Can very light neutral vector bosons exist?*

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Models which unify strong, electromagnetic, and weak interactions contain many neutral gauge fields. A generalization of the Pati-Salam model with fractional-charge quarks is used as an example to investigate whether all gauge fields, except those coupled to the $SU(2)\times U(1)$ subgroup of Weinberg-Salam, must develop large masses and become unobservable. The possible existence of a neutral boson with a mass much smaller than that of $W^{\pm}$ is considered. Constraints from known properties of neutral currents together with the assumption of an explicit symmetry-breaking mechanism produce a model with two neutral vector bosons, neither of which has the Weinberg-Salam structure. One particle is very light and decoupled from left-handed neutrinos, and the other is heavy and responsible for observed neutral-current interactions. A brief comparison with experimental data does not exclude this model. The existence of a very light, neutral, purely charm-changing neutral vector boson is considered.

I. INTRODUCTION

The Weinberg-Salam model of weak and electromagnetic interactions predicted the existence of a massive, neutral intermediate vector boson.[1] Subsequent experimental work has confirmed the existence of neutral-current weak-interaction processes.[2] All presently existing data is in apparent agreement with the simple Weinberg-Salam model. However, there has been considerable interest in generalizations of the basic gauge model which attempt to unify strong, electromagnetic, and weak interactions.[3-7] A characteristic feature of the generalized models is the existence of a plethora of gauge vector bosons. Most of the unwanted ones are removed by invoking the miraculous Higgs mechanism to give them very large masses and thereby render them unobservable. The universality of the charged-current weak interactions provides good evidence for the existence of a single charged intermediate vector boson (IVB). In the absence of similar evidence for neutral-current events, we ask the question: Can there exist more than one "light" neutral vector meson? By light we mean one having a mass of order that of the charged IVB or lighter. In particular we ask whether present knowledge is sufficient to exclude the existence of neutral IVB's which are light compared to the charged one.

In this paper we start with a unified model which is a generalization of the Pati-Salam model[6] to $SU(4)_L\times SU(4)_R\times SU(4)_c$ and investigate whether the photon and the Weinberg-Salam neutral vector field are the only survivors of the fifteen neutral color-singlet gauge fields present in the exact symmetry limit. In this model leptons are a fourth color of quark. Each quark comes in four flavors including charmed. We show that, if the symmetry-breaking process is such that there is a single light neutral gauge field in addition to the photon, it must couple to the Weinberg-Salam current. The mass of the vector boson is arbitrary since it depends on the symmetry-breaking mechanism. Experiment suggests its mass is near that predicted by the Weinberg-Salam model.[2,8] Inasmuch as it is difficult to distinguish a two-neutral-current model in which both IVB's have masses on the order of the charged $W^{\pm}$ from a model with a single neutral current, we concentrate on the possible existence of a second truly light (mass between 3 and 15 GeV/c^2) gauge field.[9] We look for conditions on its existence and for experimental situations where its effects would be dramatic. We do not discuss situations where the light boson would produce a few-percent correction to a dominant Weinberg-Salam current.

The basic result is that it is possible to construct models with a wide spectrum of light neutral currents. As many as three charm-conserving neutral IVB's are possible in the extended Pati-Salam model. In addition there can be a purely charm-changing neutral current. These models are consistent with the accepted dogma on charged- and neutral-current weak interactions. However, there are too many free parameters to permit a meaningful confrontation with experiment.

To restrict the models, we utilize a symmetry-breaking scheme in which the weak-interaction IVB's acquire their mass through the Higgs mechanism. The scalar fields which develop vacuum expectation values (VEV) transform as $(4,\bar{4},1)$ and $(4,\bar{4},15)$ under $SU(4)_L\times SU(4)_R\times SU(4)$. This choice is justified from the standpoint of economy of scalar fields and from considerations of the fermion mass matrix. The result is a model with very few undetermined parameters. We find that if there is but one neutral IVB, it must couple to a current with the Weinberg-Salam structure and have the Weinberg-Salam mass. The model with three "light" neutral IVB's is in gross conflict with neutral-current neutrino experiment since it produces at least one IVB which is very light and has unsuppressed couplings to left-handed neutrinos. If there are two charm-conserving IVB's, one of them can be very light yet effectively decoupled from left-handed neutrinos. In this two-current model with $(\overline{4}, 4, 1)$ or $(\overline{4}, 4, 15)$ symmetry breaking, the structure of the neutral weak current is uniquely specified. The model proves to be consistent with present experimental evidence.

Direct observation of the very light IVB in a two-neutral-current model is very difficult. It does compete with the photons in e^+e^- annihilation. Production cross sections and decay widths are estimated where possible. The width for decay into hadrons is large, and the width for decay into known leptons is a few percent of the hadronic width. A substantial decay width into right-handed neutrinos also occurs. A particularly interesting possibility is that of a purely charm-changing neutral current whose leptonic couplings are to right-handed neutrinos and whose hadronic couplings are such that it would be seen only in $|\Delta C| = 0$ nonleptonic processes.[10] It generates large $D^0\overline{D}^0$ mixing, where D^0 is the neutral, strangeness-zero, charmed pseudoscalar meson.

In the next section we apply theoretical and phenomenological constraints to reduce the number of possible neutral gauge bosons from fifteen to five. These five, from which the photon and as many as four neutral IVB's may be constructed, are subject to additional conditions. We show that if only one neutral IVB is present, its current must have the Weinberg-Salam structure. Such a boson cannot be very light and consistent with neutral-current data.[8] Hence a very light neutral IVB can only appear in models with more than one neutral IVB. Moreover, if there is a very light neutral IVB, its coupling to left-handed neutrinos must be suppressed relative to its other leptonic and hadronic couplings.

The third section is devoted to a discussion of the consequences of assuming a specific symmetry-breaking mechanism. The Weinberg-Salam neutral-IVB mass constraint is shown to result in models with a single neutral IVB.

In the fourth section, we extend Sec. III and consider models with more than one neutral IVB. If there are three light neutral IVB's, one must be very light and have unsuppressed coupling to left-handed neutrinos. Such a model is in gross conflict with experiment. A model with two neutral IVB's, one of them very light, is developed. The symmetry-breaking mechanism of Sec. III and the condition that a very light IVB must have suppressed neutrino coupling uniquely specifies the structure of both neutral IVB's up to the sign of one mixing angle. The model is compared with experiment and shown to be consistent with neutral-current phenomenology for semileptonic and purely leptonic processes.

The final section considers the problem of detecting very light neutral IVB's that are decoupled from left-handed neutrinos. Predictions are made for e^+e^- annihilation and $D^0\overline{D}^0$ mixing.

II. CONSTRAINING THE NEUTRAL CURRENT

The most general neutral current in the Pati-Salam model[6] has the form

$$\overline{K}^{\mu}_{L,R} = \gamma_\mu \frac{1 \pm \gamma_5}{2} \begin{pmatrix} a+b+c & 0 & 0 & \alpha \\ 0 & -a+b-c & \beta & 0 \\ 0 & \beta^* & a-b-c & 0 \\ \alpha^* & 0 & 0 & -a-b+c \end{pmatrix}, \tag{1}$$

where a, b, c are real parameters and α, β are complex. The L, R subscript on $\overline{K}^{\mu}_{L,R}$ refer to left- and right-handed currents. $\overline{K}^{\mu}_{L,R}$ is a matrix representation of an $SU(4)_L$ or an $SU(4)_R$ current. The fermions in this model have the representation

$$\psi^{L,R}_{i\alpha} = \frac{(1 \pm \gamma_5)}{2} \begin{pmatrix} u_r\, u_w\, u_b\, \nu_e \\ d'_r\, d'_w\, d'_b\, d^- \\ s'_r\, s'_w\, s'_b\, \mu^- \\ c_r\, c_w\, c_b\, \nu_\mu \end{pmatrix}. \tag{2}$$

