Doubly charged Higgs bosons and lepton-number-violating processes

Thomas G. Rizzo
*Department of Physics and Ames Laboratory,
Iowa State University,
Ames, Iowa 50011*
(Received 25 September 1981)

The discovery of doubly charged Higgs bosons may signal violation of lepton flavor by two units ($\Delta L=2$) in various processes. Although somewhat exotic for the usual $SU(2)_L \otimes U(1)$ model, such Higgs bosons do occur naturally associated with Majorana neutrinos in left-right-symmetric models with low right-handed mass scales. We examine the properties of such exotic Higgs bosons and the corresponding $\Delta L=2$ reactions within a recently proposed model of this kind. Such exotic bosons could be produced at the next round of new accelerators with significant cross sections.

I. INTRODUCTION

In the usual version of the standard Weinberg-Salam-Glashow[1] gauge model (with only a single Higgs doublet responsible for symmetry breaking), the only physical Higgs particle is a neutral boson which couples to fermions, proportional to their masses, and the usual gauge bosons. These couplings are completely determined but the physical-Higgs-boson mass is unknown although loose constraints exist.[2,3] When the Higgs sector of the standard model is extended we usually end up with a model with several Higgs doublets; this is usually done to have soft *CP* violation,[4] calculable mixing angles,[5] or some horizontal-interfamily gauge symmetry.[6]

In this last case there can be a large number of physical Higgs fields remaining after spontaneous symmetry breaking which are either neutral or singly charged. The appearance of a doubly charged scalar would then appear as something exotic in the standard picture. (This is also true in hypercolor[7] theories as well.)

The purpose of this paper is to briefly examine the possible existence of doubly charged Higgs bosons in the standard model and to examine in detail the doubly charged Higgs within the context of the left-right-symmetric model[8] based on $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ gauge group with a rather low right-handed mass scale (as discussed in previous work[9]). In such a model these exotic Higgs bosons occur *naturally* from the Majorana character of the neutrino masses with rather large (of order e) couplings to ordinary charged leptons *independently* of the mass of the charged lepton. If these Higgs bosons are relatively light (~ 100 GeV or so) then $\Delta L=2$ reactions can occur at three levels—unlike the usual one-loop lepton-number-violating processes such as $\mu \rightarrow e\gamma$—and independent of any leptonic mixing angles.

Section II contains a brief summary of how exotic, doubly charged Higgs bosons may occur in the standard $SU(2)_L \otimes U(1)$ model. Section III gives a detailed account of these Higgs bosons in the $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ model with low right-handed mass scale. We examine their production properties as well as their $\Delta L=2$ effects when exchanged between leptonic currents. Section IV gives a summary discussion as well as our conclusions based on this and our earlier work.[9] We will see that the left-right-symmetric model contains doubly charged Higgs bosons associated with the Majorana nature of neutrinos. These same Higgs bosons will lead to $\Delta L=2$ effects (as well as others) which may be observable at future accelerators.

II. DOUBLY CHARGED HIGGS BOSONS IN $SU(2)_L \otimes U(1)$

As mentioned in the Introduction the existence of doubly charged Higgs scalars (or correspondingly, pseudo-Goldstone bosons in hypercolor theories) would be quite unexpected. As is well known, the current low-energy phenomenology of the standard model is quite consistent with all present experiments[10]; this includes the now famous ρ parameter which measures the ratio of, effectively, the strength of charged- and neutral-current couplings in $\bar{\nu}$ reactions. In terms of phenomenologically relevant parameters we know

$$\rho \equiv M_Z^2 \cos^2\theta_W / M_W^2 \tag{2.1}$$

and that, for a set of Higgs bosons with weak isospin t_i and whose neutral member is t_{3i} with vacuum expectation value λ_i, we have

$$\rho \equiv \frac{2\sum_i t_{3i}^2 |\lambda_i|^2}{\sum_i [t_i(t_i+1)-t_{3i}^2]|\lambda_i|^2} . \tag{2.2}$$

Experimentally, ρ is very close to unity, which is the value given by (2.2) for a Higgs doublet excluding radiative corrections[11]:

$$\rho_{\text{exp}} = 1.00 \pm 0.02 . \tag{2.3}$$

We now can imagine two possible scenarios:

(1) In addition to the usual doublet there is (at least) one additional Higgs multiplet of high dimensionality containing exotically charged scalars. Why is this Higgs multiplet present? This is beyond answer in an electroweak theory but may result as a light remnant of a large Higgs-boson representation used to break a grand unified theory.[12] We can, however, restrict what representations of Higgs bosons are allowed if we *demand* $\rho = 1$ identically (apart from radiative corrections). This question has already been answered in an earlier work of the present author[13]—we ask for what values of t and t_3 can we obtain $\rho = 1$ independent of the value of λ_i. This leads to solving the equation

$$3t_3^2 = t(t+1) . \tag{2.4}$$

Without going into detail apart from the usual solution $(t,t_3) = (\frac{1}{2}, \pm\frac{1}{2})$ there are a host of other solutions, the simplest being $(3,\pm 2)$. This representation is a 7-plet:

$$\begin{bmatrix} \phi^+ \\ \phi^0 \\ \phi^- \\ \phi^{--} \\ \cdots \end{bmatrix} \tag{2.5}$$

and contains doubly (and higher) charged Higgs bosons.

