Electroweak scale neutrinos and Higgses

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Abstract. We present two different models with electroweak scale right-handed neutrinos. One of the models is created under the constraint that any addition to the Standard Model must not introduce new higher scales. The model contains right-handed neutrinos with electroweak scale masses and a lepton number violating singlet scalar field. The scalar phenomenology is also presented. The second model is a triplet Higgs model where again the right-handed neutrinos have electroweak scale masses. In this case the model has a rich scalar phenomenology and in particular we present the analysis involving the doubly charged Higgs.

Keywords: Neutrinos, seesaw, Higgs.
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INTRODUCTION

We present two recent models [1, 2] of electroweak scale right-handed neutrinos and their scalar phenomenology. First we describe a model based on the idea that given our current (experimental) knowledge of particle physics one should explore a "truly minimal" extension of the Standard Model (SM). We consider the possibility of having just one scale associated with all the high energy physics (HEP) phenomena. Thus we propose a minimal extension of the SM where new phenomena associated to neutrino physics can also be explained by physics at the Electroweak (EW). We then review a recent model [3] in which the RH neutrinos that participate in the seesaw mechanism are *active* in the sense that they are *electroweak nonsinglets*. If they are not too heavy, they can be produced at colliders and the seesaw mechanism could be tested. The right-handed neutrinos of [3] are members of SM doublets of mirror leptons and their Majorana masses are linked to EW scale through a coupling with a Higgs triplet that develops an EW scale VEV. In this model, the sources of the SM SSB include Higgs triplets.

MINIMAL MODEL

Based on the minimalistic constraint described above we assume

- SM particle content and gauge interactions.
- Existence of three RH neutrinos with a mass scale of EW size.
- Global $U(1)_L$ spontaneously (and/or explicitly) broken at the EW scale by a single complex scalar field.

- All mass scales come from spontaneous symmetry breaking (SSB). This leads to a Higgs sector that includes a Higgs SU(2)$_L$ doublet field Φ with hypercharge 1 (i.e. the usual SM Higgs doublet) and a SM singlet complex scalar field η with lepton number -2.

The terms of the Lagrangian relevant for Higgs and neutrino physics are $\mathscr{L}_{\nu H} = \mathscr{L}_{\nu y} - V$, with

$$\mathscr{L}_{\nu y} = -y_{\alpha i}\bar{L}_\alpha N_{Ri}\Phi - \frac{1}{2}Z_{ij}\eta \bar{N}^c_{Ri}N_{Rj} + h.c.\ , \tag{1}$$

where N_R represents the RH neutrinos, $\psi^c = C\gamma^0\psi^*$ and $\psi^c_R \equiv (\psi_R)^c = P_L\psi^c$ has left-handed chirality. The scalar potential is given by

$$\begin{aligned} V &= \mu_D^2\Phi^\dagger\Phi + \frac{\lambda}{2}\left(\Phi^\dagger\Phi\right)^2 + \mu_S^2\eta^*\eta + \lambda'\left(\eta^*\eta\right)^2 \\ &+ \kappa\left(\eta\Phi^\dagger\Phi + h.c.\right) + \lambda_m\left(\Phi^\dagger\Phi\right)\left(\eta^*\eta\right)\ . \end{aligned} \tag{2}$$

Note that the fifth term in the potential breaks explicitly the U(1) associated to lepton number.

It is useful to define the scalar mass eigenstates through

$$\mathscr{H} = \begin{pmatrix} \phi^0 \\ \rho \end{pmatrix} = \begin{pmatrix} \cos\alpha & -\sin\alpha \\ \sin\alpha & \cos\alpha \end{pmatrix}\begin{pmatrix} h \\ H \end{pmatrix}\ , \tag{3}$$

where we have used the following relations:

$$\Phi = \begin{pmatrix} 0 \\ \frac{\phi^0+v}{\sqrt{2}} \end{pmatrix} \text{ and } \eta = \frac{\rho+u+i\sigma}{\sqrt{2}}\ . \tag{4}$$

