Possible origin of a natural conservation of flavor in an interaction with neutral Higgs bosons

N. G. Ural'tsev

B. P. Konstantinov Institute of Nuclear Physics, Academy of Sciences of the USSR

(Submitted 22 January 1983)

Pis'ma Zh. Eksp. Teor. Fiz. **37**, No. 6, 281–284 (20 March 1983)

In technicolor models the masses of the neutral pseudo-Goldstone bosons which interact with quarks and leptons without flavor conservation automatically acquire an order of magnitude $M \sim (m_q \Lambda_{\rm TC})^{1/2} \sim 0.2$–1 TeV through the Yukawa interaction. As a result, an effective Lagrangian which conserves only light Higgs bosons satisfies the condition of natural flavor conservation.

PACS numbers: 12.35.Cn, 14.80.Gt

The minimal $SU(2)_L \times U(1)$ theory of electroweak interactions with a single Higgs doublet contains a single scalar boson σ, whose interaction with fermions does not

 0021-3640/83/060333-05$01.00

change their flavors. With increasing number of the Higgs fields interacting with the fermions, neutral scalar particles will generally arise; the exchange of these particles induces neutral currents with flavor conservation. The condition for the "natural" absence of such processes is known to be the requirement that the fermions with a given electric charge yield the mass of only one Higgs field.[1] In attempts to obtain a realistic mass matrix of ordinary fermions in the various technicolor models it is frequently found that the technicolor analog of this condition does not hold; i.e., the current mass of the fermions with a single electric charge arises by virtue of the vacuum condensate of several techniquarks. In this case the nonconservation of flavor in the neutral currents is particularly dangerous, since it may be induced by pseudo-Goldstone bosons with a smaller mass.

In this letter we show that when the radiative effects associated with the "expanded" technicolor interaction are taken into account in theories of the technicolor type all the light neutral bosons, which violate flavor conservation, and the corresponding charged pseudo-Goldstone bosons generally acquire a large mass. (The exchange of extended-technicolor bosons at low energies simulates the ordinary Yukawa interaction.) As a result, the effective low-energy Lagrangian of the theory contains no more than two doublet Higgs field, which give masses $Q=-1/3$ and $Q=2/3$ to the quarks and leptons, respectively; in addition, there is an arbitrary number of doublets which do not interact with quarks and leptons.

To calculate the masses of the pseudo-Goldstone bosons, we use the effective-Lagrangian method as in Ref. 2. The validity of this approach will be confirmed by means of current algebra.

We consider the simple case in which the Higgs doublets φ_1 and φ_2 give the quarks masses $Q=-1/3$, while the other doublets, φ_i, $i=3,...,n$, do not interact with the fermions:

$$-\mathcal{L}_Y = \sum_{i,j} h_1^{ij}\,[\bar{d}_R^i d_L^j \varphi_1^{(0)*} + \bar{d}_R^i u_L^j \varphi_1^{(+)*}] + \sum_{i,j} h_2^{ij}\,[\bar{d}_R^i d_L^j \varphi_2^{(0)*} + \bar{d}_R^i u_L^j \varphi_2^{(+)*}]. \tag{1}$$

Here i and j are the fermion generation indices.

The pseudo-Goldstone bosons arise because in the skeletal approximation the intrinsic Higgs potential $V(\varphi_1,...,\varphi_n)$ is invariant with respect to independent $SU(2)\times U(1)$ rotations of any doublets.[2] During the development of the (real) vacuum expectation values v_i,

$$\varphi_k = \begin{pmatrix} \varphi_k^{(+)} \\ \varphi_k^{(0)} \end{pmatrix}, \quad \langle \varphi_k^{(0)} \rangle = \frac{1}{\sqrt{2}} v_k; \qquad \varphi_k^{(0)} = \frac{1}{\sqrt{2}}(v_k + \sigma_k + iH_k), \tag{2}$$

all the charged states $\varphi_i^{(+)}$ and also the neutral states H_i are pseudo-Goldstone entities. The combinations

$$g^{(\pm)} = {}^1/_v \sum_i v_i \varphi_i^{(\pm)}, \qquad g^{(0)} = {}^1/_v \sum_i v_i H_i^{(0)}, \qquad v^2 = \sum_i v_i^2 \tag{3}$$

are true Goldstone bosons, which make $W^\pm$ and Z heavier.

Interaction (1) contributes to the single-loop Coleman-Weinberg potential[3]:

$$\delta V = -\frac{1}{2}\int \frac{d^4k}{(2\pi)^4 i}\,\mathrm{Sp}\ln\left(1 - \frac{mm^+(\varphi)}{k^2}\right) \simeq -\frac{N_c}{8\pi^2}\Lambda^2\,\{\mathrm{Tr}(h_1^+h_1)(\varphi_1^+\varphi_1)$$

$$+\ \mathrm{Tr}(h_2^+h_2)(\varphi_2^+\varphi_2) + \mathrm{Tr}(h_2^+h_1)(\varphi_1^+\varphi_2) + \mathrm{Tr}(h_1^+h_2)(\varphi_2^+\varphi_1)\}\,. \qquad (4)$$

We restrict the discussion to terms of second order in the Yakawa constants; Λ is the scale value of the ultraviolet cutoff. The trace in the last term represents a sum over only the quark generations.

