Grand unification and parity restoration at low energies. Phenomenology

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Recently, we have shown that simple grand unified theories, such as SO(10), allow the possibility of low-energy parity restoration. In this paper a detailed phenomenological analysis of the left-right-symmetric model is presented. It turns out that the theory passes all the charged- and neutral-current tests for M_{W_R} as low as 100–300 GeV, if simultaneously $\sin^2\theta_W \simeq 0.25$–$0.31$. We also estimate masses, widths, and branching ratios for the gauge bosons, as well as their production rates in pp and $\bar{p}p$ reactions.

I. INTRODUCTION

The conventional picture[1] of grand unification assumes a large desert[2] in energies above the mass scale of weak interactions, i.e., the mass of the W boson. The simplest grand unified theory, the SU(5) model of Georgi and Glashow,[3] actually *predicts*[4] no new physics in energies between M_W and the unification scale $M_U \simeq 10^{14}$ GeV. It is then important to know whether the above is necessarily true for other simple grand unified theories, such as O(10). The answer to this, as we have recently shown,[5] is *no*: O(10) grand unified theory[6] allows the possibility of a low intermediate mass scale of under a few hundred GeV, above which parity is expected to become a good symmetry. In this paper we offer a detailed phenomenological analysis of such a model, with special care paid to make sure that all the low-energy tests are satisfied.

Let us briefly recall the salient features of O(10) theory and the intermediate weak-interaction theory based on $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The often-mentioned effective features of O(10) are that it treats all the fermions symmetrically, by placing them in 16-dimensional spinorial representations, and that it predicts a naturally small neutrino mass.[7] Both of these properties are due to the fact that O(10), unlike SU(5), contains the $SU(2)_L \times SU(2)_R \times SU(4)$ model of Pati and Salam. In our discussion, its subgroup $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$[8] will play a dominant role. The models based on this group were suggested originally in order to understand parity violation in weak interactions.[8] One starts with a completely left-right-symmetric theory, but then noninvariance of the vacuum under parity conjugation, which results in a large mass for a right-handed gauge boson, leads to parity violation at low energies. The question then arises as to what the mass of W_R is. As was recently shown by Mohapatra and one of us (G.S.),[9] the right-handed neutrino in these theories is a heavy neutral Majorana lepton with $M_{\nu_R} \simeq M_{W_R}$, so it does not participate in β and μ decay. However, the constraint on M_{W_R} comes from the fact that there is also a heavy neutral gauge boson Z_2 with $M_{Z_2} \propto M_{W_R}$. A major portion of this paper is devoted to the analysis of neutral-current data,[10] which shows that M_{W_R} can be as low as 100–300 GeV, if simultaneously the value for $\sin^2\theta_W$ is increased beyond the standard-model prediction: $\sin^2\theta_W \simeq 0.25$–$0.31$. Of course, this analysis gives only a lower limit on M_{W_R}. Fortunately, if one takes seriously the idea of unification of weak, electromagnetic, and strong interactions, then the program of Georgi, Quinn, and Weinberg[4] gives the limits on mass scales beyond M_W. In a later paper we will present the following possibilities.

(1) If $M_{W_R} \gtrsim 300$ GeV, then $\sin^2\theta_W \simeq 0.23$, as in the standard model. That leads[11] to a constraint $M_{W_R} \gtrsim 10^9$ GeV, with $M_U \leq 10^{19}$ GeV. For $M_{W_R} \simeq 10^9$ GeV, $M_U \simeq 10^{19}$ GeV the proton would be effectively stable. Of course, it is possible that $M_{W_R} \simeq M_U$, in which case we have the conventional picture of a desert, with the proton lifetime estimated as in SU(5): $\tau_p \simeq 10^{31\pm2}$ yr. In any case, since M_{W_R} is so large, we would have no way of direct observation of parity restoration.

(2) In the second case: $M_{W_R} \simeq 150$–240 GeV, $M_{W_L} \simeq 65$–70 GeV, and $\sin^2\theta_W \simeq 0.28$. This provides a perfectly consistent picture of unification, with the further implication from the low-energy values of α_s that $M_U \simeq 10^{18}$–10^{19} GeV, so that $\tau_p \gtrsim 10^{42}$ yr. We should not observe proton decay, but rather restoration of parity at low energies.

We find the second possibility rather exciting, since it offers a possible existence of new thresholds at energies reachable in the near future, and yet leads to grand unification. This provides a motivation to study in detail the phenomenology of the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ theory with large $\sin^2\theta_W$ and low M_{W_R}. We find that, if the above picture is correct, the second generation of gauge bosons should be observed at ISABELLE energies. Another good test of the model is, of course, lower values of the light-gauge-boson masses compared to the standard-model predictions.

As we mentioned before, we have left a detailed discussion of the implications of grand unification for a later paper. There we also will discuss the question of the Majorana character of neutrinos and associated lepton-number violation, manifested through neutrinoless double-β decay.

We have then organized the rest of this paper in the following manner. In Sec. II we review the basic features of left-right-symmetric theories. These models have been discussed at length recently,[9] and so we will be somewhat brief. In Secs. III and IV we derive the low-energy effective Hamiltonian and neutral currents needed for the comparison of the theory with experiment. In Sec. IV a detailed phenomenological analysis is also performed. From the existing neutral-current data, we calculate the masses, total widths, and branching ratios for both charged and neutral gauge mesons and show their parameter dependence. We then calculate the production cross sections for pp and $\bar{p}p$ collisions. Finally, Sec. V is devoted to comments and a summary of basic results.

II. A LEFT-RIGHT-SYMMETRIC THEORY

As we said, the gauge group is $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The first manifestation of left-right symmetry is symmetrically placed left-handed and right-handed fermions

$$\psi_L=\begin{pmatrix}\nu_L\\ e_L\end{pmatrix}\ (\tfrac{1}{2},0,-1),\quad \psi_R=\begin{pmatrix}\nu_R\\ e_R\end{pmatrix}\ (0,\tfrac{1}{2},-1),$$
$$Q_L=\begin{pmatrix}u_L\\ d_L\end{pmatrix}\ (\tfrac{1}{2},0,\tfrac{1}{3}),\quad Q_R=\begin{pmatrix}u_R\\ d_R\end{pmatrix}\ (0,\tfrac{1}{2},\tfrac{1}{3}),\tag{2.1}$$

where the numbers in parentheses denote the quantum numbers of $SU(2)_L$, $SU(2)_R$, and $U(1)_{B-L}$ representations, respectively. It is then easy to see that the formula for the electric charge reads[12]

$$Q=I_{3L}+I_{3R}+\frac{B-L}{2}\ .\tag{2.2}$$

Since $\Delta Q=0$ and, if we are below M_{W_R}, $\Delta I_{3L}\simeq 0$, the above relation gives

$$\Delta(B-L)\simeq -2\Delta I_{3R}\tag{2.3}$$

which links the breaking of $B-L$ to the breakdown of parity. As we shall see below, the recently suggested model[9] makes full use of (2.3) by having a neutrino being a Majorana particle whose mass is related to the maximality of parity violation in weak interactions.

The Higgs sector, chosen by reasons of simplicity and the physical use of (2.3), is given by[9]

$$\phi\equiv(\tfrac{1}{2},\tfrac{1}{2},0)\,,$$
$$\Delta_L\equiv(1,0,2)\,,\quad \Delta_R\equiv(0,1,2)\,.\tag{2.4}$$

Notice that the above Higgs-boson multiplets have the same representation content as fermionic bilinears: $\phi\sim\bar{\psi}_L\psi_R(\bar{Q}_LQ_R)$, $\Delta_L\sim\psi_L^T\psi_L$, and $\Delta_R\sim\psi_R^T\psi_R$, so that the considerations presented here should be valid, at least qualitatively, in the case of dynamical symmetry breaking. Now, the symmetric Higgs potential allows the asymmetric minimum

$$\langle\Delta_L\rangle=\begin{pmatrix}0&0\\ v_L&0\end{pmatrix},\quad \langle\Delta_R\rangle=\begin{pmatrix}0&0\\ v_R&0\end{pmatrix},\quad \langle\phi\rangle=\begin{pmatrix}k'&0\\ 0&k\end{pmatrix},\tag{2.5}$$

where $k'\ll k$ in order to suppress W_L-W_R mixing and $\Delta S=2$ Higgs-particles-induced process; and $v_R\gg k$ in order to lead to heavier right-handed gauge boson; and

$$v_L\simeq\gamma k^2/v_R\tag{2.6}$$

(γ is the ratio of Higgs-boson self-couplings). To the leading order in k/v_R and neglecting W_L-W_R mixing, one obtains the following gauge bosons: $W_L^\pm$, $W_R^\pm$ in the charged sector and Z_1, Z_2 in the neutral sector with

$$M_{Z_1}{}^2\simeq\frac{M_{W_L}{}^2}{\cos^2\theta_W}\ ,\quad M_{Z_2}{}^2\simeq 2\,\frac{\cos^2\theta_W}{\cos 2\theta_W}\,M_{W_R}{}^2\,.\tag{2.7}$$

In the above, we have defined $\tan^2\theta_W=g'^2/g^2+g'^2$, as to obtain $e^2=g^2\sin^2\theta_W$ as in the standard model. In the next section we give the precise expressions for gauge-boson masses and also effective charged- and neutral-current low-energy Hamiltonians.