Note that the existence of this and other solutions of (2.4) will all contain exotic Higgs bosons while not interfering with the good result ($\rho = 1$) of the usual Higgs doublet alone. How would the standard-model phenomenology be altered by such a boson? Not much. The weak isospins of these larger representations which are solutions of (2.4) are such that they *never* couple directly to fermions but only to each other and the various gauge bosons. Thus, while we might have

$$\phi^{--} \to \phi^- W^- ,\quad W^- \searrow \mu^- \bar{\nu}_\mu \tag{2.6}$$

producing particles such as ϕ^{--} would be quite difficult. The best process would be in W decay:

$$W^- \to \phi^{--} W^+_{\text{virtual}} ,\quad W^+_{\text{virtual}} \searrow \text{fermion pair} \tag{2.7}$$

More about this kind of reaction will be presented below. Within this scenario we may be able to discover the existence of extra Higgs bosons without ever seeing them; this relies on the fact that the fermion couplings to the usual neutral Higgs boson is fixed in the standard model. If we have a $(\frac{1}{2}, \pm\frac{1}{2})$ and a $(3,\pm 2)$ representation as well we find G_F to be given by

$$\frac{G_F}{\sqrt{2}} = [4(\lambda_{1/2}^2 + 16\lambda_3^2)]^{-1} . \tag{2.8}$$

Since G_F is fixed, $\lambda_{1/2}^2$ must be smaller than in the standard model and since the coupling of the usual Higgs boson $(\frac{1}{2}, \pm\frac{1}{2})$ to fermions is

$$\frac{m_f}{\lambda_{1/2}} \bar{f} f \phi_{1/2} . \tag{2.9}$$

We would then expect these couplings to be larger than what is normally expected. If this difference is measurable then indirect evidence of a more complex Higgs structure is obtained.

(2) In this second scenario we would like to produce Majorana ν_L masses for the standard model[14] (ν_R is still absent). To do this, we must add to the usual doublet a Higgs multiplet of the triplet variety $(1, \pm 1)$:

$$\begin{bmatrix} \phi^0 \\ \phi^- \\ \phi^{--} \end{bmatrix} . \tag{2.10}$$

In this case ρ is not unity but depends on the ratio of the triplet to doublet vacuum expectation values:

$$\rho \equiv \frac{1+4x}{1+2x} , \quad x \equiv |\lambda_T/\lambda_D|^2 . \tag{2.11}$$

(Note $\rho \gtrsim 1$ for all values of x.) Obviously $x << 1$ in order for $\rho \simeq 1$, thus limiting the size of the vacuum expectation value of the Higgs triplet, λ_T. The existence of this triplet produces the couplings

$$\mathscr{L} = \frac{m_{\nu_i}}{\lambda_T}[\nu_{L_i}^T \nu_{L_i}(\lambda_T + \phi^0) + l_{L_i}^- l_{L_i}^- \phi^{--} + \ldots] \tag{2.12}$$

with m_{ν_i} being the Majorana neutrino mass of the ith generation (neglecting possible mixings). Note, however, even if $x \lesssim 10^{-2}$ and $m_{\nu_\tau} = 0.2$ GeV the coupling of $\tau^-\tau^-$ to ϕ^{--} is quite small $\simeq 10^{-2}$ in amplitude and $\sim 10^{-4}$ in rate. This compares quite unfavorably to, say, electromagnetic couplings which are $\sim 10^{-1}$ in rate (i.e., $e^2 = 4\pi\alpha \simeq 10^{-1}$). Thus unless x is exceptionally small ($x \simeq 10^{-3} \sim 10^{-4}$), we would not expect to see any appreciable effect of the existence of these exotic Higgs bosons. Even with the large couplings we have no idea what kind of masses to expect for the doubly charged Higgs scalar—in a Coleman-Weinberg[15] type of approach, we might imagine that $m_\phi = \simeq g^2$ (λ_T or λ_D), which means it may be quite light of order 10 GeV or so. As will be discussed below, recent PETRA data[16] can be used to place a lower limit of roughly $\simeq 15$ GeV on the mass of such an exotic Higgs boson.

It is clear from the above discussion that doubly charged Higgs scalars are very exotic from the standard-model point of view and should not be expected to occur on general grounds. In the situations where they could occur their coupling to fermions (as we have seen) are quite constrained either to be absent or small making their detection extremely difficult.

We shall see below, however, that these same doubly charged Higgs bosons occur in the $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ model in association with Majorana neutrino masses.

III. DOUBLY CHARGED HIGGS BOSONS IN $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$

Within the left-right-symmetric model there is a natural explanation of the lightness of the usual left-handed neutrinos provided they are of Majorana character.[8,9] The Higgs-boson representations in this model used for symmetry breaking are

$$\Delta_{L,R} = \begin{bmatrix} \Delta^{--} \\ \Delta^{-} \\ \Delta^{0} \end{bmatrix}_{L,R}, \quad (1,0,-2) \text{ and } (0,1,-2)\,,$$

$$\Phi = \begin{bmatrix} \chi^{0'} & \chi^{+} \\ \chi^{-} & \chi^{0} \end{bmatrix}, \quad (\tfrac{1}{2},\tfrac{1}{2},0)\,. \tag{3.1}$$

The numbers in the parentheses refer to left- and right-handed weak isospin, respectively, and the value of $B-L$. All charged fermions obtain their masses through coupling to Φ and its conjugate $\tilde{\Phi}$; the neutrinos, however, also obtain masses from coupling in a Majorana manner to Δ_L and Δ_R. The above scalars have the vacuum expectation values (VEV's)

$$\langle \Delta_{L,R} \rangle = \begin{bmatrix} 0 \\ 0 \\ v_{L,R} \end{bmatrix}, \quad \langle \Phi \rangle = \begin{bmatrix} k' & 0 \\ 0 & k \end{bmatrix}, \tag{3.2}$$

which satisfy[8] (from the potential minimum and phenomenological constraints)

$$v_L{}^2,\ k'^2 << k^2 \sim v_R{}^2\,. \tag{3.3}$$

In this limit, we would find a neutrino mass matrix like

$$\begin{array}{c c} & \begin{array}{cc} \nu_L & \nu_R \end{array} \\ \begin{array}{c} \nu_L \\ \nu_R \end{array} & \begin{bmatrix} 0 & m_l \\ m_l & M_N \end{bmatrix} \end{array} \tag{3.4}$$

