PRODUCTION AND DETECTION AT SSC OF HIGGS BOSONS IN LEFT-RIGHT SYMMETRIC THEORIES

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Summary

We discuss the production and detection at SSC of charged and neutral Higgs bosons of the left-right symmetric theories. The H^+, which is largely a member of a left-right "bidoublet," should be detectable. The H_2^o, a more unusual Higgs particle which, apart from mixing, is in a right-handed triplet and does not couple to quarks, may be detectable too.

I. Introduction

The left-right symmetric (LR) theories[1,2] are an attractive example of the models that predict physics which is beyond the standard model and which may be accessible at the SSC. The LR theories are appealing for a number of reasons: (i) they restore parity to the status of an exact short distance symmetry of the weak interactions,[1] and thus provide a more aesthetic description of electroweak phenomena as well as a renormalizable framework for describing the origin of parity violation; (ii) they incorporate full quark-lepton symmetry of weak interactions and give the U(1) generator of electroweak symmetry a new meaning in terms of the B-L quantum number,[3] and (iii) they lead to a natural explanation of the smallness of neutrino masses[2] by relating this smallness to the observed suppression of V+A currents in low-energy weak processes. These theories contain two W bosons, W_1 and W_2. The W_1 is the already-discovered W (82 GeV), which couples mostly to left-handed (LH) currents. The W_2, which must be much heavier than the W_1, couples mainly to RH currents. The theories also contain two Z bosons, Z_1 and Z_2, the lighter of which, Z_1, is the familiar Z (93 GeV). In the fermion sector, they contain the usual quarks and charged leptons, three light neutrino mass eigenstates, ν_1, ν_2, and ν_3, and three heavy neutrino mass eigenstates, N_1, N_2, and N_3. The ν_i, which apart from mixing effects are the familiar neutrinos ν_e, ν_μ, and ν_τ, couple predominantly to W_1, while the much heavier N_i couple largely to W_2. Finally, these theories contain a distinctive Higgs sector which includes neutral, singly-charged, and doubly-charged physical Higgs particles.

Low-energy tests of the LR models, involving such reactions as $\mu \to 3e$, $\mu \to e\gamma$, $\mu^-e^+ \to e^-\mu^+$, and K decay, have been proposed.[4] However, due to their expected large masses and their production mechanisms, the new particles, such as the W_2, which are characteristic of these models may not be observable directly until we have the very high energies which will be produced by the SSC. The possibility of searching for the W_2, the Z_2, and the N_i at the SSC has been explored previously.[5] Here, we focus on the Higgs sector of the LR models, and examine the possibility of detecting, with the aid of the SSC, some of the predicted physical Higgs particles.

In the remainder of this Introduction, we describe two illustrative versions of the basic LR model, giving special emphasis to the Higgs sector, and introducing the Higgs particles whose detectability at SSC we shall then consider. In Sec. II, we present the relevant couplings of these Higgs bosons to fermions, to gauge bosons, and to other Higgs bosons. Using these couplings, we estimate the branching fractions for the Higgs particle decays into final states which are detectable and distinctive. In Sec. III, we discuss the mechanisms and cross sections for the production of the LR Higgs bosons at the SSC. Drawing on the branching fractions from Sec. II, we then determine the event rates for Higgs production followed by decay to the especially advantageous final states. Backgrounds for these states are considered. Our conclusions concerning the detectability of the LR Higgs particles are presented. In Sec. III, we also consider the degree to which LR Higgs bosons could be mimicked by Higgs particles which are unrelated to the LR-symmetric models. Our overall conclusions are presented in a final section. (A more complete, detailed account of our analysis and its conclusions will be given in a separate publication.)

Left-Right Symmetric Model -- Two Versions

The left-right symmetric models are based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, with the quarks and leptons assigned to multiplets with quantum numbers as follows (here $Q \equiv \binom{u}{d}$, $\psi \equiv \binom{\nu}{e^-}$, and the quantum numbers are indicated in the order (I_L, I_R, B-L):

$$Q_L: \left(\frac{1}{2}, 0, \frac{1}{3}\right) ; \quad Q_R: \left(0, \frac{1}{2}, \frac{1}{3}\right)$$
$$\phi_L: \left(\frac{1}{2}, 0, -1\right) ; \quad \phi_R: \left(0, \frac{1}{2}, -1\right) \tag{1.1}$$

There are two free gauge couplings in this model: $g_L = g_R \equiv g$ for the SU(2) group and g' for the $U(1)_{B-L}$ group. The electric charge formula for this model is[3]

$$Q = I_{3L} + I_{3R} + \frac{B-L}{2} . \tag{1.2}$$

The minimal Higgs sector that leads to the symmetry breaking pattern,

$$SU(2)_L \times SU(2)_R \times U(1)_{B-L} \xrightarrow[m_R]{} SU(2)_L \times U(1)_Y \xrightarrow[m_W]{} U(1)_{EM} , \tag{1.3}$$

consists of the multiplets[2]

$$\phi \equiv \left(\frac{1}{2}, \frac{1}{2}, 0\right)$$

and

$$\Delta_L \equiv (1,0,2) \; + \; \Delta_R \equiv (0,1,2) . \tag{1.4}$$

Let us display the various components of the multiplets below:

$$\phi \equiv \begin{pmatrix} \phi_1^o & \phi_1^+ \\ \phi_2^- & \phi_2^o \end{pmatrix} ,$$

$$\Delta_L \equiv \begin{pmatrix} \Delta_L^+/\sqrt{2} & \Delta_L^{++} \\ \Delta_L^o & -\Delta_L^+/\sqrt{2} \end{pmatrix} ,$$

$$\Delta_R \equiv \begin{pmatrix} \Delta_R^+/\sqrt{2} & \Delta_R^{++} \\ \Delta_R^o & -\Delta_R^+/\sqrt{2} \end{pmatrix} . \tag{1.5}$$

The Higgs potential involving these multiplets has been studied in detail in Ref. 2, where it has been noted that there exists a range of parameters for which the minimum of the potential corresponds to the following vacuum expectation values for ϕ , Δ_L and Δ_R :

$$\langle\phi\rangle = \begin{pmatrix} \kappa & 0 \\ 0 & \kappa' \end{pmatrix}$$

$$\langle\Delta_L\rangle = \begin{pmatrix} 0 & 0 \\ v_L & 0 \end{pmatrix}$$

$$\langle\Delta_R\rangle = \begin{pmatrix} 0 & 0 \\ v_R & 0 \end{pmatrix} . \tag{1.6}$$

In terms of these vacuum expectation values (vevs) and the gauge couplings, the masses of the gauge bosons are given by

$$m_{W_1}^2 \simeq \frac{1}{2} g^2(\kappa^2 + \kappa'^2 + 2v_L^2) ,$$
$$m_{W_2}^2 \simeq \frac{1}{2} g^2(\kappa^2 + \kappa'^2 + 2v_R^2) ,$$
$$m_{Z_1}^2 \simeq \frac{g^2}{2\cos^2\theta_W} (\kappa^2 + \kappa'^2 + 4v_L^2) ,$$
$$m_{Z_2}^2 \simeq 2(g^2 + g'^2)\, v_R^2 . \tag{1.7}$$

Here, θ_W is the conventional Weinberg mixing angle, which is related to the gauge couplings of the LR model by

$$\tan \theta_W = g'/(g^2 + g'^2)^{1/2} . \tag{1.8}$$

The bosons W_1 and W_2 are mixtures of two bosons, W_L and W_R, which couple, respectively, to purely LH and RH currents. The mixing angle, ζ , is given by

$$\tan \zeta \simeq \frac{\kappa\kappa'}{v_R^2} . \tag{1.9}$$

Low-energy physics does not constrain the relative magnitudes of the vevs κ and κ' which enter symmetrically in Eqs. (1.7) and (1.9), although κ and κ' cannot be equal or the up and down quark mass matrices would be equal, contrary to observation. Thus, it is natural to make the simplifying assumption that one of these vevs dominates the other. In that case, the question of which one dominates is purely a matter of nomenclature. Therefore, we choose κ to be the dominant one, and set $\kappa' = 0$.

In order to have near maximal parity violation at low energies, as observed, we must have $\kappa \ll v_R$. Furthermore, the vevs are related by $v_L = \gamma\kappa^2/v_R$, where γ is a typical coupling parameter. Thus, $v_L \ll \kappa$, and we shall neglect it. The boson W_1 then receives its mass from the vev κ , W_2 receives its much larger mass mainly from v_R , and the W_L-W_R mixing angle ζ is zero for $\kappa' = 0$. In addition, as we see from Eq. (1.7), Z_2 is necessarily heavier than W_2. Since analysis of the K_L-K_S mass difference implies that $m_{W_2} \gtrsim 1.6$ TeV,[6] this means that $m_{Z_2} > 1.6$ TeV as well. This is the minimal version of the LR model.