A few remarks are now in order concerning leptons, especially neutrinos. The most general Yukawa couplings, consistent with gauge and left-right symmetry are

$$\mathcal{L}_Y=h_1\bar{\psi}_L\phi\psi_R+ih_2(\psi_L^TC\tau_2\Delta_L\psi_L+\ \psi_R^TC\tau_2\Delta_R\psi_R)$$
$$+\ \text{H.c.}\tag{2.8}$$

Substituting $\langle\phi\rangle\neq 0$ gives the usual Dirac mass for the electron. The situation with the neutrino is slightly more subtle and we, therefore, briefly recall it (for details see Ref. 9). As is clear from (2.8) the right-handed neutrino gets a large Majorana mass from the term

$$\mathcal{L}_Y^{\langle\Delta_R\rangle}=h_2v_RN^TCN+\text{H.c.}\,,\tag{2.9}$$

where $N\equiv C(\bar{\nu}_R)^T$. Therefore, $M_N\simeq h_2v_R\simeq h_2/gM_{W_R}$. Through the mixing terms between ν_L and ν_R (couplings h_1 and h_2), the left-handed neutrino in turn gets a small Majorana mass ($\nu\equiv\nu_L$)

$$m_\nu=\text{const}\times\frac{m_e{}^2}{m_N}\ .\tag{2.10}$$

We cannot predict the constant in (2.9), but under reasonable assumptions it is of order one.

An important comment is in order. It turns out[9] that the ν-N mixing is proportional to m_ν/m_N. Since $m_N \simeq M_{W_R} \simeq 100$ GeV (or so), clearly $m_{\nu_\mu}/m_N \leq 10^{-6}$ for ν_μ and even smaller for ν_e; we will, in subsequent discussions, neglect all the mixings between light and heavy Majorana neutrinos.

Finally, we wish to point out that if lower bounds for M_{W_R} are saturated, as the phenomenological discussion in Sec. IV and unification constraints which we treat in the subsequent publication allow, then for $m_N \simeq 100$ GeV we would obtain[9]

$$\begin{aligned} m_{\nu_e} &\simeq 10 \text{ eV},\\ m_{\nu_\mu} &\simeq 100 \text{ keV},\\ m_{\nu_\tau} &\simeq 100 \text{ MeV}. \end{aligned} \tag{2.11}$$

The large values for $m_{\nu i}$ and low values for m_{Ni} lead to lepton-number violation.[9] We will discuss in detail the properties of ν and N neutrinos in a later paper after we show that $M_{W_R} \simeq 150$–250 GeV is to be expected on the basis of grand unification (unless $M_{W_R} \leq M_U$).

III. PARAMETRIZATION OF THE LEFT-RIGHT-SYMMETRIC MODEL

With the Higgs representation described in the last section we can immediately write down the mass matrix for both the charged and neutral gauge bosons; we first introduce the following parameters:

$$\eta_L \equiv \frac{v_L^2}{k^2+k'^2}, \quad \eta_R = \frac{k^2+k'^2}{v_R^2}, \quad z=\frac{2k'^2}{k^2+k'^2} \tag{3.1}$$

with the allowed ranges

$$0 \leq \eta_{L,R},\ z \leq 1. \tag{3.2}$$

The physical meaning of the parameters in (3.1) is the following: η_R measures the amount of breaking of $SU(2)_R$ (small η_R means strong breaking of parity), η_L measures the presence of a triplet Δ_L in neutral-current phenomena, and z clearly tests the amount of W_L-W_R mixing [see (3.3) below].

We find for the charged-gauge-bosons mass matrix

$$M_W^2 = \tfrac{1}{2} g^2 v_R^2 \times \begin{bmatrix} \eta_R(1+\eta_L) & -\eta_R[1-(1-z)^2]^{1/2} \\ -\eta_R[1-(1-z)^2]^{1/2} & 1+\eta_R \end{bmatrix} \tag{3.3}$$

with eigenvalues

$$M_W^2 = \tfrac{1}{4} g^2 v_R^2 \left(1+2\eta_R+\eta_R\eta_L \pm \{(1-\eta_L\eta_R)^2 + 4\eta_R^2[1-(1-z)^2]\}^{1/2}\right) \tag{3.4}$$

the lighter (heavier) of which we call W_1 (W_2). In the limit of vanishing mixing, i.e., $z \to 0$, W_1 and W_2 become W_L and W_R, respectively.

For the neutral gauge bosons the mass matrix is

$$M_Z^2 = \tfrac{1}{2} g^2 v_R^2 \begin{bmatrix} \eta_R(1+2\eta_L) & -\eta_R & -2\epsilon\eta_L\eta_R \\ -\eta_R & 2+\eta_R & -2\epsilon \\ -2\epsilon\eta_L\eta_R & -2\epsilon & 2(1+\eta_L\eta_R)\epsilon^2 \end{bmatrix}, \tag{3.5}$$

where $\epsilon \equiv g'/g$. With $e = g\sin\theta_W$ ($x_W = \sin^2\theta_W$) we find the eigenvalues to be

$$M_Z^2 = \tfrac{1}{2} g^2 v_R^2 (1+\eta_L\eta_R)\frac{1-x_W}{1-2x_W}\left\{1+\frac{\eta_R}{1+\eta_L\eta_R}\,\frac{1-2x_W}{1-x_W} \pm \left[1+(1-2x_W)\left(\frac{\eta_R^2(1-2x_W)-2\eta_R(1+\eta_L\eta_R)x_W-4\eta_L\eta_R}{(1+\eta_L\eta_R)^2(1-x_W)^2}\right)\right]^{1/2}\right\}. \tag{3.6}$$

We also note the relations between $\sin\theta_W$ and (g,g'):

$$\sin^2\theta_W = \frac{g'^2}{g^2+2g'^2}, \quad \cos^2\theta_W = \frac{g^2+g'^2}{g^2+2g'^2}. \tag{3.7}$$

The matrix (3.3) can be diagonalized immediately by an orthogonal rotation through an angle ξ defined by

$$\tan 2\xi = \frac{2\eta_R[1-(1-z)^2]^{1/2}}{(1-\eta_R\eta_L)}. \tag{3.8}$$

Since W_L couples to J_L and W_R to J_R we can easily write down the currents coupling to the eigenstates W_1 and W_2:

$$\mathcal{L}_{CC} = \frac{g}{2\sqrt{2}}\left[(J_L\cos\xi + J_R\sin\xi)W_1 + (-J_L\sin\xi + J_R\cos\xi)W_2\right]. \tag{3.9}$$

The currents coupling to $Z_{1,2}$ are much more complex; we find

$$J_\mu^{1,2}=\frac{g}{2\cos\theta_W}[\gamma_\mu(1+\gamma_5)O_L^{1,2}+\gamma_\mu(1-\gamma_5)O_R^{1,2}] \quad (3.10)$$

with

$$O_L^1=\cos\phi(T_{3L}-Q\sin^2\theta_W)+\frac{\sin\phi\sin^2\theta_W}{(\cos 2\theta_W)^{1/2}}(T_{3L}-Q),$$
$$O_R^1=-Q\sin^2\theta_W\cos\phi+\frac{\sin\phi}{(\cos 2\theta_W)^{1/2}}(T_{3R}\cos^2\theta_W-Q\sin^2\theta_W); \quad (3.11)$$

for $O_{L,R}^2$ we merely make the transformations

$$\begin{pmatrix}\cos\phi\to-\sin\phi\\ \sin\phi\to\cos\phi\end{pmatrix} \quad (3.12)$$

in Eq. (3.11). For any particular fermion, $T_{3L(R)}$ is the third component of left- (right-) handed weak isospin and Q the charge of that fermion. The angle ϕ is found through the relation

$$\tan\phi=\frac{1+2\eta_L-(M_{Z_1}{}^2\cos^2\theta_W/\frac{1}{2}g^2v_R{}^2\eta_R)}{(\cos 2\theta_W)^{1/2}(1-2\eta_L\sin^2\theta_W/\cos 2\theta_W)}.$$

