Low scale leptogenesis and dark matter candidates in an extended seesaw model

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Abstract. We consider a variant of the seesaw mechanism by introducing extra singlet neutrinos and a singlet scalar boson, and show how low scale leptogenesis is successfully realized in this scenario. We examine whether the newly introduced neutral particles, either singlet Majorana neutrinos or singlet scalar bosons, can be dark matter candidates. We also discuss the implications of dark matter detection through scattering off the nucleus of the detecting material on our scenarios for dark matter. In addition, we study the implications for the search for invisible Higgs decay at the Large Hadron Collider, which may serve as a probe for our scenario for dark matter.

Keywords: dark matter detectors, dark matter, neutrino properties, baryon asymmetry
1. Introduction

Two important unsolved issues in particle physics and cosmology are why there is more matter than antimatter in the present Universe and what the origin of dark matter is. In this paper, we propose a model for addressing both of those problems and show that they can be solved by means of a common origin.

One of the most popular models for accommodating the right amount of baryon asymmetry in the present Universe is so-called leptogenesis [1], which is realized in the context of the seesaw mechanism [2], and thus smallness of neutrino masses and baryon asymmetry can be simultaneously achieved. However, typical leptogenesis demands a rather large scale of the right-handed Majorana neutrino mass, which makes it impossible to probe in present experimental laboratories. We have recently proposed a variant of the seesaw mechanism to lower the scale of leptogenesis [3] and showed that the required leptogenesis is allowed at a low scale—even a few TeV scale—without imposing the tiny mass splitting between two heavy Majorana neutrinos required in resonant leptogenesis [4].

The main point of our model proposed previously is to introduce an equal number of gauge singlet neutrinos in addition to the heavy right-handed singlet neutrinos. In our scenario, there exist new Yukawa interactions mediated by a singlet Higgs sector which is coupled with the extra singlet neutrinos and the right-handed singlet neutrinos. As shown in [3], the new Yukawa interactions may play a crucial role in enhancing the lepton asymmetry, such that low scale leptogenesis can be achieved.

On the other hand, it is worthwhile to examine whether those new particles, either the newly introduced singlet neutrinos or the singlet scalar bosons, can be dark matter candidates because all kinds of neutral, stable and weakly interacting massive particles (WIMPs) can be regarded as good dark matter candidates. While light singlet neutrinos with mass of order MeV or keV have been considered as warm dark matter candidates [5], heavy singlet neutrinos with mass of order 100 GeV have not been considered as dark matter candidates in much detail. While the cold dark matter (CDM) models
supplemented with a cosmological constant are in a good agreement with the observed structure of the Universe on large scales, the cosmic microwave background anisotropies and type Ia supernova observations for a given set of density parameters, there exists a growing wealth of observational data which are in conflict with the CDM scenarios [6]. To remedy the difficulties of the CDM models on galactic scales, a self-interacting dark matter candidate has been proposed [7]–[9], and it has been discussed that a gauge singlet scalar coupled to the Higgs boson, leading to an invisible decaying Higgs, is a good self-interacting dark matter candidate [10]–[12].

In this paper, we will show that either the new gauge singlet neutrinos or the singlet scalar bosons, which play an important role in enhancing the lepton asymmetry with the result that the low scale leptogenesis is realized, could be a good dark matter candidate. For our purpose, we will first present how the enhancement of the lepton asymmetry through the mediation of the newly introduced particles in our framework can be achieved and then we will show that the new gauge singlet particles with a parameter space consistent with low scale leptogenesis can be satisfied with the criteria for the dark matter candidate. We will investigate possibilities for dark matter detection through the scattering off the nucleus of the detecting material. In addition, we will study how we can probe our scenarios at high energy colliders and present why the search for the invisible Higgs decay may serve as a probe of dark matter properties.

2. Extended seesaw model and low scale leptogenesis

We begin by explaining what the extended seesaw model is and examine how low scale leptogenesis can be realized in this context. The Lagrangian that we propose is given in the charged lepton basis as

\[ \mathcal{L}_f = Y_{\nu ij} \bar{\nu}_i H N_j + M_{R ij} N_i N_i + Y_{S ij} \bar{N}_i \Phi S_j - m_{S ij} S_i S_j + h.c., \]  

(1)

where \( \nu_i, N_i, S_i \) stand for the \( SU(2)_L \) doublet, right-handed singlet and newly introduced singlet neutrinos, respectively. \( Y_{\nu ij}, Y_{S ij}, M_{R ij} \) represent the Dirac–Yukawa coupling matrix, singlet Yukawa coupling matrix and heavy Majorana neutrino mass matrix, respectively. And \( H \) and \( \Phi \) denote the \( SU(2)_L \) doublet and singlet Higgs scalars. Here, we impose \( Z_2 \) symmetry under which \( S_i \) and \( \Phi \) are odd and all other particles even, which makes this model different from the extended double-seesaw model proposed in [3] even though the contents of particles are the same. The immediate consequence of the exact \( Z_2 \) symmetry is that the lightest \( S_i \) or \( \Phi \) can be a candidate for providing the cold dark matter of the Universe.

