Summary

This is the report of the Higgs and Technicolor Summer Study. The official members of the subgroup were Charles Baltay, Estia Eichten, Peter Igo-Kemenes, and Kenneth Lane. In addition, we received wisdom on many matters from Ian Hinchliffe, Harris Kang, Gordy Kane, Lawrence Littenberg, Martin Perl, Michael Peskin, and other members of the various collider working groups.

We review the properties of elementary Higgs mesons occurring in SU(2) \(\otimes U(1)\) electroweak models with one or more Higgs doublets and of technipions expected in models in which the symmetry breaking is dynamical in origin. We discuss their couplings to ordinary matter, paying special attention to similarities and differences between couplings of elementary Higgses and technipions to specific channels. We also stress which couplings are reliably model-independent and which are not. We use these results to discuss the most likely production and detection modes of these mesons in e^+e^-, hadron-hadron, and ep colliders and we show how to distinguish experimentally among the various scenarios for the scalar sector.

I. Introduction

The structure of the scalar sector of the electroweak interactions is perhaps the greatest mystery of high-energy physics. Even assuming that the gauge and fermion sectors are correctly described by the standard SU(2) \(\otimes U(1)\) model of Glashow, Weinberg, and Salam, we are still completely ignorant of the precise mechanism by which SU(2) \(\otimes U(1)\) is spontaneously broken down to electromagnetic U(1). Of course, we all believe that the so-called Higgs mechanism is responsible, that it gives masses both to electroweak gauge bosons W and Z and to quarks and charged leptons. But, with only one rather indirect exception, we haven’t a shred of experimental evidence that this is true nor, if it is, of what the Higgs sector consists and how it works dynamically. Setting these issues, we believe, is the most important task of high-energy physics experiments in this decade.

Currently, there are two general and quite distinct scenarios for the Higgs sector. In the first, elementary (i.e., pointlike) scalar mesons are assigned to one or more doublets of the weak SU(2) group. The electrically neutral components of these doublets are assumed to develop vacuum expectation values. Just why this happens no one knows; but, once it does, SU(2) \(\otimes U(1)\) breaks to U(1)EM while fermions coupled to the Higgs mesons acquire mass. In the earliest discussions of SU(2) \(\otimes U(1)\), it was assumed that only one Higgs doublet was present, and much of electroweak phenomenology (not to mention planning for experiments) has incorporated this assumption. In fact, there can be any number of Higgs doublets. The questions of how many there are, what are the masses of the physical scalars that survive the Higgs mechanism, what are their couplings to ordinary matter and so on can only be settled experimentally.

The second scenario goes under the general rubric of dynamical symmetry breaking, or "technicolor" for short. This program attempts to give a dynamical explanation, in terms of new strong interactions at \(\sim 1\,\text{TeV}\), for the breakdown of SU(2) \(\otimes U(1)\), and the appearance of quark and lepton masses. Here, spinless mesons composed of technifermion-antifermion pairs play the role of the pointlike mesons of elementary Higgs models. The lightest and most accessible of these are called "technipions". While the idea of dynamical symmetry breaking is very attractive and natural, no phenomenologically acceptable technicolor model has been constructed. Thus, many important details of the scalar sector of technicolor are not settled theoretically — a failing they share with multi-doublet models of the elementary Higgs scenario. In any case, whether electroweak symmetry breaking is dynamically induced at 1 TeV or not is a question that must be answered. If the breaking is dynamical, mapping out the spectrum of technipions can determine their fundamental constituents, just as was done for the ordinary strong interactions.

Fortunately, the experiments of this decade can explore the scalar sector, distinguish between the two scenarios and, possibly, select among the various possibilities within each scenario. The purpose of this report is to spell out how the job can be done. We hope this report is useful to experimentalists, the intended audience.

In Section II we give a lightning review of the standard model with a single elementary Higgs doublet. This model has one observable scalar, H^0. This summary of its properties and techniques for its discovery is included for completeness and as a contrast to the discussion of multi-doublet and technicolor models in Sections III and IV.

Section III is devoted to a detailed discussion and comparison of the general properties of charged and neutral Higgses and technipions. We begin with some of the motivations for going beyond the standard one-doublet model. Then we catalog the types of technipions and other technihadrons that may occur, using a toy model due to Farhi and Susskind as an example. Finally, we discuss what is known about the couplings of Higgses and technipions to ordinary matter, with special attention to the differences that will distinguish between them. We have taken extra care to emphasize which properties are model-independent, hence reliably known, and which are not and based merely on educated guesses.

The production and detection of the electroweak scalars in e^+e^-, hadron-hadron, and ep colliders is discussed in Section IV. The couplings to ordinary matter tabulated in Section III are used to focus on the most promising means for discovering them and differentiating between the two scenarios. There we have touched only lightly on the important question of the character and estimates of backgrounds to various processes. We feel this can be well done only by the concerned experimentalists and the reader is referred to such discussions elsewhere in these proceedings.

Our main conclusions are these: (1) An e^+e^- collider with luminosity \(\geq 2 \times 10^{31}\) at toponium and the \(2\) can, with relative ease, discover most scalars existing below \(\sim 45\,\text{GeV}\) in mass and resolve the elementary versus dynamical issue. Ultrahigh energy and luminosity e^+e^-
colliding linacs copiously produce all charged techni- pions. (2) pp and pp collider experiments can produce and detect only the limited class of techni pions coupling strongly to gluons and, probably, none of the elementary Higgs mesons. However, given the time scale for completion of various machines, TEV I may well make the first discovery of a techni pion - if it exists. 

(3) ep colliders are limited to production of heavy colored techni pions which couple to a quark-lepton pair or to a gluon plus photon. Observable rates probably require $s \sim 1$ TeV.

Much of what we say here is not new. This report has benefitted from a number of previous reviews, specifically those in Ref. 10 on the standard Higgs, $H^0$, and charged Higgses and those in Ref. 11 on techni pions. Also, as noted in the text, we have relied on various preliminary reports of the $e^+e^-$ and hadron-hadron working groups at the Summer Study. This report largely may be viewed as a condensate of all those more detailed ones.

Finally, we should mention those other main currents of theoretical thought which we shall not cover explicitly. These include the scalar sector of "standard" grand-unified models, left-right symmetric models and supersymmetric electroweak or grand-unified models. To a large extent, the low-energy scalar sector of these models is phenomenologically similar to the elementary Higgs cases discussed here. Of course, we refer the reader to the appropriate group reports in the proceedings.

II. The Standard Model With a Single Higgs Doublet
A. Properties of $H^0$

The electroweak gauge group $SU(2) \times U(1)$ with a single complex $SU(2)$-doublet of elementary (pointlike) Higgs mesons $\phi$ is sufficient to describe practically all known weak and electromagnetic phenomena. In this model, three of the four real component fields of $\phi$ combine with the $W^\pm$, $W^0$ and $Z^0$ as they become massive, leaving behind only a single neutral Higgs scalar, $H^0$.12 Because this minimal model is so simple, the elementary couplings of $H^0$ to ordinary matter-quarks, leptons and weak gauge bosons - are precisely known. The relevant couplings we shall need for production and decay of $H^0$ are those to fermions,

$$f_{H^0ff} = - \frac{3}{2} g_f^{\phi \phi} \sum_{\text{fermions,f}} m_f \bar{\psi}_f \psi_f H^0, \quad (1)$$

where $m_f$ is the so-called current-algebra mass of quark or lepton $f$, and those to electroweak bosons,

$$f_{H^0WWZZ} = \left[ g W^\mu W^{\mu*} + g Z^{\mu*} Z^{\mu} + g S^2 + g^\prime g^\prime S^2 \right] \frac{1}{4} \sum_{\mu} \mu \mu H^0, \quad (2)$$

where $g = e / \sin \theta_W$, $(g^2 + g^\prime)^2 = 2 e / \sin^2 \theta_W$, $\sin^2 \theta_W = 0.22$, and $m_W$ and $m_Z$ are the $W^\pm$ and $Z^0$ masses, respectively. Note that there is no elementary coupling of $H^0$ to photons.

Unfortunately, the mass of $H^0$ is a free parameter of the model and, so, is completely unknown. Various theoretical and phenomenological arguments put the following very loose bounds on $m_H$:13

$$\text{(15 MeV)} \leq m_H \leq (1 \text{ TeV}) \quad (3)$$

In discussing the production and decay of $H^0$, therefore, the best one can do is use Eqs. (1) and (2), together with calculations of induced $H^0$ couplings to photons, etc., to give expectations of what will happen for various ranges of $m_H$. In this section, we shall assume that $m_H \geq 5$ GeV.

The branching ratios for $H^0$-decay are shown in Fig. 1 for $5 \text{ GeV} \leq m_{H^0} \leq 50 \text{ GeV}^{14}$ (assuming $m_t = 20$ GeV). If $m_{H^0} = 10$ GeV, mixing with the $3 P_0$-states of the $T$-system can become important.15 This gives a complicated dependence of decay modes on mass shown in Fig. 2. The
important thing to remember is that $H^0$ couples to fermion mass, $\Gamma(H^0 - f) \approx c_m^2 m_f^2$, with a multiplicative factor of 3 if $f$ is a color-triplet quark. Unless there are very heavy quarks or leptons coupling to $H^0$, the modes $H^0 \rightarrow WW$ and $Z^0 H^0$ become dominant for $m_H > 4 m_t \approx 200$ GeV. From all this, it is clear that $H^0$ is best searched for by missing-mass techniques.

B. Production and Detection of $H^0$ in $e^+e^-$ Colliders

1. Quarkonium - $H^0 + \gamma$

Since $H^0$ couples to fermion mass (Eq. (1)), the decay rate to $H^0 q\bar{q}$ of a heavy $qq$ bound state, $V_q$, can be appreciable. Relative to the electronic width, the rate is

$$\frac{\Gamma(q\rightarrow H^0 \gamma)}{\Gamma(q\rightarrow e^+e^-)} \approx \frac{G_F^2 m_q^2}{4\pi m_V^2} \left(1 - \frac{m_H^2}{2m_V^2}\right), \quad (4)$$

where $m_V \approx 2 m_q$. Consider toponium, $\zeta = V_q$, for example ($m_\zeta > 400$ GeV). So long as $m_H < 0.5 m_\zeta$ and there are no charged Higgses more massive than the t-quark (see Sec. III), Eq. (4) implies that $\Gamma(\zeta \rightarrow H^0 \gamma) \approx 0.2 \Gamma_\zeta$. With good photon detection and energy resolution ($E_\gamma / E_\zeta \approx 0.1$), there should be no problem reaching this limit.\(^5\) Either the $H^0$ can be found at toponium, or a very interesting lower bound on its mass on its mass will be set.

2. $e^+e^- \rightarrow Z^0 \rightarrow H^0 \nu\bar{\nu}$

As can be seen from Eq. (2), there is no elementary $H^0 Z^0\nu\bar{\nu}$-coupling, so $e^+e^- \rightarrow Z^0\gamma \rightarrow H^0 \gamma$ is not a very good way to search for $H^0$. However, the relatively large $H^0 Z^0\nu\bar{\nu}$-coupling makes $e^+e^- \rightarrow Z^0 \rightarrow H^0 \pi^0$ a viable way to search for $H^0$.\(^6\) The rate relative to $Z^0 \rightarrow \mu^+\mu^-$ is plotted in Fig. 3\(^8\) as a function of $H^0$ mass. Assuming an integrated luminosity yielding $\sim 10^7$ $Z^0$/year, we see that this method is sensitive to $m_H \leq 43$ GeV (corresponding to $\sim 30$ events/year). The best resolution on $m_H$ will come from concentrating on $e^+e^- \rightarrow e^+e^-$, for which $\Delta E_{e^+e^-} / E_{e^+e^-} \approx 0.1 E_{e^+e^-}$. This implies an acceptably small error ($\leq 25\%$) on $m_H$ for $m_H > 20$ GeV. Furthermore, the backgrounds to this process, $Z^0 \rightarrow q\bar{q} \rightarrow e^+e^-\nu\bar{\nu}$, are expected to be easily cut away. For $m_H < 20$ GeV, $\zeta \rightarrow H^0 \nu\bar{\nu}$ probably is a better way to find $H^0$. See M. Goldberg's report, and references therein, for further details.\(^9\)

3. $e^+e^- \rightarrow Z^0 \rightarrow H^0 \nu\bar{\nu}$

Finally, if $m_H > 45$ GeV, the only resort is to boost the energy well above the $Z^0$ mass and again take advantage of the large $H^0 Z^0\nu\bar{\nu}$-coupling. The final $Z^0 \rightarrow \nu\bar{\nu}$ is probably the best detected in its $e^+e^-\nu\bar{\nu}$-decay mode. For $s \leq 175$ GeV and $Z^0 \times 2 \times 10^{31}$ cm$^{-2}$ sec$^{-1}$, this method is sensitive to Higgs masses $\leq 60$ GeV.