Using the abbreviations

$$c\equiv\cos\phi,\quad s\equiv\sin\phi,\quad d\equiv(1-2\sin^2\theta_W)^{-1/2}, \quad (3.14)$$

we can now write down the various fermion couplings:

$$O_L^1(\nu)=\tfrac{1}{2}(c+sdx_W),\qquad O_L^2(\nu)=\tfrac{1}{2}(-s+cdx_W), \quad (3.15)$$

$$O_R^1(N)=\tfrac{1}{2}sd(1-x_W),\qquad O_R^2(N)=\tfrac{1}{2}cd(1-x_W), \quad (3.16)$$

$$\begin{aligned}O_L^1(e)&=c(-\tfrac{1}{2}+x_W)+\tfrac{1}{2}sdx_W,\\ O_R^1(e)&=cx_W+\tfrac{1}{2}sd(-1+3x_W),\\ O_L^2(e)&=-s(-\tfrac{1}{2}+x_W)+\tfrac{1}{2}cdx_W,\\ O_R^2(e)&=-sx_W+\tfrac{1}{2}cd(-1+3x_W),\end{aligned} \quad (3.17)$$

$$\begin{aligned}O_L^1(u)&=c(\tfrac{1}{2}-\tfrac{2}{3}x_W)-\tfrac{1}{6}sdx_W,\\ O_R^1(u)&=-\tfrac{2}{3}cx_W+\tfrac{1}{2}sd(1-\tfrac{7}{3}x_W),\\ O_L^2(u)&=-s(\tfrac{1}{2}-\tfrac{2}{3}x_W)-\tfrac{1}{6}cdx_W,\\ O_R^2(u)&=\tfrac{2}{3}sx_W+\tfrac{1}{2}cd(1-\tfrac{7}{3}x_W),\end{aligned} \quad (3.18)$$

$$\begin{aligned}O_L^1(d)&=c(-\tfrac{1}{2}+\tfrac{1}{3}x_W)+\tfrac{1}{6}sdx_W,\\ O_R^1(d)&=\tfrac{1}{3}cx_W+\tfrac{1}{2}sd(-1+\tfrac{5}{3}x_W),\\ O_L^2(d)&=-s(-\tfrac{1}{2}+\tfrac{1}{3}x_W)+\tfrac{1}{6}cdx_W,\\ O_R^2(d)&=-\tfrac{1}{3}sx_W+\tfrac{1}{2}cd(-1+\tfrac{5}{3}x_W).\end{aligned} \quad (3.19)$$

Given the gauge-boson masses and couplings we can now proceed with the phenomenological discussion of the model.

IV. PHENOMENOLOGY OF THE MODEL

In discussing the phenomenological details of the left-right-symmetric model we first must turn to the familiar low-Q^2 effective Hamiltonians for both neutral- and charged-current interactions. To do this, we need to define the Fermi constant G_F in terms of the model parameters. In the limit wherein mixings between the ν_i and N_i states are neglected (these may be smaller than 10^{-5}) the leptonic $\nu_l l$ charged currents are purely left-handed. We find then from μ decay

$$\frac{G_F}{\sqrt{2}}\equiv\tfrac{1}{8}g^2\left(\frac{\cos^2\xi}{M_1{}^2}+\frac{\sin^2\xi}{M_2{}^2}\right), \quad (4.1)$$

where M_i are the masses of the charged gauge bosons and ξ is the relevant mixing angle given by

$$\tan 2\xi=\frac{2\eta_R}{(1-\eta_L\eta_R)}[1-(1-z)^2]^{1/2}. \quad (4.2)$$

Using the eigenvalues $M_i{}^2$, we thus are lead to

$$\frac{G_F}{\sqrt{2}}\equiv\frac{1}{4(k^2+k'^2)}\frac{1+\eta_R}{1+\eta_L+\eta_L\eta_R+\eta_R(1-z)^2}, \quad (4.3)$$

where we have introduced the parameters discussed in the previous section. Taking this definition, the general charged-current–current Hamiltonian is

$$\mathcal{H}_{CC}=\frac{G_F}{\sqrt{2}}[J_L^\dagger J_L+\alpha(J_L^\dagger J_R+J_R^\dagger J_L)+\beta J_R^\dagger J_R] \quad (4.4)$$

with

$$\alpha\equiv\frac{-\eta_R}{1+\eta_R}[1-(1-z)^2]^{1/2},\qquad \beta\equiv\frac{\eta_R(1+\eta_L)}{(1+\eta_R)}, \quad (4.5)$$

and the currents $J_{L,R}$, for any fermion pair $f_{1,2}$, are given by

$$J_{L,R}=\bar{f}_1\gamma_\mu(1+\gamma_5)f_2. \quad (4.6)$$

For purely leptonic interactions, only the $J_L^\dagger J_L$ term is relevant and, hence, the presence of right-handed currents in this model leaves μ decay unaffected. For semileptonic decays, such as π decay there is also a contribution from the term proportional to α from the right-handed quark currents. This modifies the rates for π and K decay but *not* the polarization of the outgoing lepton since this comes only from the form of the leptonic current which is unmodified. It may be remarked, however, that for reasonably small values of $|\alpha|$ ($\lesssim 0.1$) we can always redefine the values of $f_{\pi,K}$:

$$\Gamma(\pi, K\to \mu\bar{\nu}) \propto (f_{\pi,K})_0^2(1-\alpha)^2 \equiv (f_{\pi,K})^2 . \quad (4.7)$$

Obviously, this reduces the true value of $f_{\pi,K}$ but keeps the ratio f_K/f_π invariant; there is some evidence[13] based on the Goldberger-Treiman relation that f_π should indeed be somewhat smaller than $f_\pi \simeq 93$ MeV which is obtained when $\alpha = 0$. There are effects of similar magnitude for β decay as well as K_{e3} decay; for nonleptonic interactions all three terms in the Hamiltonian contribute but the phenomenological implications become obscured.

A full discussion of the effect of these right-handed quark currents is beyond the scope of this paper and will be discussed in a future work.

Let us now turn to the various neutral-current processes; the *effective Hamiltonian for both the SLAC asymmetry experiment*[14] *and atomic parity violation can be expressed as*

$$\mathcal{H}_{PV} = \mathcal{H}^S_{PV} AB , \quad (4.8)$$

where we define

$$A = \frac{1+\eta_R(1-z)^2+\eta_L(1+\eta_R)}{1+\eta_L(2+\eta_R)} ,$$
$$B = \frac{1-\eta_L\eta_R}{1+\eta_R} , \quad (4.9)$$

and $\mathcal{H}^S_{PV}$ is the same as in the standard model

$$\mathcal{H}^S_{PV} = \frac{G_F}{\sqrt{2}}[\bar{e}\gamma_\mu(-1+4x_W)e\bar{q}\gamma^\mu\gamma_5 T_3 q - e\gamma_\mu\gamma_5 e\bar{q}\gamma^\mu(T_3-2x_W Q)q] .$$

Thus

$$Q_w = -AB[N - Z(1-4\sin^2\theta_W)] \quad (4.10)$$

and, using the usual notation, we find

$$C_1^u = \frac{AB}{2}(-1+\tfrac{8}{3}\sin^2\theta_W) ,$$
$$C_1^d = \frac{AB}{2}(1-\tfrac{4}{3}\sin^2\theta_W) , \quad (4.11)$$
$$C_2^u = -C_2^d = \frac{AB}{2}(-1+4\sin^2\theta_W) .$$

For neutrino-hadron scattering we find the following Hamiltonian:

$$\mathcal{H}^{\nu N} = \frac{G_F}{\sqrt{2}}\bar{\nu}\gamma_\mu(1+\gamma_5)\nu[\epsilon_L(q)\bar{q}\gamma_\mu(1+\gamma_5)q + \epsilon_R(q)\bar{q}\gamma_\mu(1-\gamma_5)q] ,$$

where q stands for any quark species and

$$\epsilon_L(u) = \frac{A}{2}\left(\frac{1}{2}\frac{\eta_R+2}{\eta_R+1} - \tfrac{4}{3}\sin^2\theta_W\right) ,$$
$$\epsilon_R(u) = \frac{A}{2}\left(\frac{1}{2}\frac{\eta_R}{\eta_R+1} - \tfrac{4}{3}\sin^2\theta_W\right) ,$$
$$\epsilon_L(d) = \frac{A}{2}\left(-\frac{1}{2}\frac{\eta_R+2}{\eta_R+1} + \tfrac{2}{3}\sin^2\theta_W\right) , \quad (4.12)$$
$$\epsilon_R(d) = \frac{A}{2}\left(-\frac{1}{2}\frac{\eta_R}{\eta_R+1} + \tfrac{2}{3}\sin^2\theta_W\right) .$$

For $\overset{(-)}{\nu}_\mu e$ interactions we obtain

$$\mathcal{H}^{\nu e} = \frac{G_F}{\sqrt{2}}\bar{\nu}\gamma_\mu(1+\gamma_5)\nu\bar{e}\gamma_\mu(g_v+g_A\gamma_5)e ,$$

where

$$g_V = -\tfrac{1}{2}(1-4\sin^2\theta_W)A ,$$
$$g_A = -\frac{1}{2}\frac{A}{\eta_R+1} . \quad (4.13)$$