Using these definitions the Lagrangian becomes

$$\begin{aligned} \mathscr{L}_{\nu y} &\supset -y_{\alpha i}\bar{\nu}_{L\alpha}N_{Ri}\frac{\phi^0}{\sqrt{2}} - \frac{1}{2}Z_{ij}\frac{(\rho+i\sigma)}{\sqrt{2}}\bar{N}^c_{Ri}N_{Rj} + h.c. \\ &= \left(-\frac{y_{\alpha i}}{\sqrt{2}}\bar{\nu}_{L\alpha}N_{Ri}(c_\alpha\ h - s_\alpha\ H) + h.c.\right) - \left(\frac{i}{2\sqrt{2}}Z_{ij}\bar{N}^c_{Ri}N_{Rj}\sigma + h.c.\right) \\ &- \left(\frac{1}{2\sqrt{2}}Z_{ij}\bar{N}^c_{Ri}N_{Rj}(s_\alpha\ h + c_\alpha\ H) + h.c\right)\ . \end{aligned} \tag{5}$$

We are interested in EW scale RH neutrinos. The Dirac part on the other hand will be constrained from the seesaw. Writing the neutrino mass matrix as

$$m_\nu = \begin{pmatrix} 0 & m_D \\ m_D & M_M \end{pmatrix}\ , \tag{6}$$

where $(m_D)_{\alpha i} = y_{\alpha i}v/\sqrt{2}$. As an example lets consider the third family of SM fields and one RH neutrino, thus Eq.(6) becomes a 2×2 matrix. Assuming $m_D << M_M$ we obtain the eigenvalues $m_1 = -m_D^2/M_M$ and $m_2 = M_M$ and by requiring $m_1 \sim$ O(eV) and $m_2 \sim (10-100)$ GeV and using $v = 246$ GeV we obtain an upper bound estimate for the coupling $y_{\tau i} \leq 10^{-6}$.

The mass eigenstates are denoted by ν_1 and ν_2 and are such that

$$\begin{aligned} \nu_\tau &= \cos\theta\ \nu_{L1} + \sin\theta\ \nu_{R2} \\ N &= -\sin\theta\ \nu_{L1} + \cos\theta\ \nu_{R2}\ , \end{aligned} \tag{7}$$

where $\theta = \sqrt{m_D/m_2} \approx 10^{-(5-6)}$.

The relevant terms in the Lagrangian become

$$\begin{aligned} \mathscr{L} \supset & \left[h\bar{\nu}^c_{L1}\nu_{L1}\left(-\frac{Z}{2\sqrt{2}}s^2_\theta s_\alpha\right) + h\bar{\nu}^c_{R2}\nu_{R2}\left(-\frac{Z}{2\sqrt{2}}c^2_\theta s_\alpha\right) + h.c.\right] \\ + & \ h\bar{\nu}_{L1}\nu_{R2}\left(\frac{y_\nu}{\sqrt{2}}(s^2_\theta - c^2_\theta)c_\alpha\right) + h\bar{\nu}_{R2}\nu_{L1}\left(\frac{y_\nu}{\sqrt{2}}(s^2_\theta - c^2_\theta)c_\alpha\right)\ , \end{aligned} \tag{8}$$

where $y^*_\nu = y_\nu$ and $Z \equiv Z_{11}$.

In this work we are interested in presenting the results for the Higgs decays to neutrinos and their signatures in this model. Using Eq. (8) we compute the following decay widths [1]:

$$\Gamma(h \to \bar{\nu}_1\nu_1) = \frac{m_h}{64\pi}|Z|^2 s^4_\theta s^2_\alpha\ , \tag{9}$$

$$\Gamma(h \to \bar{\nu}_2\nu_2) = \frac{m_h}{64\pi}|Z|^2 c^4_\theta s^2_\alpha \left(1 - \frac{4m_2^2}{m_h^2}\right)^{3/2}\ , \tag{10}$$

$$\Gamma(h \to \bar{\nu}_1\nu_2) = \frac{m_h}{16\pi}y^2_\nu(s^2_\theta - c^2_\theta)^2 c^2_\alpha \left(1 - \frac{m_2^2}{m_h^2}\right)^2\ . \tag{11}$$

