Tests of a composite model of quarks and leptons

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Recent studies of radiative corrections to electroweak processes restrict the particle content of dynamical-symmetry-breaking schemes. The minimal anomaly-free model with one doublet of Dirac fermions $F \equiv (U, D)$ having charges $Q = \pm \frac{1}{2}$ is consistent with such restrictions. These fermions may be regarded as the constituents of a strongly interacting Higgs sector, so that the longitudinal W and Z and higher-spin excitations (including vector mesons, called ρ_T , with isospin $I=1$, and ω_T , with $I=0$) are all composites of F 's. It is then natural to ask if these fermions can be the sole carriers of weak $SU(2)$, so that quarks and leptons would also have to be composed of them. Some tests of this hypothesis in hadronic and e^+e^- collisions are suggested.

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I. INTRODUCTION

The mechanism whereby electroweak symmetry is broken remains the central mystery of an otherwise outstandingly successful theory. The theory starts out with an $SU(2) \times U(1)$ symmetry, which is broken at a scale of several hundred GeV to the $U(1)$ of electromagnetism. Two manifestations of this symmetry breaking appear in the spectrum. First, the W and Z bosons are massive, while the photon is massless. Second, the quarks and the leptons e , μ , and τ , are massive. Mass terms in the Lagrangian for either the gauge bosons or the quarks and leptons break $SU(2) \times U(1)$.

The dynamical generation of W and Z masses [1] can be accomplished by having their longitudinal components be the Nambu-Goldstone bosons of a spontaneously broken chiral symmetry. These spinless objects are the analogue in some "technicolor" theory of the pions in quantum chromodynamics with two quark flavors. The analogues of the quarks are "technifermions" F , whose properties are as yet unspecified. The analogue of QCD is a new (presumably asymptotically free) interaction whose characteristic scale is about 2500 times that of QCD: 250 GeV instead of 100 MeV.

Quark and lepton masses are harder to generate dynamically. A further interaction is usually postulated [2] in which the ordinary quarks and leptons are pointlike and communicate with the technifermions F via the exchange of "extended technicolor" gauge bosons. The number of technifermions is generally increased substantially beyond the minimal set needed to give W 's and Z 's their masses. As the result of spontaneous breaking of

large chiral-symmetry groups, explicit models of such interactions generally imply the existence of numerous low-mass spinless bosons, for which there is no evidence at present. Moreover, it becomes difficult to avoid flavor-changing neutral currents in such models [3].

A new, powerful constraint on models such as extended technicolor has come on the scene recently. It has been pointed out [4–7] that electroweak radiative corrections are sensitive to effects on W and Z wave-function renormalizations which can become significant when there are large numbers of new particles carrying weak isospin and hypercharge. Analyses of electroweak data [5,8–10] now turn out to severely restrict the number of such new particles. In particular, a model [11] with a single family of colored techniquarks and colorless technileptons already seems to be ruled out.

In the present article we explore an approach to the generation of gauge-boson masses by which we hope to avoid some of the problems of extended technicolor schemes. We assume that the technifermions discussed above are the only carriers of electroweak $SU(2)$, and hence that quarks and leptons are composed of these technifermions and some other $SU(2)$ -singlet matter. The approach is not new, having strong connections with various schemes proposed over the years [12–18]. Although we have by no means solved all of the problems inherent in such a picture, we do put forward a model that is sufficiently specific that several direct experimental tests for compositeness are suggested. It is these tests which form the main subject of this paper. The tests are variants of those proposed by Eichten, Lane, and Peskin, [19] and are applied to both electron-positron and hadron-hadron collisions. Generalization to electron-hadron collisions is straightforward.

A major point of principle [20] must be addressed before a composite model of quarks and leptons can be considered successful. What keeps quarks and leptons light in comparison with the scale of the dynamics? This scale

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is at least a couple of TeV, and could be higher. It has been suggested [20] (see Ref. [21] for a review) that nearly unbroken chiral symmetry is what protects quark and lepton masses from becoming large. If the symmetry is spontaneously broken (leading to the Nambu-Goldstone bosons which become longitudinal W 's and Z 's in the presence of the electroweak gauge interactions), then it is presumably unavailable for protecting quark and lepton masses. We do not have a complete solution to this problem, but we do propose one form for electroweak symmetry breaking which is capable of generating gauge-boson masses while preserving enough symmetry to leave quarks and leptons massless. We leave for the future the question of how this symmetry can be slightly broken so that the quarks and leptons get masses that are small compared to the TeV scale of compositeness.

The tests we shall propose are not restricted to composite models of quarks and leptons. Indeed, since our main focus will be on learning the Yukawa couplings of quarks and leptons to new heavy vector bosons in the 1–2-TeV mass range, our discussion will also have relevance to technicolor theories. The Yukawa couplings of technivector mesons to quarks and leptons have been discussed extensively (see, e.g., Ref. [22]), but there does not seem to be consensus as to their magnitudes.

We find that the possibility of technivector-meson exchange in box diagrams contributing to K - $\bar{K}$ mixing leads to strong constraints on technivector mesons which couple to both right-handed and left-handed quarks. Consequently, the specific predictions we shall present are all based on the assumption of purely left-handed couplings [22]. They can be easily modified for vector couplings if desired.

We discuss properties of technifermions and their bound states in Sec. II. The model itself and one variant of it are presented in Sec. III. Implications for electron-positron and hadron-hadron collisions are explored in Secs. IV and V. Limits on vector Yukawa couplings of quarks and leptons to bound states of technifermions, based on present limits of standard electroweak couplings and right-handed W bosons, are mentioned in Sec. VI. Section VII concludes. An Appendix is devoted to symmetry breaking by means of composite fields transforming in various ways.

II. TECHNI-FERMIONS AND THEIR BOUND STATES

A. Anomalies and technifermion properties

In QCD, the lowest-lying particle whose decay provides “microscopic” information is the neutral pion. The decay $\pi^0 \rightarrow \gamma\gamma$ occurs as a result of the Adler-Bell-Jackiw anomaly [23], which is generated from the coupling of the pion to quark loops. The amplitude is proportional to the combination $S \equiv N_c(Q_u^2 - Q_d^2)$, where N_c denotes the number of colors of quarks and Q_u, Q_d are their charges. The experimental decay rate allows one to conclude that, for $Q_u = \frac{2}{3}$, $Q_d = -\frac{1}{3}$, one must have $N_c = 3$.

Aside from the properties of π^0 , one must study physics

at scales of 1 GeV or higher to learn about properties of hadronic subconstituents. Nearly all hadronic physics below 1 GeV can be described in terms of an effective Lagrangian for pion-pion scattering, suitably unitarized and with crossing symmetry restored to the best of one's ability [24]. The only other “microscopic” information below 1 GeV in addition to that provided by $\pi^0 \rightarrow \gamma\gamma$ involves couplings of the ω meson, whose mass is $783 \text{ MeV}/c^2$ [25].

One might have hoped that the two-photon decay of the corresponding neutral pseudo Nambu-Goldstone boson composed of techniquarks could provide similar useful information about techniquarks and their multiplicity. However, there is an important difference between technicolor and QCD. The neutral technipion Π^0 plays the role of the longitudinal Z boson. Hence, in order to avoid an anomalous theory, the decay $\Pi^0 \rightarrow \gamma\gamma$ must be forbidden. Thus, we do not learn anything about the number of technifermions. We do, however, learn that the charges of the technifermions must be such as to cancel the anomaly.

The minimal anomaly-free model [26,27] for the fermions responsible for dynamical electroweak symmetry breaking thus consists of two Dirac fermions $F \equiv (U, D)$ with charges $Q_U = \frac{1}{2}$, $Q_D = -\frac{1}{2}$. It is convenient to label these objects by left-handed and right-handed isospins I_{3L} and I_{3R} , as well as by the vector sum $\mathbf{I} \equiv \mathbf{I}_{3L} + \mathbf{I}_{3R}$, whose third component will be their charge:

$$Q(F) = I_3 \equiv I_{3L} + I_{3R}. \quad (2.1)$$

The gauge bosons of $SU(2)_L$ are the ones which have been seen so far. It is not necessary to gauge $SU(2)_R$. The left-handed versions of the fermions F have $I_L = \frac{1}{2}$, with $I_{3L} = \pm \frac{1}{2}$ and $I_R = 0$, while the right-handed versions have $I_L = 0$ and right-handed isospin $I_R = \frac{1}{2}$. The technifermions then have a definite value of $I = \frac{1}{2}$, just like the u and d quarks out of which pions and nucleons are made.

B. Spontaneous symmetry breaking

Following the ideas that enter the technicolor model [1], we assume that the fermions F are subject to strong interactions that have a typical momentum scale, analogous to Λ_{QCD} , on the order of a TeV. These interactions are assumed to leave invariant the $SU(2)_L \times SU(2)_R$ chiral symmetry of the technifermions, as in QCD. One assumes that this chiral symmetry is spontaneously broken to the vector $SU(2)$ subgroup. This spontaneous symmetry breaking leads to an isotriplet of massless Nambu-Goldstone bosons Π^+ , Π^- , Π^0 . We return to the question of the nature of the symmetry breaking, and thus of the technipions, in Sec. III, when we discuss the quarks and leptons.

The interactions of the technipions Π with one another are uniquely specified at low energies. They are characterized by amplitudes proportional to scalar products of momenta divided by the electroweak scale factor $v \equiv 2^{-1/4} G_F^{-1/2} = 246 \text{ GeV}$. In like manner, pion-pion scattering amplitudes involve similar scalar products di-

vided by the pion decay constant $f_\pi = 93$ MeV. Thus, the physics of low-energy technipion interactions is expected to resemble that of low-energy pion-pion scattering, but with energies scaled up by $v/f_\pi \simeq 2650$.

C. Bound states involving technifermions alone

Because of their strong interactions, the technifermions can form bound states other than the technipions. The lightest vector particles are expected to be an isotriplet and an isosinglet which we shall call ρ_T and ω_T . The field that destroys these particles can, for instance, have the form $V^\mu(x) = Z^{-1/2} \bar{F}(x) \gamma^\mu F(x)$. Various estimates [2,27,29–32] suggest that their masses are of the order of $2\pi v$, or in the range of 1.5–2 TeV/ c^2 . We shall assume for definiteness in what follows that $M(\rho_T) = M(\omega_T) = 2\pi v = 1547$ GeV/ c^2 .

