Probing the *WHZ* vertex in E_6 models at colliders

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We consider the production of charged Higgs scalar bosons at e^+e^- and hadron colliders as a probe of the WHZ_i vertex which can be nonzero at the tree level in E_6 models. Substantial cross sections at both kinds of colliders are possible but backgrounds are very substantial. These backgrounds can, in principle, be dealt with. e^+e^- annihilation with a light Z_2 may offer the best hope of probing this new interaction.

During the past few years there has been a great deal of interest in "superstring-inspired" E_6 models which promise a rich phenomenology beyond that of the standard model (SM) (Ref. 1). Not only do such models predict a whole host of new particles (e.g., exotic fermions, new gauge bosons, and their supersymmetric partners) which may be discovered in future accelerator searches but they also predict that new interactions among the more "conventional" particles may also occur. In this paper we consider one such interaction involving the gauge bosons and the Higgs field.

There are many extensions of the SM in which the Higgs sector is enlarged to include two (or more) doublets for a variety of reasons.[2] Such models necessarily lead to the existence of physical charged Higgs bosons ($H^{\pm}$) but it has been shown quite some time ago that the tree-level $W^{\pm}H^{\mp}Z$ coupling vanishes even if there are an arbitrary number of Higgs doublets and singlets.[3] It is possible, however, that if the gauge sector of the SM is also extended a tree level $W_iH_jZ_k$ vertex (with the indices labeling mass eigenstates) may arise due to mixing.[4] The simplest model of this kind is the rank-5 E_6 "superstring-inspired" model based on the group $SU(3)_C\times SU(2)_L\times U(1)_Y\times U(1)_{Y'}$, which has been extensively studied in the literature.[1] In this paper we examine how the production of charged Higgs scalars can be used to probe the *WHZ* vertex within the context of this class of models. Our analysis can be easily extended to more complex models involving additional *W*, *Z*, or *H* fields. Let us first review the basic aspects of this model.

The Higgs sector of this model contains three multiplets of Higgs fields[5] which we can write as

$$\Phi_1=\begin{bmatrix}\phi_1^0\\ \phi_1^-\end{bmatrix},\quad \Phi_2=\begin{bmatrix}\phi_2^+\\ \phi_2^0\end{bmatrix},\quad \Phi_3=\phi_3^0\ . \tag{1}$$

Under the $SU(2)_L\times U(1)_Y\times U(1)_{Y'}$ group these fields have the quantum numbers $(T_{3L},Y,Y')=(\pm\frac{1}{2},-1,Y_1')$, $(\pm\frac{1}{2},1,Y_2')$, and $(0,0,Y_3')$, respectively. This same group has the covariant derivative

$$D_\mu=\partial_\mu-ieQA_\mu-i\frac{g}{c_W}(T_{3L}-x_WQ)Z_\mu -i\frac{g}{c_W}\sqrt{x_w}Y'Z'_\mu \quad -\frac{ig}{\sqrt{2}}(T_+W_\mu^-+T_-W_\mu^+) \tag{2}$$

with $W^{\pm}$, Z being the usual massive gauge bosons of the SM and where $c_W=\cos\theta_W$, $x_W=\sin^2\theta_W$, and $T_\pm$ are isospin-raising and -lowering operators. Z' is the new neutral gauge boson coming from $U(1)_{Y'}$. The neutral members ϕ_i^0 of the above multiplets obtain vacuum expectation values $\langle\phi_i^0\rangle=v_i/\sqrt{2}$ which breaks the symmetry down to $U(1)_{em}$. Table I shows the quantum numbers of the members of the **27** representation of E_6; the superpartners of the fields N, N^c, and S^c act as the Higgs fields $\phi_i^0(i=1,2,3)$ in this model, respectively. Thus, $Y_1'=\frac{1}{6}$, $Y_2'=\frac{2}{3}$, and $Y_3'=-\frac{5}{6}$. The fields Z and Z' in general mix to form the mass eigenstates $Z_{1,2}$ with masses $M_{1,2}$:

$$Z_1=Z\cos\theta+Z'\sin\theta,\quad Z_2=Z'\cos\theta-Z\sin\theta\ . \tag{3}$$

Defining $\tan\beta\equiv v_2/v_1$ (with $v_1^2+v_2^2=v_{SM}^2$) one finds

$$\tan\theta=2\sqrt{x_W}(-Y_1'\cos^2\beta+Y_2'\sin^2\beta)M_Z^2/(M_2^2-M_Z^2) \tag{4}$$

with M_Z being the SM Z mass. A short calculation shows that there is no *WHZ* coupling resulting from (1) and (2) but a WHZ' coupling *does* exist. In the mass-eigenstate basis the $W^{\pm}H^{\mp}Z_{1,2}$ coupling is then found to be given by

$$\mathcal{L}_{WHZ_{1,2}}=2eM_Z\sin\beta\cos\beta(Y_1'+Y_2')\begin{bmatrix}\sin\theta\\ \cos\theta\end{bmatrix}\times(W^+H^-Z_{1,2}+\text{H.c.})\ . \tag{5}$$

Note both coupling terms are absent if either $\beta=0$, 90° (v_1 or $v_2=0$) or $Y_1'=-Y_2'$ and the WHZ_1 coupling is also absent when $\sin\theta=0$. As in the usual two-doublet model there is no induced $WH\gamma$ coupling. The *WHZ* couplings are completely specified by the values of $\tan\beta$ and M_2 which are both constrained by existing experimental data:[5,6] $\tan\beta$ is forced to be small for light M_2.

