Signatures for technicolor

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The present experimental constraints on the existence of light charged and neutral electroweak scalar particles are summarized. Within the context of an explicit extended technicolor model, we discuss prospects for the detection of light scalars in $\bar{p}p$ collisions at energies up to 2 TeV, in e^+e^- collisions at energies up to the Z^0 mass, and in fixed-target hadron-hadron collisions. In addition the production cross sections of the other expected scalars are computed and decay signatures are explored.

I. THE SCALAR SECTOR OF ELECTROWEAK THEORY

Now that the intermediate bosons $W^{\pm}$ and Z^0 have been found,[1] the discovery and study of the scalar sector of the electroweak interaction is among the most pressing problems facing experimental high-energy physics. The value of the ratio

$$\rho \equiv M_W{}^2/M_Z{}^2\cos^2\theta_W \;, \tag{1.1}$$

determined from low-energy neutral-current observables or from the measured intermediate boson masses, is an important constraint on the mechanism of spontaneous symmetry breaking. It is determined to be unity within errors of a few percent. Although all that is required to ensure $\rho=1$ is a global "custodial SU(2)" symmetry,[2] the condition is most gracefully satisfied if the scalar sector consists only of $\mathrm{SU}(2)_L$ weak-isospin doublets. However, even for this simplest case we have no experimental information about the number of doublets and we do not know whether the scalars are elementary Higgs bosons[3] or the composite particles that arise in technicolor models.[4] Moreover, the masses and principal decay modes of the scalars are unknown or highly model dependent.

As discussed elsewhere,[5,6] the full unraveling of the structure of electroweak dynamics probably must await experimentation at multi-TeV hadron colliders, such as the proposed Superconducting Super Collider. However, much can be learned about the scalar sector at the considerably lower energies of the CERN S$\bar{p}p$S Collider and the Fermilab Tevatron. In addition, important work remains to be done at electron-positron colliders at the Z^0 energy and below. In this paper we focus on the physics of electroweak scalars which will be relevant to machines of the current generation. The specific analysis we present is motivated by the technicolor approach to dynamical symmetry breaking, but many of the results have a considerably broader range of applicability.

The basic technicolor idea must be augmented[7,8] to allow for the generation of quark and lepton masses. In these extended technicolor (ETC) models, light spin-zero particles known as technipions are associated with approximate global symmetries in the technicolor sector. The least massive of these technipions carry no color, baryon number, or lepton number, but may be electrically charged. The neutral technipions P^0 and $P^{0'}$ have masses between approximately 2 and 40 GeV/c^2, while the masses of the charged technipions $P^{\pm}$ lie between about 8 and 40 GeV/c^2 (Ref. 8). These technipions are dynamical analogs of the elementary pseudoscalars which arise in multiple-Higgs-doublet models of the scalar sector of the electroweak Lagrangian.[9] At energies far below the electroweak scale ($\sim G_F{}^{-1/2}$), the technicolor and multiple-scalar-doublet theories are quite similar.[10] Since the masses of the resultant spinless particles may be quite low, these particles are, in principle, accessible at energies now within reach.

In Sec. II of this paper we review the general features of extended technicolor models and introduce needed terminology. Then in Sec. III we summarize current limits on the existence of the lightest electroweak scalars, calculate the production cross sections for S$\bar{p}p$S and Tevatron collider experiments, and discuss detection prospects in hadron-hadron and electron-positron collisions.

In Sec. IV we turn our attention to the colored technipions expected in extended technicolor models at a mass scale of hundreds of GeV. We discuss production cross sections and decay signatures for the color-triplet leptoquarks (denoted P_3) with baryon number $\frac{1}{3}$ and lepton number -1, and for the color-octet techni-η (denoted P_8) with no net baryon or lepton number. The $(^3S_1)$ techni—vector mesons are expected to be too massive to be studied with S$\bar{p}p$S and Tevatron colliders.

We give a brief summary of our results and conclusions in Sec. V. Throughout we remark on the applicability of our results to models other than technicolor.

II. THE IDEA OF TECHNICOLOR

In the standard electroweak model, the $SU(2)_L \otimes U(1)_Y$ local gauge symmetry is spontaneously broken to $U(1)_{em}$ through the medium of auxiliary, elementary scalar fields known as Higgs bosons. The self-interactions of the Higgs scalars select a vacuum, or minimum-energy state, which does not manifest the full symmetry of the Lagrangian. In so doing, they endow the gauge bosons and the elementary fermions of the theory with masses. Three of the four auxiliary scalars introduced in the minimal one-Higgs-doublet model become the longitudinal components of W^+, W^-, and Z^0. The fourth emerges as the physical Higgs boson. In multiple-Higgs-doublet models (in which responsibilities for giving masses to the gauge bosons and to the fermions may be divided), additional physical spinless particles—both charged and neutral—arise. For example, the Peccei-Quinn mechanism[11] for avoiding *CP* nonconservation in the strong interactions requires at least two Higgs doublets. Also, the minimal supersymmetric extension of the Weinberg-Salam theory requires at least two scalar doublets.[12] In both cases there will be charged scalars in addition to the familiar neutral Higgs boson, and extra neutral scalars as well. The expectation of additional spinless particles is thus a rather general feature of extensions to the standard model.

The desire to go beyond the standard model is motivated in part by the arbitrariness and ambiguity of the elementary scalar solution. The main criticisms of the standard model concern the multitude of arbitrary parameters associated with the Higgs potential and the Yukawa couplings that generate fermion masses. One hopes for a better, more restrictive solution, with greater predictive power.

The dynamical-symmetry-breaking approach, of which technicolor theories are exemplars, is modeled upon our understanding of another manifestation of spontaneous symmetry breaking in nature, the superconducting phase transition. The macroscopic order parameter of the Ginzburg-Landau phenomenology[13] corresponds to the wave function of superconducting charges. It acquires a nonzero vacuum expectation value in the superconducting state. The microscopic Bardeen-Cooper-Schrieffer theory[14] identifies the dynamical origin of the order parameter with the formation of bound states of elementary fermions, the Cooper pairs of electrons. By analogy, the dynamics of the fundamental technicolor gauge interactions generate scalar bound states, and these play the role of the Higgs fields. The technipion decay constant F_π corresponds to the vacuum expectation value of an elementary Higgs field.

In building a technicolor model, we postulate a new set of asymptotically free gauge interactions based on the technicolor gauge group G_{TC}. The technicolor forces act on massless technifermions which transform according to a complex representation of G_{TC}. It is convenient for illustrative purposes to choose $G_{TC} = SU(N)_{TC}$, and to assign the technifermions to the fundamental **N** representation. [Another simple possibility for a technicolor group is $O(N)$. We will mention this option when it leads to qualitative differences in the phenomenology.] With these assignments the technicolor Lagrangian exhibits an exact chiral $SU(n_f)_L \otimes SU(n_f)_R$ symmetry, where n_f is the number of techniflavors. At an energy scale of order $\Lambda_{TC} \simeq 1$ TeV, the technicolor interactions become strong and the chiral symmetry is spontaneously broken down to (vector) $SU(n_f)$, the flavor group of the technifermions, through the formation of condensates of technifermions.

The simplest possibility for the electroweak interactions of the technifermions is to assign them the same $SU(2)_L \otimes U(1)_Y$ transformation properties as ordinary fermions, as left-handed doublets and right-handed singlets under $SU(2)_L$, with weak hypercharges chosen to ensure the absence of anomalies in all gauge currents. A technicolor model with n left-handed weak-isospin doublets and $2n$ right-handed weak-isospin singlets all transforming according to the same irreducible representation of color (and of technicolor) will contain $4n^2$ color-singlet pseudoscalar bound states. Of these, one is the analog of the η' in QCD. It couples to an anomalously divergent current, and so is expected to acquire a mass on the order of several hundred GeV/c^2. Three more become the longitudinal components of the electroweak gauge bosons W^+, W^-, and Z^0. The remaining $4n^2-4$ technipions are approximate Goldstone bosons associated with the spontaneous breakdown of the chiral symmetry. These occur as (n^2-1) weak-isospin triplets, with members P_k^+, P_k^0, P_k^- $(k=1,\ldots,n^2-1)$, and (n^2-1) weak-isospin singlets, denoted $P_k^{0'}$. If technifermions are assigned to several distinct color representations (for example, in analogy with the usual quarks and leptons) $\mathbf{r}_1, \mathbf{r}_2, \ldots,$ the enumeration of technipion states goes through with the replacement $n^2 \to \sum n_i{}^2$, where n_i is the number of $SU(2)_L$ doublets assigned to color representation $\mathbf{r}_i$. A similar generalization applies for the case of several technicolor representations.

The interactions of technipions with the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge bosons occur dynamically through technifermion loops. At energies well below the characteristic technicolor scale, the technipions may be regarded as pointlike, so their effective couplings to gauge bosons may be calculated reliably using effective Lagrangian[15] or current-algebra[16] methods. The effective Lagrangian in the general case has been treated by Chadha and Peskin.[15] For the particular case of interest here, the terms involving the coupling of technipions to gauge bosons are given by

$$\mathscr{L}_{\rm eff}^{(\rm gauge)}=\sum_{k=1}^{n^2-1}\{(\partial_\mu P_k^-)(\partial^\mu P_k^+)-M_{+k}{}^2P_k^-P_k^+ +\tfrac{1}{2}[(\partial_\mu P_k^0)^2-M_{0k}{}^2P_k^0P_k^0]\}$$
$$+\sum_{k=1}^{n^2-1} ie(A^\mu+Z^\mu\cot 2\theta_W)P_k^-\overleftrightarrow{\partial}_\mu P_k^+ +\sum_{k=1}^{n^2-1} ig_W W^{+\mu}P_k^-\overleftrightarrow{\partial}_\mu P_k^0+\text{H.c.}$$
$$+\frac{e^2}{2}(A^\mu A_\mu-Z^\mu Z_\mu)\sum_{k=1}^{n^2-1}P_k^-P_k^+ +O(P^3)\ , \tag{2.1}$$

while the $P_k^{0'}$ fields do not interact because they are electroweak singlets. Couplings between technipions and gauge bosons induced by anomalies are not included in Eq. (2.1), but will be discussed in Sec. III C.

In the absence of extended technicolor interactions the masses of all neutral technipions are zero.[8] The charged technipions receive a common mass contribution from electroweak gauge-boson exchange after symmetry-breaking effects are included. In an $SU(N)_{TC}$ theory this mass term is given by[8,15,16]

$$M_{\rm EW}{}^2(P^\pm)=\frac{3\alpha_{\rm em}}{4\pi}m_Z{}^2\ln(\Lambda^2/m_Z{}^2)\ . \tag{2.2}$$

Choosing the technicolor scale parameter $\Lambda=(G_F/\sqrt{2})^{-1/2}$ we obtain $M_{\rm EW}(P^\pm)\approx 6$ GeV/c^2. For an $SO(N)_{TC}$ theory, there is an additional electroweak contribution[15,16] to the mass. As an example, in the $SO(16)_{TC}$ model,[15] the total electroweak contribution is

$$M_{\rm EW}{}^2(P^\pm)=\frac{3\alpha_{\rm em}}{4\pi}m_Z{}^2\left[\ln(\Lambda^2/m_Z{}^2)+\frac{M^4}{12\pi^2F_\pi{}^4}\ln(\Lambda^2/\alpha_s M^2)\right]\ , \tag{2.3}$$

where $M\approx 500$ GeV/c^2 is the scale of the one-photon-exchange contribution and $F_\pi\approx 124$ GeV is the technipion decay constant. Choosing $\alpha_s\approx 0.1$, we now estimate $M_{\rm EW}(P^\pm)\approx 8$ GeV/c^2. These contributions represent an approximate lower bound on the mass of the lightest charged technipion which is insensitive to details of models.

An elementary Higgs-boson model with $d=n^2$ complex weak-isospin doublets provides the corresponding linear realization of these chiral symmetries. In general such a model implies $(d-1)$ charge ± 1 spinless Higgs particles and $2d-1$ neutral Higgs bosons. The masses of these elementary Higgs bosons are not constrained by the Goldstone theorem. If in the elementary scalar model we set to $O(1$ TeV/$c^2)$ the masses of the $d=n^2$ neutral scalars which are the generalizations of σ field in the usual σ model, then the low-energy interactions of the remaining $3(n^2-1)$ Higgs bosons are identical to those of the technipion case given in Eq. (2.1).

