Type II radiative seesaw model of neutrino mass with dark matter

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Abstract

We consider a model of neutrino mass with a scalar triplet (ξ^{++}, ξ^+, ξ^0) assigned lepton number $L = 0$, so that the tree-level Yukawa coupling $\xi^0 \nu_i \nu_j$ is not allowed. It is generated instead through the interaction of ξ and ν with dark matter and the soft breaking of L to $(-1)^L$. We discuss the phenomenological implications of this model, including ξ^{++} decay and the prognosis of discovering the dark sector at the Large Hadron Collider.

1 Introduction

Nonzero neutrino mass is necessary to explain the well-established phenomenon of neutrino oscillations in many experiments. Theoretically, neutrino masses are usually assumed to be Majorana and come from physics at an energy scale higher than that of electroweak symmetry breaking of order 100 GeV. As such, the starting point of any theoretical discussion of the underlying theory of neutrino mass is the effective dimension-five operator [1]

$$\mathcal{L}_5 = -\frac{f_{ij}}{2\Lambda}(\nu_i\phi^0 - l_i\phi^+)(\nu_j\phi^0 - l_j\phi^+) + H.c., \tag{1}$$

where $(\nu_i, l_i), i = 1, 2, 3$ are the three left-handed lepton doublets of the standard model (SM) and (ϕ^+, ϕ^0) is the one Higgs scalar doublet. As ϕ^0 acquires a nonzero vacuum expectation value $\langle\phi^0\rangle = v$, the neutrino mass matrix is given by

$$\mathcal{M}_{ij}^\nu = \frac{f_{ij}v^2}{\Lambda}. \tag{2}$$

Note that $\mathcal{L}_5$ breaks lepton number L by two units.

It is evident from Eq. (2) that neutrino mass is seesaw in character, because it is inversely proportional to the large effective scale Λ. The three well-known tree-level seesaw realizations [2] of $\mathcal{L}_5$ may be categorized by the specific heavy particle used to obtain it: (I) neutral fermion singlet N, (II) scalar triplet (ξ^{++}, ξ^+, ξ^0), (III) fermion triplet $(\Sigma^+, \Sigma^0, \Sigma^0)$. It is also possible to realize $\mathcal{L}_5$ radiatively in one loop [2] with the particles in the loop belonging to the dark sector, the lightest neutral one being the dark matter of the Universe. The simplest such example [3] is the well-studied "scotogenic" model, from the Greek 'scotos' meaning darkness. The one-loop diagram is shown in Fig. 1. The new particles are a second scalar doublet (η^+, η^0) and three neutral singlet fermions N_R. The dark Z_2 is odd for (η^+, η^0) and N_R, whereas all SM particles are even. This is thus a Type I radiative seesaw model. It is of course possible to replace N with Σ^0, so it becomes a Type III radiative seesaw model [4]. What then about Type II?
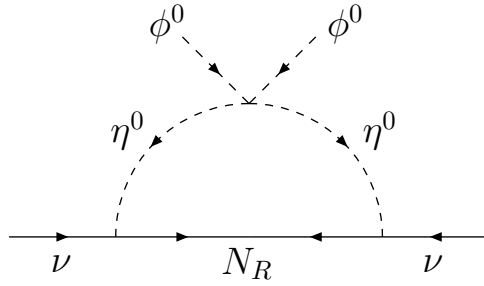
Figure 1: One-loop Z_2 scotogenic neutrino mass.

Since $\mathcal{L}_5$ is a dimension-five operator, any loop realization is guaranteed to be finite. On the other hand, if a Higgs triplet (ξ^{++}, ξ^+, ξ^0) is added to the SM, a dimension-four coupling $\xi^0 \nu_i \nu_j - \xi^+ (\nu_i l_j + l_i \nu_j)/\sqrt{2} + \xi^{++} l_i l_j$ is allowed. As ξ^0 obtains a small vacuum expectation value [5] from its interaction with the SM Higgs doublet, neutrinos acquire small Majorana masses, i.e. Type II tree-level seesaw. If an exact symmetry is used to forbid this dimension-four coupling, it will also forbid any possible loop realization of it. Hence a Type II radiative seesaw is only possible if the symmetry used to forbid the hard dimension-four coupling is softly broken in the loop, as recently proposed [6].