In closing this discussion, we mention that the development and implementation of vertex detectors which can find the decay kinks of $\tau^+$, $\nu_e$'s and $b$'s results from $H^0$-decay can enhance substantially the probability of finding $H^0$ in each of these $e^+e^-$ production mechanisms.

C. Production and Detection of $H^0$ in Hadron-Hadron Colliders

Discovery of the $H^0$ at existing or proposed pp and $p\bar{p}$ colliders will be a difficult, if not impossible, task. Three of the more promising production modes mention in the literature are discussed here.

1. pp or $p\bar{p} \rightarrow H^0 + X$ via Quark Fusion; $H^0 \rightarrow \mu^+\mu^-$

Ellis, Gaillard and Nanopoulos\(^10\) proposed production of $H^0$ from valence and sea quarks in the nucleon and detection via the $\mu^+\mu^-$ decay mode of $H^0$. They find

$$\frac{\sigma(pp \rightarrow H^0 \rightarrow \mu^+\mu^- X)}{\sigma(pp \rightarrow \mu^+\mu^- X)} \approx 10^{-3} \frac{m_H \Gamma(H^0 \rightarrow \mu^+\mu^-)}{\Delta m_H \Gamma(H^0 \rightarrow A_{11})}, \quad (5)$$

where $\Delta m_H$ is the $\mu$-pair invariant mass resolution at the Higgs mass. $\Delta m_H$ is unlikely to be much larger than 10. The factor of $10^{-3}$ comes from assuming an SU(3)-invariant sea and is due to the smallness of $G_F m_F^2 / m_H \approx 175$ MeV. Thus, Eq. (5) is a drastic over-estimate for production in the case of pp collisions because the Drell-Yan cross section in the denominator is dominated by valence quark-antiquark annihilation.

From Eq. (1) and Fig. 1, it is clear that $B(H^0 \rightarrow \mu^+\mu^- X) < 3 \times 10^{-4}$ (for $m_H \approx 50$ GeV), and it falls to $8 \times 10^{-6}$ for $m_H > 2 m_t$, assuming a 20 GeV top-quark mass. Thus, this method amounts to looking for a peak for which the signal-to-background is at most $10^{-6}$.

2. $H^0$ From Gluon-Gluon Fusion

To avoid the suppression factor $G_F m_F^2$ in $H^0$-production due to its direct coupling to light quarks, Georgi, et al., proposed a gluon-fusion mechanism.\(^20\) Here, $H^0$ is produced through its coupling to a pair of gluons, a coupling which is not suppressed when it is induced by heavy quark loops for which $m_V \approx 2 m_q$. For quark masses satisfying this inequality, the quark-mass dependence

![Fig. 3. Decays rates for $\Gamma(Z^0 \rightarrow H^0 + \gamma) / \Gamma(Z^0 \rightarrow \mu^+\mu^-)$ and $\Gamma(Z^0 \rightarrow H^0 + \gamma) / \Gamma(Z^0 \rightarrow \mu^+\mu^-)$ for different values of $m_H / m_Z$ taken from Ref. 18.](image-url)
drops out and the production cross section is proportional to $C_s^2 G p N_f$, where $N_f$ is the number of contributing heavy quark flavors. The authors calculated the production cross section in $pp$ collisions up to $\sqrt{s} = 400$ GeV. It is $\sim$ 30 pb for $m_N = 5$ GeV, falling to 1 pb for $m_N = 70$ GeV. (See Fig. 4.) The problem still is: How is one to detect $H^0$? Even ignoring the difficulty of making invariant mass plots of $\tau^+\tau^-$, $c\bar{c}$, $b\bar{b}$ and so on, each of these Higgs decay modes must lie on a horrendous background. We believe this method will not work, but defer to the judgement of anyone who has given this mechanism more careful thought.

3. Associated Production of $H^0$ With $Z^0$ or $W^\pm$\textsuperscript{21}

This method, which relies on the relatively large $H^0Z^0Z^0$ and $H^0W^+W^-$ couplings, is the hadronic collider analog of the $H^0$ search in $e^+e^- \rightarrow Z^0H^0e^+e^-$, and it seems the most promising. The $pp$ and $pp$ cross sections for $H^0Z^0$ and $H^0W^+W^-$ are shown in Figs. 5 and 6. For the interesting mass range $m_N \geq 10$ GeV, the cross section is small and falling rapidly with increasing $m_N$.

For $m_N < m_Z$, L.-L. Chau Wang has suggested concentrating on $Z^0 \rightarrow H^0Z^0 \rightarrow H^0Z^0$ and looking for a bump at $m_{H^0} = m_Z - m_H$ in the dilepton invariant mass distribution.\textsuperscript{22} For $m_N = 15$ GeV and Isabellle parameters, $\sqrt{s} = 800$ GeV and $\int dt = 10^5$ cm$^{-2}$ sec$^{-1}$, she predicts an excess of $\sim 1$ event in the bump between 60 and 80 GeV. (The excess is $\sim 1$ event for Tevatron parameters.) To reduce backgrounds due to $Y\gamma$ and $Z^0$, she suggests triggering on a third, low-energy lepton resulting from semileptonic decay of Higgs decay products. From Fig. 5b, it is clear that this proposal is useful only for $m_N < 20$ GeV. For $20 \text{ GeV} \leq m_N \leq 2m_Z \geq 200$ GeV, discovery of $H^0$ by associated production will require detection and reconstruction of its decay products in addition to triggering on $Z^0$ or $W^\pm$. For $m_N > 200$ GeV, $H^0 \rightarrow Z^0Z^0$ and $W^+W^-$ are its major decay modes. Then one can attempt to trigger on two or three elec-

![Fig. 4. $d\sigma/dm_H = 0$ as a function of $H$ mass, taken from Ref. 20. Each shaded band represents a different center-of-mass energy, $\sqrt{s}$ = 27.4 GeV (slash up to the right), $\sqrt{s}$ = 60 GeV (dot), and $\sqrt{s}$ = 400 GeV (slash up to the left).](image)

![Fig. 5. Rate of associated production of the Higgs meson with $W^\pm$ or with $Z$, versus $M_H$, expressed as a fraction of total $W^\pm$ or $Z$ production (from Ref. 21): (a) In $pp$ collisions at $\sqrt{s} = 540$ GeV. (b) In $pp$ collisions at $\sqrt{s} = 800$ GeV. Production with $W^\pm$ is indicated by the dotted bands, with $Z$ indicated by slashes. Bands are shown for $\mu_W = 60$ GeV ($\mu_Z = 77$ GeV) [lower curves] and for $\mu_W = 90$ GeV ($\mu_Z = 99$ GeV) [upper curves]. Bands indicate the range of variation due to different quark-distribution-function parametrizations.](image)
Fig. 6. Rate of associated production of the Higgs meson with $\mathcal{W}^\pm$, $\mathcal{Z}^0$, or $\mathcal{Z}$, versus energy $\sqrt{s}$, expressed as a fraction of total $\mathcal{W}^\pm$, $\mathcal{W}^0$, or $\mathcal{Z}$ production (from Ref. 21). (a) In $pp$ collisions. (b) In $pp$ collisions. Rates are shown for several $M_H$ values, all using $\mu_G = 75$ GeV ($\mu_g = 86.6$ GeV).

III. Technicolor, Charged Higgses and All That

A. Motivations for Going Beyond the Standard Model

For as long as people have been investigating gauge models of the electroweak interactions, they have been unable to resist the temptation to generalize, complicate if you will, the standard model containing a single elementary Higgs doublet. Even within the strict confines of an $SU(2) \otimes U(1)$ gauge group and all the experiments supporting it, there is much room for generalization. In particular $SU(2) \otimes U(1)$ with any number of elementary Higgs doublets is consistent with all experiments.3

Two particularly good reasons for introducing more than one doublet have to do with CP-nonconservation. In the first, proposed by Lee23 and Weinberg24, the observed weak CP-violation is supposed to originate in the Higgs sector, either by spontaneous breakdown of CP-symmetry or by virtue of CP-nonconserving Higgs meson self-couplings. These mechanisms require two or three doublets at least. The second reason for introducing two or more doublets has to do with the problem of strong CP-violation induced by instantons.25

Peccei and Quinn26 showed that this problem can be eliminated by having at least two doublets related by a global $U(1)$ symmetry. As Weinberg and Wilczek27 observed, this Peccei-Quinn symmetry implies the existence of an almost massless pseudoscalar, the axion, with couplings to ordinary matter comparable to those of $\mathcal{H}^0$ (Eqs. (1) and (2)). This, apparently, introduces another problem, but we will not get into that here. Suffice it to say that no one has ever proposed a different, clearly workable solution to the strong CP-violation problem.

An $SU(2) \otimes U(1)$ model with two elementary Higgs doublets has eight spinless fields — four charged and four neutral. One neutral and two charged mesons are eaten by $\mathcal{Z}^0$ and $\mathcal{W}^\pm$ in the Higgs mechanism, leaving behind as physical scalars $\mathcal{H}^0$, $\mathcal{H}^\pm$, $\mathcal{H}^{\mp}$ and $\mathcal{H}^\ast$.28 A model with three or more doublets has an even richer spectrum. The existence of one or more extra doublets implies additional freedom, in the form of mixing angles, in the couplings of $\mathcal{H}^0$ and the $\mathcal{H}^\pm$ to ordinary matter so that, unfortunately, the precision of Eqs.(1) and (2) no longer holds. Rather than discuss them here, it is illuminating to compare and contrast these couplings with those of the technipions which will be introduced shortly. The discussion appears in Sec. III-C.

Finally, as in the case of the one-doublet model, the masses of Higgses in the multi-doublet models are almost completely unconstrained. On theoretical grounds, they must be $< 1$ TeV.10 The fact that b-quark decays are not dominated by $b \to H^-c$ or $H^-u$ implies that charged Higgses must be more massive than $m_b \approx 5$ GeV. That is all one can say.

In large part, it is this theoretical sloppiness with Higgs masses that has motivated the more ambitious attempts to go beyond the standard model — those associated with the terms dynamical symmetry breaking (technicolor/hypercolor) and supersymmetry (which we won’t discuss here; see I. Hinchliffe’s report in these proceedings).

Very briefly, “dynamical symmetry breaking” refers to the theoretical possibility of generating spontaneous breakdown of the electroweak gauge group $SU(2) \otimes U(1)$ to electromagnetic $U(1)$5,6 without the introduction of elementary Higgs multiplets. In this scenario, unlike elementary-Higgs models, the size of the Fermi constant, $G_F = 300$ GeV, has a simple dynamical origin: There is a new strong gauge interaction, technicolor, whose characteristic scale $\Lambda_T$ is analogous to $\Lambda_{QCD}$. Technifermions, analogous to quarks, interact via TC-boson exchange. Just like quarks, they acquire mass dynamically and, as a consequence of Goldstone’s theorem, massless pseudoscalar bound states of technifermions are formed. If technifermions couple to $SU(2) \otimes U(1)$ in the same way that quarks and leptons do, i.e., one chirality in doublets, the other in singlets, then three of the massless "technipions" are eaten by $\mathcal{W}^\pm$ and $\mathcal{Z}^0$. The gauge bosons acquire mass $\mu_G = \frac{g}{\sqrt{2}} G_F$ and $\mu_g = \frac{g^2}{8} \mu_F^2$, $F_T$ is the technipion decay constant, analogous to $F_L = 95$ MeV for ordinary pions. Thus, $G_F/\sqrt{2} = \frac{g^2}{8} F_T = 1/2 F_T$ so that $F_T = 250$ GeV = 2500 $f_T$. In many respects, then, technicolor is just a scaled-up version of QCD; in particular, $\Lambda_T \approx 2500 \Lambda_{QCD} = 0.5 - 1.0$ TeV.

At this point, there is a dramatic parting of the ways of elementary Higgs models and technicolor (TC) models. In the former there generally are just three massless Goldstone bosons, the unphysical ones eaten by $\mathcal{W}^\pm$ and $\mathcal{Z}^0$. The $\mathcal{H}^\pm$ and $\mathcal{H}^0$‘s left behind can, as noted, have almost any mass because free parameters in the Higgs self-interaction can be adjusted at will. In TC models, likewise, there are physical technipions left behind. Their mass is to a large extent precisely predicted because there are, in principle, no free parameters to adjust. Let us see how this comes about.

In all quasirealistic TC models constructed so far, there are at least four flavors of technihadrons of each chirality. This implies that there will be more than three massless Goldstone technipions formed when the technifermions acquire their dynamical mass. Only three are eaten; the rest remain as physical technipions which are strictly massless in the neglect of interactions which are weak compared to technicolor at scales of order $\Lambda_T$.8 Those weaker interactions do give mass to the physical technipions, that mass is calculable, and we believe we know what these interactions are. That is the basis of the statement that TC models have no free parameters.

