New supersymmetric left-right gauge model: Higgs-boson structure and neutral-current analysis

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If the particle content of a supersymmetric SU(3)×SU(2)×SU(2)×U(1) gauge model is extended to include new quarks and leptons as given by the **27** representation of E_6, an unconventional left-right assignment is possible, resulting in a number of desirable features, one of which is a simpler Higgs-boson structure. Also, a conserved multiplicative quantum number can be defined as a generalization of the usual R parity in most models of supersymmetry. This prevents W_L-W_R mixing and makes possible a massless Dirac neutrino whose right-handed component is effectively inert. From neutral-current constraints based on present experimental data, lower bounds of about 180 and 210 GeV are obtained for the second W and Z bosons of this model, respectively.

I. INTRODUCTION

The standard SU(3)×SU(2)×U(1) gauge model is remarkably successful so far in describing all of particle physics, but it may not be a complete theory, and other phenomena outside the purview of the standard model may become observable at energies below a TeV. One such possible extension is supersymmetry which helps to stabilize the electroweak mass scale against large radiative corrections. Another is left-right symmetry which is a natural extension of the standard gauge group. Each can be considered separately or both can be combined to form a single theory, but there does not seem to be any qualitatively new physics in their combination which is not already present in their separate manifestations.

Then came superstring theory[1] and the discovery of the $E_8 \times E_8$ heterotic string,[2] whose low-energy manifestation[3] may well consist of matter supermultiplets belonging to the **27** representation of E_6 interacting with one another as well as with gauge bosons and fermions belonging to a subgroup G of E_6. The most popular choice of G appears to be SU(3)×SU(2)×U(1)×U(1), and many studies have been made[4-7] in its name. On the other hand, if G turns out to be SU(3)×SU(2)×SU(2)×U(1), then a new supersymmetric left-right model is possible,[8] with qualitatively rather different physics beyond the standard model than in conventional extensions.

In Sec. II the model is described in some detail. In Sec. III the part of the Higgs potential relevant to the spontaneous breaking of the gauge symmetry is analyzed. In particular, it is shown that an upper bound of $\sqrt{2}\, M_W$ exists on the mass of one of the neutral physical Higgs bosons. In Sec. IV present experimental data are compared against the effective neutral-current interactions of this model, and lower bounds on the masses of the second W and Z bosons are obtained. Finally, in Sec. V there are some concluding remarks.

II. DESCRIPTION OF MODEL

The fundamental **27** representation of E_6 can be decomposed according to its E_5 [=SO(10)] and E_4 [=SU(5)] contents, namely,

$$\mathbf{27}=(16,\bar{5})+(16,10)+(16,1)+(10,\bar{5})+(10,5)+(1,1)\ . \tag{2.1}$$

At the E_6 level, the only allowed Yukawa coupling is for the product $\mathbf{27}\times\mathbf{27}\times\mathbf{27}$. If one assumes that all couplings forbidden at the E_6 level are also forbidden at a lower level,[5] then all Yukawa interactions must be contained in the terms

$$\begin{aligned}&(16,\bar{5})(16,10)(10,\bar{5})\ ,\\&(16,\bar{5})(16,1)(10,5)\ ,\\&(16,10)(16,10)(10,5)\ ,\\&(10,\bar{5})(10,5)(1,1)\ .\end{aligned} \tag{2.2}$$

Two things are immediately apparent from the above. First, a multiplicative quantum number can be assigned as follows: even (odd) for the $(16,\bar{5})$,(16,10),(16,1) fermions (bosons) and the $(10,\bar{5})$,(10,5),(1,1) bosons (fermions). Second, if $(16,\bar{5})$ is interchanged with $(10,\bar{5})$ and (16,1) with (1,1), the allowed Yukawa interactions remain the same. Hence there are two possibilities and they are summarized in Table I.

The standard-model quarks and leptons can be identified with the $(16,\bar{5})$ and (16,10) fermions of option (A) or the $(10,\bar{5})$ and (16,10) fermions of option (B). The standard-model Higgs bosons are then contained in the $(10,\bar{5})$ and (10,5) bosons of option (A) or the $(16,\bar{5})$ and (10,5) bosons of option (B). Therefore, the aforementioned multiplicative quantum number is nothing but an extension of the usual R parity in most models of super-

TABLE I. Possible assignments of R parity.