2 Type II Radiative Seesaw Neutrino Masses

The symmetry used to forbid the hard $\xi^0 \nu \nu$ coupling is lepton number $U(1)_L$ under which $\xi \sim 0$. The scalar trilinear $\bar{\xi}^0 \phi^0 \phi^0$ term is allowed and induces a small $\langle \xi^0 \rangle$, but ν remains massless. To connect ξ^0 to $\nu\nu$ in one loop, we add a new Dirac fermion doublet (N, E) with $L = 0$, together with three complex neutral scalar singlets s with $L = 1$. The resulting one-loop diagram is shown in Fig. 2. Note that the hard terms $\xi^0 N N$ and $s \bar{\nu}_L N_R$ are allowed by L conservation, whereas the ss terms break L softly by two units to $(-1)^L$. A dark Z_2 parity, i.e. $(-1)^{L+2j}$, exists under which N, E, s are odd and ν, l, ξ are even. Hence the lightest s is a possible dark-matter candidate. The three s scalars are the analogs of the
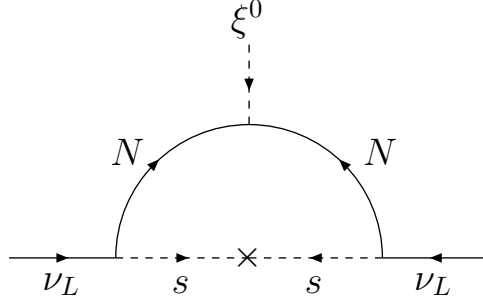
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Figure 2: One-loop neutrino mass from $L = 0$ Higgs triplet.

three right-handed sneutrinos in supersymmetry, and $(N, E)_{L,R}$ are the analogs of the two higgsinos. However, their interactions are simpler here and less constrained.

The usual understanding of the Type II seesaw mechanism is that the scalar trilinear term $\mu \xi^\dagger \Phi \Phi$ induces a small vacuum expectation value $\langle \xi^0 \rangle = u$ if either μ is small or m_ξ is large or both. More precisely, consider the scalar potential of Φ and ξ.

$$
\begin{aligned}
V &= m^2 \Phi^\dagger \Phi + M^2 \xi^\dagger \xi + \frac{1}{2}\lambda_1 (\Phi^\dagger \Phi)^2 + \frac{1}{2}\lambda_2 (\xi^\dagger \xi)^2 + \lambda_3 |2\xi^{++}\xi^0 - \xi^+ \xi^+|^2 \\
&+ \lambda_4 (\Phi^\dagger \Phi)(\xi^\dagger \xi) + \frac{1}{2}\lambda_5 [|\sqrt{2}\xi^{++}\phi^- + \xi^+ \bar{\phi}^0|^2 + |\xi^+ \phi^- + \sqrt{2}\xi^0 \bar{\phi}^0|^2] \\
&+ \mu(\bar{\xi}^0 \phi^0 \phi^0 + \sqrt{2}\xi^- \phi^0 \phi^+ + \xi^{--}\phi^+ \phi^+) + H.c.
\end{aligned}
\tag{3}
$$

