Baryogenesis via leptogenesis from asymmetric dark matter and radiatively generated neutrino mass

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We propose an extension of the standard model (SM) by including a dark sector comprised of three generations of heavy right-handed neutrinos, a singlet scalar, and a singlet Dirac fermion, where the latter two particles are stable and are viable candidates of dark matter (DM). In the early Universe, the *CP*-violating out-of-equilibrium decay of heavy right-handed neutrinos to a singlet Dirac fermion and scalar in the dark sector generates a net DM asymmetry. The latter is then transported to the visible sector via a dimension-eight operator which conserves $B - L$ symmetry and is in thermal equilibrium above the sphaleron decoupling temperature. An additional light singlet scalar is introduced which mixes with the SM Higgs and paves a path for annihilating the symmetric components of the DM candidates. We discuss the constraints on singlet-doublet Higgs mixing from invisible Higgs decay, signal strength at the LHC, and the direct search of DM at terrestrial laboratories. At tree level, the neutrinos are shown to be massless since the symmetry of the dark sector forbids the interaction of right-handed neutrinos with SM particles. However, at the one-loop level, the neutrinos acquire sub-eV masses as required by the oscillation experiments.

I. INTRODUCTION

The evidence from the galaxy rotation curve, gravitational lensing, and large scale structure of the Universe irrefutably proves the existence of dark matter (DM) in a large scale ($\gtrsim$ a few kpc) [1]. However, the microscopic picture of DM is hitherto not known. The only piece of information that we know about the DM is its relic abundance which is precisely measured by the satellite borne experiments WMAP [2] and PLANCK [3] to be $\Omega_{\rm DM} h^2 = 0.1199 \pm 0.0027$. However, a little is known about the underlying mechanism of generating the relic abundance of DM. The most considered scenario is that DM is a weakly interacting massive particle (WIMP) [4]. The latter gets thermalized in the early Universe due to its weak interaction property. As the temperature falls below its mass scale, the DM gets decoupled from the thermal bath and its density in a comoving volume remains constant

and is what we measure today. This is usually referred to as the WIMP miracle.

A curious observation about DM is that its relic density is about 5 times larger than the baryon density of the present Universe, i.e., $\Omega_{\rm DM} \approx 5\Omega_B$. This implies that the relic density of DM can be generated in a similar way that the baryon asymmetry of the Universe has been generated. See, e.g., [5–9]. The observed baryon asymmetry, usually reported in terms of the baryon to photon ratio, $\eta = n_B/n_\gamma$, is given as [10],

$$5.8 \times 10^{-10} \leq \eta \leq 6.6 \times 10^{-10} \quad (BBN) \quad (95\%CL), \quad (1)$$

where $\eta = 7.04 Y_B$ with $Y_B \equiv n_B/s$. Similarly, the observed DM abundance can be expressed as

$$Y_{\rm DM} \equiv \frac{n_{\rm DM}}{s} = 4 \times 10^{-10} \left(\frac{1 \text{ GeV}}{M_{\rm DM}}\right) \left(\frac{\Omega_{\rm DM} h^2}{0.11}\right). \quad (2)$$

This implies that $Y_{\rm DM}/Y_B \approx \mathcal{O}(1)$ if $M_{\rm DM} \sim 5$ GeV. However, it can vary from a GeV to TeV depending on the magnitude of *CP* violation in the visible and dark sectors. See, for instance, [8].