The parameters relevant for the *various asymmetries in* $e^+e^-\to\mu^+\mu^-$ are now given by

$$h_{VV} = \tfrac{1}{4}A\frac{1+2\eta_R+\eta_L\eta_R}{1+\eta_R}(1-4\sin^2\theta_W)^2 ,$$
$$h_{AA} = \tfrac{1}{4}A\frac{1+\eta_L\eta_R}{1+\eta_R} , \quad (4.14)$$
$$h_{VA} = \tfrac{1}{4}A\frac{1-\eta_L\eta_R}{1+\eta_R}(1-4\sin^2\theta_W) ,$$

and the Hamiltonian is given by

$$\mathcal{H}^{PV} = \frac{G_F}{\sqrt{2}}[h_{VV}(\bar{e}\gamma_\mu e+\bar{\mu}\gamma_\mu\mu)(\bar{e}\gamma^\mu e+\bar{\mu}\gamma^\mu\mu) + 2h_{VA}(\bar{e}\gamma_\mu e+\bar{\mu}\gamma_\mu\mu)(\bar{e}\gamma^\mu\gamma_5 e+\bar{\mu}\gamma^\mu\gamma_5\mu) + h_{AA}(\bar{e}\gamma_\mu\gamma_5 e+\bar{\mu}\gamma_\mu\gamma_5\mu)(\bar{e}\gamma^\mu\gamma_5 e+\bar{\mu}\gamma^\mu\gamma_5\mu)] .$$

Given these various phenomenological parameters we now can ask for what values of $\sin^2\theta_W$, z, η_L, and η_R the model has predictions consistent with the present neutral-current data. To make the comparison with experiment we will make use of the work of Kim, Langacker, Levine, and Williams[15] (KLLW) who have done a detailed analysis of the neutral-current data to extract the range of values of these phenomenological parameters allowed. Our procedure is then as follows; we vary $\sin^2\theta_W$, z, η_L, and η_R over a wide range and search for agreement with the values of the phenomenological parameters extracted from the KLLW analysis (by "agreement" we mean that the values of the parameters chosen lead to results which are consistent with the experimental data within $\frac{3}{2}\sigma$ or 2σ limits).

We first note the following restrictions on z and η_R from the value of α in the charged-current Hamiltonian; for a given limit on the value of α

only a certain region of the (z, η_R) plane can be physically realized. Figure 1 shows the allowed regions, below and to the left of the curves in the (z, η_R) plane for three values of $|\alpha|$. Thus, we see that for large η_R, z must be small and vice versa. Since $|\alpha|$ must be small[16] ($|\alpha| \lesssim 0.06$) in order for right-handed currents to have gone undetected experimentally and that η_R must be $\gtrsim 0.1$ to make the model clearly distinguishable from the standard one ($\eta_R = 0$) we are led to theories in which z is relatively small $\lesssim 0.1$. Although there are many models which pass the neutral-current constraints these charged-current considerations force us to consider small-z values ($\lesssim 0.2$) only.

Similarly η_L is found to be small, almost always $\lesssim 0.1$ for all values of $\sin^2\theta_W$, z, and η_R from the neutral-current normalization; η_L essentially measures the relative contribution of the left-handed Higgs triplet Δ_L relative to ϕ. In the standard model the appearance of a Higgs triplet would upset the value of the well-known ρ parameter which is experimentally close to unity; if the ratio of triplet to doublet vacuum expectation values (VEV's) was sufficiently large the good agreement of the standard model with the data would be destroyed. Similarly in the left-right-symmetric model the Higgs-triplet VEV must still be small relative to values of k, k' in the $(\frac{1}{2}, \frac{1}{2})$ representation. In the sample models presented we will thus consider small η_L, specifically $\eta_L = 0$ or 0.1. Table I presents a sample set of models which pass the neutral-current tests with low values of η_L and z; as is usual we note the general decrease of W_1 and Z_1 gauge-boson masses as $\sin^2\theta_W$ increases. We also note that with the increase of η_R for any fixed $\sin^2\theta_W$ the values of W_2 and Z_2 masses decrease drastically. Thus for large η_R and $\sin^2\theta_W$ values we have a set of very light "second-generation" gauge-boson masses.

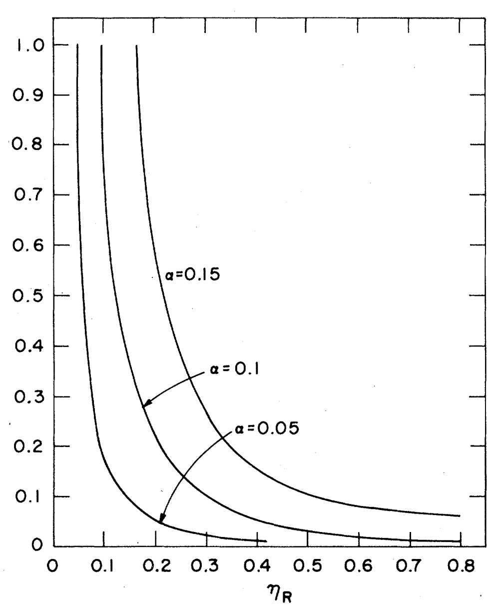


FIG. 1. Bounds on the allowed regions of the (z, η_R) plane for various values of the parameter α. The allowed region is below and to the left of the relevant curve.

Figures 2–5 show the allowed region in the (η_R, $\sin^2\theta_W$) plane for four values of z and η_L using both $\frac{3}{2}\sigma$ and 2σ limits on the ranges allowed for the various phenomenological parameters. The irregular shape of the regions presented results only from the fact that the step size of η_R increment in our analysis was 0.1.

Before turning to a discussion of the expected production rates of the various gauge bosons, we would like to say something about their properties; the masses of these particles in terms of the model parameters are given by Eq. (3.4) for the charged bosons and Eq. (3.6) for the neutral bosons. The widths of the two neutral Z's are given by (assuming $2M_N > M_Z$)

$$\Gamma^{Z_1} \simeq \frac{0.65}{\sin^2\theta_W} \frac{M_{Z_1}}{100\text{ GeV}},$$

$$\Gamma^{Z_2} \simeq \frac{0.48}{\sin^2\theta_W} \frac{M_{Z_2}}{100\text{ GeV}}.$$

For the charged bosons $W^{\pm}_{1,2}$, the widths are given by (assuming we can neglect ξ^2 terms)

$$\Gamma^{W_1} \simeq \frac{0.73}{\sin^2\theta_W} \frac{M_{W_1}}{100\text{ GeV}},$$

$$\Gamma^{W_2} \simeq \frac{0.55}{\sin^2\theta_W} \frac{M_{W_2}}{100\text{ GeV}}.$$

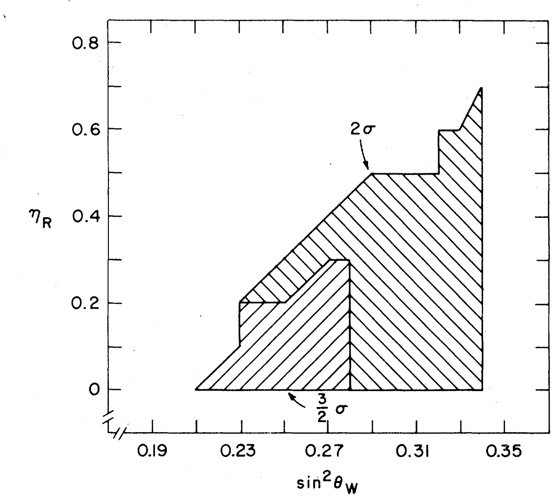


FIG. 2. Allowed region in the η_R, $\sin^2\theta_W$ plane for models passing the neutral-current tests with $z = \eta_L = 0$ (for $\frac{3}{2}\sigma$ or 2σ bounds).

TABLE I. Gauge-boson masses in GeV for various models which pass the neutral-current tests (a) $\frac{3}{2}\sigma$; (b) 2σ.

(a)

$\sin^2\theta_W$	z	η_L	η_R	M_{W_1}	M_{W_2}	M_{Z_1}	M_{Z_2}
0.23	0	0	0.1	77.74	257.82	87.47	420.44
0.23	0	0	0.2	77.74	190.42	86.30	301.34
0.23	0.1	0	0.1	77.67	260.30	88.24	424.14
0.23	0	0.1	0.1	77.74	245.83	91.71	401.02
0.25	0	0	0.2	74.56	182.64	84.08	295.71
0.25	0	0.1	0.2	74.56	174.14	88.48	282.15
0.27	0	0	0.2	71.75	175.75	82.21	291.94
0.28	0	0	0.1	70.46	233.67	82.23	406.97
0.28	0	0	0.2	70.46	172.58	81.39	290.74
0.28	0	0	0.3	70.46	146.66	80.51	240.00
0.28	0.1	0	0.2	70.23	175.93	82.71	295.46
0.28	0.2	0	0.1	70.34	237.98	83.61	413.79
0.28	0.1	0	0.1	70.39	235.92	82.95	410.52

(b)

$\sin^2\theta_W$	z	η_L	η_R	M_{W_1}	M_{W_2}	M_{Z_1}	M_{Z_2}
0.25	0	0	0.3	74.56	155.21	83.00	244.59
0.28	0	0	0.4	70.46	131.81	79.59	210.26
0.28	0	0.1	0.5	70.46	116.35	83.69	181.75
0.28	0	0.2	0.3	70.46	133.89	88.56	219.69
0.30	0	0	0.5	68.07	117.89	77.56	188.89
0.30	0	0.1	0.8	68.07	97.35	80.65	147.72
0.30	0.1	0.0	0.2	67.85	169.96	81.25	294.44
0.32	0	0	0.6	65.90	107.62	75.98	173.96
0.32	0	0.1	0.5	65.90	108.84	81.53	180.49

where we have assumed that $M_N > M_{W_{1,2}}$ to be definite. [The exact form of these expressions is given below in Eq. (4.22).] This last assumption need not be true for W_2 since the mass scale of N and W_2 are set by the same vacuum expectation value, v_R. The branching fractions in the W_1 case are very similar to those of the standard model; those for W_2 are, again, similar except that we assume the decay $W_2 \to N\bar{l}$ cannot occur. This essentially multiplies the other branching ratios by a factor of $\frac{4}{3}$. In the Z-boson case the various couplings are quite sensitive to the model parameter $\sin^2\theta_W$, η_L, η_R, and z and no simple result is representative.