The weaker interactions giving mass to technipions are the familiar color $SU(3)$ and the electromagnet $SU(2) \otimes U(1)$ and a postulated new interaction known as...
"extended technicolor" (ETC)\(^7\) or "sideways".\(^8\) The ETC interaction must be introduced to give ordinary quarks and leptons their observed nonzero current-algebra masses. This is accomplished by introducing a gauge coupling which connects quarks and leptons to technifermions. The ETC gauge group is spontaneously broken at a scale of order 100 TeV (\(^\uparrow\)) to a subgroup containing technicolor and ordinary color. (The value \(\sim 100\) TeV is set by the TC scale, 1 TeV, and the magnitude of quark and lepton masses.) The contributions of these interactions to the mass of various technipions can be well-estimated using current-algebraic techniques.\(^29\) The main uncertainty comes from the fact that no completely satisfactory TC model has yet been built, so that the ETC couplings are not precisely known. We shall summarize the mass estimates in the next subsection. For now, it is enough to say that the technipions corresponding most nearly to \(H^0\) and some of the \(H^\pm\)s are predicted to have masses \(\sim 5-40\) GeV, much less than the characteristic 1 TeV scale of technicolor.

We should also add here that ETC interactions are primarily responsible for the coupling of technipions to fermion-antifermion pairs. Thus, there is a certain amount of ignorance in calculating technipion decay branching ratios, just as there is for \(H^0\) and the \(H^\pm\). Technipion couplings to electroweak gauge bosons are determined by the more well-known couplings of their constituent technifermions to \(w^\pm\), \(Z^0\) and \(\gamma\). All this will be detailed soon.

Finally, there will be technicolor-neutral "technihadrons" analogous to the hadrons \(p, w, \epsilon\) and so on in QCD. The typical technihadron mass is \(\sim 1\) TeV. Indeed, the precise analog of the neutral Higgs(es) which gives mass to quarks and leptons in elementary Higgs models is such a technihadron and, so, is unreachable in colliders of this decade.

B. The Technicolor Zoo—Particle Types, Charges and Masses

The number and quantum numbers of light technipions and heavy technihadrons is fixed, as in QCD, by the number of technifermion constituents and how they transform under technicolor. Obviously, this is a completely model-dependent issue, one which is far from being settled. In particular, technifermions may also carry ordinary color or not. Each technifermion belonging to an N-dimensional representation of color SU(3) counts as N flavors since color SU(3) is weak compared to technicolor.

To sample the richness possible in the TC-singlet spectrum, we shall consider a specific toy model proposed by Farhi and Susskind\(^9\) and examined in some detail by Dimopoulos,\(^30\) Peskin,\(^31\) and Preskill.\(^32\) The elementary technifermions in this model are a pair of techniquarks, \(Q = (U,D)\), and a pair of technileptons, \(L = (N,E)\). Both chirality components of \(U, D, N, E\) are assumed to transform according to the same unspecified, but nonreal, \(N_{TC}\)-dimensional representation of the technicolor gauge group, \(G_{TC}\). Under (SU(3), SU(2), U(1)), they transform as follows:

\[
\begin{align*}
Q_L &= U_D^L : (3,2,Y) \\
U_R &= (1,1,Y + 2); D_R &= (1,1,Y - 2) \\
L_L &= N_E^L : (1,2,-3Y) \\
N_R &= (1,1,-3Y + 2); E_R &= (1,1,-3Y - 2)
\end{align*}
\]

The weak hypercharges assigned here guarantee that \(G_{TC} \otimes SU(3) \otimes SU(2) \otimes U(1)\) gauge currents have no anomalous divergences. As usual, the electric charge is \(Q_{EM} = T_3_{EM} + Y_{EW}\).

This specific model makes definite predictions of the number, charge and (approximate) mass of technipions and technihadrons. It is important that the reader understand that some, but not all, of these particles will exist in any other TC model.\(^8\) In particular, the lightest ones, called \(F^\pm, P^0\) and \(P^0\), below, occur in all models. Colored technipions will not exist in models in which all technifermions are ordinary color SU(3) singlets. With this caveat, let us get on to cataloging the denizens of the zoo.

This version of the Farhi-Susskind model has an \(SU(8) \otimes SU(8)\) chiral symmetry which is spontaneously broken down to \(SU(8)\) by the dynamical technifermion masses. This implies \((N,E) = 63\) Goldstone bosons, of which 3 are eaten by \(w^\pm\) and \(Z^0\). The remaining 60 physical technipions consist of 4 color singlets \(F^\pm, P^0\) and \(P^0\), 24 color triplets \(P_Q\) called leptoquarks, and 32 color octets \(P^0, P^\pm\) and \(P^0\). The color-singlet technihadrons are most nearly analogous to the \(H^0\) and \(H^\pm\). \(P^0\) is often referred to as \(\pi^T\) in the literature. The technifermion constituents, electric charge and mass of these spinless mesons are summarized in Table 1.\(^33\) As far as their strong TC interactions and their coupling to gauge bosons are concerned, these particles are all pseudoscalars. This is not generally true of their couplings to light fermionic matter, as we shall see.

\[
\begin{array}{c|c|c}
\text{Tecnipion/Technifermion Content} & Q & T_3 \\
\hline
\text{Color Singlets, } F & \left| F^0 \right| = \frac{1}{\sqrt{2}} \left[ \left[ F^+ \right] + \left[ F^- \right] \right] & 1 & 0 & 7-40 \\
\text{Color Triplet Leptoquarks, } P_Q (a = 1, 2, 3) & \left| P_{aQ} \right| = \frac{1}{\sqrt{2}} \left[ \left[ P_{aQ} \right] \right] & \pm (4Y + 1) & \pm (4Y - 1) & \pm (4Y + 1) \\
\text{Color Octets, } P_{aQ} (a = 1, \ldots, 8) & \left| P_{aQ} \right| = \frac{1}{\sqrt{2}} \left[ \left[ P_{aQ} \right] \right] & \pm (4Y + 1) & \pm (4Y - 1) & \pm (4Y + 1) \\
\end{array}
\]

\text{TABLE 1. Technipions in the Farhi-Susskind model (Refs. 9,30,31,32).} Q is the electric charge and \(T_3\) the weak isospin of the technipion. Y is the weak hypercharge of the QL-doublet. For unit normalization, the states should be divided by \(\sqrt{N_{TC}}\).
We have emphasized that these technipions are massless in the neglect of color SU(3), electroweak SU(2) \& U(1) and ETC interactions. The color-singlet \( P^+ \) acquire mass only from the latter two interactions and so they are the lightest technipions. The electroweak contribution can be calculated very reliably as it is completely model-independent. \(^8,31,32,33\) As with all such current-algebraic computations, one calculates the square of the mass. For \( P^+ \), the result is

\[ m_{P^+}^2 = \frac{2}{4\pi^2} \frac{\mu_{P^+}^2}{\mu_{P^+}^2} \left[ (7.0 - 8.5 \text{ GeV})^2 \right] \]

for \( \Lambda_{TC} = 0.5 - 1.0 \text{ TeV} \). For \( P^0 \) and \( P^0' \),

\[ m_{P^0}^2 \equiv m_{P^0'}^2 \equiv \left( \frac{2}{4\pi^2} \frac{\mu_{P^0}^2}{\mu_{P^0}^2} \right) \left[ (2 - 40 \text{ GeV})^2 \right] \]

While the ETC contribution to \( m_{P^0}^2 \) can, in principle, be fairly precisely calculated in a specific model, we believe we should give here a range of masses reflecting our ignorance of the "true model". If \( \Lambda_{TC} \approx 10 - 40 \text{ GeV} \),

\[ \left( m_{P^0}^2 \right)_{ETC} \approx \left( m_{P^0}^2 \right)_{EW} \]

Using \( m_P = \left[ (m_{P}^2)_{EW} + (m_{P}^2)_{ETC} \right]^{\frac{1}{2}} \), the expected range of the lightest technipion masses is

\[ m_{P^+} \approx m_{P^0} \approx m_{P^0'} \approx 10 - 40 \text{ GeV} \]

if \( \left( m_{P^0}^2 \right)_{ETC} \gg \left( m_{P^0}^2 \right)_{EW} \),

or

\[ m_{P^+} \approx 7 - 9 \text{ GeV} \]
\[ m_{P^0} \approx m_{P^0'} \approx 2 - 3 \text{ GeV} \]

if \( \left( m_{P^0}^2 \right)_{ETC} \ll \left( m_{P^0}^2 \right)_{EW} \).

The contribution of electroweak and ETC interactions to the masses of color triplet and octet technipions should be comparable to those just given. This is much less than the contributions from color SU(3), which have been estimated to be\(^30,31,32\)

\[ m_{LQ}^0 \approx 160 \sqrt{\frac{\Lambda_{TC}}{N_{TC}}} \text{ GeV} \]
\[ m_{\pi^0} \approx 240 \sqrt{\frac{4}{N_{TC}}} \text{ GeV} \]

where \( N_{TC} \) is the dimensionality of the TC representation of the technifermions. Any model containing these technipions should give mass estimates not very different from these.

In Table 2, we have listed the \( g^2 = 64 \) ground-state technihadrons with spin one. They are all approximately degenerate at \( \sim 900 \sqrt{\frac{\Lambda_{TC}}{N_{TC}}} \text{ GeV} \). They are the most interesting technihadrons because some of them can be produced in the very high energy colliders now being contemplated and because they provide a copious source of technipions.\(^36\) The neutral color singlet \( \omega_2 \) and \( \omega_1 \) have nonzero couplings to the photon and so appear as nearly degenerate s-channel resonances in \( e^+e^- \) annihilation. The color octet \( \omega_6 \) mixes with color-SU(3) gluons and should be produced in \( pp \) and \( pp \) collisions at \( \sqrt{s} \geq 2 \text{ TeV} \). In each of their respective settings, these neutral mesons provide the principal way to access the heavy technipion pairs \( P^+ P^0 \) and \( P^0 P^0' \).

Before closing this discussion, we mention that a seemingly slight modification of the Farhi-Susskind model leads to a very different population of its technicolor zoo. Suppose, instead, that we had assigned the technifermions \( U, D, N, E \) to a \( \overline{3} \) representation of the TC group whose symmetric product contains the TC-singlet representation.\(^37\) Then there would be a total of 135 technipions - the 63 already mentioned and 72 more. The additional technipions consist of 6 color-singlet dileptons (LL) of mass 50 - 80 GeV, 24 color triplet leptoquarks (LQ) and 6 color triplet diquarks (QQ), all of mass \( \sim 100 \text{ GeV} \). Similar complications occur in the spin-one sector. The point, of course, is that the technipion spectrum can be very rich and that details are highly model-dependent.

**TABLE 2.** Lowest-lying spin-one technihadrons in the Farhi-Susskind model. All have mass \( \sim 900 \sqrt{\frac{\Lambda_{TC}}{N_{TC}}} \text{ GeV} \). \( W^\pm \) and \( Z^0 \) are longitudinal components of gauge bosons. For unit renormalization, states should be divided by \( \sqrt{\lambda_{TC}} \).

<table>
<thead>
<tr>
<th>Vector Technihadron/Technifermion Content</th>
<th>Major Decay Modes</th>
</tr>
</thead>
<tbody>
<tr>
<td>Color Singlets</td>
<td></td>
</tr>
<tr>
<td>( \left[ \frac{1}{2} \right] \left[ \frac{3}{2} \right] ) (</td>
<td>U \rangle )</td>
</tr>
<tr>
<td>( \left[ \frac{1}{2} \right] \left[ \frac{1}{2} \right] ) (</td>
<td>D \rangle )</td>
</tr>
<tr>
<td>( \left[ \frac{1}{2} \right] \left[ 1 \right] ) (</td>
<td>N \rangle )</td>
</tr>
<tr>
<td>( \left[ \frac{1}{2} \right] \left[ 0 \right] ) (</td>
<td>E \rangle )</td>
</tr>
<tr>
<td>Color Triplet Leptoquarks</td>
<td></td>
</tr>
<tr>
<td>( \left[ \frac{1}{2} \right] \left[ \frac{3}{2} \right] ) (</td>
<td>U \rangle )</td>
</tr>
<tr>
<td>( \left[ \frac{1}{2} \right] \left[ 1 \right] ) (</td>
<td>E \rangle )</td>
</tr>
<tr>
<td>( \left[ \frac{1}{2} \right] \left[ 0 \right] ) (</td>
<td>N \rangle )</td>
</tr>
<tr>
<td>Color Octets</td>
<td></td>
</tr>
<tr>
<td>( \left[ \frac{1}{2} \right] \left[ 0 \right] ) ( \overline{3} )</td>
<td>( \pi )</td>
</tr>
</tbody>
</table>

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C. Couplings of Technipions and Higgses to Ordinary Matter

In most of the discussion that follows, we will not assume the scalar meson content of any particular elementary Higgs or technicolor model. In accord with our ignorance of the "correct" model, we shall allow for any number of neutral and singly-charged color-singlet scalar mesons. In the few instances that we specialize to the Farhi-Susskind model, that assumption will be made clear.