In the technicolor scenario as we have described it so far, there is no analog of the Yukawa couplings between Higgs fields and quarks or leptons which generate fermion masses in the standard model. Extended technicolor[7,8] provides a mechanism for endowing the ordinary quarks and leptons with masses. This is accomplished by embedding the technicolor gauge group $G_{\rm TC}$ into a larger extended technicolor gauge group $G_{\rm ETC}\supset G_{\rm TC}$ which couples quarks and leptons to the technifermions. It is assumed that the breakdown $G_{\rm ETC}\to G_{\rm TC}$ occurs at the energy scale

$$\Lambda_{\rm ETC}\sim 30-300\ \text{TeV}\ . \tag{2.4}$$

Then, the exchange of an ETC gauge boson of mass $M_{\rm ETC}$ generates bare-quark and lepton masses of order

$$m\sim g_{\rm ETC}{}^2\Lambda_{\rm TC}{}^3/M_{\rm ETC}{}^2=\Lambda_{\rm TC}{}^3/\Lambda_{\rm ETC}{}^2\ . \tag{2.5}$$

Unfortunately, no one has succeeded in constructing an extended technicolor model which is at all realistic.

The inclusion of ETC interactions has three effects on the interaction Lagrangian (2.1): First, the exchange of massive ETC gauge bosons contributes to the masses for both charged and neutral scalars. This contribution will be discussed in Sec. III. Second, the neutral isovector states P_k^0 and the isoscalar states $P_k^{0'}$ will be mixed by ETC interactions. The charged-current interaction in (2.1) will be modified to

$$\sum_{j,l=1}^{n^2-1} ig_W W^{+\mu}P_j^-\overleftrightarrow{\partial}_\mu \mathscr{U}_{jl}^\dagger\left[\sum_{k=1}^{n^2-1}\mathscr{O}_{lk}P_k^0+\sum_{k=n^2}^{2n^2-2}\mathscr{O}_{lk}P_{k-n^2+1}^{0'}\right], \tag{2.6}$$

where $\mathscr{U}$ is the $(n^2-1)\times(n^2-1)$ special unitary matrix that diagonalizes the (mass)^2 matrix of the charged pseudoscalars, whereas $\mathscr{O}$ is the $2(n^2-1)\times 2(n^2-1)$ real orthogonal matrix that diagonalizes the (mass)^2 matrix of the neutral pseudoscalars having electroweak interactions. Third, a very small $Z^0P^0P^{0'}$ coupling will be induced by electroweak symmetry breaking.

The couplings of technipions to ordinary fermions are model dependent and generally quite complicated. The same may be said for the couplings between fermions and the elementary scalars in multiple Higgs models. In the absence of a standard ETC model, we shall simply parametrize these couplings and indicate how they may be constrained by experimental data. We write

$$\mathscr{L}_{\text{eff}}^{(\text{fermions})}=i\sum_{\substack{\text{generations}\\ i,j}}\sum_{k}\left[\bar{q}\,'_{iL}\hat{U}^{\,k}_{ij}u_{jR}\begin{bmatrix}(P_k^{0'}+P_k^0)/\sqrt{2}\\ P_k^-\end{bmatrix}+\bar{q}\,'_{iL}\hat{D}^{\,k}_{ij}d_{jR}\begin{bmatrix}P_k^+\\ (P_k^{0'}-P_k^0)/\sqrt{2}\end{bmatrix}\right.$$

$$\left.+\bar{l}_{iL}\hat{L}^{\,k}_{ij}e_{jR}\begin{bmatrix}P_k^+\\ (P_k^{0'}-P_k^0)/\sqrt{2}\end{bmatrix}\right]+\text{H.c.}\,, \tag{2.7}$$

where primes on quark fields denote the electroweak flavor eigenstate fermion fields which are related to the (unprimed) mass eigenstates by the Kobayashi-Mashkawa matrix K, and the matrices $\hat{U}$, $\hat{D}$, and $\hat{L}$ have dimensions of mass. To obtain a canonical form for the coupling matrices, we reexpress them as

$$\hat{U}=V_{LU}^{\dagger}UV_{RU},\;\; \hat{D}=V_{LD}^{\dagger}DV_{RD},\;\; \hat{L}=V_{LL}^{\dagger}LV_{RL}\;, \tag{2.8}$$

where U, D, and L are diagonal matrices and the V's are special unitary matrices. For the case of three quark and lepton generations, each diagonal matrix has three complex diagonal elements.

To make contact with the common assumption that the scalar-boson couplings are proportional to the fermion mass, we let

$$U^k=\begin{bmatrix}\lambda_{Uu}^k m_u & 0 & 0\\ 0 & \lambda_{Uc}^k m_c & 0\\ 0 & 0 & \lambda_{Ut}^k m_t\end{bmatrix}\frac{1}{F_\pi} \tag{2.9}$$

(and similarly for D^k and L^k), where the λ's are in general complex numbers and F_π is the pseudoscalar decay constant of the technipions. For the simple case of mass-proportional couplings,

$$\lambda_{Xl}^k=1 \;\text{ and }\; V_{\chi X}=1\;. \tag{2.10}$$

In general the unitary matrices $V_{\chi X}$ can be expressed in terms of six independent Euler angles if we redefine the phases of the left-handed and right-handed fermions to absorb into the diagonal matrices U, D, and L two phases of $V_{\chi X}$. We shall indicate below how experiments may be used to constrain some of the matrix elements. However, the sample calculations we give will generally be based on the simplest possibility (2.10).

III. COLOR-SINGLET ELECTROWEAK SCALARS

A. Theoretical expectations

To discuss the properties expected of the technipions, it is convenient to specialize to the simplest quasirealistic model, introduced by Farhi and Susskind,[17] in which the elementary technifermions are a doublet of color-triplet techniquarks (U,D) and a doublet of color-singlet technileptons (N,E) (Ref. 18). Similar results may be expected in any extended technicolor model of the class discussed in Sec. II.

In the absence of extended technicolor interactions the neutral technipions P^0 and $P^{0'}$ remain massless, while the charged technipions P^+ and P^- acquire electroweak masses of a few GeV/c^2. When ETC interactions are included, the technipion masses have been estimated as[8,16]

$$\begin{aligned}&8<M(P^{\pm})<40\text{ GeV}/c^2\,,\\ &2<M(P^0,P^{0'})<40\text{ GeV}/c^2\,.\end{aligned} \tag{3.1}$$

The estimates are not rigorous bounds but represent an educated guess; however, they are rather insensitive to details of the model. It is worth emphasizing that the upper end of the range 40 GeV/c^2 is considerably higher than the value of 14 GeV/c^2 sometimes cited in the literature.[19] The lower value is the basis for experimental claims[20] that technicolor has been ruled out, a verdict we regard as premature.

The extended technicolor interaction couples technifermions to quarks and leptons, and so governs the decays of technipions into ordinary matter. For light color-singlet technipions, these are the principal decay models. If, like the neutral Higgs boson of the minimal one-doublet model, the technipions couple to mass (γ_5 coupling), the decays occur at a (semiweak) rate of

$$\Gamma(P\to f_i\bar{f}_j)=\frac{G_F p[M(P^{0'})^2-(m_i+m_j)^2]}{4\pi}C_{ij}\,|\mathscr{M}_{ij}|^2\,, \tag{3.2}$$

where G_F is the Fermi constant, p is the momentum of the products in the technipion rest frame, and C_{ij} is a color factor which is equal to 3 for the decay of a color singlet into quarks and 1 otherwise. The quantity $|\mathscr{M}_{ij}|^2$ is the square of a model-dependent dimensionless matrix element. For up quarks,

$$\mathscr{M}_{ij}=\frac{(\hat{U}+\hat{U}^{\dagger})_{ij}}{\sqrt{2}}\frac{1}{G_F{}^{1/2}M(P^{0'})}\;.$$

The only possible exception to the dominance of $f\bar{f}$ modes is the decay of $P^{0'}$ into two gluons, for which the partial width in SU(N)$_{\text{TC}}$ is

$$\Gamma(P^{0'}\to gg)\approx\frac{\alpha_s{}^2}{6\pi^3}\frac{M(P^{0'})^3}{F_\pi{}^2}\left[\frac{N}{4}\right]^2\,. \tag{3.3}$$

This becomes comparable to the rate for uninhibited $P^{0'}\to b\bar{b}$ decay for $M(P^{0'})\approx 40$ GeV/c^2. In the Farhi-Susskind model, the technipion decay constant is $F_\pi=(8G_F/\sqrt{2})^{-1/2}\approx 124$ GeV.

Because of the anticipated importance of $f\bar{f}$ decay modes, it is of great interest to know what the extended technicolor interactions are. It is precisely in this respect that the existing models fail to provide reliable guidance.

TABLE I. Principal decay modes of color-singlet technipions if $Pf_1\bar{f}_2$ couplings are proportional to fermion mass.

Technipion	Principal decay modes
P^+	$t\bar{b}, c\bar{b}, c\bar{s}, \tau^+\nu_\tau$
P^0	$b\bar{b}, c\bar{c}, \tau^+\tau^-$
$P^{0'}$	$b\bar{b}, c\bar{c}, \tau^+\tau^-; gg$

What can be said in general about technipion couplings to ordinary fermions is that they are parity violating,[8] and possibly *CP* violating as well.[21] This fact may lead to many interesting investigations if technipions are detected. We put aside such questions and focus on the initial search.

To the extent that the couplings of technipions to quarks and leptons are unknown, *nothing can be said about the expected branching ratios for technipion decay.* We can make reasonable guesses, but a meaningful search must cover all the allowed possibilities. According to the conventional wisdom, which is inspired by analogy with the minimal electroweak model, the technipions couple essentially to fermion mass. Bearing in mind that this tendency to couple to mass can be evaded even in the models with two or more elementary Higgs doublets, we list in Table I the expected major decay modes of technipions.

B. Experimental constraints

1. General remarks

Restrictions on the spectrum and interactions of color-singlet technipions currently are derived from direct searches in e^+e^- annihilations and from limits on rare decays that might exist in ETC models. We shall review each of these categories in turn. A potentially stringent lower bound on the mass of the charged technipion $P^\pm$ has long been recognized, but merits restatement. If the decay

$$t \to P^+ + (b \text{ or } s \text{ or } d) \tag{3.4}$$

is kinematically allowed, it proceeds semiweakly at a rate given by

$$\Gamma(t\to P^+q) \approx \frac{|\mathcal{M}_{tq}|^2[m_t^2+m_q^2-M(P^+)^2]p}{4\pi F_\pi^2} , \tag{3.5}$$

where

$$p=\frac{[m_t^2-(m_q+M(P^+))^2]^{1/2}[m_t^2-(m_q-M(P^+))^2]^{1/2}}{2m_t} \tag{3.6}$$

is the momentum of the products in the t-quark rest frame and $\mathcal{M}_{tq}=(F_\pi/2m_t)(\hat{U}^\dagger K)_{tq}$. In the Farhi-Susskind model, the technipion decay constant is

$$F_\pi=(8G_F/\sqrt{2})^{-1/2} , \tag{3.7}$$

so that

$$\Gamma(t\to P^+q) \approx \frac{G_F\sqrt{2}}{\pi}|\mathcal{M}_{tq}|^2(m_t^2+m_q^2-m_P^2)p . \tag{3.8}$$

If the P^+ has more or less "conventional" couplings to quarks and leptons, then $|\mathcal{M}_{tb}|=O(1)$ and this semiweak decay will swamp the normal weak decays. Furthermore, P^+ will decay only rarely to $e^+\nu_e$ and $\mu^+\nu_\mu$, making the semileptonic decays

$$t\to be^+\nu_e, b\mu^+\nu_\mu , \tag{3.9}$$

unobservably rare. *Not* seeing these decays with significant and equal branching fractions is evidence for a P^+ with conventional couplings.

The semileptonic branching ratio is in principle measurable in hadron collider experiments as follows. From the observed number of $W\to e\nu$ decays one may calculate for a given t-quark mass the number of produced $W\to t\bar{b}$ events. A comparison with the number of observed

$$W\to t\bar{b},\quad t \to bl^+\nu \tag{3.10}$$

events then yields the semileptonic branching ratio of the t quark. In the standard model, in the absence of nonleptonic enhancement, the expected semileptonic branching ratio is

$$B(t\to bl^+\nu)_{SM}=\Gamma(t\to bl^+\nu)/\Gamma(t\to \text{all})|_{SM} = \tfrac{1}{9} . \tag{3.11}$$

If the only nonstandard decays of top are of the class (3.4) and P^+ does not decay appreciably into $l^+\nu$, then a measurement of the branching ratio implies a measurement of the $t\to P^+b$ decay rate of

$$\frac{\Gamma(t\to P^+b)}{\Gamma(t\to bl^+\nu)}=\frac{1}{B(t\to bl^+\nu)}-\frac{1}{B(t\to bl^+\nu)_{SM}} . \tag{3.12}$$

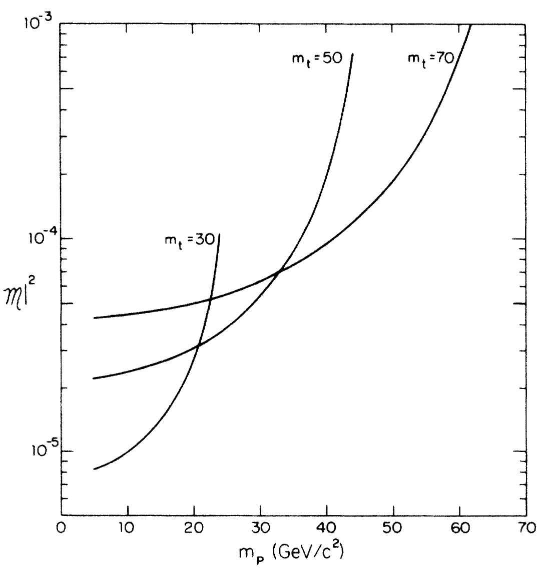

FIG. 1. Technipion coupling to fermions for which the rate for the decay $t\to P^+b$ given by (3.8) equals the rate $\Gamma(t\to bl^+\nu)=G_F^2m_t^5/192\pi^3$.