The degeneracy of the ρ and ω in QCD is a consequence of the small coupling of the ω to three-gluon states which would otherwise break this degeneracy (as in the case of pseudoscalar mesons). We shall assume a similar degeneracy holds in the case of ρ_T and ω_T . The degeneracy will be important when we come to discuss consequences of ρ_T and ω_T exchange in high-energy reactions.

The existence of ρ_T is expected to be a rather general feature of models with dynamical electroweak symmetry breaking, much as the existence of the ρ meson is a consequence of the low-energy pion-pion interactions, analyticity, and crossing symmetry [24]. If the technicolor interaction is based on an $SU(N_T)$ gauge group analogous to $SU(3)$ of color, it has often been estimated [3,27,33] that one should have $M(\rho_T) \simeq (2 \text{ TeV}/c^2)(3/N_T)^{1/2}$.

It is possible to anticipate the production of ρ_T in hadron collisions on the basis of two mechanisms: ρ_T - W mixing and gauge-boson fusion [31,32]. One expects cross sections which put the discovery of a ρ_T within the reach of multi-TeV colliders. There are some differences of opinion about whether the lower center-of-mass energy of the proposed Large Hadron Collider (LHC) at CERN (16 TeV) will be adequate, and generally more confidence that the Superconducting Super Collider (SSC) energy (40 TeV) will suffice.

It has been pointed out [32] that the ω_T , if it could be discovered, would provide more information about basic properties of the subconstituents F , much as the couplings of the ordinary ω to $\rho\pi$ or 3π are governed by physics sensitive to quark substructure [26]. As one example [32], the mixing of the ω_T with the gauge boson B of the electroweak theory coupling to weak hypercharge is proportional to $2Q_U - 1$, and vanishes for the minimal choice $Q_U = \frac{1}{2}$. As a result of this and several other circumstances (including a favored decay mode consisting of W^+W^-Z), the potential for discovery of this particle, even at the SSC, is marginal. One must turn to other possibilities for learning about the properties of the technifermions F . We shall see that studies of possible quark and lepton compositeness could, in principle, provide such information.

III. A COMPOSITE MODEL OF QUARKS AND LEPTONS

A. Building blocks

We shall assume in what follows that the technifermions F of the minimal model, with $Q_U = \frac{1}{2}$ and $Q_D = -\frac{1}{2}$, are the only subconstituents of matter carrying weak $SU(2)$ quantum numbers. The quarks and leptons must then be made of these objects and something else which is $SU(2)$ neutral.

The charges of quarks and leptons may be expressed [34] as

$$Q = I_{3L} + I_{3R} + \frac{B - L}{2}, \quad (3.1)$$

where B and L are the baryon and lepton numbers. We can identify the first two terms in (3.1) as due to the charges (2.1) of the fermions F , provided the chiralities of the quarks and leptons are the same as those of the fermions F they contain. The remaining contribution to the charge in (3.1) should then be provided by an object which does not affect the quark or lepton spin, and which distinguishes a quark from a lepton. A simple solution is to take this subconstituent to be a scalar particle:

$$\begin{aligned} \bar{S}_q^a & \text{ for a quark } [Q(\bar{S}_q^a) = \frac{1}{6}], \\ \bar{S}_l & \text{ for a lepton } [Q(\bar{S}_l) = -\frac{1}{2}]. \end{aligned} \quad (3.2)$$

Here a is a color index ($a = 1, 2, 3$). Thus, the first family of quarks and leptons may be represented as $F\bar{S}$ bound states:

$$u = U\bar{S}_q, \quad d = D\bar{S}_q, \quad \nu_e = U\bar{S}_l, \quad e^- = D\bar{S}_l. \quad (3.3)$$

The corresponding antiquarks and antileptons are to be $\bar{F}\bar{S}$ bound states.

We assume the interaction is such that $F\bar{S}$ and $\bar{F}S$ are allowed bound states, but not FS or $\bar{F}\bar{S}$. This solution is realizable, for example, if both F and S belong to the fundamental N_T -dimensional representation of an $SU(N_T)$ technicolor group, with $N_T \geq 3$, and if the technicolor group is confining in the same manner as QCD. This restriction, by the way, also forbids the existence of FF or $\bar{F}\bar{F}$ bound states, whose presence in the spectrum would probably be undesirable.

In principle, there might exist $S\bar{S}$ bound states. If they were light, one would be able to see them in decays of Z 's, since both S_l and S_q have nonzero weak hypercharge. However, one expects that these states will be heavy, since they are not protected from getting a TeV-scale mass either by a chiral symmetry or by being Nambu-Goldstone bosons.

The scalar $\bar{S}$ should be coupled to the fermion F in such a way that the total spin of the system is $\frac{1}{2}$. The simplest solution is to take the quarks and leptons to correspond to S -wave bound states of $\bar{S}$ and F particles. Thus, the field to destroy an up quark would be $u(x) \propto U(x)\bar{S}_q(x)$. We must next provide an explanation for the existence of more than one family of quarks and leptons.

B. The problem of families

Two proposals have been made for the multiplicity of quark and lepton families in the model just proposed (or variants of it). One can either assume that the scalars $\bar{S}$ (or, for that matter, the fermions F) carry a family index, or one can postulate the existence of a level structure.

Let us first consider the possibility that the quark and lepton families are the result of a quantum-mechanical level structure. Thus, for example, the electron, muon, and tau would be bound states of the same D and $\bar{S}$ constituents, but with different wave functions. The states would be analogous to the $1S$, $2S$, and $3S$ levels of the hydrogen atom, except that the constituents would be highly relativistic instead of nonrelativistic. This point of view was advocated, for example, in Ref. [15].

As mentioned in the Introduction, we take the view that a chiral symmetry prevents the quarks and leptons from having a mass on the order of $\Lambda \sim \text{TeV}$, where $1/\Lambda$ is the size of the bound state. One lacks a detailed model of the dynamics, but one may suppose that more than one bound state of the two constituents may be produced. Then, in the symmetry limit, the allowed bound states would all be massless. Then a small symmetry breaking would give different masses to the quarks and leptons of different families.

Since bound-state wave functions for different levels (i.e., families) composed of the same constituents are orthogonal to one another, it is natural for neutral-current gauge-boson couplings to be flavor diagonal at $q^2=0$. For $q^2 \neq 0$, one might expect flavor-changing neutral-current effects of order $q^2 \langle r^2 \rangle$, where r is the radius of the bound state.

Charged-current gauge-boson couplings involve replacing a fermion F by another differing from it by a unit of charge. Now, the wave functions for quarks and leptons containing U technifermions will differ slightly from ones containing D technifermions, if only as a result of different electromagnetic interactions between the technifermion and the scalar S . Thus, a nontrivial Cabibbo-Kobayashi-Maskawa matrix will be generated as a result of nonzero overlap integrals among wave functions. Such a matrix should have large diagonal elements and smaller off-diagonal ones, as is observed.

Some care must be taken to ensure that the above requirements are self-consistent. The mere orthogonality of wave functions for different states may not always ensure the absence of flavor-changing neutral currents. As one example, one can consider the two quark-model states Σ^0 and Λ , composed of u , d , and s . These are orthogonal but the magnetic dipole transition $\Sigma^0 \rightarrow \Lambda + \gamma$ is allowed. The ud system in such a transition changes from one of spin 1 to one of spin 0. In the present model, however, the situation is different. The subconstituents participating in the transition either have spin zero (and, hence, do not generate transitions between S states) or spin $\frac{1}{2}$ (so that the dominant transitions between S states are magnetic dipole, and are again highly suppressed for different S states).

An alternative to the identification of quark and lepton families with nS states was discussed in Ref. [35]. If one

wishes to identify such families with internal degrees of freedom (such as spin configurations of a system of several spin- $\frac{1}{2}$ fermions, as in the Σ^0 and Λ cases discussed above), one can also arrange for the absence of flavor-changing neutral currents. Two solutions proposed in Ref. [35] were (1) the choice of suitable charges of the subconstituents, and (2) the construction of configuration-mixed states to suppress the undesirable transitions. We do not explore these alternatives further here, but they might be viable for suitable bound-state dynamics.

The picture of families resulting from a level structure is quite attractive, especially since one does not need too many fundamental particles at the scale Λ . However, this picture has a serious problem if one wants to have the compositeness scale Λ be on the order of 1 TeV, as we do in this paper. The problem is that one must guard against flavor-changing transitions at a very low level, for example, in the decay $\mu \rightarrow e\gamma$. By writing the amplitude for the $\mu \rightarrow e\gamma$ transition in terms of the operator $\exp(iq \cdot x)$ evaluated between infinite-momentum wave functions, as in Ref. [36], one can conclude that the $\mu \rightarrow e\gamma$ amplitudes should be proportional to the square of a charge radius $\langle r^2 \rangle \sim 1/\Lambda^2$ times the square of the momentum transfer $|q|^2 \sim m_\mu^2$. The present upper limit [37] on the $\mu \rightarrow e\gamma$ branching ratio is $\leq 5 \times 10^{-11}$. If the rate for $\mu \rightarrow e\gamma$ is proportional to the square of the amplitude obtained above, $\Gamma(\mu \rightarrow e\gamma) \sim (\alpha/\pi) m_\mu (m_\mu/\Lambda)^4$, we must have $\Lambda \gtrsim 200 \text{ TeV}$.

We thus adopt in this paper the alternative view that the family structure is due to a family quantum number carried by the scalars S . We are then free to suppose that there is an exact $U(1)^3$ symmetry that results in the conservation of the family quantum number. This model is similar to the model of Ref. [13], in which the particle which we call the scalar $\bar{S}$ is actually a two-fermion composite. One of the fermions (called c) carries a color index running from 1 to 4, with the fourth color denoting a lepton. The other (called h) carries the family, or “horizontal” index. Thus, in the present language, $\bar{S}$ would carry both (extended) color and a family label. (The fermions we call F , carrying a weak index, are called w in the model of Ref. [13].)

