaseren, CERN Report TH-2961 (1980) and Erratum!

$m_{\pi'^{\pm}}$ (GeV)	$\Gamma(\pi'^{\pm} \to \pi'^{o} + \ell^{\pm}\nu_e)$ (a)	$\Gamma(\pi'^{\pm} \to \gamma + \ell^{\pm}\nu_e)$ (b)
10.0	3.3×10^{-3} eV	1.2×10^{-5} eV
20.0	11.4 eV	1.6×10^{-3} eV
30.0	0.167 KeV	2.8×10^{-2} eV
40.0	0.91 KeV	0 .23 eV
50.0	3.3 KeV	1.23 eV

TABLE 1

(a) The entries correspond to assuming $m_{\pi'^o} = 8$ GeV.

(b) The numbers correspond to assuming 2 Techniflavors x 3 Colors x 4 Technicolors in the evaluation of the triangle diagram.

Figure Captions

Fig. 1: Lowest order electroweak contribution to the mass of the π'^{+}.

Fig. 2: Fermion mass generation mechanism in the Extended Hypercolor Theories.

Fig. 3: Feynman diagrams for the couplings of the pseudo Goldstone bosons with electro weak gauge bosons.

Fig. 4: Triangle diagram contributing to the $v^i v^j \pi'$ couplings in the hypercolor theories. The loop involves hyperfermions.

Fig. 5: Feynman diagrams contributing to the processes (3.11) - (3.14) involving a pseudo Goldstone boson, π', in hypercolor models and the processes (3.11') - (3.14') involving the canonical Higgs field in the standard electroweak theory.

Fig. 6: The ratio $\Gamma(Z \to \pi'^{0}\mu^+\mu^-)/\Gamma(Z \to \varphi^0 \mu^+\mu^-)$ for equal values of m_{φ^0} and $m_{\pi'^0}$. n_F, is the number of hyperflavor doublets and N_c, the number of hypercolors (from ref. 12).

Fig. 7: The dimuon invariant-(mass)^2 distribution from the decays $Z \to \pi'^{0}\mu^+\mu^-$ and $Z \to \varphi^0\mu^+\mu^-$. Both the distributions are normalized to the same area $(= \Gamma(Z \to \varphi^0\mu^+\mu^-))$ for equal m_{φ^0} and $m_{\pi'^0}$. The relative scales can be read off fig. 6 (from ref. 12).

Fig. 8: The ratio $\sigma(e^+e^- \to \pi'^{0}Z)/\sigma(e^+e^- \to \varphi^0 Z)$ for $\sqrt{s}$ = 140 GeV, 170 GeV and 200 GeV and equal values of m_{φ^0} and $m_{\pi'^0}$, $\sin^2\theta_W = 0.20$ and m_Z = 94 GeV (from ref. 12).

Fig. 9: Lowest order contribution of a pseudo Goldstone boson exchange to the $\Delta S = 2$ $s\bar{d} \leftrightarrow \bar{s}d$ transition.

Fig. 10: The rate for the semiweak decay of toponium, J_T, as a function of the charged hyperpion mass. n_F, is the number of hyperflavor doublets (from ref. 12).

Fig. 11: The generic structure of the Extended Hypercolor Theories.

Fig. 12: a) Lowest order contribution of an Extended Hypercolor boson exchange to the flavor changing neutral current transitions.
b) Box diagram contributing to the FCNC transitions

Fig. 13: Feynman diagrams contributing to the decay of a charged pseudo Goldstone boson
a) $\pi'^{+} \to f\bar{f}'$ in the Extended Hypercolor Theories
b) Lowest order semileptonic weak decays

$$\pi'^{+} \to \pi'^{0} + (l^+\nu_l,\ q\bar{q}')$$

c) Higher order semileptonic weak decays

$$\pi'^{+} \to \gamma + (l^+\nu_l,\ q\bar{q}')$$

Fig. 14: Comparison of the normalized hadron thrust distribution from the process $e^+e^- \to \pi'^{+}\pi'^{-} \to (e,\mu)$ + hadrons with the hadron thrust distribution from $e^+e^- \to \tau^+\tau^- \to (e,\mu)^{\pm}$ + hadrons (from ref. 17).