The four quarks are μ, d', s', and c. The subscripts r, w, b indicate color. The d' and s' quarks contain a Cabibbo rotation: $d' = d\cos\theta_C + s\sin\theta_C$, $s' = s\cos\theta_C - d\sin\theta_C$, and d and s are the quarks of definite isospin and strangeness. The α subscript on $\psi_{i\alpha}$ is the color index. Note that this model explicitly contains right-handed neutrinos. The explicit form of the left-handed current is

$$K^{\mu}_{L} = \overline{\psi}^{L}_{i\alpha} \overline{K}^{\mu}_{ij} \psi^{L}_{j\alpha}, \tag{3}$$

where repeated indices are summed over. In addition to the fourteen currents given by (1) and (3), there is the current which is the 15th component of color:

$$V_{15}^{\mu} = \bar{\psi}_{i\alpha}^{L} \bar{V}_{15\alpha\beta}^{\mu} \psi_{i\beta}^{L} + \bar{\psi}_{i\alpha}^{R} \bar{V}_{15\alpha\beta}^{\mu} \psi_{i\beta}^{R} \ . \tag{4}$$

The matrix $\bar{V}_{15}^{\mu}$ acts in color space and is given by

$$\bar{V}_{15}^{\mu} = \frac{\gamma_\mu}{2\sqrt{6}} \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & -3 \end{pmatrix} . \tag{5}$$

Implicit in our calculations is the fact that there is a symmetry-breaking stage characterized by a large mass scale which separates the quark and lepton sectors in color space. In a model with fractionally charged quarks, the SU(3) color subgroup which acts in quark space is unbroken. There are three gauge coupling constants in the unified theory: g_L, g_R, and f, the color coupling constant. We could ignore the existence of the $\mathrm{SU}(4)_c$ gauge group and regard (5) as a U(1) generator which couples with strength f and differentiates between quarks and leptons.

The accumulated wisdom of recent years places constraints on the form of the neutral current.

Condition 1. There is a single charged gauge boson W, which is left handed. All other charged gauge fields in $\mathrm{SU}(4)_L$ and $\mathrm{SU}(4)_R$ must develop large masses in the symmetry-breaking process. The W picks up a mass in what we call the "final" stage of symmetry breaking. Although all symmetries are broken simultaneously, the "final" stage is characterized by a mass scale small compared to all others. We are interested in neutral currents which couple to gauge fields that are massless but for this final stage of symmetry breaking. The generators of the current coupled to the W, $K^{W^\pm}$ are generators of an $\mathrm{SU}(2)_L$ subgroup of $\mathrm{SU}(4)_L$:

$$[K^{W^+}, K^{W^-}] = K^{N_1}/\sqrt{2} \ , \tag{6}$$

where

$$K^{W^+} = \tfrac{1}{2} \begin{pmatrix} 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 1 & 0 \end{pmatrix} . \tag{7}$$

If the spontaneous-symmetry-breaking mechanism leaves W massless, this SU(2) subgroup must remain unbroken. Hence the direct consequence of condition 1 is that the gauge field coupled to $K_L^{N_1}$ must also remain massless. This generator is the term proportional to the constant c in (1):

$$K^{N_1} = \frac{1}{2\sqrt{2}} \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & -1 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix} . \tag{8}$$

Condition 2. The d' and s' quarks in (2) contain the Cabibbo rotation. The neutral-current interaction conserves strangeness.[2] We require that both left- and right-handed currents have $\Delta S = 0$. In addition the neutral currents must contain no direct $e \longleftrightarrow \mu$ coupling, or else $\mu \to ee\bar{e}$ decays would be comparable to $\mu \to e\nu\bar{\nu}$ decays. Together these constraints imply condition 2 that $\beta = 0$ and $a = b$ in (1).

Condition 3. In order that W be the only left-handed charged IVB, $\mathrm{SU}(4)_L$ must be broken to $\mathrm{SU}(2)_L \times G$, where G is a subgroup of $\mathrm{SU}(4)_L$ containing no charged currents. Explicit commutation of the charged-current generator (7) with (1) shows that the commutator is proportional to another charged current unless $a = b = 0$, $\alpha = \beta$. Since $\beta = 0$ from condition 2, we see that only c can be nonzero, corresponding to $K_L^{N_1}$, the $\mathrm{SU}(2)_L$ neutral generator. Thus G is a null group and $K_L^{N_1}$ is the only possible neutral left-handed current. Henceforth, all discussion of the parameters of (1) will be in reference to the right-handed currents belonging to $\mathrm{SU}(4)_R$.

Condition 4. We assume the neutral currents conserve CP. CP violation is due to either to a separate superweak interaction[11] or to mediation by gauge bosons so heavy as to be irrelevant for our purposes.[12] As a result α and β in (1) must be real.

Condition 5. The recent discovery of charmed particles shows that the D^0 decays to $K\pi$ but not $\pi\pi$ final states.[13] Thus, the neutral current cannot produce $\Delta C = 1$, $\Delta S = 0$ nonleptonic interactions. In the absence of miraculous cancellations that remove such transitions, we conclude that a charm-changing neutral current cannot have a charm-conserving piece. The only CP-invariant charm-changing current belongs to $\mathrm{SU}(4)_R$ and is denoted by K_R^{cc} (cc is for charm-changing)

$$K_R^{cc} = \tfrac{1}{2} \begin{pmatrix} 0 & 0 & 0 & 1 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 1 & 0 & 0 & 0 \end{pmatrix} . \tag{9}$$

If this current occurs in the weak interactions, it must be coupled to a distinct gauge boson Z^{cc} to preclude mixing with $\Delta C = 0$ currents.

Conditions 1–5 imply that charm-conserving neutral currents may contain components of only one additional current other than $K_L^{N_1}$, $K_R^{N_1}$, and V_{15}. This current is denoted by K_R^N:

$$K_R^N = \frac{1}{2}\begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix} . \tag{10}$$

At this stage we are left with five neutral currents proportional to $K_L^{N_1}$, $K_R^{N_1}$, K_R^N, K_R^{cc}, and V_{15}. The K_R^{cc} current is unmixed. It induces $u \longleftrightarrow c$ and $\nu_e \longleftrightarrow \nu_\mu$ transitions, where the neutrinos are right-handed. The K_R^{cc} current would be very difficult to detect experimentally. We will consider this current again in the final section. One linear combination of the remaining four currents is the electromagnetic current coupled to the photon. For fractional quark charges this current is

$$J = \sqrt{2}\, e(K_L^{N_1}/g_L + K_R^{N_1}/g_R - V_{15}/f\sqrt{3}) . \tag{11}$$

(The gauge coupling constants g_L, g_R, and f are divided out to produce the correct coupling to the photon.) The photon remains massless through all stages of symmetry breaking. In addition there are three currents orthogonal to (11), and we wish to know whether one of them can couple to a vector boson whose mass is much less than that of the W. Is there a choice of constants c_L, c_R, c_Z, and c_{15} such that the current

$$K_1 = c_L K_L^{N_1} + c_R K_R^{N_1} + c_Z K_R^N + c_{15} V_{15} \tag{12}$$

couples to a light-mass vector boson? If such a light IVB exists (denoted by Z_1), then it would contribute to neutrino-induced processes with a strength S characterized by

$$S = \frac{G_F}{2\sqrt{2}} \left\{ \left(\frac{m_W}{m_{Z_1}} \right)^2 \left(c_L + \frac{\sqrt{3}\, f}{g_L} c_{15} \right) \right\} . \tag{13}$$

Experiments on neutrino-induced neutral-current processes indicate that the factor in curly brackets is of order unity. By assumption $m_W{}^2/m_{Z_1}{}^2$ is large. Hence, we have the following.