C. Yukawa couplings of bound states involving fermion F

1. Pseudoscalar mesons

The interactions of the Nambu-Goldstone bosons Π with quarks and leptons must be such as to be the only surviving couplings of massive W 's or Z 's when the $SU(2)$ and $U(1)$ gauge constants g and g' are turned off [31]. For example, the charged-current weak interaction which takes the low-energy form

$$\mathcal{H}_{\text{int}} = \frac{G_F}{\sqrt{2}} J^\dagger J_\mu \quad (3.4)$$

becomes instead, in the gaugeless limit,

$$\mathcal{H}_{\text{int}} = \frac{G_F}{\sqrt{2}} (q^\mu J_\mu^\dagger) \frac{1}{q^2} (q^\nu J_\nu), \quad (3.5)$$

where J_μ is the charged weak current:

$$J_\mu = (\bar{u} \bar{c} \bar{t})_L \gamma_\mu V \begin{pmatrix} d \\ s \\ b \end{pmatrix}, \quad (3.6)$$

and V is the Cabibbo-Kobayashi-Maskawa (CKM) matrix. Thus, Eq. (3.5) describes the interaction of a spinless particle coupled to the divergence to the weak current, normalized by $G_F/\sqrt{2}=1/(2v^2)$. Now, the divergence of the current introduces terms proportional to quark masses, so all of the Yukawa couplings of the Π bosons are proportional to ratios of quark masses to the parameter v . Moreover, these Yukawa couplings must exactly follow the pattern of the CKM matrix elements, in the sense that they are all generated from the same current J_μ which involves this matrix.

2. Vector mesons

The interactions of ρ_T and ω_T with ordinary quarks and leptons are subject to constraints from the absence of flavor-changing neutral currents. We shall assume that these constraints are satisfied by having diagonal interactions of ω_T and ρ_T^0 , and interactions of $\rho_T^\pm$ which follow the pattern of the CKM matrix.

We shall further assume that ω_T and ρ_T^0 couplings are equal in magnitude and either equal or opposite in sign, depending of the sign of the charge of the fermion F in the quark or lepton. We thus assume a universal Yukawa coupling g_Y of the meson $(D\bar{D})_V \equiv (\omega_T - \rho_T^0)/\sqrt{2}$ to all particles containing D technifermions. The same Yukawa coupling also characterizes the coupling of $(U\bar{U})_V \equiv (\omega_T + \rho_T^0)/\sqrt{2}$ to all particles containing U technifermions. These universality properties appear to be characteristic of the couplings of the ordinary ω and ρ to hadrons.

One might expect the TeV-scale strong interactions to be parity conserving, by analogy with quantum chromodynamics. In this case, one would expect the technirho states to couple to the quarks and leptons with vector couplings. However, as we discuss in Sec. VI, limits from $K - \bar{K}$ mixing disfavor vector couplings which are large enough to be visible in the processes considered in this paper. These stringent limits do not, however, apply to left-handed couplings. Hence, although we shall present matrix elements for both vector and left-handed couplings, numerical results will be given only for left-handed couplings.

Are left-handed couplings reasonable? Let us point out that we know virtually nothing about the form of the TeV-scale strong interactions (if they exist at all). They may be parity conserving, like QCD, or they may be maximally parity violating, like the $SU(2)_L$ weak interactions. Let us entertain the possibility that they violate parity. Then one might have bound technirho states made of left-handed technifermions, $\bar{F}_L(x)\gamma^\mu F_L(x)$, but it could be that the right-handed technifermions do not bind, or that their bound states have substantially greater masses. In this case, the left-handed quarks and leptons $F_L(x)\bar{S}(x)$ would have enhanced couplings to the tech-

nirhos compared to the right-handed quarks and leptons. As an approximation, one could neglect the right-handed coupling at low energies.

D. Models for the symmetry breaking

We now return to a discussion of the nature of the spontaneous breaking of the $SU(2)_L \times SU(2)_R$ chiral symmetry. We adopt a notation in which the doublet (U, D) is denoted F^i with an $SU(2)$ index $i=1,2$.

In a parity-conserving model, as in QCD, it seems natural that there be strong attractive forces to bind the technifermions in the scalar channel. Then, we suppose, the field

$$\phi_{im}(x) = \bar{F}_R^m(x) F_L^i(x) \quad (3.7)$$

acquires a vacuum expectation value:

$$\langle \phi_{im}(x) \rangle \propto \delta_{im}. \quad (3.8)$$

This breaks the $SU(2)_L \times SU(2)_R$ symmetry to $SU(2)_V$ and leads to three Nambu-Goldstone bosons. In this picture, these technipions can be viewed as bound states of a technifermion and an antitechnifermion, just as the QCD pions are bound states of a quark and an antiquark. The technipions then combine with the $SU(2)_L \times U(1)$ gauge bosons to give masses to the W and Z .

This rather standard picture is quite attractive. However, it leaves us at a loss to explain why the $F\bar{S}$ bound states, the quarks and leptons, are nearly massless compared to the TeV scale of the technifermion interactions. The chiral symmetry of the theory would have provided a natural explanation, but this symmetry has been broken.

One can recover an explanation for having nearly massless quarks and leptons if one supposes that the symmetry-breaking pattern is different. The field $\bar{F}_R^m(x) F_L^i(x)$ used above transforms according to the $(\frac{1}{2}, \frac{1}{2})$ representation of $SU(2)_L \times SU(2)_R$. Suppose, instead, a field ϕ_{im} that transforms according to the $(1,1)$ representation acquires a vacuum expectation value $\langle \phi_{im}(x) \rangle \propto \delta_{im}$. Then one will still get three Nambu-Goldstone bosons. As we show in the Appendix, with this symmetry-breaking pattern the mass matrix for the W and Z bosons remains the same as in the standard model. However, now the vacuum is left invariant under the $SU(2)_L$ transformation $\mathcal{U}_L = -1$ [i.e., $F_L^i(x) \rightarrow -F_L^i(x)$]. This discrete Z_2 subgroup of the original chiral symmetry is enough to forbid mass terms $\bar{q}(x)q(x) = \bar{S}(x)\bar{F}(x)F(x)\bar{S}(x)$ in the low-energy effective Lagrangian.

This symmetry-breaking picture is more complicated, and thus less attractive, than the QCD-like picture. In particular, one loses the simple picture of the technipions as bound states of a technifermion and an antitechnifermion. However, it does provide a symmetry reason for having nearly massless quarks and leptons.

We mentioned at the beginning of Sec. II A that the decay $\pi^0 \rightarrow \gamma\gamma$ provides information about the constituents of the nucleon. Pion decay is really a low-energy phenomenon, with a 100-MeV energy scale. At this energy scale, one can make use of the low-energy effective

theory of pions and nucleons, the σ model. The pion decay amplitude is determined by the anomaly [23] in the three-point function $\langle J_5^\mu(x) J^\alpha(y) J^\beta(z) \rangle$, where $J^\alpha(x)$ is the electromagnetic current and $J_5^\mu(x)$ is the nucleon contribution to the I_3 component of the weak-SU(2) axial-vector current. Within this model,

$$J_5^\mu(x) = \bar{p} \gamma_5 \gamma^\mu p - \bar{n} \gamma_5 \gamma^\mu n, \quad (3.9)$$

and calculation of the anomaly gives the experimentally observed strength $S = (Q_p^2 - Q_n^2) = 1$. One can, however, also calculate this anomaly in the short-distance theory, QCD, in which the nucleons are composed of quarks. The strength of the anomaly matches:

$$S = N_c (Q_u^2 - Q_d^2) = 3[(\frac{2}{3})^2 - (\frac{1}{3})^2] = 1. \quad (3.10)$$

Following 't Hooft [20], one may ask whether this and other similar “anomaly-matching conditions” can be applied once again at a yet smaller distance scale, to constrain the theory at the TeV level. The answer is, unfortunately, “no,” as far as we can see. In the case of the σ model as a low-energy approximation to QCD, $J_5^\mu(x)$ is the generator of a chiral symmetry of both the long-distance and the short-distance Lagrangians. Thus, there are two ways to calculate the same anomaly, and consistency demands that we get the same answer either way. In the present case there is a long-distance theory QCD and a presumed new short-distance theory at the 1-TeV scale. The current $J_5^\mu(x)$ is the generator of a chiral symmetry for QCD, but there is no corresponding symmetry for the short-distance theory. There is an $SU(2)_L \times SU(2)_R$ chiral symmetry in the short-distance theory, but the axial part of this symmetry corresponds to the gauge group of the weak interactions and is anomaly-free in both QCD and the short-distance theory, as discussed in Sec. II A. In the long-distance theory, this symmetry group rotates both the quarks and the leptons at the same time. The short-distance theory does not have a chiral symmetry corresponding to $J_5^\mu(x)$, which rotates just the quarks. Instead, the long-distance theory possesses a larger symmetry group than the short-distance theory. With the spontaneous symmetry-breaking scheme proposed above for the short-distance theory, quark mass terms are forbidden in the long-distance effective Lagrangian. Operators such as $[\bar{q}(x)q(x)]^2$ that violate chiral symmetry are not forbidden. However, such operators are dimension six instead of dimension four, so they do not contribute to the limiting long-distance effective action.

In this way, we arrive at a long-distance effective theory, QCD, that has chiral symmetry. This theory also has chiral symmetry breaking in its Green's functions because of the axial anomaly (in addition to the spontaneous chiral symmetry breaking). One can view the anomaly as a long-distance phenomenon, as in Ref. [38]. More conventionally, the anomaly is calculated using the short-distance divergences of QCD, since it is at short distances that the one-loop triangle graph suffices for the calculation. That is, within the picture proposed here, one uses the short-distance properties of the long-distance effective theory QCD to calculate the anomaly.

It appears to be an open question whether this is mathematically consistent—that is, whether a chiral symmetry with an anomaly can arise in an effective long-distance theory from a fundamental theory that does not have the symmetry.

A real dynamical model for either form of symmetry breaking mentioned above still has not been found, despite many efforts [21] stimulated by the work of Ref. [20]. Nonetheless, with either symmetry-breaking picture, and with the specific assignment of quantum numbers proposed in Sec. II, the implications for experiments in the energy range much below 1 TeV are the same. It is to these implications that we now turn.