The standard model (SM), which is based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$, is a successful theory of fundamental particles and their interactions. However, it does not explain either the DM abundance or baryon

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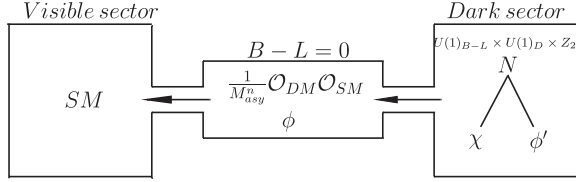‡nsahu@iith.ac.in
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FIG. 1. Pictorial presentation of a dark sector being in thermal contact with the visible sector via Higgs portal coupling as well as higher-dimension operators, which conserve $B - L$ symmetry and are in thermal equilibrium above sphaleron decoupling temperature.

asymmetry of the Universe. Moreover, it cannot explain the nonzero masses of active neutrinos. In this paper, we make an attempt to solve these problems simultaneously in a beyond SM framework. We extend the SM by including a dark sector, as shown in Fig. 1, comprised of three generations of heavy right-handed neutrinos, a singlet scalar ϕ', and a singlet Dirac fermion χ. These particles are charged under an additional symmetry, $U(1)_{B-L} \times U(1)_D \times Z_2$, while remaining inert with respect to the SM gauge group. The $U(1)_{B-L}$ is a gauge symmetry, which is broken spontaneously by the vacuum expectation value of an additional scalar ϕ_{B-L} at a high scale, say around 10^{10} GeV, and gives Majorana masses to right-handed neutrinos, while $U(1)_D$ is a global symmetry and is allowed to break softly because of the higher-dimension operators. Moreover, the $U(1)_D$ symmetry provides a distinction between the dark sector fermions N_R and χ, which have the same charge under $U(1)_{B-L} \times Z_2$ symmetry. Due to $U(1)_D$ symmetry, both the singlet fermion (χ) and scalar (ϕ') are stabilized and become viable candidates of dark matter.

The Majorana mass of heavy right-handed neutrinos breaks $B - L$ symmetry by two units. Therefore, the CP-violating out-of-equilibrium decay of heavy right-handed neutrinos to $\chi\phi'$ in the early Universe generates a net $B - L$ asymmetry [11,12]. The latter is then transferred to the visible sector by a dimension-eight operator [13–15], $\mathcal{O}_8 = \frac{1}{M_{\mathrm{asy}}^4} \bar{\chi}^2 (LH)^2$, which is in thermal equilibrium above the sphaleron decoupling temperature. Note that the operator $\mathcal{O}_8$ breaks $U(1)_D$ symmetry softly, while it conserves $B - L$ symmetry. As a result, the $B - L$ asymmetry produced by the decay of right-handed neutrinos will be distributed between the dark and visible sectors. When the DM χ decouples from the thermal bath, the asymmetry in the two sectors gets segregated. Thus, we get a net $B - L$ asymmetry in the visible sector proportional to the $B - L$ asymmetry in the dark sector. The $B - L$ asymmetry in the visible sector gets transferred to a net baryon (B) asymmetry via the sphaleron transitions, while the $B - L$ asymmetry in χ remains intact. The asymmetry in χ and ϕ' combined gives rise the present day relic density of DM. An additional singlet scalar ϕ is introduced which mixes with the SM Higgs H and paves a path for annihilating the

symmetric components of χ and ϕ'. The abundance of the singlet scalar ϕ will not be present in the current Universe due to its decay to standard model particles through Higgs mixing.

Note that N_R is odd under the Z_2 symmetry. As a result, it does not have a tree-level coupling with left-handed lepton doublets as in the type-I seesaw model [16]. However, the dimension-eight operator $\mathcal{O}_\nu = \frac{1}{\Lambda^4} (\overline{N_R} LH)^2$ is allowed, where Λ is the scale of symmetry breaking. As we discuss in Sec. II C, this generates a Majorana mass of the light neutrinos at the one-loop level. Notice that the operator $\mathcal{O}_\nu$ also breaks the $U(1)_D$ symmetry softly.

The paper is organized as follows. In Sec. II, we introduce the model. Section II C explains the neutrino masses. The generation of DM asymmetry is explained in Sec. III. The transfer of DM asymmetry to the visible sector is discussed in Sec. IV. In Sec. V, we describe the condition for annihilation of symmetric components of the DM. In Sec. VI, we demonstrate the constraints on model parameters from invisible Higgs decay, signal strength of a SM-like Higgs, the requirement of correct relic abundance of DM, and its direct detection. We conclude in Sec. VII.