It should be pointed out that the above widths will change (only slightly) if the masses of the right-handed Majorana neutrinos are lighter than the various gauge bosons.

Let us now consider the production[17] of these gauge bosons in pp and $p\bar{p}$ collisions relevant at CERN and ISABELLE using the Drell-Yan[18] mechanism. The cross section for $a+b\to B+X$, where B is the gauge boson of interest, can be written as

$$\frac{d\sigma}{dx_F} = \frac{2\pi}{3M_B^2}\,\frac{x_a x_b}{(x_F^2+4\tau)^{1/2}} \times\sum_i (g_L^2+g_R^2)_i\,[\,q_a^i(x_a)\,\bar{q}_b^{i'}(x_b)+\bar{q}_a^{i'}(x_a)\,q_b^i(x_b)]\,, \tag{4.15}$$

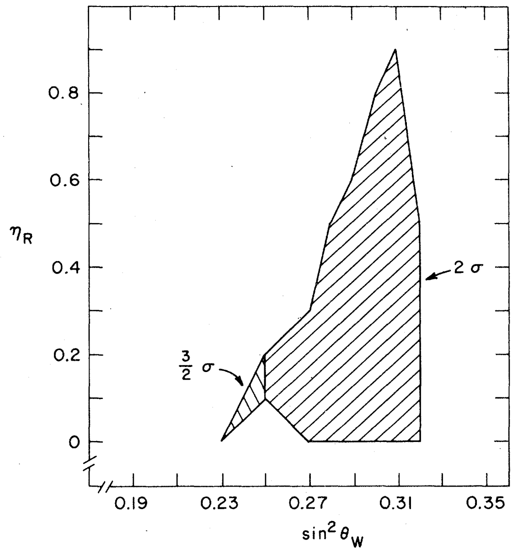


FIG. 3. Same as Fig. 2 but with $\eta_L=0.1$.

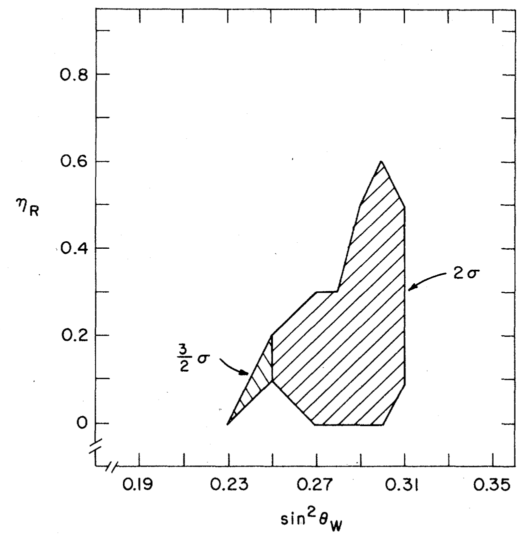


FIG. 5. Same as Fig. 2 but with $\eta_L=z=0.1$.

where i and i' label the relevant quark flavors ($i=i'$ for Z bosons and are separated by one unit of charge for W bosons) $g_{L,R}$ are the relevant chiral couplings to the gauge bosons in question.[19] For neutral-current couplings to the Z_i we refer the reader to the previous section; for charged currents we need the chiral couplings to W_1 and W_2. We may always write the charged-current-W_i Lagrangian as

$$\mathcal{L}=\frac{g}{2\sqrt{2}}(J_L W_L+J_R W_R) \tag{4.16}$$

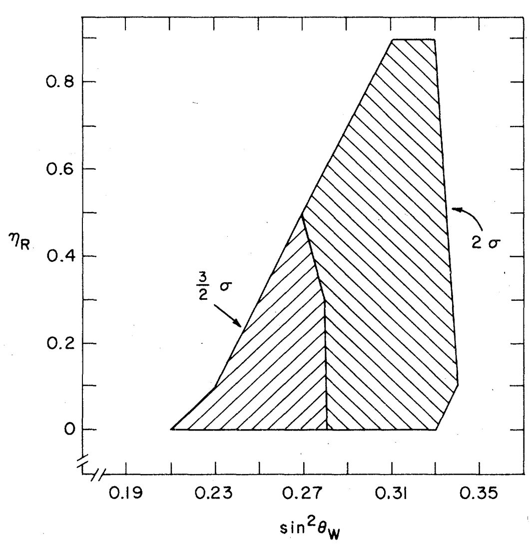


FIG. 4. Same as Fig. 2 but with $z=0.1$.

which in terms of the mass eigenstates W_i is given by

$$\mathcal{L}=\frac{g}{2\sqrt{2}}[(J_L\cos\xi+J_R\sin\xi)W_1 + (-J_L\sin\xi+J_R\cos\xi)W_2] \tag{4.16$'$}$$

with $J_{L,R}$ being the chiral currents (4.6). Note that once both J_L and J_R are present for quarks the combination $g_L{}^2+g_R{}^2$ is equal to unity for any two members of an isodoublet in the weak-eigenstate basis. We need only to include the usual weak, Cabibbo-type angles in order to perform the calculation. Since the heavy-quark content is irrelevant for this calculation we limit ourselves to the usual Cabibbo angle (i.e., the four-quark limit).

In Eq. (4.15) we have defined

$$\begin{aligned}&\tau=M_B{}^2/s\,,\\&x_{a,b}=\tfrac{1}{2}[(x_F{}^2+4\tau)^{1/2}\pm x_F]\,,\end{aligned} \tag{4.17}$$

with x_F being the standard Feynman variable and s is the center-of-mass energy of the pp or $p\bar{p}$ machine. The quark distributions $q_a^i(x_a)$ are the ones relevant for a particle of type a; we take the usual Field-Feynman distributions for purposes of demonstration[20]

$$\begin{aligned}&xu(x)=1.14(1-x)^3\,,\\&xd(x)=2.9(1-x)^4\,,\\&x\bar{u}(x)=0.17(1-x)^{10}\,,\\&x\bar{d}(x)=0.17(1-x)^7\,,\\&xs(x)=x\bar{s}(x)=0.1(1-x)^{10}\,,\\&xc(x)=x\bar{c}(x)=0.03(1-x)^{10}\,.\end{aligned} \tag{4.18}$$

To obtain the total cross section we merely integrate (4.15) over the allowed range of x_F which for a fixed τ is

$$-(1-\tau)\leq x_F \leq (1-\tau)\,. \tag{4.19}$$

Since the calculations are meant to be suggestive only we will make the assumption that these distributions scale and neglect the effects of quantum chromodynamics.

Next, since the usual signal for a Z-type gauge boson is a dimuon $(\mu^+\mu^-)$ peak at the Z mass while that for a W-type gauge boson is to look for high-energy prompt muons from the $\mu\bar{\nu}$ decay, we must calculate the branching ratios for these particular decays. In the Z case, we use the couplings presented in the previous section while for the W boson we take the couplings from (4.16′); in this calculation we will *assume* that *all* of the three heavy Majorana neutrinos are heavier than all of the gauge bosons. This assumption increases the desired branching ratios; however, it may not be valid in general. The heavy Majorana scale *is* set by v_R and, hence, our assumption is most probably correct for W_1 and Z_1 but questionable for W_2 and Z_2 since their masses are also set by the same scale. With this assumption only J_L will effectively contribute for leptons since decays such as $W_2\to N\bar{l}$ and $Z_2\to N\bar{N}$ are kinematically forbidden. We note here the couplings of the W_i to the leptons neglecting ν-N mixings:

$$\mathcal{L}_l=\frac{g}{2\sqrt{2}}\,[\bar{\nu}\gamma_\mu(1+\gamma_5)l\cos\xi+\bar{N}\gamma_\mu(1-\gamma_5)l\sin\xi]W_1 + \frac{g}{2\sqrt{2}}[-\bar{\nu}\gamma_\mu(1+\gamma_5)l\sin\xi+\bar{N}\gamma_\mu(1-\gamma_5)l\cos\xi]\,W_2\,. \tag{4.20}$$