IV. IMPLICATIONS FOR ELECTRON-POSITRON COLLISIONS

A. Coherence of vector-meson contributions

We wish to consider electron-positron annihilation reactions which are affected by the presence of Yukawa couplings to the vector mesons ρ_T^0 and ω_T . These couplings lead to amplitudes which can be depicted in graphical [39] form as shown in Fig. 1. Only the amplitudes for production of charged leptons (e , μ , and τ) and $I_3 = -\frac{1}{2}$ quarks (d , s , and b) are affected. In these reactions, with our assumption about the relation between Yukawa couplings of ρ_T^0 and ω_T , these two vector mesons contribute constructively and equally to the amplitude. The ρ_T^0 and ω_T contributions to production of neutrinos and u -type quarks in e^+e^- annihilations will cancel one another. One cannot write diagrams such as those of Fig. 1 for $e^+e^- \rightarrow (u\bar{u} \text{ or } \nu\bar{\nu})$ since neutrinos and u -type quarks are assumed to contain U technifermions, while electrons and positrons are taken to contain D technifermions.

We shall consider two types of Yukawa couplings: vector and left handed. It will turn out in Sec. VI that limits from $K-\bar{K}$ mixing disfavor vector couplings which are large enough to be visible in the processes considered here. Hence, although we shall present matrix elements for both cases, numerical results will be given in this section and the next only for left-handed couplings.

The parts of the Lagrangian describing the interactions between fermion field ψ and photons, Z 's, and technivector mesons are

$$\mathcal{L}_{\text{int,em}} = -e \bar{\psi} \gamma^\mu A_\mu Q \psi, \quad (4.1)$$

where e is the positron charge ($e > 0$) and Q is a diagonal

FIG. 1. Diagrams denoting contributions of technivector mesons to (a) $e^+e^- \rightarrow \mu^+\mu^-$, (b) $e^+e^- \rightarrow d\bar{d}$.

matrix measuring the charge of the fermion:

$$\mathcal{L}_{\text{int},Z} = -\sqrt{g^2 + g'^2} \bar{\psi} \gamma^\mu Z_\mu (I_{3L} - Q \sin^2 \theta) \psi, \quad (4.2)$$

where $g \equiv e/\sin\theta$ and $g' \equiv e/\sin\theta$ are SU(2) and U(1) coupling constants of the electroweak theory, and θ is the electroweak mixing angle; and

$$\mathcal{L}_{\text{int},V} = -g_Y \bar{\psi}_1 \gamma^\mu V_\mu \psi_2 \quad (\text{vector}) \quad (4.3a)$$

or

$$\mathcal{L}_{\text{int},V} = -(g_Y/2\sqrt{2}) \bar{\psi}_1 \gamma^\mu (1 - \gamma_5) V_\mu \psi_2 \quad (\text{lefthanded}). \quad (4.3b)$$

Note the normalizations in (4.3). They are adopted by analogy with electromagnetism for the vector coupling and W exchange for the left-handed coupling. By the discussion in Secs. II and III (see also Fig. 1), the normalizations are appropriate for both charged and neutral technivector-meson exchange.

We then find that the exchange of $(D\bar{D})_V$ in Fig. 1 leads to an amplitude for $e^+e^- \rightarrow f\bar{f}$ of the form

$$\mathcal{M}|_{\text{techni}} = 4\pi\alpha_Y J^\mu(f) \frac{1}{M_V^2 - q^2} J_\mu(e) \quad (\text{vector}) \quad (4.4a)$$

or

$$\mathcal{M}|_{\text{techni}} = 2\pi\alpha_Y J_L^\mu(f) \frac{1}{M_V^2 - q^2} J_{\mu L}(e) \quad (\text{left handed}), \quad (4.4b)$$

where $J_L^\mu \equiv \bar{u} \gamma^\mu [(1 - \gamma_5)/2] u$ is the left-handed current. The amplitude (4.4a) or (4.4b) is to be added to the amplitude for photon exchange:

$$\mathcal{M}|_\gamma = 4\pi\alpha Q_f Q_e J^\mu(f) \frac{1}{-q^2} J_\mu(e), \quad (4.5)$$

and a corresponding amplitude for Z exchange. Here J_μ denotes the vector current, while $\alpha_Y \equiv \xi g_Y^2/(4\pi)$ is the analogue of the fine-structure constant α . The parameter ξ allows us to relax the universality assumption (which would correspond to $\xi = +1$); we shall examine the consequences of $\xi = \pm 1$.

B. Cross-section formulas

The differential cross section for e^+e^- annihilation to a fermion-antifermion pair $f\bar{f}$ may be written as

$$\frac{d\sigma}{d(\cos\theta^*)}(e^-e^+ \rightarrow f\bar{f}) = \frac{\pi\alpha^2}{8s} N_c [(|f_{LL}|^2 + |f_{RR}|^2)(1 + \cos\theta^*)^2 + (|f_{LR}|^2 + |f_{RL}|^2)(1 - \cos\theta^*)^2], \quad (4.6)$$

where α is the electromagnetic fine-structure constant, s is the squared center-of-mass (c.m.) energy, θ^* is the c.m. scattering angle, $N_c = 1$ for leptons and 3 for quarks, and the f 's are amplitudes for scattering in definite helicity states. The helicities are those of the incident e^- and final fermion f (a d -type quark or negatively charged lepton). The helicities of the remaining particles are specified by the form of the interaction. (We are neglecting initial and final fermion masses.) For processes in which the exchange of the meson $(D\bar{D})_V$ is assumed to contribute, one can write

$$\begin{aligned} f_{(L,R)L} &\simeq -Q_f + \frac{(I_{3(L,R)} - Q\bar{x})(-1/2 + \bar{x})}{\bar{x}(1 - \bar{x})} \frac{s}{s - M_Z^2 + iM_Z\Gamma_Z} \\ &\quad + \frac{\alpha_Y}{\alpha} \frac{s}{s - M_V^2} \quad (\text{vector}), \end{aligned} \quad (4.7a)$$

$$\begin{aligned} f_{(L,R)R} &\simeq -Q_f + \frac{I_{3(L,R)} - Q\bar{x}}{1 - \bar{x}} \frac{s}{s - M_Z^2 + iM_Z\Gamma_Z} \\ &\quad + \frac{\alpha_Y}{\alpha} \frac{s}{s - M_V^2} \quad (\text{vector}), \end{aligned} \quad (4.7b)$$

where $\bar{x} \equiv \sin^2\theta_W$ as measured at the Z pole. The corresponding expressions for left-handed couplings are

$$\begin{aligned} f_{LL} &\simeq -Q_f + \frac{(I_{3L} - Q\bar{x})(-1/2 + \bar{x})}{\bar{x}(1 - \bar{x})} \frac{s}{s - M_Z^2 + iM_Z\Gamma_Z} \\ &\quad + \frac{\alpha_Y}{2\alpha} \frac{s}{s - M_V^2} \quad (\text{left handed}), \end{aligned} \quad (4.8a)$$

$$\begin{aligned} f_{RL} &\simeq -Q_f + \frac{(I_{3R} - Q\bar{x})(-1/2 + \bar{x})}{\bar{x}(1 - \bar{x})} \frac{s}{s - M_Z^2 + iM_Z\Gamma_Z} \\ &\quad (\text{left handed}), \end{aligned} \quad (4.8b)$$

$$\begin{aligned} f_{(L,R)R} &\simeq -Q_f + \frac{I_{3(L,R)} - Q\bar{x}}{1 - \bar{x}} \frac{s}{s - M_Z^2 + iM_Z\Gamma_Z} \\ &\quad (\text{left handed}). \end{aligned} \quad (4.8c)$$

The approximation consists in neglecting the widths of the technivector mesons. This will be satisfactory for the present case, in which we consider e^+e^- c.m. energies $\sqrt{s} \ll M_V$. In fact, one expects the ρ_T to be considerably

broad than the ω_T , so it is not sufficient merely to put in terms corresponding to a single width when evaluating Eqs. (4.4) near the ρ_T and ω_T poles. A more detailed treatment, analogous to the discussion of ρ - ω mixing, would be needed.

The forward-backward asymmetries for these processes may also be computed. If we define

$$\sigma_F \equiv \int_0^1 d(\cos\theta^*) \frac{d\sigma}{d(\cos\theta^*)},$$

$$\sigma_B \equiv \int_{-1}^0 d(\cos\theta^*) \frac{d\sigma}{d(\cos\theta^*)},$$
(4.9)

then the forward-backward asymmetry $A_{FB} \equiv (\sigma_F - \sigma_B)/(\sigma_F + \sigma_B)$ is given in terms of the helicity amplitudes by

$$A_{FB} = \frac{3}{4} \frac{|f_{LL}|^2 + |f_{RR}|^2 - |f_{LR}|^2 - |f_{RL}|^2}{|f_{LL}|^2 + |f_{RR}|^2 + |f_{LR}|^2 + |f_{RL}|^2}. \quad (4.10)$$

C. Results

We take $\alpha = \frac{1}{128}$ and $\bar{\alpha} = 0.23$ to correspond to coupling strengths at the Z pole. Variations for other values of s are ignored. The values of the vector-meson masses are

FIG. 2. Cross sections for e^+e^- production affected by the presence of the meson $(D\bar{D})_\nu$ in the direct channel. Solid lines: $\alpha_Y=0$; dashed lines: $\alpha_Y=\pm 0.1$ (labeled by values); dotted lines: $\alpha_Y=\pm 1$ (labeled by values). (a) b -quark pairs, (b) charged-lepton pairs, (c) hadrons (five flavors). Hadronic cross sections are presented without the QCD correction factor $(1 + [\alpha_s/\pi])$.

taken to be $M_V = 2\pi v = 1547 \text{ GeV}/c^2$, as noted earlier. The results are then shown for cross sections in Fig. 2 and for forward-backward asymmetries in Fig. 3. We present results for b quarks because their flavor is readily identified through their semileptonic decay.

Cross-section measurements may be able to spot the presence of Yukawa couplings $\alpha_Y = \pm 1$ at the highest energies reached at the CERN e^+e^- collider LEP. Limits which can be interpreted as excluding such large couplings for vector interactions have already been presented [40]. Values of $\alpha_Y = \pm 0.1$ (corresponding to the dotted curves in Fig. 2) are probably not accessible to cross-section measurements either below or above the Z at present or foreseen levels of precision.