To show what sort of constraints on technipion masses and couplings can be obtained in this way, we display in Fig. 1 the value of $|\mathcal{M}_{tb}|^2$ for which

$$\Gamma(t\to P^+b)=\Gamma(t\to bl^+\nu)_{\rm SM}\ , \tag{3.13}$$

for various values of m_t and $M(P)$, and with $m_b=5$ GeV/c^2. Analogous results apply for $|\mathcal{M}_{ts}|^2$ and $|\mathcal{M}_{td}|^2$. One might guess that establishment of the t quark by the chain (3.10) will require

$$B(t\to bl^+\nu)\gtrsim \tfrac{1}{2}B(t\to bl^+\nu)_{\rm SM}\ , \tag{3.14}$$

for which

$$\Gamma(t\to P^+b)\lesssim 10\Gamma(t\to bl^+\nu)_{\rm SM}\ , \tag{3.15}$$

with the corresponding implications for technipion masses and couplings. If the $bl^+\nu$ mode is observed at the canonical rate, one may still wish to verify that it occurs via the canonical path by checking e-μ-τ universality and the details of decay distributions, before concluding that technicolor has been ruled out for $M(P^\pm)<m_t-m_b$.

If the t-quark mass exceeds the intermediate-boson mass, the technipion masses and couplings may be constrained in a comparison of the branching ratios for the semiweak decays $t\to P^+q$ and $t\to W^+b$. The latter occurs at a rate

$$\Gamma(t\to W^+b)=\frac{G_F}{4\pi\sqrt{2}}\frac{[(m_t^2-m_b^2)^2+m_W^2(m_t^2+m_b^2)-2m_W^4]p}{m_t^2}\ , \tag{3.16}$$

where p is the momentum of the products (W and b) in the t-quark rest frame. In this circumstance, there is no built-in calibration of the number of t quarks produced.

2. Electron-positron annihilations

Several searches for charged spinless particles have been carried out in electron-positron annihilations into hadrons. The most elementary method is to search for an excess yield of hadronic events, beyond what is expected from pair production of the known quarks. The contribution to be expected,

$$\frac{\sigma(e^+e^-\to P^+P^-)}{\sigma(e^+e^-\to\mu^+\mu^-)}=\frac{[1-4M(P)^2/s]^{3/2}}{4} \tag{3.17}$$

is small [asymptotically contributing $\frac{1}{4}$ unit to $R=\sigma(e^+e^-\to\text{hadrons})/\sigma(e^+e^-\to\mu^+\mu^-)$], and turns on only slowly above threshold. Existing measurements[22] are unable to exclude charged scalars as light as about 10 GeV/c^2. One must therefore design searches which are sensitive to specific final states.

The results of three such investigations have been published. The JADE Collaboration[23] has looked for charged scalars in the chains

$$e^+e^-\to P^+P^-\to(\tau\nu)(\tau\nu)\ \text{or}\ (\tau\nu)\ (\text{hadrons})\ . \tag{3.18}$$

They are sensitive to hadronic decays into "light" $u\bar d$, $u\bar s$, $c\bar d$, $c\bar s$ quark pairs, and (above $c\bar b$ threshold) to $c\bar b$ pairs as well. Their search places upper limits on the branching ratio $B(P\to\tau\nu)$ of 4–11% in the mass range $4<m_{P\pm}<12$ GeV/c^2. The TASSO Collaboration[24] has searched for four-jet events characteristic of the process

$$e^+e^-\to P^+P^-\to(\text{hadrons})(\text{hadrons})\ . \tag{3.19}$$

They are sensitive to light-quark pairs (the case of $c\bar s$ is treated explicitly), but apparently have no sensitivity to $c\bar b$ pairs. Taken together with the JADE limits, the TASSO results rule out a charged scalar decaying into $\tau\nu$ or hadrons in the mass range

$$5\ \text{GeV}/c^2\le M(P^\pm)\le c\bar b\ \text{threshold}$$
$$\approx 7\text{–}8\ \text{GeV}/c^2\ . \tag{3.20}$$

For masses up to 13 GeV/c^2, these experiments exclude charged scalars decaying into $\tau\nu$ or light quarks. It is worth emphasizing that in this regime, the $c\bar b$ decay mode might well be dominant, and that this possibility has not been tested by any experiment. Finally, the MARK-J Collaboration[25] quotes a bound of

$$M(P^\pm)>17\ \text{GeV}/c^2\quad\text{for}\ B(P\to\tau\nu)>\tfrac{1}{4}\ . \tag{3.21}$$

The experimental situation is summarized in Fig. 2.

If the couplings are controlled by mass alone, then the decay $P^+\to c\bar b$ will dominate if $M(P^\pm)>8$ GeV/c^2. The other channels will have small branching ratios. This is precisely the case which is not ruled out by existing data. Clearly, a more thorough search is warranted.

The techniques discussed in this section can also be applied to the searches planned at the CERN Large Electron Positron (LEP) Collider, and the Stanford Linear Collider

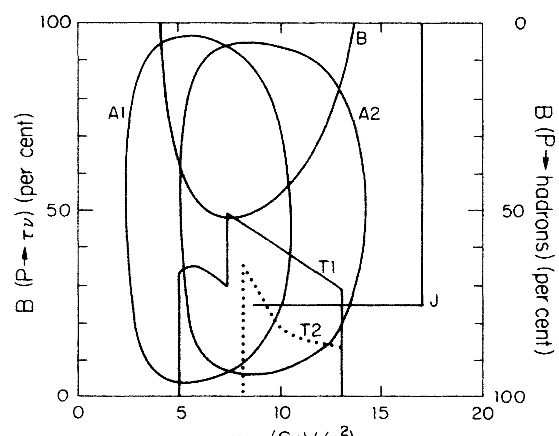

FIG. 2. Limits on the branching ratio $B(P^\pm\to\tau\nu)$ as a function of charged scalar mass. Curves $A1$, $A2$, and B represent the JADE limits (Ref. 23) for the final states $(\tau\nu)$ (light hadrons), $(\tau\nu)$ (hadrons), and $(\tau\nu)$ $(\tau\nu)$, respectively. Curves $T1$ and $T2$ depict the TASSO (Ref. 24) limits for the $(c\bar s)(\bar cs)$ final state and for equal mixtures of $c\bar s$ and $c\bar b$, respectively. The Mark J limit (Ref. 25) is shown as curve J.

(SLC). We will discuss the capabilities of these devices and their complementarity to experiments at hadron colliders in Sec. IV.

3. *Rare processes*

The experimental absence of flavor-changing neutral-current (FCNC) processes imposes strong constraints on the possible couplings of technipions to quarks and leptons. The analysis of these constraints is complicated in any realistic model by the existence of a whole spectrum of light technipions P_k^+, P_k^0, P_k^-, and $P_k^{0'}$ whose effects must be disentangled.

In extended technicolor (ETC) models there is an additional complication. The leading contribution to FCNC processes arising from the exchange of massive ETC gauge bosons is, at least formally, the same order of magnitude as the effects induced by the exchange of neutral technipions. This is easy to show, as follows. The interaction [shown in Fig. 3(a)] associated with P^0 exchange gives rise to an amplitude

$$\mathscr{M}_{\rm PGB} \sim \frac{1}{M(P^0)^2} \frac{(m_1+m_2)(m_3+m_4)}{4F_\pi^2}(A_{12}A_{34}^\dagger) , \quad (3.22)$$

where the mixing matrix elements A_{ij} are naturally of order unity. The corresponding amplitude driven by ETC exchange [Fig. 3(b)] is

$$\mathscr{M}_{\rm ETC} \sim g_{\rm ETC}{}^2/M_{\rm ETC}{}^2 . \quad (3.23)$$

Now recalling[8] that

$$M(P^0)^2F_\pi{}^2 \sim g_{\rm ETC}{}^2\Lambda_{\rm TC}{}^6/M_{\rm ETC}{}^2 \quad (3.24)$$

and that

$$m_i \sim g_{\rm ETC}{}^2\Lambda_{\rm TC}{}^3/M_{\rm ETC}{}^2 , \quad (2.5')$$

we have

$$\frac{m_i{}^2}{M(P^0)^2F_\pi{}^2} \sim \frac{M_{\rm ETC}{}^2}{g_{\rm ETC}{}^2\Lambda_{\rm TC}{}^6}\frac{g_{\rm ETC}{}^4\Lambda_{\rm TC}{}^6}{M_{\rm ETC}{}^4} \sim \frac{g_{\rm ETC}{}^2}{M_{\rm ETC}{}^2} , \quad (3.25)$$

which establishes that (3.22) and (3.23) are the same order.

Because of this similarity, the constraints on technipion properties inferred from experimental limits on rare processes are highly model dependent. Furthermore the lack of an obvious analog of the Glashow-Iliopoulos-Maiani (GIM) mechanism[26] is precisely the feature of all known ETC models that makes them phenomenologically problematical.[8,27] Recently several attempts have been made to construct a GIM-like mechanism for ETC theories.[28,29] However, no proposal has yet been a complete success. Another possibility is to raise the ETC scale itself.[30,31]

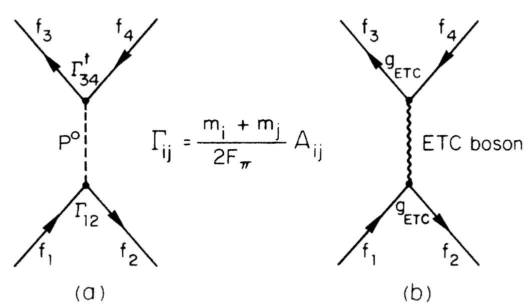

FIG. 3. Contribution of (a) neutral technipion exchange (b) ETC gauge-boson exchange to fermion-antifermion scattering.

With these warnings given, we will proceed to make some simplifying assumptions which will lead to rough estimates of the constraints experiments impose on technipion masses and couplings. In so doing, we are *assuming* that a solution can be found to the problem of FCNC process mediated by ETC gauge bosons. Then, ignoring such contributions, we consider only the contributions of a single set of light scalars, which we imagine to be the lightest charged and neutral technipions in the spectrum.

A lower bound $M(P^\pm) > 6$ GeV/c^2 on the masses of the charged technipions is implied by the well-understood contributions due to electroweak symmetry breaking in the standard model,[8,14,16] which are independent of any ETC model. Consequently the $P^\pm$ is too massive to be a decay product of any of the established quarks (u,d,s,c,b) or leptons $(\nu_e,\nu_\mu,\nu_\tau,e,\mu,\tau)$. Transitions which could be mediated by virtual $P^\pm$ may also proceed by the usual charged current weak interactions. No significant limits on the mass or couplings of the charged technipions can be derived from the decay rates of the established quarks and leptons. Only the semiweak decay of the top quark, discussed in Sec. III B 1, would provide useful constraints.

For neutral technipions, on the other hand, rare processes do impose significant restrictions on the properties of the lightest scalar, P^0. This situation comes about for two reasons. First, the mass of P^0 arises solely from ETC interactions,[8] so no model-independent (theoretical) lower bound can be given. Indeed, the P^0 mass could be as small as a few GeV/c^2 (Ref. 32). Second, standard-model backgrounds to FCNC processes are highly suppressed by the GIM mechanism.

The interactions (2.7) of the lightest neutral pseudo-Goldstone boson (PGB) with quarks and leptons can be reexpressed as

$$\begin{aligned}\mathscr{L}_I = i & \left[\frac{m_{u_i}+m_{u_j}}{2F_\pi}\right]\bar{u}_{iL}A_{ij}u_{jR}P^0 \\ & + i\left[\frac{m_{d_i}+m_{d_j}}{2F_\pi}\right]\bar{d}_{iL}B_{ij}d_{jR}P^0 \\ & + i\left[\frac{m_{e_i}+m_{e_j}}{2F_\pi}\right]\bar{e}_{iL}C_{ij}e_{jR}P^0 + \text{H.c.} , \end{aligned} \quad (3.26)$$

where A, B, and C are dimensionless mixing matrices.