Option (A)	Fermions	Bosons
(16,$\bar{5}$),(16,10),(16,1)	+	−
(10,$\bar{5}$),(10,5),(1,1)	−	+
Option (B)	**Fermions**	**Bosons**
(10,$\bar{5}$),(16,10),(1,1)	+	−
(16,$\bar{5}$),(10,5),(16,1)	−	+

symmetry, which is conserved as long as only those neutral scalar bosons with even R parity have nonzero vacuum expectation values. In fact, the two options are equivalent[8] if the low-energy gauge group is the standard SU(3)×SU(2)×U(1) or its U(1) extension. However, if it is the left-right extension, the two options are not equivalent[8] and in this paper, we will only be concerned with option (B). The particle assignments under SU(3)×SU(2)×SU(2)×U(1) are

$$(u,d)_L:(3,2,1,\tfrac{1}{6}) , \tag{2.3}$$

$$d_L^c:(\bar{3},1,1,\tfrac{1}{3}) , \tag{2.4}$$

$$(h^c,u^c)_L:(\bar{3},1,2,-\tfrac{1}{6}) , \tag{2.5}$$

$$h_L:(3,1,1,-\tfrac{1}{3}) , \tag{2.6}$$

$$\begin{bmatrix} \nu_e & E^c \\ e & N_E^c \end{bmatrix}_L :(1,2,2,0) , \tag{2.7}$$

$$(e^c,n)_L:(1,1,2,\tfrac{1}{2}) , \tag{2.8}$$

$$(\nu_E,E)_L:(1,2,1,-\tfrac{1}{2}) , \tag{2.9}$$

$$N_L^c:(1,1,1,0) , \tag{2.10}$$

where the superscript c denotes the charge-conjugate state; hence ψ_L^c can be interpreted as ψ_R. In the above, the R parity is even for u,d,ν_e,e, and odd for h,E,ν_E,N_E^c,n, with the opposite assignments for their supersymmetric scalar partners. In particular, the scalar bosons $\tilde{\nu}_E$, $\tilde{N}_E^c$, and $\tilde{n}$ all have even R parity; hence, the spontaneous breaking of the gauge symmetry through their vacuum expectation values does not disturb its conservation. Fermion masses are generated as follows: m_u comes from $\langle \tilde{N}_E^c \rangle$; m_d and m_e come from $\langle \tilde{\nu}_E \rangle$; m_h and m_E come from $\langle \tilde{n} \rangle$; ν_e combines with N^c to form a massive Dirac neutrino through $\langle \tilde{N}_E^c \rangle$ if the (10,$\bar{5}$)(10,5)(1,1) term is present; and ν_E, N_E^c, and n form a 3×3 mass matrix through $\langle \tilde{n} \rangle$, $\langle \tilde{\nu}_E \rangle$, and $\langle \tilde{N}_E^c \rangle$. If the (10,$\bar{5}$)(10,5)(1,1) term is absent, then the R parity of N^c has to be odd instead of even, resulting in a massless Dirac neutrino. The role of N^c here is rather analogous to that of a possible right-handed neutrino ν_R in the standard SU(3)×SU(2)×U(1) model. It is inert with respect to the gauge interactions, and it is only needed if the neutrino has a Dirac mass. However, if we go beyond SU(3)×SU(2)×SU(2)×U(1), N^c will have nontrivial interactions. In fact, flux breaking[9] of E_6 can be used[10] to take it down to SU(3)×SU(2)×SU(2)×U(1)×U(1), but to break U(1)×U(1) to just U(1), $\tilde{N}^c$ must develop a large vacuum expectation value. Therefore, if R parity is to be maintained at this higher level, N^c should be odd. This indicates that there may be a connection between the smallness of the neutrino mass, if it exists, and the nonconservation of R parity at the mass scale $\langle \tilde{N}^c \rangle$.

As for the gauge bosons and fermions, R parity is even and odd, respectively, *except* for the second W boson which links particles of opposite R parity; hence, its own R parity must be odd and that of its supersymmetric partner even. Consequently, there is no mixing between the two W bosons of this model, and that is why a massless Dirac neutrino is possible.[8] Phenomenologically, the astrophysical limit on the effective total number of neutrinos, namely,[11] $N_\nu < 4$, is not violated by the existence of N^c's because they are inert. The mass of the second W boson is not constrained by the K_L-K_S mass difference[12] because W_2 does not couple to d and s quarks; it is also not constrained by polarized μ^+ decay[13] because n is presumably heavy. Actually, there is a potential problem here with m_n. Because of left-right symmetry, the 3×3 mass matrix spanning ν_E, N_E^c, and n is of the form

$$\begin{bmatrix} 0 & m_E & m' \\ m_E & 0 & m_e \\ m' & m_e & 0 \end{bmatrix} , \tag{2.11}$$

where $m' = m_E \langle \tilde{N}_E^c \rangle / \langle \tilde{n} \rangle$. This means that m_n is roughly given by $2m_e \langle \tilde{N}_E^c \rangle / \langle \tilde{n} \rangle$ which is less than 1 MeV. On the other hand, a diagonal mass for n is possible through radiative corrections. For example, the diagram of Fig. 1 is finite if there are soft-supersymmetry-breaking terms of the form $\tilde{E}\tilde{E}^c\tilde{n}$. A rough estimate for m_n is then $10^{-2} m_E$ which may well be a few GeV. Note that R parity is not broken by such a mass term.