with m_l the 3×3 charged-lepton mass matrix and M_N a 3×3 mass matrix for the Majorana (right-handed) neutrinos. Neglecting intergeneration mixing we find

$$m_{\nu_{Li}} = \frac{m_{l_i}{}^2}{M_{N_i}} \quad (i=1,2,3)\,,$$

$$M_{N_i} = h_R^i v_R \quad (i=1,2,3)\,. \tag{3.5}$$

(Below, N will stand for the right-handed Majorana neutrinos.) This all results from the couplings

$$\mathscr{L}=h_L^i(\nu,l)_{Li}C\tau_i\begin{bmatrix}\nu\\ l\end{bmatrix}_{Li}\Delta_L+(L\to R)$$
$$+h(\nu,l)_{Li}\begin{bmatrix}0&0\\ 0&k\end{bmatrix}\begin{bmatrix}\nu\\ l\end{bmatrix}_{Ri}+\text{H.c.}\ , \qquad (3.6)$$

where C is the charge-conjugation matrix and the h's in (3.6) are *a priori* unknown Yukawa couplings. Equation (3.5) can be rewritten as

$$m_{\nu_{Li}}\simeq K_i m_{l_i}{}^2/M_{W_R}\ ,$$
$$M_{N_i}\simeq K_i{}^{-1}M_{W_R}\ . \qquad (3.7)$$

The K_i are unknown numerical constants of order unity thus making M_{N_i} roughly degenerate with masses of $\simeq$100 GeV each.

In this order of approximation we see that the couplings of Δ_L^{--} to fermions (h_L^i) are not defined by the phenomenological masses of any of the fermions and are completely arbitrary. One might imagine, however, that they are relatively small, since the relevant mass scale is that of left-handed neutrinos. The couplings of Δ_R^{--} are, however, completely fixed (see Fig. 1):

$$\mathscr{L}=\kappa_i l_i^-(1+\gamma_5)l_i^-\Delta_R^{--},\quad \kappa_i=\frac{M_{N_i}}{2V_R}\ . \qquad (3.8)$$

Note that

$$\kappa^i\simeq 0.16\left[\frac{M_{N_i}}{100\text{ GeV}}\right] \qquad (3.9)$$

or roughly $\simeq\frac{1}{2}e$ for $M_N\simeq$100 GeV. We will concentrate on Δ_R^{--} below (hereafter called Δ), but whatever we say can be taken over to the Δ_L^{--} case as well. (We remind the reader that, in general, Δ_L^{--} and Δ_R^{--} will mix and not be mass eigenstates—we will neglect this mixing for the purposes of this discussion.)

The coupling κ^i being large, we might imagine that the interaction (3.8) can lead to $\Delta L=2$ reactions with possibly observable rates—this would depend, however, on the mass of Δ which is set by V_R (just as is W_R whose mass is $\simeq$150 GeV). We now turn to an examination of various processes in which Δ's could either be produced or have their existence felt indirectly.

The first place one might see doubly charged Higgs scalars is in e^+e^- reactions where pair production of Δ's is possible above threshold. In the one-photon-exchange limit we already know the cross section for Δ production

$$\frac{d\sigma_\Delta}{d\cos\theta}=\frac{\pi\alpha^2}{3s}Q_\Delta{}^2\left[1-\frac{4m_\Delta{}^2}{s}\right]^{3/2}(1-\cos^2\theta)\ ,$$
$$\sigma_\Delta=\frac{\pi\alpha^2}{3s}Q_\Delta{}^2\ . \qquad (3.10)$$

As in the more well-known case of singly charged scalars, we see the usual β^3 turn-on and $\sin^2\theta$ angular distribution. The new feature here is that $Q_\Delta{}^2=4$ so that as $\beta\to 1$, Δ's contribute a full unit to R. The present data from PETRA tells us only that $m_\Delta \gtrsim 15$ GeV— not an unsurprising result.

At higher e^+e^- energies we must include the contributions of both Z^0 bosons in the pair production cross sections for Δ's (see Fig. 2). The, at this point, unknown coefficients C_i (for Z_i^0) can be calculated directly from the gauge-model parameters as we will see below; note that $C_0\equiv -2\ \sin\theta_W$ since $Q_\Delta=-2$. We find

$$\frac{d\sigma_\Delta}{d\cos\theta}=\frac{g^2}{64\pi s}\beta^3(1-\cos^2\theta)\sum_{i,j}C_iC_j(v_i^ev_j^e+a_i^ea_j^e)P_iP_j\quad (i,j=0,1,2) \qquad (3.11)$$

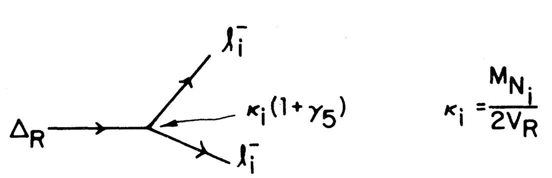

FIG. 1. Coupling of Δ^{--} to l^-l^- pairs in the left-right-symmetric model.