One draws a similar conclusion from the predictions for forward-backward asymmetries shown in Fig. 3. A value of $\alpha_Y = \pm 1$ could have been missed up to now, but would show up around $\sqrt{s} = 150 \text{ GeV}$. There is no hope of seeing effects of $|\alpha_Y| \leq 0.1$ in e^+e^- forward-backward asymmetries.

$$\frac{d\sigma}{d(\cos\theta^*)}(d\bar{u} \rightarrow e^- \bar{\nu}_e) = \frac{\pi\alpha^2}{8N_c s} [(|f_{LL}|^2 + |f_{RR}|^2)(1 + \cos\theta^*)^2 + (|f_{LR}|^2 + |f_{RL}|^2)(1 - \cos\theta^*)^2], \quad (5.1)$$

where the helicity amplitudes are now

$$f_{LL} \simeq \frac{1}{2\bar{x}} \frac{s}{s - M_W^2 + iM_W \Gamma_W} + \frac{\alpha_Y}{\alpha} \frac{s}{s - M_V^2} \quad (\text{vector}), \quad (5.2a)$$

$$f_{LR} = f_{RL} = f_{RR} \simeq \frac{\alpha_Y}{\alpha} \frac{s}{s - M_V^2} \quad (\text{vector}), \quad (5.2b)$$

or

$$f_{LL} \simeq \frac{1}{2\bar{x}} \frac{s}{s - M_W^2 + iM_W \Gamma_W} + \frac{\alpha_Y}{2\alpha} \frac{s}{s - M_V^2} \quad (\text{lefthanded}), \quad (5.3a)$$

$$f_{LR} = f_{RL} = f_{RR} = 0 \quad (\text{lefthanded}). \quad (5.3b)$$

Here the vector meson V is the charged ρ_T . As in Eqs. (4.4), we have neglected its width. The Yukawa coupling strength is normalized in the same way for $\rho_T \equiv D\bar{U}$ exchange as for the exchange of the $D\bar{D}$ meson; a single dual graph with unit Clebsch-Gordan coefficient contributes in either case.

B. Production of jets

Processes such as those shown in Fig. 1 can also contribute to quark-quark scattering, and hence to the production of jets at high transverse momenta. In most vari-

V. IMPLICATIONS FOR HADRON COLLISIONS

A. Production of lepton pairs

(1) Pairs of charged leptons. The Drell-Yan process for production of pairs of charged leptons (e.g., e^+e^-) by pairs of d -type quarks is given by an expression similar to (4.6). One must merely replace N_c by $1/N_c$ when considering the annihilation of quarks in color-singlet hadrons. For present purposes we shall assume that the Yukawa couplings of $(D\bar{D})_V$ to $d\bar{d}$, $s\bar{s}$, and $b\bar{b}$ are all equal. There are additional contributions from annihilation of u -type quarks which only involve the photon and Z poles.

(2) Single-charged-lepton production. There is also a Drell-Yan process for production of lepton pairs involving a single charged lepton. In the standard model, this process occurs via the W pole. In the present model, the charged ρ_T also contributes, as shown in Fig. 4. The subprocess cross section is of the form (for the annihilation of quarks in color-singlet hadrons)

ants of the present model the cross sections should not be appreciably larger than for the Drell-Yan lepton-pair production processes, except that both $u\bar{u}$ and $d\bar{d}$ collisions can contribute. However, backgrounds are expected to be considerably greater, since the process of interest is now competing with gluon exchange rather than exchange of electroweak gauge bosons. Consequently, we have not calculated cross sections for such processes, though the methods for doing so are straightforward.

C. Results

As in Sec. IV, we take $\alpha = \frac{1}{128}$ and $\bar{x} = 0.23$. For e^+e^- productions, we consider contributions of $d\bar{d}$, $s\bar{s}$, and $b\bar{b}$ annihilations, while for $e^- \bar{\nu}_e$ production, we consider the contributions of both $d\bar{u}$ and $s\bar{c}$ collisions. We assume quark distributions in the proton to be described by the structure functions of Harriman, Martin, Roberts, and Stirling (set B) [41]. We consider proton-antiproton collisions at present Fermilab Tevatron energies ($\sqrt{s} = 1.8 \text{ TeV}$). It is likely that the Tevatron can be run at 2 TeV in several years, adding to the ability to probe higher parton-parton subenergies.

(1) Pairs of charged leptons. Predicted cross sections for charged-lepton-pair production are shown in Fig. 5. We show effects of both positive and negative signs for α_Y . As mentioned earlier, the positive sign is the one expected if the couplings are universal. The present CDF data, with several events per pb of cross section, can probably exclude values as large as $|\alpha_Y| = 1$ immediately, but even experiments with hundreds of events per pb will

FIG. 3. Forward-backward asymmetries in e^+e^- production affected by the presence of the meson $(D\bar{D})_\nu$ in the direct channel. (a) b quark pairs, (b) charged-lepton pairs. Labels as in Fig. 2.

FIG. 4. Diagrams for the subprocess $d\bar{u} \rightarrow e^-\bar{\nu}_e$. (a) Standard model, (b) ρ_T contribution.

FIG. 5. Predicted cross sections for $\bar{p}p \rightarrow e^-e^+$ with the transverse momentum of each lepton greater than $p_T^{\min}$. Curves are labeled as in Fig. 2. Of each pair, the lower curve corresponds to the case of positive α_Y .

have a hard time seeing the effects of $|\alpha_Y|=0.1$.

(2) Single-charged-lepton production. Predicted cross sections are shown in Fig. 6. The conclusions are qualitatively similar to those for charged-lepton-pair production. Present CDF data see no evidence for departure from the standard model out to transverse momenta where present statistics become a limitation [42]. One can thus exclude values of $|\alpha_Y|=1$ but not yet $|\alpha_Y|=0.1$. Slightly greater sensitivity to smaller values of α_Y appears in this process than in charged-lepton-pair production.

FIG. 6. Predicted cross sections for $\bar{p}p \rightarrow e^-\bar{\nu}_e$, with transverse momentum of each lepton greater than $p_T^{\min}$. Curves are labeled as in Fig. 2. Of each pair, the lower curve corresponds to the case of positive α_Y .

VI. LOW-ENERGY LIMITS ON VECTOR YUKAWA COUPLINGS

A. Interaction properties

The exchange of a charged ρ_T should behave in many respects like the exchange of a linear combination of left-handed and right-handed W bosons. We can write the effective Lagrangian for the vector interaction of ρ_T with a pair of weak eigenstates (for example, ν_e and e^-) as

$$\begin{aligned}\mathcal{L}_{\text{eff}}^T &= -g_Y [\bar{\nu}_e \gamma^\mu e] \rho_{T\mu}^{(+)} + \text{H.c.} \\ &= -(g_Y/2) [\bar{\nu}_e \gamma^\mu (1 - \gamma_5) e] \rho_{T\mu}^{(+)} \\ &\quad - (g_Y/2) [\bar{\nu}_e \gamma^\mu (1 + \gamma_5) e] \rho_{T\mu}^{(+)} + \text{H.c.},\end{aligned}\quad (6.1)$$

to be compared with the standard-model interaction

$$\mathcal{L}_{\text{eff}}^W = -\frac{g}{2\sqrt{2}} [\bar{\nu}_e \gamma^\mu (1 - \gamma_5) e] \rho_{T\mu}^{(+)} + \text{H.c.}, \quad (6.2)$$

where g is the electroweak $\text{SU}(2)$ coupling constant, $g^2/(4\pi) = \alpha/\bar{x} \simeq 1/29.4$.

B. Tree-level processes

The exchange of a ρ_T with mass M_V should behave for processes which involve only left-handed neutrinos as an additional contribution to the Fermi interaction. At low values of momentum transfer, one then expects a modification

$$\frac{G_F}{\sqrt{2}} = \frac{g^2}{8M_W^2} \rightarrow \frac{g^2}{8M_W^2} + \frac{g_Y^2}{4M_V^2} = \frac{G_F}{\sqrt{2}} \left[1 + \frac{g_Y^2 v^2}{2M_V^2} \right], \quad (6.3)$$

where we have used the relation $M_W = gv/2$. With the nominal value $M_V = 2\pi v$, which we have been assuming, Eq. (6.3) reduces to a modification

$$\frac{G_F}{\sqrt{2}} \rightarrow \frac{G_F}{\sqrt{2}} \left[1 + \frac{\alpha_Y}{2\pi} \right]. \quad (6.4)$$

As long as no right-handed neutrinos exist which are sufficiently light, nothing about this modification can be learned from the study of leptonic process alone. In particular, experiments [43] on the electron spectrum in muon decays which have been used to search for a right-handed W would not be affected. It becomes clear here that the present model would be quite uncomfortable with light right-handed neutrinos.

One may compare purely leptonic processes with semileptonic ones [to which the right-handed part of (6.2) would also contribute] to see the effects of ρ_T exchange. The fact that the CKM matrix appears unitary to a high degree of accuracy [44], with $|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 0.998 \pm 0.002$, provides a constraint. The precise measurements of CKM elements, from which the unitarity constraints are derived, are based on properties of *vector* currents, whose strengths are regulated by conservation laws [45,46].