The "rare process" which leads to the most stringent bounds on these couplings is the K_L-K_S mass difference. The requirement that the P^0-exchange contribution evaluated using vacuum saturation for the hadronic matrix element[33] not exceed the observed mass difference leads to the constraint

$$\frac{M(P^0)}{|\,\mathrm{Re}B_{sd}\,|} \gtrsim 8 \text{ TeV}/c^2 \quad (3.27)$$

for $m_s = 140$ MeV/c^2 and $F_\pi = 124$ GeV. Since we expect $M(P^0) \lesssim 40$ GeV/c^2, this constraint in turn implies

$$|\mathrm{Re}B_{sd}| \lesssim 5\times 10^{-3} . \tag{3.28}$$

Other constraints involving B_{sd} arise from limits on FCNC semileptonic decays of kaons. The most potent of these derives from the limit[34]

$$\frac{\Gamma(K^+\to\pi^+\mu^+e^-)}{\Gamma(K^+\to\mathrm{all})} \lesssim 7\times 10^{-9} , \tag{3.29}$$

which implies

$$\frac{M(P^0)^2}{|B_{sd}|\,|C_{\mu e}|} \gtrsim (32\ \mathrm{GeV}/c^2)^2 . \tag{3.30}$$

Another process of potential interest is the decay $K_L\to\mu e$. The limit currently accepted by the Particle Data Group[35] is[36]

$$\frac{\Gamma(K_L\to\mu e)}{\Gamma(K_L\to\mathrm{all})} \lesssim 6.3\times 10^{-6} , \tag{3.31}$$

which implies the constraint

$$\frac{M(P^0)^2}{|\mathrm{Re}B_{sd}|\,|\mathrm{Re}C_{\mu e}|} \gtrsim (160\ \mathrm{GeV}/c^2)^2 . \tag{3.32}$$

Given the earlier result (3.27), this is unlikely to lead to a meaningful constraint on $|\mathrm{Re}C_{\mu e}|$. A stronger published limit[37] of

$$\frac{\Gamma(K_L\to\mu e)}{\Gamma(K_L\to\mathrm{all})} < 2\times 10^{-9} \tag{3.33}$$

would increase the right-hand side of (3.32) to (1.2 $\mathrm{TeV}/c^2)^2$, and an experiment in progress[38] at Brookhaven with projected sensitivity of 10^{-10} would further increase the bound to $(2.5\ \mathrm{TeV}/c^2)^2$.

Additional restrictions on the couplings of P^0 to leptons are implied by limits on rare leptonic transitions. The bound[39]

$$\frac{\Gamma(\mu^-\to e^-e^+e^-)}{\Gamma(\mu^-\to\mathrm{all})} < 2.4\times 10^{-12} \tag{3.34}$$

corresponds to the restriction

$$\frac{M(P^0)^2}{|C_{\mu e}|\,|C_{ee}|} \gtrsim (280\ \mathrm{GeV}/c^2)^2 . \tag{3.35}$$

The upper limit[40]

$$\frac{\Gamma(\mu\to e\gamma)}{\Gamma(\mu\to\mathrm{all})} < 1.7\times 10^{-10} \tag{3.36}$$

requires only that

$$\frac{M(P^0)^2}{|C_{\mu\mu}|\,|C_{\mu e}|} \gtrsim (23\ \mathrm{GeV}/c^2) . \tag{3.37}$$

The atomic transition $\mu^- + {}^{32}\mathrm{S}\to e^- + {}^{32}\mathrm{S}$ does not give a strong bound on P^0 couplings because the (scalar and pseudoscalar) interaction does not act coherently. In principle, constraints may be deduced for the couplings of P^0 to c, b, and τ. At present such restrictions are not very severe because the experimental limits on rare processes involving these fermions are not yet strong.

Direct searches for P^0 in vector-meson decays[41] $V^0\to\gamma P^0$ provide the most stringent limits on the diagonal couplings of P^0 to heavy quarks. The bound from the measured inclusive photon spectrum is dependent on photon energy. It is reasonable[42] to assume

$$\frac{\Gamma(\psi\to\gamma P^0)}{\Gamma(\psi\to\mathrm{all})} \lesssim 10^{-4} , \tag{3.38}$$

which implies a constraint

$$|A_{cc}|\,[1-M(P^0)^2/M(\psi)^2]^{1/2} \lesssim 0.55 . \tag{3.39}$$

Similarly, the bound from the inclusive photon spectrum on the Υ resonance suggests[43] that

$$\frac{\Gamma(\Upsilon\to\gamma P^0)}{\Gamma(\Upsilon\to\mathrm{all})} \lesssim 10^{-3} , \tag{3.40}$$

which implies that

$$|B_{bb}|\,[1-M(P^0)^2/M(\Upsilon)^2]^{1/2} \lesssim 0.88 . \tag{3.41}$$

Bounds on the flavor-changing neutral couplings of heavy quarks follow from the nonobservation of heavy-meson decays involving real P^0's. The partial width for the decay of a heavy quark Q into P^0 and a (massless) light quark q is

$$\Gamma(Q\to P^0 q) = \frac{(|A_{Qq}|^2 \text{ or } |B_{Qq}|^2)[m_Q{}^2 - M(P^0)^2]^2}{64\pi F_\pi{}^2 m_Q} . \tag{3.42}$$

The agreement of the semileptonic branching ratios or lifetimes of c and b mesons with standard model predictions may then be used to bound the decay rates (3.42). (Similar arguments were put forward for the t-quark case in Sec. III B 1.) If we require that

$$\Gamma(c\to P^0 u) < 10\Gamma(c\to s l^+\nu)_{\mathrm{SM}} , \tag{3.43}$$

we conclude that

$$|A_{cu}|\,[1-M(P^0)^2/m_c{}^2] \lesssim 2\times 10^{-3} . \tag{3.44}$$

Similarly, the condition

$$\Gamma(b\to P^0 q) < 10\Gamma(b\to c l^-\bar\nu)_{\mathrm{SM}} \tag{3.45}$$

leads to the bounds[44]

$$\left\{\begin{matrix} |B_{bs}| \\ |B_{bd}| \end{matrix}\right\} \times [1-M(P^0)^2/m_b{}^2] \lesssim 4.2\times 10^{-4} . \tag{3.46}$$

In summary, rare processes provide no meaningful constraints on $P^\pm$ masses and couplings. For P^0, the flavor-changing coupling $|\mathrm{Re}B_{sd}|$ must be very small ($\lesssim 5\times 10^{-3}$) for any mass in the expected range. The couplings $|B_{bs}|$ and $|B_{bd}|$ are restricted for $M(P^0)\lesssim 5$ GeV/c^2, and $|A_{cu}|$ is restricted in the unlikely event that $M(P^0)\lesssim 1.5\ \mathrm{GeV}/c^2$. The only leptonic coupling combination which is constrained is $(|C_{\mu e}||C_{ee}|)^{1/2} \lesssim 0.14$. Other couplings are not seriously constrained by current data.

C. Prospects for detection in $\bar pp$ collisions

The principal sources of light, color-singlet technipions in $\bar pp$ collisions are (1) the production of the weak-isospin-singlet states $P_k^{0'}$ by the gluon fusion mechanism, (2) pair production of $P_j^{\pm}P_k^0$ through the production of real or virtual $W^{\pm}$ bosons, (3) production of $P_k^+P_k^-$ pairs by the Drell-Yan mechanism, especially near the Z^0 pole, and (4) production of $P_k^{\pm}$ in semiweak decays of heavy quarks. We shall consider each of these in turn, adopting for definiteness the one-generation Farhi-Susskind model[17] in numerical examples.

The production of a single technipion P is governed by its coupling to a pair of a gauge bosons B_1 and B_2. This coupling arises from a triangle (anomaly) graph containing technifermions, analogous to the graph responsible for the decay $\pi^0 \to \gamma\gamma$. The amplitude for the PB_1B_2 coupling is[45]

$$\mathscr{A}_{PB_1B_2} = \frac{S_{PB_1B_2}}{8\pi^2\sqrt{2}F_\pi} \epsilon_{\mu\nu\lambda\rho}\epsilon_1^\mu \epsilon_2^\nu p_1^\lambda p_2^\rho , \tag{3.47}$$

where the triangle anomaly factor is

$$S_{PB_1B_2} = \frac{g_1 g_2}{2} \mathrm{Tr}(Q_P\{Q_1, Q_2\}) . \tag{3.48}$$

Here g_1 and g_2 are the gauge coupling constants, Q_1 and Q_2 are the gauge charges, or generators, corresponding to the gauge bosons, and Q_P is the chiral generator[46] of the technipion. The contributions from different gauge-boson helicity states are summed separately in the trace. The anomaly factors in the Farhi-Susskind model are given in Table IV of Ref. 5. These results lead to the following approximate decay rates valid when the product masses are negligible:

$$\Gamma(B_1 \to PB_2) = \frac{M(B_1)^3}{96\pi} \left| \frac{S_{PB_1B_2}}{8\pi^2\sqrt{2}F_\pi} \right|^2 , \tag{3.49}$$

$$\Gamma(P \to B_1B_2) = (1+\delta_{B_1B_2}) \frac{M(P)^3}{32\pi} \left| \frac{S_{PB_1B_2}}{8\pi^2\sqrt{2}F_\pi} \right|^2 . \tag{3.50}$$

The PB_1B_2 coupling is of experimental interest only for the production of neutral technipions, because the charged technipions are more easily produced in pairs. Among the light neutrals, only $P^{0'}$ couples to gluon pairs. The anomaly factor for this coupling is $S(P^{0'}g_a g_b) = g_s{}^2(N/\sqrt{6})\delta_{ab}$, where N refers to $\mathrm{SU}(N)_{\mathrm{TC}}$. Ignoring any mixing with P^0, we may write the differential cross section as

$$\frac{d\sigma}{dy}(ab \to P^{0'} + \text{anything}) = \frac{\pi^2 \Gamma(P^{0'} \to gg)}{8M^3} \tau f_g^{(a)}(x_a, M^2) f_g^{(b)}(x_b, M^2) , \tag{3.51}$$

where we have abbreviated $M(P^{0'})$ as M and as usual in the parton model

$$x_a = \sqrt{\tau} e^y, \quad x_b = \sqrt{\tau} e^{-y} , \tag{3.52}$$

with

$$\tau = M^2/s .$$

The parton distribution $f_i^{(a)}(x_a, Q^2)$ is the number density of partons of species i seen with momentum fraction x_a of hadron a by a probe characterized by Q^2. In numerical work we use Set 2 of the Eichten-Hinchliffe-Lane-Quigg (EHLQ) structure functions.[5]

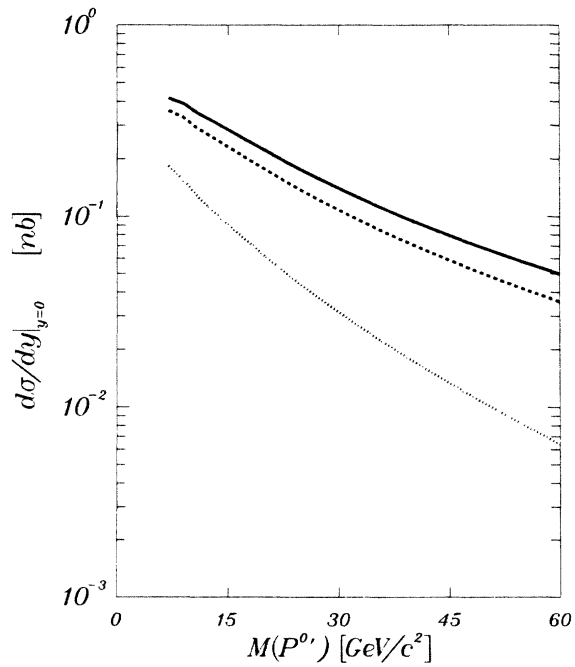

FIG. 4. Differential cross section for production of the color-singlet technipion $P^{0'}$ at $y=0$ in pp or $\bar{p}p$ collisions, for $\sqrt{s} = 2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve).

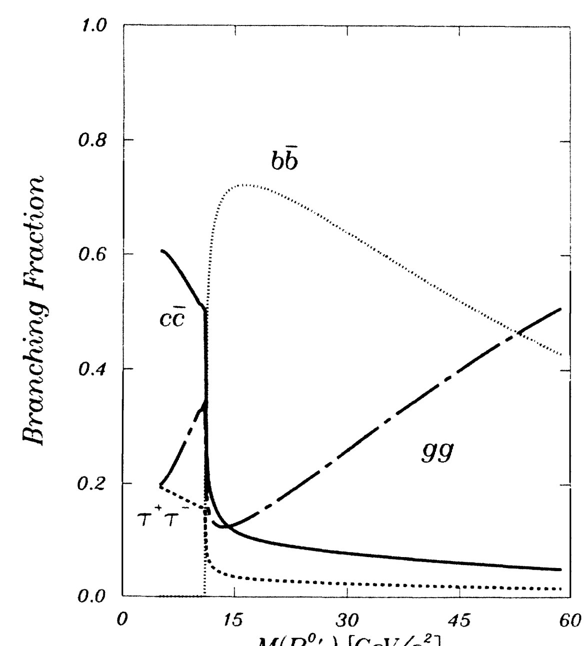

FIG. 5. Approximate branching ratios for $P^{0'}$ decay. In Eq. (3.3) we choose $N=4$ and use the running coupling constant $\alpha_s(M(P^{0'})^2)$ given by Eq. (2.42) of Ref. 5.