The phenomenological constraint on m_{Z_2} is much weaker than that on m_{W_2}, requiring only that $m_{Z_2} \gtrsim 275$ GeV.[7] Therefore we shall also consider a somewhat modified version of the LR model in which the Higgs sector has been slightly enlarged so as to make it possible for Z_2 to be much lighter than W_2. In this version, one breaks $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ in two stages:

$$SU(2)_L \times SU(2)_R \times U(1)_{B-L} \xrightarrow[m_{W_R}]{} SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L} \xrightarrow{M_{B-L}} SU(2)_L \times U(1)_Y$$

The new Higgs multiplets that we add to ϕ and $\Delta_{L,R}$ in order to implement this scheme are

$$\Sigma_L \equiv \begin{pmatrix} \sigma_L^o & \sqrt{2}\sigma_L^+ \\ \sqrt{2}\sigma_L^- & -\sigma_L^o \end{pmatrix} ,$$

and

$$\Sigma_R \equiv \begin{pmatrix} \sigma_R^o & \sqrt{2}\sigma_R^+ \\ \sqrt{2}\sigma_R^- & -\sigma_R^o \end{pmatrix} . \tag{1.10}$$

The multiplets Σ_L and Σ_R are triplets with the quantum numbers (1,0,0) and (0,1,0), respectively. We assign to them the vevs

$$\langle\Sigma_L\rangle = 0 \quad \text{and} \quad \langle\Sigma_R\rangle = \begin{pmatrix} \sigma_R & 0 \\ 0 & -\sigma_R \end{pmatrix} , \tag{1.11}$$

with $\sigma_R \gg v_R \gg \kappa$. All the formulae of the minimal version of the model remain unchanged except for those which give the mass of W_2 and the left-right mixing angle ζ . These now become

$$m_{W_2}^2 \simeq \frac{1}{2} g^2(\kappa^2 + \kappa'^2 + 2v_R^2 + 2\sigma_R^2)$$

and

$$\tan \zeta \simeq \frac{\kappa\kappa'}{v_R^2 + \sigma_R^2} . \tag{1.12}$$

The Z_2 continues to obtain its mass from the vev v_R as in the minimal model, but the W_2 now derives most of its mass from the much larger vev σ_R . Thus, we can now have $m_{Z_2} \ll m_{W_2}$.

II. The Physical Higgs Bosons, their Couplings, and their Decays

In this section we wish to display the physical Higgs bosons of the model, their couplings, and their consequent decay patterns. Let us first consider the minimal version of the model. To simplify the discussion, we set $\kappa' = v_L = 0$. In this limit, the unphysical Higgs bosons that become the longitudinal components of the massive gauge bosons are given by $(\kappa \ll v_R)$:

$$W_2^{long} = \Delta_R^+ - \frac{\kappa}{\sqrt{2}v_R}\phi_1^+ ,$$

$$W_1^{long} = \phi_2^+ ,$$

$$Z_2^{long} = \text{Im}\Delta_R^o - \frac{\kappa}{v_R}\frac{\cos 2\theta_W}{2\sin^2\theta_W} \cdot \text{Im}\phi_1^o ,$$

$$Z_1^{long} = \text{Im}\phi_1^o + \frac{\kappa}{2v_R}\frac{\cos 2\theta_W}{\sin^2\theta_W}\text{Im}\Delta_R^o . \tag{2.1}$$

The remaining Higgs bosons are the ten physical Higgs particles. They are the quanta of the fields

$$H^+ = \phi_1^+ + \frac{\kappa}{\sqrt{2}v_R}\Delta_R^+ ,$$

$$H_1^o = \sqrt{2}\left[\cos\theta^o\ \text{Re}\phi_1^o + \sin\theta^o\ \text{Re}\Delta_R^o\right] ,$$

$$H_2^o = \sqrt{2}\left[-\sin\theta^o\ \text{Re}\phi_1^o + \cos\theta^o\ \text{Re}\Delta_R^o\right] ,$$

$$\Delta_R^{++}, \Delta_L^{++}, \Delta_L^+, \sqrt{2}\text{Re}\Delta_L^o, \sqrt{2}\text{Im}\Delta_L^o, \sqrt{2}\text{Re}\phi_2^o, \sqrt{2}\text{Im}\phi_2^o \tag{2.2}$$

Here the $\text{Re}\phi_1^o$-$\text{Re}\Delta_R^o$ mixing angle θ^o is given by

$$\tan 2\theta^o = \frac{\alpha_H\kappa v_R}{\rho_H v_R^2 - \lambda_H\kappa^2} , \tag{2.3}$$

in which α_H, ρ_H, and λ_H are linear combinations of the scalar boson self-couplings in the model. The masses of the Higgs bosons depend on the symmetry scale as well as on the scalar self couplings. The latter being rather arbitrary, we cannot predict the masses of the physical Higgs bosons but will instead assume various illustrative values in our computations.

Our analysis will focus on the H^+ and the H_2^o . The H^+ should not be too difficult to produce and detect at the SSC. The H_2^o will be less easy to observe, but will be quite distinctive. The remaining Higgs particles appear to be difficult to study at the SSC. (When $\sin\theta^o$ is small, the H_1^o is the LR-model analogue of the neutral Higgs boson of the standard model, whose detection at the SSC has been discussed in great detail by a number of authors.[8] The only difference here is that H_1^o can decay to a pair of heavy neutrinos N through the $\sin\theta^o$ mixing term in Eq. (2.2).)

The couplings of the physical Higgs scalars to the gauge bosons are dictated by the covariant derivatives of the Higgs fields and the kinetic energy term. These couplings have been carefully analyzed by F. Olness, and A. Mendez, and are listed in Table I.

TABLE I. Couplings of the Higgs particles to gauge bosons.

Combination of fields in Lagrangian	Coefficient
$H^+W_2^-Z_1$	$\frac{-g^2\kappa}{\sqrt{2}\cos\theta_W}$
$H^+W_1^-Z_1$	$\frac{-g^2\kappa^2\kappa'}{v_R^2\sqrt{\cos\theta_W}} \simeq 0$
$H^+W_1^-H_1^o$	$\simeq 0$
$H_2^oZ_1Z_1$	$-\sin\theta^o\ \frac{g^2\kappa}{4}\frac{1}{\cos^2\theta_W}$
$H_2^oZ_2Z_2$	$\cos\theta^o\sqrt{2}g^2v_R\frac{\cos^2\theta_W}{\cos 2\theta_W}$
$\Delta_L^{++}W_1^-W_1^-$	$g^2v_L \simeq 0$
$\Delta_R^{++}W_2^-W_2^-$	g^2v_R

$H_1^oW_1W_1$	$g^2\cos\theta^o\kappa/2\sqrt{2}$
$H_1^oZ_2Z_2$	$\sin\theta^o\sqrt{2}\, g^2 v_R \frac{\cos^2\theta_W}{\cos 2\theta_W} +$
	$\cos\theta^o \frac{g^2\kappa}{2\sqrt{2}} \frac{\cos 2\theta_W}{\cos^2\theta_W}$
$H_1^oZ_1Z_1$	$g^2\kappa\, \cos\theta^o/\sqrt{2}\, \cos\theta_W$
$H_1^oW_2W_2$	$\frac{g^2}{2\sqrt{2}}\left(\cos\theta^o\kappa + \sin\theta^o v_R\right)$

We turn next to the couplings of the Higgs bosons to fermions. These arise from the gauge invariant Yukawa interaction

$$\mathcal{L}_Y = h_{ij}^u \bar{Q}_{iL}\phi Q_{jR} + h_{ij}^d \bar{Q}_{iL}\tilde{\phi}\, Q_{jR} + h_{ij}^\nu \bar{\psi}_{iL}\phi\psi_{jR} + h_{ij}^\ell \bar{\psi}_{iL}\tilde{\phi}\psi_{jR} + h_{ij}^M\left(\psi_{iL}^T C\tau_2\Delta_L\psi_{jL} + \psi_{iR}^T C\tau_2\Delta_R\psi_{jR}\right) + \text{h.c.} \quad (2.4)$$

In this expression, the h_{ij} are couplings constants, i and j run over the three generations, $\tilde{\phi} \equiv \tau_2\phi^*\tau_2$, and C is the Dirac charge-conjugation matrix. We can estimate the coupling constants using the fact that h_Y is the source of the fermion masses. Indeed, from Eqs. (1.6) and (2.4) we find (when $\kappa' = 0$) that

$$h^u = \kappa^{-1} M^u, \quad h^d = \kappa^{-1} M^d,$$
$$h^\nu = \kappa^{-1} M^{\nu^D}, \quad h^\ell = \kappa^{-1} M^\ell. \quad (2.5)$$

Here h^u is the matrix of couplings h_{ij}^u, etc., and the matrices M^u, M^d, M^{ν^D}, and M^ℓ are. respectively, the up-quark mass matrix, the down-quark mass matrix, the Dirac part of the neutrino mass matrix, and the charged-lepton mass matrix. Neglecting mixing, M^u, M^d, and M^ℓ are diagonal matrices whose diagonal elements are the quark and charged lepton masses. Considerations based on SO(10) suggest that $M^{\nu^D} \simeq \frac{1}{3} M^u$, and we shall use this relation to estimate h^ν. To estimate h^M, we recall that after symmetry breaking, the h^M and h^ν terms in h_Y lead to the neutrino mass matrix

$$M_i = \begin{array}{c|cc} & \nu & N \\ \nu & 2h_{ii}^M v_L & h_{ii}^\nu\kappa \\ N & h_{ii}^\nu\kappa & 2h_{ii}^M v_R \end{array}. \quad (2.6)$$

Here we are neglecting mixing between the generations, so that Eq. (2.6) gives the mass matrix for the light and the heavy neutrino in the i'th generation. Neglecting v_L, it follows from this matrix that the two neutrino mass eigenstates in this generation have masses[9]

$$m_{N_i} \simeq 2h_{ii}^M v_R, \quad (2.7a)$$

and

$$m_{N_i} \simeq m_{\nu_i D}^2/m_{N_i}. \quad (2.7b)$$

In obtaining the "see-saw relation," Eq. (2.7b), we have used the relation $h_{ii}^\nu\kappa = m_{\nu_i D}$ in Eq. (2.5).