Due to the exact \( Z_2 \) symmetry, the singlet scalar field \( \Phi \) cannot drive a vacuum expectation value. Thus, in this model, light neutrino masses are generated by a typical seesaw mechanism, which makes this model different from the model in [3]. After integrating out the right-handed heavy neutrino sector \( N_R \) in the above Lagrangian, the light neutrino masses are given by

\[ m_\nu = \left( \frac{Y_{\nu i} v_{\text{EW}}}{4 M_R} \right)^2, \]  

(2)

where we omitted the indices of the mass matrix and Dirac–Yukawa coupling, and \( v_{\text{EW}} = 246 \text{ GeV} \) is the Higgs vacuum expectation value. Although the absolute values of
the three neutrino masses are unknown, their masses are expected to be of the order of \(\sqrt{\Delta m_{\text{atm}}} \approx 0.05\) and \(\sqrt{\Delta m_{\text{sol}}} \approx 0.01\) eV, provided that the mass spectrum of neutrinos is hierarchical. There is also a bound on neutrino masses coming from WMAP observation, which is \(m_\nu < 0.23\) eV. Thus it is interesting to see how the neutrino masses of the order of 0.01–0.1 eV can be obtained in our scenario. Such light neutrino masses can be generated through the seesaw formula, equation (2), if we take, as an example, \(Y_\nu\) and \(M_R\) to be of order \(10^{-6}\) and \(10^4\) GeV, respectively.

Now, let us consider how low scale leptogenesis can be achieved by the decay of the lightest right-handed Majorana neutrino in our framework. Right-handed heavy Majorana neutrinos are even under \(Z_2\) symmetry, so they can decay into a pair of the singlet particles, \(\Phi\) and \(S\), or the standard model Higgs and the lepton doublet. Since the Yukawa couplings \((Y_S)_{2(3)i}\) are taken to be large, the processes involving \(N_{2,3}\) remain in thermal equilibrium even at \(T \simeq M_{R_1}\), and thus the decays of \(N_{2,3}\) cannot lead to the desired baryon asymmetry. However, the decay processes of the lightest right-handed Majorana neutrino \(N_1\) depart from thermal equilibrium at \(T \lesssim M_{R_1}\), and thus lead to the desired baryon asymmetry.

It will be shown that there exists a new contribution to the lepton asymmetry which is mediated by the extra singlet neutrinos \(S_i\) and scalar boson \(\Phi\), and successful leptogenesis can be realized with rather light right-handed Majorana neutrino masses which can escape the gravitino problem encountered in the supersymmetric standard model. Without loss of generality, we can rotate and re-phase the fields to make the mass matrices \(M_{R_{ij}}\) and \(m_{S_{ij}}\) real and diagonal. In this basis, the elements of \(Y_\nu\) and \(Y_S\) are in general complex.

The lepton number asymmetry required for baryogenesis is given by

\[
\varepsilon_1 = - \sum_i \left[ \frac{\Gamma(N_1 \rightarrow l_i \overline{H}) - \Gamma(N_1 \rightarrow l_i H)}{\Gamma_{\text{tot}}(N_1)} \right],
\]

where \(N_1\) is the lightest right-handed neutrino and \(\Gamma_{\text{tot}}(N_1)\) is the total decay rate. Thanks to the new Yukawa interactions, there is a new contribution of the diagram which corresponds to the self-energy correction of the vertex which has arisen due to the new Yukawa couplings with singlet neutrinos and Higgs scalars. Figure 1 shows the structure of the diagrams contributing to \(\varepsilon_1\). Assuming that the masses of the Higgs scalars and the newly introduced singlet neutrinos are much smaller compared to that of the right-handed neutrino, to leading order, we have

\[
\Gamma_{\text{tot}}(N_i) = \frac{(Y_\nu Y_\nu^\dagger + Y_S Y_S^\dagger)_{ii}}{4\pi} M_{R_i},
\]

and so

\[
\varepsilon_1 = \frac{1}{8\pi} \sum_{k \neq 1} \left[ \left( g_V(x_k) + g_S(x_k) \right) T_{k1} + g_S(x_k) S_{k1} \right],
\]

where \(Y_\nu\) and \(Y_S\) are given in the basis where \(M_R\) and \(M_S\) are diagonal, \(g_V(x) = \sqrt{x}\{1 - (1 + x)\ln[(1 + x)/x]\}, \quad g_S(x) = \sqrt{x_k/(1 - x_k)}\) with \(x_k = M_{R_k}^2/M_{R_1}^2\) for \(k \neq 1\),

\[
T_{k1} = \frac{\text{Im}[(Y_\nu Y_\nu^\dagger)_{k1}^2]}{(Y_\nu Y_\nu^\dagger + Y_S Y_S^\dagger)_{11}}
\]
Figure 1. Diagrams contributing to the lepton asymmetry.

and

\[ S_{k1} = \frac{\text{Im}[(Y_{\nu}^* Y_{\nu}'\phi)_{k1}(Y_{\nu}' Y_{\nu})_{11}]}{(Y_{\nu} Y_{\nu}')_{11} + Y_{S} Y_{S}')_{11}}. \]  

(7)

Notice that the term proportional to \( S_{k1} \) comes from the interference of the tree-level diagram with the new contribution.