Fig. 15: Angular distribution $\frac{1}{N}\frac{d\sigma}{d\cos\theta_\mu}$ from the process $e^+e^- \to \chi'^+\chi'^- \to \mu^\pm + X$ and a heavy pair production process $e^+e^- \to L^+L^- \to \mu^\pm + X$

a) $\sqrt{s}$ = 35 GeV, $m_{\chi'^\pm} = m_{L^\pm}$ = 8 GeV

b) $\sqrt{s}$ = 100 GeV, $m_{\chi'^\pm} = m_{L^\pm}$ = 25 GeV

(from ref. 17).

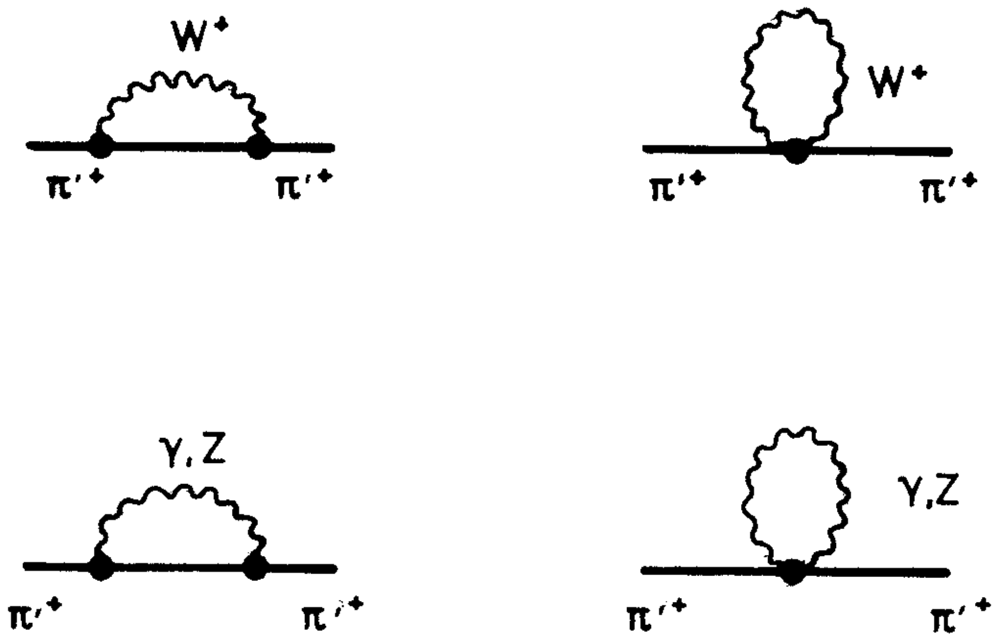

Fig. 1

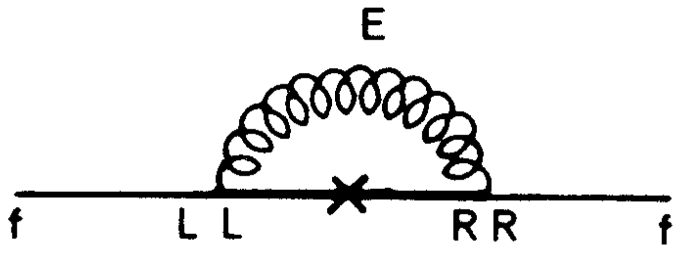

Fig. 2

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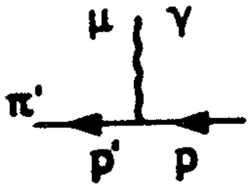
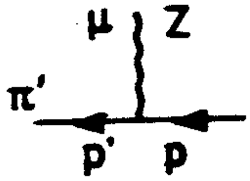
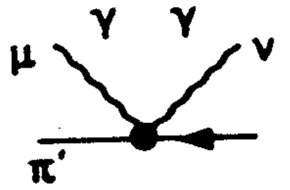


