TECHNICOLOUR - OASIS OR MIRAGE?

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1. INTRODUCTION

By now the standard SU(3) x SU(2) x U(1) model\(^1\) is the generally accepted theory of fundamental physics at energies up to about 100 GeV. The model seems to pose few remaining problems of principle, and most work on it now concentrates on fiddling technical questions, such as does QCD confine quarks and do the \(W^+\) and \(Z^0\) exist? It is not surprising that some adventurous (or foolhardy) souls are already riding off into the sunrise, looking for the deeper theory beyond the standard model. There is no consensus about which direction to ride in, nor even how far one must ride until the next landmark is reached. Perhaps there is a great desert to be crossed first, as advocated by the ascetic, minimalist devotees\(^2\) of grand unified theories (GUTs) who suggest that there is no new iteration scale below \(10^{15}\) GeV. As illustrated in Fig. 1 they extrapolate simple-mindedly the logarithmic variations of the known coupling constants until they come together in a GUT at a scale not far removed from the Planck energy scale of \(10^{19}\) GeV at which quantum gravity effects must become important. In making this enormous extrapolation they assume that there are no cases of new physics intervening in the great desert. In particular, they assume that the known apparently "elementary" fields - gauge bosons, quarks and leptons and Higgses - remain structureless and point-like down to distance scales of order \(10^{-29}\) cm or smaller.

Many physicists consider this grand scenario either a dangerous manifestation of hubris, technically unsound or boring. They replace this scenario with that illustrated in Fig. 2 where the great desert is riddled with cases due to the onset of new types of strong interactions such as techni-, extended techni-, meta-, hyper-, super- or heavy colour.
These theories are certainly less boring at low energies than are the minimal GUTs of Fig. 1. They share the general feature that some or all of the presently "elementary" fields will appear to have structure and be composite on an energy scale of one to 160 TeV, (i.e., $\geq 10^{-19}$ cm). This compositeness is postulated on various grounds, a common motivation being the mad proliferation of bosons, quarks and leptons which seem too numerous to be truly fundamental. There are other motivations in particular for the compositeness of scalar (Higgs) fields: elementary scalars have the unfortunate property that their $(mass)^2$ renormalization

$$S m_s^2 \sim \Lambda^2 \text{ (cut-off)}$$

If this correction to the $(mass)^2$ is to obey the "naturalness" condition that

$$|S m_s^2| \leq |m_s^2| \leq O(1) \text{ TeV}^2$$

for Higgses associated with the SU(2) $\times$ U(1) Weinberg-Salam model, then it must be that the cut-off $\Lambda < O(1)$ TeV, and this cut-off may well be provided by the scale of compositeness of scalar fields. Because the renormalization of a fermion mass

$$|S m_f| \sim \alpha m_f \ln (\Lambda m_f)$$

is only logarithmic, the "naturalness" requirement for fermion masses is much less stringent, requiring only that

$$\Lambda \leq m_f \exp (\frac{1}{\alpha}) \sim O(\text{Planck mass})?$$

It is therefore reasonable to start exploring models where Higgses alone are composite on a scale of $O(1)$ TeV.
In this talk I review scenarios for new strong interactions popu-
lating the desert, concentrating on models exhibiting this "precocious
compositeness" of scalar fields\textsuperscript{4}. I also discuss new ideas\textsuperscript{5} for pro-
tecting scalar masses by using new strong interactions which are super-
symmetric so that the quadratic divergences of type (1) actually cancel.
Also mentioned is the proposal\textsuperscript{6} that weak SU(2) interactions may in fact
be strong, with the low energy four-fermion interactions being mediated
by composite gauge bosons. Models of composite quarks and leptons are
left to the talk in these proceedings by Harari\textsuperscript{7}.

Most of this talk concerns composite Higgs fields, starting off
with the elementary ideas of dynamical symmetry breaking and techni-
colour\textsuperscript{4}, and then discussing the need to extend the technicolour\textsuperscript{8} in-
teractions in order to give masses to fermions. Some technical problems
arising in determining the pattern of dynamical symmetry breaking and
the criterion of the "most attractive channel" (MAC)\textsuperscript{9} will be touched
upon but not elaborated. Most of the talk emphasizes phenomenological
challenges for technicolour and extended technicolour theories. These
arise from two sources: the dangerously large magnitude of flavour-
changing neutral interactions in such models\textsuperscript{10} and the proliferation of
light spin-zero bosons\textsuperscript{11} which have not yet been observed. These phe-
nomenological difficulties have been among the motivations for the
recently proposed\textsuperscript{5} supersymmetric technicolour or supercolour theories.
These theories also have interesting phenomenological implications at
low energies: for example there should exist spin 1/2 supersymmetric
partners of the gluons, called gluinos, which may have masses of a few
GeV. There are also startling experimental implications of the "strong
weak interactions" idea\textsuperscript{6}: the $W^+$ and $Z^0$ gauge bosons would be consider-
dably heavier than in the standard model.

Certainly these new strong interaction theories generate plenty
of interesting predictions testable at present-day accelerators!
2. HOW HIGGS Bosons WORK – AND WHY THEY SHOULD NOT

Before seeing how we can replace elementary Higgs fields by composite spin-zero bosons which are bound states of fermions with a new type of strong interaction, we should first remind ourselves how elementary Higgses work, for example in the minimal Weinberg-Salam model. There one has a complex doublet of fields \( (\phi^+, \phi^-) \) and its complex conjugate \( (\bar{\phi}^0) \). Spontaneous symmetry breaking is achieved via an expectation value for the symmetric combination of neutral fields

\[
\langle 0 | \phi^0 | 0 \rangle = \frac{v}{\sqrt{2}} = \langle 0 | \bar{\phi}^0 | 0 \rangle \neq 0
\]

When this happens, three spin-zero fields remain massless:

\[
(\phi^+, \frac{1}{\sqrt{2}} (\phi^- - \bar{\phi}^0), \phi^-)
\]

and one combination is massive, the physical Higgs

\[
H = \frac{1}{\sqrt{2}} (\phi^+ + \bar{\phi}^0) - \nu^-
\]

The three weak gauge bosons \( W^\pm \) and \( Z^0 \) acquire masses by eating the three fields (6) and exploiting them as their third longitudinal helicity states. This mechanism for mass generation is illustrated in Fig. 3: the gauge-covariant derivative \( D_\mu \phi \) provides the direct \( W-\phi \) coupling \( i p_\mu (g_\nu/2) \) illustrated in Fig. 3a. A propagating \( W \) can then emit and reabsorb \( \phi \) fields an indefinite number of times, as in Fig. 3b. The sum of all these diagrams is a geometric series:

\[
\frac{1}{q^2} + \frac{1}{q^2} \left( \frac{g_\nu}{2} \right)^2 q^2 \left( \frac{g_\nu}{2} q^2 \right) \frac{1}{q^2} + \left( \frac{g_\nu}{2} \right)^3 (q^2)^3 + \cdots
\]

\[
= \frac{\frac{1}{q^2}}{1 - \frac{g_\nu^2 q^2}{4 q^2}} = \frac{1}{q^2} m_W^2 \quad \text{where} \quad m_W = \frac{g_\nu}{2}
\]
Mass generation for the $Z^0$ works similarly apart from the minor complication of mixing between the neutral member of the SU(2) triplet of gauge bosons and the U(1) boson. With three out of the four $\phi$ degrees of freedom eaten in this way, only the physical Higgs $H$ (7) remains as an observable particle whose couplings are fixed:

$$g_{H\phi\phi} = \frac{m_H}{\nu} = 2^{\frac{3}{2}} G_F \sqrt{2} m_H$$

(9)

but whose mass is almost arbitrary, except that it should be less than about 1 TeV if it is not to be strongly interacting$^{12}$.