In particular, for \( x \approx 1 \), the vertex contribution to \( \varepsilon_1 \) is much smaller than the contribution of the self-energy diagrams and the asymmetry \( \varepsilon \) is resonantly enhanced, and we do not consider this case. To see to what extent the new contribution may be important in this case, for simplicity, we consider a particular situation where \( M_{R_1} \approx M_{R_2} < M_{R_3} \), so that the effect of \( N_3 \) is negligibly small. In this case, the asymmetry can be written as

\[ \varepsilon_1 \approx -\frac{1}{16\pi} \left[ \frac{M_{R_2} \text{Im}[(Y_{\nu}^* m_{\nu} Y_{\nu})_{11}]}{v^2 (Y_{\nu} Y_{\nu}')_{11} + Y_{S} Y_{S}')_{11}} + \sum_{k \neq 1} \frac{\text{Im}[(Y_{\nu} Y_{\nu}')_{k1}(Y_{S} Y_{S}')_{11}]}{(Y_{\nu} Y_{\nu}')_{11} + Y_{S} Y_{S}')_{11}} \right] R, \]  

(8)

where \( R \) is a resonance factor defined by \( R \equiv |M_{R_1}|/(|M_{R_2}| - |M_{R_1}|) \). For successful leptogenesis, the size of the denominator of \( \varepsilon_1 \) should be constrained by the out-of-equilibrium condition, \( \Gamma_{N_1} < H|_{T=M_{R_1}} \) with the Hubble expansion rate \( H \), from which the corresponding upper bound on the couplings \( (Y_{S})_{11} \) reads \( \sqrt{\sum_i |(Y_{S})_{i1}|^2} < 3 \times 10^{-4} \sqrt{M_{R_1}/10^9} \text{ (GeV)} \). Then, the first term of equation (8) is bounded, as \( (M_{R_2}/16\pi v^2)\sqrt{\Delta m_{\text{atm}}^2 R} \). So if the first term of equation (8) dominates over the second one, \( R \sim 10^{6-7} \) is required to achieve TeV scale leptogenesis, which implies severe fine-tuning. However, since the size of \( (Y_{S})_{21} \) is not constrained by the out-of-equilibrium condition, a large value of \( (Y_{S})_{21} \) is allowed for which the second term of equation (8) can dominate over the first one and thus the size of \( \varepsilon_1 \) can be enhanced. For example, if we assume that \( (Y_{\nu})_{21} \) is aligned with \( (Y_{S})_{21} \), i.e. \( (Y_{S})_{21} = \kappa (Y_{\nu})_{21} \) with constant \( \kappa \), the upper limit of the second term of equation (8) is given in terms of \( \kappa \) by \( \kappa^2 M_{R_2} \sqrt{\Delta m_{\text{atm}}^2 R}/16\pi v^2 \).
and then we can achieve successful low scale leptogenesis by taking rather large values of \( \kappa \), instead of imposing very tiny mass splitting between \( M_{R1} \) and \( M_{R2} \). The right amount of the asymmetry, \( \varepsilon_1 \sim 10^{-6} \), can be obtained for \( M_{R1} \sim 10^4 \) GeV, provided that \( \kappa = (Y_S)_{2i}/(Y_\nu)_{2i} \sim 10^3 \) and \( M_{R2}/M_{R1} \sim 10 \). We emphasize that such a requirement for the hierarchy between \( Y_\nu \) and \( Y_S \) is much less severe than the fine-tuning of the mass splitting between two heavy Majorana neutrinos required to achieve successful leptogenesis at low scale.

The generated \( B-L \) asymmetry is given by \( Y^*_B = -\eta \varepsilon_1 Y^*_N \), where \( Y^*_N \) is the number density of the right-handed heavy neutrino at \( T \gg M_{R1} \) in thermal equilibrium given by \( Y^*_N \approx (45/\pi^4)(\zeta(3)/g_s k_B)^{1/2} \) with Boltzmann constant \( k_B \) and the effective number of degrees of freedom \( g^* \). The efficient factor \( \eta \) can be computed through a set of coupled Boltzmann equations which take into account processes that create or wash out the asymmetry. To a good approximation the efficiency factor depends on the effective neutrino mass \( m_1 \) given in the presence of the new Yukawa interactions with the coupling \( Y_S \) by

\[
\hat{m}_1 = \frac{Y_\nu Y^*_\nu + Y_S Y^*_S}{M_{R1}} v^2. \tag{9}
\]

In our model, the new process of type \( S\Phi \rightarrow IH \) will contribute to wash-out of the \( B-L \) asymmetry produced. The process occurs mainly through virtual \( N_{2,3} \) exchanges because the Yukawa couplings \( (Y_S)_{2(3)i} \) are taken to be large in our model and the rate is proportional to \( M_{R1} |Y_S Y^*_\nu /M_{R_{2,3}}|^2 \). The effect of the wash-out can be easily estimated from the fact that it looks similar to the case for the typical seesaw model if \( M_{R1} \) is replaced with \( M_{R1} (Y_S/Y_\nu)^2 \). It turns out that the wash-out factor for \( (Y_S)_{3i} \sim (Y_\nu)_{1i} (Y_S)_{2i}/(Y_\nu)_{2i} \sim 10^3 \) and \( M_{R1} \sim 10^4 \)GeV is similar to that for the case of the typical seesaw model with \( M_{R1} \sim 10^4 \)GeV and \( \hat{m}_1 \sim 10^{-3} \) eV, and it is estimated such that \( \varepsilon_1 \sim 10^{-6} \) can be enough to explain the baryon asymmetry of the Universe provided that the initial lightest right-handed neutrinos are thermal [13].