We have computed the branching ratios for the Higgs decays and the results are presented in Figure 1. In each plot we have included the results for three values of $\cos\alpha$ (0.1, 0.5 and 0.9). The two graphs correspond to the values of $m_2 = 60$ and 100 GeV respectively. Only the dominant contributions are shown for clarity, i.e. $h \to \nu_2\bar{\nu}_2$, $b\bar{b}$ and $\tau\bar{\tau}$. It is interesting to note that for the whole range where it is possible, the decay $h \to \nu_2\bar{\nu}_2$ dominates in all three cases. This is a clear distinctive signature of our model. In order to study the specific signatures that would be observed in this scenario, we consider the ν_2 decays. In Table 1 we present the possible signatures of these decays.

Since we are interested in a Higgs mass in the natural window of $100-200$ GeV, and in neutrino masses such that they can appear in Higgs decays, we will consider neutrino masses of order $10-100$ GeV, therefore we need to consider the 3-body decays $\nu_2 \to \nu_1 + V^*(\to f\bar{f}')$, where $V^* = W^*, Z^*$:

$$\Gamma = \frac{m_2^5}{256\pi^3}\frac{5}{16}\frac{(B^2+C^2)(a_f^2+b_f^2)}{M_V^4}\ , \tag{12}$$

[1] All SM decay widths will have an extra factor of c^2_α

TABLE 1. Signatures for the Higgs decays considered in the text.

Higgs decay	$\nu_2 \to \nu_1 Z^*$	$\nu_2 \to l W^*$	$\nu_2 \to \nu_1 \gamma$
$h \to \nu_1 \nu_2$	$l^+ l^- + \text{inv.}$	$l + l' + \text{inv.}$	$\gamma + \text{inv.}$
	$q\bar{q} + \text{inv.}$	$l + q\bar{q}' + \text{inv.}$	
$h \to \nu_2 \nu_2$	$l^+ l^- + l^+ l^- + \text{inv.}$	$l + l' + l'' + l''' + \text{inv.}$	
	$l^+ l^- + q\bar{q} + \text{inv.}$	$l + l' + l'' + q\bar{q} + \text{inv.}$	$\gamma + \gamma + \text{inv.}$
	$q\bar{q} + q\bar{q} + \text{inv.}$	$l + l' + q\bar{q} + q\bar{q} + \text{inv.}$	
$h \to \nu_1 \nu_1$	-	-	-

where

$$(V = W) \quad \to \quad \begin{cases} a_f = -b_f \equiv a = \frac{g}{2\sqrt{2}} \\ B = -C = a\, s_\theta \end{cases} \qquad (V = Z) \to \begin{cases} a_f = \frac{g}{2c_w}(T_f^3 - 2Q_f s_w^2) \\ b_f = -\frac{g}{2c_w} T_f^3 \\ B = a_\nu\, c_\theta s_\theta \\ C = b_\nu\, c_\theta s_\theta \end{cases}$$

The branching ratios for these processes are presented in table 2. We show the results for $m_2 = 100$ GeV as the results are similar in all the m_2 range considered in this paper. We find that the dominant contributions are the ones associated to the W^* decay process.

TABLE 2. Branching ratios for the ν_2 three body decays discussed in the text. The results correspond to $m_2 = 100$ GeV and do not depend strongly on the value of m_2.

$m_2(GeV)$	$\nu\, l^+\, l^-$	$\nu\, \nu\, \nu$	$\nu\, q_u\, \bar{q}_u$	$\nu\, q_d\, \bar{q}_d$	$l^\pm\, l^\pm\, \nu$	$l^\pm\, q\, \bar{q}'$
100	0.008	0.015	0.018	0.034	0.308	0.617

MODEL WITH HIGGS TRIPLETS

We now review the basic structure of the second model. The full description of the scalar sector involving the triplet fields can be found in [4, 6, 7, 8], here we briefly review the extension of the basic model to include electroweak neutrinos.