If such couplings exist, how can they be probed? One possibility which has been discussed[5,7] is the decay $Z_2\to W^{\pm}H^{\mp}$ since $Z_1\to W^{\pm}H^{\mp}$ is kinematically forbidden. But what if $H^{\pm}$ is too heavy to participate in Z_2 decays or we are at energies below the threshold for Z_2 production? We then must, e.g., examine $e^+e^-\to W^{\pm}H^{\mp}$ off resonance which proceeds via the diagrams in Fig. 1. The cross section for this process is found to be (summing over W^+H^- and W^-H^+ final states)

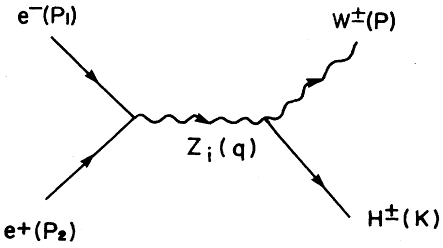

FIG. 1. Diagrams contributing to the process $e^+e^- \to W^\pm H^\mp$.

$$\frac{d\sigma}{dz}=\frac{25\pi\alpha^2 M_Z^2}{72x_W(1-x_W)}\left[\frac{t^2}{1+t^2}\right]$$
$$\times\left[\left[1+\frac{M_W^2-M_H^2}{s}\right]^2-\frac{4M_W^2}{s}\right]^{1/2}$$
$$\times\left[(1+z^2)+\frac{E_W^2}{M_W^2}(1-z^2)\right]$$
$$\times\sum_{i,j} c_i c_j (v_i^e v_j^e + a_i^e a_j^e)P_{ij}\ , \tag{6}$$

where $t\equiv\tan\beta$, $z=\cos\theta$, $\sqrt{s}$ is the center-of-mass energy, and where

$$P_{ij}\equiv\frac{(s-M_i^2)(s-M_j^2)+(\Gamma_i M_i)(\Gamma_j M_j)}{[(s-M_i^2)^2+(\Gamma_i M_i)^2][(s-M_j^2)^2+(\Gamma_j M_j)^2]}\ ,$$
$$E_W/M_W=\frac{s+M_W^2-M_H^2}{2\sqrt{s}M_W},\quad c_{1,2}=\begin{cases}\sin\theta\\ \cos\theta\end{cases} \tag{7}$$

with the v_i^e, a_i^e couplings normalized to

$$\mathcal{L}=\frac{g}{2c_W}\bar{e}\gamma_\mu(v_i^e-a_i^e\gamma_5)eZ_i^\mu\ . \tag{8}$$

Γ_i are the decay widths of the Z_i bosons. In our numerical work, we take $\alpha^{-1}=128.0$, $M_W=81$ GeV, $M_1\simeq M_Z=92$ GeV, $x_W=0.230$, $\Gamma_1=2.42$ GeV, and $\Gamma_2=1.5M_2/M_1$ GeV (Ref. 5). Note that the shape of the angular distribution is controlled by the E_W^2/M_W^2 factor in (6): for purely longitudinal *W*'s we obtain a $(1-z^2)$ behavior while purely transverse *W*'s produce a $(1+z^2)$ behavior but longitudinal *W*'s clearly dominate at high energies ($E_W^2/M_W^2\gg 1$). Figures 2(a) and 2(b) show the integrated cross section for $e^+e^-\to W^\pm H^\mp$ at $\sqrt{s}=200$ GeV as a function of m_H for different values of $\tan\beta$ with $M_2=150$ and 250 GeV, respectively. We limit ourselves to $\tan\beta\lesssim 0.6$ for such light Z_2 masses.[5] Note that for larger M_2 values, σ decreases $\sim M_2^{-4}$ since it is dominated by Z_2 exchange and because the mixing angle θ is also tightly constrained as a function of M_2.

If $\tan\beta\simeq\frac{1}{2}$ then for either value of the M_2 mass $\sigma\simeq 0.1-0.2$ pb which would yield $\simeq 500-1000$ events for an integrated luminosity of 5×10^{39} cm^{-2}sec^{-1}. This process should be easily identified via leptonic *W* decay ($W\to l\nu$) and $H\to t\bar{b}$ (assuming $m_H>m_t+m_b$). The largest backgrounds are from $e^+e^-\to W^+W^-$ and $2Z^0$ but these are strongly peaked forward and/or backward while $W^\pm H^\mp$ production is roughly flat in z at these energies. $e^+e^-\to H^+H^-$ might also contribute to the background from subsequent decay but its cross section falls off as β^3 and has a $\sin^2\theta$ behavior as can easily be seen from the expression

$$\frac{d\sigma}{dz}(e^+e^-\to HH^-)=\frac{\pi\alpha^2}{64x_W^2(1-x_W)^2}\frac{1}{s}\beta^3(1-z^2)$$
$$\times\sum_{ij} c_i c_j (v_i^e v_j^e + a_i^e a_j^e)s^2 P_{ij}\ . \tag{9}$$

Note that $\sigma(H^+H^-)$ will now have a weak $\tan\beta$ dependence due to $Z-Z'$ mixing as shown in Figs. 3(a) and 3(b). Also note that, e.g., $\sigma(H^+H^-)\simeq\sigma(W^\pm H^\mp)$ for $m_H\simeq 80-90$ GeV and $\tan\beta\simeq 0.3-0.5$ depending on M_2 so that these two cross sections are quite comparable over the mass region of interest.

The separation of the $W^\pm H^\mp$ signal from the W^+W^- background may be performed as follows. Although W^+W^- production occurs at a large rate ($\simeq 17$ pb at $\sqrt{s}=200$ GeV) it is mostly peaked forward due to the t-channel, ν-exchange graph. Making a cut such as $-0.9\le z\le 0$ (or using polarized beams) can remove a major part of this background reducing it by an order of magnitude ($\simeq 2$ pb) while not drastically affecting the signal. One now looks for events with *W*'s in them by using the $W\to$leptons mode. Once a *W* is observed on one side, the other side of the event is examined. If decays such as $W/H\to t\bar{b}$ are allowed kinematically they should be easily identified using conventional heavy-quark triggers. However, the branching fraction for $W\to t\bar{b}$ will be reasonably small due to phase-space suppression, i.e., if $m_t=60-70$ GeV then the branching fraction is $\simeq 0.02-0.05$ or so. This will reduce the background from W^+W^- production by another factor of 20–50

TABLE I. Quantum numbers of the particles in the **27** representation of E_6 under $SU(3)_C\times SU(2)_L\times U(1)_Y\times U(1)_{Y'}$.