The different cross section for $P^{0'}$ production at $y=0$ in $p^{\pm}p$ collisions is shown as a function of the technipion mass in Fig. 4. According to (3.2) and (3.3), the principal decays will be

$$P^{0'} = \begin{cases} gg\ , \\ b\bar{b}\ , \\ \tau^+\tau^-\ , \end{cases} \tag{3.53}$$

with branching ratios indicated in Fig. 5. Comparing with the two-jet mass spectra in Figs. 93 and 94 of Ref. 5, we see that there is no hope of finding $P^{0'}$ as a narrow peak in the two-jet invariant-mass distribution. The background from $b\bar{b}$ pairs is shown in Fig. 6 (Ref. 47). It is 2—3 orders of magnitude larger than the anticipated signal, for any realistic resolution. The background to the $\tau^+\tau^-$ mode is the Drell-Yan process, for which the cross section is indicated in Fig. 7. Even when the small ($\sim 2\%$) branching ratio into τ pairs is taken into account, the signal is approximately equal to or an order of magnitude larger than the background. The signal-to-background ratio is crucially dependent upon the experimental resolution in the invariant mass of the pair. It seems questionable that τ's can be identified with high efficiency and measured with sufficient precision to make this a useful signal, but it may set an interesting target for detector development.

Because the assumption that couplings for technipion decays into fermion pairs are proportional to fermion masses is not on firm theoretical ground, it is important to keep in mind the possibility of unconventional decays and to design experiments which are sensitive to a variety of modes. For example, the $\mu^+\mu^-$ decay mode might occur with a significant branching ratio. Such a felicitous circumstance would make observing the $P^{0'}$ easy.

If the $P^{0'}$ is very light, then it may be produced at a sufficient rate in fixed-target experiments at Fermilab that a search may be carried out even with "conventional" branching ratios proportional to m_f^2. Below the threshold for decay into $b\bar{b}$ pairs, the cross section times muonic branching ratio is about 3.5 times larger for $P^{0'}$ than for the standard-model Higgs boson. In 800-GeV/c pp collisions, we estimate $B(P^{0'}\to\mu^+\mu^-)d\sigma/dy\,|_{y=0}$ of about 1.9×10^{-5} nb for a 5-GeV/c^2 technipion and about 7.5×10^{-7} nb for a 10-GeV/c^2 technipion. At 1 TeV/c, these cross sections increase to 2.3×10^{-5} nb and 1.2×10^{-6} nb.

It is appropriate here to note that the production of a neutral technipion in association with an intermediate boson is suppressed by 4—5 orders of magnitude compared with the corresponding standard-model process involving an elementary Higgs boson. This is because the technipion coupling is only through the anomaly, whereas the Higgs coupling is direct. The detection in these processes of a neutral scalar at the level anticipated in the standard model would seem to rule out the technicolor scenario. The cross sections for Higgs-boson—gauge-boson associated production are shown for reference in Figs. 8 and 9. They are probably too small to be observed with anticipated luminosities.

We now discuss the production of pairs of color-singlet

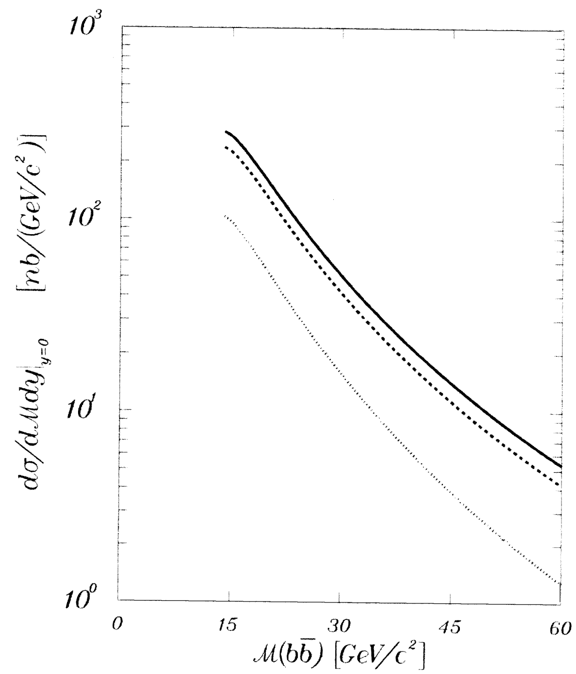

FIG. 6. Cross section $d\sigma/d\mathcal{M}\,dy\,|_{y=0}$ for the production of $b\bar{b}$ pairs in $\bar{p}p$ collisions, at $\sqrt{s}=2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve).

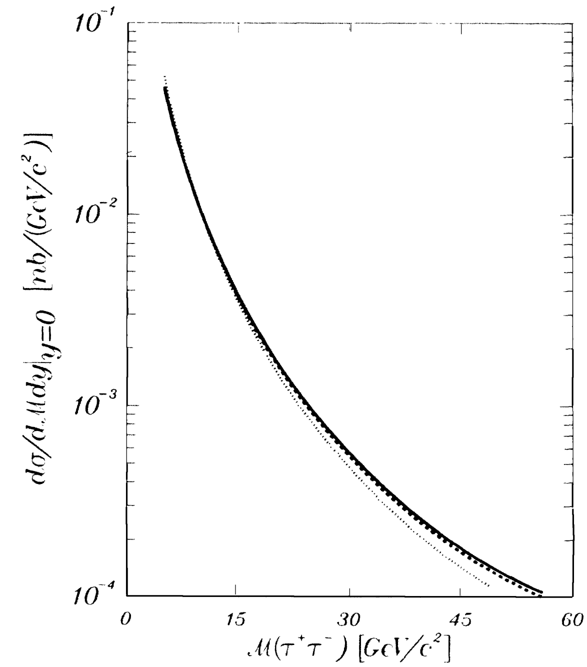

FIG. 7. Cross section $d\sigma/d\mathcal{M}\,dy\,|_{y=0}$ for the production of $\tau^+\tau^-$ pairs in proton-antiproton collisions at $\sqrt{s}=2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve). The contributions of virtual-photon and Z^0 intermediate states are included.

technipions through the chains

$$\bar{p}p \to W^{\pm} + \text{anything} \\ \hookrightarrow P^{\pm}P^{0} \tag{3.54}$$

and

$$\bar{p}p \to Z^{0} + \text{anything} \\ \hookrightarrow P^{+}P^{-}\,, \tag{3.55}$$

where the intermediate bosons may be real or virtual.[48] From Ref. 5 the branching ratios in the Farhi-Susskind model are

$$\frac{\Gamma(W^{\pm}\to P^{\pm}P^{0})}{\Gamma(W^{\pm}\to \text{all})} = \frac{\alpha M_W \left[1-\frac{[M(P^0)+M(P^{\pm})]^2}{m_W{}^2}\right]^{3/2}\left[1-\frac{[M(P^0)-M(P^{\pm})]^2}{m_W{}^2}\right]^{3/2}}{48x_W\Gamma(W^{\pm}\to\text{all})}$$

$$\approx 0.02\left[1-\frac{[M(P^0)+M(P^{\pm})]^2}{m_W{}^2}\right]^{3/2}\left[1-\frac{[M(P^0)-M(P^{\pm})]^2}{m_W{}^2}\right]^{3/2}, \tag{3.56}$$

$$\frac{\Gamma(Z^0\to P^+P^-)}{\Gamma(Z^0\to\text{all})} = \frac{\alpha m_Z(1-2x_W)^2\left[1-\frac{4M(P^{\pm})^2}{m_Z{}^2}\right]^{3/2}}{48x_W(1-x_W)\Gamma(Z^0\to\text{all})}$$

$$\approx 0.01\left[1-\frac{4M(P^{\pm})^2}{m_Z{}^2}\right]^{3/2}, \tag{3.57}$$

where $x_W=\sin^2\theta_W$. The estimate (3.57) for the technipion branching ratio of Z^0 is model independent because charged scalars transform purely as isovectors. The estimate (3.56) holds in the absence of isospin-breaking ETC interactions that induced $P^0-P^{0'}$ mixing. These results are identical to those which obtain in a two-doublet Higgs model, with the transcriptions $P^{\pm}\to H^{\pm}$, $P_i^0\to H_i^0$. In contrast, although the decay $Z^0\to H^0H^{0'}$ is allowed, with

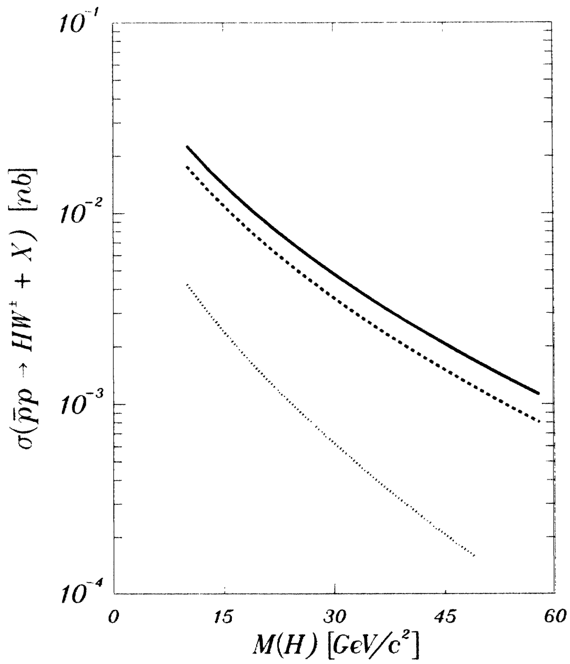

FIG. 8. Integrated cross sections for associated $HW^{\pm}$ production in $\bar{p}p$ collisions at $\sqrt{s}=2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve). The cross sections for $P^0W^{\pm}$ and $P^{0'}W^{\pm}$ production vanish in the Farhi-Susskind model.

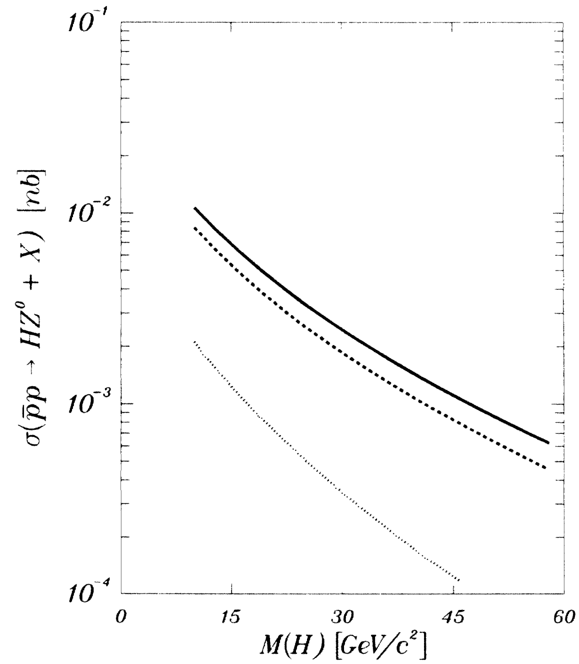

FIG. 9. Integrated cross sections for associated HZ^0 production in $\bar{p}p$ collisions at $\sqrt{s}=2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve). The cross sections for P^0Z^0 and $P^{0\prime}Z^0$ production are 4—5 orders of magnitude smaller.

$$\frac{\Gamma(Z^0 \to H^0 H^{0'})}{\Gamma(Z^0 \to \text{all})} = \frac{\alpha M_Z \left[1 - \frac{(M_0 + M_0')^2}{M_Z{}^2}\right]^{3/2} \left[1 - \frac{(M_0 - M_0')^2}{M_Z{}^2}\right]^{3/2}}{48 x_W (1 - x_W) \Gamma(Z^0 \to \text{all})} , \tag{3.58}$$

the decay of Z^0 into pairs of neutral technipions does not occur.[49]

The elementary process corresponding to reaction (3.54) is

$$q_i \bar{q}_j \to W^{\pm} \to P^+ P^0 . \tag{3.59}$$

The differential cross section is

$$\frac{d\hat{\sigma}_{ij}}{d\hat{t}} = \frac{\pi \alpha^2 |K_{ij}|^2}{12 s^2 x_W{}^2} \frac{[\hat{u}\hat{t} - M(P^+)^2 M(P^0)^2]}{(\hat{s} - m_W{}^2)^2 + m_W{}^2 \Gamma_W{}^2} , \tag{3.60}$$

where $\hat{s}$, $\hat{t}$, and $\hat{u}$ are the Mandelstam variables for the parton-parton collision, and K_{ij} is an element of the Kobayashi-Maskawa quark mixing matrix. The total cross section is

$$\hat{\sigma}_{ij} = \frac{\pi \alpha^2 \mathscr{S}^3 |K_{ij}|^2}{72 \hat{s}^2 x_W{}^2 [(\hat{s} - m_W{}^2)^2 + m_W{}^2 \Gamma_W{}^2]} , \tag{3.61}$$

with

$$\mathscr{S} = \{\hat{s} - [M(P^+) + M(P^0)]^2\}^{1/2} \{\hat{s} - [M(P^+) - M(P^0)]^2\}^{1/2} . \tag{3.62}$$