### 3. Investigation of possible dark matter candidates

Now, let us examine whether either the newly introduced singlet Majorana neutrino or the singlet scalar boson can be a dark matter candidate. For our purpose, in addition to the Lagrangian \( \mathcal{L}_t \) given in equation (1), we allow quartic scalar interactions for the scalar sectors. Then, the Lagrangian that we consider is given by

\[
\mathcal{L} = \mathcal{L}_t + \frac{1}{2} (\partial_\mu \Phi)^2 - \frac{1}{2} m^2 \Phi^2 - \frac{\lambda_s}{4} \Phi^4 - \lambda H^\dagger H \Phi^2. \tag{10}
\]

Note that the self-interacting coupling of the singlet scalar \( \lambda_s \) is largely unconstrained and thus can be chosen arbitrarily. But, we assume that it should not be so large that the perturbation breaks down. As mentioned before, in order to guarantee the stability of the dark matter candidate, we impose \( Z_2 \) discrete symmetry under which all the standard model (SM) particles are even, whereas the singlet Majorana neutrinos \( S_i \) and singlet scalar boson \( \Phi \) are odd. In addition, we demand that the scalar potential is bounded from below so as to guarantee the existence of a vacuum, and the minimum of the scalar potential must spontaneously break the electroweak gauge group, \( \langle H^0 \rangle \neq 0 \), but must not
break the $Z_2$ symmetry imposed above. After breaking the electroweak gauge symmetry, the singlet scalar $\Phi$-dependent part of the scalar potential is given by

$$V = \frac{1}{2}(m_{\Phi^0}^2 + \lambda v_{EW}^2)\Phi^2 + \frac{\lambda}{4}\Phi^4 + \lambda v_{EW}\Phi^2 h + \frac{\lambda}{2}\Phi^2 h^2,$$

where we have adopted $\sqrt{2}H^\dagger = (h, 0)$ and shifted the Higgs boson $h$ via $h \rightarrow h + v_{EW}$, where $v_{EW} = 246$ GeV is the Higgs vacuum expectation value. The physical mass of $\Phi$, $m_\Phi$, is then given by $m_{\Phi^0}^2 + \lambda v_{EW}^2$. We assume that the spectrum of the singlet neutrinos $S_i$ is not degenerate and that the two heavier ones $S_{2,3}$ are so much heavier than $S_1$ that they could not be dark matter candidates. Here, we notice that only the lightest odd particle (LOP) can be a dark matter candidate in our scenario because the next to lightest odd particle (NOP) can decay into a LOP by cascade decay, as shown in figure 2.

Since there are two kinds of odd particles under $Z_2$ in our scenario, we can classify two possible dark matter candidates according to which particle is the LOP. One case is that where the singlet Majorana neutrino $S_1$ is the LOP and the other is that where $\Phi$ is the LOP. As will be clearer later, the primordial abundance of the dark matter candidate is predicted as a function of masses and coupling $\lambda$. Thus, imposing the preferred values of $\Omega_{DM}h^2$ observed from WMAP, $0.094 < \Omega_{DM}h^2 < 0.128$ [14,15], we can get a strong relation between the mass of the dark matter candidate and the coupling $\lambda$.

### 3.1. Singlet neutrino $S$ as a dark matter candidate

The singlet neutrino $S_1$ can be a dark matter candidate, provided that $m_{S_1} \lesssim m_\Phi$. We omit the generation index 1 of $S_1$ hereafter. To estimate the relic abundance of the singlet neutrino $S$ at a freeze-out temperature $x_f = m_S/T_f$, we need to know what the annihilation processes are [16]. There are two possible annihilation processes for the singlet neutrino: one happens at loop level and the other is mediated by $\Phi$. But, both possibilities are irrelevant to fitting the required present relic density of the dark matter because the annihilation cross section for the dark matter candidate $S$ is too small; therefore, these processes would predict too much relic abundance of $S$. However, it is well known that if the mass of the NOP is close to that of the LOP, it would not be decoupled from thermal equilibrium at the freeze-out temperature of the LOP and thus influences the relic abundance of the LOP [17]. We call this coannihilation [17] and this mechanism lowers the present day relic density of $S$ in this scenario. It turns out that in our scenario if $\delta m = m_\Phi - m_S \approx T_f$, processes of annihilation of the NOP into a pair of the SM particles through the s-channel, as shown in figure 3, can significantly affect the relic abundance of $S$ to make it appropriately reduced.