Since W_2 has odd R parity, it is produced only in pairs or in association with another particle of odd R parity. At present, no firm experimental lower limit on M_{W_2} is known. However, the Higgs-boson structure relates M_{W_2} to M_{Z_1} and M_{Z_2}, so that neutral-current constraints can be used to limit both M_{Z_2} and M_{W_2}. In the next section the Higgs-boson structure is worked out; then in Sec. IV the effective neutral-current interactions are compared against the data.

III. HIGGS-BOSON STRUCTURE

In conventional left-right-symmetric gauge models, the minimal Higgs sector usually consists of the multiplets $\phi(1,2,2,0)$, $\Delta_L(1,3,1,1)$, and $\Delta_R(1,1,3,1)$. The complex

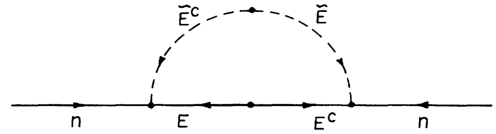

FIG. 1. Radiative mass term for n in the presence of soft-supersymmetry breaking ($\tilde{E}\tilde{E}^c\tilde{n}$) and spontaneous gauge-symmetry breaking ($\langle \tilde{n} \rangle \neq 0$).

triplet Δ_R is used to break $SU(2)_R$ and supply a large Majorana mass to the right-handed neutrino. Together with a small Dirac mass coming from ϕ, the seesaw mechanism[14] then assures a very small Majorana mass for the observed left-handed neutrino. However, there are no triplets in the **27** representation of E_6. So we must consider

$$\phi(1,2,2,0)=\begin{bmatrix}\phi_{11} & \phi_{12}\\ \phi_{21} & \phi_{22}\end{bmatrix}=\begin{bmatrix}\tilde{E}^c & \tilde{N}_E^c\\ \tilde{\nu}_e & \tilde{e}\end{bmatrix}, \tag{3.1}$$

$$H_L(1,2,1,-\tfrac{1}{2})=\begin{bmatrix}H_{L1}\\ H_{L2}\end{bmatrix}=\begin{bmatrix}\tilde{\nu}_E\\ \tilde{E}\end{bmatrix}, \tag{3.2}$$

and

$$H_R(1,1,2,\tfrac{1}{2})=\begin{bmatrix}H_{R1}\\ H_{R2}\end{bmatrix}=\begin{bmatrix}\tilde{e}^c\\ \tilde{n}\end{bmatrix}. \tag{3.3}$$

The nonzero vacuum expectation values are

$$\langle \tilde{N}_E^c\rangle = v\ , \tag{3.4}$$

$$\langle \tilde{\nu}_E\rangle = v_L\ , \tag{3.5}$$

$$\langle \tilde{n}\rangle = v_R\ . \tag{3.6}$$

The mass of the neutrino comes from the term $\nu_e N^c\langle \tilde{N}_E^c\rangle$ which may or may not be zero, but if it is zero, then R parity will ensure that it remains zero, at least at the level of $SU(3)\times SU(2)\times SU(2)\times U(1)$. Note that in conventional left-right models, both neutral components of $\phi(1,2,2,0)$ must acquire nonzero vacuum expectation values. Here, it is possible[8] to maintain $\langle\tilde{\nu}_e\rangle=0$, again because of R parity.

As discussed before, there is no mixing between the two W bosons of this model. Their masses are given by

$$M_{W_1}^2=\tfrac{1}{2}g_L^2(v^2+v_L^2)\ , \tag{3.7}$$

$$M_{W_2}^2=\tfrac{1}{2}g_R^2(v^2+v_R^2)\ , \tag{3.8}$$

Since it is assumed that v_R is only a few times greater than v_L and v, g_R should only differ very slightly from g_L. For simplicity, $g_R=g_L$ will be used from now on. Let

$$x=\frac{e^2}{g_L^2}=\sin^2\theta_W\ , \tag{3.9}$$

$$y=\frac{v^2}{v^2+v_L^2}\ , \tag{3.10}$$

then the 2×2 mass matrix spanning the standard Z boson and the extra D boson is given by

$$\mathcal{M}_{Z,D}^2=\begin{bmatrix}(1-x)^{-1}M_{W_1}^2 & \left[\dfrac{x}{1-x}-y\right]\dfrac{M_{W_1}^2}{\sqrt{1-2x}}\\ \left[\dfrac{x}{1-x}-y\right]\dfrac{M_{W_1}^2}{\sqrt{1-2x}} & \left[\dfrac{1-x}{1-2x}\right]M_{W_2}^2+\left[\dfrac{x}{1-x}-2y\right]\dfrac{xM_{W_1}^2}{1-2x}\end{bmatrix}. \tag{3.11}$$