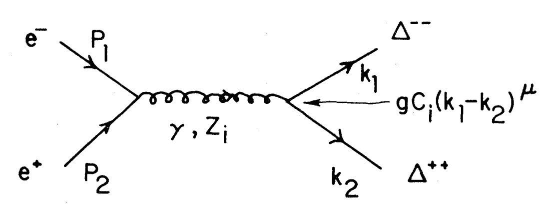

FIG. 2. Production of $\Delta\bar{\Delta}$ pairs in e^+e^- via γ, Z_1^0, and Z_2^0 exchange.

with $\beta=(1-4m_\Delta{}^2/s)^{1/2}$ and where v_i^e (a_i^e) are the vector-vector (axial) coupling of the electron to the ith Z^0 boson ($i=0$ is the photon term). The P_i's are the normalized Z^0-boson propagators:

$$P_i \equiv \frac{s}{s-M_{Z_i}{}^2} \quad (i=0,1,2) \tag{3.12}$$

(note $P_0=1$). v_i^e and a_i^e are normalized such that in the standard Weinberg-Salam (WS) model

$$v_{\rm WS}^e=\frac{e}{2\sin\theta_W\cos\theta_W}(-\tfrac{1}{2}+2\sin^2\theta_W)\ ,$$
$$a_{\rm WS}^e=\frac{-e}{4\sin\theta_W\cos\theta_W}\ . \tag{3.13}$$

Sitting on a Z_1^0 resonance we may also pair produce Δ's (Fig. 3):

$$\Gamma(Z_i\to\bar\Delta\Delta)=\frac{g^2M_{Z_i}}{48\pi}C_i{}^2(1-4m_\Delta{}^2/M_{Z_i}{}^2)^{3/2}\ . \tag{3.14}$$

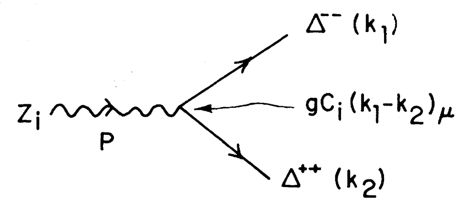

FIG. 3. The decay of Z_1^0 or Z_2^0 into $\Delta\bar\Delta$ pairs.

To determine the couplings C_i, we note that in terms of gauge eigenstates the Δ's couple to

$$\Delta_R:\ -g[W_{3R}+(g'/g)B]\ ,$$
$$\Delta_L:\ -g[W_{3L}+(g'/g)B]\ . \tag{3.15}$$

With $W_{3R(L)}$ being the third component of $SU(2)_{R(L)}$ and B the boson corresponding to $(B-L)/2$. Straightforward algebra then gives us Δ_L and Δ_R couplings:

$$gC_0(\Delta_L,\Delta_R)=-2e\ ,$$
$$gC_1(\Delta_R)=\frac{g}{\cos\theta}\left[2\cos\phi\sin^2\theta_W-\sin\phi(\cos2\theta_W)^{1/2}\left[1-\frac{\sin^2\theta_W}{\cos2\theta_W}\right]\right]\ ,$$
$$gC_2(\Delta_R)=\frac{g}{\cos\theta}\left[-2\sin\phi\sin^2\theta_W-\cos\phi(\cos2\theta_W)^{1/2}\left[1-\frac{\sin^2\theta_W}{\cos2\theta_W}\right]\right]\ , \tag{3.16}$$
$$gC_1(\Delta_L)=\frac{g}{\cos\theta}\left[\sin\phi\frac{\sin^2\theta_W}{(\cos2\theta_W)^{1/2}}-\cos2\theta_W\cos\phi\right]\ ,$$
$$gC_2(\Delta_L)=\frac{g}{\cos\theta}\left[\cos\phi\frac{\sin^2\theta_W}{(\cos2\theta_W)^{1/2}}+\cos2\theta_W\sin\phi\right]\ ,$$

with[8]

$$\tan\phi\equiv\frac{1-M_{Z_1}{}^2\cos^2\theta_W\sin^2\theta_W/2\pi\alpha(k^2+k'^2)}{(\cos2\theta_W)^{1/2}}\ . \tag{3.17}$$

Numerically, we find all the C_i are of order unity:

$$C_1(\Delta_R)\simeq0.59\ ,\quad C_1(\Delta_L)\simeq-0.48\ ,$$
$$C_2(\Delta_R)\simeq-0.40,\quad C_2(\Delta_L)\simeq0.52\ . \tag{3.18}$$

If the Δ is below 30–40 GeV in mass then sitting on the Z_1^0 many Δ's can be produced since the branching ratio into Δ pairs is of order a few percent (2–3 % for each Δ). How do we know when Δ's are produced?

The primary decay mode of the Δ is into an identical pair of charged leptons $\Delta\to l^-l^-$ as shown in Fig. 4. The decay width is given by

$$\Gamma(\Delta\to l^-l^-)=\frac{\kappa^2}{4\pi}m_\Delta\left[\frac{m_l}{m_\Delta}\right]^2\left[1-\left[\frac{4m_l{}^2}{m_\Delta{}^2}\right]\right]^{1/2} \tag{3.19}$$

with the usual helicity suppression due to chiral

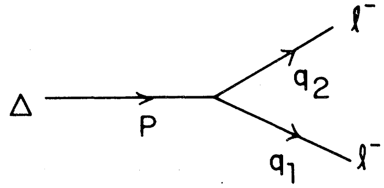

FIG. 4. Primary decay mode of the Δ, $\Delta \to l^- l^-$.

couplings. With all N's roughly degenerate Δ's will like to decay into $\tau^-\tau^-$ pairs predominantly with a rather small width ($\simeq 200$ keV for $m_\Delta = 35$ GeV). With $Z_1 \to \bar{\Delta}\Delta$ and subsequent Δ decay we would effectively see $Z_1 \to \tau^+\tau^+\tau^-\tau^-$ with *like-sign* pairs having identical invariant masses equal to the Δ mass. This would be a clear signal for lepton-number-violating couplings at the tree level.