![Figure 2. Cascade decays of the NOP to the LOP: (a) the case for $m_S < m_\Phi$, (b) the case for $m_\Phi < m_S$.](image-url)
With the help of the standard formulae for calculating the relic abundance of s-wave annihilation [16],

$$\Omega_S h^2 = \frac{(1.07 \times 10^9) x_f}{g_s^{1/2} M_{Pl} \int_{x_f}^{\infty} \langle \sigma v \rangle_{rel} x^{-2} dx},$$

we can estimate the present relic density of singlet Majorana neutrinos $S$. Here $g_s$ is the degree of freedom in equilibrium at annihilation, and $x_f = m_{LOP}/T_f$ is the inverse freeze-out temperature in units of $m_{LOP}$, which can be determined by solving the equation

$$x_f \simeq \ln \left( 0.038 g_s (g_s x_i)^{-1/2} M_{Pl} m_{LOP} \langle \sigma v \rangle_{rel} \right),$$

where $\langle \rangle$ means the relevant thermal average, $g_{eff} = \sum_i g_i (m_i/m_{LOP})^{3/2} e^{-(m_i-m_{LOP})/T}$ with the number of degrees of freedom $g_i$ for $i = (S, \Phi)$, and $\langle \sigma v \rangle_{eff}$ is the effective cross section defined in [17]. On calculating $\Omega_S h^2$, we have used the micrOMEGAs 2.0 program [18]. The dominant contribution to $\sigma_{eff}$ in this case is the pair annihilation cross section of the heavier particle $\Phi$. Since it is for $T_f \sim 0.04 m_{\Phi}$ that the annihilation of $\Phi$ is important, what we need to know is the non-relativistic annihilation cross section. In the non-relativistic limit, s-channel annihilation of $\Phi$ via Higgs exchange is given by [10]

$$\sigma_{ann}v_{rel} = \frac{8 \lambda^2 \mu_{EW}^2}{(4 m_{\Phi}^2 - m_h^2)^2 + m_h^2 m_{12}^2} F_X,$$

where $\Gamma_h$ is the total Higgs decay rate, and $F_X = \lim_{m_h \to 2 m_{\Phi}} (\Gamma_h/m_h)$ with the partial rate for decay $h \to X$, for a virtual Higgs $h$. Requiring $\Omega_S h^2$ to be in the region measured from WMAP, $0.094 < \Omega h^2 < 0.128$ [14,15], we can obtain a relation between the coupling $\lambda$ and the mass of the scalar boson $m_{\Phi}$.

For the case of $\Omega_S h^2 = 0.128$, in figure 4 we represent the relation between $\lambda$ and $m_S$ for $\delta m = 5$ GeV. Figure 4(a) corresponds to the Higgs mass $m_h = 120$ GeV, whereas figure 4(b) corresponds to $m_h = 200$ GeV. The shadowed region is forbidden due to the breakdown of perturbation, so the region $57 \leq m_S \leq 74$ (190) GeV for $m_h = 120$ (200) GeV and $\Omega h^2 = 0.128$ is not relevant to our scenario. In the case of $\Omega h^2 = 0.128$, it is $57 \leq m_S \leq 75$ (196) GeV for $m_h = 120$ (200) GeV.

We notice from figure 4 that there exist kinematically special regions, such as the Higgs threshold ($2 m_{\phi} \simeq m_H$) and the two-particle threshold in the final states ($m_{\phi} \simeq m_Z$ or $m_{\phi} \simeq m_H$, and so on). We see from figure 4 that for $50 \leq 74$ GeV $m_S \leq 234$ (245) GeV for $m_h = 120$ (200) GeV and $\Omega h^2 = 0.128$ except for the regions corresponding to the poles and particle thresholds, the abundance constraint that has arisen from WMAP results requires $\lambda \sim O(0.5(0.3) - 1)$. This result indicates that we do not need any fine-tuning or special choice of the parameters in order to achieve the right amount of

\[ \text{Figure 3. Diagram for processes of annihilation of the singlet scalar bosons $\Phi$.} \]
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Figure 4. Relationship between $\lambda$ and $m_S$ that has arisen from the constraints $\Omega_S h^2 = 0.128$ and 0.094 corresponding to the upper and lower limits of $\Omega_{DM} h^2$ measured from WMAP, respectively. Here the mass difference $\delta m = m_\Phi - m_S$ has been taken to be 5 GeV and the Higgs mass $m_h$ to be (a) 120 GeV and (b) 200 GeV. Here the shadowed region is forbidden due to the breakdown of perturbation.

relic abundance of dark matter candidates. It also turns out that if $m_S$ is lighter than 50 (74) GeV or heavier than 234 (245) GeV for $m_h = 120$ (200) GeV and $\Omega h^2 = 0.128$, the relic abundance of $S$ is incompatible with WMAP results for the relic density. It is worth noticing that the coupling $\lambda$ gets significantly suppressed down to the level of $10^{-2}$ near the Higgs pole. This is because the Higgs resonance is quite narrow, which in turn considerably enhances the scalar boson $\Phi$ annihilation rate, especially if $2m_\Phi$ is slightly smaller than $m_h$ [17, 10].

3.2. Singlet scalar boson $\Phi$ as a dark matter candidate

The singlet Higgs scalar $\Phi$ can be a dark matter candidate, provided that $m_\Phi \lesssim m_S$. Since $\Phi$ is the lightest particle involved in the interaction term $Y_S N S \Phi$, this particle cannot decay into other particles, and thus the annihilation processes relevant to a successful
dark matter candidate can occur through the Higgs interaction term $\lambda \Phi^2 H^\dagger H$. In this case, the singlet scalar bosons annihilate into the SM particles mediated by the Higgs boson $h$.