Since spontaneously broken gauge theories are renormalizable, everything looks fine and dandy, except for an embarrassing technicality$^{3,4}$. Higher order loop corrections to $m_H^2$ and $m_W^2$ are quadratically divergent because of diagrams like those in Fig. 4 which yield a result of the form (1). People find it "unnatural" that the corrections to $m_H$ should be larger than the physical $m_H$ itself, which we saw earlier should be $< 0(1)$ TeV. This suggests that we should seek a cut-off $\Lambda = O(1)$ TeV.

Two ways of finding one have been pursued: one is to go to a supersymmetric theory$^5$ in which such diagrams as Fig. 4 are fated to cancel - we will return later to this possibility. The line we shall follow for the time being is that of "dissolving" the diagrams by making the Higgses composite$^4$ on a scale of 1 TeV.

3. THE BASIC TECHNICOLOUR IDEA

We postulate$^4$ a new set of non-Abelian gauge interactions based on a group such as SU$(N)_{TC}$ or SO$(N)_{TC}$ which are asymptotically free and become strong at an energy scale $\Lambda_{TC} \leq O(1)$ TeV. We further postulate$^4$ some massless technifermions which "feel" and are confined by these technicolour interactions: the simplest example would be

$$\begin{pmatrix} U_L \\ D_L \end{pmatrix}_{1, \ldots, N}, \begin{pmatrix} U_R \end{pmatrix}_{1, \ldots, N}, \begin{pmatrix} D_R \end{pmatrix}_{1, \ldots, N}$$

(10)
It is then expected that the vacuum will contain a condensate of technifermions in the same way as the QCD vacuum contains a condensate of quarks:

\[
\langle 0 | \overline{F}_R F_L | 0 \rangle = \langle 0 | \overline{F}_L F_R | 0 \rangle = O(\Lambda_{QCD}^3) \quad \text{cf} \quad \langle 0 | \overline{F}_R q_L | 0 \rangle = O(\Lambda_{QCD}^3) \quad (F = u, d)
\]

corresponding to a dynamical breakdown of chiral symmetry: \( SU(2)_L \times SU(2)_R \rightarrow SU(2)_{L+R} \) in the case of the two techniflavour model (10).

The spectrum of the technicolour theory will then contain massless spin 0 Goldstone bosons \( \pi_T \), one for each spontaneously broken generator, of which there are three in the toy model (10):

\[
(\pi_T^+, \pi_T^0, \pi_T^-)
\]

These are directly analogous to the pions of QCD which are believed to be light because they are almost the Goldstone bosons of an analogously broken chiral symmetry, having

\[
m_{\pi_T}^2 = O(m_u, d, \Lambda_{QCD}) + (\text{radiative corrections})
\]

so that they would be strictly massless if the quark masses were zero and non-strong interactions were switched off.

We are now in a position to replay the Higgs trick of Section 2, this time with composite fields. The \((\pi_T^+, \pi_T^0, \pi_T^-)\) (12) can be eaten by the \( W^+ \) and \( Z^0 \) in the same way as the \((\phi^+, 1/\sqrt{2} (\phi^0 - \phi^-)), \phi^-)\) (6) of Section 2 because there is a \( W^- \pi_T \) coupling, directly analogous to the \( W^- \pi \) coupling \( f_{\pi} \) which allows the usual \( \pi^+ \) to decay weakly:

\[
iq_{\mu} (\frac{gF_{\pi}}{2}) \quad \text{cf} \quad iq_{\mu} (\frac{gf_{\pi}}{2})
\]

The coupling (14) replaces the \( W^- \phi \) coupling of Fig. 3a, and generates a geometric series of diagrams (cf. Fig. 3b) which give the \( W^+ \) masses just like Eq. (8):
\[ m_w = \frac{g F_\pi}{2} \]  

To get \( m_w \approx 80 \text{ GeV} \) as required, the quantity \( F_\pi \) must \( = 250 \text{ GeV} \), compared with the conventional QCD \( f_\pi \) value of 93 MeV. Since one expects

\[ \frac{F_\pi}{\Lambda_{\text{TC}}} = O\left(\frac{f_\pi}{\Lambda_{\text{QCD}}}\right) \]  

(16)

to within minor numerical factors, this specifies \( \Lambda_{\text{TC}} \) to be \( O(1) \) TeV.

Developing further the analogy with the elementary Higgs model, the condensate (11) replaces the Higgs vacuum expectation value (5), while corresponding to the physical Higgs (7) one expects to find a massive \( O(1) \) TeV scalar technimeson, the \( e_T \) which is an \( FF \) bound state analogous to the familiar \( e \) (500 or 600 or 700 or 1000 or 1300 or \( ? \) MeV) meson of QCD. This scheme looks relatively simple and elegant, but there are problems.

4. FERMION MASSES AND THE NEED TO EXTEND TECHNICOLOUR

Quarks and leptons have masses, and the way they acquire them in the elementary Higgs theory\(^1\)\(^2\) is by coupling to Higgs fields:

\[ g_f \bar{f}_R \phi \bar{f}_L + \text{(herm. conj.)} \]  

(17)

as illustrated in Fig. 5a. Then when the Higgs acquires a vacuum expectation value (5) the fermions acquire masses

\[ m_f = g_f \frac{v}{\sqrt{2}} \]  

(18)

In a composite Higgs model, our knee-jerk reaction would be to couple the fermions to our composite Higgs \( e_T \):

\[ g_e \bar{f}_R f_L e_T + \text{(herm. conj.)} \]  

(19)
as suggested in Fig. 5b. But since the \( e \) is an \( \bar{F}F \) composite, the coupling (19) involves a four-fermion interaction (Fig. 5c):

\[
g_e(\bar{f}f)(\bar{F}F)\tag{20}
\]

which is necessarily dimensional: \([g_e] = M^{-2}\). Such a four-fermion coupling is non-renormalizable and hence illegal, unless we can generate it indirectly. This can be done\(^9\) by the exchange of a heavy gauge vector boson \( E \) (Fig. 5d) in the same way as we now believe the Fermi four-fermion weak interaction is generated by \( W^\pm \) exchange:

\[
g_e = \frac{g_E^2}{m_E^2} \quad \quad c_E = g_F^2 \quad \quad \frac{g_F^2}{8m_W^2} \tag{21}
\]