Condition 6. The coupling to left-handed neutrinos is suppressed:

$$c_L + \frac{\sqrt{3}\, f}{g_L} c_{15} \approx 0 . \tag{14}$$

We choose for now to ignore $m_{Z_1}{}^2/m_W{}^2$ effects and to make this condition exact. (In Sec. IV we will consider consequences of relaxing condition 6.) This condition means that a very light neutral gauge field can exist only if it does not couple to left-handed neutrinos. Condition 6 and orthogonality to the electromagnetic current (11) specify that the structure of K_1 is given by

$$K_1 = c_L \left[K_L^{N_1} - \frac{g_L}{\sqrt{3}\, f} V_{15} - \frac{g_R}{g_L}\left(1 + \frac{1}{3}\frac{g_L{}^2}{f^2}\right) K_R^{N_1} \right] + c_Z K_R^N . \tag{15}$$

For comparison the Weinberg-Salam neutral current is given by[7]

$$K_{\text{WS}} = \frac{[1 + (\lambda\omega)^2] K_L^{N_1} - \lambda^2 \omega K_R^{N_1} + \lambda V_{15}}{[(1 + \lambda^2\omega^2)(1 + \lambda^2 + \omega^2\lambda^2)]^{1/2}}$$

$$= \cos\theta_W K_L^{N_1} - \sin\theta_W \frac{\lambda\omega K_R^{N_1} - V_{15}}{(1 + \lambda^2\omega^2)^{1/2}} , \tag{16}$$

where we have introduced the parameters $\lambda = \sqrt{3}\, f/g_L$ and $\omega = g_L/g_R$. In a model where the very light current of (15) exists, the current responsible for observed neutral-current events cannot be given by (16), since it is not orthogonal to (15) unless $c_L = 0$. Thus, although Z_1 may not participate in those reactions, models containing a Z_1 should be distinguishable from the conventional model. However, without further information there are at least two effective parameters that must be fixed to specify the current coupled to neutrinos; both c_L and the mass of the gauge field or fields must be determined. For this reason we use a definite symmetry-breaking mechanism in the next section to fix all parameters.

We conclude this section by returning to (12) and showing that if all gauge bosons but one are given large masses, the remaining one must have the Weinberg-Salam structure of (16). If there is to be but one neutral IVB, then at the stage of symmetry breaking where the W is massless there must be two neutral massless gauge fields. By condition 1, one of them must couple to $K_L^{N_1}$. Since the photon couples to (11), the other massless combination couples to

$$K' = \frac{\lambda\omega K_R^{N_1} - V_{15}}{(1 + \lambda^2\omega^2)^{1/2}} . \tag{17}$$

After the final stage of symmetry breaking, we have the electromagnetic current (11) and an orthogonal combination of (17) and $K_L^{N_1}$. This is exactly the current given in (16). The mass of the field coupled to (16) depends on the explicit symmetry-breaking mechanism used in the final stage.

III. SYMMETRY BREAKING

In the absence of a convincing mechanism for dynamical symmetry breaking, all gauge fields coupled to broken symmetries develop a mass through the Higgs mechanism. To break the symmetry from $SU(4)_L \times SU(4)_R \times SU(4)_c$ to $U(1) \times SU(3)_c$, we need a number of different multiplets of scalar fields. We are specifically concerned with the "final" stage of symmetry breaking in which the IVB's mediating the weak interactions acquire their masses. These masses are less than or of order m_W and are much lighter than the masses produced by "previous" stages of symmetry breaking.

The smallest representations for Higgs fields which breaks the $SU(2)_L$ symmetry generated by the charged current, gives mass to neutral right-handed vector bosons, and preserves electromagnetic and color symmetry are $(4,\bar{4},1)$ and $(\bar{4},4,1)$. Their effects are identical and we could choose $(4,\bar{4},1)$ as the basis of the most economical model for the "final" stage of symmetry breaking. This representation is used in Ref. 6 to give masses to the Weinberg-Salam W and Z. Moreover, since the fermion multiplets transform as $(4,1,\bar{4})$ and $(1,4,\bar{4})$, Higgs field transforming as $(4,\bar{4},1)$ and $(4,\bar{4},15)$ are needed to give the fermions mass.[14] In particular, a $(4,\bar{4},15)$ is needed to differentiate between quarks and leptons. However, charge and color conservation restrict the $(4,\bar{4},15)$ vacuum expectation value to one which transforms like V_{15} in color space. The effect of such a $(4,\bar{4},15)$ on the masses of the IVB is identical to that produced by a $(4,\bar{4},1)$. Thus, our model of symmetry breaking uses a single VEV transforming as a $(4,\bar{4},1)$.

The most general VEV $\langle a\rangle$ that belongs to $(4,\bar{4},1)$ and conserves charge is given by

$$\langle a\rangle_{ij,\alpha\beta}=\delta_{\alpha\beta}\begin{pmatrix} a_1 & 0 & 0 & a_2\\ 0 & a_2 & a_4 & 0\\ 0 & a_5 & a_6 & 0\\ a_7 & 0 & 0 & a_8\end{pmatrix} . \tag{18}$$

Calculation of the vector-meson mass matrix is a familiar procedure in gauge theories. The result for the mass of the W is

$$m_W{}^2=\frac{g_L{}^2}{8}\sum_{i=1}^{8}(a_i)^2=\frac{g_L{}^2}{8}r . \tag{19}$$

We identify the neutral gauge fields by the same symbol as their currents and choose a set of basis fields $V_1=V_{15}$, $V_2=(N_L^1+\omega N_R^1)/(1+\omega^2)^{1/2}$, $V_3=(-\omega N_L^1+N_R^1)/(1+\omega^2)^{1/2}$, $V_4=N_R$, $V_5=Z^{cc}$. Again we use $\omega=g_L/g_R$. The mass term in the vector-meson Lagrangian is

$$\mathcal{L}^{\text{mass}}=\tfrac{1}{2}V_i m_{ij}{}^2 V_j , \tag{20}$$

where

$$m^2=\frac{g_R{}^2}{4}\begin{pmatrix} 0 & 0 & 0 & 0 & 0\\ 0 & 0 & 0 & 0 & 0\\ 0 & 0 & \eta^2 r & \eta t & \eta u\\ 0 & 0 & \eta t & s & 0\\ 0 & 0 & \eta u & 0 & s\end{pmatrix} , \tag{21}$$

with $\eta=[(1+\omega^2)/2]^{1/2}$. The quantities r,s,t,u are functions of the a_i:

$$r=\sum_{i=1}^{8}a_i{}^2 , \tag{22a}$$

$$s=a_1{}^2+a_2{}^2+a_7{}^2+a_8{}^2 , \tag{22b}$$

$$t=a_1{}^2-a_2{}^2+a_7{}^2-a_8{}^2 , \tag{22c}$$

$$u=2(a_1a_2+a_7a_8) . \tag{22d}$$

The physical vector mesons diagonalize $\mu^2=M^2+m^2$, where M^2 is that part of the total mass matrix that introduces the superheavy mass scale. The photon is a zero-mass eigenstate of M^2 and m^2.

In the preceding section we argued that prior to the "final" step of symmetry breaking, the fields coupled to $K_L^{N_1}$ and K^1 [given by (17)] must be massless. In a model with a single neutral current, these generate the only unbroken neutral symmetries. According to (11) the electromagnetic current is just

$$J=\sin\theta_W K_L^{N_1}+\cos\theta_W K^1 , \tag{23}$$

where the Weinberg angle is defined implicitly in (16) [$\sin\theta_W=\lambda/(1+\lambda^2+\lambda^2\omega^2)^{1/2}$]. The Weinberg-Salam current is given in (16). Using the basis V_i, we have

$$Z_{\text{WS}}=\frac{\sin\theta_W}{(1+\lambda^2\omega^2)^{1/2}}V_1+\frac{\cos\theta_W}{(1+\lambda^2\omega^2)(1+\omega^2)^{1/2}}V_2-\frac{\omega}{(1+\omega^2)^{1/2}\cos\theta_W}V_3 . \tag{24}$$

To lowest order the mass of Z_{WS} is

$$m_{Z_{\text{WS}}}{}^2=(Z_{\text{WS}},m^2Z_{\text{WS}})=\frac{g_R{}^2}{4}\frac{\omega^2\eta^2}{(1+\omega^2)\cos^2\theta_W}=\frac{m_W{}^2}{\cos^2\theta_W} . \tag{25}$$

Hence, in a single-current model, the $(4,\bar{4},1)$ symmetry-breaking mechanism reproduces the Weinberg-Salam model of unified weak and electromagnetic interactions. Moreover, inclusion of any number VEV's that transform as a $\underline{4}$ under $SU(4)_L$ and conserve charge [e.g., a $(4,\bar{4},15)$] yields the Weinberg-Salam mass constraint.