1. Elementary Higgses and Light Technipions

It is useful to discuss H± and H°'s together with p± and p°'s because their couplings to ordinary matter - quarks, leptons, electroweak bosons and gluons - are expected to be quite similar in some cases and completely different in others. The reason for this is that technipion couplings to light matter always occur through their constituent technifermions. Elementary Higgses, on the other hand, can have couplings to ordinary matter at the Lagrangian level, as in Eqs. (1) and (2), and some of these are expected to be much larger than the corresponding technipion couplings. If relatively light charged and neutral scalars are discovered, these differences should make it possible to determine whether they are elementary or composite.

The effective interaction for the coupling of the ith charged and neutral technipions, P± and P°, to ordinary fermions has the general form (assuming no light right-handed neutrinos)\(^{38}\)

\[
\mathcal{L}_{\text{eff}} = \sum_{i,s} \left[ \bar{U}_s \gamma^\mu P_i^+ \gamma_\mu P_i^0 - \bar{U}_s \gamma^\mu P_i^0 \gamma_\mu P_i^+ \right] + \text{h.c.} \tag{13}
\]

The sums in Eq. (13) extend over all flavors r,s of quarks and leptons; there is also an implicit sum over quark colors. The dimensionless couplings \(A_{rsi} \), etc. are \(\propto F^{-1}_{rsi} \) times a mass which is related in a complicated and model-dependent way to the current algebra mass of \(q_r, \bar{q}_s\) or of \(\tau_r, \tau_s\). Much of this complication is contained in unknown mixing-angle factors that arise from diagonalizing the mass matrices of quarks, leptons, technifermions and technipions. In general, \( A_{rsi} \neq B_{rsi} \) and so on.\(^{39}\) Thus, even though technipions are pseudoscalars so far as their strong TC interactions are concerned, their light-fermion couplings are parity-violating. Moreover, it is generally true that

\[
F_{rsi} A_{rsi} = F_{rsi} B_{rsi} = \sum m_{r_s} \text{ mass of the current algebra masses of } m_{r_s} \text{ and } m_{r_s}
\]

Very similar remarks apply to the couplings of H± and H°'s to fermions. They depend on the number of Higgs doublets in the model and how these are chosen to couple to the fermion doublets and singlets. Models in which the Higgses are expected to be less massive than several hundred GeV generally are constructed according to the "natural flavor conservation" guidelines of Glashow and Weinberg.\(^{40}\) In this class of models, all fermions having the same electric charge acquire their masses from the vacuum expectation value of only one of the Higgs doublets. (While a corresponding condition is not obviously necessary in TC theories, we shall effectively assume it for purposes of practical calculations of decay branching ratios.) Thus, in models with natural flavor conservation, the Higgs-fermion interaction looks like Eq. (13), but one can relate the coupling constants to fermion masses as follows:

\[
F_{rsi} A_{rsi} = (M K)_{rsi} A_{rsi}^+, \quad F_{rsi} B_{rsi} = (M K)_{rsi} B_{rsi}^-, \tag{14}
\]

and

\[
F_{rsi} A_{rsi}^0 = F_{rsi} B_{rsi}^0 = m_{r_s} A_{rsi}^0 \delta_{rsi}, \quad F_{rsi} C_{rsi}^+ = F_{rsi} D_{rsi}^+ = m_{r_s} C_{rsi}^0 \delta_{rsi}, \quad F_{rsi} B_{rsi}^0 = F_{rsi} Y_{rsi} = m_{r_s} B_{rsi} \delta_{rsi}. \tag{15}
\]

In Eqs. (14) and (15),

\[
K = \begin{pmatrix} K_{ud} & K_{us} & K_{ub} \\ K_{cd} & K_{cs} & K_{cb} \\ K_{td} & K_{ts} & K_{tb} \end{pmatrix} = \begin{pmatrix} 0.97 & -0.22 & -0.068 \\ 0.22 & 0.85 & 0.48 \\ 0.046 & 0.48 & 0.88 \end{pmatrix} \tag{16}
\]

is the usual Kobayashi-Maskawa\(^{41,42}\) matrix,

\[
M_u = \begin{pmatrix} m_u & m_c & m_t \end{pmatrix} = \begin{pmatrix} 5 \text{ MeV} \\ 1.2 \text{ GeV} \\ 20,25 \text{ GeV (assumed)} \end{pmatrix} \tag{17}
\]

are the diagonalized quark current algebra mass matrices\(^{43}\) and \(m_{r_s}\) is the mass of charged lepton \(\ell_r\). The
factors $A_i, B_i, \ldots, \beta_i$ are ratios of functions of Higgs-meson mixing angles. There are completely unknown. Because of these factors, the usual presumption that the Higgs coupling to $f + f'$ goes like the mass of the heavier fermion of the pair is little more than a guess. The presumption is greater still for technipions. Nevertheless, it is a reasonable one, worth hanging on to until we are proven wrong.

To give some idea of the ranges one might expect for Higgs/technipion branching ratios, we have calculated them for four different "models" for the couplings $A_{rsi}$, etc. The first two models are based on the elementary Higgs couplings just discussed. We use Eqs. (13)-(15) with $A^0 = C^0 = \beta_i = 1$ and

$A^\pm = B^\pm = \alpha_i = 1$ (Model 1) (18a)

$A^\pm = -B^\pm = \alpha_i = 1$ (Model 2) (18b)

The second class of models is a (necessarily) crude attempt to represent simply the general calculations of technipion-fermion couplings in Ref. 38. We choose all $P^0$ couplings and those of $P^\pm$ to $f^\pm u_d$ as in models 1 and 2; $P^+ - u_d d_s$ amplitudes are taken to be

$A^\pm_{rsi} = B^\pm_{rsi} = \sqrt{\frac{m_u - m_d}{m_r}} / F$ (Model 3) (19b)

$A^\pm_{rsi} = -B^\pm_{rsi} = \sqrt{\frac{m_r - m_s}{m_r}} / F$ (Model 4) (19b)

For handy reference, the rate for $P^+_1 - u_d d_s$ is

$\Gamma(p_1^+ - u_d d_s) = \frac{3p_{rsi}}{16m_r} \left( |A^\pm_{rsi} + B^\pm_{rsi}|^2 \left( 1 - \frac{(m_r - m_ds)^2}{m_p^2} \right) 
+ |A^\pm_{rsi} - B^\pm_{rsi}|^2 \left( 1 - \frac{(m_r - m_ds)^2}{m_p^2} \right) \right), (20)$

where $p_{rsi}$ is the quark momentum and the factor of 3 is due to the sum over quark colors. The factor of 3 is absent for decay to $f^\pm u_d$.

The branching ratios for a range of charged and neutral Higgs/technipion masses are listed in Tables 3 and 4. There we see the general tendency, built into the models, for $P^+/H^\pm$ to decay into the heaviest possible quark-antiquark pair. Note, however, that in models 2 and 4 this tendency is sometimes foiled by phase space limitations. For $m_\pm < m_t$, the branching ratio for $P^+ - r^+ u_d$ is fairly large, 10-20%; it generally drops to $\sim 1-4\%$ once decay to $t\bar{t}$ is allowed. We stress that

### TABLE 3. Branching ratios for $P^+/H^\pm$ decay to fermion + antifermion. The four values for each decay mode and mass refer to the four "models" for the couplings described in the text.

<table>
<thead>
<tr>
<th>$P^+/H^\pm$ Decay Mode ($m_\pm = 20$ GeV)</th>
<th>Branching Ratio (%) for $m_\pm = 7.5$ GeV</th>
<th>10 GeV</th>
<th>20 GeV</th>
<th>30 GeV</th>
<th>40 GeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\tau^+ \nu$</td>
<td>23.0, 26.5; 15.5, 16.0; 1.0, 1.0; 0.5, 0.5</td>
<td>18.5, 20.0; 9.0, 9.5; 1.5, 2.5; 1.0, 1.0</td>
<td>14.5, 23.5; 11.0, 14.0; 9.0, 9.5; 1.5, 2.5; 1.0, 1.0</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$c_d$</td>
<td>24.5, 28.5; 19.5, 21.0; 1.0, 1.0; 0.5, 0.5</td>
<td>15.5, 16.0; 9.0, 9.5; 1.5, 2.5; 1.0, 1.0</td>
<td>20.0, 22.5; 16.0, 16.0; 1.0, 1.0; 0.5, 0.5</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$c_b$</td>
<td>50.5, 43.0; 61.0, 57.5; 3.5, 4.0; 3.5, 4.0</td>
<td>60.0, 57.5; 57.5, 60.0; 3.5, 4.0; 3.5, 4.0</td>
<td>79.0, 66.0; 86.0, 80.0; 87.0, 86.5; 13.0, 23.5; 8.5, 10.5</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$t\bar{t}$</td>
<td>21.0, 26.5; 21.0, 23.0; 3.0, 5.0; 3.0, 4.0</td>
<td>21.0, 26.5; 21.0, 23.0; 3.0, 5.0; 3.0, 4.0</td>
<td>73.5, 66.5; 75.5, 73.0; 82.0, 68.0; 87.5, 83.5</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>$P^+/H^\pm$ Decay Mode ($m_\pm = 25$ GeV)</th>
<th>Branching Ratio (%) for $m_\pm = 30$ GeV</th>
<th>40 GeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\tau^+ \nu$</td>
<td>2.5, 3.0; 8.0; 0.5, 0.5; 1.0, 1.5</td>
<td></td>
</tr>
<tr>
<td>$c_d$</td>
<td>2.5, 3.0; 1.5, 3.0; 0.5, 0.5; 0.5, 0.5</td>
<td></td>
</tr>
<tr>
<td>$c_b$</td>
<td>10.5, 14.0; 36.5, 75.5; 2.0, 2.0; 9.5, 14.0</td>
<td></td>
</tr>
<tr>
<td>$t\bar{t}$</td>
<td>29.5, 40.0; 3.0, 6.0; 21.5, 24.0; 3.0, 4.5</td>
<td></td>
</tr>
<tr>
<td>$t\bar{b}$</td>
<td>54.0, 38.5; 55.5, 7.5; 75.0, 72.0; 86.0, 79.5</td>
<td></td>
</tr>
</tbody>
</table>
TABLE 4. Branching ratios for $\PZ^0/\H^0$ decay to fermion + antifermion. The choice of signs in Eq. (15) makes no significant difference on the values obtained. Note that, for $m_\PZ > 2m_t$, $\PZ^0/\H^0 \to \PZ^+\PZ^-$ is expected to be the dominant decay mode.

<table>
<thead>
<tr>
<th>$\PZ^0/\H^0$ Decay Mode</th>
<th>$m_\PZ = 2.5$ GeV</th>
<th>10 GeV</th>
<th>20 GeV</th>
<th>30 GeV</th>
<th>40 GeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\mu^+\mu^-$</td>
<td>11</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$\tau^+\tau^-$</td>
<td></td>
<td>41</td>
<td>5</td>
<td>5</td>
<td>5</td>
</tr>
<tr>
<td>$\mu^-\mu^+$</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$c\bar{c}$</td>
<td>89</td>
<td>58</td>
<td>7</td>
<td>7</td>
<td>6</td>
</tr>
<tr>
<td>$b\bar{b}$</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

this $\tau^+\PZ^-\PZ^+$ mode, which is an important signal for charged scalars, is sizable because we have taken $\alpha$ comparable in magnitude to $A_T^\PZ$ and $B_T^\PZ$. The real world may be quite different! Finally, the total width of a charged scalar is $\sim 0.01 - 0.1$ MeV for $m_\PZ < m_t$ and $\sim 1 - 10$ MeV for $m_\PZ > m_t$. Neutral scalar widths are comparable.

Couplings to Electroweak Gauge Bosons and QCD Gluons. Elementary Higgs mesons couple directly to $\PZ$, $\PW^\pm$ and $\PZ^0$ in the electroweak Lagrangian and so these interactions are fairly large. The most important terms for $H^\pm$ and $H^0$ production are trilinear in the Higgs and electroweak fields. For any model containing two or more Higgs doublets, the nonvanishing elementary two-Higgs production amplitudes are:

\[
A(\PZ(q) \to H_1^+(p) + H_2^+(k)) = i e R_{ij} \epsilon(q) \cdot (p-k)
\]
\[
A(\PW^0(q) \to H_1^+(p) + H_2^+(k)) = i e R_{ij} \cos \theta_\PZ \epsilon(q) \cdot (p-k)
\]
\[
A(\PW^0(q) \to H_1^0(p) + H_2^0(k)) = -i e \sin \theta_\PZ \epsilon(q) \cdot (p-k)
\]
\[
A(\PW^+(q) \to H_1^0(p) + H_2^0(k)) = i e \frac{1}{2} R_{ij} \epsilon(q) \cdot (p-k)
\]

Here, $\epsilon(q)$ is the gauge boson's polarization 4-vector, and $R_{ij}^{\PZ}$ and $R_{ij}^{\PW}$ are functions of the Higgs meson mixing angles. $R_{ij}^{\PZ}$ vanishes if $H_j^0 = H_i^0$. Note that, apart from the mixing angle factors, the coupling constants in Eq. (21) are just the electromagnetic, neutral weak and charged weak current charges $Q$, $T_3$ - $\sin^2 \theta_\PZ Q$ and $g/2$ of the appropriate Higgs mesons.