Potential (4), which violates the invariance with respect to independent unitary rotations of φ_1 and φ_2, leads to an identical mass M^2 for the states[2]:

$$\chi^{(+)} = (v_2\varphi_1^{(+)} - v_1\varphi_2^{(+)})/v_{12}\,, \qquad \chi^{(0)} = (v_2H_1 - v_1H_2)/v_{12}\,, \qquad (5)$$

$$M^2 = \frac{N_c}{8\pi^2}\,|\,\mathrm{Tr}\,h_1^+h_2\,|\,\Lambda^2(x_{12} + 1/x_{12}), \quad x_{12} = v_1/v_2, \quad v_{12}^2 \simeq v_1^2 + v_2^2\,. \qquad (6)$$

To evaluate M^2 in order of magnitude we set $|\mathrm{Tr}h_1^+h_2| \simeq m_b^2/v_1v_2$, $N_c = 3$, $x_{12} + x_{12}^{-1} = 2$, $v_{12} = 150$ GeV, and $m_b \simeq 4.5$ GeV. In the technicolor model the ultraviolet cutoff is[2] $\Lambda \sim \Lambda_{\mathrm{TC}}(m_q^{-1}\Lambda_{\mathrm{TC}})^{1/2}$. With $\Lambda_{\mathrm{TC}} \simeq 1$ TeV we then find $M \simeq 175$ GeV. If the Yukawa constant corresponds to the mass of the heaviest quark, for example, $m_q \simeq 70$ GeV, then we find $M \simeq 700$ GeV.

The remaining $n-2$ massless neutral Goldstone bosons must be orthogonal with respect to the "heavy pseudo-Goldstone boson" (5), so that the components of the fields φ_1 and φ_2 must appear in them in the same combination, $v_1\varphi_1 + v_2\varphi_2$, as in (3) for the true Goldstone bosons $g^{\pm}$ and $g^{(0)}$. The interaction of the light neutral pseudo-Goldstone bosons with fermions will thus be diagonal in the flavors, as can also be seen directly.

Let us determine the new doublet fields

$$\Phi = (v_1\varphi_1 + v_2\varphi_2)/v_{12}, \qquad \chi = (v_2\varphi_1 - v_1\varphi_2)/v_{12}\,. \qquad (7)$$

The interaction with fermions in accordance with (5) leads to a large mass, (6), of the components of the field χ. On the other hand, the vacuum expectation value of the field χ is obviously zero. From the standpoint of light pseudo-Goldstone bosons, therefore, there is only a single Higgs field Φ, which gives the quarks masses $Q = -1/3$, and $n-2$ doublets, which do not interact with fermions.

The mass M^2 in (6) is far larger than the ordinary masses of pseudo-Goldstone bosons.[2] Accordingly, if an interaction of type (1) is not the sole source of their masses, the mixing of the components of the field χ with the light pseudo-Goldstone bosons will still be small, $\sim m^2/M^2$, where m is the "ordinary" mass of light pseudo-Goldstone bosons.

It is easy to see that again the general case, with an arbitrary number of scalar fields giving masses $Q = -1/3$ and $Q = 2/3$ to the quarks and leptons, all combinations of neutral pseudo-Goldstone bosons which change flavors, along with the corresponding charged pseudo-Goldstone bosons, acquire a large mass, on the order of that in (6). The only states which may remain light are states of the type Φ in (7), which clearly conserve flavor upon their exchange. The low-energy phenomenology of pseudo-Goldstone bosons will be described by the model with natural flavor conservation,[1] mentioned above, which contains no more than three doublets of Higgs fields which interact with quarks and leptons. We wish to emphasize that it is quite possible that (for example) a minimum version with only a single Higgs field may also arise.

At low energies, heavy pseudo-Goldstone bosons may arise by virtue of the appearance of neutral currents which do not conserve flavor. The most severe restriction on the masses of the corresponding bosons comes from $K_L^0 - K_S^0$ splitting. The pseudoscalar P, which has the sd coupling

$$\xi\, m_s /2v\, \bar{s} i \gamma_5 d P + h.c., \quad v = 250 \text{ GeV}, \tag{8}$$

must have a mass $m_P > |\xi|$ 1.8 TeV (a weaker limitation, $m_S > |\xi|$ 0.5 TeV is found on the scalar coupling). The mass in (6) is thus literally insufficient for the necessary suppression of $K - \overline{K}$ transitions. However, the uncertainties in both the mass in (6) and the resulting off-diagonal couplings, (8), rules out a final conclusion. For the other decays $(\mu^+ \to e^+ e^- e^+, K_L^0 \to \mu^+ \mu^-, K^+ \to \pi^+ e^- \mu^+, \mu \to e\gamma)$, etc., a restriction no worse than $m \gtrsim 150$ GeV is found for the masses of the corresponding bosons for natural values of the coupling of the type in (8). A possible admixture of heavy pseudo-Goldstone bosons in the light pseudo-Goldstone bosons (see the discussion above) in processes of the type $K \to \overline{K}$, where the change in flavor should occur twice, makes a negligible contribution, while in decays with a single change in flavor the contribution may be of the same order of magnitude as in the exchange of a heavy boson considered here.