If $2m_\Delta > M_Z$ but $M_Z > m_\Delta$ we can imagine the process shown in Fig. 5 which is also lepton-number violating. The width for this process is

$$\frac{d\Gamma}{d\chi}(Z_i \to \bar{\Delta} l^- l^-) = \frac{M_{Z_i}\alpha}{3\pi^2 \sin^2\theta_W} \kappa^2 C_i^{\,2} (x^2 - \delta)^{3/2} \times \left[\frac{1 - 2\chi + \delta_i}{(1-2\chi)^2} \right], \quad (3.20)$$

where χ is E_Δ / M_Z and $\delta_i = m_\Delta^{\,2}/M_{Z_i}^{\,2}$; χ ranges between $\delta_i^{\,1/2}$ and $\frac{1}{2}(1+\delta_i)$. The branching ratio for this process, unlike (3.14), is $\simeq 10^{-4}$ and probably unobservable unless we were using a Z^0 factory.

Besides decaying into $\tau^-\tau^-$ pairs the Δ can also decay into a lighter singly charged Higgs boson and a W as shown in Fig. 6 if $m_\Delta > m_H + M_W$. For Δ_R, some straightforward manipulations yield

$$\Gamma(\Delta_R \to H^- W_R^-) = \frac{\alpha}{4\sin^2\theta_W} \frac{m_\Delta^{\,3}}{M_R^{\,2}} \times [(1+R-H)^2 - 4R]^{1/2} \times [(1-H-R)^2 - 4HR] , \quad (3.21)$$

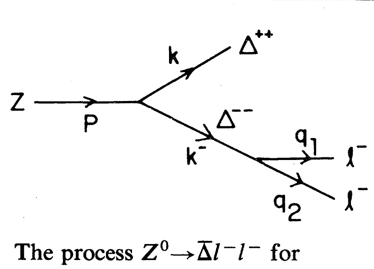

FIG. 5. The process $Z^0 \to \bar{\Delta} l^- l^-$ for $2m_\Delta > M_{Z_i} > m_\Delta$.

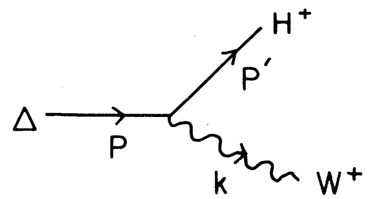

FIG. 6. Diagram for the Δ decay $\Delta_R \to H^- W_R^-$.

where

$$R = M_R^{\,2}/m_\Delta^{\,2} , \quad H = m_H^{\,2}/m_\Delta^{\,2} . \quad (3.22)$$

There are two variants of this result; first we might have $m_H + M_R > m_\Delta$ but $m_\Delta > m_H$ in which case we have the process of Fig. 7, $\Delta_R \to H + \bar{f}f$; for a given species of $\bar{f}f$ we find

$$\frac{d\Gamma}{d\chi}(\Delta_R \to H\bar{f}f) = \frac{\alpha^2 m_\Delta}{3\pi x_W^{\,2}} N_C \frac{(\chi^2 - H)}{(1 - 2\chi + H - R)^2} \quad (3.23)$$

with $N_C = 1$ if $\bar{f}f$ are leptons and $N_c = 3$ for $\bar{f}f$ being quarks. The range of χ in (3.23) is

$$H^{1/2} \le \chi \le \tfrac{1}{2}(1+H) . \quad (3.24)$$

Similarly, we might have $m_H + M_R > m_\Delta$ but $m_\Delta > M_R$ leading to the process shown in Fig. 8, $\Delta_R \to W_R \bar{f}f$; for a given $\bar{f}f$ species we find the following decay rate (with the same notation as above):

$$\frac{d\Gamma}{d\chi}(\Delta_R \to W_R \bar{f}f) = N_C \frac{m_\Delta \alpha}{2\pi^2 x_W} \frac{m_\Delta^{\,2}}{M_R} (\chi^2 - R)^{3/2} \times \frac{(1-2\chi)(a^2 - b^2)}{(1 - 2\chi + R - H)^2} \quad (3.25)$$

with

$$R^{1/2} \le \chi \le \tfrac{1}{2}(1+R) . \quad (3.26)$$

In the above, we have taken the $\bar{f}fH$ coupling to be $a + b\gamma_5$.

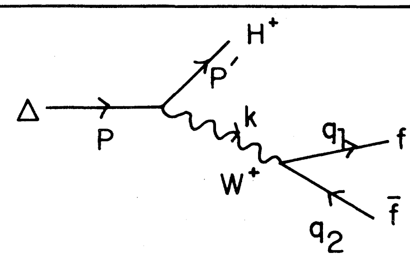

FIG. 7. The process $\Delta_R \to H f \bar{f}$ (for $M_R > m_\Delta > m_H$).

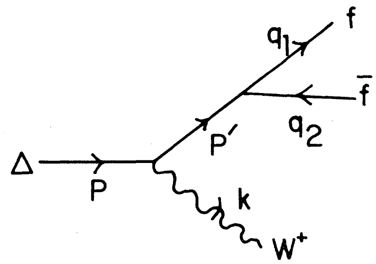

FIG. 8. Similar to Fig. 7 is the process $\Delta_R \to W_R f \bar{f}$ (for $m_H > m_\Delta > M_R$).

As a final example of the interplay between Δ's and gauge bosons we can imagine Δ production in Δ decay through a diagram like that in Fig. 9 for the Δ_R. If the $W_R W_R \Delta_R$ coupling is

$$\mathscr{L}_{eH} = \lambda M_R{}^2 W_R W_R \Delta_R \ , \tag{3.27}$$

then

$$\frac{d\Gamma}{dX}(W_R \to \Delta_R + \bar{f} f) = N_C \frac{\alpha \lambda^2 M_R{}^3}{144\pi^2 x_W}(X^2 - \delta)^{1/2} \times \frac{3 - 6X + X^2 + 2\delta}{(2X - \delta)^2} \tag{3.28}$$

with $\delta^{1/2} \le X \le \frac{1}{2}(1+\delta)$. λ is found to be given by

$$\lambda \equiv \left[\frac{4\sqrt{2} G_F \eta_R}{(1+\eta_R)^2} \right]^{1/2} \tag{3.29}$$

with $\eta_R = 0.3.$[8] (η_R is essentially $k^2/v_R{}^2$.) This leads, unfortunately, to a branching ratio which is quite small of order (of order $\simeq 10^{-9}$) and probably unobservable.