Equation (2.7a) relates the coupling h_{ii}^M to the (unknown) masses of N_i and of W_2 or Z_2. We can obtain lower bounds on the m_{N_i} from the see-saw relation combined with an estimate of the $m_{\nu_i D}$ and with the experimental upper bounds on the m_{ν_i}: $m_{\nu_e} < 18$ eV,[10] $m_{\nu_\mu} < 250$ keV,[11] and $m_{\nu_\tau} < 70$ MeV.[11] Depending on whether we take $m_{\nu_i D} = \frac{1}{3} m_{u_i}$, as suggested by SO(10), or $m_{\nu_i D} = m_{\ell_i}$, we obtain the following bounds:

	$m_{\nu_i D} = m_{\ell_i}$		$m_{\nu_i D} = \frac{1}{3} m_{u_i}$	
$m_{N_e} >$	14 GeV	or	200 GeV	
$m_{N_\mu} >$	45 GeV	or	1 TeV	
$m_{N_\tau} >$	45 GeV	or	?	(2.8)

It may also be worth pointing out[12] that vacuum stability arguments lead to an upper bound $m_N \lesssim m_{W_R}$.

The dominant fermion-fermion-Higgs boson couplings which are of interest to us have been obtained from the $\mathcal{L}_Y$ of Eq. (2.4) and its relation to the fermion masses, and are given in Table II. In constructing this table, we have used the relation $\kappa^{-1} = 2^{3/4} G_F^{1/2}$ which follows from Eq. (1.7).

TABLE II. Couplings of the Higgs particles to fermions. The quantity V^R is the right-handed quark mixing matrix, which in accord with LR symmetry we assume to be equal to its left-handed analogue, the ordinary Kobayashi-Maskawa matrix.

Combination of fields in Lagrangian	Coefficient
$H^+\bar{t}_L b_R$	$2^{3/4}G_F^{1/2}m_t$
$H^+\bar{t}_R b_L$	$-2^{3/4}G_F^{1/2}m_b$
$H^+\bar{t}_L s_R$	$2^{3/4}G_F^{1/2}m_t V_{ts}^R$

$H^+\bar{c}_L b_R$	$2^{3/4}G_F^{1/2}m_c V^R_{cb}$
$H^+\bar{c}_R b_L$	$-2^{3/4}G_F^{1/2}m_b V^R_{cb}$
$H^+\bar{c}_L s_R$	$2^{3/4}G_F^{1/2}m_c$
$H^+\bar{c}_R s_L$	$-2^{3/4}G_F^{1/2}m_s$
$H^+\bar{c}_L d_R$	$2^{3/4}G_F^{1/2}m_c V^R_{cd}$
$H^+\bar{\nu}_{\tau L}\tau_R$	$2^{3/4}G_F^{1/2}\frac{m_t}{3}$
$H^+\bar{N}_{\tau R}\,\tau_L$	$-2^{3/4}G_F^{1/2}m_\tau$
$H_2^o\bar{N}_\ell\,N_\ell$	$\cos\theta^o\,\frac{m_{N_\ell}}{2\sqrt{2}\,v_R}$
$\Delta^{++}_{R,L}\,\ell^T_{R,L}\,C\ell_{R,L}$	$\frac{m_{N_\ell}}{2v_R}$

From Table II, we see that when $m_{H^+} \gg m_t + m_b$, the relative strengths of the fermionic decays of the H^+, including color factors, are as given in Table III.

TABLE III. Relative strengths of fermionic decays of a heavy H^+.

Mode	Relative strength
$t\bar{b}$	$3(m_t^2 + m_b^2)$
$t\bar{s}$	$3m_t^2(V^R_{ts})^2$
$c\bar{s}$	$3(m_c^2 + m_s^2)$
$c\bar{d}$	$3m_c^2(V^R_{cd})^2$
$c\bar{b}$	$3(m_b^2 + m_c^2)(V^R_{cb})^2$
$\tau\nu_\tau$	$(m_t/3)^2$
τN_τ	m_τ^2

Finally, we turn to the couplings of Higgs particles to other Higgs particles. The couplings of this type which will be important to us are those which can contribute significantly to the decays of the H_2^o. In the limit of no $\Delta_R^o - \phi_1^o$ mixing, these couplings involve the Δ_R^o. Imposing a discrete symmetry to simplify the Higgs potential,[13] we are left with three terms that will contribute to Δ_R^o decays:

$$\alpha_1 \mathrm{Tr}\phi^\dagger\phi\,\mathrm{Tr}\Delta_R^\dagger\Delta_R + \alpha_2 \mathrm{Tr}\Delta_R^\dagger\phi^\dagger\phi\Delta_R + \alpha_2' \mathrm{Tr}\Delta_R^\dagger\tilde{\phi}^\dagger\tilde{\phi}\Delta_R \ . \quad (2.9)$$

These lead (in the absence of mixing) to the couplings given in Table IV for the decays $H_2^o \to H^+H^-$ and $H_2^o \to H_1^oH_1^o$.

TABLE IV. Couplings of H_2^o to other Higgs bosons.

Combination of fields in Lagrangian	Coefficient
$H_2^oH^+H^-$	$\sqrt{2}\,(\alpha_1 + \alpha_2)\,v_R$
$H_2^oH_1^oH_1^o$	$\frac{1}{\sqrt{2}}(\alpha_1 + \alpha_2')\,v_R$

We shall not consider the additional physical Higgs bosons which occur in the enlarged version of the LR model which includes the $\Sigma_{L,R}$ Higgs triplets. However, we note that when we are dealing with that version of the model, where $m_{Z_2} \ll m_{W_2}$, we must relate the vacuum expectation value v_R in the tables for the various couplings to m_{Z_2}, rather than to m_{W_2}, using Eq. (1.7).

In order to discuss H_2^o production by Z_2Z_2 fusion, and to discuss the decay $H_2^o \to Z_2Z_2$ followed by Z_2 decay, we will lead the neutral current couplings of the Z_2. Neglecting Z_1-Z_2 mixing and κ/v_R, these are given by

$$Z_2 = \frac{g}{\cos\theta_W}\sin\theta_W Z_2^\mu \sum_f \bar{f}\gamma_\mu\left(\alpha I_{3R} - \frac{1}{2\alpha}(B-L)\right) f \ . \quad (2.10)$$

Here f runs over all chiral fermions, and $\alpha = \sqrt{\cos 2\theta_W}/\sin\theta_W$.

Decay Rates and Branching Fractions

We see from Eq. (2.4) that the Higgs particles $\Delta^{++}_{R,L}$ have a very distinctive feature: they do not couple to quarks. This property and their interesting decays to e^+e^+, $\mu^+\mu^+$, etc., make them very special indeed. Unfortunately, the absence of a coupling to quarks also makes them very difficult to produce in pp collisions.

We see further from Eq. (2.4) that, in the absence of Δ_R^o-ϕ_1^o mixing, the H_2^o $(\sim\Delta_R^o)$ also does not couple to quarks. This particle has a distinctive decay to $N_\ell N_\ell$. Despite the absence of a coupling to quarks, it could perhaps be produced in pp collisions via Z_2Z_2 fusion if the Z_2 is light. Alternatively, if Δ_R^o and ϕ_1^o do mix appreciably, then the H_2^o could perhaps be produced via the coupling of its ϕ_1^o component to quarks. Of course, when there is large mixing, the H_2^o can decay to quarks as well, losing some of its unique character.

If we assume $m_{H_2^o} \lesssim 1$ TeV, then H_2^o cannot decay into Z_2Z_2 in the minimal LR model where $m_{Z_2} > m_{W_2}$. Its major decays are then to H^+H^-, $H_1^oH_1^o$, and N_iN_i. From the couplings given previously, we find for these modes the decay rates

$$\Gamma(H_2^o \to H^+H^-) = \frac{(\alpha_1 + \alpha_2)^2}{8\pi} \cos^2\theta^o \frac{v_R^2}{m_{H_2^o}} \left[1 - \frac{4m_{H^+}^2}{m_{H_2^o}^2}\right]^{1/2} \qquad (2.11)$$

$$\Gamma(H_2^o \to H_1^oH_1^o) = \frac{(\alpha_1 + \alpha_2')^2}{16\pi} \cos^2\theta^o \frac{v_R^2}{m_{H_2^o}} \left[1 - \frac{4m_{H_1^o}^2}{m_{H_2^o}^2}\right]^{1/2} ,$$

and

$$\Gamma(H_2^o \to N_iN_i) = \frac{1}{32\pi} \cos^2\theta \left(\frac{m_{N_i}}{v_R}\right)^2 m_{H_2^o} \left(1 - \frac{4m_{N_i}^2}{m_{H_2^o}^2}\right)^{3/2} .$$

When $m_{Z_2} > m_{W_2}$ and $(m_{N_i}/v_R)^2 \ll 1$, the decays into Higgs final states will dominate if the couplings α_1, α_2, α_2' are not too small.