An important observation to make from $\mathcal{L}_l$ is that as $\xi\to 0$ (i.e., $z\to 0$) W_2 decouples from $\bar{\nu}l$ and, hence, for $z=0$ the decay $W_2\to\bar{\nu}l$ is forbidden. For three generations, neglecting fermion masses, we find the branching ratios to be

$$B_1=\frac{\Gamma(W_1\to\mu\nu)}{\Gamma(W_1\to \text{all})}=\frac{\cos^2\xi}{9+3\cos^2\xi}\,,\qquad B_2=\frac{\Gamma(W_2\to\mu\nu)}{\Gamma(W_2\to \text{all})}=\frac{\sin^2\xi}{9+3\cos^2\xi} \tag{4.21}$$

and the total widths to be

$$\Gamma^1_{\text{tot}}=\frac{M_{W_1}}{6\pi}\left(\frac{g}{2\sqrt{2}}\right)^2(9+3\cos^2\xi)\,,\qquad \Gamma^2_{\text{tot}}=\frac{M_{W_2}}{6\pi}\left(\frac{g}{2\sqrt{2}}\right)^2(9+3\sin^2\xi)\,. \tag{4.22}$$

Since $\xi\sim z$ we do not expect the widths and branching ratios of W_1 to be very different from the W of the standard model; this is born out by explicit calculations. Similarly, we might expect the width and branching ratios of Z_1 to be quite similar to the usual Z boson; this, too, has been verified by explicit calculations. In addition, an analysis over a wide class of models indicates that Z_2 has branching ratios similar to Z_1; in particular we find the following branching ratios as quite typical:

$$B(W_1\to\mu\nu)\simeq 8\%\,,\quad B(Z_1\to\mu^+\mu^-)\simeq 3\%\,,\quad B(Z_2\to\mu^+\mu^-)\simeq 2.5\text{–}3\,\%\,. \tag{4.23}$$

Detecting these gauge bosons proceeds in the usual manner; the problem occurs for W_2 which, in general, has a very tiny (if not zero) branching ratio into the $\mu\nu$ final state. It may be possible to detect W_2's by some other means; Figs. 6–11 give the values of σB for Z_2 and σ for W_2^+ in both pp and $p\bar{p}$ reactions for various models. We have not plotted the corresponding results for $W_1^\pm$ and Z_1 since these are quite similar to those expected in the standard model. The results for W_2^- are comparable but slightly lower by a factor of a few than those for W_2^+.

A most interesting set of models are those with relatively large values of $\sin^2\theta_W$ and η_R; these models have extremely low values of the masses of the second generation of gauge bosons. It should also be noted that as the heavier gauge bosons get sufficiently light, the first set of gauge bosons become significantly lighter (by more than 7 or 8 GeV) than those expected in the standard model. Since radiative corrections seem to increase the value of gauge-boson masses[21] observation of such light W and Z's would clearly indicate further structure beyond the standard $SU(2)_L\times U(1)$ group.

Taking σB for the Z in the standard model to be $\sim 5\times 10^{-35}$ cm^2 at ISABELLE ($\sqrt{s}=800$ GeV) energies and 3×10^{-34} cm^2 at the CERN $\bar{p}p$ machine ($\sqrt{s}=540$ GeV) we see that for an ISABELLE luminosity of 10^{33} sec^{-1} we expect $\sim 4.4\times 10^3$ events/day whereas for a CERN luminosity of 10^{30} sec^{-1} we expect $\sim$26 events/day. Let us compare these numbers with a sample set of Z_2's and W_2^+'s for various models; we quote the number of dimuon events expected from the Z_2 peak and the *total* W_2 production rate since the $\mu\bar{\nu}$ signal is suppressed for this particle. The results can be seen in Table II; we have assumed the same luminosities as above. We see immediately that ISABELLE has a chance to detect the lighter Z_2's while CERN seems not to have sufficient luminosity. The production rates for W_2^+ are quite large but may be quite difficult to detect at either high-energy facility since we do not have a $\mu\bar{\nu}$ final-state signal. However, it may be possible to see the jets produced from the heavy quarks in the decay $W_2^+ \to t\bar{b}$.

The situation is different, of course, if N_i, or at least one of them, are lighter than W_2. First, we point out that the predictions for various branching ratios given before are altered in a trivial way (it is obviously a matter of counting only). Since cross-section estimates using the Drell-Yan formula are no better than a factor of 2, small changes of order a few percent due to $W_2 \to \mu N_i$, etc., are clearly negligible within our approxima-

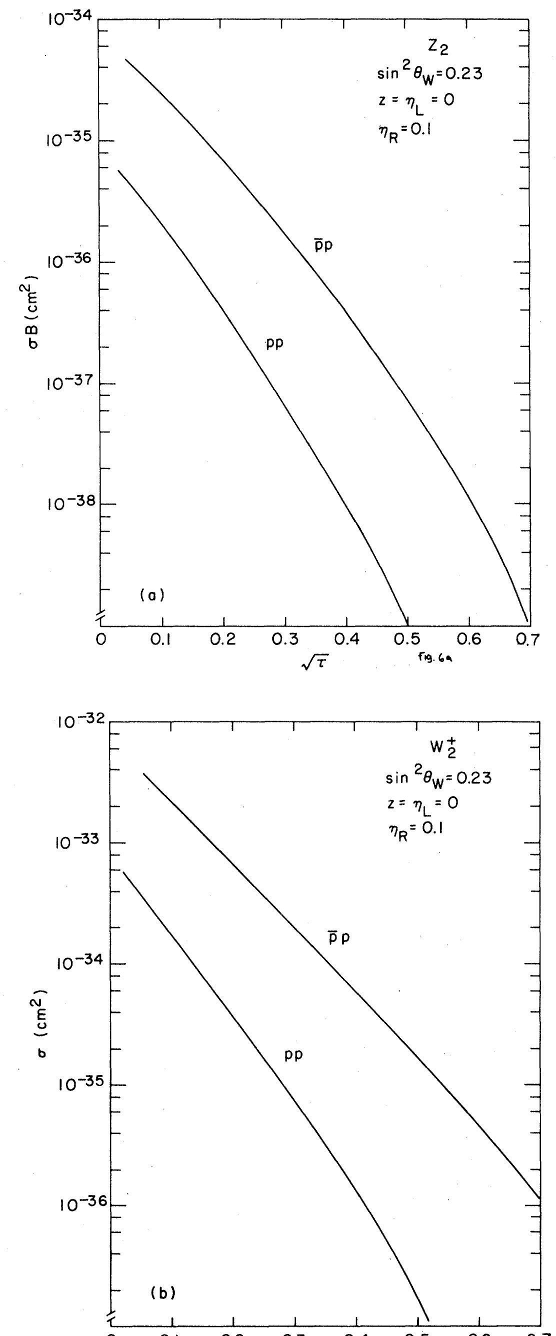


FIG. 6. (a) Z_2 production cross sections for $\bar{p}p$ and pp collisions with $\sin^2\theta_W = 0.23$, $\eta_R = 0.1$, and $\eta_L = z = 0$. (b) Same as (a) but for $W_2^\pm$ production.

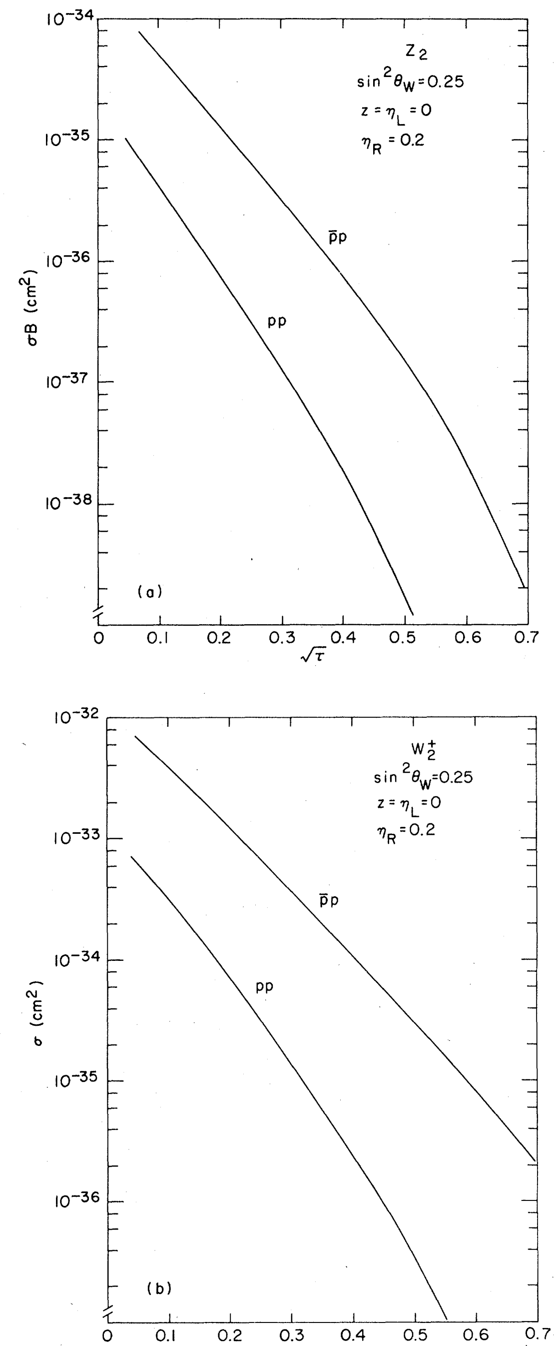


FIG. 7. (a) Z_2 production cross section for $\bar{p}p$ and pp collisions with $\sin^2\theta_W = 0.25$, $z = \eta_L = 0$, and $\eta_R = 0.2$. (b) Same as (a) but for $W_2^\pm$ production.

tions. However, this possibility would provide a clear signal for W_2 production. If this whole picture is correct and we will argue in paper II that it is favored by grand unification, then this could be one of the best probes of the model.