In [8], it has been proposed that a stable, strongly self-coupled scalar field can solve the problems of cold dark matter models for structure formation in the Universe, as regards galactic scales. Thus, the singlet scalar boson in our model can serve as a self-interacting scalar dark matter candidate. If the mass difference between $m_\Phi$ and $m_S$ is very large, the new Yukawa interactions $Y_S N S \Phi$ will affect the relic density of $\Phi$ negligibly, so the behaviour of $\Phi$ as a dark matter candidate is the same as that of the self-interacting scalar dark matter candidate. If the mass difference between $S$ and $\Phi$ is very small, the particles $S_i$ are thermally accessible and they are as abundant as $\Phi$. In this case, the formulae for the relic abundance of the particle $\Phi$ and its freeze-out temperature $x_f$ are given by the same forms of equations (12) and (13), respectively. In fact, however, the annihilation cross sections associated with the heavier particles $S_i$ turn out to be negligibly small in this scenario, so $\sigma_{\text{eff}} \simeq \sigma_\Phi g^2 / g_{\text{eff}}^2$, where $\sigma_\Phi$ is the pair annihilation cross section of $\Phi$ and $g$ is the internal degree of freedom of $\Phi$. Inserting $\sigma_{\text{eff}}$ into $x_f$, we see that the value of $x_f$ for $\delta m = m_S - m_\Phi = 5$ GeV and $m_\Phi = 500$ GeV in this scenario is about 4% lower than that in the singlet Higgs model without the particles $S_i$, which in turn increases the relic abundance of $\Phi$.

Since several important features of our results are quite similar to those of [10], here we just present what the differences between [10] and our model are. In figure 5, we show the relation between $\lambda$ and $m_\Phi$ for $m_h = 120$ (200) GeV and $\delta m = m_S - m_\Phi = 5$ GeV, which is generated by requiring $\Omega h^2 = 0.128$ and 0.094. We see from figure 5 that $m_\Phi$ should be less than 551 GeV (571 GeV) for $m_h = 120$ (200) GeV and $\Omega h^2 = 0.128$ in our model, so as to be consistent with the relic abundance constraint without breaking down perturbation (i.e. $\lambda \lesssim 1$). This upper limit on $m_\Phi$ is much more restrictive than what is obtained in [10]. Also, we see in figure 5(a) that when $m_\Phi$ lies between 80 and 551 GeV, the value of $\lambda$ in our model is much larger than that given in [10].

4. Implication for dark matter searches

For directly detecting dark matter, a typical proposed method is that of detecting the scattering of dark matter off the nucleus of the detecting material. Since the scattering cross section is expected to be very small, the energy deposited by a dark matter candidate on the detector nucleus is also very small. In order to measure this small recoil energy, typically of order keV, of the nucleus, a very low threshold detector condition is required. Since the sensitivity of detectors to a dark matter candidate is controlled by their cross section for elastic scattering off the nucleus, it is instructive to examine how large the size of the elastic cross section could be. First, to estimate the cross section for elastic scattering off the nucleus, we need to know the relevant matrix element for slowly moving spin-$J$ nuclei, which is approximately given by [10]

$$\frac{1}{2J + 1} \sum_{\text{spins}} \left| \langle n' | \sum_f y_f \bar{f} f | n \rangle \right|^2 \simeq \frac{|A_n|^2}{(2\pi)^6},$$

(15)
Figure 5. Relationship between $\lambda$ and $m_\Phi$ corresponding to $\Omega_\Phi h^2 = 0.128$ (the lower solid line) and 0.094 (the upper solid line). The mass difference $\delta m = m_S - m_\Phi$ has been taken to be 5 GeV and $m_h$ to be (a) 120 GeV and (b) 200 GeV. Here the shadowed region is forbidden due to breakdown of perturbation and the dashed line corresponds to the prediction of the model with self-interacting dark matter [10] for $\Omega_\Phi h^2 = 0.128$.

where $n$ denotes nucleons and $|A_n|$ is determined as

$$A_n = g_{Hnn} \approx \frac{190 \text{ MeV}}{v_{\text{EW}}},$$

by following the method given in [10] and taking the strange quark mass to be 95 MeV and $\langle n|s\bar{s}|n\rangle \sim 0.7$.

Now, let us estimate the sizes of the elastic scattering cross sections in each dark matter candidate case.

(i) Case for $m_S < m_\Phi$:

The Feynman diagram describing the scattering of the singlet Majorana neutrino $S$ off nucleons and nuclei is given by t-channel Higgs and heavy Majorana neutrino exchange, as shown in figure 6(a). In this case, the non-relativistic spin-independent quasi-elastic
scattering cross section is approximately given by
\[
\sigma_{el} \approx \frac{Y^2 Y_n^2 |A_n|^2}{128 \pi^3} \left( \frac{m_n^2 m_S^2}{m_h^2 m_N^2 m_S} \right) \times m_{Pl},
\] (17)
where \( m_{Pl} \) is the Planck mass.

However, the size of \( \sigma_{el} \) turns out to be less than \( 10^{-71} \) cm\(^2\) due to the small neutrino mass \( m_\nu \) as well as the small Yukawa couplings \( Y \) and \( Y_\nu \), which are much smaller than the current bound from dark matter experiments. Thus, it is extremely difficult to detect the signal for the singlet Majorana neutrino \( S \) at dark matter detectors.