Let $\langle \phi^0 \rangle = v$, then the conditions for the minimum of V are given by [5]

$$
m^2 + \lambda_1 v^2 + (\lambda_4 + \lambda_5)u^2 + 2\mu u = 0, \tag{4}
$$

$$
u[M^2 + \lambda_2 u^2 + (\lambda_4 + \lambda_5)v^2] + \mu v^2 = 0. \tag{5}
$$

For $\mu \neq 0$ but small, u is also naturally small because it is approximately given by

$$
u \simeq \frac{-\mu v^2}{M^2 + (\lambda_4 + \lambda_5)v^2}, \tag{6}
$$

where $v^2 \simeq -m^2/\lambda_1$. The physical masses of the $L = 0$ Higgs triplet are then given by

$$
m^2(\xi^0) \simeq M^2 + (\lambda_4 + \lambda_5)v^2, \tag{7}
$$

$$m^2(\xi^+) \simeq M^2 + (\lambda_4 + \frac{1}{2}\lambda_5)v^2, \tag{8}$$

$$m^2(\xi^{++}) \simeq M^2 + \lambda_4 v^2. \tag{9}$$

Since the hard term $\xi^0 \nu \nu$ is forbidden, u by itself does not generate a neutrino mass. Its value does not have to be extremely small compared to the electroweak breaking scale. For example $u \sim 0.1$ GeV is acceptable, because its contribution to the precisely measured ρ parameter $\rho_0 = 1.00040 \pm 0.00024$ [7] is only of order 10^{-6}. With the soft breaking of L to $(-1)^L$ shown in Fig. 2, Type II radiative seesaw neutrino masses are obtained. Let the relevant Yukawa interactions be given by

$$\mathcal{L}_Y = f_s s \bar{\nu}_L N_R + \frac{1}{2} f_R \xi^0 N_R N_R + \frac{1}{2} f_L \xi^0 N_L N_L + H.c., \tag{10}$$

together with the allowed mass terms $m_E(\bar{N}N + \bar{E}E)$, $m_s^2 s^* s$, and the L breaking soft term $(1/2)(\Delta m_s^2)s^2 + H.c.$, then

$$m_\nu = \frac{f_s^2 u r x}{16\pi^2} [f_R F_R(x) + f_L F_L(x)], \tag{11}$$

where $r = \Delta m_s^2/m_s^2$ and $x = m_s^2/m_E^2$, with

$$F_R(x) = \frac{1+x}{(1-x)^2} + \frac{2x \ln x}{(1-x)^3}, \tag{12}$$

$$F_L(x) = \frac{2}{(1-x)^2} + \frac{(1+x)\ln x}{(1-x)^3}. \tag{13}$$

Using for example $x \sim f_R \sim f_L \sim 0.1$, $r \sim f_s \sim 0.01$, we obtain $m_\nu \sim 0.1$ eV for $u \sim 0.1$ GeV. This implies that ξ may be as light as a few hundred GeV and be observable, with $\mu \sim 1$ GeV. For $f_s \sim 0.01$ and m_E a few hundred GeV, the new contributions to the anomalous muon magnetic moment and $\mu \to e\gamma$ are negligible in this model.

In the case of three neutrinos, there are of course three s scalars. Assuming that the L breaking soft terms $|(\Delta m_s^2)_{ij}| << |m_{s_i}^2 - m_{s_j}^2|$ for $i \neq j$, then the 3×3 neutrino mass matrix is diagonal to a very good approximation in the basis where the s mass-squared matrix is

diagonal. This means that the dark scalars s_j couples to $U_{ij}l_i$, where U_{ij} is the neutrino mixing matrix linking e, μ, τ to the neutrino mass eigenstates $\nu_{1,2,3}$.

3   Doubly Charged Higgs Production and Decay

The salient feature of any Type II seesaw model is the doubly charged Higgs boson ξ^{++}. If there is a tree-level $\xi^{++}l_i^- l_j^-$ coupling, then the dominant decay of ξ^{++} is to $l_i^+ l_j^+$. Current experimental limits [8] on the mass of ξ^{++} into $e\mu$, $\mu\mu$, and ee final states are about 490 to 550 GeV, assuming for each a 100% branching fraction. In the present model, since the effective $\xi^{++}l_i^- l_j^-$ coupling is one-loop suppressed, $\xi^{++} \rightarrow W^+ W^+$ should be considered [9] instead, for which the present limit on $m(\xi^{++})$ is only about 84 GeV [10]. A dedicated search of the $W^+ W^+$ mode in the future is clearly called for.