The nonzero elementary amplitudes for bremsstrahlung of a Higgs meson from an electroweak boson are (cf. Eq. (22)):

\[
A(Z^0(q) \to Z^0(p) + H_1^0(k)) = \frac{2 e u_{ij}^{\PZ} \sum_{\PZ} R_{ij}^{\PZ} v_{ij}^{\PZ} \epsilon(q) \cdot \epsilon(p)}{\sin^2 \theta_\PZ}
\]
\[
A(\PW^0(q) \to \PW^0(p) + H_1^0(k)) = \frac{e u_{ij}^{\PW} \sum_{\PW} R_{ij}^{\PW} v_{ij}^{\PW} \epsilon(q) \cdot \epsilon(p)}{\sin \theta_\PZ}
\]

where $v_{ij}/2$ is the vacuum expectation value of the $i$th unmixed Higgs doublet $\phi_i$, appearing in the Lagrangian, $R_{ij}^{\PZ}$ is the matrix diagonalizing the neutral-Higgs mass matrix and $F_\PZ = \sqrt{\frac{\mu_\PZ^2}{v_{ij}^{\PZ}}} \approx 250$ GeV. Apart from this sum rule on the $v_{ij}$, they are arbitrary. For making numerical estimates, we shall assume they are roughly equal in magnitude and not less than $\sim 100$ GeV (corresponding to $\leq 6$ Higgs doublets). Note that there is no elementary coupling of $\PW^\mu$, $\PW^\nu$, and $\PW^\tau$ to physical charged Higgses. Thus, while $H_1^0 \to H_2^0 + \gamma$ is expected to be the dominant decay mode.

We turn now to the technipions. All of their interactions with electroweak and QCD gauge bosons proceed through technifermion loops. At energies well below their characteristic scale, technipions are essentially point-like and these interactions are reliably calculated using well-known current-algebraic methods and the known technifermion-gauge boson couplings in, e.g., Eq. (6). (An example of the form factor effects that appear at production energies $\sim m_\PZ$ will be given in Secs. IV A, B.) Current algebra gives the $\gamma$, $\PW^\mu$, and $\PW^\nu$ couplings to a pair of technipions as the charges $eQ$, $eT_3 - \sin^2 \theta_\PZ Q$, and $g/2$, respectively, of their constituent technifermions. Thus,

\[
A(\PZ(q) \to \PW^\mu_1 + \PW^\mu_2) = i e \delta_{ij} \epsilon(q) \cdot (p-k)
\]
\[
A(Z^0(q) \to \PW^\mu_1 + \PW^\mu_2) = i e \delta_{ij} \cos \theta_\PZ \epsilon(q) \cdot (p-k)
\]
\[
A(Z^0(q) \to \PW^\nu_1 + \PW^\nu_2) = 0
\]
\[
A(W^+(q) \to \PW^\mu_1 + \PW^\mu_2) = i e \frac{1}{2} \sin \theta_\PZ \delta_{ij} \epsilon(q) \cdot (p-k)
\]

The first two amplitudes are the same as in Eqs. (21) because of the assumed electroweak quantum numbers of technifermions (Eq. (6)). Insofar as their assignment to SU(2), doublets and singlets is required for a phenomenologically consistent TC theory, these amplitudes are model-independent. So is the third of Eqs. (23), which is our first important difference with elementary Higgs models. The fact that it vanishes in the current-algebra approximation means that it is suppressed by powers of $(m_\PZ^0/\Lambda_{\text{TC}})^2$ relative to other amplitudes in...
Eq. (23). The quantity $R^j_1$ in the last equation is a model-dependent mixing angle factor. In the simplest Farhi-Susskind model, $R^j_1 = 1$ for $\rho^0$ and it vanishes for $\rho^1$.

More striking departures from the elementary Higgs scenarios occur for the coupling of two gauge bosons to one technipion. The calculation of these are similar to that of the triangle anomaly graph for $\sigma^0 - \gamma\gamma$. The coupling of $\sigma$ to gauge bosons $B_1$ and $B_2$ is given by

$$\sqrt{2} S_{PB_{1B_2}}^{PB_{1B_2}} = \frac{g_{PB_{1B_2}}^2}{8\pi^2} e^{\mu
u\lambda\delta} \epsilon^{\alpha\beta\gamma\delta} \frac{1}{2} \bar{q}^\alpha p^\beta \gamma^\nu \sigma^\mu \gamma^\delta \epsilon^{\gamma\delta\epsilon\zeta} \epsilon_{\epsilon\zeta\lambda\mu}$$

(24)

where the triangle anomaly factor $S_{PB_{1B_2}}^{PB_{1B_2}}$ is

$$S_{PB_{1B_2}}^{PB_{1B_2}} = \frac{1}{2} g_{PB_{1B_2}}^2 \text{Tr} \left( Q_p \left( Q_{1}, Q_{2} \right) \right) .$$

(25)

Here, $g_{PB_{1B_2}}$ are the $B_1 B_2$ gauge couplings and $Q_{1,2}$ their gauge charges. $Q_p$ is the chiral charge of technipion $P$. The contributions from different gauge boson helicity states are summed separately in the trace. For the Farhi-Susskind model, we have listed the more important values of $S_{PB_{1B_2}}^{PB_{1B_2}}$ in Table 5. For more complicated models, the anomaly factors are multiplied by different group-theoretic numbers and functions of technipion mixing angles. In any case, the values in Table 5 will not seriously mislead us.

From Table 5 and Eq. (22), we estimate

$$\left| \frac{A_{PB_{1B_2}}^{\sigma^0 - \gamma\gamma}}{A_{PB_{1B_2}}^{\sigma^0 - \gamma\gamma}} \right|^2 \approx 10^{-4} - 10^{-5} .$$

(26)

Folding this into the rate for $e^+ e^- \rightarrow \gamma\gamma$ in Fig. 3, it is clear that the $\rho^0$ cannot be produced this way. To put it bluntly, observation of a narrow neutral scalar in this process rules out technicolor.

### 2. Heavy Technipions, $P_{LQ}$ and $P_8$

Now that we have lavished so much time on light technipion couplings and, I hope, struck the correct note of uncertainty, we can deal with the heavy technipions in short order. We will discuss only the color-triplet leptoquarks $P_{LQ}$ and octets $P_8$ of the simplest Farhi-Susskind model. We are even more unsure of their couplings to fermions (not to mention which fermion channels lie open to them, nor their very existence) than we are for their lighter cousins, and there seems little point in making more conventional-wisdom guesses. This applies all the more strongly to the di-leptons and diquarks found in the real-TC-representation version of the model, so we will not discuss them here. The interested reader can find them discussed in Refs. 30 and 11.

In view of their color SU(3) quantum numbers, we expect that the most important fermionic couplings of $P_8$ and $P_{LQ}$ are of the form $A_8^{\alpha \beta \gamma \delta} \bar{q}_{\gamma} \gamma^\lambda a_{\alpha} r^\alpha_{\beta} v^\delta_{\gamma} \delta_{\lambda}$. Where $A_8^{\alpha \beta \gamma \delta}$ is color-SU(3) matrix, and $A_{LQ}^{\alpha \beta \gamma \delta} \bar{q}_{\gamma} \gamma^\lambda a_{\alpha} r^\alpha_{\beta} v^\delta_{\gamma} \delta_{\lambda}$. Note that the latter coupling requires that $Y = 1/6$ in Eq. (6). The coefficients $A_8^\gamma_{\alpha}$ are expected to be of order (heavier fermion mass $x F_{m}^{-1}$). As one almost certainly too naive example,

<table>
<thead>
<tr>
<th>Vertex $PB_{1B_2}$</th>
<th>Anomaly Factor $S_{PB_{1B_2}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\rho^0 \gamma\gamma$</td>
<td>$e^2 \frac{4}{\sqrt{6}} N_{TC}$</td>
</tr>
<tr>
<td>$Z_0^0 \gamma\gamma$</td>
<td>$e^2 \frac{2 \cos^2 \theta_w}{\cos^2 \theta_w} \frac{1}{\sqrt{6}} N_{TC}$</td>
</tr>
<tr>
<td>$Z_0^0 \gamma\gamma$</td>
<td>$e^2 \frac{2(1 - 4 \sin^2 \theta_w)}{\sin 2 \theta_w} \frac{1}{\sqrt{6}} N_{TC}$</td>
</tr>
<tr>
<td>$W^+ W^-$</td>
<td>0</td>
</tr>
<tr>
<td>$\rho^0 \gamma\gamma$</td>
<td>$g_{8}^\delta_{ab} N_{TC}$</td>
</tr>
<tr>
<td>$Z_0^0 \gamma\gamma$</td>
<td>$e g_{8} \cot^2 \theta_w \delta_{ab} N_{TC}$</td>
</tr>
<tr>
<td>$\rho^0 \gamma\gamma$</td>
<td>$g_{8}^\delta_{abc} N_{TC}$</td>
</tr>
<tr>
<td>$Z_0^0 \gamma\gamma$</td>
<td>$e g_{8} \delta_{ab} N_{TC}$</td>
</tr>
<tr>
<td>$\rho^0 \gamma\gamma$</td>
<td>$g_{8} \delta_{ab} N_{TC}$</td>
</tr>
<tr>
<td>$Z_0^0 \gamma\gamma$</td>
<td>$e g_{8} \delta_{ab} N_{TC}$</td>
</tr>
</tbody>
</table>

Barbiellini, et al.11 take as the $P_8$ couplings appropriate to the Farhi-Susskind model:

$$f \sum_{n} \sum_{r,s} \sum_{8a} \sum_{8b} \left\{ \frac{1}{2} \sum_{r,s} \sum_{8a} \sum_{8b} \left( \frac{1}{2} \right) \frac{1}{2} \frac{1}{2} \frac{1}{2} \right\}$$

(27)
are \( \rho_8^0 - (\xi \xi)_8 \), \( \rho_8^+ - (\xi \xi)_8 \) and so on. Likewise, if \( Y = 1/6 \), \( h_{TB} = -\gamma T \), \( h_{BB} = -\nu_3 T \), etc. are presumed to be the largest leptoquark-fermion couplings. Since leptoquarks are color triplets these, in fact, are their dominant decay modes.

The couplings of a single gauge boson to a pair of heavy technipions are calculated using current algebras and gauge invariance, as before. We have listed the most interesting ones in Table 6 for the Farhi-Susskind model. However, the reader should note that these amplitudes generally are not directly useful for calculating production rates. At the center-of-mass energies required to pair-produce them, the amplitudes are expected to be dominated by spin-one technihadrons (see Sec. IV).

The color octet technipions can couple to a pair of gauge bosons. The amplitude is given in Eq. (24) and the anomaly factors \( S_{BB}B_1B_2 \) included in Table 5. The decay rate for \( g_8 - B_1B_2 \) is \[ \Gamma(g_8 - B_1B_2) = \frac{\alpha_s}{128\pi} \left( \frac{2 S_{BB}B_1B_2}{4\pi^2 f_g} \right)^2. \] (28)

We can see from this that the \( P_{g8}^8 \)-octet has an appreciable coupling to two gluons. Furthermore, using Eq. (24), we have \[ \frac{\Gamma(P_{g8}^8 \rightarrow GG)}{\Gamma(P_{g8}^8 \rightarrow (\xi \xi)_8)} = \frac{\alpha_s}{3} \left( \frac{2 m_g^2}{m_t} \right)^2 \frac{2 N_{TC}}{4} \approx 0.25 \frac{N_{TC}}{4}, \] (29)

where, again, \( N_{TC} \) is the dimensionality of the TC representation of techniquarks and we used \( \alpha_s(m_t) \approx 0.1 \) and \( m_g \approx m_t = 12 \sqrt{\frac{4}{N_{TC}}}. \) Thus, \( P_{g8}^8 \) (a.k.a. \( \eta_g \)) decays predominantly to \( t\bar{t} \) and to two gluons. Since \( N_{TC} \) is unknown, we cannot be more specific than that.

This concludes our discussion of technipion and Higgs couplings to ordinary matter. In the next section we use these results to estimate the most likely means for their production and detection.