Since Δ's induce lepton-number-violating reactions perhaps this should be exploited in the search for such particles. One of the best ways to observe the effects of a Δ is to have an $e^- e^-$ collider since normally the only possible reactions must have $e^- e^-$ in the final state. To lowest order in electroweak coupling the basic expected reaction is simple elastic scattering: $e^- e^- \to e^- e^-$ (see Fig. 10). The existence of Δ allows an s-channel-exchange contribution (Fig. 11) which is, unfortunately, chirally suppressed:

$$\frac{d\sigma}{d\cos\theta}(e^- e^- \to l^- l^-) = \frac{\kappa^4}{8\pi s} \left[\frac{m_e{}^2}{s} \right] \left[\frac{m_l{}^2}{s} \right] \frac{s^2}{(s - m_\Delta{}^2)^2 + m_\Delta{}^2 \Gamma_\Delta{}^2} \ . \tag{3.30}$$

Δ exchange however allows $l^- \neq e^-$, in fact, favors $e^- e^- \to \tau^- \tau^-$. Note that if the Δ had coupled as $(a + b\gamma_5)$ we would have instead obtained

$$\frac{d\sigma}{d\cos\theta}(e^- e^- \to l^- l^-) = \frac{(a^2 - b^2)_e (a^2 - b^2)_l}{16\pi s} \frac{s^2}{(s - m_\Delta{}^2) + m_\Delta{}^2 \Gamma_\Delta{}^2} \ , \tag{3.30'}$$

which may be significant large depending on the

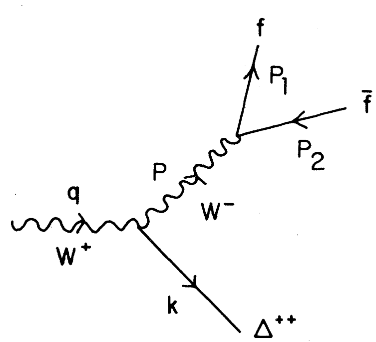

FIG. 9. Exotic decay mode of W_R, $W_R \to \Delta_R f \bar{f}$ $(M_R > m_\Delta)$.

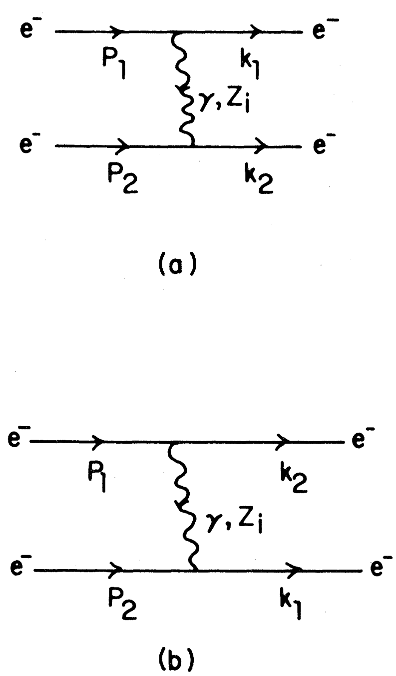

FIG. 10. The process $e^- e^- \to e^- e^-$ with γ and Z_i exchanges.

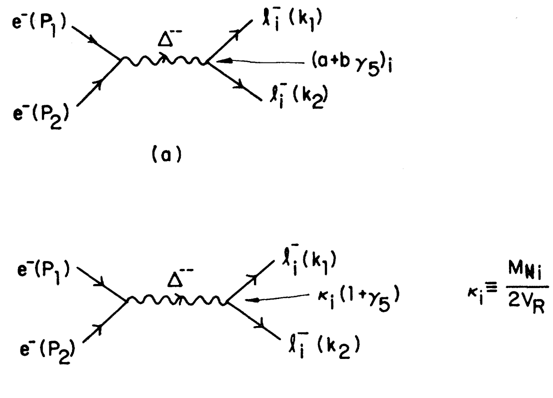

FIG. 11. The s-channel Δ contribution to $e^-e^- \rightarrow l^-l^-$.

couplings.

To get around this chiral suppression, consider the process $e^-e^- \rightarrow \Delta+\gamma$ through the digrams of Fig. 12; we will neglect Fig. 12(c) in the calculation since it is chirally supressed. We find

$$\frac{d\sigma}{d\cos\theta}(e^-e^- \rightarrow \Delta+\gamma)=\frac{\alpha\kappa^2}{s}\left[1-\frac{m_\Delta{}^2}{s}\right] \times\left[\frac{1+\cos^2\theta}{1-\cos^2\theta}\right]. \quad (3.31)$$

Above threshold this cross section has a coefficient of order 0.6 of one unit of R; the signal is quite distinct:

$$e^-e^- \rightarrow \Delta+\gamma \quad , \quad (3.32)$$
$$\searrow \tau^-\tau^-$$

i.e., a hard photon and two same-sign leptons (not e^-'s).