$$= \; i\,(p+p')_\mu \; eQ$$

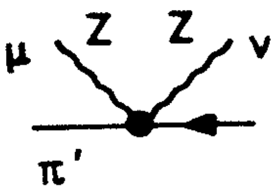


$$= \; i\,(p+p')_\mu \left(\frac{e}{\sin 2\theta_W}\right)(I_3 - 2\sin\theta_W Q)$$



$$= \; 2\,ig^{\mu\nu}\; e^2 Q^2$$

$$= \; 0$$



$$= \; -2\,ig^{\mu\nu}\frac{e^2 Q}{\cos^2\theta_W}(I_3 - \sin^2\theta_W Q)$$

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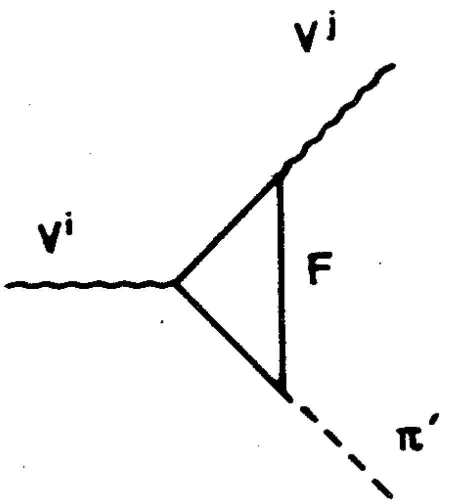
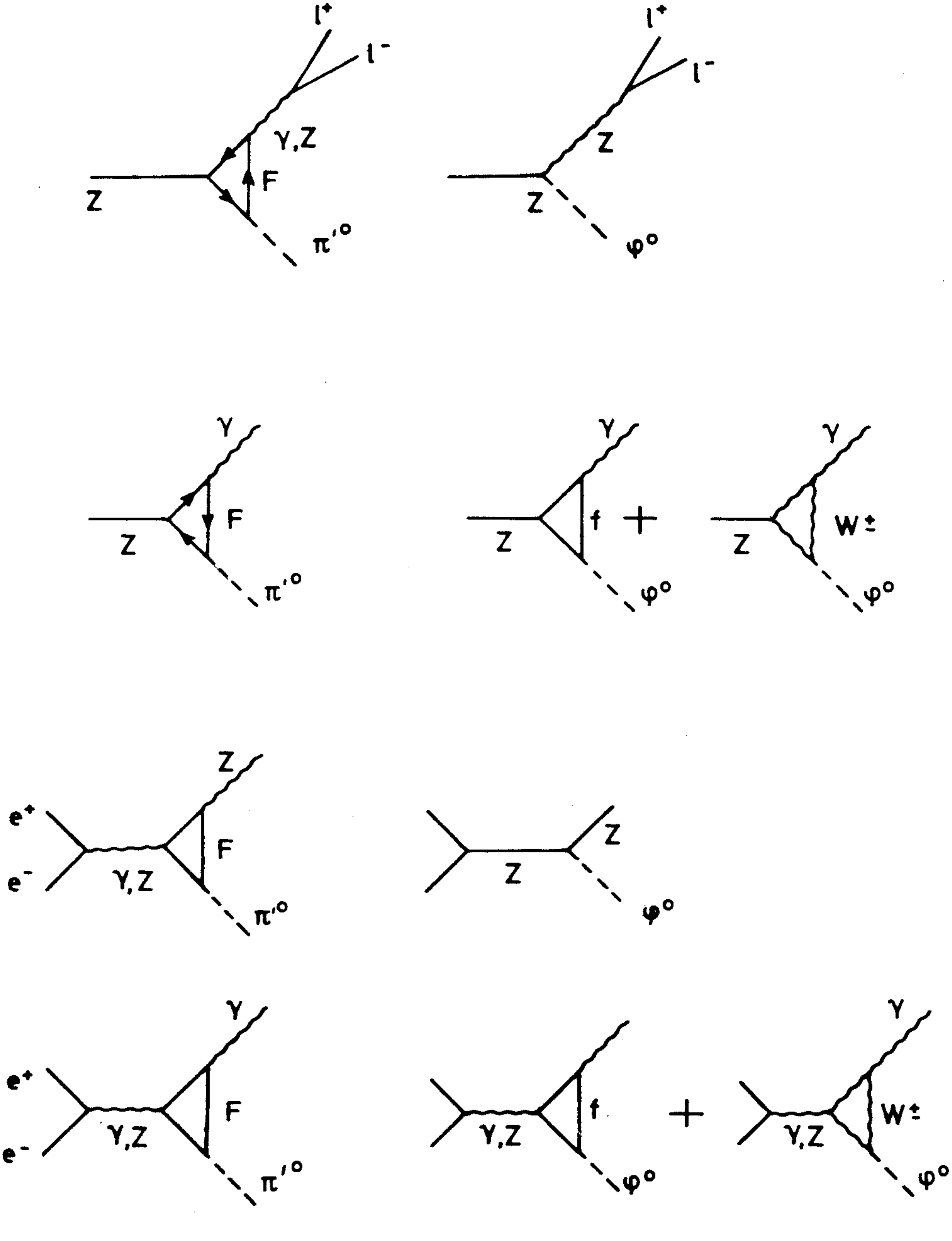
Fig. 3