	$SU(3)_C$	Q	T_{3L}	Y'
u	3	$\frac{2}{3}$	$\frac{1}{2}$	$-\frac{1}{3}$
d	3	$-\frac{1}{3}$	$-\frac{1}{2}$	$-\frac{1}{3}$
u^c	$\bar{3}$	$-\frac{2}{3}$	0	$-\frac{1}{3}$
d^c	$\bar{3}$	$\frac{1}{3}$	0	$\frac{1}{6}$
ν_e	1	0	$\frac{1}{2}$	$\frac{1}{6}$
e	1	-1	$-\frac{1}{2}$	$\frac{1}{6}$
e^c	1	1	0	$-\frac{1}{3}$
ν_e^c	1	0	0	$-\frac{5}{6}$
h	3	$-\frac{1}{3}$	0	$\frac{2}{3}$
h^c	$\bar{3}$	$\frac{1}{3}$	0	$\frac{1}{6}$
N	1	0	$\frac{1}{2}$	$\frac{1}{6}$
E	1	-1	$-\frac{1}{2}$	$\frac{1}{6}$
E^c	1	1	$\frac{1}{2}$	$\frac{2}{3}$
N^c	1	0	$-\frac{1}{2}$	$\frac{2}{3}$
S^c	1	0	0	$-\frac{5}{6}$

since $H^+ \rightarrow t\bar{b}$ will be by far the dominant Higgs-boson decay. The WW background has now been reduced to the level of the signal (or below it) by using angular cuts and triggering on the $t\bar{b}$ final state. In addition, if the W^+ and H^+ masses are sufficiently different ($\gtrsim 10-20$ GeV, say) one may also be able to separate the events based on the difference in the invariant mass of the jets produced by the $t\bar{b}$ system. More over, one might also

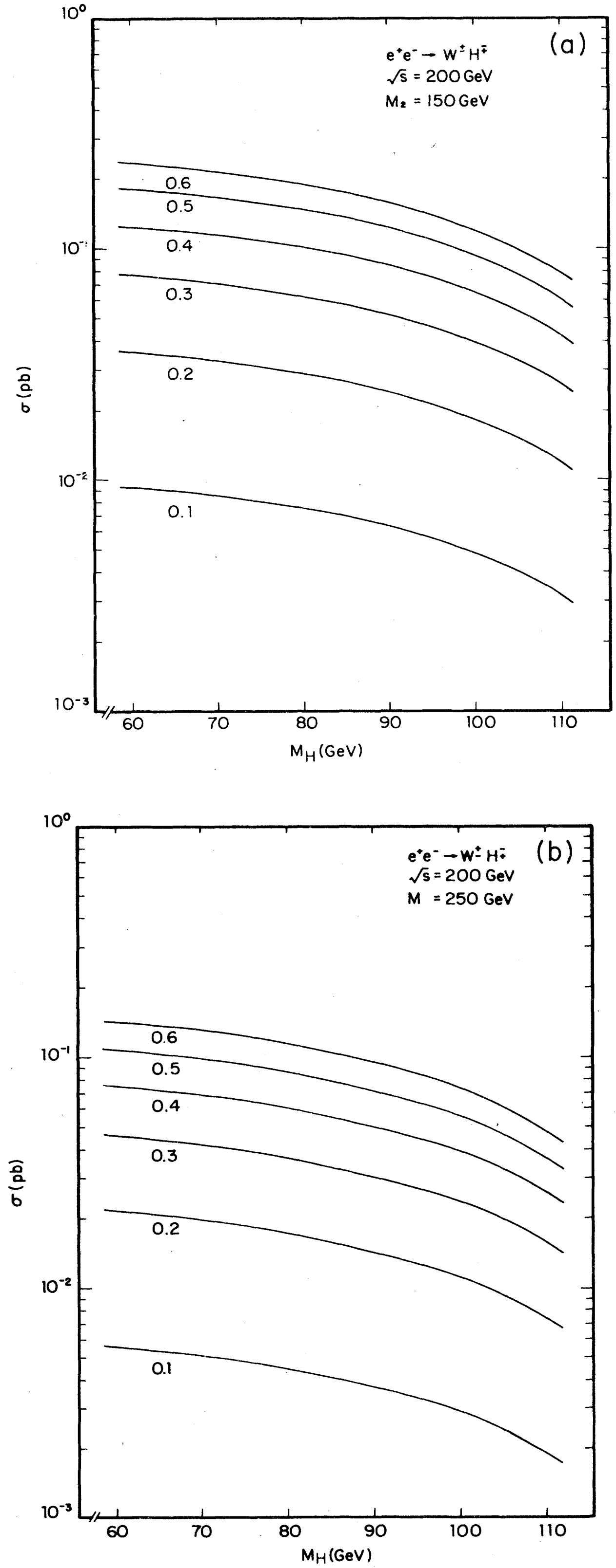

FIG. 2. Cross section for $e^+e^- \rightarrow W^{\pm}H^{\mp}$ at $\sqrt{s} = 200$ GeV as a function of M_H for different values of $\tan\beta$ with (a) $M_2 = 150$ GeV, (b) $M_2 = 250$ GeV.