The only background reaction, as mentioned above, is elastic scattering $e^-e^- \rightarrow e^-e^-$ as shown in Fig. 10. This reaction has not been studied

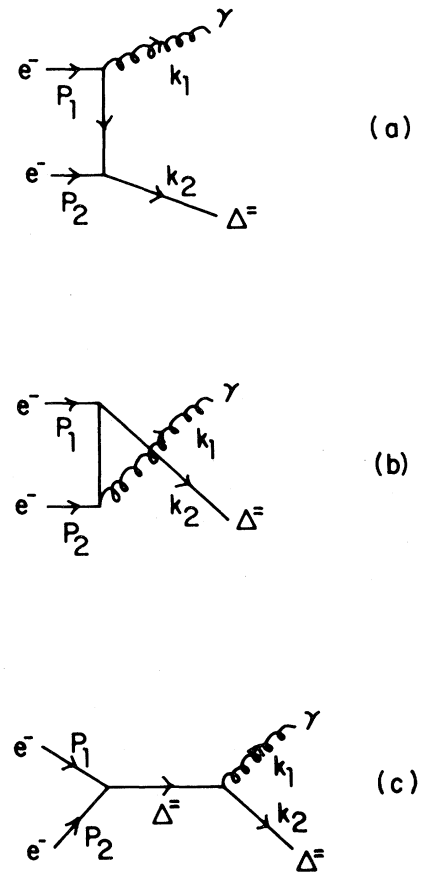

FIG. 12. The diagrams contributing to $e^-e^- \rightarrow \Delta+\gamma$.

thoroughly in the literature so we note in passing the differental cross section for this process in a general model with n Z^0 bosons ($z=\cos\theta$):

$$\frac{d}{dz}(e^-e^- \rightarrow e^-e^-)=\frac{|M|^2}{4\pi s} \ , \quad (3.33)$$

where

$$|M|^2=\sum_{i,j} P_-^i P_-^j \left\{(v_iv_j+a_ia_j)^2\left[1+\frac{(1+z)^2}{4}\right]+(v_ia_j+a_iv_j)^2\left[1-\frac{(1+z)^2}{4}\right]\right\}$$
$$+\sum_{i,j} P_+^i P_+^j \left\{(v_iv_j+a_ia_j)^2\left[1+\frac{(1-z)^2}{4}\right]+(v_ia_j+a_iv_j)^2\left[1-\frac{(1-z)^2}{4}\right]\right\}$$
$$+2\sum_{i,j} P_-^i P_+^j [(v_iv_j+a_ia_j)^2+(v_ia_j+a_ia_j)^2] \quad (3.35)$$

and

$$P_\pm^i=\left[(1\pm z)+\frac{2M_{Z_i}^2}{s}\right]^{-1} . \quad (3.35)$$

As usual v_i and a_i are the electron vector and axial-vector coupling constants for the Z_i^0 boson.

Another channel where the effects of the Δ might be expected to show up is $e^+e^- \rightarrow e^+e^-$

(Fig. 13) where Δ can occur as a t exchange. With general couplings $a+b\gamma_5$ we find

$$\frac{d\sigma}{dz}=\frac{(a^2-b^2)^2}{2\pi s}\frac{(1+z)^2}{[(1+z)+4m_\Delta{}^2/s]^2} \tag{3.36}$$

but for our chiral couplings $(a^2-b^2)^2 \to \kappa^2\,(m_e{}^2/s)^2$ and so this cross section is highly suppressed and unobservable.

Obviously, we have not completely exhausted all the possible reactions and channels in which a Δ may show some effect. We have, however, given sufficient grounds to show that if the left-right-symmetric theory with low right-handed mass scale is correct, some evidence of $\Delta L=2$ reactions should be seen in future accelerators. We note that the reaction $e^-e^-\to l^-l^-+\gamma$ ($l\neq e$) is very important in establishing the existence of the Δ since there is no background to speak of. It may be possible to collide two electron beams using the SLAC collider, hence, opening up the e^-e^- channel for exploration and possibly producing Δ's at appreciable rates.

IV. DISCUSSION AND CONCLUSIONS

As we have seen, the existence of doubly charged Higgs bosons in the standard model is quite unnatural and remote. On the other hand, the existence of both left- and right-handed Majorana neutrinos in the left-right-symmetric model demands the existence of (at least) two doubly charged Higgs bosons both with masses set by the "right-handed" vacuum expectation value v_R. The Majorana coupling structure of the fermion to these Higgs bosons naturally lead to lepton number violating processes, e.g., $e^-e^-\to\tau^-\tau^-+\gamma$ with appreciable cross sections (or branching ratios). Unlike the one- or two-loop lepton-number-violating processes in the standard model, such as $\mu\to e\gamma$, these processes are allowed at the tree level.

As discussed in our earlier work the embedding of $SU(2)_L \otimes SU(2) \otimes U(1)_{B-L}$ into SO(10) can occur in two ways: First it is possible that the right-handed mass scale is much larger then M_W ($\sim$80 GeV) (in fact, we have shown that if

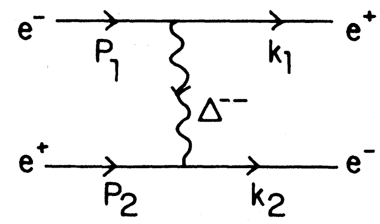

FIG. 13. t-channel Δ contribution to $e^+e^-\to e^+e^-$.

$M_R >> M_W$ then $M_R > 10^{6-7}$). If this is so, exchange of right-handed gauge bosons do not disturb the apparent low-energy weak interaction group of $SU(2)_L \otimes U(1)_Y$ with $\sin^2\theta_W \simeq 0.22$. It is also possible, and perhaps more appealing, that M_R may be comparable to M_W ($M_{W_R}\sim 2M_{W_L}$, say). In this case, the effective low-energy gauge group is left-right symmetric and $\sin^2\theta_W\simeq 0.27$. In this case there is no real distinction between left- and right-handed mass scales and so all the "new" particles of such a model (Δ's, N's, W_R, and Z_2) are relatively light (say $\lesssim$250 GeV). This greatly enriches the phenomenology above M_W (the "oasis" in the desert) which filters down into low-energy phenomenology such as $\mu\to e\gamma$ and neutrinoless double-β decay.