$V^{i,j} = \gamma, W^{\pm}, Z$

Fig. 4



l⁺
l⁻
γ,Z
F
Z
π'⁰
l⁺
l⁻
Z
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W±
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Fig.5

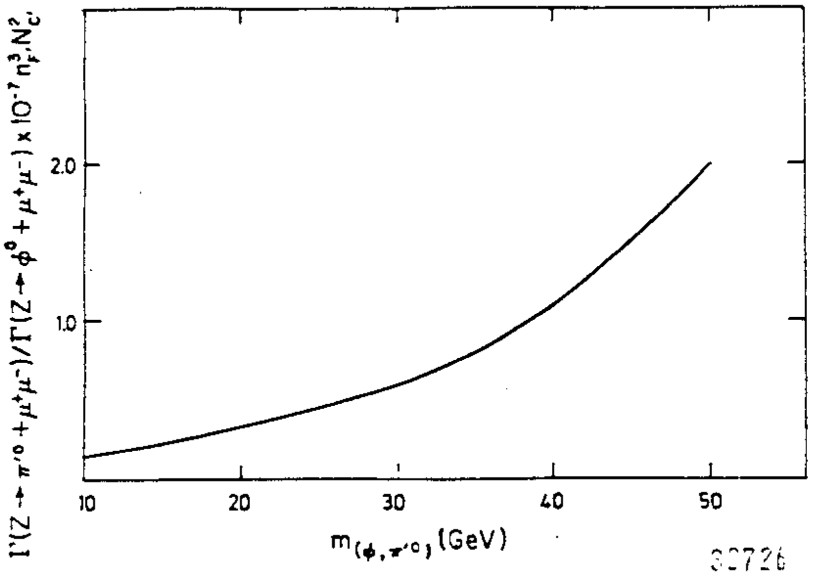

Fig.6