In the extended model where Z_2 can be light, H_2^o decay into Z_2Z_2 becomes possible. From the coupling in Table I, we find for the corresponding decay rate

$$\Gamma(H_2^o \to Z_2Z_2) = \frac{g^2}{8\pi} \cos^2\theta^o \frac{\cos^2\theta_W}{\cos 2\theta_W} \frac{m_{Z_2}^2}{m_{H_2^o}} \left(1 - \frac{4m_{Z_2}^2}{m_{H_2^o}^2}\right)^{1/2} \times$$

$$\times \left\{2 + \frac{m_{H_2^o}^4}{4m_{Z_2}^4} \left(1 - \frac{2m_{Z_2}^2}{m_{H_2^o}^2}\right)^2\right\} . \qquad (2.12)$$

When Z_2 is light, this can be an important decay.

The dominant decays of H^+ are into fermion pairs. From Table III, we see that the decay into $t\bar{b}$ dominates all else if $m_{H^+} \gg m_t + m_b$. As we shall discuss, detection of the H^+ may be easiest through observation of its decay into $\tau\ \nu_\tau$. From Table III, we see that the branching fraction for $H^+ \to \tau\ \nu_\tau$ is actually independent of unknown couplings and masses (given our assumption that $M^{\nu^D} \simeq \frac{1}{3} M^u$), and is given by

$$\frac{\Gamma(H^+ \to \tau\ \nu_\tau)}{\Gamma_{tot}(H^+)} \simeq \frac{1}{28} \simeq 3\% . \qquad (2.13)$$

If H^+ is too light to decay into $t\bar{b}$, but heavy enough to decay into $c\bar{b}$ or lighter pairs, the $\tau\nu_\tau$ decay will dominate. So long as $m_t > 30$ GeV, we will have

$$\frac{\Gamma(H^+ \to \tau\ \nu_\tau)}{\Gamma_{tot}(H^+)} \gtrsim 90\% . \qquad (2.14)$$

III. Phenomenology

A. Charged Higgs $H^\pm$:

We begin by discussing the production and detection of $H^\pm$ assuming it to be much lighter than W_2. In this case, the decay $W_2 \to H^+Z_1$ would be approximately 1/4 the branching ratio to a typical $\ell\nu$ mode. This implies, for instance that for $m_{W_2} = 2$ TeV, with inclusive cross-section of about 1 pb, we would obtain about 12 H^+Z_1 events in which the Z_1 is detected in its e^+e^- or $\mu^+\mu^-$ mode.

Turning now to more direct $H^\pm$ production and decay, we note that if $m_{H^+} < m_t$, then $t\bar{t}$ production followed by t-decay would be a copious source of charged Higgs. For the couplings listed in Table II, (assuming $m_{W_1} < m_t$)

$$\frac{\Gamma(t \to H^+b)}{\Gamma(t \to W_1^+b)} = \frac{p_{H^+}}{p_{W_1^+}} \frac{m_t^2(m_t^2-m_{H^+}^2)}{(m_t^2 + 2m_{W_1}^2)(m_t^2-m_{W_1}^2)} \qquad (3.1)$$

where p_{H^+} and $p_{W_1^+}$ are the rest frame momenta of the respective decays. Obviously, the H^+-decay mode is fully competitive with the W_1^+-mode. In fact, if $m_{H^+} < m_t < m_{W_1}$, then the t-quark will decay almost entirely into H^+. The H^+ and H^- coming from t and $\bar{t}$ in turn decay almost 100% of the time into $\tau\nu_\tau$ final states. The final state signature would consist of 2 jets $+ \tau^+ + \tau^- + p_T^{miss}$. The τ's could be observed in their single charged particle modes. Because of the strong production rate it is difficult to imagine competitive backgrounds, especially if a τ vertex trigger is available. For instance, if $m_t = 0.1$ TeV and the $t \to H^+b$ branching ratio is .5, then the cross section times branching ratio for the above final state is $\gtrsim 10^4$ pb.

A more probable scenario, however, is one in which the charged Higgs is heavier than the top. The primary production mechanism is that of bt fusion.[14] This mechanism is studied in ref. 14. Without going into detail let us summarize by saying that a lower bound to the bt fusion production of $H^\pm$ can be obtained by computing the cross section for the $2 \to 2$ process

$$bg \to H^\pm t , \qquad (3.2)$$

where the top mass effects are retained. Over the top and Higgs mass range we shall consider the true cross section can be up to a factor of 3 larger. Results for the $2 \to 2$ process are given in fig. 1. Note that the cross section is quite substantial, especially if the top is heavy. We must now consider whether a charged Higgs produced in this manner can be detected.

In the range being considered the $H^\pm$ may decay to both the bt and $\tau\nu_\tau$ channels. From eqn. (2.13), we find that

$$BR(H^{\pm} \to \tau\nu_\tau) \gtrsim 0.03$$

$$BR(H^{\pm} \to bt) \lesssim 0.97 \ , \qquad (3.3)$$

depending upon the precise relative $H^{\pm}$ and t masses. Note that the $\tau\nu_\tau$ BR is considerably larger than that expected in a simple two-doublet version of the standard model, although such a large value for this branching ratio can be approached in supersymmetric two-doublet models.[20] Let us examine this "rare" mode first.

In ref. 14 the observability of a charged Higgs in its $\tau\nu_\tau$ mode, given a large branching ratio, such as eqn. (3.3) is considered. We summarize the conclusions here. We imagine searching for the τ in one of its single charged particle decay modes: $\tau \to e\nu\nu$, $\tau \to \mu\nu\nu$, $\tau \to \pi\nu$, or $\tau \to \rho\nu_\tau$, with combined branching ratio of ≈ 0.67. There are two critical ingredients in overcoming backgrounds.

1. First, in order to use the $e\nu\nu$ and $\nu\nu\nu$ modes, final states containing energetic τ's must be separable, presumably via vertex detection. Otherwise, backgrounds from $W \to e\nu, \mu\nu$ will generally swamp spectra from $H^{\pm}$ decay, and only the $\pi\nu$ and $\rho\nu$ modes of τ decay would be useable, with consequent loss of effective event rate.

2. Secondly, a trigger on the spectator t quark produced in association with the $H^{\pm}$ from reaction (3.2) must be implemented. This is accomplished by triggering on the secondary leptons coming fromthe t decay. Using a p_T cut of order 10 GeV reduced standard model backgrounds by a factor of order 70 while retaining approximately 45% of the charged Higgs signal.[14] Thus a net improvement of signal/background by a factor of 30 is possible.

Let us examine two special cases in more detail. First we imagine that $m_{H^{\pm}} \approx m_W$, and $m_T = 40$ GeV. The cross section for $H^{\pm}$ production via reaction (3.2) is $\gtrsim 300$ pb. In comparison, the cross section for single W production is of order 10^5 pb. The branching ratio for $W \to \tau\nu_\tau$ decay is of order .08. However, the W effective event rate is reduced via the stiff-lepton associated-t-quark trigger, discussed above, by a factor of 70. Thus in a standard SSC operating year of 10^4 pb^{-1} we obtain 10^6 events from the W background. The $H^{\pm}$ event number, after including the roughly 50% efficiency of the stiff lepton trigger, and the branching ratio (3.3), is of order 5×10^4 -- a statistically significant effect. However, the W cross section must be accurately normalized using other W decay channels in order to see a 5% effect in the $\tau\nu_\tau$ channel. Coupling this enhancement with differences in the spectra of the τ's from W vs. $H^{\pm}$ decay might make detection of such an $H^{\pm}$ possible.

The second scenario we consider is that of $m_{H^{\pm}} = 1$ TeV. The cross section from Fig. 1 is $\gtrsim 0.3$ pb, at $m_t = 40$ GeV, and could be above 1 pb if the top is substantially heavier. Combined with the branching ratio (3.3) and the 50% stiff-lepton trigger efficiency, we obtain a yearly event rate for $H^{\pm}$ production followed by $\tau\nu_\tau$ decay of $\gtrsim 50$. This is reduced by a factor of $\approx 2/3$ if we imagine looking only at single charged particle τ decay modes. Again, the primary background is from SM W production followed by (off-shell) decay to the $\tau\nu_\tau$ final state. THe spectra of the single charged particles from τ decay arising from this background were studied in ref. 15, as a function of p_T of the charged decay product. Integrating over the region of $p_T \geq 200$ GeV (which includes roughly 40% of the π^+ signal), and summing over the e, μ, π and ρ single particle channels, yields a W background yearly event rate of order 700. However, this is before imposing the stiff-trigger requirement which reduces this event rate by another factor of 70 to about 10. In all we have a ratio of 14/10 for signal/background, for the worst case $m_t = 40$ GeV choice. Relaxing the p_T cut on the charged decay products increases the number of events from the background more rapidly than for the $H^{\pm}$ signal, but might be advantageous. In practice the $p_T \leq 100$ GeV region would be used to normalize the background, and the region above $p_T = 100$ GeV would be examined for the $H^{\pm}$ enhancement. Further study is required to decide whether a very heavy $H^{\pm}$ can be detected in this way.

At intermediate masses, the conclusions are generally more favorable than for either of the above extreme cases. The background from W production drops more rapidly with increasing $m_{H^{\pm}}$ than does the H^+ signal rate. In this range the enhancement in the e, μ, π and ρ single particle decay spectra, coming from $H^{\pm}$ production and decay is statistically significant as well as sizeable on a percentage basis.