IV. MULTI-NEUTRAL-CURRENT MODELS

If there is a single light neutral current which couples to a symmetry that is broken only in the "final" stages, weak-interaction models with $(4,\bar{4},1)$ symmetry breaking are identical to the Weinberg-Salam model. Thus, we must consider models with at least two charm-conserving neutral currents. In addition, we require that there be at least one very light gauge field after all symmetries are broken. Such an IVB, the Z_1, must couple to the current given by (15) in order that it decouple from left-handed neutrinos and be orthogonal to the photon. If the Z_1 is massless up to the "final" stage of symmetry breaking, $M^2 Z_1 = 0$, where M^2 is that portion of the mass matrix responsible for the superheavy mass scale. Moreover, if $m_{Z_1} \ll m_W$, the symmetry coupled to Z must be approximately conserved by the VEV's in (18). In particular we must have

$$\langle Z_1, m^2 Z_1\rangle \approx 0 \ , \tag{26}$$

where m^2 is given by (21). We make this approximate equality exact and write Z in the form

$$\begin{aligned} Z_1 &= c_L\left[N_L^1 - \frac{1}{\lambda}V_{15} - \frac{1}{\omega}\left(1+\frac{1}{\lambda^2}\right)N_R^1\right] + c_Z N_R \\ &= -\cos\phi\,\frac{\omega\lambda V_1 + [\omega/(1+\omega^2)^{1/2}]V_2 + [(1+\lambda^2+\omega^2\lambda^2)/(1+\omega^2)^{1/2}]V_3}{[(1+\lambda^2)(1+\lambda^2+\lambda^2\omega^2)]^{1/2}} + \sin\phi\, V_4 \ . \end{aligned} \tag{27}$$

Equation (26) constrains the elements of the mass matrix (21). We have

$$r\xi^2\cos^2\phi - 2t\xi\cos\phi\sin\phi + s\sin^2\phi = 0 \ , \tag{28}$$

where

$$\xi = \left[\frac{1+\lambda^2+\omega^2\lambda^2}{2(1+\lambda^2)}\right]^{1/2} . \tag{29}$$

One solution to (28) is $\cos\phi = 0$ and $s=0$. If $\cos\phi = 0$, the very light IVB is pure N_R. In addition if $s=0$, then $t=u=0$ from (22). When N_R is decoupled from the other fields, the arguments at the ends of the preceding two sections can be used to show that if there is a single additional neutral gauge field (a two-current model) that the second field has the Weinberg-Salam mass and couples to the Weinberg-Salam neutral current. We prove below that within the context of $(4,\bar{4},1)$ symmetry breaking, a two-current model is the only possibility. The decoupled N_R does not interact directly with electrons, muons, or left-handed neutrinos. It has no effect on neutrino-induced interactions or on e^+e^- annihilation. This exotic particle would be important for the explanation of parity-violating effects in $\Delta S = 0$ nuclear transitions. However, detailed calculation of these effects is beyond the scope of this paper.

If $\cos\phi \neq 0$ and $s \neq 0$ in (28), then (28) becomes a quadratic equation for $\tan\phi$ with solution $\tan\phi = \xi[t \pm (t^2 - rs)^{1/2}]/s$. Explicitly use of (22) shows that $t^2 \leq rs$. Therefore, $\tan\phi$ is real only if $t^2 = rs$. (As a parenthetical remark we note that this condition is approximately satisfied when the vacuum expectation values are constrained to generate a physical fermion mass matrix.) When the conditions on the a_i are fully exploited, we find $r = s = |t|$, $u=0$. The mixing angle ϕ is fixed to be

$$\tan\phi = \pm\,\xi = \pm\left[\frac{1+\lambda^2+\omega^2\lambda^2}{2(1+\lambda^2)}\right]^{1/2} . \tag{29}$$

The structure of Z_1 is uniquely determined in this model of symmetry breaking. The vanishing of the parameter u guarantees that the charm-changing field $V_5 = Z_{cc}$ decouples from the charm-conserving fields.

In the basis spanned by the V_i there can be two IVB's in addition to the photon and Z_1. If both of the additional ones are light, their masses must not receive any contributions from the superheavy mass matrix. Rather the masses must be entirely generated by m^2 given in (21) with $u=0$ and $r=s=|t|$. However, explicit calculation with m^2 shows that it has just one nonzero eigenvalue. Thus, one of the two additional IVB's must also be massless, yet have substantial coupling to left-handed neutrinos. Such an IVB is in gross conflict with experiment. Hence, one of the two must be superheavy, and there is a single neutral massive IVB in addition to Z_1 and the photon.

The structure of Z_h is uniquely determined if we require that it be orthogonal to Z_1, the photon, and Z_{sh}, where Z_{sh} is that linear combination of V_{15}, N_L^1, N_R^1, and N_R that develops a superheavy mass. Z_{sh} is fixed to be that state which is orthogonal to Z_1 and the photon and has no projection onto N_L^1. The absence of N_L^1 components in Z_1 is a consequence of the fact that the $SU(2)_L$ symmetry is preserved by the superheavy symmetry-breaking mechanism. The structure of Z_{sh} is

$$Z_{sh} = \omega\lambda V_{15} + N_R^1 \pm\sqrt{2}\, N_R \ . \tag{30}$$

The ± sign in (30) is the same sign ambiguity that appears in the expression for $\tan\phi$ in (29). Given (20) it is straightforward to calculate the state responsible for neutrino-induced neutral-current events:

$$Z_h = \frac{-3\lambda V_{15} - (3+\omega^2\lambda^2)N_1^L + \omega\lambda^2 N_1^R \pm\sqrt{2}\,\omega\lambda^2 N_R}{\{(3+\omega^2\lambda^2)[3(1+\lambda^2)+\omega^2\lambda^2]\}^{1/2}} . \tag{31}$$

When this state is transformed to the basis V_i, its mass can be calculated from (21) with $r=s=\pm t$, $u=0$. The result is

$$m_{Z_h}{}^2 = (Z_h, m^2 Z_h) = \frac{g_R{}^2}{4}\,\frac{r\omega^2}{2}\left[\frac{3(1+\lambda^2)+\omega^2\lambda^2}{3+\omega^2\lambda^2}\right]^{1/2}$$
$$= m_W{}^2\left[\frac{3(1+\lambda^2)+\omega^2\lambda^2}{3+\omega^2\lambda^2}\right]. \tag{32}$$

If the $SU(4)_L \times SU(4)_R \times SU(4)$ theory is embedded in a truly unified theory based on a single simple or semisimple group and if coupling-constant renormalization is neglected, we have $\omega^2=1$, $\lambda^2=\frac{3}{2}$, and $m_{Z_h}{}^2 = 2m_W{}^2$.[7] This result is to be compared with $m_Z{}^2 = 1.6\, m_W{}^2$ in the Weinberg-Salam theory with $\sin^2\theta_W = \frac{3}{8}$.