Let us consider the effective Hamiltonian for semilep-

tonic interactions of u quarks:

$$\begin{aligned}\mathcal{H}_{\text{eff}} &= \frac{g^2}{8M_W^2} [\bar{u} \gamma^\mu (1 - \gamma_5) d'] [\bar{l} \gamma_\mu (1 - \gamma_5) \nu_l] \\ &\quad + \frac{g_Y^2}{M_V^2} [\bar{u} \gamma^\mu d'] [\bar{l} \gamma_\mu \nu] + \text{H.c.}\end{aligned}\quad (6.5)$$

As long as neutrinos are left handed, we can rewrite the second term of (6.5) as

$$\begin{aligned}\Delta\mathcal{H}_{\text{eff}} &= \frac{g_Y^2}{4M_V^2} [\bar{u} \gamma^\mu (1 - \gamma_5) d' + \bar{u} \gamma^\mu (1 + \gamma_5) d'] \\ &\quad \times [\bar{l} \gamma_\mu (1 - \gamma_5) \nu] + \text{H.c.},\end{aligned}\quad (6.6)$$

where we have expressed the left-handed neutrino field as $\nu = (1 - \gamma_5)\nu/2$. The effective hadronic vector current V^μ then appears corrected by a term

$$V^\mu \rightarrow V^\mu \left[1 + \frac{8M_W^2}{g^2} \frac{g_Y^2}{2M_V^2} \right] = V^\mu \left[1 + \frac{\alpha_Y}{\pi} \right] \quad (6.7)$$

if we assume as above that $M_V = 2\pi v$. The net effect, when the overall renormalization of G_F in Eq. (6.4) is taken into account, is a correction to the amplitude of order $\alpha_Y/(2\pi)$. Since this correction should be $\lesssim 0.2\%$ at the 2σ level [44], one has $|\alpha_Y| \lesssim 0.013$.

Nonleptonic decay processes, in principle, could provide limits on the Yukawa coupling g_Y if we were able to precisely evaluate the effects of strong interactions. Our very rudimentary understanding of the K_S lifetime indicates that this goal is still in the future.

C. Effects on processes in the full $\text{SU}(2) \times \text{U}(1)$ theory

In the unified $\text{SU}(2) \times \text{U}(1)$ theory of weak and electromagnetic interactions, the Fermi constant G_F is related to the fine-structure constant, the W mass, and the value of $\bar{x}$ measured at the Z pole. This relation is

$$G_F = \frac{\pi\alpha}{\sqrt{2}\bar{x}M_W^2}. \quad (6.8)$$

The left-hand side of (6.8) is $1.16637(2) \times 10^{-5} \text{ GeV}^{-2}$. Taking $\alpha^{-1} = 128.02 \pm 0.12$, $M_W = 80.14 \pm 0.31 \text{ GeV}/c^2$ (see the discussion in Ref. [47]), and $\bar{x} = 0.232 \pm 0.003$ from recent measurements of forward-backward asymmetries in lepton pair production at LEP [48], we find the right-hand side is equal to $1.175(18) \times 10^{-5} \text{ GeV}^{-2}$. Thus, there is agreement to a couple of percent. As a result, we find there is still room for a value of $|\alpha_Y|$ of up to about 0.1 in Eq. (6.4). This bound is less stringent than that in Sec. IV B. It will improve as the experimental errors on M_W and $\bar{x}$ decrease.

It is conceivable that new physics (such as W wavefunction renormalization) could affect the relation (6.8) by a percent or two [5,8–10]. Our motivation has been explicitly to avoid large effects of this sort, but it is certainly possible that they could conspire to camouflage a larger value of $|\alpha_Y|$. As we shall see, there is one more constraint on α_Y which appears to be quite powerful.

D. K - $\bar{K}$ mixing

The absence of large K - $\bar{K}$ mixing effects was one of the motivations for the introduction of the charmed quark, since the contribution of box diagrams such as those shown in Fig. 7(a) is then brought under control. However, diagrams in which one of the left-handed W bosons is replaced by a right-handed one W_R have a vastly enhanced matrix element, and contribute too much to the mixing term unless [49,50] $M(W_R) \geq 1.4$ – 2.5 TeV. This bound is tailor made for our purposes. We are assuming the existence of a new boson whose mass is just at this limit, and whose couplings have the space-time behavior of $W_L + W_R$, as indicated in Eq. (6.1). The bounds of Refs. [49,50] apply to a W_R coupling equal to g [the weak SU(2) coupling strength]. Thus, comparing (6.1) and (6.2), we conclude that one must have

$$g_Y \leq g/\sqrt{2}, \quad \alpha_Y \leq \left[\frac{1}{2} \right] \left[\frac{1}{29.4} \right] \simeq 0.017. \quad (6.9)$$

A more detailed investigation of this bound has been performed [51]; one ends up with a restriction comparable to that in Sec. VI B. It appears to be a serious constraint on the size of possible Yukawa couplings. We have assumed here a pattern of CKM couplings for ρ_T which mirrors those of the ordinary W , in accord with our discussion in Sec. III. To alter this pattern so as to weaken the bounds, as was done for the W_R in Ref. [50], seems to us very dangerous, as it opens up all sorts of possibilities for flavor-changing neutral currents.

One can escape the strictness of the bound (6.9) if the ρ_T couples only to left-handed currents. Since the tests proposed in Secs. IV and V are sensitive to both left-handed and right-handed couplings of the technivector mesons, they are more general than the bound just mentioned.

E. The role of $S\bar{S}$ mesons

We alluded in Sec. III A to the apparent absence of low-lying $S\bar{S}$ mesons. If the lightest of such mesons is, in fact, heavy (in the 1–2-TeV mass range), one may be able to anticipate some of its properties by reference to the dual graphs of Fig. 8. Thus, information on the Yukawa couplings of technivector mesons to ordinary quarks and leptons, which govern processes such as that shown in Fig. 8(a), may be relevant to Yukawa couplings of $S\bar{S}$ mesons [Fig. 8(b)].

FIG. 7. Examples of box diagrams for K - $\bar{K}$ mixing. (a) Exchange of two (left-handed) W bosons; (b) exchange of a left-handed W boson and a charged ρ_T .

FIG. 8. Diagrams illustrating duality between exchange of (a) technivector and (b) $S\bar{S}$ mesons.

VII. CONCLUSIONS

We have discussed a composite model for quarks and leptons which envisions them as S -wave bound states of a fermion F and a scalar $\bar{S}$. The fermions have charge $Q = \pm \frac{1}{2}$ and carry weak SU(2)_L and SU(2)_R quantum numbers. The scalars distinguish quarks from leptons; for a quark $Q(\bar{S}_q) = \frac{1}{6}$, while for a lepton $Q(\bar{S}_l) = -\frac{1}{2}$. The scalars also distinguish one family from another.

A condensate of a composite formed from the fermion fields F and $\bar{F}$ is assumed to be the source of electroweak dynamical symmetry breaking. Aside from pseudoscalar Nambu-Goldstone bosons, which play the role of longitudinal components of W and Z bosons, there are no other bound states of the F 's expected below 1 TeV. The next lowest in mass are expected to be vector mesons, whose mass should lie in the 1–2 TeV range. We have explored some tests for exchanges of these mesons which rely upon their (as yet specified) Yukawa couplings g_Y to familiar quarks and leptons.

Both in e^+e^- reactions at LEP energies and in proton-antiproton collisions at 1.8 TeV, we find the eventual limit of sensitivity to be around $\alpha_Y \equiv g_Y^2/(4\pi) = 0.1$. This is close enough to the largest Yukawa coupling of the Higgs boson to matter (i.e., the coupling of the Higgs boson to the top quark, assuming $m_t = 100$ – 200 GeV/ c^2) that searches for such effects are worthwhile. We stress that there is little theoretical guidance on the magnitude of g_Y and on the question of its universality. With some assumptions about universality, we have found that the absence of large K - $\bar{K}$ mixing effects limits α_Y in a theory with vector couplings to be at most 1 or 2 %. Such a value is not accessible at present energies, and requires higher-energy e^+e^- or hadron collisions to be seen.

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APPENDIX

We want to consider symmetry breaking by means of a composite field that transforms under the (j_L, j_R) repre-

sensation of $SU(2)_L \times SU(2)_R$. Thus, we consider a composite field $\phi_{im}(x)$ that transforms according to $\phi \rightarrow \hat{\phi}$ where

$$\hat{\phi}_{im} = U_L(\mathbf{g}_L)_{ij} U_R(\mathbf{g}_R)^*_{mn} \phi_{jn} . \quad (\text{A1})$$

Here U_L is the j_L representation of the group element $\mathbf{g}_L$ of $SU(2)_L$ and U_R is the j_R representation of the group element $\mathbf{g}_R$ of $SU(2)_R$.

We consider two choices of (j_L, j_R) . The first is the standard choice for compositeness, $(j_L, j_R) = (\frac{1}{2}, \frac{1}{2})$. As we shall see, the standard model can be considered to be a special case of this choice [52]. The second possibility considered is $(j_L, j_R) = (1, 1)$. In this case, the symmetry breaking leaves a discrete chiral symmetry that keeps the quarks and leptons massless. (We do not discuss schemes for breaking the chiral symmetry and thus giving masses to the quarks and leptons in this paper.) In the case $(j_L, j_R) = (\frac{1}{2}, \frac{1}{2})$, $\phi_{im}(x)$ is a complex field, while in the case $(j_L, j_R) = (1, 1)$, we take $\phi_{im}(x)$ to be a real field.

1. The vector-boson mass matrix

Although ϕ is a composite field for us, we investigate the group theory by treating it as an elementary field. The term in the Lagrangian that generates the vector-boson masses is

$$\mathcal{L} = \frac{1}{2} [\mathcal{D}_\mu \phi]_{im}^* [\mathcal{D}^\mu \phi]_{im} . \quad (\text{A2})$$

(If ϕ is really complex, one usually omits the $\frac{1}{2}$ in this equation, so that the present normalization is a bit non-standard.) Here

$$\begin{aligned} [\mathcal{D}_\mu \phi]_{im} = & [\partial_\mu \delta_{ij} \delta_{mn} - i g t_{ij}^a W_\mu^a(x) \delta_{mn} \\ & + i g' \delta_{ij} t_{nm}^3 B_\mu(x)] \phi_{jn}(x) . \end{aligned} \quad (\text{A3})$$

Here t_{ij}^a is the generator for the j_L ($= j_R$) representation of $SU(2)$. The hypercharge matrix $\frac{1}{2} Y_{mn} = -(t_{mn}^3)^* = -t_{nm}^3$ generates the $U(1)$ part of $SU(2)_L \times U(1)$. It is related to the electric charge by

$$Q_{im,jn} = t_{ij}^3 \delta_{mn} + \delta_{ij} \frac{1}{2} Y_{mn} . \quad (\text{A4})$$

Now we suppose that ϕ acquires a vacuum expectation value

$$\langle \phi_{im}(x) \rangle = v_{im} . \quad (\text{A5})$$

The symmetry breaking should leave the electric charge conserved:

$$Q_{im,jn} v_{jn} = 0 . \quad (\text{A6})$$

That is,

$$t_{ij}^3 v_{jm} = v_{in} t_{nm}^3 . \quad (\text{A7})$$

In addition, in the composite models we consider, it should leave the global $SU(2)_V$ symmetry unbroken. We shall see how these conditions can be satisfied.