Before turning to the $H^{\pm} \to bt$ search mode, we make several remarks. First, we remind the reader that all $H^{\pm}$ cross sections used above were lower bounds and that the true cross sections will probably turn out to be more than a factor of 2 larger. Secondly, we note that if $m_t \geq m_{H^{\pm}}$ then the $H^{\pm} \to \tau\nu_T$ branching ratio is of order 97%. Detection of a singly produced $H^{\pm}$ would then be straightforward.

Searches for $H^{\pm}$ in the tb decay mode will encounter enormous backgrounds from QCD 2-jet production. For instance, at $m_{H^{\pm}} = .5$ TeV we have, from fig. 1, an $H^{\pm}$ cross section of $\gtrsim 3$ pb. In comparison the two jet cross section at this same jet-jet invariant mass is of order $d\sigma/dM_{jj} \approx 2 \times 10^3$ pb/GeV. For a mass resolution of 15% we obtain an effective cross section of 1.5×10^5 pb. Of this total, approximately 2% are gt or $g\bar{t}$ final states. If we imagine that a highly selective top quark jet trigger can be constructed, without sacrificing the 15% mass resolution (a somewhat questionable assumption given the results of ref. 10), then our effective background is of order 3×10^3 pb, some 1000 times larger than the signal. No further gain is possible using the stiff-lepton trigger on the t quark produced in association with the $H^{\pm}$, since gt production also occurs with an associated spectator t quark. Thus we would need to discriminate g jets from b jets at the level of 1/1000 in order to detect $H^{\pm}$ in the bt mode. No technique for differentiation has yet achieved such a factor.

A final question concerns distinguishing the $H^{\pm}$ of a left-right symmetric model from one that appears in two-doublet versions of the standard model and supersymmetric versions thereof. If $m_t > m_{H^{\pm}}$ then the very large branching ratio for $H^{\pm} \to \tau\nu_T$ is a sure signal for a left-right symmetric model, and is directly related to the Dirac mass matrix of the neutrinos in such theories. However, if the tb decay mode is allowed, the $H^{\pm} \to \tau\nu_T$ branching ratio is not especially different from that which occurs in some supersymmetric extensions of the standard model,[14] though it is much larger than would be obtained in simple two-doublet versions of the SM. In general, the $H^{\pm}$ cannot alone be used to determine the nature of the new physics sector.

B. Charged Triplet Members -- Δ_R^{++}, Δ_L^{+} and Δ_L^{++} .

Without a doubt, these are the most unique members of the Higgs menagerie of the left-right symmetric model. A careful review of the couplings given earlier shows that they cannot be singly produced via gauge boson fusion, in the absence of W_L-W_R mixing, in the limit where m_{W_R} is large and $v_L = 0$. They also cannot be produced by gg fusion since they do not couple to quarks. The only production possibility appears to be via Drell-Yan type processes, in which pairs of these particles are made. Given the many parameters in the Lagrangian, we find that these charged Higgs could be either light or heavy. Presumably, however, they are too heavy to be made at PEP and PETRA. Unfortunately, Drell-Yan pair production drops rapidly with increasing mass. The cross section for a doubly charged pair is the same as for a l^+l^- pair as a function of the mass of the produced particle. The latter cross sections can be found in EHLQ.[8] From the graphs there it is clear that only charged Higgs with masses less than about 100 GeV will be produced at a rate which exceeds 10 pair events in an SSC year of 10^4 pb^{-1}. However, for the doubly-charged Higgs pairs the signature is extraordinary; the only allowed decays of a doubly-charged triplet are to like-sign lepton pairs. The final state is fully reconstructable and 10 events should suffice for discovery. If the charged-triplet Higgs are heavier than 100 GeV they become extremely difficult to produce and will probably not be detected except possibly as decay products of neutral Higgs, as we shall discuss shortly.

C. Neutral Higgs

As discussed in sec. II, the neutral Higgs bosons H_1^o and H_2^o are linear superpositions of $\mathrm{Re}\phi_1^o$ and $\mathrm{Re}\Delta_R^o$, with mixing angle given by θ^o . We will consider three limiting cases (a) $\theta = 0$ or $\pi/2$ and (b) $\theta^o = \pi/4$. In general, we have for the $\mathrm{Re}\phi_1^o$-$\mathrm{Re}\Delta_R^o$ mass matrix

$$M^2 = \begin{pmatrix} 2\lambda_H\kappa^2 & -\alpha_H v_R\kappa \\ -\alpha_H v_R\kappa & 2\rho_H v_R^2 \end{pmatrix} \tag{3.4}$$

leading to

$$\tan 2\theta_o = \alpha_H v_R\kappa/(\rho_H v_R^2 - \lambda_H\kappa^2) \ . \tag{3.5}$$

Production

As is well known, the dominant production mechanisms for massive neutral Higgs are:

1. gg fusion via a quark loop; and

2. gauge boson fusion, including in the present case W_1W_1, Z_1Z_1, Z_1Z_2, and Z_2Z_2.

We discuss these in turn.

Regarding gg fusion, we recall from earlier sections that group quantum number considerations forbid couplings between the Δ_R^o and quarks. This implies that in the limit where $\theta^o = 0(\theta^o = \pi/2)$ H_2^o (H_1^o) cannot be produced via gg fusion, while H_1^o (H_2^o) production via gg fusion will be the same as in the SM. In the case where $\lambda_H = 0.5(\lambda_{max} + \lambda_{min})$, both H_1^o and H_2^o couple to quarks with $1/\sqrt{2}$ of the normal coupling and have gg fusion production cross sections that are 1/2 of the values for a SM Higgs. We should also keep in mind that the size of the gg fusion cross section is very dependent upon m_t, becoming dominant over WW/ZZ fusion, in SM Higgs production, for Higgs masses below 1 TeV if $m_t \gtrsim 150$ GeV.[15] At $m_t = 40$ GeV gg fusion falls below WW/ZZ fusion at about $m_H = 0.3$ TeV. Thus the role of gg fusion in H_1^o and H_2^o production depends critically upon m_t,λ_H , and the masses $m_{H_1^o}$ and $m_{H_2^o}$. We will consider several special situations later.

Let us turn next to the gauge boson fusion contributions. In case (a) the cross section for H_1^o production receives contributions from W_1W_1, Z_1Z_1, and Z_2Z_2 fusion (see Table I), since $\theta^o = 0$. In this limit the left-hand sector gauge particles couple with SM strength, and there is an additional similar coupling to Z_2Z_2. All contributions are added coherently, including interference effects where appropriate. (Interference is neglected in the effective W,Z approximation, and in all the calculations in the literature.) The present calculation obtains cross sections by numerical integration of exact matrix elements. The Z_2 couplings to quarks given earlier are used. Structure functions are the updated EHLQ ones.[17] The results for the H_1^o cross section in case (a) are presented in fig. 2 for two different m_{Z_2} values. We see that the Z_2Z_2 contribution is small, compared to the SM-like contributions, even for m_{Z_2} as low as 0.2 TeV. The cross section is about a factor of 2 to 3 lower than that given in ref. 14 and a factor of 2 lower than the exact (without interference) results of ref. 18. We have not attempted to trace this difference.

These results for the dominantly ϕ_1^o Higgs boson can be compared to those in fig. 3 for H_1^o production. In the limit of case (a) only Z_2Z_2 fusion contributes. We note that for $m_{Z_2} = 0.2$ TeV the H_2^o cross section is only a factor of 2 to 3 below the gauge boson fusion cross section for H_1^o at equivalent mass values. However, as m_{Z_2} rises to more probable values the H_2^o cross section vanishes as $1/m_{Z_2}^2$.

Because of the symmetry between the $\theta^o = 0$ and $\theta^o = \pi/2$ limits it is apparent that the results of figs. 2 and 3 apply regardless of whether it is the dominantly Δ_R^o or dominantly ϕ_1^o Higgs that is most massive -- fig. 2 should always be used for ϕ_1^o and fig. 3 for Δ_R^o , whenever mixing is very small. With this in mind, it is clear that there is a scenario in which we will be unable to see either neutral Higgs boson at the SSC, even if one is fairly light. This occurs when the dominantly Δ_R^o Higgs is light, but m_{Z_2}

is large, while the dominantly ϕ_1^o boson is heavier than a few TeV. Then the cross sections for both are probably beyond the range of observability, though, of course, one is entering the range of parameter space where the vector boson sector becomes strongly interacting. We shall return later to consideration of less extreme cases.

Turning now to case (b) it is clear from Table I that the Z_2Z_2 coupling differentiates between H_1^o and H_2^o. (Recall that, in this "maximal mixing" case, H_1^o always refers to the lighter of the two Higgs bosons.) For the positive α_H sector we shall consider, the H_1^o cross section is always larger than that for H_2^o, at equivalent mass values. This along with the m_{Z_2} dependence of the cross sections is illustrated in fig. 4. Note that by $m_{Z_2} = 0.6$ TeV the Z_2Z_2 fusion contributions have become negligible, and the H_1^o and H_2^o cross sections have become equal to one-half that of a SM Higgs. This is true, not only for these gauge-boson fusion contributions, but also for the gg fusion contributions discussed previously.

We are now in a position to turn to the decays and signatures for these neutral Higgs bosons.