In the same way, the differential cross section for the elementary process

$$q_i \bar{q}_i \to \gamma \text{ or } Z^0 \to P^+ P^- \tag{3.63}$$

is

$$\frac{d\hat{\sigma}_i}{d\hat{t}} = \frac{4\pi\alpha^2}{3\hat{s}^2} \left[\frac{e_i{}^2}{\hat{s}^2} + \frac{e_i (L_i + R_i)(1 - 2x_W)(\hat{s} - m_Z{}^2)}{8 x_W (1 - x_W) \hat{s} [(\hat{s} - m_Z{}^2)^2 + m_Z{}^2 \Gamma_Z{}^2]} + \frac{(L_i{}^2 + R_i{}^2)(1 - 2x_W)^2}{64 x_W{}^2 (1 - x_W)^2 [(\hat{s} - m_Z{}^2)^2 + m_Z{}^2 \Gamma_Z{}^2]} \right]$$
$$\times [\hat{u}\hat{t} - M(P^+)^4] , \tag{3.64}$$

where e_i is the quark charge in units of the proton charge, $L_i = (\tau_3 = 2 e_i x_W)$ and $R_i = -2 e_i x_W$. The total cross section is

$$\hat{\sigma}_i = \frac{2\pi\alpha^2 \hat{s}}{9} [1 - 4M(P^+)^2/\hat{s}]^{3/2} \left[\frac{e_i{}^2}{\hat{s}^2} + \frac{e_i (L_i - R_i)(1 - 2x_W)(\hat{s} - m_Z{}^2)}{8 x_W (1 - x_W) \hat{s} [(\hat{s} - m_Z{}^2)^2 + m_Z{}^2 \Gamma_Z{}^2]} + \frac{(L_i{}^2 + R_i{}^2)(1 - 2x_W)^2}{64 x_W{}^2 (1 - x_W)^2 [(\hat{s} - m_Z^2)^2 + m_Z{}^2 \Gamma_Z^2]} \right] . \tag{3.65}$$

We show in Fig. 10 the cross section for $P^{\pm} P^0$ pairs produced in $\bar{p}p$ collisions. Both $P^{\pm}$ and P^0 are required to have rapidities $|y_i| < 1.5$. The cross section is appreciable only for $[M(P^{\pm}) = M(P^0)] < m_W/2$, for which the rate is determined by the decays of real $W^{\pm}$. In this regime, an experiment with sensitivity 10^{37} cm^{-2} will record a potential signal of $\sim 10^3$ events, to be reconstructed in channels such as $t\bar{b}b\bar{b}$, $c\bar{b}b\bar{b}$, etc. (Compare Table I.)

In the same way, the cross section for production of $P^+ P^-$ pairs is only of experimental interest at currently attainable energies if $M(P^{\pm}) < m_Z/2$. This as shown in Fig. 11. Again, cross sections approaching 0.03 nb are predicted, populating channels such as $t\bar{b}\bar{t}b$ and $c\bar{b}\bar{c}b$.

As we have remarked in Sec. II, the couplings of $P^{\pm}$ and P^0 to fermion pairs can be complicated by (ETC) model-dependent mixing matrices. For this reason, the search for scalar particles from $W^{\pm}$ and Z^0 decay must be as broad and thorough as practicable. As an extreme example, we may conceive of models in which the lightest technipions couple predominantly to the lightest generation of quarks or leptons, instead of having couplings proportional to fermion mass.[50]

The $W^+ \to P^+ P^0$ decay mode results in four hadronic jets when the technipions decay into light quark-antiquark pairs. The major background to these four-jet events comes from ordinary QCD processes. Before cuts are imposed, this background may be 2 orders of magnitude larger than the anticipated signal. Significant couplings of P^0 to $e^+ e^-$ or $\mu^+ \mu^-$ pairs would also improve the prospects for detection. Whether the background can be adequately suppressed in the generic four-jet mode to allow detection of the technipions in this case requires further study.

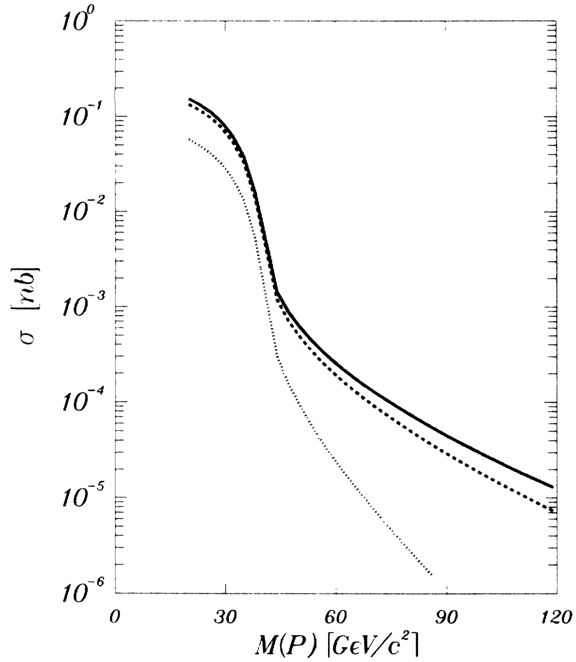

FIG. 10. Cross section for the production of P^+P^0 and P^-P^0 (summed) in $\bar{p}p$ collisions as a function of the common (by assumption) mass of the technipions, for $\sqrt{s}=2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve). Both $P^{\pm}$, P^0 must satisfy $|y|<1.5$.

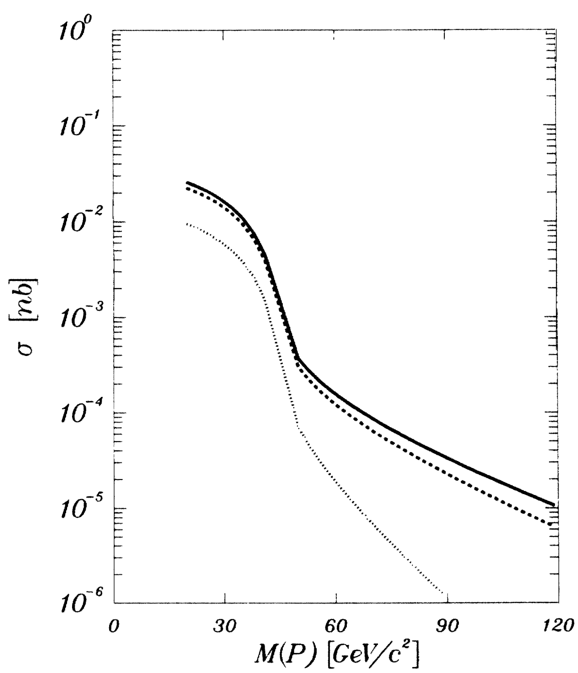

FIG. 11. Cross section for the production of P^+P^- pairs in $\bar{p}p$ collisions as a function of the mass of P^+, for $\sqrt{s}=2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve).

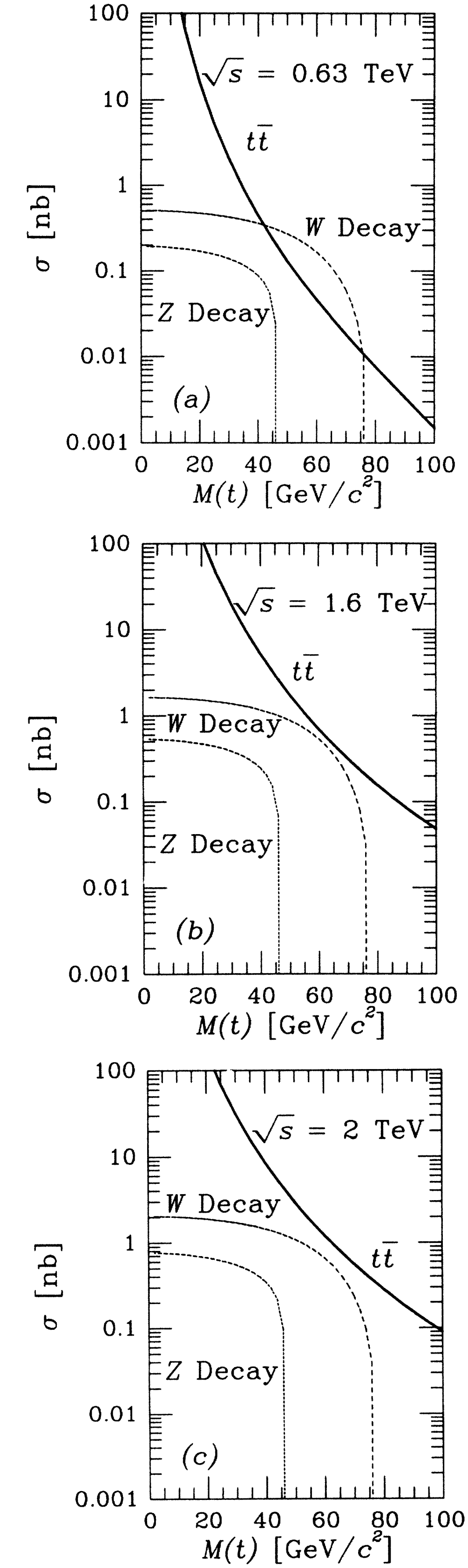

FIG. 12. Cross sections for the production of t or $\bar{t}$ quarks in $\bar{p}p$ collisions as a function of the mass of the heavy quark. (a) $\sqrt{s}=630$ GeV; (b) $\sqrt{s}=1600$ GeV; (c) $\sqrt{s}=2000$ GeV.

The rate of production of P^+P^- pairs is so low that it is rather unlikely that this channel could be detected at a hadron collider in the near future. The situation is somewhat better at "Z^0 factories." For an integrated luminosity of 10^{38} cm^{-2} we can expect approximately 4×10^6 Z^0's, which will yield 4×10^4 P^+P^- pairs. The signal will be four jets containing heavy flavors as in the hadronic case, with a background from four-jet events. The experiments are likely to be at least as challenging as the searches at PEP and PETRA energies discussed in Sec. III B 2.

The final source of light technipions to be discussed is the semiweak decays of heavy quarks, which we have already mentioned in Sec. III B 1. All that remains to be investigated is the number of events to be expected in $\bar{p}p$ collisions. We show in Fig. 12 the cross sections for the production of heavy quarks in the reactions

$$\bar{p}p \to W^{\pm} + \text{anything} \quad \hookrightarrow t\bar{b} \ , \tag{3.66}$$

$$\bar{p}p \to Z^0 + \text{anything} \quad \hookrightarrow t\bar{t} \ , \tag{3.67}$$

and

$$\bar{p}p \to t\bar{t} + \text{anything} \ , \tag{3.68}$$

for which we consider only the $gg\to t\bar{t}$ contribution. Roughly speaking, for $m_t \approx 40$ GeV/c^2, we expect about 1 nb of single-t production from W decay and 1 nb of $t\bar{t}$ pair production from the gluon-fusion process. An experiment with a sensitivity of 10^{37} cm^{-2} may therefore expect approximately 10^4 events of each type. It is then overwhelmingly likely that each t quark will yield a charged technipion, provided the decay

$$t \to P^+ b \tag{3.69}$$

is kinematically allowed. If the subsequent decay of the technipion is into the $c\bar{b}$ channel, identifying the P^+ requires reconstructing the $c\bar{b}$ or $\bar{c}b$ invariant mass in $\bar{b}c\bar{b}b$ and $c\bar{b}\,\bar{b}b\bar{c}b$ final states.

We have seen that the rates for single $P^{0'}$ pair production and for $P^{\pm}P^0$ pair production are sufficiently large that we may look forward to meaningful searches at the Tevatron collider. The rate for production of P^+P^- pairs occurs at a smaller rate, so that incisive searches may only be possible on the Z^0 peak at LEP or the SLC.

IV. COLORED TECHNIPIONS

A. Theoretical expectations

For definiteness, we continue to discuss the consequences of the Farhi-Susskind model[17] with a pair (U,D) of color-triplet techniquarks and a pair (N,E) of color-singlet technileptons. From these constituents we may build, in addition to the color-singlet states discussed in Sec. III, the following ${}^1S_0(F\bar{F})$ states: an isospin triplet ${P_3}^1$, ${P_3}^0$, ${P_3}^{-1}$ of color-triplet technipions; an isospin-singlet color-triplet state P_3'; the corresponding color-antitriplet states; an isospin triplet P_8^+, P_8^0,P_8^- of color-octet technipions; an isoscalar color-octet state $P_8^{0'}$, also known as η_T.