II. THE MODEL

The model under consideration is based on the symmetry SM $\times U(1)_{B-L} \times U(1)_D \times Z_2$, where $U(1)_{B-L}$ is a local gauge symmetry and is broken spontaneously at a high scale by the vacuum expectation value of a singlet scalar ϕ_{B-L}, whereas $U(1)_D$ is a global symmetry and is allowed to break softly due to higher-dimension operators, as we discuss below. In addition to that, we extend the SM particle content by introducing a dark sector comprised of three generations of heavy right-handed neutrinos N_{iR}, $i = 1, 2, 3$, a Dirac fermion χ, and a singlet scalar ϕ'. An additional singlet scalar ϕ is also introduced, which mixes with the SM Higgs H. The particle content of the model, along with the quantum numbers, is given in Table I. Under the discrete symmetry Z_2, which remains unbroken, both N_R and χ particles are odd. As a result, the lightest Z_2 odd particle χ is stable and is a viable candidate of DM. In addition to that, we assume $\langle \phi' \rangle = 0$. This implies ϕ' is also stable due to $U(1)_D$ symmetry. As a result, the relics of χ and ϕ' constitute the DM content of the present Universe.

TABLE I. Particles of the dark sector and their quantum numbers under the imposed symmetry.

Fields	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$	$U(1)_D$	Z_2
N_R	1	1	0	−1	1	−
χ	1	1	0	−1	1/3	−
ϕ	1	1	0	0	0	+
ϕ'	1	1	0	0	2/3	+
ϕ_{B-L}	1	1	0	+2	−2	+

The corresponding Lagrangian can be given as

$$
\begin{aligned}
\mathcal{L} \supset\, & \overline{N_{Rj}} i\gamma^\mu D_\mu N_{Rj} + \bar{\chi} i\gamma^\mu D_\mu \chi + \frac{1}{2}(\partial_\mu \phi)(\partial^\mu \phi) \\
& + (\partial_\mu \phi')^\dagger (\partial^\mu \phi') + (D_\mu \phi_{B-L})^\dagger (D^\mu \phi_{B-L}) \\
& + M_\chi \bar{\chi}\chi + \lambda_{B-L}\phi_{B-L}\overline{(N_{Ri})^c}N_{Rj} + \lambda_{\mathrm{DM}}\bar{\chi}\chi\phi \\
& + y_i \overline{N_{Ri}}\chi\phi' + \mathrm{H.c.} - V(H,\phi,\phi'),
\end{aligned}
\tag{3}
$$

where

$$
D_\mu = \partial_\mu + i g_{B-L} Y_{B-L} (Z_{B-L})_\mu
$$

and

$$
\begin{aligned}
V(H,\phi,\phi') =\, & -\mu_H^2 H^\dagger H + \lambda_H (H^\dagger H)^2 + \frac{1}{2}M_\phi^2 \phi^2 \\
& + \frac{1}{4}\lambda_\phi \phi^4 + M_{\phi'}^2 \phi'^\dagger \phi' + \lambda_{\phi'}(\phi'^\dagger \phi')^2 \\
& + \frac{1}{2}\lambda_{H\phi}(H^\dagger H)\phi^2 + \mu_\phi \phi(H^\dagger H) + \mu'_\phi \phi(\phi'^\dagger \phi') \\
& + \lambda_{H\phi'}(H^\dagger H)(\phi'^\dagger \phi') + \frac{\lambda_{\phi\phi'}}{2}\phi^2(\phi'^\dagger \phi').
\end{aligned}
\tag{4}
$$

In Eq. (4), we assume that $U(1)_{B-L}$ gauge symmetry is broken spontaneously by the vev of ϕ_{B-L} at a high scale, $\langle\phi_{B-L}\rangle = v_{B-L} \sim 10^{10}$ GeV (say). Therefore, ϕ_{B-L} does not play any role in the low-energy electroweak phenomenology. However, the vev of ϕ_{B-L} gives super heavy masses to right-handed neutrinos as well as neutral gauge boson Z_{B-L}. The $B-L$ quantum numbers of N_R and χ are the same and are taken to be -1. However, they are distinguishable by their $U(1)_D$ quantum numbers.