### Production and Detection of Higgs Mesons and Technipions in Colliders

It is obvious that exploration of the scalar sector of electroweak interactions requires high-energy colliders. It is equally important that these colliders have consistently high luminosity. In this section we discuss in turn \( e^+e^- \) colliders, both at and below the \( Z^0 \) and colliding linacs at ultrahigh energies, the hadron-hadron colliders Isabelle and TEV I and ep colliders. We assume luminosities of \( 2 \times 10^{31} \) for \( e^+e^- \) at the \( Z^0 \) and \( 10^{33} \) for colliding linacs, \( 10^{32} \) for pp at \( s = 800 \) and \( 10^{32} \) for ep colliders at \( s = 300-1000 \) GeV. Except in a few cases, we shall not dwell much on the all-important question of backgrounds; the reader is referred to the reports of the appropriate collider groups for details.

These are the principal conclusions of our study:

1. Overall, \( e^+e^- \) colliders are the most effective and versatile. The existing and proposed machines can discover all scalars, elementary or composite, lighter than \( \sim 45 \) GeV. Taken together, studies at toponium and the \( Z^0 \) can decide between the elementary and technicolor scenarios. Ultrahigh energy colliding linacs copiously produce spin-one technihadrons as s-channel resonances and these decay exclusively to technipions, including the heavy colored ones.

2. Hadron colliders such as Isabelle and TEV I produce at detectable levels only composite scalars, specifically those which couple strongly to gluons. They do the best (and soonest) job of singly-producing \( P_{g8}^8 \) and, possibly, \( P_0^0 \). Colored technipion pairs may be produced at observable rates at TEV I, though not at Isabelle. Development of efficient high-resolution vertex detectors are essential for success in all these searches.

3. ep colliders are limited to single-production of heavy colored technipions such as \( P_{g8}^8 \) and \( P_0^0 \). However, observable rates probably require \( s \gg 1 \) TeV.

4. In almost all cases, accurate determination of scalar meson masses requires good jet identification, discrimination and mass resolution. The reason, of course, is that \( P^+ \) and \( H^0 \) are expected to decay mainly to \( c\bar{b} \) or \( t\bar{t} \), \( P_0^0 \) to \( t\bar{t} \), etc. and these heavy quark jets will have to be distinguished from all others.

#### A. e^+e^- Colliders

The entire event structure is useful for detecting scalar mesons produced in \( e^+e^- \) annihilation. This fact makes \( e^+e^- \) the simplest and surest method for discovering them. In the following we discuss Higgs and technipion production in continuum \( e^+e^- \) annihilation, at toponium, at the \( Z^0 \) and at ultrahigh energies. We spell out how searches at toponium and the \( Z^0 \) can distinguish between the elementary and composite scenarios for the scalar sector and among various possibilities within each scenario.

### Continuum Production

We noted in Sec. III that technipions are essentially point-like below \( s \approx M_{TC} \). Furthermore, the \( \gamma Z^0 - P_{P}^{T^*} \) and \( \gamma Z^0 - H^0 \) amplitudes are identical and model-independent (Eqs. (21) and (23)). Thus, the charged-pair production cross section appropriate to this decade's generation of \( e^+e^- \) collider is

\[
\frac{d\sigma(e^+e^- \rightarrow P_{P}^{T^*} \text{ or } H^0)}{d(cos \theta)} = \frac{\alpha_s^2}{8s} l_p^2 \sin^2 \theta \left[ |f_R(s)|^2 + |f_L(s)|^2 \right]^2
\] (30)
where $\beta = \sqrt{1 - 4m_t^2/s}$ is the meson velocity, $\theta$ the c.m. production angle and

$$f_R(s) = 1 - \frac{\cos 2\theta}{2} \frac{s}{s - \mu_{+}^{2} + \mu_{-}^{2}} \Gamma_{\pm}$$

$$f_L(s) = 1 + \frac{\cos 2\theta}{2} \frac{s}{s - \mu_{+}^{2} + \mu_{-}^{2}} \Gamma_{\pm}$$

(31)

We use $\mu_{\pm} = 93.0$ GeV and $\Gamma_{\pm} = 2.92$ GeV in calculations. Two points that stand out in Eq.(31) are:

1. The continuum cross section is rather small; e.g., $R_\pm(s) = (\sigma(e^+e^- \rightarrow P^-P^+\pm/\sigma(e^+e^- \rightarrow \mu^+\mu^-)) \approx 0.14$ for $m_t = 10$ GeV and $s = 35$ GeV. It should be remembered, however, that each distinct charged pair ultimately contributes 1/4 unit to $R_\pm$. (2) The angular distribution peaks at 90°. This fact is useful in searching for striking, though rare, signals and in event reconstruction.

The most striking signature of $P^-\tau^-/H^-\tau^-$ production should be $P^-\tau^-$-hadrons ($c\bar{s}$, $c\bar{c}$, etc.) together with $P^- \rightarrow \tau^-\nu_\tau$. From Table 3, we would expect this class of decays to occur 20-60% of the time for $m_\pi < m_t$ and 1-10% of the time for $m_\pi > m_t$.

In these events, at least ~3/8 of the total energy is missing in neutrinos. The degree of jettiness of the hadronic decay products of $P^-$ depends both on $\beta$ and the decay channels open to it. This fact requires shifting through different event samples for high- and low-mass searches. In 35% of these events, there will be a relatively low-energy, wide-angle isolated $e$ or $\mu$ from $\tau^-$-decay. Such events, though rare, are practically background-free. (This will be discussed in more detail when we take up the $\tau^+$.) In another 20% of these events, a single charged pion results from $\tau^+ \rightarrow \mu^+\nu_\mu$ and $\nu_\mu$. This, too, gives an isolated charged track if $\beta$ is high enough.

Several groups at PETRA and PEP have already searched for charged scalars in both the $(\tau^\nu)(qq)$ and $(\tau\nu)(c\bar{s})$ modes. Upper limits on $B(P^\pm/H^\pm \rightarrow \tau^\nu\nu)$ ranging from 4 to 11% were obtained for 4 GeV $< m_\pi < 12$ GeV (see Fig. 7). Only model 3 of Table 3 falls under these limits. Nevertheless, it is important to look for the $(qq)(qq)$ modes to completely rule out $m_\pi < 12$ GeV. It is worth remarking here that the all-hadrons mode most likely contains four charmed quarks, another striking signal if the technology exists to detect their decays.

Finally, continuum production of $P^0$ or $H^0$ is negligible. These neutral mesons are best searched for at toponium and the $Z^0$, as we shall see.

Production at Toponium. Decays of the toponium ground state, $\zeta = 1 S_1(t\bar{t})$, may be a copious source of both charged and neutral scalars. If $m_t > m_\pi + m_\tau$, the semiweak transition $t \rightarrow P^\mu/P^\tau + b$ is expected to be the dominant decay mode of the top quark. The decay rate is given by (for each charged meson)

$$\Gamma(t \rightarrow P^\pm b) = \frac{p}{64\pi \mu_t^2} \partial_{b}^{2}\left[\left(\text{m}_t - \text{m}_b\right)^2 - \text{m}_{\pm}^2\right]$$

$$+ \left[\partial_{b}^{2} - \text{m}_t^2\right] \left[\left(\text{m}_t + \text{m}_b\right)^2 - \text{m}_{\pm}^2\right]$$

(32)

For $m_t = 25$ GeV and $m_\tau = 15$ GeV, the four "models" of Sec. III (Eqs.(14)-(19)) give $\Gamma(t \rightarrow P^\pm b) \beta > 80$ keV and 460 keV. The width of $\zeta$ expected in the standard quarkonium model is $\approx 60$ keV. Thus, the first unmistakable signal that $m_t > m_\tau + m_\pi$ will be the extra large toponium width. Further, a large fraction of $\zeta$ decays will be

$$\zeta \rightarrow P^\pm b\bar{b} \rightarrow 4 \text{ charmed quarks} + \tau^\pm$$

(33)

This would be impossible to miss with adequate vertex detectors.

Toponium is a good place to search for elementary $P^0$'s, in the decay $\zeta \rightarrow H^0\tau^-$. If the technicolor scenario is correct, it is the only place $P^0$'s can be found. (Remember that $m_{P^0} < 40$ GeV is almost certainly less than $m_\tau$.) The rate, relative to $\zeta \rightarrow e^+e^-$, is given by Eq.(4) times unknown mixing angle factors. Not all these factors can be much less than unity. Thus, if $P^0$'s exist, they should be found in its radiative decays unless the branching ratio is much reduced by the existence of light charged scalars. We should always be faced with such alternatives!

Production at $Z^0$. Two crucial tests for technicolor come at the $Z^0$. First, if charged technipions exist, then $\mu_\tau > 2\mu_{\pm} + \mu_\tau$, and they should be pair-produced rather copiously there. For each distinct charged pair of scalars, Eqs.(21),(23), and (30) give
neutral technipions cannot (see Eq. (26)). Assuming no
\[ \sigma(e^+ e^- \to Z^0 \to P^+ P^- \text{ or } H^+ H^-) \]
\[ = \frac{2.3 \cos^2 \theta_W}{6\pi} \frac{(1 - 4 \sin^2 \theta_W + 8 \sin^4 \theta_W)}{\sin^2 \theta_W} \]
\[ = 428 \beta^3 \text{pb} \] \tag{34a}
and
\[ \Gamma(Z^0 - P^+ P^- \text{ or } H^+ H^-) = \frac{\alpha_\mu^3 \beta^3}{12 \cot^2 \theta_W} = 0.89 \times 10^{-2} \beta^3 \Gamma_z. \] \tag{34b}
The angular distribution is proportional to $\sin^2 \theta$.

As in the case of continuum production of charged scalars, the simplest signal to look for is $e$ or $\mu$ + hadrons + missing energy. Unless the scalar is quite heavy, $\sim 40$ GeV, the hadrons should form one or two distinct jets and the lepton should be well isolated. Assuming $10^7$ $Z^0$'s per year, Table 7 gives the expected annual yield of $P^+ P^-$ or $H^+ H^-$ and the number of these that end up in the signal mode. The $\tau^+ \tau^-$ branching ratios used are an average of the appropriate ones in Table 3. The principal backgrounds to the $e/\mu$ + hadrons signal are discussed by Kagan.\textsuperscript{51} They are $Z^0 - \tau^+ \tau^-$, two-photon processes and $Z^0 - \ell^+ \ell^-$. By requiring that the charged hadron multiplicity be $\geq 6$, that the missing momentum point into the detector, that $e$ or $\mu$ not be within $\sim 60^\circ$ of another particle and other kinematic cuts, Kagan finds that these backgrounds are reduced below 0.03 pb, 0.1 pb and 0.1 pb, respectively. At most 30-50% of the signal is lost by these cuts. Furthermore, $m_\tau$ can be determined to $\sim 10\%$ with existing detector technology.\textsuperscript{52} Thus, the $Z^0$ can be used to discover $P^\pm$ or $H^\pm$ up to 40 GeV in mass.

If $m_\tau > m_e + m_\mu$, another copious source of charged scalars will be $Z^0 - \tau^+ \tau^- e^+ e^-$. Of course, this mode is much less clean than the simple $Z^0 - P^+ P^-$ decay. But, since it could yield $\sim 10^6$ $P^+ P^-$ events per year, detecting $P^\pm$ in this way clearly warrants further careful study.

The second crucial test for technicolor involves the search for neutral scalar mesons in $Z^0$ decays. As we have stated before, neutral elementary Higgses arising from either single- or multi-doublet models can be found by missing-mass techniques in $e^+ e^- \to Z^0 - e^+ e^- X$; neutral technipions cannot (see Eq. (26)). Assuming no drastic suppression by mixing-angle factors, this search mode is sensitive to neutral Higgs masses up to $\sim 45$ GeV and, so, is the method of choice at the $Z^0$.\textsuperscript{19}

Another decay which leads to neutral Higgses, but not technipions, is $Z^0 \to H^0 H^0$ ($i \neq j$). From Eq. (21), the rate is
\[ \Gamma(Z^0 - H^0 H^0)^{1/2} \left( \frac{\beta^3}{3 \sin^2 \theta_W} \right) = 0.23 \left( \frac{\beta^3}{3 \sin^2 \theta_W} \right) \] \tag{35a}
A search for this decay would be more difficult than for $Z^0 - P^+ P^- H^+ H^-$. Naively, we expect the major H decay modes to be $\ell^+ \ell^-$ (if allowed), $b \bar{b}$, $c \bar{c}$ and $\tau^+ \tau^-$.\textsuperscript{33} Probably the best signal is $H^0 - \tau^+ \tau^-$, leading to $e^2$ or $\mu^2 +$ hadrons + missing energy ($R_{\mu\mu} = 3/16 \mu_\mu$). I am not aware of any careful study of the best signals and attendant backgrounds for this two-Higgs decay of $Z^0$; one should be carried out.

To sum up, scalar meson searches at the $Z^0$ can distinguish between the various elementary and composite scenarios for the Higgs sector. If neutral scalars are found in either of the two decay modes just discussed, then technicolor is definitely ruled out. The same is true if charged scalar pairs are not found in $Z^0$ decays. Observation of $Z^0 - H^0 H^0$ tells us that there are 2 or more elementary Higgs doublets. Finally, observation of charged, but not neutral, scalars at the $Z^0$ very strongly suggests that the technicolor scenario is correct. We can take it as proved, moreover, if a neutral scalar also is seen in radiative toponium decays.