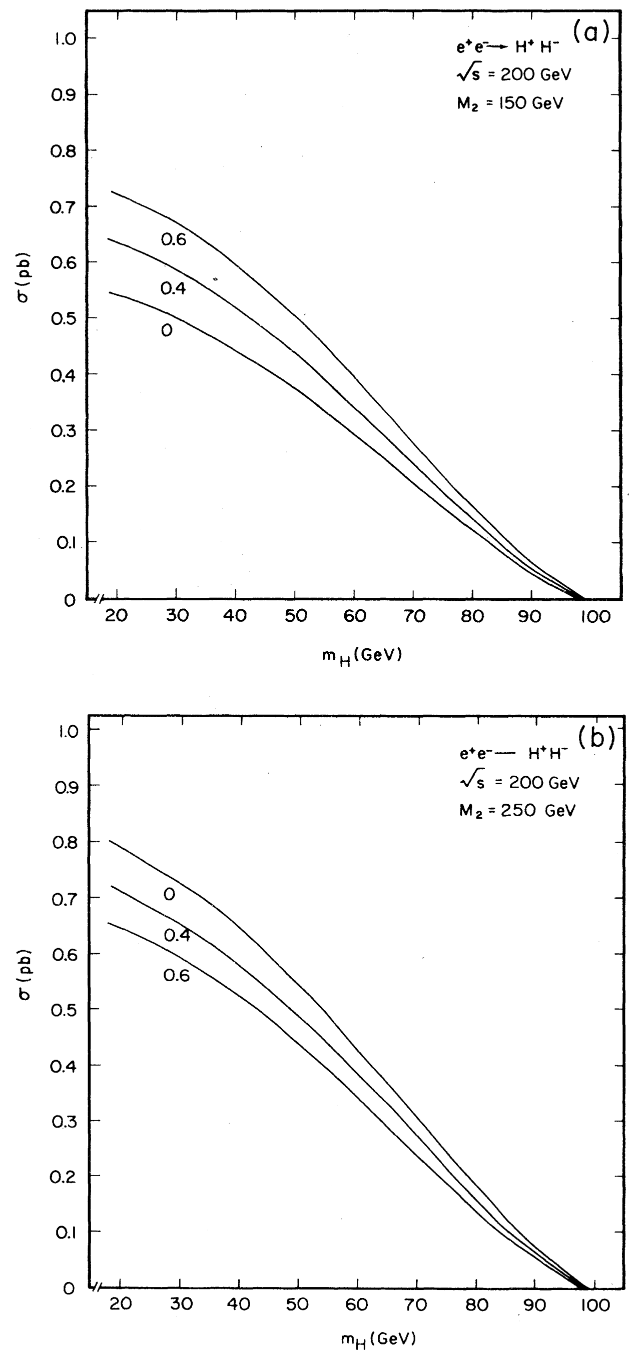

FIG. 3. Cross section for $e^+e^- \rightarrow H^+H^-$ at $\sqrt{s} = 200$ GeV as a function of M_H for different values of $\tan\beta$ with (a) $M_2 = 150$ GeV, (b) $M_2 = 250$ GeV.

make use of the difference in the threshold behaviors of the WH and WW final states (since the background-to-signal ratio is quite sensitive to the value of $\sqrt{s}$) to separate the two classes of events. The $2Z$ final state is more easily dealt with since the total cross section is smaller than that for W^+W^- by an order of magnitude and one looks for events with lepton pairs and not lepton plus missing energy. Even if one Z is misidentified as W, the decay $Z\to t\bar{t}$ is probably kinematically forbidden so there should be no $2Z$ events containing t quarks.

If $m_H \gtrsim 100$ GeV or so, $\sigma(W^\pm H^\mp)$ essentially vanishes at $\sqrt{s}=200$ GeV so one might ask if one could do better at a $\sqrt{s}=1$ TeV collider. At these energies, $\sigma(W^\pm H^\mp)$ is found to be significantly damped $\sim s^{-2}$ in comparison to the more typical $\sim s^{-1}$ type of cross-section behavior unless there is a very nearby Z_2 resonance even though larger values of $\tan\beta$ are now possible. We find that at such energies and luminosities, the $W^\pm H^\mp$ cross section is too small to be observable.

Single-$H^\pm$ production is also possible via gauge-boson fusion,[8] e.g., $e^+e^-\to e^+\nu Z_i W^-\to H^- e^+\nu$ as shown in Fig. 4. Assuming the cross section is dominated by longitudinal W's and Z_i's we find that

$$\sigma_L^i(e^+e^-\to H^\pm l\nu) = \frac{G_F^2 M_W^2 M_Z^2}{\pi^2}(v_i^{e2}+a_i^{e2})\frac{\Gamma_L^i(H^+\to W^+Z_i)}{m_H^3} \times[(1+m_H^2/s)\ln s/m_H^2-2(1-m_H^2/s)]\ , \quad (10)$$

where $(\delta_W\equiv m_W/m_H,\ \delta_i\equiv m_i/m_H)$

$$\Gamma_L^i=\frac{G_F m_H^3}{8\sqrt{2}\pi}\frac{c_{Wi}^2}{M_i^2}(1-\delta_W^2-\delta_i^2)^2 \times[1-2(\delta_W^2+\delta_i^2)+(\delta_W^2-\delta_i^2)^2]^{1/2}\ . \quad (11)$$

Here the WHZ_i coupling is given by gc_{Wi} where the values of c_{Wi} are then obtainable from (5). As M_2 increases, the cross section decreases rapidly since $\Gamma_L \sim M_2^{-2}$ and is further phase-space suppressed. Figure 5 shows σ_L as a function of $\sqrt{s}$ for $M_2=150$ GeV and $\tan\beta=0.6$ with $M_H=0.5$ or 1.0 TeV. Thus, σ_L^i is found to be quite small and dominated by Z_2 as expected. For $m_H=500$ GeV, $\sqrt{s}=1$ TeV, $\tan\beta=0.6$, $M_2=150$ GeV, $\sigma_L=\sigma_L^1+\sigma_L^2$ is found to be only 1.6×10^{-5} pb which is substantially smaller than that expected for neutral-Higgs-boson production at these energies. This small cross section results from $M_2^2>M_1^2$, $\tan\beta$ suppression and the smallness of the couplings of v_2^e and a_2^e. It thus seems very unlikely that this process could yield substantial information about the charged-Higgs-boson couplings.