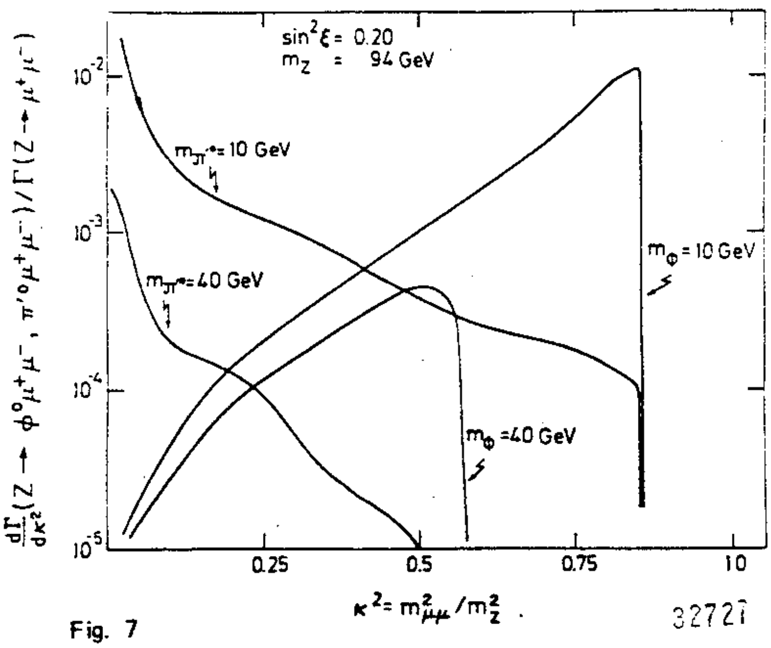

Fig. 7

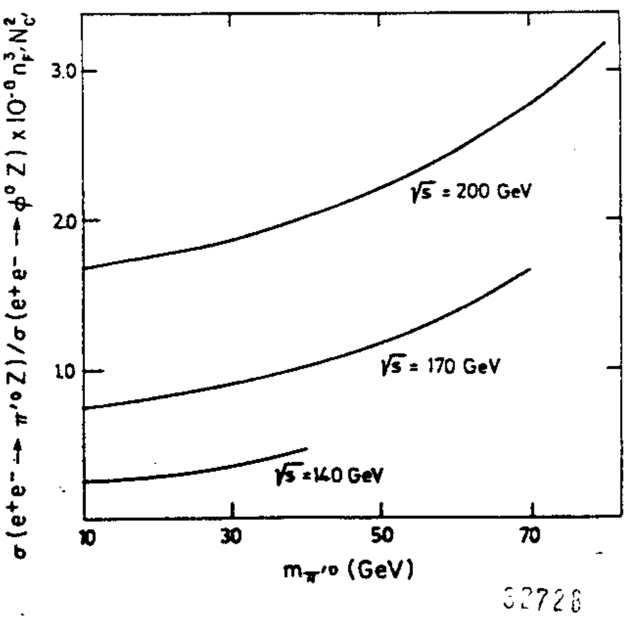

Fig.8

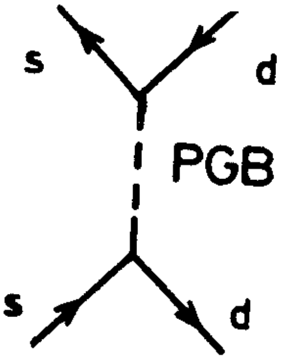

Fig. 9

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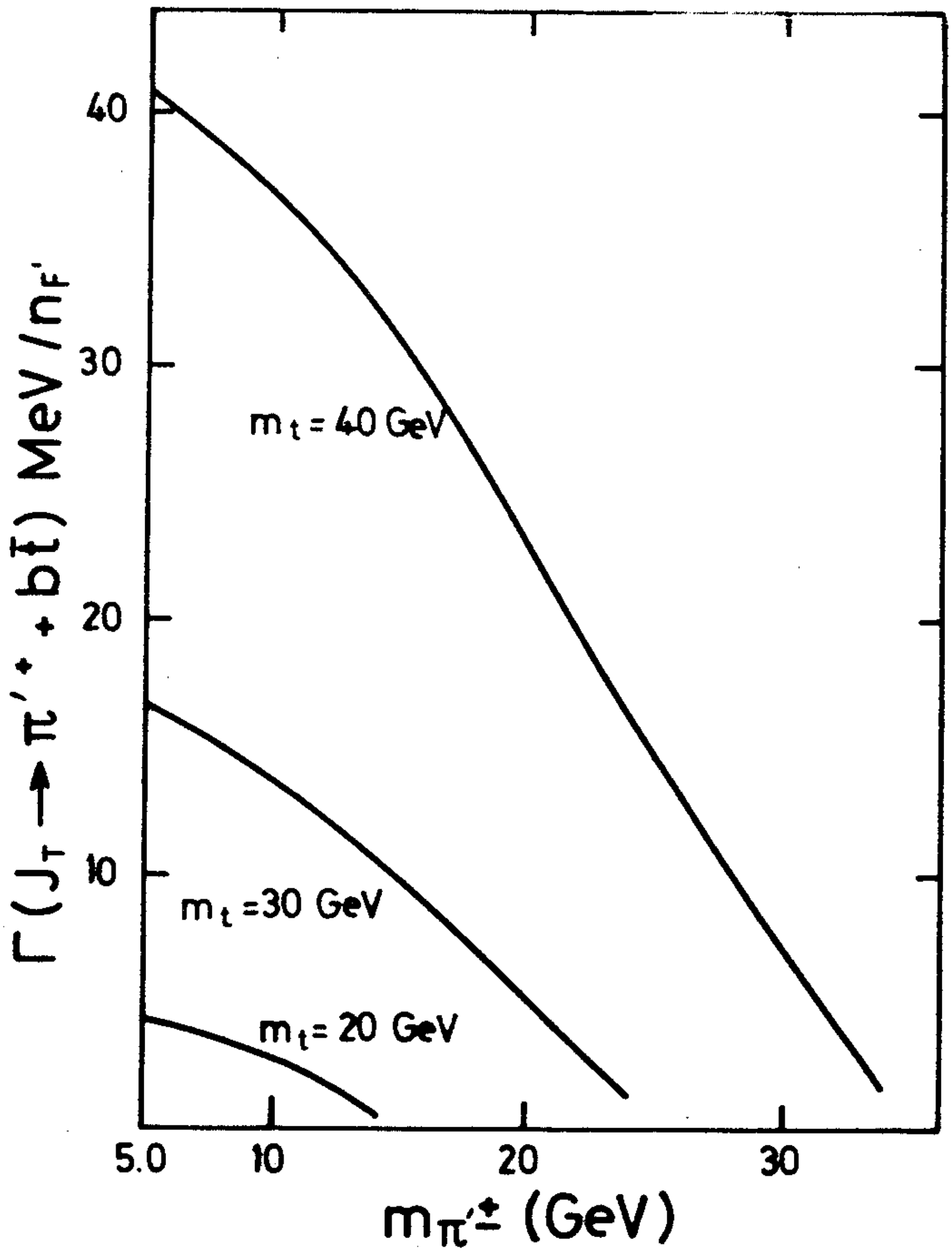