**Production at Ultrahigh Energies.** We assume for this discussion the eventual existence of an $e^+ e^-$ colliding linac capable of covering the energy range $\sqrt{s} = 0.5 - 1.5$ TeV with an average luminosity of $10^{33}$ cm$^{-2}$ sec$^{-1}$ for $10^7$ sec per year. One unit of $R$ at 1 TeV amounts to 865 events per year.\textsuperscript{54} It is likely that only the technicolor scenario can provide a spectacularly obvious and simple signal of scalar mesons in such a machine, namely the production of technihadrons as s-channel resonances and their subsequent decay to technipion pairs. Moreover, this is the only significant mechanism for producing heavy colored technipions in $e^+ e^-$ annihilation. For definiteness, we restrict our discussion to the mesons listed in Tables 1 and 2 for the Farhi-Susskind model. Other models should give at least qualitatively similar results.

At energies of order $\Lambda_T$, technipions can no longer be considered point-like, and form factors enter their production amplitudes. In ultrahigh energy $e^+ e^-$ annihilation, these form factor effects should be adequately described by vector-technimeson dominance. The lowest lying technihadrons coupling to $e^+ e^-$ through the $\gamma$ and $Z^0$ are $P^2$ and $W^3$ of Table 2. Note that they are nearly degenerate at $m_{W^3} \approx 900 \sqrt{\frac{M}{\Lambda_T}}$ GeV. Each of these decay to a pair of charged technipions and a pair of logitudinal $W^3$'s, $W^3 W^3$. The pair-production cross section has been calculated by Peskin.\textsuperscript{52} Relative to the point cross section, $4\pi \alpha/3$, he obtained
\[ R_{TC}(s) = \left( \frac{\alpha - \beta^2}{4} \right) \frac{1}{I_{TC}^2} \times \left( \frac{1}{s - m_{W^3}^2} \right)^2 + \left( \frac{g_s^2}{\Lambda^2} \right) \frac{1}{I_{W^3}^2} \left( \frac{1}{s - m_{W^3}^2} \right). \] \tag{37}
N \ F = 8 \text{ technifermion flavors. Notice that Peskin has used an s-dependent width in Eq.}(37).

In closing, we remind the reader that the results presented here were for one specific model—one which we know cannot be correct in detail. Thus, our discussion should be viewed as only exemplary of what to expect if technicolor is correct. If it is, the details may be quite different, but there must be spin-one resonances in the TeV region of e+e- annihilation and these will decay to technipion and \( \nu \bar{\nu} \) pairs.

B. Hadron-Hadron Colliders

Exploration of the scalar sector by pp and \( \bar{p}p \) colliders is limited to those mesons which couple relatively strongly to gluons. As we have stated and will re-emphasize here, this means that these colliders cannot see any elementary Higgs mesons and they will have to depend on the essential correctness of the technicolor scenario to make a positive discovery here. The most promising cases for Isabelle and TEV I are two-gluon fusion of \( \rho_0 \) and possibly, \( \rho_8 \). Furthermore, production of heavy colored technipion pairs may be observable at TEV I, though probably not at Isabelle. Let us discuss single- and pair-production of technipions in turn.

**P\( ^0_8 \) Production.** In principal, of course, hadron colliders can produce the elementary H\( ^0 \) and the technipions \( P0^0 \) and \( P0^8 \), and it is instructive to compare the three. The differential cross section for production of a scalar meson S\( ^0 \) of mass \( m_S \) by gluon fusion is\(^{10,35} \)

\[
\frac{d\sigma}{dy}(\text{pp or } \bar{p}p \rightarrow S^0 + X) = \frac{\pi^2}{8} \left( \frac{G(x)^2}{m_S^3} \right) \frac{1}{m_S^3} \frac{1}{m_S^3} \frac{1}{m_S^3} \sim \frac{1}{m_S^3}
\]

where \( G(x) \) is the gluon distribution function, \( x_1 \) and \( x_2 \) are the fractions of longitudinal momentum carried by the gluons and \( y \) is the rapidity of S\( ^0 \):

\[
y = \frac{1}{2} \ln \left( \frac{E_0 + P_{0L}}{E_0 - P_{0L}} \right) = \frac{1}{2} \ln \frac{x_1}{x_2};
\]

\[
x_1 = \frac{m_S^3}{m_S^3} e^y, \quad x_2 = \frac{m_S^3}{m_S^3} e^{-y}.
\]

Ignoring mixing-angle factors that would arise in any but the simplest models, the two-gluon widths are (see Ref. 20, Eq. (18) and Table 5):

\[
\Gamma(S^0 \rightarrow \rho_8 \rho_8) = \frac{G(x)^2}{m_S^3} \frac{1}{m_S^3} \frac{1}{m_S^3} \frac{1}{m_S^3} \sim \frac{1}{m_S^3}
\]

For purposes of comparison, we take \( m_{\rho^0} = m_{\rho^8} = 20 \text{ GeV} \).
and $\alpha_s(20 \text{ GeV}) = 0.25$; $m_{P_{8}'} = 240 \text{ GeV}$ and $\alpha_s(240 \text{ GeV}) = 0.1$; and $N_{f} = 3$, $N_{TC} = 4$. Using the popular choice $C(x) = 3(1-x)^5$, we get

$$\frac{d\sigma(\ell^0)}{dy} \bigg|_{y=0} = \frac{9 \alpha_s^2}{64\pi F_{\tau}^2} \left( 1 - \frac{m_{\ell^0}}{\sqrt{s}} \right)^{10} \approx \begin{cases} 14 \text{ pb at } \sqrt{s} = 800 \text{ GeV} \\ 16 \text{ pb at } \sqrt{s} = 2000 \text{ GeV} \end{cases}$$

(43)

$$\frac{d\sigma(\ell^0)}{dy} \bigg|_{y=0} = \frac{3 \alpha_s^2}{4\pi F_{\tau}^2} \left( 1 - \frac{m_{\ell^0}}{\sqrt{s}} \right)^{10} \approx \begin{cases} 74 \text{ pb at } \sqrt{s} = 800 \text{ GeV} \\ 86 \text{ pb at } \sqrt{s} = 2000 \text{ GeV} \end{cases}$$

(44)

$$\frac{d\sigma(\ell^0)}{dy} \bigg|_{y=0} = \frac{15 \alpha_s^2}{2\pi F_{\tau}^2} \left( 1 - \frac{m_{\ell^0}}{\sqrt{s}} \right)^{10} \approx \begin{cases} 4.3 \text{ pb at } \sqrt{s} = 800 \text{ GeV} \\ 42 \text{ pb at } \sqrt{s} = 2000 \text{ GeV} \end{cases}$$

(45)

In Sec. II we were pessimistic about detecting $H^0$ in its $\ell\ell$, $b\bar{b}$ and $\tau^+\tau^-$ decay modes for cross sections as small as those in Eq. (43). For the same mass, the $P_{8}'$ cross section is $16/3$ times larger and this may be enough to warrant a serious study of detecting it in these modes. We urge that this be done (and apologize in advance if we are ignorant of someone's already having done it). The cross sections for $P_{8}'$ production are also fairly small. Here, however, the expectation that $P_{8}'$ decays predominantly to $\ell\ell$ ($B_{\ell\ell} > 30\%$) is cause for optimism, as we discuss next.

Figure 9 shows the total $P_{8}'$ cross section at $\sqrt{s} = 800$ GeV and 2000 GeV. They differ by a factor of $\sim 30$ in the interesting mass range $m_{P_{8}'} = 200$ - 300 GeV. However, the assumed Isabelle and TEV I luminosities, 10$^{33}$ and 10$^{30}$ cm$^{-2}$ sec$^{-1}$ respectively, turn the tables. For $m_{P_{8}'} = 240$ GeV, one expects $\sim 75,000 P_{8}'$ produced per year at Isabelle, only $\sim 3000$ at TEV I.

These events are to be searched for in the $\ell\ell$ decay mode. The nonleptonic decays of $t$ and $\bar{t}$, comprising $\sim 45\%$ of all $tt$ events, generally give two narrow jets each with little missing energy. Assuming a finely segmented hadronic calorimeter (say, $5^\circ \times 5^\circ$ cells) and large angles between the $t$- and $\bar{t}$-jets, Baltay estimated a 4-jet invariant mass resolution of $\delta m/m \approx 3\%$ at $m \approx 200$ GeV. This estimate may be optimistic now, but one may hope that it will become realistic as experience is gained in single- and multi-jet mass reconstruction.

The main backgrounds to $P_{8}' - \ell\ell$ - jets are light quark and gluon jets with high transverse momentum ($P_T$) and heavy quark ($cc, bb$ and $tt$) jets. The background levels expected for a $10^7$ second run at Isabelle are shown in Fig. 10 along with estimates of the signal for $m_{P_{8}'} = 200$ GeV and 300 GeV. To obtain a useful signal/background level requires: (1) triggering on 4 or more jets (2 or 3 on each side) in which each pair of jets has high invariant mass, of order $m_{t}$; and (2) development of efficient high-resolution vertex detectors to tag the charm decays from $t \to b \to c$. Since only $\approx 5\%$ of high-$P_T$ light quark and gluon jets contain a $cc$ pair, requiring charm on both sides reduces this background by a factor of $2.5 \times 10^{-3}$, where $C_c$ is the charm

![Fig. 9. The total $P_{8}'$ production cross section expected at Isabelle ($\sqrt{s} = 800$ GeV) and TEV I ($\sqrt{s} = 2000$ GeV); from Ref. 56.](image)

![Fig. 10. Background levels to $P_{8}' - \ell\ell$ - jets at Isabelle, with $\int dt = 10^{40}$ cm$^{-2}$. The bumps are the expected $P_{8}'$ signal for $m_{P_{8}'} = 200$ and 300 GeV. Taken from Ref. 56.](image)
detection efficiency of the vertex device. It is expected that requiring that the shape or effective mass of jet pairs be consistent with a t-quark origin will further reduce this background and the c⃗ t, b⃗ b jet backgrounds by ~1/3. A glance at Fig. 10 shows that these rejection factors are sufficient.

Figure 11 shows the results of a hypothetical search for a 300 GeV Pₜ at Isabelle. To obtain this, Baltay assumed b[(Pₜ, t)B(t, 4 jets)] = (0.85)(0.6), cₑ = 0.85 and various geometric cuts, all of which led to an estimated 10% detection efficiency for Pₜ. Folding in this estimate with the Pₜ production cross sections (Fig. 9) and the assumed Isabelle and TEV I luminosities results in Table 8 comparing the event rates at the two machines. It is clear from this table that Pₜ, if it exists, can be found at Isabelle so long as the vertex detector efficiency does not fall below ~10%; success at TEV I will require ~50% efficiency.

<table>
<thead>
<tr>
<th>m_Pₜ (GeV)</th>
<th>σ(pp→Pₜ)(cm²) (events)</th>
<th>τ_Pₜ Events Detected</th>
<th>t_Pₜ Events Detected</th>
</tr>
</thead>
<tbody>
<tr>
<td>100</td>
<td>2 x 10⁶</td>
<td>2 x 10⁵</td>
<td>2000</td>
</tr>
<tr>
<td>200</td>
<td>2 x 10⁵</td>
<td>2 x 10⁴</td>
<td>500</td>
</tr>
<tr>
<td>300</td>
<td>2 x 10⁴</td>
<td>1.5 x 10³</td>
<td>150</td>
</tr>
<tr>
<td>400</td>
<td>2 x 10³</td>
<td>60</td>
<td>60</td>
</tr>
</tbody>
</table>

Fig. 11. Results of a hypothetical search for Pₜ at Isabelle with √s = 800 GeV, ∫dt = 10⁴ cm⁻²; from Ref. 56.

Eichten finds that the pp → PₚX cross section is increased by ~50% and the pp → PₚX cross section by only 10%. His results for the PₜPₜ and PₜPₜPₜPₜ cross sections, per channel, are shown in Fig. 9 for Isabelle and Fig. 14 for TEV I.

The basic subprocesses are GG → PₚPₚ and q̅q → PₚPₚ. The latter is important only for pp collisions when √s = 1-2 TeV. The naive lowest-order graphs are shown in Fig. 12. At high enough energies, the indicated s-channel graphs in Fig. 12 should be vector-meson dominated by the coupling of Gₜ to the technihadron Wₜ. Eichten has calculated the pair production cross sections including this effect as well as that of estimated scaling violations. Unfortunately, Isabelle and TEV I energies are still too low for the Wₜ to have a dramatic effect on the subprocesses. Eichten finds that the pp → PₚX cross section is increased by ~50% and the pp → PₚX cross section by only 10%. Fig. 13 shows the cross section for PₜPₜ and PₜPₜPₜPₜ production at Isabelle (√s = 800 GeV) as a function of technipion mass; from Ref. 59.