If $M_{W_2}^2 \gg M_{W_1}^2$, then $M_{Z_1}=M_{W_1}/\cos\theta_W$ as expected. Note that there is in general some mixing between Z and D, but the effect on the mass eigenstate Z_1 is usually quite small. To a first approximation,

$$M_{Z_1}^2\simeq\frac{M_{W_1}^2}{1-x}\left[1-\left[\frac{x}{1-x}-y\right]^2\frac{M_{W_1}^2}{M_{W_2}^2}\right]. \tag{3.12}$$

Consider now the most general superpotential involving ϕ, H_L, and H_R as supermultiplets:

$$W=\lambda\phi_{ij}H_{L\alpha}H_{R\beta}\epsilon_{i\beta}\epsilon_{j\alpha}+\mu\phi_{ij}\phi_{kl}\epsilon_{ik}\epsilon_{jl}\ , \tag{3.13}$$

where the parameters λ and μ are real because of left-right symmetry. The scalar potential is then

$$\begin{aligned}V=&4\mu^2\phi_{ij}^*\phi_{ij}+2\mu\lambda(\phi_{ij}^*H_{Lj}H_{Ri}+\text{H.c.})\\ &+\lambda^2[\,|\phi_{12}H_{R2}-\phi_{22}H_{R1}|^2+|\phi_{11}H_{R2}-\phi_{21}H_{R1}|^2+|\phi_{21}H_{L2}-\phi_{22}H_{L1}|^2\\ &\quad+|\phi_{11}H_{L2}-\phi_{12}H_{L1}|^2+(H_{Li}^*H_{Li})(H_{Rj}^*H_{Rj})]\\ &+\frac{e^2}{8x}(4\,|H_{L1}^*H_{L2}+\phi_{\alpha1}^*\phi_{\alpha2}|^2+4\,|H_{R1}^*H_{R2}+\phi_{1\alpha}^*\phi_{2\alpha}|^2+|H_{L1}^*H_{L1}^*-H_{L2}^*H_{L2}+\phi_{\alpha1}^*\phi_{\alpha1}-\phi_{\alpha2}^*\phi_{\alpha2}|^2\\ &\quad+|H_{R1}^*H_{R1}-H_{R2}^*H_{R2}+\phi_{1\alpha}^*\phi_{1\alpha}-\phi_{2\alpha}^*\phi_{2\alpha}|^2)+\frac{e^2}{8(1-2x)}(H_{Li}^*H_{Li}-H_{Ri}^*H_{Ri})^2\ ,\end{aligned} \tag{3.14}$$

where all the terms arising from the gauge interactions have also been included. To this we add the most general soft-supersymmetric-breaking terms

$$\begin{aligned}V'=&\tilde{m}_1^2\phi_{ij}^*\phi_{ij}+m_2^2(\phi_{ij}\phi_{kl}\epsilon_{ik}\epsilon_{jl}+\text{H.c.})+m_3^2(H_{Li}^*H_{Li}+H_{Ri}^*H_{Ri})\\ &+\lambda A(\phi_{ij}H_{L\alpha}H_{R\beta}\epsilon_{i\beta}\epsilon_{j\alpha}+\text{H.c.})+\tilde{B}(\phi_{ij}^*H_{Lj}H_{Ri}+\text{H.c.})\ .\end{aligned} \tag{3.15}$$

The resulting stationary conditions are

$$v\left[m_1^2+\lambda^2(v_L^2+v_R^2)+\frac{e}{4x}(2v^2-v_L^2-v_R^2)\right]-\lambda A v_L v_R=0\ , \qquad (3.16)$$

$$v_L\left[m_3^2+\lambda^2(v^2+v_R^2)+\frac{e^2}{4}\left[\frac{v_L^2-v^2}{x}+\frac{v_L^2-v_R^2}{1-2x}\right]\right]-\lambda A v v_R=0\ , \qquad (3.17)$$

$$v_R\left[m_3^2+\lambda^2(v^2+v_L^2)+\frac{e^2}{4}\left[\frac{v_R^2-v^2}{x}+\frac{v_R^2-v_L^2}{1-2x}\right]\right]-\lambda A v v_L=0\ , \qquad (3.18)$$