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Fig. 10

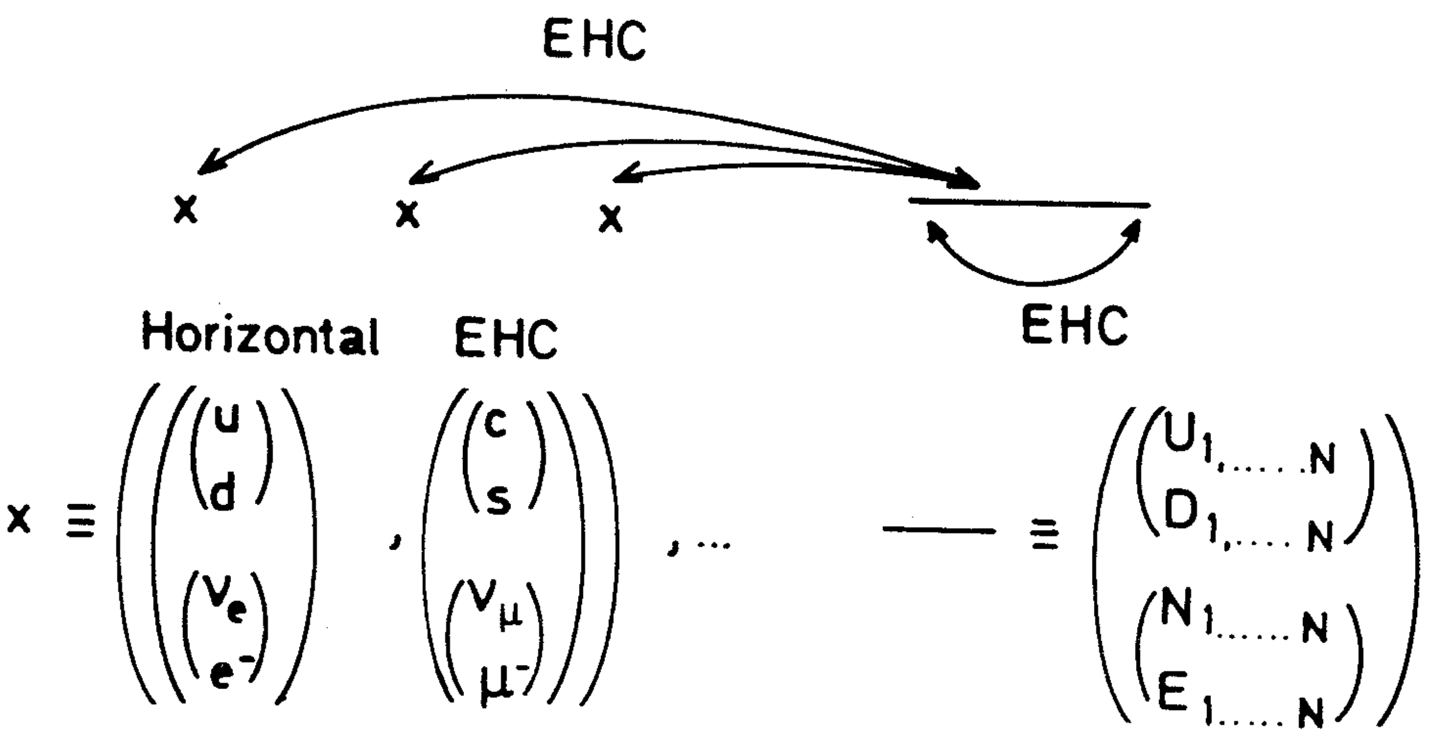

Fig. 11

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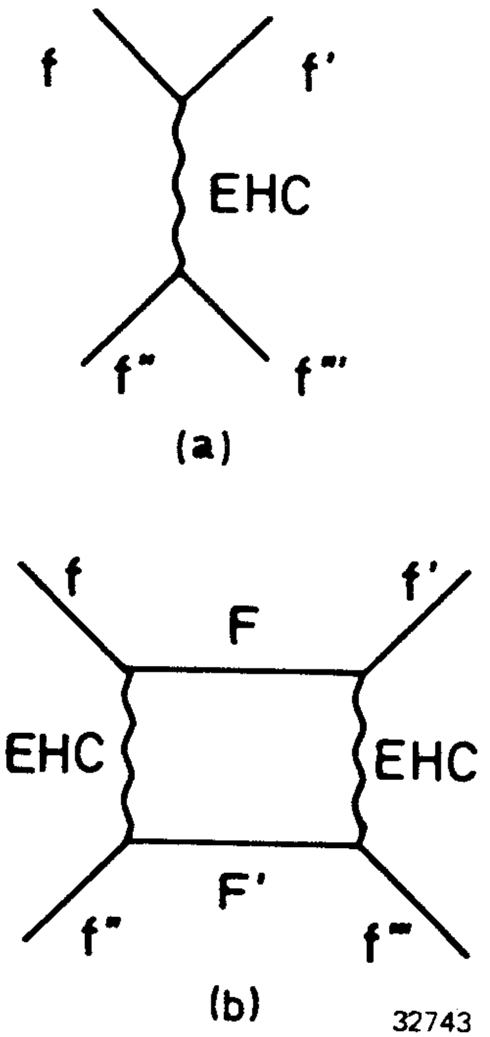