Decays and Signatures

Case a:

Once again we shall consider the special cases (a) and (b). Focusing first on case (a), in which there is no mixing between ϕ_1^o and Δ_R^o, it is clear that all the usual standard model decay widths for the dominantly ϕ_1^o will apply. However, there are two new modes of potential importance,

$$\phi^o \rightarrow \left\{ \begin{matrix} Z_2Z_2 \\ H^+H^- \end{matrix} \right. . \tag{3.6}$$

Decay modes of ϕ_1^o to charged and neutral triplet members vanish in the mixing limit. Of the two modes in (3.6), the Z_2Z_2 decay width is given by $\Gamma(\phi_1^o \rightarrow Z_2Z_2) = (\cos 2\theta_W)^2\Gamma(\phi_1^o \rightarrow Z_1Z_1)$, in the absence of significant phase space suppression. The coupling responsible for the second is specified by $2\lambda_H(2v_v^2/(2v_v^2+\kappa^2))\phi_1^oH^+H^-$. Inserting the appropriate value of λ_H for case (a) we find $\Gamma(\phi_1^o \rightarrow H^+H^-) \approx \Gamma(\phi_1^o \rightarrow Z_1Z_1)$, when $v_R >> \kappa$; in this case phase space suppression will certainly be negligible.

While both these modes are interesting, it is also clear that the scenarios developed for discovering the standard model Higgs apply, as well, to the dominantly ϕ_1^o Higgs being considered. Various surveys of SM Higgs physics are available.[10] The only difference between the SM scenario, and that for the dominantly ϕ_1^o Higgs in the left-right symmetric model with small mixing, in addition to the above extra decay channels, is the slight enhancement of the gauge boson fusion cross section, illustrated in fig. 2, for a light Z_2. Hopefully, discovery of the ϕ_1^o using SM-like techniques and modes would be possible, and then the above two exotic decays could be used to signal the left-right symmetric nature of the underlying model.

The most important question in case (a) is whether or not the dominantly Δ_R^o Higgs boson can be observed. From fig. 3 it is apparent that cross sections are small unless both Δ_R^o and Z_2 are rather light on the scale of m_{W_2}. In the small mixing limit the possible Δ_R^o decay channels are rather restricted. They include

$$\Delta_R^o \rightarrow \left\{ \begin{matrix} Z_2Z_2 \\ N_iN_i \end{matrix} \right. \tag{3.7}$$

and a variety of Higgs pair decay modes,

$$\Delta_R^o \longrightarrow \begin{matrix} H^+H^- \\ \Delta_R^{++}\Delta_R^{--} \\ \Delta_L^o\Delta_L^o \\ \Delta_L^+\Delta_L^- \\ \Delta_L^{++}\Delta_L^{--} \end{matrix} . \tag{3.8}$$

We have discarded the $\phi_2^{o}{}'\phi_2^{o}{}'$ decay mode, since the ϕ_2^o must be quite heavy to avoid flavor changing neutral currents. (By $\phi_2^{o}{}'\phi_2^{o}{}'$ we mean real part of one times the imaginary part of the other, in order to avoid conflict with bose statistics and parity.) Of the decays in eqn. (3.8), the H^+H^- mode involves a coupling proportional to the same (or similar) lagrangian parameters that are forced to be small in the small mixing limit, and a coupling that is suppressed by the factor $\kappa^2/(2v_R^2+\kappa^2)$ compared to the other modes. We will not consider it further. The other Higgs decay channels are possibly forbidden, since the masses of these Higgs are crudely of order $M^2 \sim \rho_H' v_R^2$, where the parameters ρ_H' are closely related to ρ_H which determines the mass of Δ_R^o in the small mixing limit. Thus all these Higgs could have very similar mass to Δ_R^o and the pair decays not allowed. In this case only the N_iN_i and Z_2Z_2 channels might be open. We pursue this possibility first.

Because of the large $\Delta_R^oZ_2Z_2$ coupling of table I, and the large natural widths associated with the longitudinal modes of a vector boson pair decay channel, the Z_2Z_2 mode, if allowed, will dominate all other channels (neglecting the Higgs-pair decays). The Z_2 decay must then be considered. It cannot decay to $W_1^+W_1^-$ in the absence of Z_2-Z_1 mixing. The $W_2^+W_2^-$ channel is forbidden by the very high W_2 mass. Higgs pair channels are, we shall assume, forbidden by phase space, except for the H^+H^- mode (recall that $H^\pm$ is light if α_H-like parameters are small). Thus the Z_2 will decay primarily to available fermion-antifermion pairs according to the couplings given earlier, including the N_iN_i mode. The branching ratio for an interesting overall Δ_R^o mode must be combined with the appropriate cross section read from fig. 3. A particularly amusing case would be where one $Z_2 \rightarrow e^+e^-$

$\cdot\mu^+\mu^-$ so that we can use mass reconstruction to eliminate Z_1 backgrounds, while the other decays to N_1N_1. Half the time the N_1N_1 pair will produce leptons of equal charge in their decays. The net branching ratio for such final states is about .75%. For cross sections from fig. 3 above about 0.5 pb there would be more than 40 such events in an SSC year. Given that even a very heavy SM Higgs can be detected in the mode Z_1Z_1 where one $Z_1 \to \ell^+\ell^-$ and the other $Z_1 \to \nu\bar{\nu}$,[20] using transverse mass reconstruction and a similar number of events, it would seem that the present more fully reconstructable mode would allow Δ_R^o discovery. Of course, other modes may also be accessible, such as the mixed $Z_2 \to \ell^+\ell^- - Z_2 \to q\bar{q}$ leptonic-hadronic decay. This mode encounters the backgrounds from mixed QCD-Electroweak processes discussed in ref. 21. Its branching ratio of 4% is, however, somewhat larger than that of a SM Higgs due to the absence of W-pair decay modes for Δ_R^o. It too is likely to be marginally viable above the 0.5 pb level of fig. 3.

A more likely and exciting scenario is one in which the Z_2 is heavy enough that only the N_1N_1 decay mode of Δ_R^o is allowed, neglecting for the moment the Higgs pair modes. A large fraction of the time the final state would then contain exactly two energetic same sign leptons plus quark jets. The major background would likely arise from single real and virtual Z_2 production followed by decay to N_1N_1. We plan to examine this in the near future.

The entire situation changes if the modes of eqn. (3.8) are allowed. There are even choices of parameters in the Lagrangian for which ρ_H' is zero and all the Higgs listed are massless. Strictly speaking, this is ruled out for charged pairs by the PEP-PETRA data, as discussed earlier. But the mass limits are not large and for the present discussion we neglect phase space suppression and other corrections due to finite masses. In this case, the zero mixing limit determines all couplings and we find

$$\Gamma(\Delta_R^o \to \Delta_L^o\Delta_L^o) = \Gamma(\Delta_R^o \to \Delta_L^+\Delta_L^-) = \Gamma(\Delta_R^o \to \Delta_L^{++}\Delta_L^{--})$$

$$= 4\Gamma(\Delta_R^o \to \Delta_R^{++}\Delta_R^{--}) = \Gamma(\Delta_R^o \to Z_2Z_2) \qquad (3.9)$$

when the Z_2Z_2 decay is not phase space suppressed. Clearly the Δ_R^o would have to be found in the Higgs pair decays. There are several amusing possibilities. For instance, the Δ_L^{--} and Δ_R^{--} prefer to decay almost entirely into $\ell^-\ell^-$ pairs, and similarly for the charge conjugates. This produces a unique and fully reconstructable final state for $\ell = e,\mu$ with large branching ratio. Of course, as discussed earlier, the pair states $\Delta_L^{--}\Delta_L^{++}$ and $\Delta_R^{--}\Delta_R^{++}$ are also made via Drell-Yan pair production. However, the present Δ_R^o decay could yield a larger pair cross section, and would, in any case, cause peaking in the pair mass spectrum.

Case b:

The maximal mixing case (b) is quite different from the small mixing case. Both H_1^o and H_2^o will decay, in the mass region being considered, largely to $W_1^+W_1^+$, Z_1Z_1 and possibly Z_2Z_2 and Z_2Z_1 pairs, with widths to the first two channels being one-half those of the SM Higgs. Because of the v_R term in the H_1^o and H_2^o couplings to Z_2Z_2, this channel can also be quite large when allowed. Finally, charged and neutral Higgs pair decay modes are possible. The modes are the same as listed in eqn. (3.8), and, in addition, we could have the decay $H_2^o \to H_1^oH_1^o$. If phase space allowed, these modes will, as in the case of Δ_R^o decay, be very substantial. The overall width, when all channels are allowed, is often larger than for a SM Higgs.

To illustrate the possibilities, we first focus on the channels other than the Higgs-pair modes. We have plotted in fig. 5 the decay widths for H_1^o and H_2^o into the relevant channels, exhibiting m_{Z_2} dependence where appropriate. In fig. 5a we present those decay widths that are independent of Higgs type (1 or 2). The Z_1Z_1, W_1W_1 and $t\bar{t}$ widths, presented there, are also independent of m_{Z_2}, whereas the N_1N_1 width decreases with increasing m_{Z_2} as $1/m_{Z_2}^2$. The $t\bar{t}$ and N_1N_1 widths do, of course, grow with increasing m_t and m_{N_1}, respectively, until phase space suppression sets in. Note that, unless m_{N_1} is larger than the 70 GeV choice of fig. 5a, and the Z_2 is quite light, the N_1N_1 mode will have very small branching ratio. In fig. 5b we present the combined widths of the Z_2Z_1 and Higgs states. As expected, the importance of these modes decreases very rapidly as m_{Z_2} increases. We note that all the results in figs. 5a and 5b depend only on the mass of the Higgs whose width is presented, and not on the mass of the other Higgs as well. From these graphs it is apparent that it would be beneficial to employ the Z_2 modes in Higgs searches, if m_{Z_2} is small. We have already described the important features and branching ratios for the Z_2Z_2 mode; the Z_1Z_2 mode, which is generally small in comparison to Z_2Z_2 when either is significant, can be analyzed similarly. If the Higgs pair channels are absent, it would seem that both H_1^o and H_2^o can be found via the usual techniques[21] and obvious extrapolations thereof.