If $m(\xi^{++}) > 2m_E$, then the decay channel $\xi^{++} \rightarrow E^+ E^+$ opens up and will dominate. In that case, the subsequent decay $E^+ \rightarrow l^+ s$, i.e. charged lepton plus missing energy, will be the signature. The present experimental limit [11] on m_E, assuming electroweak pair production, is about 260 GeV if $m_s < 100$ GeV for a 100% branching fraction to e or μ, and no limit if $m_s > 100$ GeV. There is also a lower threshold for ξ^{++} decay, i.e. $m(\xi^{++})$ sufficiently greater than $2m_s$, for which ξ^{++} decays through a virtual $E^+ E^+$ pair to ssl^+l^+, resulting in same-sign dileptons plus missing energy.

In Fig. 3 we plot the pair production cross section of $\xi^{++}\xi^{--}$ at the Large Hadron Collider (LHC) at a center-of-mass energy of 13 TeV. We assume that ξ^+ and ξ^0 are heavier than ξ^{++} so that we can focus only on the decay products of $\xi^{\pm\pm}$. The $W^{\pm}W^{\pm}$ mode is always possible and should be looked for experimentally in any case. However, as already noted, a much more interesting possibility is the case $m(\xi^{++}) > 2m_E$, with the subsequent decay $E^+ \rightarrow l^+ s$. This would yield four charged leptons plus missing energy, and depending on the
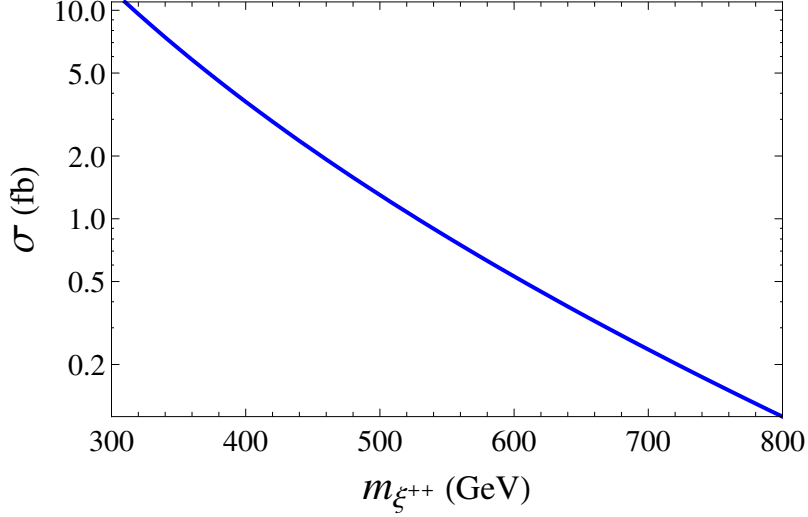
Figure 3: LHC Production cross section of $\xi^{++}\xi^{--}$ at 13 TeV.

linear combination of charged leptons coupling to s, there could be exotic final states which have very little SM background, becoming thus excellent signatures to search for. Suppose s_1 is the lightest scalar, and $s_{2,3}$ are heavier than E^+, then E^+ decays to $s_1 \sum U_{i1} l_i^+$. Hence the decay of $\xi^{++}\xi^{--}$ could yield for example $e^+ e^+ \mu^- \mu^-$ plus four s_1 (missing energy) in the final state.