It thus appears quite possible to produce the second set of gauge bosons predicted by the left-right-symmetric model with sufficient rates at either CERN or ISABELLE.

Before ending this section we would like to make a few more phenomenological remarks concerning

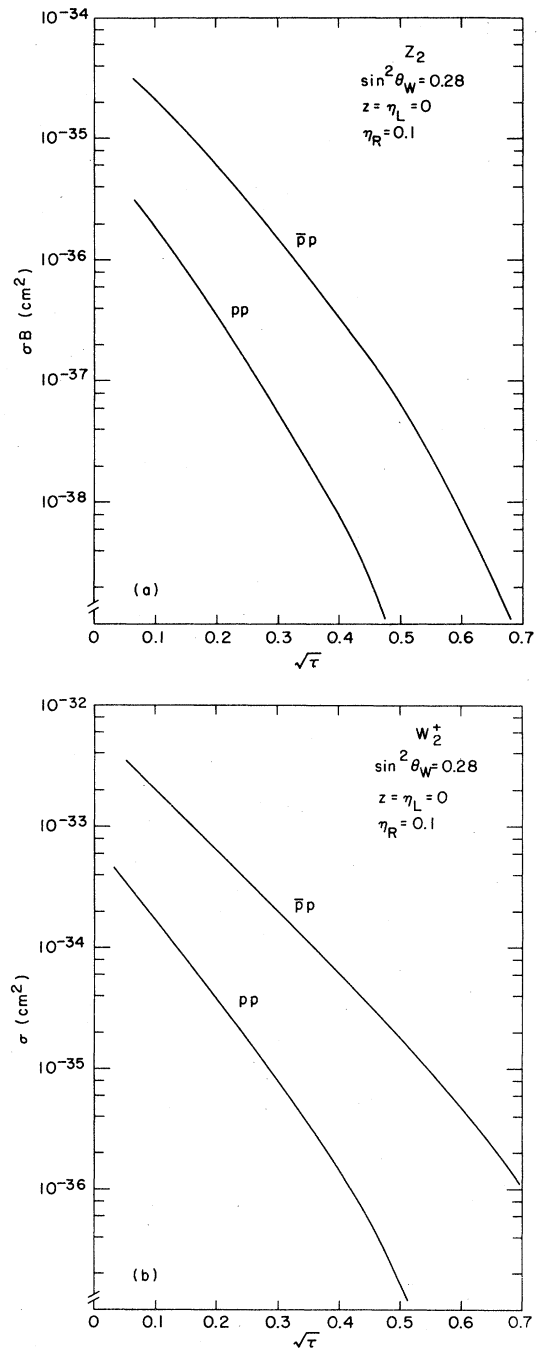


FIG. 8. (a) and (b) Z_2 and W_2^+ production cross sections in $\bar{p}p$ and pp collisions with $\sin^2\theta_W=0.28$, $z=\eta_L=0$, $\eta_R=0.1$.

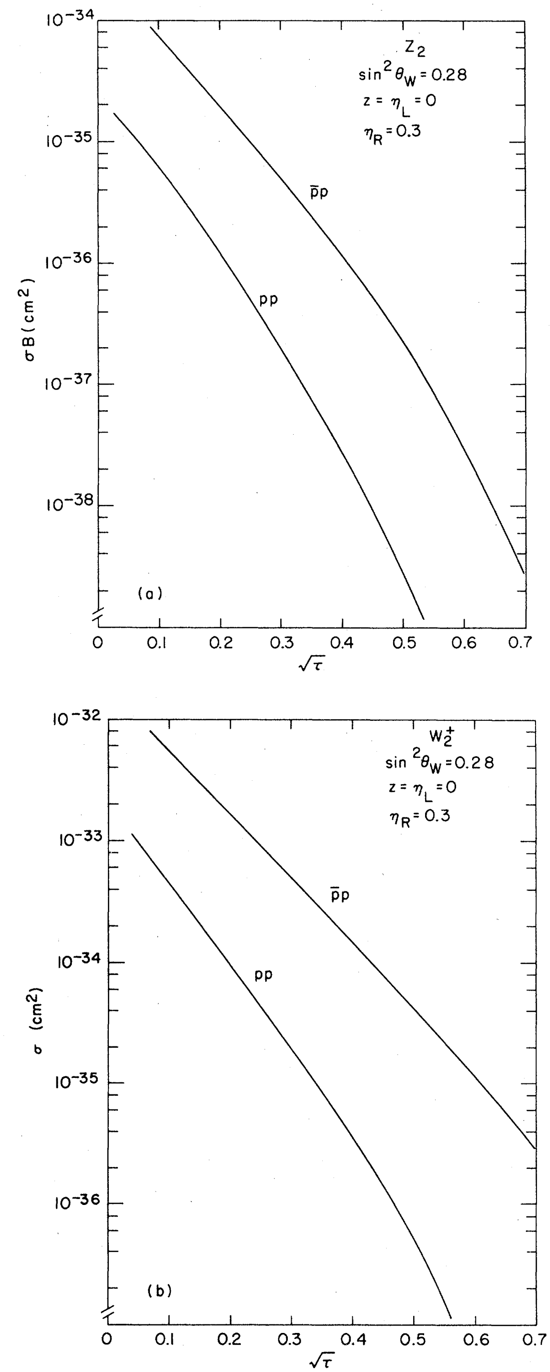
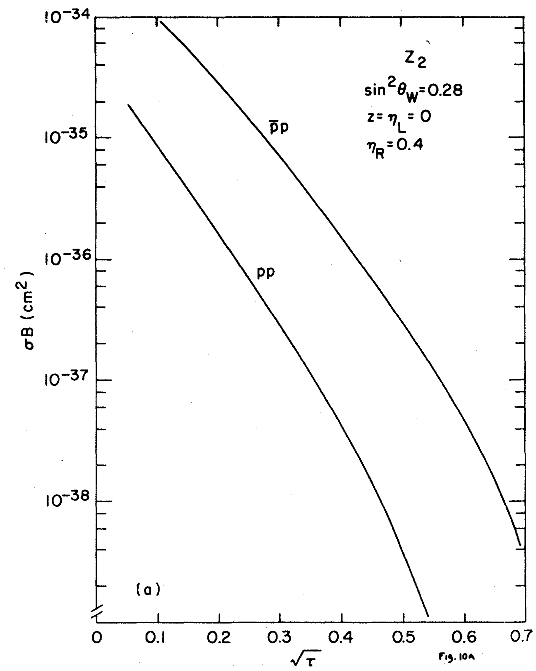


FIG. 9. (a) and (b) Same as Fig. 8 but with $\eta_R=0.3$.

a new set of particles in the theory. Besides the usual left-handed Majorana neutrinos and the usual quarks and leptons, the left-right-symmetric model contains three, heavy, right-handed Majorana neutrinos, N_i $(i=1,2,3)$. These heavy Majorana particles are the major contributions to processes such as neutrinoless double-β decay[22] and lepton-number-violating processes such as $\mu \to e\gamma$. Estimates for the ratios of these processes have been made earlier.[23] The N_i themselves form an interesting set of particles; the mass scale at which they get their masses is v_R and, hence, they may