and

$$Bv_Lv_R-2m_2^2v=0\ , \qquad (3.19)$$

where we have defined

$$m_1^2=\tilde{m}_1^2+4\mu^2\ , \qquad (3.20)$$

$$B=\tilde{B}+2\mu\lambda\ . \qquad (3.21)$$

Now it appears that with only three variables, i.e., v, v_L, and v_R, it is not possible to satisfy all four equations of constraint in general. However, there is no cause for alarm because R-parity invariance requires μ, m_2^2, and $\tilde{B}$ to be all zero, so that Eq. (3.19) is eliminated trivially. Without R-parity invariance, $\phi_{21}=\tilde{\nu}_e$ must develop a nonzero vacuum expectation value which would then be the fourth variable. The absence of the μ term in the superpotential of Eq. (3.13) is of course natural within the context of E_6 because 27×27 does not make a singlet.

Consider first the charged scalar bosons. If we define

$$\tan\theta_{L,R}=\frac{v}{v_{L,R}}\ , \qquad (3.22)$$

then

$$G_L^-=\cos\theta_L H_{L2}-\sin\theta_L\phi_{11}^*\ , \qquad (3.23)$$

$$G_R^-=\cos\theta_R H_{R1}^*-\sin\theta_R\phi_{22}\ , \qquad (3.24)$$

and their charge conjugates are the would-be Goldstone bosons absorbed by the $W_{1,2}$ bosons, while the two corresponding orthogonal combinations have physical masses given by

$$m_{H_1^-}^2=\left[\frac{\lambda A v_R}{v v_L}+\frac{e^2}{2x}-\lambda^2\right](v^2+v_L^2)\ , \qquad (3.25)$$

$$m_{H_2^-}^2=\left[\frac{\lambda A v_L}{v v_R}+\frac{e^2}{2x}-\lambda^2\right](v^2+v_R^2)\ . \qquad (3.26)$$

Note that H_1^- and H_2^- have opposite R parity, so they do not mix.

Consider next the imaginary components of the four neutral scalar fields. Two are eaten up by the $Z_{1,2}$ bosons and we are left with again two physical scalar bosons of opposite R parity. Let

$$\psi_1^0=\sqrt{2}\,\mathrm{Im}\phi_{21}\ , \qquad (3.27)$$

$$\psi_2^0=\frac{\sqrt{2}\,\mathrm{Im}(\phi_{12}+\tan\theta_L H_{L1}+\tan\theta_R H_{R2})}{(1+\tan^2\theta_L+\tan^2\theta_R)^{1/2}}\ . \qquad (3.28)$$

Then their masses are given by

$$m_{\psi_1^0}^2=\frac{\lambda A v_L v_R}{v}+\frac{e^2}{2x}(v_L^2+v_R^2-2v^2)-\lambda^2(v_L^2+v_R^2)\ , \qquad (3.29)$$

$$m_{\psi_2^0}^2=\lambda A\left[\frac{v_L v_R}{v}+\frac{v v_R}{v_L}+\frac{v v_L}{v_R}\right]\ . \qquad (3.30)$$

Consider finally the real components of the four neutral scalar fields, all of which are physical. The one with odd R parity is $H_1^0=\sqrt{2}\,\mathrm{Re}\phi_{21}$, and it has a mass equal to that of ψ_1^0. Hence, ϕ_{21} remains a complex neutral scalar boson. The other three, i.e., $\sqrt{2}\,\mathrm{Re}\phi_{12}$, $\sqrt{2}\,\mathrm{Re}H_{L1}$, and $\sqrt{2}\,\mathrm{Re}H_{R2}$, have even R parity, and their mass matrix is given by

$$\mathcal{M}_H^2=\begin{bmatrix} \dfrac{\lambda A v_L v_R}{v}+\dfrac{e^2v^2}{x} & -\lambda A v_R+\left[2\lambda^2-\dfrac{e^2}{2x}\right]v v_L & -\lambda A v_L+\left[2\lambda^2-\dfrac{e^2}{2x}\right]v v_R \\ -\lambda A v_R+\left[2\lambda^2-\dfrac{e^2}{2x}\right]v v_L & \dfrac{\lambda A v v_R}{v_L}+\dfrac{e^2}{2x}\left[\dfrac{1-x}{1-2x}\right]v_L^2 & -\lambda A v+\left[2\lambda^2-\dfrac{e^2}{2-4x}\right]v_L v_R \\ -\lambda A v_L+\left[2\lambda^2-\dfrac{e^2}{2x}\right]v v_R & -\lambda A v+\left[2\lambda^2-\dfrac{e^2}{2-4x}\right]v_L v_R & \dfrac{\lambda A v v_L}{v_R}+\dfrac{e^2}{2x}\left[\dfrac{1-x}{1-2x}\right]v_R^2 \end{bmatrix}\ . \qquad (3.31)$$