The rich phenomenology of a theory such as this, which might be viewed as a positive aspect, may also be its downfall. If any one of the new phenomena predicted by the theory turns out not to exist, the entire theory might have to be discarded. Our purpose then is to examine as fully as possible *all* of the possible implications of the low right-handed mass-scale scenario. The doubly charged Higgs boson and its associated lepton-number-violating reactions is just one example of the rich phenomenology implied by this model.

ACKNOWLEDGMENTS

The author would like to thank W. Marciano for encouraging a deeper look at doubly charged Higgs bosons during the ISABELLE workshop. This work was supported in part by the U. S. Department of Energy, Contract No.W-7405-ENG-82, Division of High Energy Physics and Nuclear Physics.

[1]S. Weinberg, Phys. Rev. Lett. 19, 1264 (1967); Phys. Rev. D 5, 1412 (1972); A. Salam, in *Elementary Particle Theory: Relativistic Groups and Analyticity (Nobel Symposium No. 8),* edited by N. Svartholm (Almqvist and Wiksell, Stockholm, 1968), p. 367; S. L. Glashow, J. Iliopoulos, and L. Maiani, Phys. Rev. D 2, 1285 (1970).

[2]M. Veltman, Acta Phys. Pol. B8, 475 (1978); Phys. Lett. B70, 253 (1977); B. W. Lee, C. Quigg, and H. B. Thacker, Phys. Rev. D 16, 1519 (1979); Phys. Rev. Lett. 38, 883 (1977); T. Appelquist and R. Shanker, Nucl. Phys. B158, 317 (1979); T. Appelquist and C. Bernard, Phys. Rev. D 22, 200 (1980); H. E. Haber and G. L. Kane, Nucl. Phys. B144, 525 (1978).

[3]P. J. Steinhorst, Phys. Lett. 97B, 147 (1980); A. Guth and E. J. Weinberg, Phys. Rev. Lett. 45, 1131 (1980); E. Derman, Phys. Lett. 95B, 369 (1980); H. Huffel and G. Pocsik, Z. Phys. 8, 13 (1981).

[4]For a recent review on soft *CP* violation, see G. Senjanović, *High Energy Physics—1980,* proceedings of the XX International Conference, Madison, Wisconsin, edited by L. Durand and L. G. Pondrom (AIP, New York, 1981).

[5]See, for example, R. Gatto, G. Morchio, and F. Strocchi, Phys. Lett. 83B, 348 (1979); G. Sartori, *ibid.* 82B, 255 (1979); M. Suzuki, *ibid.* 85B, 91 (1979); G. Ecker, Vienna Report No. UWThPh-80-34 (unpublished); A. C. Rothman and K. Kang, Phys. Rev. D 23, 2657 (1981); D 24, 167 (1981); T. G. Rizzo, *ibid.* 23, 181 (1981).

[6]T. Yanagida, Phys. Rev. D 22, 1826 (1980); G. Ecker, W. Grimus, and W. Konetschny, Phys. Lett. 94B, 381 (1980); C. L. Ong, Phys. Rev. D 22, 2886 (1980); A. Davidson, M. Koca, and K. C. Wali, Phys. Rev. Lett. 43, 92 (1979); F. Wilczek and A. Zee, *ibid.* 42, 421 (1979); R. D. Peccei, Max Planck Report No. MPI-PAE/PTh 28/80, 1980 (unpublished).

[7]See, for example, P. Sikivie, CERN Report Nos. CERN-TH-2951, 1980 (unpublished) and CERN-TH-3083, 1981 (unpublished); E. Fahri, and L. Susskind, Phys. Rep. 74, 277 (1981); J. Ellis, M. K. Gaillard, D. V. Nanopoulos, and P. Sikivie, Nucl. Phys. B182, 529 (1981); S. Dimopoulos, *ibid.* B182, 505 (1981); V. Elias, Phys. Rev. D 22, 2879 (1980); S. Dimopoulos, S. Raby, and G. L. Kane, Michigan Report No. UM-HE-80-22, 1980 (unpublished); G. J. Gounaris and A. Nicoloaides, Phys. Lett. 102B, 144 (1981).

[8]J. C. Pati and A. Salam, Phys. Rev. D 10, 275 (1974); R. N. Mohapatra and J. C. Pati, *ibid.* 11, 566 (1974); 2558 (1974); G. Senjanovíc and R. N. Mohapatra, *ibid.*, 12, 1502 (1975); Phys. Rev. Lett. 44, 912 (1980).

[9]T. G. Rizzo and G. Senjanovíc, Phys. Rev. Lett. 46, 1315 (1981); Phys. Rev. D 24, 704 (1981); 25, 235 (1982); Brookhaven Report No. BNL-29716, 1981 (unpublished).

[10]See, for example, the recent review by J. E. Kim, P. Langacker, M. Levine, and H. H. Williams, Rev. Mod. Phys. 53, 211 (1981).

[11]See, for example, W. Marciano and A. Sirlin, Phys. Rev. D 22, 2695 (1980).

[12]For a review of grand unified theories, see P. Langacker, Phys. Rep. 72, 185 (1981).

[13]T. G. Rizzo, Phys. Rev. D 21, 1404 (1980).

[14]T. P. Cheng and L. F. Li, Phys. Rev. D 22, 2860 (1980).

[15]S. Coleman and E. J. Weinberg, Phys. Rev. D 7, 1888 (1973).

[16]For a recent review of PETRA results see P. Dunker, DESY Report No. DESY 81-012, 1981 (unpublished).

[17]See, however, R. Gastman and Y. Van Ham, Phys. Rev. D 10, 3629 (1974); C. H. Llewellyn Smith and D. V. Nanopoulos, Nucl. Phys. B48, 207 (1974).