In fig. 6 we turn to the Higgs pair widths. We have adopted the Lagrangian parameter choice in which all the Higgs in eqn. (3.4) have small mass, and determined all remaining parameters as functions of $m_{H_1^o}$ and $m_{H_2^o}$. For a different organization of Lagrangian parameters, however, all Higgs pair modes could be phase space forbidden. In this case, the widths of the remaining modes are not altered, being determined by the H_1^o and H_2^o masses and the maximal mixing λ_H choice. In order to present the widths graphically, it is convenient to choose some fixed value for the mass of the Higgs whose width is not being plotted. Results do depend upon this choice. We have arbitrarily taken $m_{H_1^o} = 0.1$ TeV when considering H_2^o and $m_{H_2^o} = 1.1$ TeV when considering H_1^o. We also combine the widths of $\Delta_L^o\Delta_L^o$, $\Delta_L^+\Delta_L^-$ and $\Delta_L^{++}\Delta_L^{--}$, since they are all equal. We denote the combined channels by $\Delta_L\Delta_L$. In fig. 6a we exhibit the $H_2^o \to H_1^oH_1^o$ decay width. It is large when allowed. In fig. 6b we exhibit the H^+H^- widths. This width is always big for the H_1^o and becomes increasingly important for H_2^o , as the H_2^o mass increases. The $\Delta_R^{++}\Delta_R^{--}$ and $\Delta_L\Delta_L$ widths of figs. 6c and 6d are very similar in behavior. They are very large for H_1^o at m_{Z_2} = 0.2 TeV but decrease rapidly with increasing m_{Z_2}.

From the results presented in fig. 6, it is clear that if any of the Higgs pair channels are allowed by phase space considerable reassessment of Higgs discovery techniques is required. We discuss here only two examples. Consider first the case of H_1^o when m_{Z_2} is of order 0.2 TeV. The combined branching ratio for $\Delta_R^{++}\Delta_R^{--} + \Delta_L^{++}\Delta_L^{--}$ is larger than 0.1 for all $m_{H_1^o} \lesssim 1$ TeV, and approaches 0.5 at the lower $m_{H_1^o}$ values. Both final states have the unique double leptonic decays discussed earlier, in which we have two like sign leptons in one hemisphere and two leptons of the opposite sign in the other hemisphere. This final state is fully reconstructable and would make H_1^o discovery trivial. In another extreme m_{Z_2} is very large, but $m_{H^\pm}$ is small; the only Higgs-pair mode of importance is then H^+H^-. It has branching ratio BR $\gtrsim$ 0.85 in H_1^o decay for $m_{H_1^o} \lesssim 0.5$ GeV. The search for all but a fairly heavy H_1^o would be impossible unless $H^\pm$ were lighter than the top and decayed only into $\tau\nu_\tau$. It is remarkable how much less severely the H_2^o decays are impacted by the Higgs-pair modes. Only if the Z_2 is very light is there significant dimunition of the Z and W modes of the H_2^o , by $H_2^o \rightarrow H_1^oH_1^o$ decay. Overall, it seems that both H_1^o and H_2^o discovery should prove possible in the maximally mixed scenario, aside from the special H_1^o case described above. However, in the large m_{Z_2} limit, the major decay channels of the H_1^o and H_2^o are not necessarily dissimilar from those in a two doublet version of the standard model. In this case, the other Higgs of the left-right symmetric models might provide the first evidence as to the nature of the underlying theory.

D. The ϕ_2^o and Δ_L^o Higgs

We consider first the Δ_L^o . Strictly speaking, it has both real and imaginary degrees of freedom; but we lump these together for purposes of simplifying the discussion. The Δ_L^o , being a pure triplet member, has no quark-antiquark couplings, and thus cannot be produced via gg fusion. Since it is unmixed, quantum number considerations forbid couplings to charged gauge boson pairs. Since $v_L \approx 0$, its couplings to neutral gauge boson pairs are also very suppressed compared to those of the SM Higgs. Thus the Δ_L^o is extremely difficult to produce in the limit of $\kappa' = 0$ that we are considering here. The primary mechanism is pair production via neutral gauge boson fusion. The cross section will obviously be small compared to those discussed for the mixing sector. In addition, this type of final state is difficult to analyze with respect to backgrounds. Thus we shall not consider it further here.

The ϕ_2^o has quark-anti-quark couplings. Up to an undetermined CKM-type matrix they are completely determined by the quark and lepton masses in the standard way. Thus gg fusion could provide a substantial production cross section. The gauge boson pair couplings to ϕ_2^o are zero since $\kappa' = 0$, and, as for Δ_L^o , gauge boson fusion can only produce ϕ_2^o in pairs, with consequent small cross section levels. Thus only single ϕ_2^o production via gg fusion is potentially of interest. The ϕ_2^o will decay primarily to the maximally coupled fermion-antifermion pair(s) allowed by phase space. The couplings are analogous to those in the charged Higgs sector, implying that $t\bar{t}$ and $\tau^+\tau^-$ pairs are likely to dominate. In the former mode QCD background will probably prevent its observation. The latter mode has much larger branching ratio than in the SM and could prove interesting. However, all this is irrelevant if the CKM matrix referred to above is such as to allow FCNC transitions. If it is analogous to the standard CKM matrix then the mass of ϕ_2^o must be above about 3-10 TeV, i.e. outside the SSC range.

IV. Conclusion

In summary, we have looked at the detection possibility for the Higgs bosons of the left-right symmetric model. The most promising case for production appears to be for the singly charged boson H^+ whose decay via the $\tau\nu_\tau$ decay mode provides a way for detection up to mass range of one TeV. Searches via the $t\bar{b}$ decay are difficult due to QCD background. As far as the neutral bosons go, we expect two light neutral bosons H_1^o and H_2^o , which are admixtures of the standard model Higgs boson and a lepton number carrying Higgs boson typical of LR models. Their decay mode via Z_2-pairs may lead to detectable signatures. If, however, their mass is bigger than $2M_{\Delta^{++}}$, then this decay mode provides a cleaner signature.

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Figure Captions

1. We plot the cross section for $H^{\pm}$ production via the process $bg \to H^{\pm}t$ as a function of $m_{H^{\pm}}$ for two extreme values of m_t: $m_t = 40$ GeV and $m_t = 150$ GeV.

2. We present the full gauge boson fusion cross section for lighter Higgs H_1^o, as a function of its mass, for two different values of m_{Z_R}, in the case $\lambda_H = \lambda_{min}$. $(Z_R \equiv Z_2)$.

3. We present cross sections for H_2^o production, in the $\lambda_H = \lambda_{min}$ case, as a function of m_{Z_R}, for several representative $m_{H_2^o}$ values.

4. We plot the cross sections for H_1^o and H_2^o as a function of the mass, for two different values of m_{Z_R}, in the maximal mixing case (b).

5. We present $H_{1,2}^o$ decay widths for the modes: a) $Z_LZ_L + W_LW_L$, $t\bar{t}$, and N_RN_R; and b) $Z_RZ_L + Z_RZ_R$. Plots are as a function of the mass of either H_1^o or H_2^o, $(Z_R \equiv Z_2;\ Z_L \equiv Z_1;\ W_L \equiv W_1)$.

6. We present decay widths for the four channels: a) $H_2^o \to H_1^oH_1^o$; b) $H_{1,2}^o \to H^+H^-$; c) $H_{1,2}^o \to \Delta_R^{++}\Delta_R^{--}$; and d) $H_{1,2}^o \to \Delta_L^o\Delta_L^o + \Delta_L^+\Delta_L^- + \Delta_L^{++}\Delta_L^{--}$. Plots are as a function of m_H, the mass of H_1^o or H_2^o. The mass of the Higgs, whose width is not being plotted in a given curve, is chosen as stated in the text. The Lagrangian parameters have been constrained so that all charged and L-triplet Higgs have zero mass.