(ii) Case for \( m_\phi < m_S \):
In this case, the Feynman diagram relevant to scalar–nucleon elastic scattering is presented in figure 6(b); it has already been considered in [10]. Then, the non-relativistic elastic scattering cross section is given by [10]
\[
\sigma_{el} = \frac{\lambda^2 v_{EW}^2 |A_n|^2}{\pi} \left( \frac{m_\phi^2}{m_h^2 m_\phi^2 \lambda} \right),
\] (18)
where \( m_\phi = m_\phi m_n/(m_\phi + m_n) \) is the reduced mass for the collision. Substituting equations (16) into (18),
\[
\sigma_{el}(\text{nucleon}) \approx \frac{1}{\pi} \left( \frac{\Lambda 190 \text{ MeV}}{m_h^2} \right)^2 \left( \frac{m_p}{m_\phi} \right)^2 \left( \frac{m_p}{m_\phi} \right)^2 \left( \frac{\Omega_\Phi h^2}{0.128} \right)
= 2\lambda^2 \left( \frac{100 \text{ GeV}}{m_h} \right)^4 \left( \frac{50 \text{ GeV}}{m_\phi} \right)^2 \left( 4.47 \times 10^{-42} \text{ cm}^2 \right),
\]
where the mass \( m_p \) is the mass of a proton. In figure 7, we plot the predictions for the elastic scattering cross section as a function of the scalar mass \( m_\phi \) for \( m_h = 120 \) and 200 GeV, and the mass difference (\( \delta m = m_S - m_\phi \)) is taken to be 5 GeV. The lower line corresponds to \( \Omega_\Phi h^2 = 0.128 \), whereas the upper line corresponds to \( \Omega_\Phi h^2 = 0.094 \). In calculating \( \sigma_{el} \), we used the relationship between \( \lambda \) and \( m_\phi \) which is obtained through the constraint from WMAP results as before.

In figure 7, we plot the new 90% C.L. upper bound for the WIMP–nucleon spin-independent cross section as a function of \( m_\phi \) obtained from the XENON10 dark matter experiment [20]. As one can see from figure 7, when \( \Omega h^2 = 0.128 \), the region \( m_\phi < 48 \) (71) GeV for \( m_h = 120 \) (200) GeV and \( \delta m = 5 \) GeV is excluded by the XENON10 dark matter experiment [20]. In particular, the elastic cross section for \( m_\phi \sim 75 \) GeV in the case of \( m_h = 120 \) GeV nearly reaches the current upper bound from XENON10 [20].
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Figure 7. Plots of the elastic cross section $\sigma_{el}$ as a function of $m_\Phi$ for (a) $m_h = 120$ GeV and (b) for $m_h = 200$ GeV. The dotted line shows the spin-independent WIMP–nucleon cross section upper limits (90% C.L.) from the XENON10 dark matter experiment [20]. Here the dashed line corresponds to the prediction of the model with self-interacting dark matter [10] for $\Omega h^2 = 0.128$.

Comparing our results with those from the model of self-interacting scalar dark matter [10], we see that our prediction for the elastic cross section in the region $80 \ (110) \ GeV < m_\Phi < 551 \ (571) \ GeV$ for $m_h = 120 \ (200) \ GeV$, and $\delta m = 5 \ GeV$ is about 1.4 \ (1.5) to 4.3 \ (4.4) times larger than that estimated in [10]. This indicates that our scenario for scalar dark matter is distinguishable from the original model of the self-interacting scalar dark matter.

5. Implication for Higgs searches at LHC

Now we investigate the implications of our scenarios for Higgs searches in collider experiments. The singlet scalar boson will not directly couple to ordinary matter, only to the Higgs fields. Although the presence of the singlet scalars will not affect electroweak phenomenology in a significant way, it will affect the phenomenology of the Higgs boson. Due to the large values for the coupling $\lambda \sim O(0.1–1.0)$ required by relic abundance
constraints, real or virtual Higgs production may be associated with the singlet Higgs $\Phi$ production, as discussed in [10]. We see from equation (11) that if $2m_\Phi < m_h$, the real Higgs boson can decay into a pair of singlet scalars, whereas if $2m_\Phi > m_h$, the singlet scalar bosons cannot be produced by real Higgs decays, but will arise only via virtual Higgs exchange. We know that any singlet scalar bosons produced are not expected to interact inside the collider; thus they only give rise to strong missing energy signals.