Since these new interactions couple technifermions to ordinary fermions they are called "extended technicolour" or ETC. That they give fermion masses can be seen by substituting the condensate value (11) into the four-fermion interaction (20):

\[
m_f = g_e \langle 0 | \bar{F}F | 0 \rangle = \frac{g_E^2}{m_E^2} \langle 0 | \bar{F}F | 0 \rangle = \frac{g_E^2}{m_E^2} O(N_{TC}^3) \tag{22}
\]

as illustrated in Fig. 5e. Notice that because different fermions have a wide range of different masses, from \( 1/2 \) MeV for the electron to \( \geq 20 \) GeV for the top quark, there must be a large number of different ETC bosons \( E_\ell \) with different masses:

\[
m_{E_\ell}^2/g_E^2 = \frac{O(N_{TC}^3)}{m_f} \tag{23}
\]

Since we know the scale of \( \langle 0 | \bar{F}F | 0 \rangle \) from the \( W^\pm \) masses, and the \( m_f \) from experiment, we can estimate the masses of the bosons \( E_\ell \). The picture can be complicated by mixing, but the simple ansatz (23) gives some
insight into the mass scales involved, as displayed in the Table\textsuperscript{10} below for quark masses assuming $g_{Fq}^2 = 0(1)$.

<table>
<thead>
<tr>
<th>Quark</th>
<th>$u$</th>
<th>$d$</th>
<th>$s$</th>
<th>$c$</th>
<th>$b$</th>
<th>$t$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Assumed mass</td>
<td>6 MeV</td>
<td>10 MeV</td>
<td>200 MeV</td>
<td>1.5 GeV</td>
<td>5 GeV</td>
<td>30 GeV</td>
</tr>
<tr>
<td>$m_{Fq}^2$</td>
<td>2200 TeV\textsuperscript{2}</td>
<td>1200 TeV\textsuperscript{2}</td>
<td>60 TeV\textsuperscript{2}</td>
<td>8 TeV\textsuperscript{2}</td>
<td>$\frac{2}{2}$ TeV\textsuperscript{2}</td>
<td>$\frac{1}{2}$ TeV\textsuperscript{2}</td>
</tr>
</tbody>
</table>

Table: ETC masses

In a simple gauge group, the existence of gauge bosons coupling $f_1$ to $F$ and $F$ to $f_2$ will entail\textsuperscript{8,10} the existence of bosons coupling $f_1$ to $f_2$ with similar masses. If $f_1$ and $f_2$ are quarks or leptons of the same charge from different generations, these new gauge interactions $f_1 f_2$ are of the "horizontal" type, and so we\textsuperscript{10} call them "horizontal ETC". They are unavoidable unless one multiplies the sets of technifermions, and will subsequently cause serious problems with flavour changing neutral interactions.

5. WHERE DO THE ETC MASSES COME FROM?

It seems that we need a big ETC group containing bosons with a sequence of different masses suggestive of successive stages of spontaneous symmetry breaking called "tumbling"\textsuperscript{9}:

$$
\varphi_{\text{ETC}} \rightarrow \varphi_{\text{ETC}}' \rightarrow \varphi_{\text{ETC}}'' \rightarrow \cdots \rightarrow \varphi_{\text{TC}}
$$

In line with our philosophy of replacing elementary Higgses by fermion composites, we would like each stage of symmetry breaking in the sequence (24) to be caused by condensates of fermions $\langle 0 | \bar{F}_R P_L | 0 \rangle \neq 0$. 
Since we are in the business of breaking $G_{ETC}$, this condensate must have a non-trivial $G_{ETC}$ representation content, and to get that we must have different representations $R_{ETC}$ for the left- and right-handed fermion fields

$$R_{ETC}(F_R) \neq R_{ETC}(F_L)$$

(25)

whereas their TC representation content should be identical [cf. Eq. (10)] as we do not want to break technicolour itself. We need a criterion for deciding the sequence of channels through which tumbling occurs, i.e., which fermion condensates $\bar{F}_RF_L$ form in which sequence. It has been proposed that the appropriate criterion is that of the "Most Attractive Channel" (MAC) computed by determining the channel in which the one-gluon exchange (OGE) potential is most attractive. This requires maximizing

$$C_i + C_j - C_{ij}$$

(26)

where $C_i$ and $C_j$ are the Casimirs of the fermions $F_i$ and $F_j$ which are candidates for condensation, and $C_{ij}$ is the Casimir of a representation in the product $\bar{F}_iF_j$. When this channel condenses one expects to find a corresponding mass for the gauge bosons breaking ETC $\rightarrow$ E'TC, etc.: 

$$\langle 0 | \bar{F}_i F_j | 0 \rangle = O(\Lambda_{ij}^3) \Rightarrow m_{E_{ij}} = O(\Lambda_{ij})$$

(27)

The procedure is then repeated for E'TC $\rightarrow$ E"TC, etc., until at last the exact symmetry group of technicolour is reached. In this way all masses scales: $m_{E'}$, $\Lambda_{TC}$, $m_W$ and $m_\ell$ could perhaps be generated dynamically.

Unfortunately, there are technical questions about the validity of the MAC criterion. For one thing, one expects gauge boson condensates as in QCD:
\[ \langle 0 | F_\mu^i F_\nu^j | 0 \rangle = O(\Lambda^{4}_{QCD}) \]  

which superficially look very attractive. Perhaps they condense before the fermions and screen the interactions so that the MAC criterion is inapplicable? Another problem\(^{15}\) is that condensation of a certain \(F_i F_j\) combination may give rise to masses for (and hence screen) some gauge bosons crucial to the condensation of the self-same \(F_i F_j\) channel. This may make another channel more attractive, or may imply that several different stages of symmetry breaking occur at essentially the same scales \(\Lambda_{ij}\). It was never clear anyway that the MAC criterion could give energy scales \(\Lambda_{ij}\) sufficiently widely separated to accommodate the range of different \(m_{E_q}\) needed to fit the \(m_q\) in the Table. Another problem\(^{16}\) is that a naive application of the MAC criterion (26) often suggests that channels with Lorentz vector quantum numbers might condense before Lorentz scalar channels, probably implying an unacceptable violation of Lorentz invariance.

These questions about the MAC criterion have been investigated\(^{17}\) in two-dimensional models, with the conclusion that there the MAC criterion fails, but for reasons specific to two dimensions. Sometimes the MAC criterion would lead to a pattern of symmetry breaking involving the generation of massless Goldstone bosons, which are forbidden in two dimensions\(^{18}\), but not of course in four. Even if the MAC criterion can avoid this pitfall, often it does not respect the distinction between left-moving and right-moving fermions specific to two dimensions. It is worrying that the MAC criterion fails in models where OGLE confines, whereas OGLE does not in four dimensions. But the two-dimensional models may well be unrealistic guides. In any case, we do not need specifics of tumbling or MAC in the rest of this talk, just the general features of the ETC scenario shown in Fig. 6.
6. PSEUDO-GOLDSTONE BOSONS

The generic model of Fig. 6 contains a complete technigeneration\textsuperscript{11}. If we tried to get by with any fewer technifermions, we would need ETC bosons with non-trivial SU(3) \times SU(2)_{L} \times U(1) transformation properties. We would then need to unify ETC and SU(3) \times SU(2) \times U(1) interactions at some energy scale \( \lesssim \) \( 10 \) to \( 100 \) TeV, which seems difficult since the ETC interactions are strong and the others relatively weak. We do not assume more than one technigeneration because this would complicate life, in particular with many more light spin-zero bosons, which give severe problems even in the minimal scenario as we will soon see. Also in Fig. 6 we assume that all technifermions including the technineutrino \( N \) come in both left and right helicity states. Here I will finesse the problem of neutrino masses, which has been faced up to by Holdom\textsuperscript{19}.