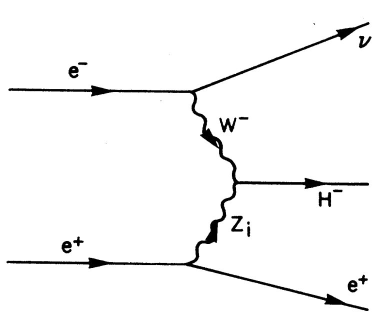

FIG. 4. Diagrams responsible for $e^+e^-\to H^\pm e\nu$.

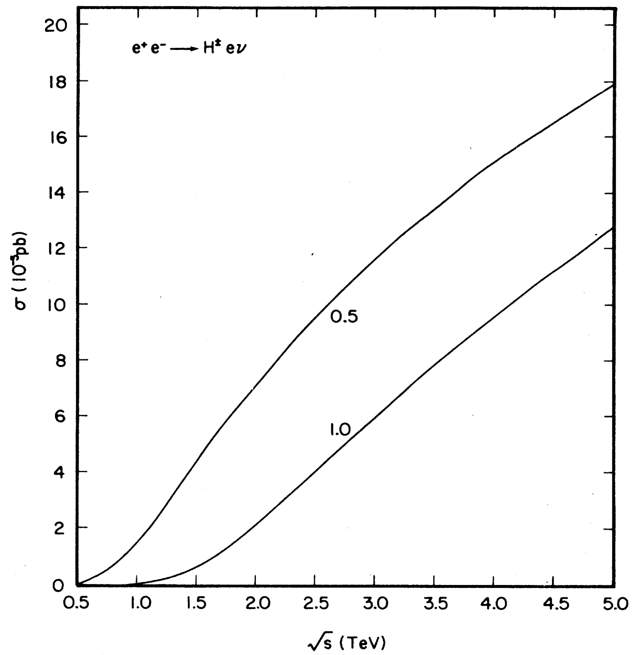

FIG. 5. Cross section for $e^+e^-\to H^\pm e\nu$ as a function of $\sqrt{s}$ for $M_H=0.5$ and 1.0 TeV, $M_2=150$ GeV, and $\tan\beta=0.6$.

We now turn to the production of $H^\pm$ in hadronic collisions. If $H^\pm$ are relatively light the cross section for associated production via $q\bar{q}\to Z_i\to W^\pm H^\mp$ may be significant and might allow for a test of the WHZ vertex. The cross section for this process is (for pp collisions)

$$\frac{d\sigma}{dM}=\sum_q \frac{2M}{s}\int_{-Y}^{Y}dy[q(x_1,M^2)\bar{q}(x_2,M^2) +\bar{q}(x_1,M^2)q(x_2,M^2)] \times\int_{-Z_0}^{Z_0}dz\frac{d\hat{\sigma}_q}{dz}\ , \quad (12)$$

where $q,\bar{q}$ are the appropriate parton distribution function. $d\hat{\sigma}_q/dz$ is obtainable from (6)–(8) above by the replacement $e\to q$, $s\to M^2$ and dividing by a color factor of 3. We also define

$$x_{1,2}\equiv\frac{M}{\sqrt{s}}e^{\pm y},\quad Z_0\equiv\min[\beta^{-1}\tanh(Y-|y|),1]\ , \quad (13)$$

with

$$\beta\equiv\left[\left(1-\frac{M_W^2+M_H^2}{M^2}\right)-\frac{4M_W^2M_H^2}{M^4}\right]^{1/2} \times\left(1+\frac{M_W^2-M_H^2}{M^2}\right)^{-1} \quad (14)$$

and Y is the rapidity cutoff (we will take $Y=2.5$ in our numerical analysis below). Figures 6(a)–6(c) show $d\sigma/dM$ as a function of M for different values of M_H, $\tan\beta$, and M_2 (Ref. 8). As can be easily seen, the Z_2 peak