Complicated as it might seem, one of the eigenvalues of this matrix is actually bounded from above by $2M_{W_1}{}^2$. This can be seen as follows. Let

$$\tan\beta = \frac{v_L v_R (v_R{}^2 + v_L{}^2)^{1/2}}{v(v_R{}^2 - v_L{}^2)} , \tag{3.32}$$

$$\tan\gamma = \frac{v_L}{v_R} , \tag{3.33}$$

and

$$H_2^0 = \sqrt{2}\mathrm{Re}[\cos\beta\phi_{12} + \sin\beta(\cos\gamma H_{L1} - \sin\gamma H_{R2})] . \tag{3.34}$$

Then using the identity

$$\lambda A v = \left[\lambda^2 - \frac{e^2}{4x(1-2x)}\right] v_L v_R , \tag{3.35}$$

the diagonal mass term for H_2^0 is calculated to be

$$m_{H_2^0}{}^2 = \left[\frac{e^2 v^2}{x}\right]\cos^2\beta \le 2M_{W_1}{}^2 . \tag{3.36}$$

Now H_2^0 is not necessarily a mass eigenstate, but any mixing into the other states can only reduce its mass. Hence, Eq. (3.36) represents an absolute upper bound of $\sqrt{2}M_{W_1}$ on one of the neutral physical Higgs bosons of this model. This bound is saturated only in the limit $v_L/v = 0$, for which it is equal to 116 GeV. If $v_L/v = 1$, then it is reduced by half to only 58 GeV. Finally we have checked that all mass eigenvalues can be positive, so that v, v_L, v_R do indeed correspond to a minimum of the Higgs potential.

IV. NEUTRAL-CURRENT ANALYSIS

The interaction Hamiltonian of this model for neutral currents is given by[15]

$$H_{\mathrm{int}} = \frac{eZ}{\sqrt{x(1-x)}}(J_3^{(1)} - xJ_{\mathrm{em}}) + \frac{eD}{\sqrt{x(1-x)(1-2x)}}[xJ_3^{(1)} + (1-x)J_3^{(2)} - xJ_{\mathrm{em}}] , \tag{4.1}$$

where $J_3^{(1)}, J_3^{(2)}$ are the neutral currents associated with the $SU(2)_L, SU(2)_R$ gauge groups, and J_{em} is the electromagnetic current. As far as the known quarks and leptons are concerned, their couplings to the standard Z boson are as usual given by

$$\begin{aligned} J_3^{(1)} - xJ_{\mathrm{em}} = {} & \tfrac{1}{2}\bar{\nu}_e\gamma\left[\frac{1-\gamma_5}{2}\right]\nu_e \\ & + \tfrac{1}{2}(-\tfrac{1}{2}+2x)\bar{e}\gamma e - \tfrac{1}{2}(-\tfrac{1}{2})\bar{e}\gamma\gamma_5 e \\ & + \tfrac{1}{2}(\tfrac{1}{2}-\tfrac{4}{3}x)\bar{u}\gamma u - \tfrac{1}{2}(\tfrac{1}{2})\bar{u}\gamma\gamma_5 u \\ & + \tfrac{1}{2}(-\tfrac{1}{2}+\tfrac{2}{3}x)\bar{d}\gamma d - \tfrac{1}{2}(-\tfrac{1}{2})\bar{d}\gamma\gamma_5 d , \end{aligned} \tag{4.2}$$

where the Lorentz index of the Dirac γ matrix has been suppressed for notational simplicity. The corresponding couplings to the D boson are new and different from those of the conventional left-right model; they are given by

$$\begin{aligned} xJ_3^{(1)} + (1-x)J_3^{(2)} - xJ_{\mathrm{em}} = {} & -\tfrac{1}{2}(1-2x)\bar{\nu}_e\gamma\left[\frac{1-\gamma_5}{2}\right]\nu_e + \tfrac{1}{2}(-1+\tfrac{5}{2}x)\bar{e}\gamma e - \tfrac{1}{2}(-\tfrac{1}{2}x)\bar{e}\gamma\gamma_5 e \\ & + \tfrac{1}{2}(\tfrac{1}{2}-\tfrac{4}{3}x)\bar{u}\gamma u - \tfrac{1}{2}(-\tfrac{1}{2}+x)\bar{u}\gamma\gamma_5 u + \tfrac{1}{2}(\tfrac{1}{6}x)\bar{d}\gamma d - \tfrac{1}{2}(-\tfrac{1}{2}x)\bar{d}\gamma\gamma_5 d . \end{aligned} \tag{4.3}$$