Corresponding to each of these pseudoscalars is a hyperfine partner ${}^3S_1$ techni—vector meson. The technifermion wave functions of all these states are given in Table II of Ref. 5.

Like the color-singlet technipions, the colored technipions receive electroweak and ETC contributions to their masses, but these are much smaller than the expected QCD contribution[51]

$$M(P_3) = 160 \text{ GeV}/c^2 \left[\frac{4}{N}\frac{n_f}{8}\right]^{1/2} , \tag{4.1}$$

$$M(P_8) = 240 \text{ GeV}/c^2 \left[\frac{4}{N}\frac{n_f}{8}\right]^{1/2} , \tag{4.2}$$

where N refers to the technicolor gauge group $SU(N)_{TC}$ and n_f is the number of "flavors" of technifermions assigned to the fundamental **N** representation. For the case at hand, $n_f = 2\times 3 + 2 = 8$. The estimates (4.1) and (4.2) are of course specific to the color-SU(3) representations, to the choice of an $SU(N)_{TC}$ technicolor group, and to the techniflavor symmetry group

$$G_f = SU(n_f)_L \otimes SU(n_f)_R \otimes U(1)_V \ . \tag{4.3}$$

To discuss the decays of the colored technifermions into ordinary quarks and leptons, we must specify the weak hypercharge assignments of the technifermions. If under $SU(3)_c \times SU(2)_L \otimes U(1)_Y$ the technifermions transform as

$$\begin{aligned} Q_L = (U,D)_L : (3,2,Y), \quad & U_R : (3,1,Y+1) \ , \\ & D_R : (3,1,Y-1) \ , \\ L_L = (N,E)_L : (1,2,-3Y), \quad & N_R : (1,1-3Y+1) \ , \\ & E_R : (1,1-3Y-1) \ , \end{aligned} \tag{4.4}$$

the absence of anomalies is guaranteed for all gauge currents. For the choice $Y=\frac{1}{3}$, the techniquark and technilepton charges are those of ordinary quarks and leptons.

The electric charges of the color-triplet isovector technipions are then $2Y+1, 2Y, 2Y-1$, while the isoscalar state carries charge $2Y$. If the weak hypercharge Y satisfies

$$2Y = \tfrac{2}{3} + \text{integer} \ , \tag{4.5}$$

as it will for the canonical choice $Y=\frac{1}{3}$, then the color-triplet technipions will decay as

$$P_3 \to q\bar{l} + \cdots \quad \text{or } \bar{q}\,\bar{q} + \cdots \ . \tag{4.6}$$

If condition (4.5) is not met, the lightest P_3's will be absolutely stable. The decay rates for P_3 and P_8 into fermion-antifermion pairs may be estimated using (3.2). The color-octet states may also decay into a pair of gauge bosons, with widths given by (3.50). Of these decays the most prominent is that of $P_8^{0'}$ into two gluons, for which the decay rate is

TABLE II. Principal decay modes of colored technipions if $Pf_1\bar{f}_2$ couplings are proportional to fermion mass.

Technipion	Principal decay modes
P_3, P_3'	$t\tau^+, t\bar{\nu}_\tau, b\tau^+, \ldots$ or $\bar{t}\bar{b}, \ldots$
P_8^+	$(t\bar{b})_8$
P_8^0	$(t\bar{t})_8$
$P_8^{0'}$	$(t\bar{t})_8; gg$

$$\Gamma(P_8^{0'} \to gg) \approx \frac{5\alpha_s^2}{24\pi^3} \frac{M(P_8^{0'})^3}{F_\pi^2} \left[\frac{N}{4}\right]^2 . \quad (4.7)$$

What we may expect to be the principal decay modes of the colored technipions are listed in Table II.

In O(N)-based technicolor models there is a richer spectrum of possible technicolor-singlet pseudoscalar states. These include ditechniquark P_6 and P_{3^*} (QQ) color sextet and antitriplet states, $P_3'(QL)$ color triplets, and color-singlet dilepton-number $P_L(LL)$ states, with masses of approximately[14] $M(P_6) \approx 270$ GeV/c^2, $M(P_{3^*}) \approx M(P_3') \approx 160$ GeV/c^2, and $M(P_L) \approx 70$ GeV/c^2.

B. Prospects for detection in $\bar{p}p$ collisions

Because the masses of the colored technipions are expected to be in the range of a few hundred GeV/c^2, production and detection of these states is likely to be quite challenging for the current generation of hadron colliders. The principal sources of the colored technipions in $\bar{p}p$ collisions are (1) the production of the weak-isospin-singlet state $P_8^{0'}$ by the gluon-fusion mechanism and (2) the production of $(P_3\bar{P}_3)^0$ or $(P_8P_8)^0$ pairs in $q\bar{q}$ or gg fusion. We shall consider these in turn, always returning to the one-generation Farhi-Susskind model for numerical examples.

The differential cross section for $P_8^{0'}$ production may be calculated using (3.51) with the partial width $\Gamma(P' \to gg)$ replaced by (4.7). The differential cross section (summed over the eight color indices) at $y=0$ is shown as a function of the technipion mass in Fig. 13. The dominant decay modes are commonly expected to be

$$P_8^{0'} \to \begin{cases} gg \, , \\ t\bar{t} \, . \end{cases} \quad (4.8)$$

The expected branching ratios depend upon the top-quark mass. Representative estimates are shown in Fig. 14. The background expected from $t\bar{t}$ production by conventional mechanisms is plotted in Fig. 15, for top quark masses of 30, 40, 50, 60, and 70 GeV/c^2. When the branching ratios are taken into account, the signal and background are roughly comparable in bins appropriate to experimental resolution, but the expected number of events is small. At $\sqrt{s}=2$ TeV, and for $M(P_8)=240$ GeV/c^2, an experiment with a sensitivity of 10^{36} cm^{-2} can expect a signal of a few $t\bar{t}$ pairs. At $\sqrt{s}=630$ GeV, the signal is nonexistent.

In the two-gluon channel, the signal will be comparable

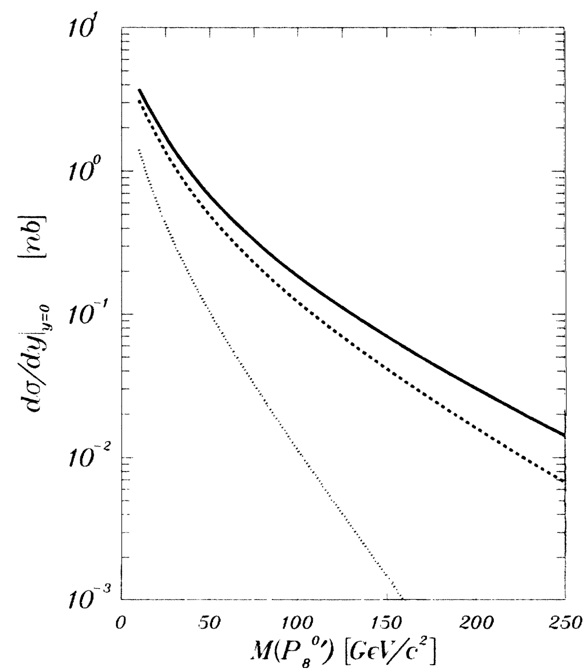

FIG. 13. Differential cross section for the production of color-octet technipions $P_8^{0'}$ at $y=0$ in $p^{\pm}p$ collisions, for $\sqrt{s}=2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve). The expected mass, according to (4.2), is approximately 240 GeV/c^2.

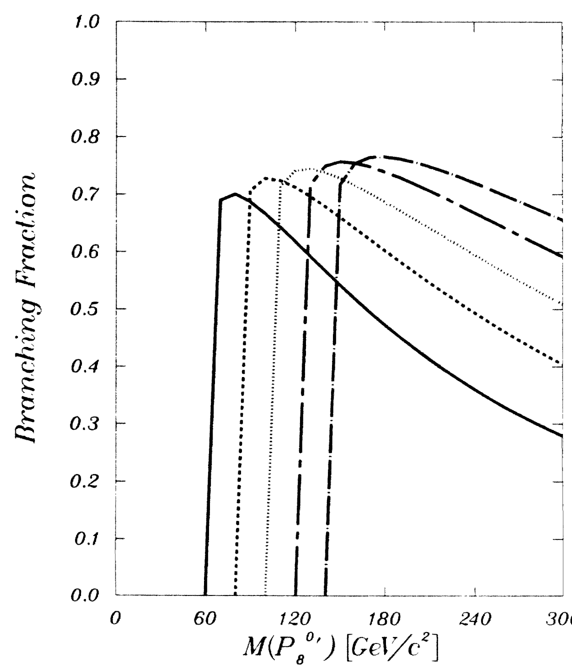

FIG. 14. Branching fractions for $P_8^{0'} \to t\bar{t}$. The remaining decays are into the two-gluon channel. In Eq. (3.3) we choose $N=4$ and use the running coupling constant $\alpha_s(M(P_8^{0'})^2)$ given by Eq. (2.42) of Ref. 5. The top-quark mass is 30 (solid), 40 (dashed), 50 (dotted), 60 (chain-dashed), or 70 (chain-dotted) GeV/c^2.

to the $t\bar{t}$ signal or somewhat larger, depending upon the mass of the top quark. The expected background, which may be judged from Figs. 93 and 94 of Ref. 5, is very large compared to the signal.

The γg decay of $P_8^{0'}$ has been suggested[52] as a low-background channel in which to reconstruct the color-octet technipion. The branching ratio for γg decay may be computed from (3.50) as

$$\frac{\Gamma(P_8^{0'}\to g\gamma)}{\Gamma(P_8^{0'}\to gg)}=\left[\frac{\alpha_{\rm em}}{\alpha_s}\right]\frac{1}{30}\ . \tag{4.9}$$

For the cross sections presented in Fig. 13, this leads to a negligible event rate at the S$\bar{p}p$S and Tevatron colliders. If the mass of $P_8^{0'}$ is lower than our estimate, the prospects for detection improve.

The elementary processes for pair production of colored technipions are depicted in Fig. 16. The differential cross sections for neutral channels are

$$\frac{d\sigma}{d\hat{t}}(q\bar{q}\to PP)=\frac{2\pi\alpha_s{}^2}{9\hat{s}^2}T(R)\beta^2|X|^2(1-z^2) \tag{4.10}$$

and

$$\frac{d\sigma}{d\hat{t}}(gg\to PP)$$
$$=\frac{2\pi\alpha_s{}^2T(R)}{\hat{s}^2}\left[\left[\frac{T(R)}{d(R)}-\frac{3}{32}\right](1-2V+2V^2)+\frac{3}{32}\beta^2z^2(|X|^2-2V\,\mathrm{Re}X+2V^2)\right]\ , \tag{4.11}$$

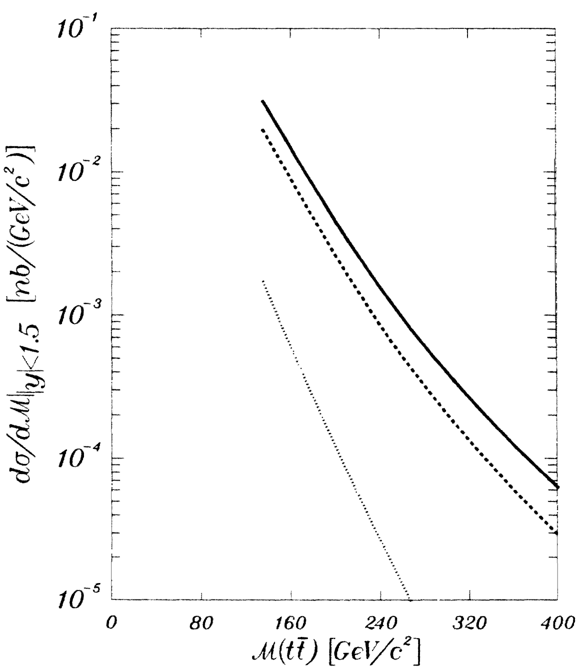

FIG. 15. Mass spectrum of $t\bar{t}$ pairs produced in proton-antiproton collisions at $\sqrt{s}=2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve). The rapidity of each produced quark is constrained to satisfy $|y_t|<1.5$.