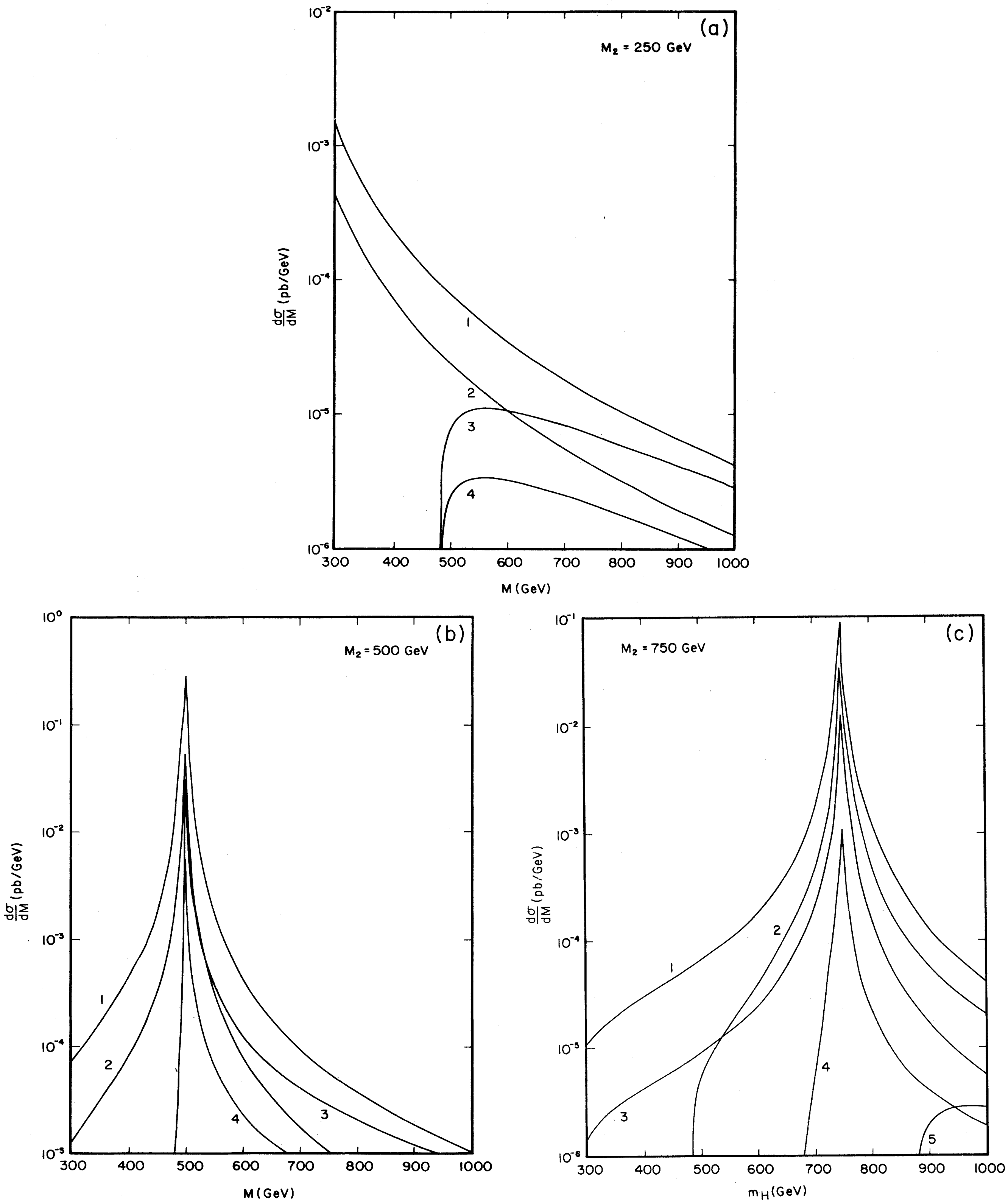

FIG. 6. Differential cross section $d\sigma/dM$ for $pp \to W^{\pm}H^{\mp}+X$ at $\sqrt{s}=40$ TeV as a function of M for different model parameters (a) $M_2=250$ GeV: (1) $M_H=200$ GeV, $\tan\beta=0.6$; (2) $M_H=200$ GeV, $\tan\beta=0.3$; (3) $M_H=400$ GeV, $\tan\beta=0.6$; (4) $M_H=400$ GeV, $\tan\beta=0.3$. (b) $M_2=500$ GeV: (1) $M_H=200$ GeV, $\tan\beta=0.9$; (2) $M_H=200$ GeV, $\tan\beta=0.3$; (3) $M_H=400$ GeV, $\tan\beta=0.9$; (4) $M_H=400$ GeV, $\tan\beta=0.3$; (c) $M_2=750$ GeV: (1) $M_H=200$ GeV, $\tan\beta=1.2$; (2) $M_H=400$ GeV, $\tan\beta=1.2$; (3) $M_H=200$ GeV, $\tan\beta=0.3$; (4) $M_H=600$ GeV, $\tan\beta=0.3$; (5) $M_H=800$ GeV, $\tan\beta=1.2$.

may play a very important role in probing the *WHZ* vertex. In Fig. 6(a), $M_H \geq 200$ GeV so that we are kinematically constrained to be above the Z_2 resonance region. In this case the cross sections are quite small $\sim 10^{-3}\sigma(pp \to W^+W^- + X)$ or less yielding a few thousand events/yr assuming an integrated luminosity of 10^4 pb^{-1}/yr and favorable values of $\tan\beta$ and M_H. $\sigma(pp \to W^\pm H^\mp + X)$ is roughly a factor of 5 below the SM cross section for $pp \to W^\pm H^0 + X$ for identical Higgs-boson masses. This situation improves drastically when the Z_2 resonance can make a significant contribution as shown in Figs. 6(b) and 6(c). The differential cross sections fall off drastically on either side of the Z_2 resonance since it is quite narrow. The cross section also falls off quite rapidly with increasing M_H. It would thus appear from Figs. 6(a)–6(c) that unless we are somewhat lucky in the choice of parameters it will not be a straightforward task to use associated production to test the vertex of interest. The possibility of observing $Z_2 \to W^\pm H^\mp$ has also been discussed in Ref. 5. Things are, however, a bit more optimistic than they appear.

Even if the reconstruction of the Z_2 peak is made difficult by the missing p_T from $W \to$lepton+neutrino decay, the resulting $W^\pm H^\mp$ final state should still be observable with significant rates. Taking $\tan\beta = 0.6$ and $M_H = 100$ GeV, for purposes of demonstration, we find the total $W^\pm H^\mp$ production cross section to be 8.6 (1.2) pb for $M_2 = 250$ (500) GeV, which yields 8.6×10^4 (1.2×10^4) events/yr at the Superconducting Super Collider (SSC). The major backgrounds include W^+W^- and $W^\pm$ plus heavy-quark jets. The W^+W^- background ($\simeq 90$ pb) can be handled in a manner similar to e^+e^- where the observation of the $t\bar{b}$ from the second W decay drastically reduces the W^+W^- contamination of the signal. Also, since the range of interest for M_H is significantly larger in the pp case the $t\bar{b}$ jets need to reconstruct to a value which is substantially different from M_W. The rate for $W^\pm$ plus a pair of heavy-quark jets with an invariant mass near M_H is not expected to be large in comparison to the signal. Including the leptonic branching fractions of the W (not including τ's) for the values of the parameters above one should expect $1.9(0.26) \times 10^4$ events/yr of the type $l^\pm$+heavy-quark jets+missing p_T at the SSC. Such a signal should be observable but further studies of the QCD $W^\pm$-plus-heavy-jet backgrounds are still necessary.

It should be noted that $W^\pm H^\mp$ production can also occur via gg fusion at the one-loop level via triangle and box diagrams even if there is no tree-level *WHZ* coupling. For the production of high-mass charged Higgs bosons where gluon fusion dominates this will be an irreducible background for any observation of tree-level *WHZ* coupling. It is expected, however, that such processes should yield only small event rates.[9] These calculations are presently in progress.