Of course if Z_1 is actually massless, it would compete with the photon and introduce parity-violating effects where they are not wanted. Moreover, a theory in which Z_h is the only neutral vector boson coupled to neutrinos is in apparent conflict with conventional neutrino phenomenology.[15] Therefore, we relax both the constraint on the coupling of Z_1 to left-handed neutrinos and the constraint on its mass. Equation (14) is replaced by

$$c_L + \frac{\sqrt{3f}}{g_L}\, c_{15} = \epsilon c_L , \tag{33}$$

and (26) is replaced by

$$\langle Z_1, m^2 Z_1\rangle = m_W{}^2\delta . \tag{34}$$

The parameters ϵ and δ are presumed to be small. The structure and mass of Z_h is only slightly perturbed by this change. However, Z_1 now contributes to neutrino-induced interactions with a strength proportional to ϵ/δ, a quantity which can be of order unity even when ϵ and δ are infinitesimal. The light gauge field now has the representation

$$Z_1 = c_L\left[N_L^1 - \frac{1-\epsilon}{\lambda}V_{15} - \frac{1}{\omega}\left(1+\frac{1-\epsilon}{\lambda^2}\right)N_R^1\right] + c_Z N_R$$
$$= -\frac{\cos\phi}{\Delta}\left[\omega\lambda(1-\epsilon)V_1 + \frac{\omega(1-\epsilon)}{(1+\omega^2)^{1/2}}V_2 + \frac{(1-\epsilon+\lambda^2+\omega^2\lambda^2)}{(1+\omega^2)^{1/2}}V_3\right]$$
$$+ \sin\phi\, V_4 . \tag{35}$$

The factor in square brackets is normalized by Δ. When (35) is used in (34), we find that $\tan\phi$ is shifted from the value in (29) by terms of order ϵ and δ. However, the precise magnitude of the shift is not fixed. The contribution of Z_1 to neutrino-induced interactions does not depend on the shift in ϕ to order ϵ. The $(4,\bar{4},1)$ symmetry-breaking mechanism does not uniquely fix the structure of Z_1 when its mass is not identically equal to zero.

We conclude this section by considering the phenomenology of neutral-current interactions when two neutral IVB's contribute. One is Z_h with the structure of (31) and the mass in (32). The other is Z_1 with the structure of (35) with $\tan\phi$ from (29). The mass of Z_1 and the parameter ϵ in (35) are small, but undetermined. The parameters λ and ω are fixed when the Pati-Salam model is embedded in a larger group. In the absence of renormalization $\lambda^2=\frac{3}{2}$ and $\omega^2=1$. If we renormalize the coupling constants consistent with a symmetry-breaking scheme in which either

$$G \to SU(2)_L \times U(1)\times U(1)\times SU(3)_{color}$$
$$\to U(1)_{em}\times SU(3)_{color} , \tag{36a}$$

or

$$G \to SU(4)_L \times SU(4)_R \times SU(4)_{color}$$
$$\to SU(2)_L \times U(1)\times U(1)\times SU(3)_{color}$$
$$\to U(1)_{em}\times SU(3)_{color} , \tag{36b}$$

then in the limit that the quark-gluon coupling constant is large compared with the electromagnetic coupling constant, we have $\lambda^2=\frac{1}{2}$ and $\omega^2=3$. The V_{15} coupling constant f must be distinguished from the quark-gluon coupling constant f_s although they are equal before renormalization. The two U(1) groups in (36a) or (36b) are needed to accommodate the neutral gauge fields which ultimately mix with the neutral member of $SU(2)_L$ to form Z_1, Z_h, and the photon.

We outline the steps necessary to calculate the effects of a neutral current which has the general structure

$$K = c_L K_L^{N_1} + c_R K_R^{N_1} + c_Z K_R^N + c_{15}V_{15} . \tag{37}$$

The effective Hamiltonian for neutrino interactions is

$$\mathcal{H}_{eff} = \frac{G_F}{\sqrt{2}}[\bar{\nu}\gamma_\lambda(1+\gamma_5)\nu]$$
$$\times\{\bar{\mu}\gamma^\lambda[c_+(1+\gamma_5)+c_-(1-\gamma_5)]\mu + g_V V_3^\lambda + g_A A_3^\lambda + g_V^1 V_0^\lambda + g_A^1 A_0^\lambda\} . \tag{38}$$

The various parameters have the following definitions:

$$c_+ = \frac{\mathcal{S}}{2}(-c_L + \lambda c_{15}),$$

$$c_- = \frac{\mathcal{S}}{2}\left(-\frac{c_R}{\omega} + \lambda c_{15}\right),$$

$$g_V = \frac{\mathcal{S}}{2}\left(2c_L + \frac{2c_R}{\omega} + \sqrt{2}\,\frac{c_Z}{\omega}\right),$$

$$g_A = \frac{\mathcal{S}}{2}\left(2c_L - \frac{2c_R}{\omega} - \sqrt{2}\,\frac{c_Z}{\omega}\right),$$

$$g_V^1 = \frac{\mathcal{S}}{2}\left(-\tfrac{4}{3}\lambda c_{15} + \sqrt{2}\,\frac{c_Z}{\omega}\right),$$

$$g_A^1 = \frac{\mathcal{S}}{2}\left(-\sqrt{2}\,\frac{c_Z}{\omega}\right). \tag{39}$$

Each term in (39) has the common factor $\mathcal{S} = (c_L + \lambda c_{15}) m_W^2/m_Z^2$. The quark currents V_3^λ and V_0^λ are isovector and isoscalar vector currents while A_3^λ and A_0^λ are isovector and isoscalar axial vector currents. In Tables I and II, the parameters in (39) are tabulated for both Z_1 and Z_h. Table I utilizes bare values of λ and ω, while renormalized values are used in Table II. All parameters for Z_1 are proportional to ϵ/δ. The tabulated parameters are used in a standard calculation of the cross-section ratios $R = \sigma(\nu N \to \nu x)/\sigma(\nu N \to \mu^- x)$ and $\overline{R} = \sigma(\overline{\nu} N \to \overline{\nu} x)/\sigma(\overline{\nu} N \to \mu^+ x)$. We assume $r = \sigma(\overline{\nu} N \to \mu^+ x)/\sigma(\nu N \to \mu^- x) = 0.4$. The explicit expressions for R and $\overline{R}$ in a quark-parton model calculation are[2,9]

$$R = \tfrac{1}{8}\{(G_A + G_V)^2 + (G_A^1 + G_V^1)^2 + r[(G_A - G_V)^2 + (G_A^1 - G_V^1)^2]\}, \tag{40a}$$

$$\overline{R} = \tfrac{1}{8}\{(G_A + G_V)^2 + (G_A^1 + G_V^1)^2 + r^{-1}[(G_A - G_V)^2 + (G_A^1 - G_V^1)^2]\}, \tag{40b}$$

where $G_A = g_A(Z_1) + g_A(Z_h)$, etc.

In Figs. 1 and 2 the calculated values of R and R are obtained by varying the ratio ϵ/δ. The specific value $\epsilon/\delta = 0$ corresponds to the case in which Z_1 is exactly decoupled from left-handed neutrinos. The 1974 Gargamelle data show that $R = 0.25 \pm 0.04$ and $\overline{R} = 0.39 \pm 0.06$.[2,16] These ranges are indicated by the shaded portions of Figs. 1 and 2. From the figures we see that there exist ranges of values for ϵ/δ that fit the measured values of R and $\overline{R}$ within experimental limits. If $\tan\phi = +\xi$, then $0.60 \le \epsilon/\delta \le 1.60$ in the unrenormalized theory and $-1.4 \le \epsilon/\delta \le 0.0$ in the renormalized theory. If $\tan\phi = -\xi$, we find $0.8 \le \epsilon/\delta \le 1.1$ in the bare theory and there are no consistent values in the renormalized theory. However, for intermediate amounts of renormalization the theory is consistent with either sign of $\tan\phi$.

The 1974 Gargamelle data also place limits on purely leptonic processes. We define $c_+ \equiv c_+(Z_1) + c_+(Z_h)$, $c_- \equiv c_-(Z_1) + c_-(Z_h)$. It has been shown that c_+ and c_- obey the constraints[2]

TABLE I. Neutrino-interaction amplitude coefficients for unrenormalized values of λ and ω ($\lambda^2 = \frac{3}{2}$, $\omega^2 = 1$). The upper (lower) element of a pair corresponds to $\tan\phi = +\xi$ ($-\xi$).