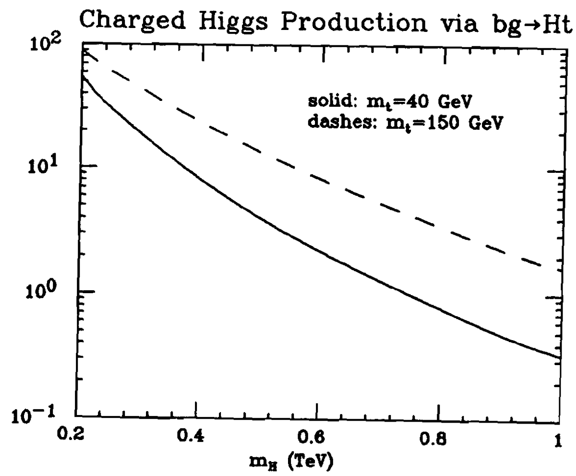


Fig. 1

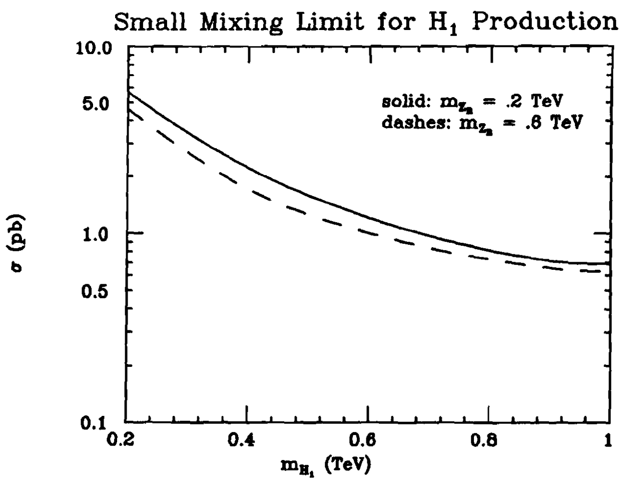


Fig. 2

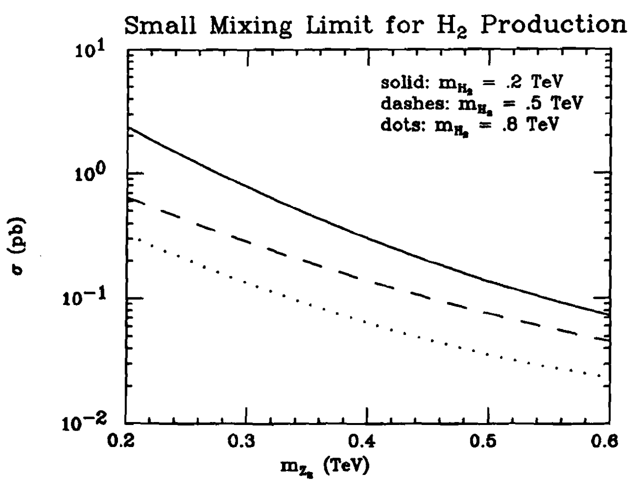
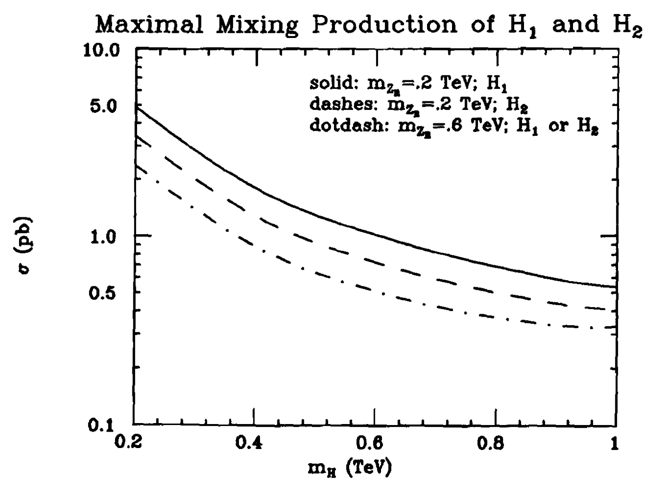


Fig. 3



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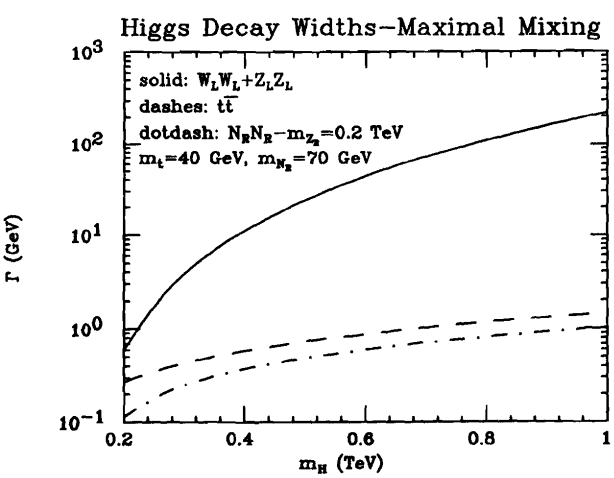

Fig. 5a

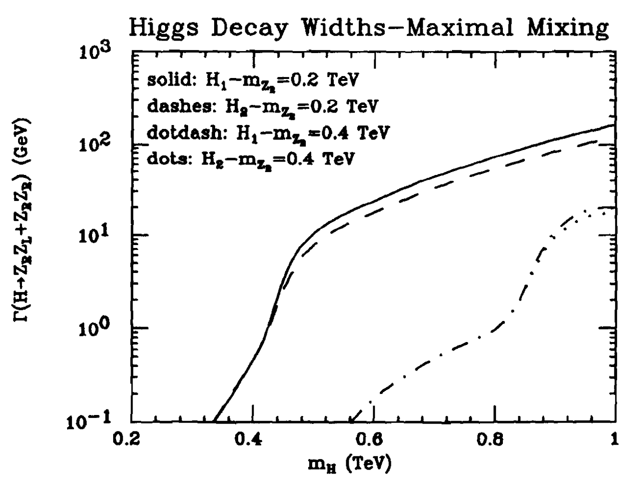

Fig. 5b

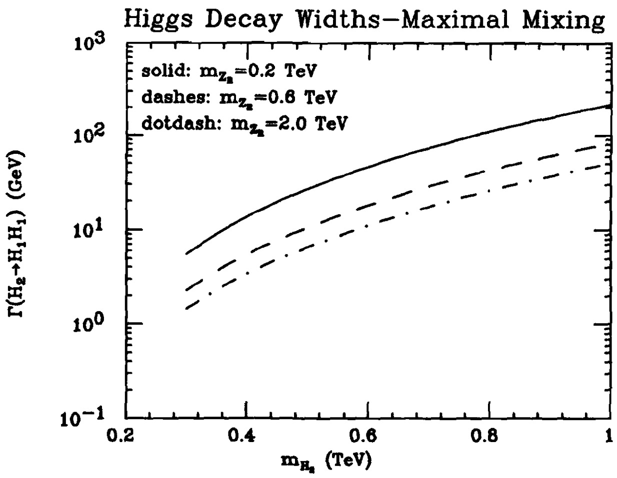

Fig. 6a

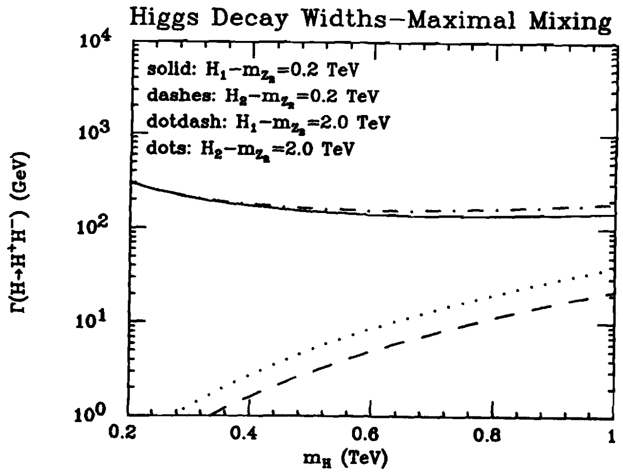

Fig. 6b

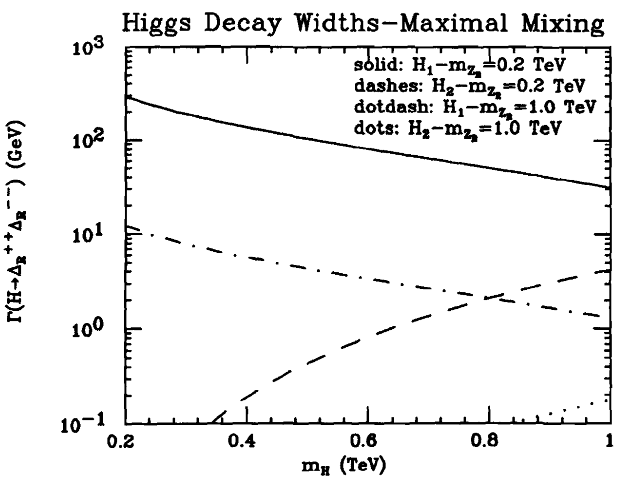

Fig. 6c

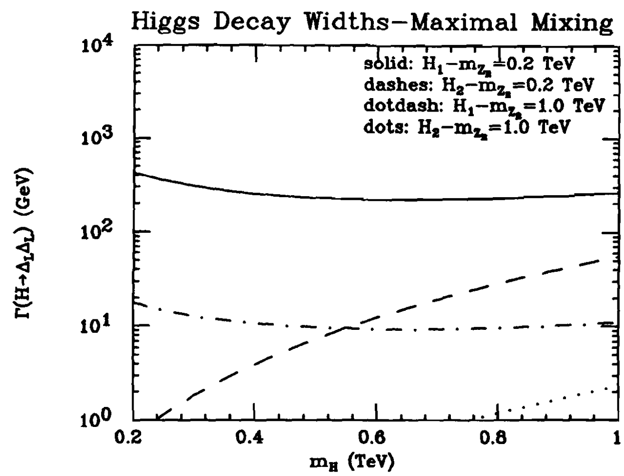

Fig. 6d