For heavier charged Higgs bosons, WZ_i fusion may lead to a substantial cross section.[10] In particular, we examine the subprocess $W^\pm Z_i \to H^\pm \to WZ_j$ to probe the *WHZ* vertex. The $W^\pm Z_i \to W^\pm Z_j$ process has several potential contributions as shown in Fig. 7. Since we are looking for a bump in the WZ_j invariant-mass distribution as a charged-Higgs-boson signal and since we assume $H^\pm$ are heavy ($M_H = 500$ GeV, say) we will make two simplifying approximations. (1) We will assume that only longitudinally polarized W's and Z_i's contribute significantly to the cross section and (2) we will throw away all diagrams other than s-channel $H^\pm$ exchange. This should be a reasonable approximation near the resonance peak. The subprocess cross section is then found to be (assuming only longitudinal gauge bosons)

$$\hat{\sigma}_L(M) = \frac{G_F^2 M_W^4}{32\pi M^2}\left[1 - \frac{2(M_W^2 + M_j^2)}{M^2} + \frac{(M_W^2 - M_j^2)^2}{M^4}\right]^{1/2} \times c_{Wi}^2 c_{Wj}^2 [(M^2 - M_H^2)^2 + \Gamma_H^2 M_H^2]^{-1} \times \frac{(M^2 - M_i^2 - M_W^2)^2 (M^2 - M_j^2 - M_W^2)^2}{M_W^4 M_i^2 M_j^2} \tag{15}$$

be rewriting Eq. (5) as

$$\mathcal{L}_{WHZ_{1,2}} = g c_{Wi} W_\mu^+ Z_i^\mu H^- + \text{H.c.} , \tag{5'}$$

where M is the boson pair invariant mass and c_{W_i} are defined above. The Higgs width Γ_H, get contributions from $W^\pm Z_i$ as given above and $t\bar{b}$:

$$\Gamma(H^+ \to t\bar{b}) = \frac{3G_F M_H^3}{4\sqrt{2}\pi} |V_{tb}|^2 \delta_t^2 (1 - \delta_t^2 - \delta_b^2) \times [1 - 2(\delta_t^2 + \delta_b^2) + (\delta_t^2 - \delta_b^2)^2]^{1/2} \times \left[(\tan^2\beta)^{-1} + \frac{\delta_b^2}{\delta_t^2}(\tan^2\beta)\right] . \tag{16}$$

Because of the small WHZ_i couplings, $t\bar{b}$ can make a substantial contribution to Γ_H. The differential cross section, with these approximations, for $pp \to H^\pm + X \to W^\pm Z_j + X$ is given by

$$\frac{d\sigma}{dM^2} = \frac{1}{s}\int_{M^2/s}^{1} \frac{d\tau}{\tau} \int_{-\eta}^{\eta} dy \sum_{a,b} q_a(\sqrt{\tau}e^{y}) q_b(\sqrt{\tau}e^{-y}) \int_{-\hat{\eta}}^{\hat{\eta}} d\hat{y} f_L^a(\sqrt{\hat{\tau}}e^{\hat{y}}) f_L^b(\sqrt{\hat{\tau}}e^{-\hat{y}}) \hat{\sigma}_L(M) , \tag{17}$$

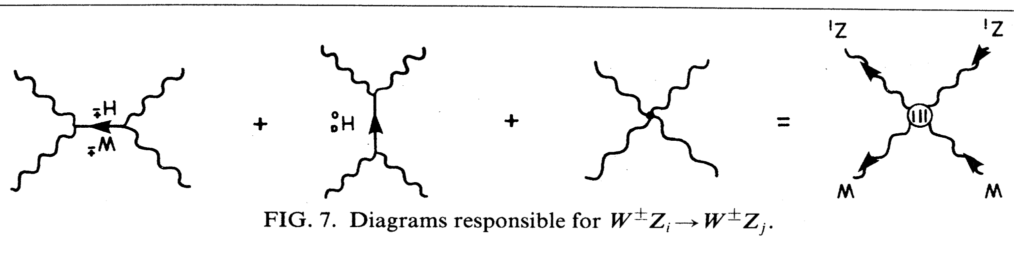

FIG. 7. Diagrams responsible for $W^\pm Z_i \to W^\pm Z_j$.