	Z_1	Z_h
c_+	$-\frac{1}{8}\frac{\epsilon}{\delta}$	0
c_-	$\frac{1}{2}\frac{\epsilon}{\delta}$	$\frac{1}{3}$
g_V	$\begin{Bmatrix}\frac{1}{12}\\ -\frac{1}{4}\end{Bmatrix}\times\frac{\epsilon}{\delta}$	$\begin{Bmatrix}\frac{1}{6}\\ \frac{1}{2}\end{Bmatrix}$
g_A	$\begin{Bmatrix}\frac{1}{6}\\ \frac{1}{2}\end{Bmatrix}\times\frac{\epsilon}{\delta}$	$\begin{Bmatrix}\frac{5}{6}\\ \frac{1}{2}\end{Bmatrix}$
g_V^1	$\begin{Bmatrix}\frac{1}{4}\\ -\frac{1}{12}\end{Bmatrix}\times\frac{\epsilon}{\delta}$	$\begin{Bmatrix}-\frac{1}{2}\\ -\frac{1}{6}\end{Bmatrix}$
g_A^1	$\begin{Bmatrix}-\frac{1}{6}\\ \frac{1}{6}\end{Bmatrix}\times\frac{\epsilon}{\delta}$	$\begin{Bmatrix}\frac{1}{6}\\ -\frac{1}{6}\end{Bmatrix}$

TABLE II. Neutrino-interaction amplitude coefficients for renormalized values of λ and ω ($\lambda^2 = \frac{1}{2}$, $\omega^2 = 3$). The upper (lower) element of a pair corresponds to $\tan\phi = +\xi$ ($-\xi$).

	Z_1	Z_h
c_+	$-\frac{1}{12}\frac{\epsilon}{\delta}$	$-\frac{1}{4}$
c_-	0	$\frac{1}{6}$
g_V	$\begin{Bmatrix}\frac{1}{12}\\ -\frac{1}{12}\end{Bmatrix}\times\frac{\epsilon}{\delta}$	$\begin{Bmatrix}\frac{7}{12}\\ \frac{3}{4}\end{Bmatrix}$
g_A	$\begin{Bmatrix}\frac{1}{12}\\ \frac{1}{4}\end{Bmatrix}\times\frac{\epsilon}{\delta}$	$\begin{Bmatrix}\frac{11}{12}\\ \frac{3}{4}\end{Bmatrix}$
g_V^1	$\begin{Bmatrix}\frac{5}{36}\\ -\frac{1}{36}\end{Bmatrix}\times\frac{\epsilon}{\delta}$	$\begin{Bmatrix}-\frac{1}{4}\\ -\frac{1}{2}\end{Bmatrix}$
g_A^1	$\begin{Bmatrix}-\frac{1}{12}\\ \frac{1}{12}\end{Bmatrix}\times\frac{\epsilon}{\delta}$	$\begin{Bmatrix}\frac{1}{12}\\ -\frac{1}{12}\end{Bmatrix}$

$$3.7c_+{}^2+1.0c_-{}^2\le 0.60\ ,$$
$$0.24c_+{}^2+1.0c_-{}^2\le 0.21\ , \qquad (41)$$
$$0.24c_+{}^2+1.0c_-{}^2\ge 0.020\ .$$

(These inequalities are obtained from the ellipses of Ref. 2, Fig. 1.) In Fig. 3, theoretical values for c_+ and c_- are obtained by varying ϵ/δ. We see once again that the curves intersect the shaded areas which delineate the empirical limits. We find $-7.8\le\epsilon/\delta\le 1.8$ in the renormalized theory and $-3.2\le\epsilon/\delta\le 2.4$ in the bare theory. These results are independent of the sign of $\tan\phi$.

The neutral-current data can be made consistent with the proposed two-current model if we use coupling-constant ratios intermediate between the unrenormalized and fully renormalized values. The analysis used here is essentially the same as that used to obtain limits on $\sin^2\theta_W$ in the Weinberg-Salam model.[17,18]

A model with a very light neutral IVB would have its most dramatic consequences in processes not involving neutrinos. We expect, for example, that such an IVB would manifest itself in parity-violation effects in $\Delta S=0$ nuclear-physics transitions. Theoretical estimates for these transitions based on charged-current interaction Hamiltonians have tended to be small when compared to experiment.[19] (For $n+p\to d+\gamma$ the most recent calculations are two orders of magnitude too small.) If the nucleon-nucleon-meson weak vertex in such processes has a contribution from Z_1 exchange, we expect this vertex to be enhanced by a factor $m_W{}^2/m_{Z_1}{}^2=\delta^{-1}$. Such enhancements would increase theoretical values for $\Delta S=0$ parity violation by at least an order of magnitude. Detailed calculations of these effects include a large number of assumptions which preclude their use as a definitive criterion for neutral-current models.[20] Clearly the definitive test for the existence of a very light neutral IVB is the direct detection of such a particle. We address this question in the following section.

V. DETECTION OF A VERY LIGHT AND/OR CHARM-CHANGING NEUTRAL IVB

In order to maintain consistency with neutral-current neutrino-interaction data the very light IVB Z_1 of the preceding section must have sup-

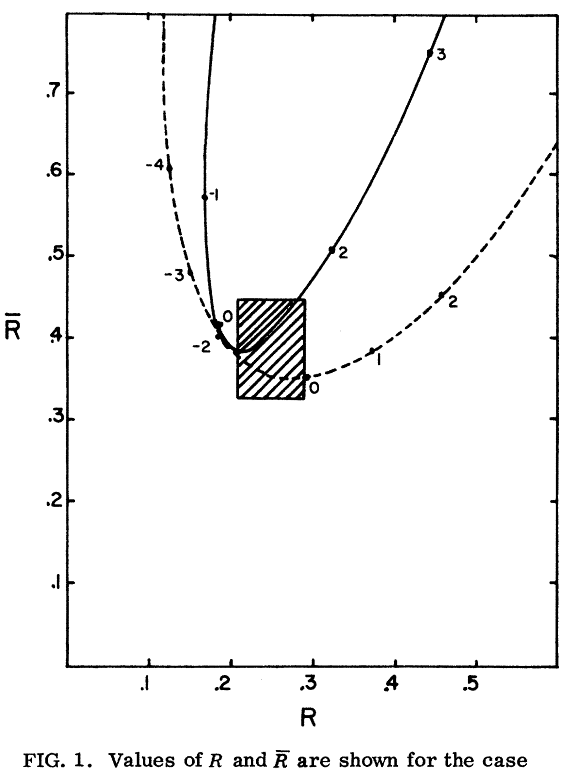

FIG. 1. Values of R and $\bar{R}$ are shown for the case $\tan\phi=+\xi$. The curves are generated by varying ϵ/δ. Some specific values of ϵ/δ are indicated on the curves. The shaded rectangle gives the experimental limits from the 1974 Gargamelle data, recently updated. The solid curve uses the unrenormalized values of g_V, g_A, g_V^1, g_A^1 from Table I; the dashed curve uses renormalized values from Table II.

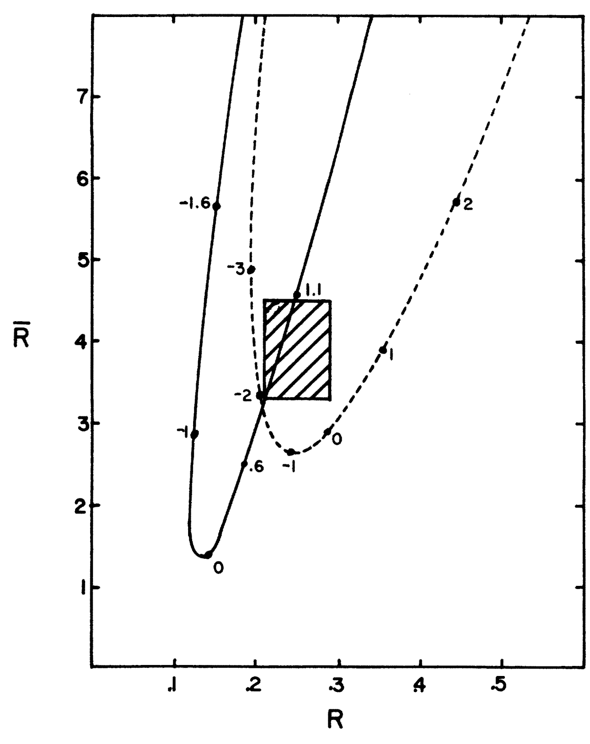

FIG. 2. The calculations generating Fig. 1 are repeated with the choice $\tan\phi=-\xi$. The solid curve uses unrenormalized coupling constants and the dashed curve uses renormalized values.