The term in the Lagrangian that generates the vector boson masses is obtained by setting ϕ equal to its vacuum

expectation value in $\mathcal{L}$ above. This gives

$$\begin{aligned} \mathcal{L}_M = & \frac{1}{2} v_{jn}^* [g t_{ji}^a W_\mu^a(x) \delta_{nm} - g' \delta_{ji} t_{mn}^3 B_\mu(x)] \\ & \times [g t_{ik}^b W_\mu^b(x) \delta_{mr} - g' \delta_{ik} t_{rm}^3 B_\mu(x)] v_{kr} . \end{aligned} \quad (\text{A8})$$

The vector-boson mass term thus has the form

$$\mathcal{L}_M = \frac{1}{2} (W_{a\mu} W_\mu^a M_{WW}^{ab} + W_{a\mu} B^\mu M_{WB}^a + B_\mu B^\mu M_{BB}) , \quad (\text{A9})$$

where (denoting $[A, B]_+ = AB + BA$)

$$M_{WW}^{ab} = g^2 \text{Tr}(v v^\dagger \frac{1}{2} [t^a, t^b]_+) , \quad (\text{A10})$$

$$M_{WB}^a = -g g' \text{Tr}(v^\dagger t^a v t^3) , \quad (\text{A11})$$

$$M_{BB} = g'^2 \text{Tr}(v^\dagger v t^3 t^3) . \quad (\text{A12})$$

Note that one can apply an $SU(2)_L$ rotation (i.e., redefine W_a) so that M_{WB}^a lies in the three-direction. Then it is the W_3 meson that mixes with the B meson. If, after thus defining the three-direction, it happens that t^3 commutes with v , as in Eq. (A7), then we have

$$\mathcal{L}_M = \frac{1}{2} \tilde{W}_{a\mu} \tilde{W}_\mu^a M_{WW}^{ab} , \quad (\text{A13})$$

where

$$\tilde{W}_1^\mu = W_1^\mu ,$$

$$\tilde{W}_2^\mu = W_3^\mu ,$$

$$\tilde{W}_3^\mu \equiv \sqrt{1 + (g'/g)^2} Z^\mu = W_3^\mu - \frac{g'}{g} B^\mu . \quad (\text{A14})$$

Then the Z meson acquires a mass, while the orthogonal combination of W_3 and B mesons, namely, the photon, is massless. The standard-model mass relation between the W and Z mesons is obtained when $M_{WW}^{ab} = \text{const} \times \delta_{ab}$. It is interesting that, in the $(j_L, j_R) = (\frac{1}{2}, \frac{1}{2})$ case, one has $\frac{1}{2} [t^a, t^b]_+ = \delta_{ab}$, so that $M_{WW}^{ab} = \text{const} \times \delta_{ab}$ is guaranteed whenever the symmetry breaking is such as to leave the photon massless.

2. The effective potential

We suppose that the field $\phi_{im}(x)$ acquires a vacuum expectation value because of the new strong interactions at the TeV scale. The vacuum expectation value v_{im} will occur at a minimum of the effective potential $V_{\text{eff}}(\phi_{im})$.

One may expect that the electroweak gauge interactions will also influence the vacuum expectation value, but we ignore this effect to start with. Thus, we should consider the effective potential created by the new strong interactions (and also the QCD gauge interactions) with the electroweak interactions turned off.

Since (we suppose) the new interactions are invariant under $SU(2)_L \times SU(2)_R$, the effective potential will also have this symmetry. Thus, V_{eff} can depend only on the $SU(2)_L \times SU(2)_R$ invariants formed from ϕ_{im} . For the case of the $(1, 1)$ representation, these invariants may be taken to be $\lambda_1, \lambda_2, \lambda_3$, where

$$\begin{aligned}
\lambda_1 + \lambda_2 + \lambda_3 &= \frac{1}{3} \text{Tr}[\phi\phi^\dagger], \\
\lambda_1^2 + \lambda_2^2 + \lambda_3^2 &= \frac{1}{3} \text{Tr}[(\phi\phi^\dagger)^2], \\
\lambda_1\lambda_2\lambda_3 &= \det(\phi\phi^\dagger).
\end{aligned} \tag{A15}$$

Thus $\lambda_1, \lambda_2, \lambda_3$ are the three eigenvalues of the Hermitian matrix $\phi\phi^\dagger$ (and also of $\phi^\dagger\phi$). For the case of the $(\frac{1}{2}, \frac{1}{2})$ representation, these invariants may be taken to be λ_1, λ_2 , and ϕ_0 , where

$$\lambda_1 + \lambda_2 = \frac{1}{2} \text{Tr}[\phi\phi^\dagger], \quad \lambda_1\lambda_2 = \det(\phi\phi^\dagger), \tag{A16}$$

and where ϕ_0 is the phase of $\det(\phi)$. [We assume that the Lagrangian is invariant under changes in the overall phase of $\phi_{im}(x)$, so that $\mathcal{V}$ does not depend on ϕ_0 .] Again, λ_1, λ_2 are the eigenvalues of $\phi\phi^\dagger$. Thus, $V_{\text{eff}}(\phi_{im}) = \mathcal{V}(\lambda_1, \lambda_2, \lambda_3)$ or $V_{\text{eff}}(\phi_{im}) = \mathcal{V}(\lambda_1, \lambda_2)$, depending on the choice of representation.

Now the minimum of $\mathcal{V}$ could occur at $\lambda_1 = \lambda_2 (= \lambda_3) = 0$, that is, at $\phi_{im} = 0$. This possibility is quite natural because it has maximum symmetry. However, it does not break any symmetries and thus leaves all of the gauge bosons massless. We will assume that nature chooses the next most symmetric possibility, that the minimum of $\mathcal{V}$ occurs at $\lambda_1 = \lambda_2 (= \lambda_3) = |a|^2 \neq 0$.

3. Vacuum expectation value and gauge-boson masses

Under this assumption about $\mathcal{V}$, the vacuum expectation value v_{im} will lie on the surface $\lambda_1 = \lambda_2 (= \lambda_3) = |a|^2$ in the space of ϕ_{im} . This surface can be characterized as follows. If v lies on this surface, then one must have $vv^\dagger = |a|^2 1$. Thus, $a^{-1}v$ is a unitary matrix in the $(\frac{1}{2}, \frac{1}{2})$ case or an orthogonal matrix in the $(1, 1)$ case:

$$v = aU, \quad U \in \text{O}(3) \text{ or } \text{SU}(2). \tag{A17}$$

The whole surface is obtained by applying an $\text{SU}(2)_L$ transformation to the matrix $v = a1$. If the $\text{SU}(2)_L$ transformation is g , then the resulting v is $v = aU(g)$, where $U(g)$ is an orthogonal matrix or $\text{SU}(2)$ matrix as appropriate. Further transformations from the whole $\text{SU}(2)_L \times \text{SU}(2)_R$ group do not move us out of the surface since they amount to left multiplying or right multiplying by further orthogonal or $\text{SU}(2)$ matrices.

These surfaces are of special interest in the case of the $(\frac{1}{2}, \frac{1}{2})$ representation. Suppose for the moment that $\phi_{im}(x)$ is an elementary field instead of a composite. Suppose further that $\phi_{im}(x)$ is constrained to have the form $\phi(x) = a(x)U(x)$ as above, and that the phase ϕ_0 is constrained to be zero. Then one can write $\phi(x)$ in the form [52]

$$\phi(x) = \begin{bmatrix} \phi_1(x) & -\phi_2(x)^* \\ \phi_2(x) & \phi_1(x)^* \end{bmatrix}. \tag{A18}$$

We then arrive at a theory of an $\text{SU}(2)_L$ -doublet field (ϕ_1, ϕ_2) coupled to $\text{SU}(2)_L \times \text{U}(1)$ gauge fields—that is, the Higgs and gauge sectors of the standard model. Thus, the $(\frac{1}{2}, \frac{1}{2})$ model with a minimum of V_{eff} at an equal- λ_i point is the same as the standard model except that the quantum $\phi(x)$ can have fluctuations off the equal- λ_i , $\phi_0 = 0$ surface. The fluctuations away from $\lambda_1 = \lambda_2 = |a|^2$ yield four Higgs bosons rather than the single Higgs boson of the standard model. (In the context of composite models considered in this paper, one expects the Higgs bosons to have masses in the TeV range and to merge with the spectrum of other heavy particles.) The fluctuations away from $\phi_0 = 0$ yield an apparent Nambu-Goldstone boson analogous to the η' in QCD. Presumably, this particle does not survive as a Nambu-Goldstone boson because of the axial $\text{U}(1)$ anomaly, as in QCD [28].

Consider now a vacuum expectation value on the surface $\lambda_1 = \lambda_2 = \lambda_3 = |a|^2$ in the $(1, 1)$ representation. As we have seen, this surface has the form $v = aU$ with U running over $\text{SU}(2)$. The surface has three dimensions, so there would be three Nambu-Goldstone bosons in the absence of gauge couplings. This is, of course, exactly what is needed to give masses to three of the gauge bosons. To see how the gauge-boson masses come out, we first redefine the W^a fields by an $\text{SU}(2)_L$ rotation so that v takes the diagonal form $v = a1$. Then one finds the gauge-boson mass matrix (at the Born level in the electroweak couplings) takes the form (A13) with $M_{WW}^{ab} = \text{const} \times \delta_{ab}$, as in the standard model.

We do not discuss in detail what happens when the minimum of the effective potential is at a point of unequal λ_i . This possibility is not favored as a model because the unequal- λ_i surfaces have more than three dimensions. If the photon is to remain massless, this means that there must be Nambu-Goldstone bosons remaining after three have been used to provide masses to the W and Z . The Nambu-Goldstone bosons will not remain massless in the full theory, but one expects that their masses will be small compared to the 1-TeV mass scale.