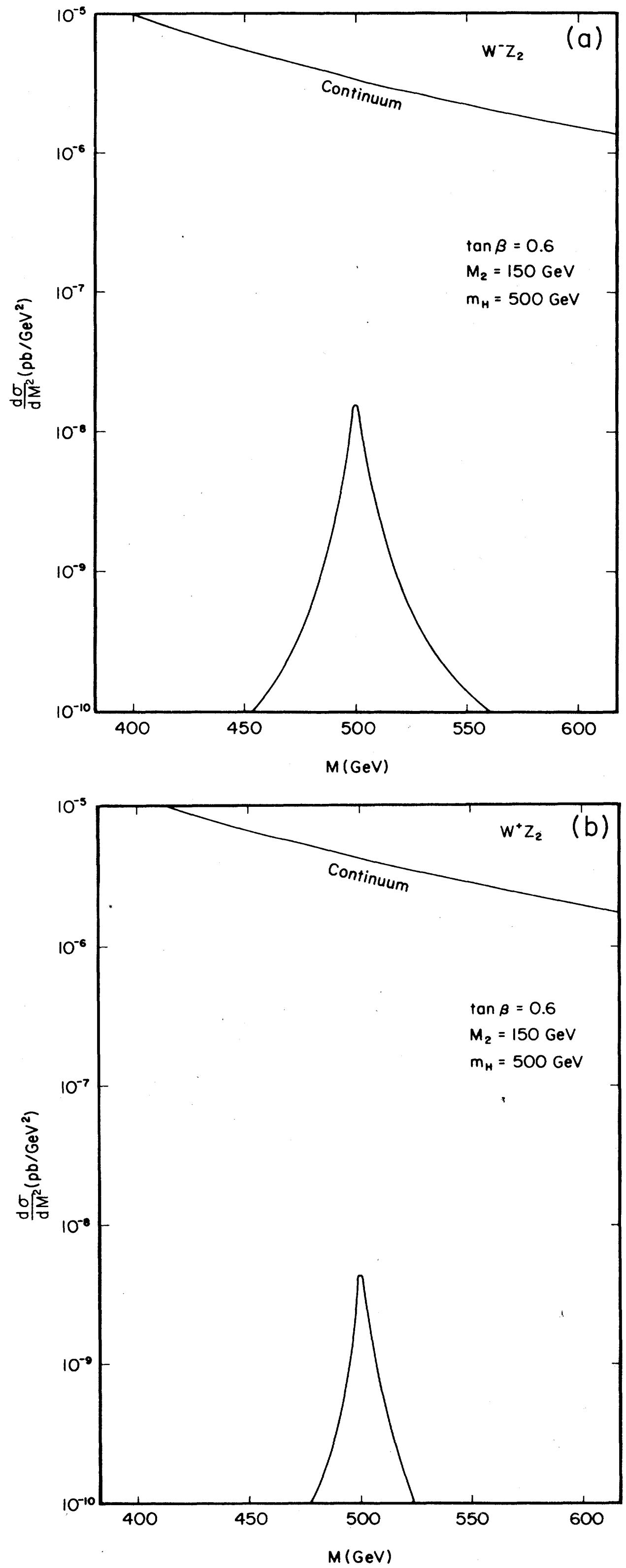

FIG. 8. Differential cross sections $d\sigma/dM^2$ for (a) $pp \to W^- Z_2 + X$ and (b) $pp \to W^+ Z_2 + X$ from the continuum $q\bar{q} \to W^\pm Z_2$ and from resonant charged-Higgs-boson production: $W^\pm Z_i \to H^\pm \to W^\pm Z_2$ at $\sqrt{s} = 40$ TeV.

where $\tau = \hat{s}/s$, $\hat{\tau} = M^2/\hat{s}$, $\eta' = -\ln(\sqrt{\hat{\tau}})$, and

$$\eta \equiv \min[Y, -\ln(\sqrt{\tau})] \ . \tag{18}$$

$q_{a,b}$ are the partons emitting the $W^\pm$ and Z_i and f_L^a is given by[11]

$$f_L^a(x) = \frac{g_V^2 + g_a^2}{4\pi^2} \frac{1-x}{x} \frac{M^2}{M^2 + M_V^2(1-X)} \ . \tag{19}$$

Here $g_{V,a}$ are the couplings of parton a to the gauge boson V. Note in (17) that a sum over all possible partons emitting $W^\pm$ and Z_i is performed. In (18), Y is the rapidity cut which we take to be $Y = 2.5$. Before evaluating $d\sigma/dM^2$ we note that there is an irreducible background coming from $q\bar{q} \to W^\pm Z_j$ via the usual s-, t-, and u-channel-exchange diagrams. A small modification of the SM result for $q\bar{q} \to W^\pm Z$ allows us to calculate the cross section for this background process which shows only a weak dependence on tanβ. As far as signal to background is concerned, the $W^\pm Z_2$ final state is best since $\sigma(q\bar{q} \to W^\pm Z_2) < \sigma(q\bar{q} \to W^\pm Z_1)$ due to mass and couplings suppression while $\sigma(W^\pm Z_i \to W^\pm Z^2) > \sigma(W^\pm Z_i \to W^+ Z_1)$ due to enhanced couplings. Even so, the $H^\pm$ signal is very far below the $W^\pm Z_2$ continuum as can be seen in Figs. 8(a) and 8(b) for a charged Higgs boson of mass 500 GeV and with an optimum choice of model parameters: tan$\beta = 0.6$ and $M_2 = 150$ GeV. Increasing M_2 and/or lowering tanβ would only decrease the $H^\pm$ resonance "signal" in comparison to the $W^\pm Z_2$ continuum "background." The situation here is comparable to (worse than?) the situation in the SM where one is looking for $H^0 \to W^+ W^-$ with a large $W^+ W^-$ background continuum since the ratio of "signal" to "background" is comparable. As is the case in the SM, the continuum is dominated by transverse-gauge-boson pairs whereas the $H^\pm$ peak consists dominantly of longitudinal-gauge-boson pairs. For heavier charged Higgs bosons, the "signal" to "background" decreases further making it unlikely that WHZ couplings could ever be probed without a judicious choice of experimental cuts.

We have examined the possibility of probing the WHZ_i couplings at e^+e^- and pp colliders within the context of the simplest E_6 superstring-inspired model. We find that at pp colliders, i.e., the SSC, invariant-mass distributions are either too small or suffer from enormous backgrounds unless we are not far from the Z_2 resonance and the $W^\pm H^\mp$ mode is kinematically allowed. The total $W^\pm H^\mp$ cross section, however, appears promising with large rates. Background analysis for this process from other sources are now in progress. A similar situation holds for heavy $H^\pm$ at e^+e^- colliders. The only situation which looks relatively hopeful is $M_H \approx 60-110$ GeV at an e^+e^- center-of-mass energy $\sqrt{s} = 200$ GeV with a nearby Z_2 ($M_2 = 150-300$ GeV). Backgrounds from W^+W^- are potentially large even in this case but angular and (to some extent) invariant-mass cuts are extremely helpful and $W^\pm H^\mp$ has a distinct decay signature. Use of the $W/H \to t\bar{b}$ mode drastically reduces this background. $W^\pm H^\mp$ production may be used as a signal for a Z_2 below threshold. The results we obtained are expected to hold, at last qualitatively, in other extended electroweak models with new neutral gauge bosons and an enlarged Higgs sector. The observation of the *WHZ* coupling at both e^+e^- and hadron colliders may thus be possible. The observation of tree level *WHZ* couplings would be a unique signal for new physics beyond the SM.

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