The generic model of Fig. 6 contains 64 spin-zero bosons with the quantum numbers of \( \bar{F}_{i} F_{j} \) where \( F_{i} \) and \( F_{j} \) can be any of the eight techni-flavours

\[
\left( U_{R,Y,B} ; D_{R,Y,B} ; N ; E \right)
\]  

(29)

These correspond to the spontaneous breakdown of chiral

\[
SU(8)_{L} \times SU(8)_{R} \times U(1)_{L-R} \Rightarrow SU(8)_{L+R}
\]  

(30)

and are analogous to the conventional pseudoscalars \( \pi, K, D, \eta, \ldots \) which would be massless if the quarks were massless and there were no \( SU(2) \times U(1) \) interactions. There would be even more spin-zero bosons corresponding to \( F_{i} F_{j} \) channels in theories with \( SO(N)_{TC} \). Of the 64 \( \bar{F}_{i} F_{j} \) mesons, the one corresponding to \( U(1)_{L-R} \) gets a mass of order 1 TeV from the same non-perturbative mechanism that gives the \( \eta' \) a mass of order 1 GeV in QCD. Of the remaining 63 bosons, three are massless and
"eaten" by the $W^\pm$ and $Z^0$ to give them masses as in Section 3, while 60 are left as physical "pseudo-Goldstone bosons" (PGBs) to be observed experimentally. The spectrum of PGBs is shown in Fig. 7.

They are called "pseudo" because none of them are strictly massless, all the Goldstone symmetries to which they correspond being broken by some type of non-TC interaction. In fact, in the simple model of Fig. 6 the ETC interactions do not contribute to the masses of the colour octets and singlets, but they do contribute to the leptoquark $P_{LO}$ masses:

$$S_{P_{LO}}^2 \approx O(50) \text{GeV}^2$$

(31)

Colour $SU(3)$ interactions contribute to the masses of the colour non-singlet PGBs

$$S_{P_6} \sim 260 \text{GeV}, \quad S_{P_8} \sim 250 \text{GeV}, \quad S_{P_3} \sim 160 \text{GeV}$$

(32)

Weak and electromagnetic interactions give relatively small contributions to the PGB masses: they are most interesting for the charged colour singlets $P^\pm$

$$m_{P^\pm} \approx \begin{cases} 5 \text{ to } 8 \text{ GeV} & \text{for } SU(N)_{TC} \\ 8 \text{ to } 14 \text{ GeV} & \text{for } SO(N)_{TC} \end{cases}$$

(33)

So far we still do not have masses for the neutral colour singlets $P^0$ and $P^3$, and to give them masses we need extra interactions coupling quarks directly to leptons. The simplest possibility is the vertical $SU(4)$ of (colour + lepton number) introduced by Pati and Salam (PS).

As shown in Fig. 8, PS gauge bosons would mediate $K^0 \rightarrow \mu e$ decay and to be compatible with the upper limit on this mode we need$^{22,23}$
\[ m_{p_0,3} \geq 300 \text{ TeV} \] (34)

The PS bosons then contribute \( P^0 \) and \( P^3 \) masses\(^{20,22-24}\)

\[ m_{p_{0,3}} \approx \frac{\alpha (\Lambda_T^2)}{m_{P_S}} \times \frac{1}{2} \text{ GeV} \times \left( \frac{300 \text{ TeV}}{m_{P_S}} \right) \times 2^{0.1} \leq 3 \text{ GeV} \] (35)

if we assume that \( SU(4)_{PS} \) commutes with ETC, as seems most plausible. The resulting masses of all the PGBs are displayed in Fig. 7. Generically one expects their couplings to fermions to be similar in magnitude to those of Higgs bosons [cf, Eq. (18)].

\[ g_{P_S \bar{f}} = O\left( \frac{\alpha^2}{m_T^2} \right) = O\left( \frac{m_0}{\Lambda_T} \right) = O\left( \frac{m_{\mu}}{m_\nu} \right) = O\left( \frac{m_f}{v} \right) \] (36)

We will return later to more specific models\(^{25}\) of these couplings and their phenomenological implications.

7. FLAVOUR-CHANGING NEUTRAL INTERACTIONS

These are strongly suppressed in the standard Glashow-Weinberg-Salam model\(^1\) as a result of some well-understood mechanisms\(^{26}\). When all fermions of the same charge have the same transformation properties under the weak gauge group [e.g., left-handed fields all \( SU(2) \) doublets, right-handed fields all \( SU(2) \) singlets] then the GIM mechanism\(^{27}\) guarantees that the gauge boson exchanges automatically\(^{26}\) conserve flavour to a very good approximation:

\[ \frac{A(\Delta F \neq 0)}{A(\Delta F = 0)} = O(\alpha_F m_T^2) = O(10^{-5}) \] (37)
Similarly, in models with relatively simple Higgs structures in which all the fermions of a given charge get their masses from the same Higgs field, the couplings of the physical neutral Higgs also conserve flavour to the very good approximation (37). This condition is fulfilled a fortiori in the minimal SU(2) × U(1) model of Section 2 in which there is just one complex Higgs doublet.

Flavour-changing neutral interactions arising through gauge bosons and PGB exchanges pose challenges for ETC theories on two fronts. It is evident from Fig. 6 that different conventional fermions of the same charge cannot have identical ETC transformation properties, and we already argued at the end of Section 4 that ETC theories necessarily contain flavour-changing horizontal HETC gauge bosons coupling e.g., d to s or u to c with essentially maximal strength. The magnitudes of these exchanges are

\[ A(\Delta F = 0) = O\left(\frac{g^2}{m_E^2}\right) \]  

(38)

which we can estimate on the basis of the Table and compare with experimental limits. We find that for the \( \Delta S = 2 \) transition operator responsible for \( K^0 - \bar{K}^0 \) mixing

\[ \text{Re} A(\Delta S = 2) \big|_{\text{Exp}} = O(10^{-12}) \text{ GeV}^{-2} \]  

(39a)

\[ |A(\Delta S = 2) |_{\text{HETC}} = O(10^{-9}) \text{ GeV}^{-2} \]  

(39b)

\[ \text{Im} A(\Delta S = 2) \big|_{\text{Exp}} = O(10^{-15}) \text{ GeV}^{-2} \]  

(39c)

while for the \( \Delta C = 2 \) transition operator which could give \( D^0 - \bar{D}^0 \) mixing which has not yet been observed:
\[ |A(\Delta C = 2)|_{\text{exp}} \leq O(3 \times 10^{-9}) \text{GeV}^{-2} \]  
\[ |A(\Delta C = 2)|_{\text{HETC}} \cong O(3 \times 10^{-9}) \text{GeV}^{-2} \]

Comparing (39a) and (39b) we find a discrepancy by three orders of magnitude, while the imaginary part (39c) is off by six orders of magnitude. There is even a two order of magnitude problem with the \( \Delta C = 2 \) amplitudes (40). In various other cases flavour-changing ETC effects graze the experimental limits\(^{28}\), but these are the clearest examples of serious difficulties\(^{29}\).