Fig. 12

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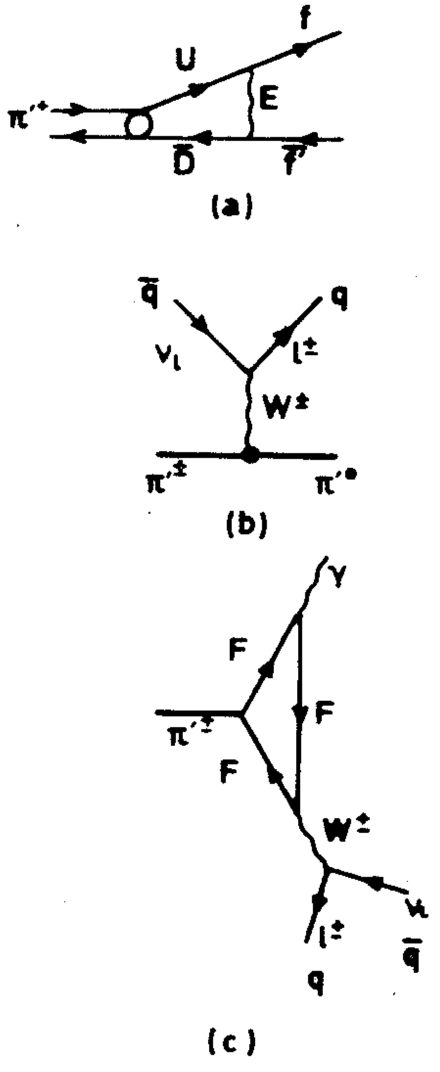
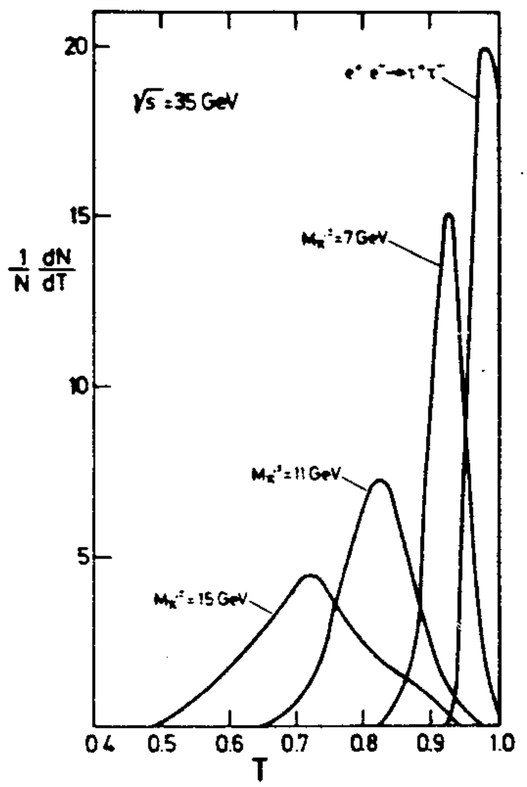


Fig. 13



Thrust Distribution

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Fig.14

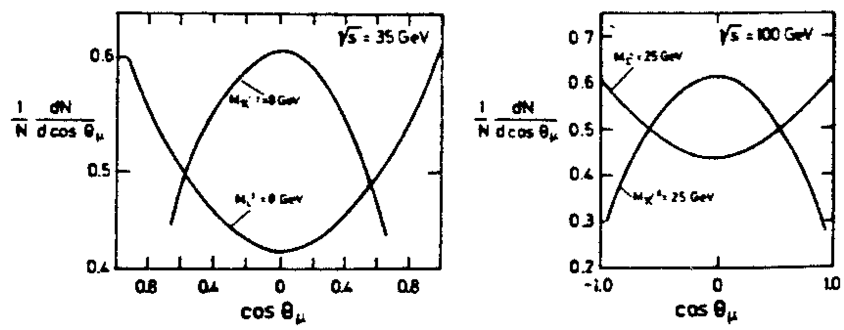

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Fig.15