Even though the mass matrix for Z and D is not diagonal in general, all one needs is its inverse to obtain the effective current-current interactions of this model at low energies. Let

$$\xi = \frac{M_{W_1}{}^2}{M_{W_2}{}^2} . \tag{4.4}$$

Then for deep-inelastic neutrino and antineutrino scattering off nucleons, the relevant quantities to be compared against data are

$$\epsilon_L^u = \frac{1}{1-\xi y^2}[\tfrac{1}{2} - \tfrac{2}{3}x + \xi(\tfrac{2}{3}x - \tfrac{1}{2}y)] , \tag{4.5}$$

$$\epsilon_R^u = \frac{1}{1-\xi y^2}[-\tfrac{2}{3}x + \xi(-\tfrac{1}{2} + \tfrac{2}{3}x + \tfrac{1}{2}y)] , \tag{4.6}$$

$$\epsilon_L^d = \frac{1}{1-\xi y^2}[-\tfrac{1}{2} + \tfrac{1}{3}x + \xi(-\tfrac{1}{3}x + \tfrac{1}{2}y)] , \tag{4.7}$$

$$\epsilon_R^d = \frac{1}{1-\xi y^2}[\tfrac{1}{3}x + \xi(-\tfrac{1}{3}x)] , \tag{4.8}$$

where $x = \sin^2\theta_W$ and $y = \sin^2\theta_L$ as given by Eqs. (3.9), (3.10), and (3.22). The standard-model results are of course recovered in the limit $\xi \to 0$. The numerical values of the above quantities have been derived from experimental data after radiative corrections. We use the recent compilation of Fogli[16] and fit against those values, taking into account the error-correlation matrix. In Fig. 2, we show the one-standard-deviation allowed region in the (x, ξ) plane for various given values of y, based on Fogli's numbers plus the constraint due to the measured values of M_{W_1} and M_{Z_1}, namely[17]

$$1 - \frac{M_{W_1}{}^2}{M_{Z_1}{}^2} = 0.218 \pm 0.022 , \tag{4.9}$$

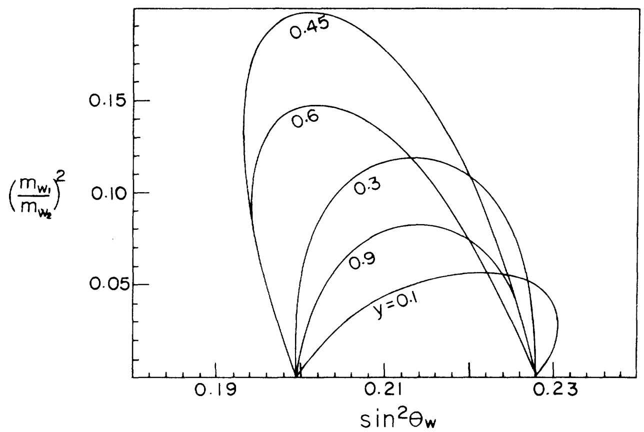

FIG. 2. One-standard-deviation allowed region in $x=\sin^2\theta_W$ and $\xi = M_{W_1}{}^2/M_{W_2}{}^2$ for various values of $y=\sin^2\theta_L$, based on neutrino deep-inelastic data and $M_{W_1}{}^2/M_{Z_1}{}^2$.

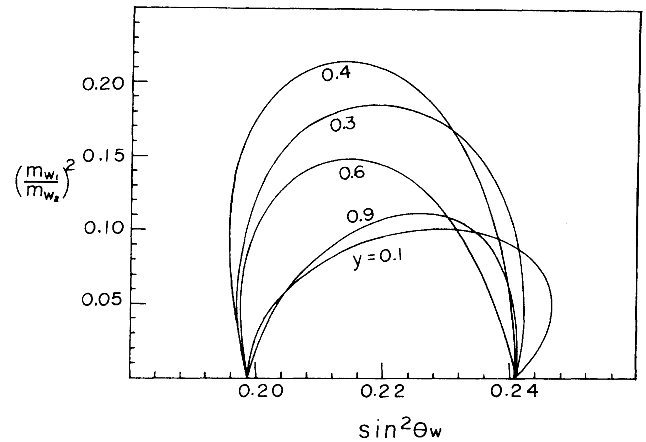

FIG. 3. Same as in Fig. 2, but based on all neutral-current data.

where $M_{Z_1}{}^2$ is the smaller of the two eigenvalues derivable from Eq. (3.11). The maximum allowed value of ξ in Fig. 2 is about 0.20.