At this point one may sense some of the ETC rats leaving the sinking ship, but they may be too hasty. Pierre Sikivie and I\(^{30}\) analyzed a simple model in which SU(8)\(_{\text{ETC}} \rightarrow SO(5)_{\text{TC}}\) and one could obtain quasi-realistic quark masses and Cabibbo mixing. In that model we found

\[ |A(\Delta S = 2)| = O(10^{-9}) \text{GeV}^{-2}, \quad |A(\Delta C = 2)| = O(10^{-9}) \text{GeV}^{-2} \]

The \( \Delta S = 2 \) amplitude behaved as expected, but the \( \Delta C = 2 \) amplitude was suppressed by unanticipated cancellations down to an acceptable level. It may be possible to find a model extending this trick to \( \Delta S = 2 \) also, using an ETC version of the old "schizion" alternative\(^{31}\) to the GIM mechanism, in which "generation-raising" and "generation-lowering" bosons do not like to mix, though we have not been able to find such a model. But it is not yet obvious to us that the HETC flavour-changing neutral interactions are necessarily fatal.

PGB exchanges are potentially troublesome\(^{10}\) because although their couplings (36) are relatively small, so also are some of their masses. If we suppose that the flavour-changing PGB couplings to fermions \( f_1 \) and \( f_2 \) are of order
\[ g_{PfF^*_1} = O\left( \frac{g_{m_Pf}}{m_w} \right) \times \Theta_{12} \]  

(42)

for some mixing angle \( \Theta_{12} \), then we have the following constraints from the \( K^0 - \bar{K}^0 \) system:

\[ \Theta_{sd} \leq \frac{m_{Pf}}{6000 \text{ GeV}} \]  

(43a)

and from the \( D^0 - \bar{D}^0 \) system:

\[ \Theta_{uc} \leq \frac{m_{Pf}}{2000 \text{ GeV}} \]  

(43b)

It is clear that the couplings of the lightest PGBs \( P^0 \) and \( P^3 \) must conserve flavour to very high accuracy. We have proposed\(^{25}\) a criterion for working this trick which is analogous to the requirement\(^{26}\) of a single Higgs coupling to all fermions of a given charge. Called "monophagy", it is the requirement that all fermions of a given charge get their mass from the same technifermion condensate. This means, for example, that one cannot have some of the charge 2/3 quarks' masses coming from \( <0|\bar{U}U|0> \) and some from \( <0|\bar{D}D|0> \):

\[ (m_u, m_c, m_t) \neq \left( \frac{\phi_{E}^2}{m_E} \right) \text{matric} \begin{pmatrix} <0|\bar{U}U|0> \\ <0|\bar{D}D|0> \end{pmatrix} \]  

(44)

It is then easy to show\(^{25}\) that the dominant \( O(g_{m_Pf}/m_w) \) couplings of the \( P^0 \) and \( P^3 \) conserve flavour, and are in fact severely constrained. If ETC commutes with \( SU(3) \times SU(2) \times U(1) \) there is a unique form that they can take:
\[
\frac{p^3}{F_\pi} \left[ (\bar{u} m_u^u \gamma_5 u + \bar{d} m_d^d \gamma_5 d) \sqrt{\frac{1}{3}} - \sqrt{\frac{1}{3}} \left( m_t^t \gamma_5 \tau \right) \right] \right]^{(45)}
\]

where \(m_u, d, \tau\) are the mass matrices for charge 2/3, charge -1/3 and charge -1 fermions respectively. However, in addition to the couplings (45) there are \(0(g^2 m_t^2/m_W^2)\) couplings which must change flavour, either for charge 2/3 quarks (u class models) or for charge -1/3 quarks (d class models), or both, even in monophagic theories. These flavour changing couplings are not sufficient to cause problems with the \(K^0 - \bar{K}^0\) and \(p^0 - \bar{p}^0\) systems, but they may make \(p^0\) and \(p^3\) conspicuous in heavy quark decays, as we will see in the next section.

3. PGB PHENOMENOLOGY

Since the PGBs are the lightest technicolour particles, and the only ones that can be produced by accelerators which now exist or are under construction, it is reasonable to concentrate on their phenomenology.

The colour octet PGBs are expected (Fig. 7) to have masses of (250 GeV), and will have a rather small production cross-sections even in high energy hadron-hadron colliding beam machines. Their dominant decay mode is expected to be \(t\bar{t}\), with \(gg\) coming in next. An experiment with very good dijet resolution may just be able to pick out an indistinct bump in \((t\bar{t})\) jet invariant masses (Fig. 9) but this looks fairly difficult.

The colour triplet leptoquarks \(P_{LQ}\) with masses 0(160) GeV can be produced in ep collisions by the mechanism shown in Fig. 10a. Here again the production cross-section is not very encouraging (Fig. 11), but at least there is a distinctive decay signature: \(P_{LQ} \rightarrow t\bar{t}\) indicated
in Fig. 10a, which may enable it to be detected at a machine like HERA. One might find an indirect signature for the existence of the $P_{LQ}$ from the decay $K^0 \rightarrow \mu e$, which can be mediated by crossed channel $P_{LQ}^\pm$ exchange as in Fig. 10b. Calculations suggest a rate close to the present experimental upper limit, but many other mechanisms exist for $K^0 \rightarrow \mu e$ decays in ETC theories, such as crossed channel PS exchange and direct channel HETC or PGB exchange, any of which might be near the present experimental limit. Thus the interpretation of any $K^0 \rightarrow \mu e$ decays seen might be delightfully ambiguous!