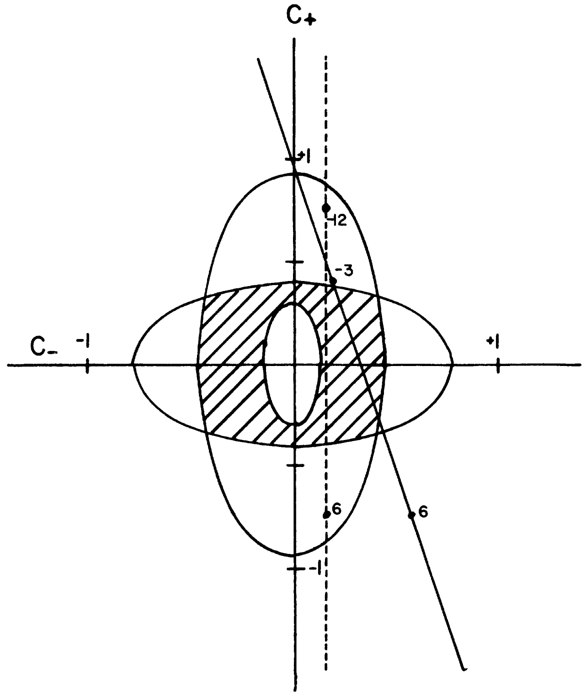

FIG. 3. Experimental results constrain c_+ and c_- to lie within the shaded area bounded by the three ellipses corresponding to Eq. (41). The solid straight line represents the values of c_+ and c_- obtained by varying ϵ/δ in the unrenormalized theory. The dashed line is calculated using renormalized coupling constants. Values of ϵ/δ are indicated along the lines. This analysis parallels the discussion by Wolfenstein in Ref. 2.

pressed coupling to left-handed neutrinos. To unambiguously detect such an IVB, we need to look beyond neutrino interactions and probe the unsuppressed couplings of Z_1. For example, when the e^+e^- center-of-mass energy is equal to the mass of $Z_1(m_{Z_1}{}^2=m_W{}^2\delta)$, the e^+e^- annihilation cross section should exhibit a resonance. In the discussion that follows we use the structure of Z_1 (27) and the mixing-angle constraint (29) to calculate the width and magnitude of such a resonance as a function of m_{Z_1}.

In terms of the fields N_L^1, N_R^1, N_R, and V_{15} the very light IVB has the structure

$$\begin{aligned}Z_1&=[\,\omega\lambda^2N_L^1-(1+\lambda^2)N_R^1-\omega\lambda V_{15}]\\&\quad\times\frac{\cos\phi}{[(1+\lambda^2)(1+\lambda^2+\omega^2\lambda^2)2]^{1/2}}+\sin\phi\, N_R\\&=c_LN_L^1+c_RN_R^1+c_{15}V_{15}+c_ZN_R\,,\end{aligned}\tag{42}$$

where $\tan\phi$ is given by (29). We choose the negative sign for $\tan\phi$, since that sign is consistent with the set of VEV's that generate the fermion mass matrix. The decay modes of Z_1 are into right-handed neutrinos (both ν_e and ν_μ), electrons, muons, and hadrons. A straightforward calculation yields the following widths[9]:

$$\Gamma_\nu=\sqrt{2}\;G_Fm_{Z_1}{}^3a_\nu{}^2/6\pi\,,\tag{43a}$$

$$\Gamma_e=\Gamma_\mu=\sqrt{2}\;G_Fm_{Z_1}{}^3a_e{}^2/12\pi\,,\tag{43b}$$

$$\Gamma_h=\sqrt{2}\;G_Fm_{Z_1}{}^3(v_3{}^2R_{3v}+a_3{}^2R_{3a}+v_0{}^2R_{0v}+a_0{}^2R_{0a})\,.\tag{43c}$$

The various parameters are given by

$$\begin{aligned}a_\nu{}^2&=\frac{1}{4\delta^2}\left[(c_L+\lambda c_{15})^2+\frac{2}{\omega^2}(c_R+\lambda\omega c_{15})^2+\frac{4}{\omega^2}c_Z{}^2\right],\\a_e{}^2&=\frac{1}{4\delta^2}\left[\left(c_L-\frac{c_R}{\omega}\right)^2+\left(c_L+\frac{c_R}{\omega}-2\lambda c_{15}\right)^2\right],\\v_3&=\frac{1}{2\delta}\left(2c_L+\frac{2c_R}{\omega}+\sqrt{2}\;\frac{c_Z}{\omega}\right),\\a_3&=\frac{1}{2\delta}\left(2c_L-\frac{2c_R}{\omega}-\sqrt{2}\;\frac{c_Z}{\omega}\right),\\v_0&=\frac{1}{2\delta}\left(-\tfrac{4}{3}\lambda c_{15}+\sqrt{2}\;\frac{c_Z}{\omega}\right),\\a_0&=\frac{1}{2\delta}\left(-\sqrt{2}\;\frac{c_Z}{\omega}\right).\end{aligned}\tag{44}$$

in (43c)[9]

$$R_{3v}=R_{3a}=\left[\frac{\sigma(e^+e^-\to\text{hadrons with }I=1)}{\sigma(e^+e^-\to\mu^+\mu^-)}\right]_{\text{off resonance}}.\tag{45}$$

The equality of R_{3v} and R_{3a} is a consequence of chiral symmetry.[9] We choose $R_{3v}=2$ based on the experimental values of R below the charm threshold. The value of $R_{0v}=R_{0a}$ is difficult to estimate.[9] At 7.4 GeV $R=R_{3v}+R_{0v}=5.9\pm0.9$.[21] This leads us to the choice $R_{0v}=R_{3v}=4$, a value consistent with an integral-quark-charge Pati-Salam model in which the colored gluons develop a mass ~1 GeV. The structure of the photon in an integral-charge model differs from that used in this paper; however, the above analysis is still correct if the gluons have masses small compared to that of Z_1.

The partial decay widths of Z_1 as a function of λ, ω, and m_{Z_1} are given in Table III.[22] Values are given for both bare and renormalized values of λ and ω. A large fraction of the total width arises from the decay into right-handed neutrinos. A Breit-Wigner resonance structure is used to estimate the contribution of Z_1 to the cross section for e^+e^- producing definite final state and to the total annihilation cross section. These results together with the integrated hadronic cross section Σ_h are also tabulated in Table III.

From Table III we see that if m_{Z_1} is between 7 and 15 GeV, we expect a peak of magnitude

TABLE III. Partial and total decay widths of Z_1, the maximum cross section $G^1_{Z_1}$, and the integrated cross section Σ_h for $e^+e^- \to Z_1 \to$ hadrons.