We have also considered other neutral-current constraints. For elastic neutrino and antineutrino scattering off the electron, we use the compilation of Kim, Langacker, Levine, and Williams[18] for

$$g_V^e = \frac{1}{1-\xi y^2}\left[-\tfrac{1}{2}+2x+\xi(1-2x-\tfrac{1}{2}y)\right] , \quad (4.10)$$

$$g_A^e = \frac{1}{1-\xi y^2}\left(-\tfrac{1}{2}+\tfrac{1}{2}\xi y\right) . \quad (4.11)$$

For parity nonconservation in atomic transitions, we use the two most recent results[19] for

$$\sum g_A^e g_V^u = \frac{1}{1-\xi y^2}\left[-\tfrac{1}{4}+\tfrac{2}{3}x+\xi(-\tfrac{1}{4}+\tfrac{2}{3}x)y\right] , \quad (4.12)$$

$$\sum g_A^e g_V^d = \frac{1}{1-\xi y^2}\left(\tfrac{1}{4}-\tfrac{1}{3}x-\tfrac{1}{3}\xi xy\right) . \quad (4.13)$$

The measurements[20] of the asymmetry in polarized-electron scattering off the deuteron are also used to fit

$$2\sum g_A^e g_V^u - \sum g_A^e g_V^d = \frac{1}{1-\xi y^2}\left[-\tfrac{3}{4}+\tfrac{5}{3}x+\xi(-\tfrac{1}{2}+\tfrac{5}{3}x)y\right] , \quad (4.14)$$

$$2\sum g_V^e g_A^u - \sum g_V^e g_A^d = \frac{1}{1-\xi y^2}\left[-\tfrac{3}{4}+3x+\xi(1-2x-y+xy)\right] . \quad (4.15)$$

In Fig. 3 we show again the one-standard-deviation allowed region in the (x,ξ) plane for various given values of y, but now using all of the above data. The maximum allowed value of ξ in Fig. 3 is about 0.21. We have not used the constraint due to measurements of the forward-backward asymmetry in $e^+e^- \to \mu^+\mu^-$, etc., for which

$$\sum g_A^e g_A^\mu = \frac{1}{1-\xi y^2}\left(-\tfrac{1}{2}\right)^2 . \quad (4.16)$$

This result is greater than that of the standard model by a factor of $1/(1-\xi y^2)$, but it does not really limit ξ, because we can always vary y to make ξy^2 small.

Using $M_{W_1}=82$ GeV, our bound $\xi < 0.21$ implies that

$$M_{W_2} > 180 \text{ GeV} \quad (4.17)$$

and

$$M_{Z_2} > 210 \text{ GeV} . \quad (4.18)$$

V. CONCLUSION

If there is physics beyond the standard $SU(3)\times SU(2)\times U(1)$ model, it may or may not show up at energies below a TeV. From an experimental point of view, it will certainly be more interesting if the new physics is indeed accessible. In this paper, we have discussed just such a model. It is a new left-right model based on the particle content of the **27** representation of E_6. It is a supersymmetric model with a conserved multiplicative quantum number which is a generalization of the usual R parity. As a result, there are several notable features in this model, beyond the well-known one that the lightest particle of odd R parity is stable. The second W boson of this model has odd R parity, so it does not mix with the standard W boson. This paves the way for the possibility of a massless Dirac neutrino. In other words, if the Dirac mass of neutrino is zero at the tree level, it remains zero to all orders because of R parity. Of course, it is also allowed to be nonzero, in which case its value is arbitrary. This model also has the interesting feature that the right-handed mass partner of the neutrino is not the same as its current partner, as in conventional left-right models. In fact, the former is inert and the latter may be heavy. This means that the astrophysical limit on the effective total number of neutrino, i.e., $N_\nu < 4$, is not violated, and that the constraint on M_{W_2} from polarized μ^+ decay is not applicable. Similarly, since W_2 does not couple to d and s quarks, the K_L-K_S mass difference is also not a constraint.

The Higgs sector of this model is surprisingly simple. Only three nonzero vacuum expectation values are needed. In addition to the gauge couplings, only four other parameters are required to specify the most general Higgs scalar potential. Since particles of opposite R parity cannot mix, there is also simplification in the mass matrices for the physical Higgs bosons. In particular, we find an absolute upper bound of $\sqrt{2}M_{W_1} = 116$ GeV on the mass of one of the neutral Higgs bosons of this model, although that value is likely to be much lower.

The neutral-current interaction of this model involves a second boson D in addition to the standard Z, and the two can mix in general. We have analyzed the constraints due to present experimental data, and we find lower bounds of 180 and 210 GeV, respectively, for M_{W_2} and M_{Z_2}. In the context of this model, prospects of many new physics discoveries exist at forthcoming and proposed high-energy accelerators.

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