The only PGBs easily accessible to present accelerators are the colour singlets $P^\pm$, $P^0$ and $P^3$. We saw in Eq. (33) that the $P^\pm$ are sufficiently light\(^{11,21}\) (5 to 14 GeV) to be pair-produced in $e^+e^-$ annihilation at PETRA or PEP:

$$\frac{\sigma(e^+e^- \rightarrow \gamma^* \rightarrow P^+ P^-)}{\sigma(e^+e^- \rightarrow \gamma^* \rightarrow \mu^+ \mu^-)} = \frac{1}{4} \beta^3$$  \hspace{1cm} (46)

From the generic couplings \(^{36}\) one expects them to decay predominantly into heavy fermion pairs. Their dominant couplings are precisely determined in the monopagic model\(^{25}\) of Section 7:

$$\frac{P^+}{F_\pi} \left[ \frac{U_{\ell \ell}}{U_{\ell \ell}} \begin{array}{c} \mu u \\frac{m_u}{2} \end{array} - m_u U_{\ell \ell} \frac{1 - \gamma_5}{2} d \sqrt{3} \\
- \sqrt{6} \begin{array}{c} \ell \end{array} \begin{array}{c} m_{\ell} \frac{1 + \gamma_5}{2} \end{array} \right]$$  \hspace{1cm} (47)

where $U_{KM}$ is the generalized Cabibbo mixing matrix of Kobayashi and Maskawa\(^{35}\), and $m_u, m_\ell$ are the fermion mass matrices net in Eq. (45). The couplings (47) fix the decay branching ratios of the $P^\pm$:
\[ \Gamma(P^+ \to \tau\nu) : \Gamma(P^+ \to c\bar{s}) : \Gamma(P^+ \to c\bar{b}) \approx 3 : |C_1C_2C_3 - S_2S_3 e^{i\delta}|^2 \sim 10 |C_1S_2S_3 + S_2S_3 e^{i\delta}|^2 \leq 1? \] (48)

These predictions look rather incompatible with the JADE 90% confidence level limits\(^{36}\) shown in Fig. 12, which apply just to the mass range expected for the $P^+$. It seems necessary to change either the prediction of the $P^+$ mass (33) or of its couplings (47), or perhaps a few more ETC rats will start leaving the sinking ship. If the $P^+$ are not yet excluded, they should show up as prominent decay products of top quarks, and even of toponium\(^{32}\) $\zeta$. From the mass (33) and the couplings (47) one estimates

\[ \frac{\Gamma(t \to P^+ + \ell)}{\Gamma(t \to b\bar{q}\bar{q})} = O(10^4) \left( \frac{204 \text{ GeV}}{m_t} \right)^2 \] (49)

so that the top threshold would be an excellent $P^+$ factory. In fact one finds the toponium decay width

\[ \Gamma(\zeta \to b + P^+ + \ell + x \quad \text{or} \quad \bar{b} + P^- + t + x) \sim 5 \text{ MeV} \] (50)

which is much larger than conventional decay modes of $\zeta$ such as three gluons. Thus the bulk of toponium and $t\bar{t}$ continuum final states should be $P^+P^- + X$. If the $P^+$ decay mostly into $\tau\nu$ as suggested by (48), then the final states will be mainly $\tau + \nu + \bar{\tau} + \bar{\nu} + X$ with relatively little hadronic energy\(^{32}\): very different from the final states usually assumed in searches for the top quarks!
Finally, we turn to the neutral colour singlet FGBs $P^0$ and $P^3$. As mentioned before (35), they should have masses $20,22-24$

$$m_{P^{0,3}} \approx \frac{O(N^2_{\text{fc}})}{m_{P^S}} \approx \frac{1}{2} \text{eV} \times \left(\frac{300 \text{GeV}}{m_{P^S}}\right) \times 2^{0.1} \lesssim 3 \text{eV} \quad (51)$$

Clearly, any improvement in the upper limit on $K^0 + \mu e$ decay translates into an improvement in the lower limit (34) on $m_{P^S}$, and hence reduces the upper limit (51) on $m_{P^{0,3}}$. Also, it was mentioned at the end of Section 7 that $P^0$ and $P^3$ should have $O(\frac{e^2}{m^2_{H^*} m^2_W})$ flavour changing coupling either to charge $\frac{2}{3}$ quarks (u class models) or to charge $-\frac{1}{3}$ quarks (d class models) or to both. Even in the absence of these couplings one expects $g G_F m^3_q / m^2_W$. These mean that the non-observation of $K^+ + \pi^+ + P^{0,3}$ imply in either case

$$m_{P^{0,3}} \gtrsim 350 \text{ MeV} \quad (52)$$

In u class models one expects decay branching ratios

$$B(D \rightarrow P^{0,3} + X) = O(10^{-2}) \quad \text{if} \quad m_{P^{0,3}} \lesssim 1/2 \text{ GeV}$$

$$B(B \rightarrow P^{0,3} + X) = O(10^{-3}) \quad \text{if} \quad m_{P^{0,3}} \lesssim 4 \text{ GeV} \quad (53a)$$

whereas in d class models one expects

$$B(D \rightarrow P^{0,3} + X) = O(10^{-10}) \quad \text{if} \quad m_{P^{0,3}} \lesssim 1/2 \text{ GeV}$$

$$B(B \rightarrow P^{0,3} + X) = O(100 \%) \quad \text{if} \quad m_{P^{0,3}} \lesssim 4 \text{ GeV} \quad (53d)$$

Could the $P^{0,3}$ have been detected if they were produced at the rates (53)? In the monophagic model of Section 7 one expects
while $P^0$ should have a smaller branching ratio into $\mu^+\mu^-$ because it has an additional $gg$ decay mode. Hadronic decay modes of the $P^0, P^3$ are probably quite messy and difficult to pick out. It is probably not yet possible to exclude $u$ class models (53u) on the basis of the $\mu^+\mu^-$ decay branching ratio (54), though forthcoming D factory experiments might be able to exclude $B(D \to P^0, P^3 + X) = O(10^{-2})$, hence ruling out $u$ class models with $m_{P^0}, m_{P^3} \leq 1.2$ GeV. More seriously, the recent upper limit from CESR

$$B(D \to P^0, P^3 + X) < 0.2 \%$$

seems to exclude $d$ class models (53d) completely. While the $P^0$ and $P^3$ are not dead yet, to use the immortal words of Monty Python, they do look rather sick and one can sense a few more ETC rats leaving the sinking ship.

Before discussing alternative theories, however, it is worth recording some other reactions in which one might imagine $P^0, P^3$ could be produced $^{25, 39}$ and contrast their production rates with the conventional physical Higgs $H$. One example is Onium + $H$ or $P^0$ or $P^3 + \gamma$, where one expects in the monophagic model $^{25}$ of Section 7:

$$B(0 \to P^0, P^3 + \gamma) \approx 3 B(0 \to H + \gamma) \approx 3 \% \times \left( \frac{m_0}{40 GeV} \right)^2$$

unless the exciting alternative decay mode (50) opens up. It has often been observed $^{12}$ that the $2^0$ is a good catalyst for making the $H$. The same is not true for $P^0$ and $P^3$:
\[ B(Z^0 \rightarrow p^0, 3 + \gamma) = O(10^{-7} \text{ to } 10^{-9}) \]  

(57)

which is much less than \( B(Z^0 \rightarrow H + \gamma) \) and unobservably small, while

\[ \frac{\sigma(e^+e^- \rightarrow Z^0 + p^0, 3)}{\sigma(e^+e^- \rightarrow \mu^+\mu^-)} \leq O(10^{-2}) \]

(58)

for \( E_{\text{c.m.}} \leq 250 \text{ GeV} \)

is also much smaller than the analogous process \( e^+e^- \rightarrow Z^0 \rightarrow H \). It is clear from Eqs. (56)-(58) that if \( p^0, 3 \) are ever found they will be distinguishable from a physical Higgs.