In addition to the SM particle content the model of [3] contains the additional fields shown in table 3. There is also an additional global $\mathrm{U}(1)_M$ symmetry under which

$$L_R^M,\, e_L^M \to e^{i\theta_M} L_R^M,\, e_L^M;\;\; \tilde{\chi} \to e^{-2i\theta_M}\tilde{\chi},\;\; \phi_S \to e^{-i\theta_M}\phi_S\,, \tag{13}$$

and all other fields are singlets. This global symmetry was invoked in order to avoid certain terms as indicated below and was explained in detail in [3].

Since ν_R is not an $\mathrm{SU}(2)_L$ singlet, it does not couple to $\bar{L}_L\tilde{\Phi}$. Instead, the Dirac neutrino mass comes from the term $\mathscr{L}_S = -g_{sl}\bar{L}_L \phi_S L_R^M + h.c.$, which leads to $M_\nu^D = g_{sl} v_s$, where $\langle\phi_S\rangle = v_S$ and thus the neutrino Dirac mass is independent of the EW scale.

RH neutrinos must have a mass $> M_Z/2$ in order not to contribute to the Z width. This is accomplished with the $Y = -2$ triplet $\tilde{\chi}$ through the term $g_M L_R^{M,T}\sigma_2\tau_2\tilde{\chi}L_R^M$, which leads to $M_R = g_M v_M$, with $\langle\chi^0\rangle = v_M$ and where $v_M = O(\Lambda_{EW})$. This allows to have

TABLE 3. Additional field content

Additional fields	SU(2)$_W$	U(1)$_Y$
$L_R^M = \left(\nu_R \ \ e_R^M\right)$	**2**	0
$\tilde{\chi} = \left(\chi^0 \ \ \chi^+ \ \ \chi^{++}\right)^T$	**3**	-2
$\xi = \left(\xi^+ \ \ \xi^0 \ \ \xi^+\right)^T$	**3**	0
e_L^M & ϕ_S	**1**	0

EW-scale masses for the right-handed neutrinos without having to fine-tune the Yukawa coupling g_M to be abnormally small.

The U(1)$_M$ symmetry is introduced in order to forbid the terms $g_L L_L^T \sigma_2 \tau_2 \tilde{\chi} L_L$ and $L_L^T \sigma_2 \tau_2 \tilde{\chi} L_R^M$ at tree level. The main consequence of this is that the Dirac mass for the neutrinos comes from v_s exclusively and the Majorana mass, M_L, for the left-handed neutrinos arises at the one-loop level and can be much smaller than M_R.

Taking all of this into consideration one obtains the following Majorana mass matrix:

$$\mathscr{M} = \begin{pmatrix} M_L & m_\nu^D \\ m_\nu^D & M_R \end{pmatrix} , \tag{14}$$

where $M_L \sim \varepsilon (m_\nu^D)^2/M_R < 10^{-2}(m_\nu^D)^2/M_R$.

We are interested in the scenario where $g_{sl} \sim$ O(g_M) and $v_M >> v_S$. In this case, the eigenvalues of $\mathscr{M}$ become $-(g_{sl}^2/g_M)(v_s/v_m)v_s(1-\varepsilon)$ and M_R, where $\varepsilon < 10^{-2}$. Now, since $v_M \sim \Lambda_{EW}$, and using the bound $m_\nu \leq 1$ eV, we have $v_S \approx \sqrt{(1\text{eV}) \times v_M} \sim$ O$(10^{5-6}$eV).