	Bare ($\lambda^2=\frac{3}{2}$, $\omega^2=1$)	Renormalized ($\lambda^2=\frac{1}{2}$, $\omega^2=3$)
$\Gamma_e=\Gamma_\mu$	$4.3\times10^{-4} m_{Z_1}$	$5.8\times10^{-4} m_{Z_1}$
Γ_h	$9.3\times10^{-3} m_{Z_1}$	$7.1\times10^{-3} m_{Z_1}$
Γ^L_ν	~ 0	~ 0
$\Gamma^R_{\nu_e}$	$2.7\times10^{-3} m_{Z_1}$	$2.3\times10^{-3} m_{Z_1}$
$\Gamma^R_{\nu_\mu}$	0	0
$\Gamma_{\rm tot}$	$1.3\times10^{-2} m_{Z_1}$	$1.0\times10^{-2} m_{Z_1}$
Σ_h	$\dfrac{7.3\times10^9\ \mathrm{MeV^2\,nb}}{m_{Z_1}}$	$\dfrac{9.3\times10^4\ \mathrm{MeV^2\,nb}}{m_{Z_1}}$
$\sigma^{\max}_{Z_1}$	$\dfrac{5.0\times10^{11}\ \mathrm{MeV^2\,nb}}{m_{Z_1}^2}$	$\dfrac{8.8\times10^{11}\ \mathrm{MeV^2\,nb}}{m_{Z_1}^2}$

10^3–10^4 nb in the e^+e^- total cross section. This is roughly the magnitude of the $\psi(3095)$ peak in the SPEAR data.[21] Hence, the Z_1 should easily be observed if it exists. The integrated hadronic cross section is of order 10^6 eV nb, several orders of magnitude larger than for $e^+e^- \to \psi(3095) \to$ hadrons. A strong, isolated resonance in e^+e^- annihilation with a width of 100 MeV or so would be a good candidate for Z_1. The large width for decay into leptons would be a characteristic feature of this particle.

In Sec. II we encountered a neutral, right-handed, charm-changing IVB. Since the current to which this IVB is coupled contains no charm-conserving pieces, this Z_{cc} would not contribute to nonleptonic $|\Delta c|=1$ decays of charmed particles. If Z_{cc} couples to a current proportional to K_{cc} in (9), it does not interact with electrons, muons, or left-handed neutrinos. If Z_{cc} exists, it would not contribute to any presently observed leptonic or semileptonic interaction. Its effects would be apparent in $|\Delta C|=2$ and $|\Delta C|=0$ nonleptonic processes. For example, Z_{cc} would lead to the mixing of the charmed-pseudoscalar-meson states D^0 and $\bar{D}^0$. Such mixing has been proposed as a possible explanation for the wrong-sign dimuon events in neutrino-nucleon scattering.[23] The charge current produces mixing in second order. The resulting mass difference Δm is too small. Too few $\bar{D}^0$ would become D^0's and decay to the wrong-sign muon. However, with Z_c the mixing becomes a first-order effect. The remainder of this section is devoted to an estimate of the mass difference Δm between D_1 and D_2, the CP eigenstates of the $D^0\bar{D}^0$ system. The calculation parallels similar estimates of $K^0\bar{K}^0$ mixing, and the results are a function of the mass of Z_{cc}.

The $D^0\bar{D}^0$ transition amplitude is calculated from the effective Hamiltonian

$$H^{cc}=\frac{G_F}{\sqrt{2}}\, m_W^2 D_F(x^2, m_{Z_{cc}}^2)[(\bar{c}u)_R+(\bar{u}c)_R]^2\ , \quad (46)$$

where $(\bar{c}u)_R=\bar{c}\gamma_\lambda(1+\gamma_5)u$. The term $[(\bar{c}u)_R(\bar{c}u)_R]$ is responsible for the $\bar{D}^0\to D^0$ transition. Thus we have

$$\langle D^0|H^{cc}|\bar{D}^0\rangle=\sqrt{2}\,G_F m_W^2\int dx\, D_F(x, m_{Z_{cc}}^2)\langle D^0|T\{[\bar{c}(x)u(x)]_R[\bar{c}(0)u(0)]_R\}|\bar{D}^0\rangle\ . \quad (47)$$

The time-ordered product is evaluated by means of a Wilson expansion.[24] The dominant term is the four-quark operator of dimension six. The renormalization group equation is used to estimate strong-interaction effects.[25] Since there are no dramatic enhancements or suppressions, we ignore renormalization effects in our estimate, and (47) reduces to

$$\langle D^0|H^{cc}|\bar{D}^0\rangle=\sqrt{2}\,G_F\frac{m_W^2}{m_{Z_{cc}}^2}\langle D^0|:[\bar{c}(0)u(0)]_R[\bar{c}(0)u(0)]_R:|\bar{D}^0\rangle\ . \quad (48)$$

The approach of Lee and Gaillard[26] is used to evaluate (48). A complete set of intermediate states is replaced by the vacuum state yielding the final result for the $\bar{D}^0D^0$ transition amplitude.

$$\langle D^0|H^{cc}|\bar{D}^0\rangle=\frac{32}{3}\frac{G_F}{\sqrt{2}}\frac{m_W^2}{m_{Z_{cc}}^2} f_D^2 m_{D^0}^2\ . \quad (49)$$

Since D^0 and $\bar{D}^0$ are degenerate before mixing, the diagonal matrix elements are

$$\langle D^0|H|D^0\rangle=\langle\bar{D}^0|H|\bar{D}^0\rangle=m_{D^0}^2/2\ . \quad (50)$$

Together (49) and (50) form the $D^0, \bar{D}^0$ mass matrix. The CP eigenstates $D_1=(D^0+\bar{D}^0)/\sqrt{2}$ and $D_2=(D^0-\bar{D}^0)/\sqrt{2}$ diagonalize the matrix. The D_1, D_2 mass difference is found to be

$$\Delta m=\frac{64}{3}\frac{G_F}{\sqrt{2}}\frac{m_W^2}{m_{Z_{cc}}^2} f_D^2 m_D\ . \quad (51)$$

If $f_D=f_\pi$, from SU(4) symmetry, and $m_{D^0}=2$ GeV,[13]

$$\Delta m=8\left(\frac{m_W^2}{m_{Z_{cc}}^2}\right)^2 \text{keV}\ . \quad (52)$$

Kingsley, Treiman, Wilczek, and Zee[23] argued that $D^0\bar{D}^0$ mixing produces sufficient wrong-sign dimuon events if $\Delta m \geq \lambda$, where λ is the mean decay rate of D^0 and $\bar{D}^0$. This rate has been estimated to be[27]

$$\lambda = \frac{(G_F \cot\theta_C)^2 m_c^5}{192\pi^3} \approx 5 \times 10^{-5} \text{ keV} , \tag{53}$$

where m_c is the mass of the charmed quark ($m_c \approx 1.5$ GeV) and θ_C is the Cabibbo angle. Hence, $m_{Z_{cc}}$ may be as large as $10^3\, m_W$ and the mixing mechanism could still be responsible for the wrong-sign dimuon events. The basic point is that with Z_{cc}, this process is first order in the weak interactions.

VI. CONCLUSIONS

We have shown that very light neutral gauge bosons are not ruled out by present experimental results on neutral-current weak interactions. Unified gauge theories of strong, electromagnetic, and weak interactions contain many neutral gauge fields. Using the generalized Pati-Salam model as a prototype, we have shown that there are definite restrictions on the nature of possible light neutral intermediate vector bosons. However, these restrictions permit a large class of models with more than one neutral current. The choice of a definite symmetry-breaking mechanism leads to a unique model which contains a normal vector boson (mass of order m_W) and a very light neutral particle of unspecified mass. This model, with essentially one free parameter, is consistent with neutral-current phenomenology. It needs to be tested in detailed calculations before we propose it as a serious alternative to the Weinberg-Salam model.

The very light neutral particle is necessarily an elusive particle to detect experimentally. Otherwise it would have been found by now. The most direct evidence for its existence would be an isolated resonance in e^+e^- annihilation with an unusually large decay rate into leptons. Our investigations have also uncovered the possibility that there might exist a very light vector boson which couples to purely charm-changing current. Such a particle would be very hard to detect experimentally, since it does not couple to charged leptons—at least in the generalized Pati-Salam model. Experiment should soon tell us whether $D^0\bar{D}^0$ mixing, if it exists, is a first-order effect. If it is not a first-order effect, this charm-changing possibility can be ruled out.

Finally we note that our analysis was based upon a particular unified model. A similar analysis could be carried out in any other model. We expect that similar results would emerge.

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