Recent LHC searches for multilepton signatures at 8 TeV by CMS [12] and ATLAS [13] are consistent with SM expectations, and are potential restrictions on our model. In particular, the CMS study includes rare SM events such as $e^+ e^+ \mu^- \mu^-$ and $e^+ e^+ \mu^-$. Due to the absence of opposite-sign, same-flavor (OSSF) $l^+ l^-$ pairs, both events are classified as OSSF0 where lepton l refers to electron, muon, or hadronically decaying tau. Leptonic tau decays contribute to the electron and muon counts, and this determines the OSSFn category. Details from CMS are shown in Table 1 for ≥ 3 leptons and $N_{\tau\text{had}} = 0$. The CMS study estimates a negligible SM background for SR1-SR3, and in our simulation we use the same selection criteria. We impose the cuts on transverse momentum $p_T > 10$ GeV and psuedo-rapidity $|\eta| < 2.4$ for each charged lepton, with at least one lepton $p_T > 20$ GeV. In order

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Selected CMS results OSSF0 $\quad N_{\tau\text{had}} = 0$, $N_b = 0$					
signal regions		$H_T > 200$ GeV		$H_T < 200$ GeV	
≥ 4 leptons	$\not{E}_T$ (GeV)	Obs.	Exp.(SM)	Obs.	Exp.(SM)
SR1	$(100, \infty)$	0	$0.01^{+0.03}_{-0.01}$	0	$0.11^{+0.08}_{-0.08}$
SR2	$(50, 100)$	0	$0.00^{+0.02}_{-0.00}$	0	$0.01^{+0.03}_{-0.01}$
SR3	$(0, 50)$	0	$0.00^{+0.02}_{-0.00}$	0	$0.01^{+0.02}_{-0.01}$
3 leptons	$\not{E}_T$ (GeV)	Obs.	Exp.(SM)	Obs.	Exp.(SM)
SR4	$(100, \infty)$	5	3.7 ± 1.6	7	11.0 ± 4.9
SR5	$(50, 100)$	3	3.5 ± 1.4	35	38 ± 15
SR6	$(0, 50)$	4	2.1 ± 0.8	53	51 ± 11

Table 1: Events observed by CMS at 8 TeV with integrated luminosity 19.5 fb^{-1}.

to be isolated, each lepton with p_T must satisfy $\sum_i p_{Ti} < 0.15 p_T$, where the sum is over all objects within a cone of radius $\Delta R = 0.3$ around the lepton direction.

We implement our model with FeynRules 2.0 [14]. Using the CTEQ6L1 parton distribution functions, we generate events using MadGraph5 [15], which includes the Pythia package for hadronization and showering. MadAnalysis [16] is then used with the Delphes card designed for CMS detector simulation. Generated events intially have 4 leptons. About half are detected as 3 lepton events, but the constraints from signal regions SR4-SR6 are less restrictive than SR1-SR3. The number of detected events in the OSSF0 ≥ 4 lepton category is almost the same as $e^\pm e^\pm \mu^\mp \mu^\mp 2s_1 2s_1^*$ with very few additional leptons from showering or initial/final state radiation.

To examine the production of $e^\pm e^\pm \mu^\mp \mu^\mp$ we take the mass of s_1 to be 130 GeV, which allows s_1 to be dark matter as discussed in the next section. We use the values $f_R = f_L = 0.1$ and $f_s = 0.01$, although the results are not sensitive to the exact values due to on-shell production and decay. The effects due to $u \sim 0.1$ GeV may be neglected.

For our model, we scan the mass range of ξ^{++} and E^+. In Fig. 4 we plot contours showing