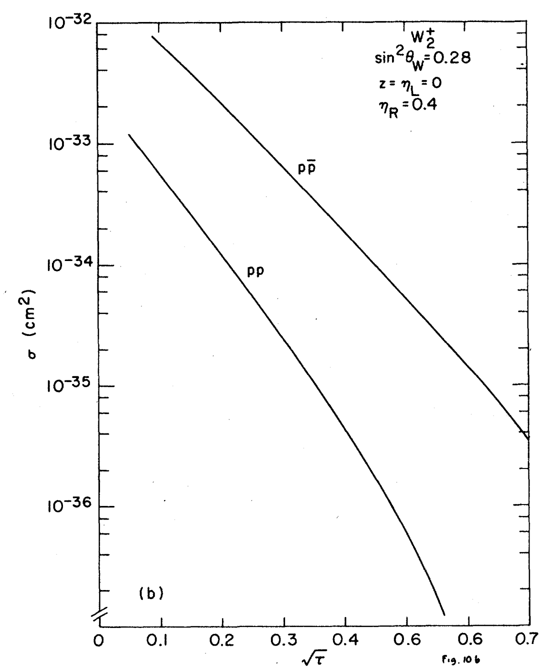
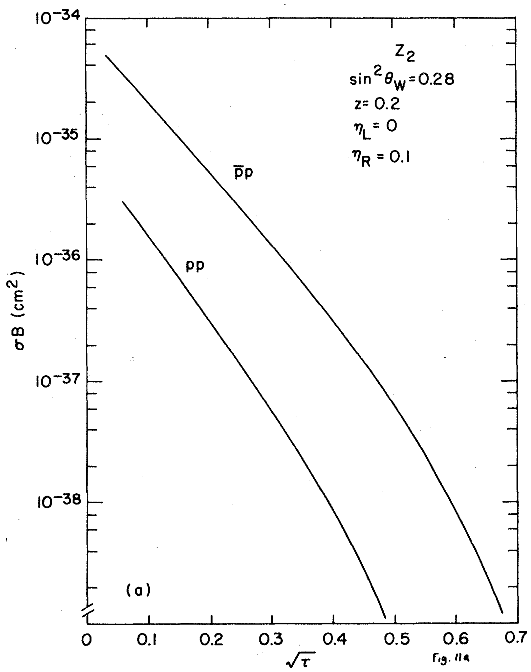


FIG. 10. (a) and (b) Same as Fig. 8 but with $\eta_R=0.4$.



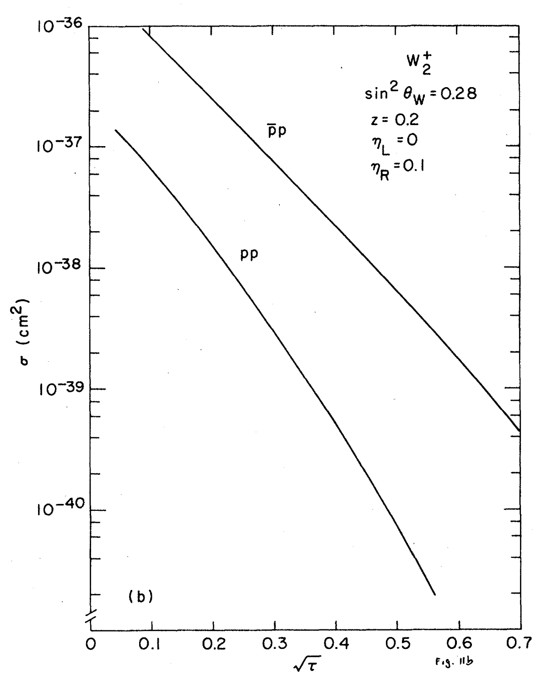


FIG. 11. (a) and (b) Same as Fig. 8 but with $z=0.2$ and $\eta_R=0.1$.

TABLE II. Predictions for gauge-boson rates (events/day) at the CERN collider ($\bar{p}p$) and ISABELLE (pp); for Z_2 decay we look for the $\mu^+\mu^-$ signature so actually σB is used.

$\sin^2\theta_W$	z / η_L / η_R	Z_2 pp	Z_2 $\bar{p}p$	$W_2^{\pm}$ pp	$W_2^{\pm}$ $\bar{p}p$
0.23	0 / 0 / 0.1	5.2×10^{-2}	10^{-6}	4.3×10^{2}	2.2×10^{-2}
0.25	0 / 0 / 0.2	3.5	5.2×10^{-3}	3.9×10^{3}	6.1×10^{-1}
0.28	0 / 0 / 0.1	4.3×10^{-2}	4.3×10^{-4}	6.9×10^{2}	6.9×10^{-2}
0.28	0 / 0 / 0.2	3.5	6.0×10^{-3}	4.8×10^{3}	24.2
0.28	0 / 0 / 0.3	17.3	4.8×10^{-2}	1.0×10^{4}	61.3
0.28	0 / 0 / 0.4	43.2	1.6×10^{-1}	1.8×10^{4}	1.1×10^{2}
0.28	0 / 0.1 / 0.5	1.2×10^{2}	5.2×10^{-1}	3.0×10^{4}	1.9×10^{2}
0.28	0.2 / 0.0 / 0.1	0.70	$<10^{-6}$	0.26	1.1×10^{-3}

have masses on the order of 100 GeV or more (as has been assumed above).

The N_i with masses in this range may have unusual decay modes; consider any N_i in the limit that mixings between the N_i as well as mixings between the usual neutrinos and the N_i can be neglected. We have the usual decays through W_1 and W_2 exchange which, in the limit $\xi\to 0$, proceeds through W_2 only. Thus for ξ small, the virtual W_2 exchange leads to the decay modes

$$N_l\to l+\sum \bar{q}q' ,\qquad N_l\to l+\sum N_{l'}l' , \tag{4.24}$$

where $N_{l'}$ labels an N lighter than N_l with decay rate characterized by the width and the sum extends over all possible contributions. The scale of the decay widths of these particles is given by (for $M_N<M_{W_2}$)

$$\Gamma_0=\frac{g^4M_N{}^5}{3.2^{11}\pi^3M_W{}^4}\left[\frac{6}{a^4}[2a-a^2+2(1-a)\ln(1-a)]-\frac{2}{a}\right], \tag{4.25}$$

where $a=M_N{}^2/M_{W_2}{}^2$. For three generations, the total semileptonic decay rate is $\simeq 9\Gamma_0$ while the purely leptonic decay depends upon which of the N_i is decaying (i.e., the N_i mass spectrum) and is either 1 or $2\Gamma_0$. If $M_N>M_{W_2}$, however, N can decay directly into a real W_2, i.e., $N\to lW_2$ with a rate given by

$$\tilde{\Gamma}=\frac{g^2}{64\pi}M_N\left(\frac{M_N}{M_{W_2}}\right)^2\left[1-3\left(\frac{M_{W_2}}{M_N}\right)^4+2\left(\frac{M_{W_2}}{M_N}\right)^6\right]\cos^2\xi . \tag{4.26}$$

This second decay mode is much more rapid due not only to the two powers of g^2 but to a numerical factor of order 10^3. Consider the possibility that the N mass in question lies in the range $M_{W_1}\le M_N\le M_{W_2}$; then N decays into a real W_1 with a rate given by (4.26) with $M_{W_2}\to M_{W_1}$ multiplied by a factor of $\sin^2\xi$. It also decays via a virtual W_2 with a rate given by Eq. (4.25); these may now competing in rate since the value of $\sin^2\xi$ may be comparable to g^2 divided by the numerical factor described above.

A detailed discussion of the N_i phenomenology is quite complicated owing to the unknown value of the N_i masses and is beyond the scope of this paper; we simply wish to point out the complicated nature of the possible decay scheme of these particles and how it depends sensitively on the N_i masses.

V. SUMMARY AND COMMENTS

Through the analysis presented in this paper, we have seen the first clue indicating the possible absence of a desert between M_W and the unification scale. We have found that the value of $\sin^2\theta_W$ need not be ~0.23, but in the case of the existence of a new mass scale not far from M_W, it could be anywhere between 0.25 and 0.31. The mass scale in question is associated with the restoration of spontaneously broken left-right symmetry, above which $SU(2)_L\times SU(2)_R\times U(1)_{B-L}$ presumably becomes a good symmetry. This picture becomes consistent with unification of all particle forces (but gravity), when one recalls that the large value of $\sin^2\theta_W$ is a necessary ingredient[11] for the existence of low intermediate mass scales in grand unified theories. This will be discussed in full detail in a subsequent publication.

Here, we have analyzed the properties of the second generation of gauge bosons, including the masses, widths, branching ratios, and production rates in pp and $p\bar{p}$ collisions. If our scenario is correct, we should be able to produce these bosons at ISABELLE energies with significant rates, thus providing a clear test of the ideas presented above. If the lower bounds on the mass of W_2 and Z_2 are saturated, as seems likely from the constraint of unification, the lighter gauge bosons W_1 and Z_1 are found to have masses 5–10 GeV smaller than in the standard model. This provides another possible probe of our alternative to the conventional picture.

The analysis makes use of the model-independent determination of neutral-current coupling constants. We have found a large family of models which satisfy all of the constraints imposed by neutral- and charged-current data.[24] It is remarkable that the predominantly charged right-handed gauge boson can be as light as 150 GeV ($\frac{3}{2}\sigma$) and even 100 GeV (2σ). If it turns out that at least one of the heavy right-handed Majorana neutrinos is lighter than W_2 then the decay $W_2 \to Nl$ provides a clean signature for its detection.

In the subsequent paper we will consider the implications of the embedding of $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ into O(10), in particular the predictions for the value of the energy scale associated with parity restoration. In addition, we discuss the detailed properties of the neutrino sector of the theory, with the special emphasis on the question of lepton-number violation and its possible observation in neutrinoless double-β decay.

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