9. SUPERSYMMETRIC TECHNICOLOUR

In view of the clouds gathering over extended technicolour theories, various people (including some of our previous rats) have gone on to consider the alternative way of making Higgses "natural" mentioned in Section 2, namely incorporating simple supersymmetry (susy) so as to cancel all the quadratic divergences of Fig. 4. The set of interactions proposed in "supercolour" or "supersymmetric technicolour" theories is

\[ SU(3)_c \times SU(2)_L \times U(1) \times SU(N)_{SC} \times SU(N)_{TC} \times \text{SUSY} \]

(59)

The immediate effect of the simple susy factor is to provide a direct connection between bosons and fermions: in the underlying theory they must occur in equal numbers and obey certain mass relations. Since the known fermions and bosons cannot be so related by susy, one must introduce new supersymmetric partners for the known particles, e.g. Majorana fermions to accompany the gauge bosons, scalars and pseudo-scalars (squarks and sleptons) to accompany the quarks and leptons.

The non-observation of such supersymmetric partners means that susy must be broken, and the suggestion is that this occurs dynamically via the formation of condensates by the new strong \( SU(N)_{SC} \) interactions:
\[ \langle 0 | \bar{f}_s f_s | 0 \rangle = C(N_{sc}^3), \quad \langle 0 | s_s^\dagger s_s | 0 \rangle = C(N_{sc}^2), \text{ etc.} \quad (60) \]

where \( f_s \) and \( s_s \) are "supercoloured" fermions and scalars respectively. It should be emphasized\(^{40}\), however, that nobody has yet proved that dynamical breaking of susy can occur in four dimensions. If it does, the unseen supersymmetric partners can now acquire mass, for example the "wino" \( \tilde{W} \) (supersymmetric partner of the \( W \)) and the "bino" \( \tilde{B} \) [supersymmetric partner of the \( B \_ \mu \) U(1) gauge field] can acquire masses\(^5\) through diagrams like those in Fig. 13a which are

\[ m_{\tilde{W}} \sim 4 \text{ TeV} \quad m_{\tilde{B}} \sim 1 \text{ TeV} \quad \text{if} \quad \Lambda_{sc} \sim 10 \text{ TeV} \quad (61) \]

while higher order diagrams (Fig. 13b) give masses\(^5\) to the supersymmetric partners of conventional fermions (the squarks and sleptons), and to the Higgs

\[ m_{sQ}, m_{sL}, m_H \sim (\text{few} \ 00) \text{ GeV} \quad (62) \]

The Higgs is protected from acquiring a larger mass by the underlying supersymmetry appearing at the scale of \( \Lambda_{sc} \). The low energy theory therefore contains light quarks, leptons, gauge bosons and a "natural" Higgs.

So far we still have \( SU(2)_L \times U(1) \) unbroken and there are no masses for the conventional quarks and leptons. One might guess that these could be achieved using the Higgs whose mass we have been protecting. Dine, Fischler and Srednicki\(^5\) claim that when this is tried, one finds unwanted \( U(1) \) symmetries with unacceptable real Goldstone bosons and/or axions. They propose instead to make most of the \( SU(2) \times U(1) \) breaking via technicolour [in parentheses in Eq. (59)], though the Higgs also develops a small vacuum expectation value. The
fermions then acquire masses not through technicolour or its extensions which were the source of our problems in the previous sections, but through conventional couplings of the fermions to the protected Higgs as in Fig. 5a. The fermion masses are then somewhat smaller than $m_w$ because $<0|\phi|0> < \Lambda_{TC}$, and the scale of $0(10)$ TeV for $\Lambda_{SC}$ was chosen with this scale of the Higgs system in mind.

In such a theory the absence of "horizontal" gauge boson interactions and the fermion mass generation by a single Higgs avoid all the flavour-changing neutral interaction problems\textsuperscript{10} discussed in Section 7. We can now get away with only one weak doublet of technifermions, so we are back to the spin-zero spectrum of the toy model in Section 3. Therefore, there are no troublesome light PGBs left uneaten by the $W^\pm$ and $Z^0$, of the types\textsuperscript{11} which caused us so much grief\textsuperscript{24,36} in Section 8. On the other hand one has to work quite hard to avoid phenomenologically unacceptable axion-like particles.

There are some interesting light particles proposed by supercolour theories which experimentalists should watch out for. One is the gluino, the spin 1/2 partner of the gluon, which gets its mass through higher order diagrams like that shown in Fig. 13c. It is estimated\textsuperscript{5} that they yield

$$m_{\tilde{g}} = O(5) \text{ GeV}$$

though this is subject to considerable uncertainties. It is expected to decay into a neutrino-like particle and a gluon with a lifetime

$$\tau_{\tilde{g}} = O(10^{-10}) \text{ sec}.$$  (64)

One could look for such a particle in multiplet jet events in the $e^+e^-$ continuum (Fig. 14a) or in anomalous toponium\textsuperscript{41} decays (Fig. 14b). In either case, the relatively long lifetime (64) should provide a distinctive signature. Another type of supersymmetric particle expected
to be relatively light is the Higgsino \( \mathcal{H} \), the spin \( 1/2 \) partner of the Higgs. It is estimated to have a mass

\[
m_H = 0(30) \text{ GeV}
\]  

(65)

which puts its pair-production comfortably within the range of LEP or the SLC.

While these supersymmetric technicolour theories are as yet relatively unexplored, they do seem to avoid the worst of the phenomenological problems with extended technicolour theories discussed in Sections 7 and 8. They also provide some interesting particles to look for at present (63), planned (65) and imagined (61), (62) particle accelerators.

10. ARE THE WEAK INTERACTIONS STRONG?

The new sets of strong interactions discussed up to now in this talk have all been attempts to address the "naturalness" problem\(^3,4\) of elementary scalar fields. Finally, I would like to discuss a daring model\(^6\) with strong interactions giving composite gauge fields, mainly in the hope of giving the SLC people\(^42\) nightmares. The conventional view is that the weak interaction gauge group \( SU(2)_L \times U(1) \) is spontaneously broken and has a weak coupling \( g^2/4\pi = 0(1/30) \), whereas the strong interaction gauge group \( SU(3) \) is unbroken and confining. Abbott and Farhi\(^6\) propose the heresy that \( SU(2)_L \) may in fact also be strong: \( g^2/4\pi = 0(1) \), unbroken and confining. In their approach the observed left-handed fermions in fact become composites\(^43\) of confined elementary fermion fields and Higgs fields:

\[
\mathcal{F}_L \rightarrow \left( \mathcal{F}_L \cdot \mathcal{H} \right)
\]  

(66a)

while the right-handed fermions are still elementary and the \( W \) bosons become composites of Higgs fields.
$W_\mu \rightarrow (\widetilde{H}^+ \widetilde{e}^- \overline{\nu}_\mu \nu_\mu)$ (66b)

as does the physical Higgs. The theory possesses a global SU(2) symmetry which is sufficient\textsuperscript{6} to guarantee the desired form of the low energy effective four-fermion interaction:

$$\frac{1}{4} \mathcal{L}_{\text{eff}} = \frac{g_F}{\sqrt{2}} \left( \frac{1}{2} \overline{\nu}_\mu J^\nu _\mu + (J^+ - J^-) - \sin^2 \Theta_W \overline{\nu}_\mu \nu_\mu \right) + \ldots$$ (67)

The over-all normalization $g_F/\sqrt{2}$ is taken from phenomenology. Remembering the classic formula

$$\frac{g_F}{\sqrt{2}} = \frac{g^2}{8 m_W^2}$$ (68)

we expect that the W mass will be larger than in the conventional Glashow- Weinberg-Salam\textsuperscript{1}, just because $g^2$ is larger than is usually believed.