The kinetic part of the Higgs Lagrangian is given by

$$\mathscr{L}_{kin} = \frac{1}{2} Tr[(D_\mu \Phi)^\dagger (D^\mu \Phi)] + \frac{1}{2} Tr[(D_\mu \chi)^\dagger (D^\mu \chi)] + |\partial_\mu \phi_s|^2 \; . \tag{15}$$

The potential (for Φ and χ)[2] to be considered is [4]

$$\begin{aligned} V(\Phi,\chi) \;&=\; \lambda_1 (Tr\Phi^\dagger\Phi - v_2^2)^2 + \lambda_2 (Tr\chi^\dagger\chi - 3v_m^2)^2 \\ &+\; \lambda_3 (Tr\Phi^\dagger\Phi - v_2^2 + Tr\chi^\dagger\chi - 3v_m^2)^2 \\ &+\; \lambda_4 (Tr\Phi^\dagger\Phi Tr\chi^\dagger\chi - 2Tr\Phi^\dagger T^i \Phi T^j \cdot Tr\chi^\dagger T^i \chi T^j) \\ &+\; \lambda_5 [3Tr\chi^\dagger\chi\chi^\dagger\chi - (Tr\chi^\dagger\chi)^2] \; . \end{aligned} \tag{16}$$

Note that this potential is invariant under $\chi \to -\chi$. When χ gets a vev $\langle\chi\rangle = diag(v_M, v_M, v_M)$ it breaks the global symmetry SU(2)$_L\times$ SU(2)$_R$ down to the custodial SU(2)$_C$. It was shown in [4, 5] that the structure of the VEV is dictated by the proper

[2] We work under the assumption that ϕ_S does not couple with the other Higgses at tree level. We choose to work with this assumption because the coupling generated at loop level, through the ϕ_S couplings to SM left-handed fermions and to mirror right-handed fermions, can be very small [9]

vacuum alignment. Now, using $\langle\Phi\rangle = v_2/\sqrt{2}$, the W and Z masses can be obtained from Eq. (15) and are given by $M_W = gv/2$ and $M_Z = M_W/\cos\theta_W$, with $v^2 = v_2^2 + 8v_M^2$, with $v \approx 246\,\mathrm{GeV}$. This gives rise to $\rho = 1$ at tree level.

A convenient parametrization can be made by defining $\cos\theta_H = c_H \equiv v_2/v$ and thus $\sin\theta_H = s_H \equiv 2\sqrt{2}v_M/v$. Using these parameters we can see that $\tan\theta_H = t_H$ characterizes the amount of the W mass coming from either the doublet or the triplet scalars.

If the potential preserves the SU(2)$_C$ then the fields get arranged in the following manner (based on their transformation properties under the custodial SU(2)):

$$\text{five}-\text{plet} \;\rightarrow\; H_5^{\pm\pm},\, H_5^{\pm},\, H_5^0 \leftrightarrow \text{ degenerate} \tag{17}$$

$$\text{three}-\text{plet} \;\rightarrow\; H_3^{\pm},\, H_3^0 \leftrightarrow \text{ degenerate} \tag{18}$$

$$2-\text{singlets} \;\rightarrow\; H_1^0,\, H_1^{0\prime} \leftrightarrow \text{ Only these can mix },\tag{19}$$

where the definitions and Feynman rules for vector boson couplings can be found in [7]. In the search for the Higgs scalars discussed in this work, it is important to know what those scalars couple to. The couplings of this extended Higgs sector can be found in [5] while the Feynman rules for scalar fermion couplings including the mirror fermions are presented in [2].

In this paper we present the results obtained for the doubly charged Higgs phenomenology. The complete numerical analysis of this model can be found in [2].

The presence of a doubly charged Higgs in this model provides with interesting phenomenology. Furthermore, the phenomenology of this model is specific and different from that of the general two triplets model due to the following observations:

- Due to the U(1)$_M$ symmetry of the model or its embedding in a Pati-Salam type of quark-lepton unification, the term proportional to $l_l^T\sigma_2\tau_2\tilde{\chi}l_L$ is not allowed and thus the decay $\Gamma(\chi^{++}\to l^+l^+)$ is not present.
- The presence of mirror fermions and ϕ_S allows for the decays $\Gamma(\chi^{++}\to l_i^M\, l_j^M)$ and $\Gamma(\chi^{++}\to l\;\phi_S\; l_M)$ or even $\Gamma(\chi^{++}\to ll\phi_S\phi_S)$.