(i) Case for $2m_\Phi < m_h$:

In this case, the Higgs boson can decay invisibly into a pair of the singlet scalar bosons. The invisible decay width is given at tree level by

$$\Gamma_{H \rightarrow \Phi\Phi} = \frac{\lambda^2 v_{\text{EW}}^2}{32\pi m_h} \sqrt{1 - \frac{4m_\Phi^2}{m_h^2}}. \quad (19)$$

Since the relic abundance constraints require a large value of the coupling $\lambda$, the width of invisible decay to the Higgs boson gets large. It is known that if the Higgs mass is less than $2M_W$, and hence the Higgs partial width in the SM particles is very small, the Higgs will decay predominantly into singlet scalar bosons. Then, LHC may yield a discovery signal for an invisible Higgs with enough reachable luminosity, for instance, $10 \text{ fb}^{-1}$ of integrated luminosity for $m_h = 120 \text{ GeV}$, in associated production with $Z$ bosons [21]. To quantify the signals for the invisible decay of the Higgs boson, we investigate the ratio $R$ defined as follows [10]:

$$R = \frac{Br_{h \rightarrow b\bar{b},c\bar{c},\tau\tau}(\text{SM + }\Phi)}{Br_{h \rightarrow b\bar{b},c\bar{c},\tau\tau}(\text{SM})} = \frac{\Gamma_{h,\text{total}}(\text{SM})}{\Gamma_{h \rightarrow \Phi\Phi} + \Gamma_{h,\text{total}}(\text{SM})}.$$

The ratio $R$ indicates how the expected signal for the visible decay of the Higgs boson can decrease due to the existence of the singlet scalar bosons.

(a) Case for singlet neutrino dark matter. In figure 8, we plot the value of $R$ as a function of $m_\Phi$ for (a) $m_h = 120 \text{ GeV}$ and (b) $m_h = 200 \text{ GeV}$. Here, we fixed the value of $\delta m$ as 5 GeV, and used the relation between $\lambda$ and $m_\Phi$ which is obtained from the relic abundance constraints as explained before. In calculating the rate of decay of Higgs particles, we have used CalcHEP 2.4.5 [19]. As can be seen in figure 8, the prediction for the value of $R$ is totally different from that in [10]. But, in this case, we can probe the invisible Higgs decay by using $R$ only for the very narrow region $55 \text{ (79) GeV} < m_\Phi < 60 \text{ (100) GeV}$ when $m_h = 120 \text{ (200) GeV}$. This is because the lower limit of $m_S$ has been determined as 50 (74) GeV for $m_h = 120 \text{ (200) GeV}$, as shown in figure 4, and the upper limit of $m_\Phi$ is constrained by the condition $2m_\Phi < m_h$. From figure 8, like in [10], it turns out that the invisible Higgs decay width dominates the total width everywhere except in the vicinity $2m_\Phi = m_H$. This means that a tremendous suppression of the observable Higgs signal may happen at the LHC.

(b) Case for singlet scalar dark matter. For this case, we plot the prediction for $R$ in figure 9 and the result turns out to be almost the same as that in [10]. Like in the above case, the invisible width also dominates the total width everywhere except in the region $2m_\Phi \simeq m_h$ [10]. We notice that when $m_h = 120 \text{ (200) GeV}$ and $\Omega h^2 = 0.128$ the parameter spaces below $m_\Phi = 48 \text{ (71) GeV}$ are excluded by the current bound obtained from the XENON10 dark matter experiment. In this case, if we impose the current bound from the XENON10 dark matter experiment [20], the mass of the scalar boson $m_\Phi$ should be larger than 46 (64) GeV for $m_h = 120 \text{ (200) GeV}$ and $\Omega h^2 = 0.128$. Then, the values of
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Figure 8. Plots of the ratio $R$ as a function of $m_\Phi$ for our model with $\Omega_\Phi h^2 = 0.128$ and 0.094 (solid lines), when (a) $m_h = 120$ GeV and (b) 200 GeV, and for the model with self-interacting dark matter [10] with $\Omega_\Phi h^2 = 0.128$ (dashed line). Here the shadowed region represents the region forbidden by figure 4, which applies only to our model.

$R$ in our model are constrained to be 0.26–1 (0.81–1) where the upper limit is determined by the condition $2m_\Phi < m_h$, as it should be.

(ii) Case for $2m_\Phi > m_h$:

In this case, the singlet scalar bosons can be produced only through virtual Higgs exchange. Like in the previous case, the singlet scalar particles produced can be detected as missing energy above an energy threshold, $E \geq 2m_\Phi$. In this case, LHC is an unlikely place for discovery of a missing energy signal, whereas a future linear collider might be a good place for detecting such a signal.

6. Conclusion

We have considered a variant of the seesaw mechanism by introducing extra singlet neutrinos and the singlet scalar boson and showed how low scale leptogenesis is realized in our scenario. We have examined whether the newly introduced neutral particles, either
Figure 9. Plots of the ratio $R$ as a function of $m_\Phi$ for our model with $\Omega_\Phi h^2 = 0.128$ and 0.094 (solid lines), when (a) $m_h = 120$ GeV and (b) 200 GeV, and for the model with self-interacting dark matter [10] with $\Omega_\Phi h^2 = 0.128$ (dashed line). Here the shadowed region represents the region forbidden by the XENON10 dark matter experiment [20], which applies only to our model.

Singlet Majorana neutrinos or the singlet scalar boson, can be dark matter candidates. We have shown that the process of coannihilation of dark matter candidates and the next lightest $Z_2$ odd particles plays a crucial role in generating the right amount of the relic density of dark matter candidates. We have also discussed the implications of dark matter detection through scattering off the nucleus of the detecting material for our scenarios for dark matter candidates. From the recent result from the XENON10 dark matter experiment, we were able to get some constraints on the mass of the singlet scalar boson. In addition, we have studied the implications for the search for invisible Higgs decay at LHC which may serve as a probe for our scenarios for dark matter.

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