The $J^3_\mu - J^{em}_\mu$ mixing term in the effective Lagrangian (67) is generated by a deviation from point-like structure in the strongly interacting $W^3 - \mathcal{W}^{+}_{L} \mathcal{W}^{-}_{L}$ vertex\textsuperscript{44}. Calculating it is an unsolved strong interaction problem, and it is not completely clear that the experimental value of $\sin^2 \Theta_W$ can in fact be reached. The dots in (67) represent an additional form of neutral current-current interaction which may be present, but it can be argued\textsuperscript{6} that it should be relatively small [experimentally its relative strength is $\leq 0(5\%)$].

Crucial experimental tests of this foolhardy scenario are of course the masses of the $\widetilde{W}^\pm$ and $Z^0$, which should be heavier than expected, up to 200 GeV. Eventually, there should be radial recurrences of them at very high energies [cf. the $\rho'$ (1600 MeV)]. At low energies one should look for the dotted deviations (67) from the conventional neutral current interactions. This strong weak interaction model looks so bizarre that one is rather surprised that it cannot be ruled out
instantaneously. The most serious phenomenological problem it faces at present may be the anomalous magnetic moment of the muon. Successful QED calculations tell us that the scale of compositeness in the muon's magnetic form factor must be $< (600 \text{ GeV})^{-1}$, whereas to get $\sin^2 \theta_W = 0(1/4)$ we need structure in its electric form factor on a scale $[0(100) \text{ GeV}]^{-1}$. In the absence of reliable strong interaction calculations one can always cross one's fingers and hope for the best, but this point is worrisome.

11. GENERAL OUTLOOK

What messages can we draw from this review of possible new strong interactions and composite models of Higgs fields in particular? It is clear that the original idea of dynamical symmetry breaking as embodied in technicolour is very elegant. On the other hand things get rather messy when one tries to give masses to fermions by extending technicolour. In addition to theoretical questions about the likely pattern of dynamical symmetry breaking in such theories, one also has to contend with serious phenomenological challenges arising from flavour changing neutral interactions in the $K^0 - \bar{K}^0$ and $D^0 - \bar{D}^0$ systems in particular. There are also problems with the plethora of light spin-zero bosons which are expected in ETC models but not yet seen. However, neither of these two arguments should be construed as proving ETC to be wrong, and one should press on with experimental tests, since the phenomenological suggestions thrown up by ETC are quite likely to be predictions of other theories as well. Therefore, one should carry on with probes of flavour changing neutral interactions such as searches for $D^0 - \bar{D}^0$ mixing, $K^0 \to \mu e$ decay and $\mu \to e$ conversion. One should also continue searches for light spin-zero bosons in $e^+ e^-$ annihilation, onium and heavy quark decays.
The new alternatives to ETC which exploit supersymmetry\textsuperscript{5} are somewhat embryonic, but promising. They avoid the worst of the problems with flavour-changing neutral interactions and light spin-zero bosons. They are worth studying further both theoretically (can supersymmetry be broken dynamically? Can one avoid unacceptable axion-like particles?) and phenomenologically (can one detect gluinos and/or Higgsinos?) Supertheories are worth watching.

Finally, we had the farout possibility that the weak interactions are in fact strong. This is of course a ridiculous idea, but we cannot yet prove it to be wrong. And won't Abbott and Farhi\textsuperscript{6} be looking smug if it turns out that the $W^\pm$ and $Z^0$ do weigh much more than 100 GeV?

However one may judge these various possible cases on the grounds of aesthetics or theoretical plausibility, one comment cannot be denied. In contrast to theories with a great desert, they do provide plenty of experimental tests to be performed at contemporary accelerators. You can prove these theories right or wrong. Good luck!
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mixing where ETC theories are close to the experimental limits on
flavour-changing neutral interactions. However, it seems likely
that the problems with helicity-flipping moment operators ($\mu^*\gamma_\mu$, $d_\mu$)
have been exaggerated in that paper as they should be suppressed by
an extra factor of $\Lambda_{\text{EC}/m_{\text{EC}}}$ - J. Preskill, private communication,
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**Fig. 1:** The ascetic, minimal picture of grand unified theories with a great desert.

**Fig. 2:** The less boring, more cluttered picture of theories with new strong interactions such as technicolour or supercolour.
Fig. 3: (a) The $W - \phi$ coupling in elementary Higgs theories, which (b) can be repeated indefinitely to give the $W$ mass.

Fig. 4: Samples of quadratically divergent one-loop diagrams which make a light elementary Higgs "unnatural".

Fig. 5: (a) A fermion-antifermion coupling to an elementary Higgs field, and (b) the technicolour analogue which requires (c) a four-fermion $(\bar{f}f)(\bar{F}F)$ coupling, provided (d) in ETC theories by heavy boson $E$ exchange, giving (e) a mass to the fermion $f$ through a condensate $\langle 0 | \bar{F}F | 0 \rangle \neq 0$ denoted by $\phi$. 
Fig. 6: A minimal ETC scenario, with one technigeneration, technicolour, extended technicolour and horizontal extended technicolour interactions.

250 GeV $\quad \bar{P}_8^{0,3,\pm}$ colour octets

160 GeV $\quad \bar{P}_{\bar{U}E, \bar{U}N, \bar{D}E, \bar{D}N}$

0(10) GeV $\quad P^{0,3,\pm}$ colour singlets

Fig. 7: The spectrum of PGBs expected in the minimal ETC model of Fig. 6 with the technifermions in a complex representation of the ETC group.

Fig. 8: A contribution to $K^0 \to \mu e$ decay from the crossed channel exchange of a Pati-Salam SU(4) gauge boson.
Fig. 9: Can one see a bump from hadron + hadron + $P_B + X$ to $P_B + t\bar{t}$? [taken from Ref. 33].
Fig. 10: (a) The leptoquark $P_{LQ}$ can be produced directly in $ep$ collisions, and (b) can mediate $K^0 \rightarrow \mu e$ decay.

Fig. 11: Cross-sections for $P_{LQ}$ production in high energy $ep$ collisions [taken from Ref. 34].
Fig. 12: Limits at 90% confidence from the JADE experiment 36) on the masses and branching ratios $B_\tau$ into $\tau + \nu$ of the conjectured charged colourless PGB $P^\pm$. 
Fig. 13: Diagrams which give masses (a) to the winos $\tilde{W}$ and binos $\tilde{B}$, (b) to the squarks $\tilde{Q}$ and sleptons $\tilde{\ell}$, and (c) to the gluino $\tilde{g}$.

Fig. 14: One can look for the gluino (a) in multiple jet events in the $e^+e^-$ continuum, or (b) in anomalous toponium decays.