the expected number of detected events in the OSSF0 ≥ 4 lepton category for 13 TeV at luminosity 100 fb^{-1} assuming a negligible background as for the 8 TeV case. Although the branching fractions of E^+ to $\tau^+ s_1$ or $\mu^+ s_1$ are comparable, we find that most of the contributions from $\tau^\pm$ decay to $e^\pm$ or $\mu^\pm$ in the ≥ 4 lepton final state are not detected. A similar analysis performed for 8 TeV at 19.5 fb^{-1} has a maximum number of detected events of 0.4 in the plot analogous to Fig. 4, which corresponds to a small estimated exclusion at the 15% confidence level.
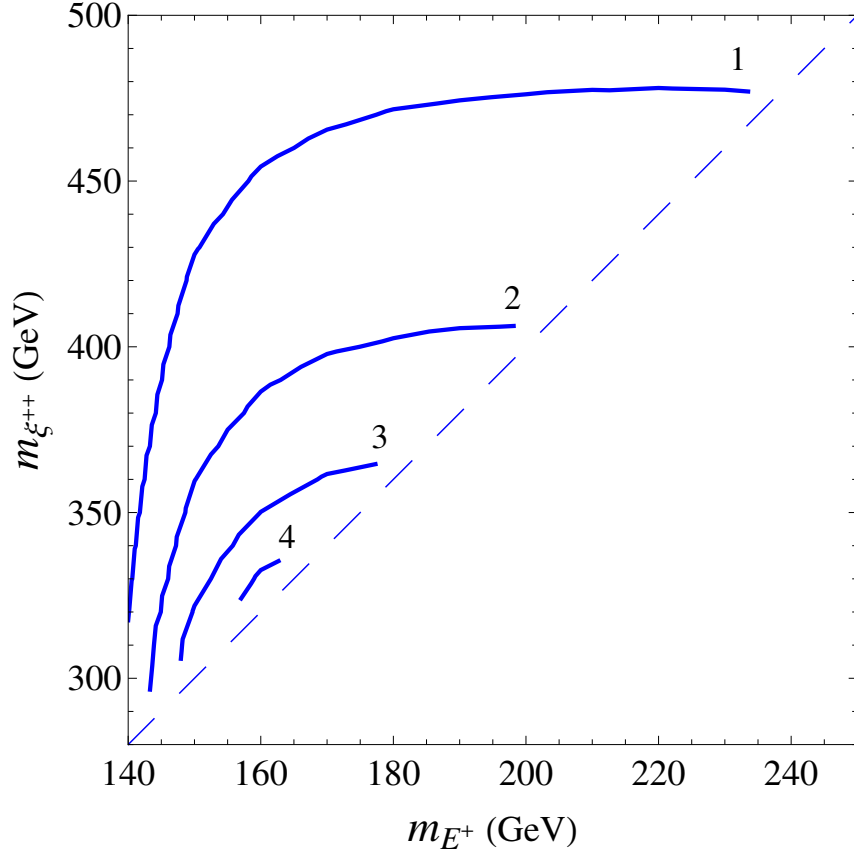

Figure 4: Number of $e^\pm e^\pm \mu^\mp \mu^\mp 2s_1 2s_1^*$ events for 13 TeV at luminosity 100 fb^{-1}.

4 Dark Matter Properties

The lightest s, say s_1, is dark matter. Its interaction with leptons is too weak to provide a large enough annihilation cross section to explain the present dark matter relic density Ω_M of the Universe. However, it also interacts with the SM Higgs boson through the usual quartic coupling $\lambda_s s^* s \Phi^\dagger \Phi$. For a value of λ_s consistent with Ω_M, the direct-detection cross section in underground experiments is determined as a function of m_s. A recent analysis [17] for a real s claims that the resulting allowed parameter space is limited to a small region near $m_s < m_h/2$.

In our model, we can evade this constraint by evoking $s_{2,3}$. The mass-squared matrix spanning $s_i^* s_j$ is given by

$$(\mathcal{M}_s^2)_{ij} = m_{ij}^2 + \lambda_{ij} v^2, \tag{14}$$

whereas the coupling matrix of the one Higgs h to $s_i^* s_j$ is $\lambda_{ij} v \sqrt{2}$. Upon diagonalizing $\mathcal{M}_s^2$, the coupling matrix will not be diagonal in general. In the physical basis, s_1 will interact with s_2 through h. This allows the annihilation of $s_1 s_1^*$ to hh through s_2 exchange, and contributes to Ω_M without affecting the s_1 scattering cross section off nuclei through h. This mechanism restores s_1 as a dark-matter candidate for $m_s > m_h$.