4. Gauge-boson masses beyond the Born level

The $\text{U}(1)$ gauge interactions break the $\text{SU}(2)_L \times \text{SU}(2)_R$ symmetry, so the tree-level relation between the W and Z masses will not hold exactly. This also happens in the standard model, but a new effect occurs here in the case of the $(1, 1)$ representation: once the vacuum expectation value v_{im} moves off the equal- λ_i surface, the W - Z mass relation is altered. However, it is plausible that this is a small effect, since the curvature of $\mathcal{V}(\lambda_1, \lambda_2, \lambda_3)$ at the minimum should be large, characterized by a TeV mass scale. Thus, electroweak radiative corrections to the effective potential should yield only a small change in the position of the minimum, and thus in v_{im} .

- [1] S. Weinberg, Phys. Rev. D **16**, 974 (1976); **19**, 1277 (1979); L. Susskind, *ibid.* **20**, 2619 (1979).
- [2] S. Dimopoulos and L. Susskind, Nucl. Phys. **B155**, 237 (1979); E. Eichten and K. Lane, Phys. Lett. **90B**, 125 (1980); E. Eichten, I. Hinchliffe, K. D. Lane, and C. Quigg, Phys. Rev. D **34**, 1547 (1986).
- [3] For reviews, see E. Farhi and L. Susskind, Phys. Rep. **74**, 277 (1981); K. Lane, in *Elementary Particle Physics and Future Facilities*, Proceedings of the DPF Summer Study, Snowmass, Colorado, edited by R. Donaldson, R. Gustafson, and F. Paige (World Scientific, Singapore, 1982), p. 222.
- [4] Mitchell Golden and Lisa Randall, Nucl. Phys. **B361**, 3 (1991).
- [5] Michael E. Peskin and Tatsu Takeuchi, Phys. Rev. Lett. **65**, 964 (1990).
- [6] B. Holdom and J. Terning, Phys. Lett. B **247**, 88 (1990); Antonio Dobado, Doménec Espriu, and Maria L. Herrero, *ibid.* **255**, 405 (1991).
- [7] Robert Johnson, Bing-Lin Young, and Douglas W. McKay, Phys. Rev. D **42**, 3855 (1990); **43**, R17 (1991).
- [8] William J. Marciano and Jonathan L. Rosner, Phys. Rev. Lett. **65**, 2963 (1990).
- [9] D. C. Kennedy and Paul Langacker, Phys. Rev. Lett. **65**, 2967 (1990).
- [10] G. Altarelli and R. Barbieri, Phys. Lett. B **253**, 161 (1991).
- [11] E. Farhi and L. Susskind, Phys. Rev. D **20**, 3404 (1979).
- [12] Jogesh C. Pati and Abdus Salam, Phys. Rev. D **10**, 275 (1974); J. C. Pati, A. Salam, and J. Strathdee, Phys. Lett. **59B**, 265 (1975).
- [13] K. Akama, Y. Chikashige, and H. Terazawa, Phys. Rev. D **15**, 480 (1977).
- [14] H. Terazawa, Phys. Rev. D **22**, 184 (1990).
- [15] O. W. Greenberg and J. Sucher, Phys. Lett. **99B**, 339 (1981).
- [16] H. Fritzsch and G. Mandelbaum, Phys. Lett. **102B**, 319 (1981).
- [17] R. Casalbuoni and R. Gatto, Phys. Lett. **103B**, 113 (1981).
- [18] O. W. Greenberg, R. N. Mohapatra, and S. Nussinov, Phys. Lett. **148B**, 465 (1984).
- [19] Estia J. Eichten, Kenneth D. Lane, and Michael E. Peskin, Phys. Rev. Lett. **50**, 811 (1983).
- [20] G. 't Hooft, in *Recent Developments in Gauge Theories*, Proceedings of the Cargèse Summer Institute, Cargèse, France, 1979, edited by G. 't Hooft *et al.*, NATO Advanced Study Institute, Series B: Physics, Vol. 59 (Plenum, New York, 1980), p. 135.
- [21] R. Peccei, in *Selected Topics in Electroweak Interactions*, Proceedings of the Second Lake Louise Institute on New Frontiers in Particle Physics, Alberta, Canada, 1987, edited by J. M. Cameron *et al.* (World Scientific, Singapore, 1987), p. 564.
- [22] R. Sekhar Chivukula, in *High Energy Physics in the 1990's (Snowmass, 1988)*, Proceedings of the Summer Study, Snowmass, Colorado, 1988, edited by Sharon Jensen (World Scientific, Singapore, 1989); p. 131; R. Sekhar Chivukula and M. Golden, Phys. Rev. D **41**, 2795 (1990). The merits of purely left-handed Yukawa interactions of technivector mesons for suppressing unwanted flavor-changing effects have been stressed by R. Sekhar Chivukula and Howard Georgi, Phys. Lett. B **188**, 99 (1987).
- [23] S. L. Adler, Phys. Rev. **177**, 2426 (1969); J. C. Bell and R. Jackiw, Nuovo Cimento A **60**, 47 (1969).
- [24] See, e.g., Lowell S. Brown and Robert L. Goble, Phys. Rev. Lett. **20**, 346 (1968); Phys. Rev. D **4**, 723 (1971).
- [25] J. Wess and B. Zumino, Phys. Lett. **37B**, 95 (1971); William A. Bardeen, Phys. Rev. **184**, 1848 (1969); Edward Witten, Nucl. Phys. **B223**, 422 (1983).
- [26] S. Dimopoulos and L. Susskind, Nucl. Phys. **B155**, 237 (1979); H. Terazawa, Phys. Rev. D **22**, 184 (1980).
- [27] Estia J. Eichten, Ian Hinchliffe, Kenneth D. Lane, and Chris Quigg, Rev. Mod. Phys. **56**, 579 (1984); **58**, 1065(E) (1986).
- [28] G. 't Hooft, Phys. Rev. Lett. **37**, 8 (1976).
- [29] Benjamin W. Lee, C. Quigg, and H. B. Thacker, Phys. Rev. D **16**, 1519 (1977).
- [30] P. Q. Hung and H. B. Thacker, Phys. Rev. D **31**, 2866 (1985).
- [31] M. S. Chanowitz and M. K. Gaillard, Phys. Lett. **142B**, 85 (1984); Nucl. Phys. **B261**, 379 (1985); M. S. Chanowitz, in *Results and Perspectives in Particle Physics*, proceedings of the Rencontres de Physique de la Vallée d'Aoste, La Thuile, Italy, 1987, edited by M. Greco (Editions Frontières, Gif-sur-Yvette, France, 1987), p. 335, and references therein; M. Chanowitz, M. Golden, and H. Georgi, Phys. Rev. Lett. **57**, 2344 (1986); Phys. Rev. D **36**, 1490 (1987); P. Q. Hung and H. B. Thacker, *ibid.* **31**, 2866 (1985); J. D. Bjorken, in *New and Exotic Phenomena*, proceedings of the 7th Moriond Workshop, Les Arcs, France, 1987, edited by O. Fackler and J. Tran Thanh Van (Editions Frontières, Gif-sur-Yvette, France, 1987), p. 1; P. Q. Hung, T. N. Pham, and T. N. Truong, Phys. Rev. Lett. **59**, 2251 (1987).
- [32] R. Rosenfeld and J. L. Rosner, Phys. Rev. D **38**, 1530 (1988); R. Rosenfeld, *ibid.* **39**, 971 (1989); **42**, 126 (1990); Int. J. Mod. Phys. A **4**, 1999 (1989).
- [33] M. A. B. Bég, H. D. Politzer, and P. Ramond, Phys. Rev. Lett. **43**, 1710 (1979).
- [34] R. Marshak and R. Mohapatra, Phys. Lett. **91B**, 222 (1980).
- [35] Jonathan L. Rosner, Prog. Theor. Phys. **66**, 1422 (1981).
- [36] Stanley J. Brodsky and Sidney D. Drell, Phys. Rev. D **22**, 2236 (1980).
- [37] R. D. Bolton *et al.*, Phys. Rev. D **38**, 2077 (1988).
- [38] A. D. Dolgov and V. I. Zakharov, Nucl. Phys. **B27**, 525 (1971).
- [39] Haim Harari, Phys. Rev. Lett. **22**, 562 (1969); Jonathan L. Rosner, *ibid.* **22**, 689 (1969).
- [40] CELLO Collaboration, Z. Phys. C **51**, 149 (1991). For a recent example of limits which include the possibility of purely left-handed couplings see TOPAZ Collaboration, I. Adachi *et al.*, Phys. Lett. B **255**, 613 (1991).
- [41] P. N. Harriman, A. D. Martin, W. J. Stirling, and R. G. Roberts, Phys. Rev. D **42**, 798 (1990).
- [42] CDF Collaboration, F. Abe *et al.*, Phys. Rev. Lett. **67**, 2609 (1991).
- [43] J. Carr *et al.*, Phys. Rev. Lett. **51**, 627 (1983); A. E. Jodido *et al.*, Phys. Rev. D **34**, 1967 (1984); **37**, 237(E) (1988); B. Balke *et al.*, *ibid.* **37**, 587 (1988).
- [44] J. J. Hernandez *et al.*, Phys. Lett. B **239**, 1 (1990).
- [45] Richard P. Feynman and Murray Gell-Mann, Phys. Rev. **109**, 193 (1958); E. C. G. Sudarshan, and R. E. Marshak, *ibid.* **109**, 1860 (1958).
- [46] M. Ademollo and R. Gatto, Phys. Rev. Lett. **13**, 264 (1965).
- [47] Jonathan L. Rosner, Phys. Rev. D **42**, 3107 (1990).
- [48] J. Pilcher, Bull. Am. Phys. Soc. **36**, 1396 (1991).
- [49] G. Beall, M. Bander, and A. Soni, Phys. Rev. Lett. **48**, 848

- (1982); G. Ecker and W. Grimus, Nucl. Phys. **B258**, 328 (1985).
- [50] Paul Langacker and S. Uma Sankar, Phys. Rev. D **40**, 1562 (1989).
- [51] Gerard Jungman and Jonathan L. Rosner, Phys. Lett. B (to be published).
- [52] P. Sikivie, L. Susskind, M. Voloshin, and V. Zakharov, Nucl. Phys. **B173**, 189 (1980).