We wish to emphasize that in this proposed mechanism for suppressing the neutral currents which alter the flavors we have used only two properties of the technicolor theories: the existence of global symmetries (exact or approximate) to within radiative corrections and the presence of a physical scale dimension for the ultraviolet cutoff. Consequently, this mechanism might prove more general—not depending on the technicolor concept.

In conclusion we will prove the validity of (4)–(6) in terms of technicolor theory. Specifically, we consider a model with a single doublet of ordinary fermions, $q = (u,d)$, and two doublets of techniquarks, $Q_1 = (U_1,D_1)$ and $Q_2 = (U_2,D_2)$; here $\langle \overline{Q}_1 Q_1 \rangle \neq 0$ and $\langle \overline{Q}_2 Q_2 \rangle \neq 0$. Let us assume that the current masses of the fermions u and d arise from the exchange of expanded-technicolor bosons E_1 and E_2 with the interaction

$$g_1 (\overline{Q}_1 \gamma_\mu q) E_{1\mu} + g_2 (\overline{Q}_2 \gamma_\mu q) E_{2\mu} + h.c., \tag{9}$$

and for simplicity we set $M_{E_1} = M_{E_2}$. The calculations in (4)–(6) actually correspond to a calculation for the diagram in Fig. 1a with the corresponding contraterms. In this case this diagram, which disrupts the SU(2)×U(1) invariance of the independent rotations Q_1 and Q_2, is as shown in Fig. 1b. A corresponding calculation leads to the local four-fermion operator

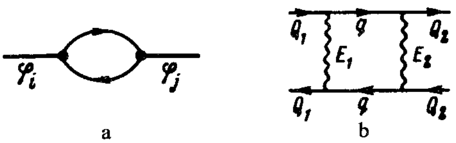

FIG. 1.

$$\delta\mathcal{H} = \frac{g_1^2 g_2^2}{16\pi^2 M_E^2}\left\{\frac{3}{4}(\bar{Q}_2^i\gamma_\mu Q_1^j)(\bar{Q}_1^j\gamma_\mu Q_2^i) - \frac{3}{2}(\bar{Q}_2^i\gamma_\mu\gamma_5 Q_1^j)(\bar{Q}_1^j\gamma_\mu\gamma_5 Q_2^i)\right\} \tag{10}$$

(i and j are the technicolor indices), which leads to the masses of the physical pseudo-Goldstone bosons on the basis of the Dashen formula[4]:

$$M^2 = \frac{9 g_1^2 g_2^2}{128\pi^2 M_E^2}(F_{\pi_1}^{-2} + F_{\pi_2}^{-2})\,|\langle 0|(\bar{Q}_1 Q_1)(\bar{Q}_2 Q_2) - (\bar{Q}_1\gamma_5 Q_1)(\bar{Q}_2\gamma_5 Q_2)|0\rangle|. \tag{11}$$

Taking into account the relationship with the effective low-energy parameters,

$$v_1 = F_{\pi_1}, \quad h v_1 = \frac{g_1^2\,|\langle\bar{Q}_1 Q_1\rangle|}{M_E^2} \quad (\text{and } 1\to 2)\,, \tag{12}$$

and using the hypothesis of a vacuum construction for the four-fermion expectation values, we can write (11) in the form

$$M^2 = \frac{9}{64\pi^2} h_1 h_2 M_E^2\left(x + \frac{1}{x}\right), \tag{13}$$

which differs from (6) only by a factor of 9/8 when we identify Λ^2 with M_E^2.

In the technicolor theories, we might note, the presence of interactions (9) should be accompanied by the existence of an extended-technicolor boson with a mass $\sim M_E$ and a coupling of the type $\bar{Q}_1\gamma_\mu Q_2$. The exchange of this boson could lead to a slightly higher mass of a pseudo-Goldstone boson, $\widetilde{M}^2 \sim (\pi/\alpha_{\mathrm{ETC}})M^2$, where M^2 is a mass of the order of (11) or (13).

I am deeply indebted to A. A.Ansel'm for interest in ths study and for useful discussions and to V. Yu. Petrov for a discussion of the results.

[1]S. L. Glashow and S. Weinberg, Phys. Rev. D **15**, 1958 (1977).
[2]A. A. Ansel'm and N. G. Ural'tsev, Zh. Eksp. Teor. Fiz. **82**, 1725 (1982) [Sov. Phys. JETP **55**, 997 (1982)].
[3]S. Coleman and E. Weinberg, Phys. Rev. D 7, 1888 (1973).
[4]R. Dashen, Phys. Rev. **183**, 1245 (1969).

Translated by Dave Parsons
Edited by S. J. Amoretty