Using the expressions for the χ^{++} decays in [2] we can compute the branching ratios. In the following analysis we have made the following assumptions:

- g_M and g_{sl} are proportional to the identity matrix and so, in each of the expressions above, g_M and g_{sl} represent numbers.
- The model requires $g_{sl}^2/g_M \sim$ O(1). We have chosen numbers of O(1) for both couplings and for the numerical results presented below they have been set to $g_M = 0.7$ and $g_{sl} = 0.8$.

Given these assumptions we compute the following branching ratios: $B(\chi^{++}\to l_M^+l_M^+)$, $B(\chi^{++}\to W^+W^+)$, $B(\chi^{++}\to H_3^+W^+)$, $B(\chi^{++}\to l^+\nu W^+)$ and $B(\chi^{++}\to l^+\phi_S l_M^+)$.

Figure 2 shows the branching ratios for three different values of $\sin\theta_H$ and for small values of the mirror fermions masses (taken to be degenerate) $m_{lM} = 50\,\mathrm{GeV}$. We can see that the dominant one always corresponds to $B(\chi^{++}\to l_M l_M)$, while the relative dominance of the other channels depends on $\sin\theta_H$.

Similar results are obtained for larger m_{lM} as can be seen in figure 3 where we show the branching ratios for $m_{lM} = 100\,\text{GeV}$.

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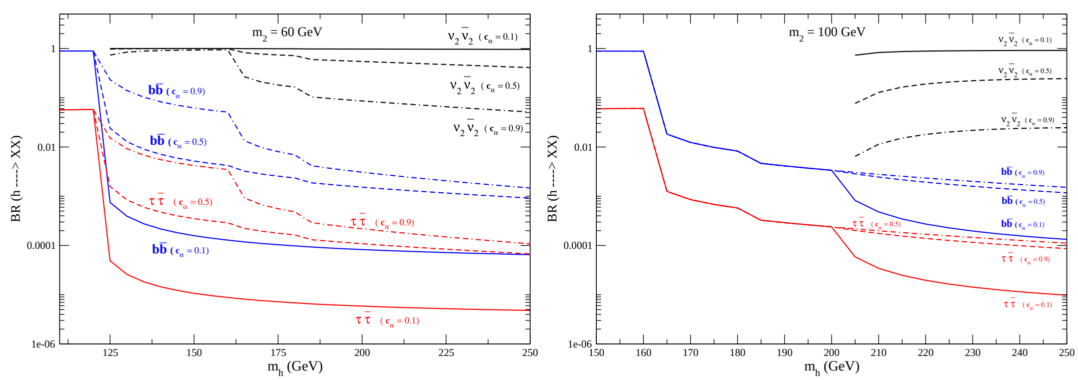

FIGURE 1. Dominant branching ratios for Higgs decays. Two cases are presented for $m_2 = 60$ and 100 GeV respectively. Each plot includes results for the three values of $\cos\theta = 0.1,\ 0.5$ and 0.9 as discussed in the text.