Fig. 12. Feynman graphs for the subprocesses (a) GG → PₚPₚ and (b) q̅q → PₚPₚ. Curly lines are gluons, dashed lines are technipions and solid lines are quarks. The fourth graph in (a) and graph (b) are dominated by the Wₜ pole at high energies.

Fig. 14. The cross section for PₜPₜ and PₜPₜPₜPₜ production at TEV I (√s = 2000 GeV) as a function of technipion mass. The dashed curves do not include Wₜ from Ref. 59.
For the technipion mass range of interest, the cross sections expected at Isabelle are:

\[
\sigma(pp \rightarrow P_L^0 P_R^0) = \begin{cases} 
2.9 \times 10^{-38} \text{ cm}^2 & (m_{P_L^0} = 150 \text{ GeV}) \\
0.5 \times 10^{-38} \text{ cm}^2 & (m_{P_L^0} = 175 \text{ GeV}) 
\end{cases}
\]  

(46a)

\[
\sigma(pp \rightarrow P_L^0 P_R^0) = \begin{cases} 
2.9 \times 10^{-39} \text{ cm}^2 & (m_{P_L^0} = 225 \text{ GeV}) \\
0.4 \times 10^{-39} \text{ cm}^2 & (m_{P_L^0} = 250 \text{ GeV}) 
\end{cases}
\]  

(46b)

The expected cross sections at TEV I are:

\[
\sigma(pp \rightarrow P_L^0 P_R^0) = \begin{cases} 
7.1 \times 10^{-36} \text{ cm}^2 & (m_{P_L^0} = 150 \text{ GeV}) \\
3.1 \times 10^{-36} \text{ cm}^2 & (m_{P_L^0} = 175 \text{ GeV}) 
\end{cases}
\]  

(47a)

\[
\sigma(pp \rightarrow P_L^0 P_R^0) = \begin{cases} 
1.0 \times 10^{-35} \text{ cm}^2 & (m_{P_L^0} = 225 \text{ GeV}) \\
0.5 \times 10^{-35} \text{ cm}^2 & (m_{P_L^0} = 250 \text{ GeV}) 
\end{cases}
\]  

(47b)

Thus, for this particular technicolor model, one can expect at most a few hundred leptoquark pairs of each species and practically no color octet pairs produced per year at Isabelle. The decay topologies are similar to those discussed for heavy technipion pair production in ultrahigh energy e+e- annihilation. Clearly, this meager sample of events will be swamped by backgrounds. The number of technipion pairs of each type will be no more than 50-100 at TEV I. There the background levels will be 1000 times smaller than at Isabelle so that, with adequate calorimetry and vertex detection, they may be observable.

In our discussion of e+e- annihilation at ultrahigh energies, we closed with a remark cautioning the reader not to rely too heavily on the detailed results of any specific technicolor model. The same applies here, of course. Even if technicolor is correct, there is no guarantee that technifermions carry ordinary color - a necessary condition for the existence of heavy colored technipions. If they do carry color, they may belong to representations other than the triplet. This possibility would lead to a richer spectrum of technipions, with different masses and production cross sections than used here. Finally, it is important to keep in mind that there is an inherent uncertainty in calculations of technipion production in hadron colliders. The correct parameterizations of the gluon distribution function and of scaling violations and the choice of \( \alpha_s \) are far from settled issues. Different choices can and do lead to appreciably different cross section estimates.

C. ep Colliders

If color triplet leptoquarks exist, they can be found in ep collisions, if the colliders have high enough energy and luminosity and adequate detectors. They may also be possible to photoproduce color octet technipions. Here we assume \( L_{ep} = 10^{32} \text{ cm}^{-2} \text{ sec}^{-1} \) for 10^7 seconds per year.

The naive expectations that leptoquarks couple to fermions like ordinary charged Higgses leads one to the conclusion that their production in ep collisions is dominated by the subprocess shown in Fig. 15. The cross section has been calculated by Rudaz and Vermaseren.\(^{60}\) They take the \( \alpha' \) coupling constant to be \( \sqrt{2} \alpha \eta_\text{ret} \), where \( \eta_\text{ret} \) is an unknown mixing angle factor. Assuming \( m_{P_L^0} = 160 \text{ GeV} \) and \( m_t \approx 20 \text{ GeV} \), they find

\[
\sigma(ep \rightarrow P_L^0 + t + X) \approx \begin{cases} 
3 |\eta_\text{ret}|^2 \text{ pb at } s = 310 \text{ GeV} \\
13 |\eta_\text{ret}|^2 \text{ pb at } s = 800 \text{ GeV} 
\end{cases}
\]  

(48)

The naive expectation also suggests \( |\eta_\text{ret}|^2 \ll 1 \). If this is correct, then prospects for leptoquark production at HERA (\( \sqrt{s} = 310 \text{ GeV} \)) are slim. Clearly, a high-energy machine, \( \sqrt{s} \approx 1 \text{ TeV} \), is required for this job. Such high ep energies are, in fact, achievable by a "cheap" modification of the e+e- colliding linacs discussed earlier.\(^{55}\)

Photoproduction of \( P_L^0 \) and \( P_R^0 \) in ep colliders has been proposed by Grifols and Méndez.\(^{61}\) The basic subprocesses they considered were \( \gamma + g \rightarrow P_L^0 \) and \( \gamma + g \rightarrow P_R^0 \), where the photon is quasireal. The \( \gamma g \) couplings were given in Table 5. For a gluon distribution function \( g(x) = 3(1-x)^5 \), they obtained

\[
\sigma(\gamma p \rightarrow P_L^0) = \frac{1}{10} \times 8 \times \left( \frac{m_{P_L^0}}{s} \right)^2 \frac{\alpha_s g(m_{P_L^0})}{2\pi F^2} \frac{N_{TC}}{6} \text{ pb} \]  

(49)

where our canonical estimate \( \alpha_s m_{P_L^0} = 0.1 \) was used. The subprocesses \( \gamma + g \rightarrow P_L^0 \) and \( \gamma + g \rightarrow P_R^0 \) were overlooked by Grifols and Méndez, but they will not increase the production rates by more than a factor of 2 at very high energies. If we assume an effective \( \gamma p \) luminosity of \( 10^{32} \text{ cm}^{-2} \text{ sec}^{-1} \), we see that one can expect only \( \sim 10 P_L^0 \) produced per year at HERA energies.

We conclude that probes of the scalar sector by ep colliders require center-of-mass energies of the order of 1 TeV.

Fig. 15. Dominant mechanism for technileptoquark production in ep collisions (taken from Ref. 60).

For the techniplon mass range of interest, the cross sections expected at Isabelle are:
Acknowledgements

In one way or another, this report could not have been written without the help, encouragement and information supplied by the following people: Charlie Baltay, Joanne Day and her marvelous colleagues, Bob Diebold, Estia Eichten, Bernie Cittelman, Kurt Gottfried, Ian Hinchliffe, Harris Kagan, Gordy Kane, Jacques Leveille, Steve Olsen, Frank Paige, Dave Pellett, Marty Perl, Michael Peskin, Abe Seiden, Maury Tigges, and Helmut Wiedemann. In other words, I am indebted to the A.P.S. Division of Particles and Fields for its wonderfully stimulating Summer Study. I am also grateful to the Theory Group at CERN for its hospitality and facilities during the initial stages of writing the report. Finally, preparing the report would have been impossible without the amazing typing skill of Phyllis Dolan and the patience and courtesy of Kathy Cox.

This work was supported in part by the Department of Energy under Contract No. AC02-76ER01545.

References and Footnotes


3. From measurement of neutral-current and charged-current deep-inelastic neutrino-nucleon cross sections, it is inferred that \( \delta = \frac{\mu_e / m_e \cos^2 \theta_W}{1.00 \pm 0.03} \) (see, e.g., J.E. Kim, et al., Rev. Mod. Phys. 53, 211 (1981)). The only natural way to have \( \delta = 1 \) in the standard model is spontaneous gauge symmetry breaking induced by elementary or composite Higgs doublets alone.


10. J. Ellis, M.K. Gaillard and D.V. Nanopoulos, Nucl. Phys. B106, 292 (1976); G. Barbiellini, et al., "The Production and Detection of Higgs Particles at LEP", DESY Report 78-06 (1979). Marty Perl, Michael Peskin, Abe Seiden, Maury Tigges, and Helmut Wiedemann. In other words, I am indebted to the A.P.S. Division of Particles and Fields for its wonderfully stimulating Summer Study. I am also grateful to the Theory Group at CERN for its hospitality and facilities during the initial stages of writing the report. Finally, preparing the report would have been impossible without the amazing typing skill of Phyllis Dolan and the patience and courtesy of Kathy Cox.

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33. To avoid the possibility of absolutely stable color-triplet leptons, it is generally assumed that \( Y = 1/6 \). While we are not sure there is a problem if \( Y = 1/6 \) (e.g., if \( Y = 1/12 \), a bound state of \( F_{Y} \) and \( \bar{U} \) is a color singlet fermion with charge 2/3), we shall use that value in practical calculations.
34. J.D. Bjorken, unpublished calculations.
36. As indicated in Table 2, color singlet technihadrons decay with full strength into \( W_{L} \) and \( Z_{L} \), the longitudinal components of the electroweak gauge bosons. The reason for this is that these components really are just the three unphysical technipions absorbed in the Higgs mechanism (see Ref. 6).
37. This, in fact, is what happens in the original version of the Farhi-Susskind model, Ref. 9.
38. K.D. Lane, unpublished calculations.
39. To have \( \alpha_{T}^{+} = \alpha_{T}^{-} \), ETC interactions would have to be parity conserving, in which case charge 2/3 and -1/3 quarks would be degenerate in pairs (Ref. 8). This is not unique to ETC models; the couplings of elementary charged Higgs mesons to fermions must be parity-violating for the same reason.
43. Two remarks are in order here: (1) It is possible that the enhancement of \( \phi_{R} \) modes brought about by the color factor of 3 in Eq. (20) is ameliorated by Giesachs in the \( A_{\mu}^{R}, B_{\mu}^{R} \) etc. For example, in one model considered in G. Barbieri, et al. (Ref. 11), couplings to quarks are reduced by a factor of 3 relative to those to leptons. In another model, however, they are increased by a relative factor of 3. We have taken a moderate middle road in our models, but it is easy to modify the numbers in Tables 3 and 4 to reflect such factors. (2) We have assumed flavor-conserving couplings of neutral technipions to quark and lepton pairs. We remind the thorough reader that it is just possible that this is not the case and that neutral technipions may prefer such decay modes as \( \bar{c}_{L} \) or \( c_{L} \), \( b_{R} \) or \( b_{L} \), etc.
44. In terms of the original unmixed fields \( \phi_{R} \) (see Ref. 28), the interaction giving the third of Eq. (21)

\[
- \frac{\sin \alpha_{T}}{2} \bar{e}^{\mu} \mathcal{Q}_{T}^{(1)} \mathcal{Q}_{T}^{(1)\mu} = \frac{3}{2} \Re \mathcal{Q}_{T}^{(1)} \mathcal{Q}_{T}^{(1)\mu} \text{Im} \mathcal{Q}_{T}^{(1)} \text{(summed over } \nu_{1})
\]

only the real parts \( \phi_{1}^{T} + \phi_{1}^{T} / 2 - v_{1} \) contribute to the interaction term.
48. I thank John Ellis for supplying me with the composite drawing (Fig. 7) of the results of the collaborations in Ref. 47.
49. See S. Olsen's discussion of toponium in the report of the e+e- Collider Group in these proceedings.
51. H. Ragan, in Report of the 100 GeV Facility Subgroup of the e+e- Collider Group, these proceedings.
52. The reader hardly needs to be reminded that one can do a much more thorough search with improved capability to detect \( P^{+} \) in all-hadron modes.
53. Note that not both neutral Higgs mesons in this decay need be more massive than 2m_{t}. If both are, \( Z_{0} \) may be kinematically forbidden.
54. See the report of the 1 TeV Facility Subgroup of the e+e- Collider Group for discussion of the parameters, physics capabilities and physics backgrounds of colliding linacs.
55. I thank M. Peskin for his speedy calculation of this cross section at the Summer Study. The importance of technihadron dominance in high-energy e+e- and hadron-hadron collisions was, as far as I know, first realized in Snowmass. (See also Ref. 59.)
56. This discussion of \( P_{0}^{0} \) production and detection is abstracted from C. Baltay's transparencies at the Summer Study. For further details, see the report of the Hadron-Hadron Collider Group in these proceedings.
57. In his discussion, Baltay assumed the nonleptonic branching ratio of the t-quark to be 3/5. Naively, it should be 2/3, but this does not affect the results of Table 8 significantly.
58. J.A. Grifols and A. Méndez, Phys. Rev. D26, 324 (1982). These authors did not include effects of \( \phi_{0}^{0} \) nor of scale violation in their calculations.
59. E. Eichten, private communication.