To demonstrate the scale of the values involved, we consider the simplifying case when $m_{s_2} = m_{s_3}$ and $\lambda_{12} = \lambda_{13}$. The additional choice $m_{s_{2,3}}^2 = m_{s_1}^2 + m_h^2$ ensures that $s_{2,3}$ are heavier than s_1, and is convenient because then the relic abundance requirement no longer depends explicitly on $m_{s_{2,3}}^2$. Taking into account that s_1 is a complex scalar, we use $\sigma \times v_{rel} = 4.4 \times 10^{-26} \mathrm{cm}^3 \mathrm{s}^{-1}$ [18] and in Fig. 5 we plot the allowed values for λ_{12} and m_{s_1} taking $\lambda_{11} = 0$ for simplicity to satisfy the LUX data.

Another possible scenario is to add a light scalar χ with $L = 0$, which acts as a mediator for s self-interactions. This has important astrophysical implications [19, 20, 21, 22, 23, 24].
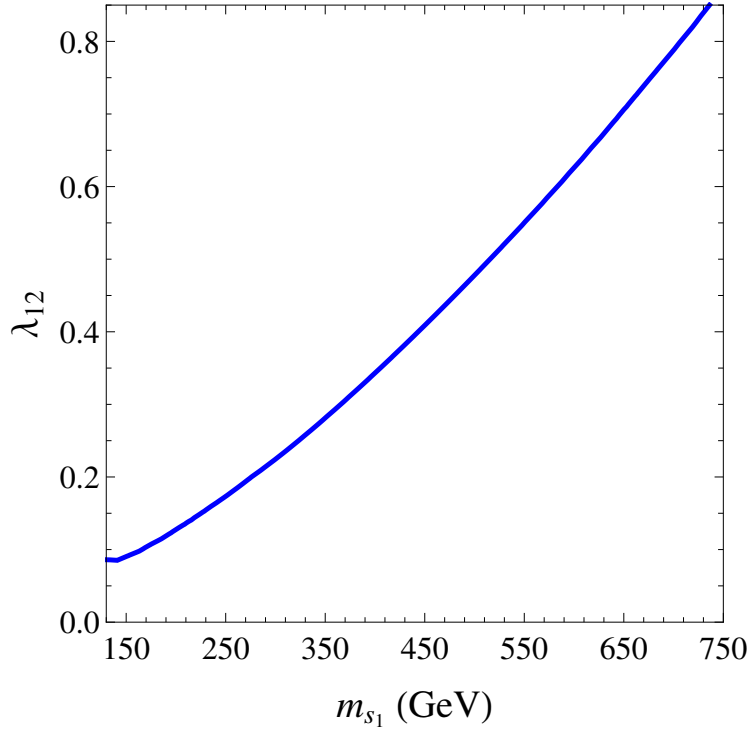
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Figure 5: Allowed values of λ_{12} plotted against m_{s_1} from relic abundance assuming $\lambda_{11} = 0$.

In this case, $s_1 s_1^*$ annihilating to $\chi\chi$ becomes possible.

5   Conclusion

We have studied a new radiative Type II seesaw model of neutrino mass with dark matter [6], which predicts a doubly charged Higgs boson ξ^{++} with suppressed decay to l^+l^+, thereby evading the present LHC bounds of 490 to 550 GeV on its mass. In this model, ξ^{++} may decay to two charged heavy fermions E^+E^+, each with odd dark parity. The subsequent decay of E^+ is into a charged lepton l^+ and a scalar s which is dark matter. Hence there is the interesting possibility of four charged leptons, such as $\mu^-\mu^-e^+e^+$, plus large missing energy in the final state. We show that the LHC at 13 TeV will be able to probe such a doubly charged Higgs boson with a mass of the order 400 to 500 GeV.

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