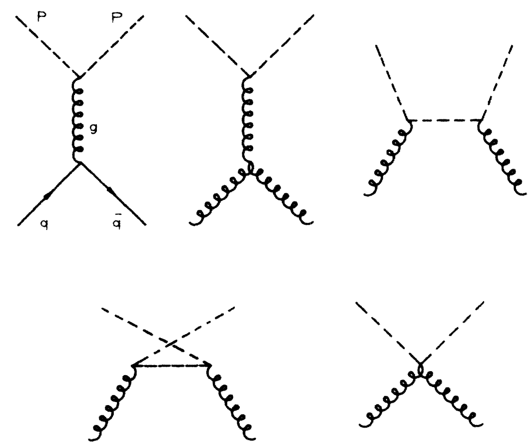

FIG. 16. Feynman graphs for the production of pairs of colored technipions. The curly lines are gluons, solid lines are quarks, and dashed lines are technipions. The graphs with s-channel gluons include the $\rho_8^{0'}$ enhancement.

where $z=\cos\theta^*$ measures the c.m. scattering angle,

$$\beta^2=1-4M(P)^2/\hat{s} \tag{4.12}$$

and

$$V=1-\frac{1-\beta^2}{1-\beta^2z}\ . \tag{4.13}$$

The existence of a color-octet, isoscalar techni—vector meson $\rho_8^{0'}$ with mass

$$M(\rho_8^{0'})\approx M(\rho)\frac{F_\pi}{f_\pi}\left[\frac{3}{N}\right]^{1/2}\approx(885\ \mathrm{GeV}/c^2)\left[\frac{4}{N}\right]^{1/2}\left[\frac{8}{n_f}\right]^{1/2} \tag{4.14}$$

(with $f_\pi=93$ GeV) gives rise (by techni—vector meson dominance) to an enhancement of the amplitudes arising from gluon exchange in the s channel. We represent this enhancement by

$$X=\frac{M(\rho_8^{0'})^2}{M(\rho_8^{0'})^2-\hat{s}-iM(\rho_8^{0'})\Gamma(\hat{s})}\ , \tag{4.15}$$

where the energy-dependent width of $\rho_8^{0'}$ is

$$M(\rho_8^{0'})\Gamma(\hat{s})=\frac{g_{\rho T}{}^2\hat{s}}{48\pi}[\beta_3{}^3\theta(\beta_3)+3\beta_8{}^3\theta(\beta_8)]\ , \tag{4.16}$$

and the techni-ρ coupling constant $g_{\rho T}$ is related to the $\rho\to\pi\pi$ coupling constant $g_\rho(g_\rho{}^2/4\pi=2.98)$ by

$$g_{\rho T}{}^2=g_\rho{}^2(3/N)\ . \tag{4.17}$$

The color factors are

$$T(R)=\begin{cases}\frac{1}{2}\ \text{for}\ P_3\ ,\\ 3\ \text{for}\ P_8\ ,\end{cases} \tag{4.18}$$

$$d(R)=\begin{cases}3 & \text{for } P_3 \, ,\\ 8 & \text{for } P_8 \, .\end{cases} \tag{4.19}$$

In writing (4.10) and (4.11) we have summed over all charges and colors. The individual charge states

$$P_3{}^1\bar{P}_3{}^1, \quad P_3{}^{-1}\bar{P}_3{}^{-1}, \quad P_3{}^0\bar{P}_3{}^0, \quad P_3'\bar{P}_3' \, , \tag{4.20}$$

or

$$P_8^+P_8^- \, , \quad P_8^0P_8^0+P_8^{0'}P_8^{0'} \tag{4.21}$$

occur with equal cross sections.

The integrated cross section for the reaction

$$\bar{p}p \to P_3\bar{P}_3 + \text{anything} \, , \tag{4.22}$$

summed over the charge states (4.20), is plotted as a function of $M(P_3)$ in Fig. 17, without the $\rho_8^{0'}$ enhancement. For the purposes of this calculation we adopted the canonical value (4.14) $M(\rho_8^{0'})=885$ GeV/c^2 of the techni-ρ mass, and evaluated the mass-dependent techni-ρ width using (4.16) with $M(P_8)$ fixed at its nominal value of 240 GeV/c^2 as given by (4.2). With these parameters, and at the energies and technipion masses under consideration, neither the techni-ρ enhancement nor the restriction to central rapidities $|y|<1.5$ is particularly effective. The cross sections are small. However, in variants of the model, the choice of a large technicolor group $SU(N)_{TC}$ implies through (4.1) and (4.2) smaller P_3 masses, for which the production rates become correspondingly larger. It is therefore worth making a few remarks about signatures.

If the technipions are stable, which will be the case if (4.5) is not satisfied, the signatures should be quite striking and essentially background-free. Each event will appear as a pair of extremely narrow jets consisting of the very massive P_3 core (plus a quark or antiquark to neutralize its color), together with relatively soft $q\bar{q}$ pairs and gluons. The decay of unstable technipions into $q+\bar{l}+\cdots$ should also provide a characteristic signature: a jet and an isolated lepton on each side of the beam. In this case the only comparable conventional background would be from the pair production of heavy quarks, with the subsequent decay

$$Q \to qW \atop \qquad \hookrightarrow l\nu \, . \tag{4.23}$$

For such events one expects equal numbers of electrons, muons, and τ's. In contrast, the technipion decays are expected to favor τ's.

We turn next to the pair production of octet technipions. The integrated cross sections for the reaction

$$\bar{p}p \to P_8\bar{P}_8 + \text{anything} \tag{4.24}$$

with $|y|<1.5$ are plotted in Fig. 18 without the $\rho_8^{0'}$ enhancement. These are typically a few times the cross sections for color-triplet technipion production at the same mass, because of the larger color factors in (4.11), and slightly smaller than the cross sections for single-$P_8^{0'}$ production. In this case, we have computed the mass-dependent $\rho_8^{0'}$ width using (4.16) with $M(P_3)$ fixed at its nominal value (4.1) of 160 GeV/c^2. The techni-ρ enhancement is less effective in the octet technipion chan-

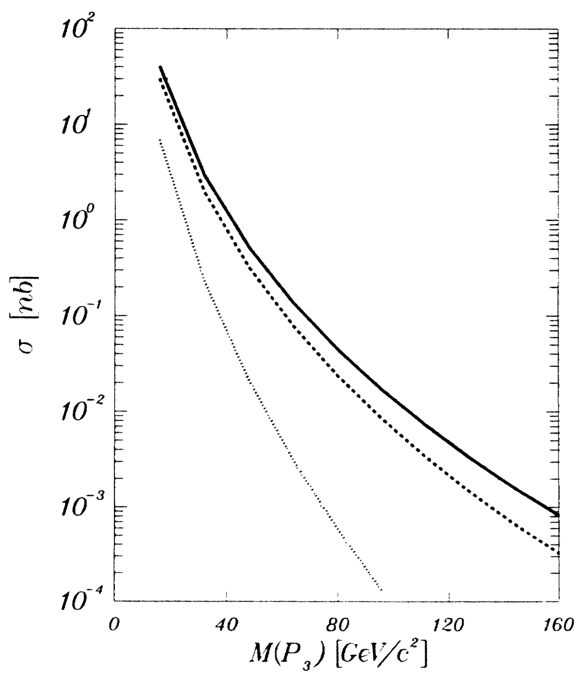

FIG. 17. Cross sections for the production of $P_3\bar{P}_3$ pairs in $\bar{p}p$ collisions at $\sqrt{s}=2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve). The canonical value of the technipion mass is $M(P_3)=160$ GeV/c^2.

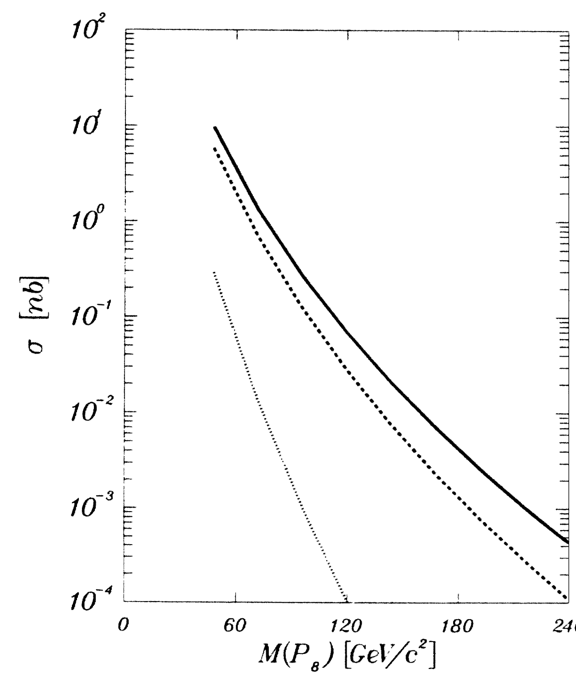

FIG. 18. Cross sections for the production of P_8P_8 pairs in $\bar{p}p$ collisions at $\sqrt{s}=2000$ GeV (solid curve), 1600 GeV (dashed curve), and 630 GeV (dotted curve). The canonical value of the technipion mass is $M(P_8)=240$ GeV/c^2.

nel because of the large color factor in the first term in (4.11). For the canonical values of the technipion masses, the cross sections are too small to be of interest.

The expected decays of octet technipions are

$$P_8^+ \to t\bar{b}, \quad P_8^0 \to t\bar{t} \tag{4.25}$$

and

$$P_8^{0'} \to \begin{cases} t\bar{t} \,, \\ gg \,, \end{cases} \tag{4.26}$$

with branching fractions given earlier in Fig. 14. The signature for the $P_8^+ P_8^-$ channel is therefore $t\bar{b}$ on one side of the beam and $\bar{t}b$ on the other. If the heavy flavors can be tagged with high efficiency, we know of no significant conventional backgrounds. Similar remarks apply for the neutral octet technipions.

V. CONCLUDING REMARKS

Gaining an understanding of the scalar sector of the electroweak interactions is one of the great challenges of elementary-particle physics. Although the complete elaboration of electroweak symmetry breaking must await colliders which can attain subprocess energies of 1 TeV or more, there is much to be learned at the present generation of $\bar{p}p$ colliders at CERN and Fermilab, as well as at the Z^0 factories nearing completion at CERN and SLAC. While we have concentrated on signals for technicolor, many of the channels discussed are generally of interest in models that go beyond the standard electroweak theory.

The spectrum and properties of low-lying particles in technicolor theories were discussed in Sec. II. All of these technipion states are spin-zero particles which may be classified by the mechanisms by which they acquire mass. The lightest states are color-singlets ($P^\pm, P^0, P^{0'}$) which acquire masses ($\lesssim 40$ GeV/c^2) only through extended technicolor interactions and electroweak interactions. The colored states P_3 and P_8 have masses in the range of 100–300 GeV/c^2 generated by QCD and electroweak interactions. The explicit spectrum of states and their properties were detailed in the Farhi-Susskind technicolor model. Light charged and neutral scalars may also occur in multiple-Higgs-doublet models, and in the $N=1$ supersymmetric generalization of the standard electroweak theory.

The color-singlet technipions were discussed in Sec. III. Experimental constraints on the masses and couplings of the lightest technipions may be summarized as follows. For the neutral technipions ($P^0, P^{0'}$), there are no direct limits, but the absence of rare (decay) processes imposes limits on flavor-changing couplings to fermion pairs. The only limits on $P^\pm$ masses and couplings come from searches in e^+e^- collisions. The ensuing mass restrictions are summarized in Fig. 2 for decay modes other than $P^+ \to c\bar{b}$, which is completely unconstrained by existing measurements.

In $\bar{p}p$ collisions, the principal sources of light technipions are single-$P^{0'}$ production in gluon-gluon collisions; $P^\pm P^0$ pair production via real $W^\pm$ bosons; and the production of $P^\pm$ in heavy-quark decays. If $P^{0'}$ decay amplitudes are proportional to fermion masses, the large hadronic two-jet background precludes detection of $P^{0'}$ in any but the $\tau^+\tau^-$ channel. Fixed-target experiments may be able to carry out interesting searches in the $\mu^+\mu^-$ channel for $P^{0'}$ masses up to about 10 GeV/c^2. At the energies of interest in this paper, pair production in $W^\pm$ decays is generally a more promising source of technipions, because of the large W cross section and the favorable (up to 2%) branching ratio. Similarly, the $O(1\%)$ branching ration for $Z^0 \to P^+P^-$ decays makes Z^0 factories an advantageous setting for $P^\pm$ searches. Finally, charged technipions should be emitted in the decays of top quarks provided that the semiweak decay $t \to P^+ b$ is energetically allowed. Unless the coupling of P^+ to top is unexpectedly suppressed, this would be the dominant decay of the top quark, and should be relatively easy to infer.

The colored technipions were treated in Sec. IV. Within the Farhi-Susskind model, their masses are so large that production rates will be small even at the Tevatron collider. While the favored decay modes lead to distinctive signatures, good fortune will be required to see these states. If the masses of P_3 and P_8 are smaller than anticipated in the Farhi-Susskind model, or if the production cross sections are enhanced, prospects are considerably brightened. Given the indefiniteness of ETC models, searches should not be unduly restricted by theoretical expectations.

In closing we emphasize that the occurrence of charged and neutral spinless particles with masses of less than $O(100$ GeV/$c^2)$ is common in theoretical attempts to go beyond the standard electroweak model. In multiple-Higgs-doublet generalizations of the standard model, or in supersymmetric $SU(2)_L \otimes U(1)_Y$, the properties expected for the light scalars are more or less those we have described for the light technipions in this article. Existing searches for these particles are decidely incomplete. Systematic studies are of crucial importance both for the narrow purpose of confirming or ruling out the technicolor approach, and also as part of the general search for clues to the nature of electroweak symmetry breaking.

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