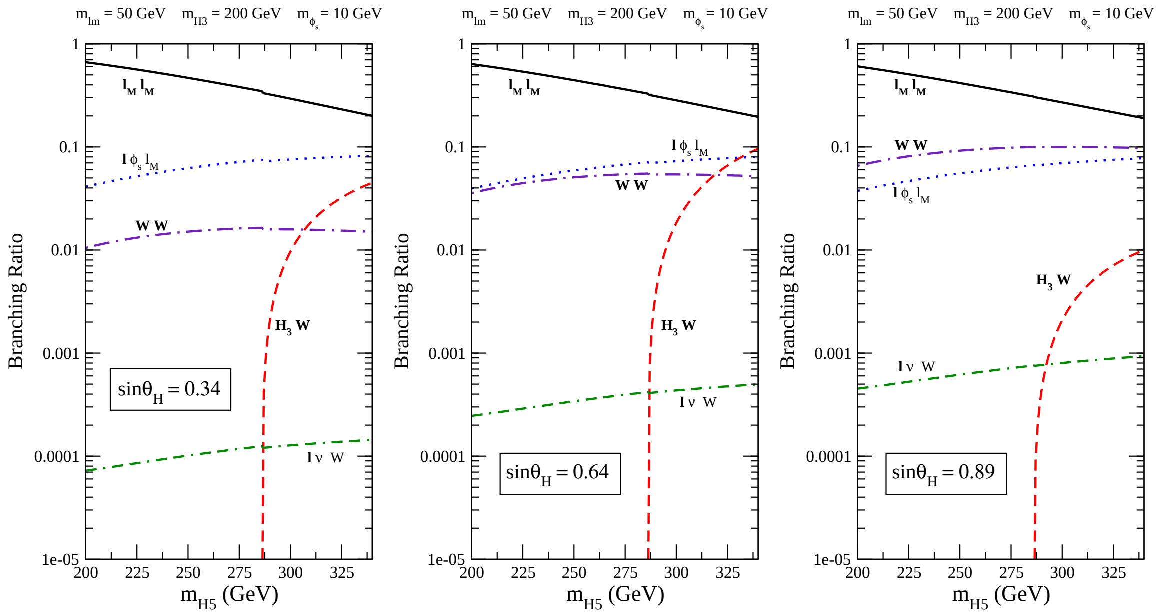

FIGURE 2. Branching ratios for χ^{++} as a function of its mass, for three different values of $\sin\theta_H$, and for a small m_{lM}.

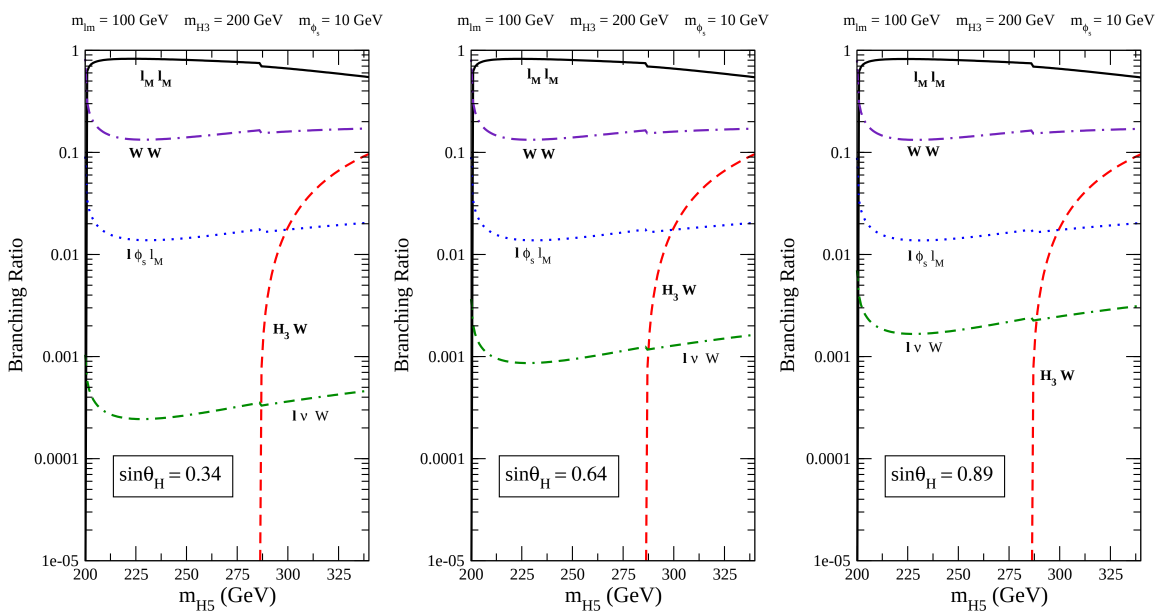

FIGURE 3. Same as before but with a heavier m_{lM}.