Since $B-L$ charges of all the SM fermions are known, it is straightforward to see that the uplifting of global $U(1)_{B-L}$ symmetry of the SM to a gauge one brings in $B-L$ anomalies. In particular, the nontrivial one per family is given by [17]

$$
\begin{aligned}
U(1)_{B-L}^3:\, & 3\left[2\times\left(\frac{1}{3}\right)^3 - \left(\frac{1}{3}\right)^3 - \left(\frac{1}{3}\right)^3\right] \\
& + [2\times(-1)^3 - (-1)^3] = -1,
\end{aligned}
$$

where the number 3 in front is the color factor. This can be exactly canceled by introducing one right-handed neutrino per family as we did in this model. Thus, the model is anomaly free. Since χ is a vectorlike fermion, it does not introduce any additional anomaly though it is charged under $U(1)_{B-L}$.

As discussed above, the mass of heavy right-handed neutrinos is $M_N \gg M_W$, while the mass of χ is $M_\chi < M_W$. The neutral gauge boson corresponding to $B-L$ symmetry acquires a large mass $M_{Z_{B-L}} \gg M_Z$. In the following, we discuss the $\phi - H$ mixing on which the annihilation of the symmetric component of DM depends.

A. Case-I

The electroweak phase transition occurs as the SM Higgs acquires a vacuum expectation value (vev) $v = \langle H \rangle$. This induces a nonzero vev to ϕ due to the trilinear term $\mu_\phi \phi(H^\dagger H)$ as given in Eq. (4). We assume that $\langle\phi\rangle = u \ll v$. Then the quantum fluctuations around the minimum can be given as

$$
H = \begin{pmatrix} 0 \\ \frac{v+h}{\sqrt{2}} \end{pmatrix} \quad \text{and} \quad \phi = u + \tilde{\phi}.
\tag{5}
$$

By minimizing the scalar potential (4), we get the vacuum expectation values,

$$
u = \frac{-\mu_\phi v^2}{M_\phi^2 + \lambda_{H\phi}v^2},
\tag{6}
$$

and

$$
v = \sqrt{\frac{\mu_H^2 - \frac{1}{2}\lambda_{H\phi}u^2 - \mu_\phi u}{2\lambda_H}}.
\tag{7}
$$

Notice that $\mu_\phi \propto u$. In the limit $u \to 0$, we recover the vev of the SM Higgs,

$$
v = \sqrt{\frac{\mu_H^2}{2\lambda_H}}.
\tag{8}
$$

The small vev u does not affect our discussions in the following sections and, hence, we set it to zero from here on.

B. Case-II

Here, we relax the ϕ vev to zero, i.e., $\langle\phi\rangle = 0$. Then the quantum fluctuation around the minimum is given by

$$
H = \begin{pmatrix} 0 \\ \frac{v+h}{\sqrt{2}} \end{pmatrix} \quad \text{and} \quad \phi = \tilde{\phi}.
\tag{9}
$$

As a result, after the electroweak phase transition, the two scalars h and $\tilde{\phi}$ mix with each other. The mass matrix is given by

$$
\begin{pmatrix} 2\lambda_H v^2 & \frac{\mu_\phi v}{\sqrt{2}} \\ \frac{\mu_\phi v}{\sqrt{2}} & M_\phi^2 + \frac{\lambda_{H\phi}}{2}v^2 \end{pmatrix}.
\tag{10}
$$

Diagonalizing the above mass matrix, we get the masses M_{h_1} and M_{h_2} corresponding to the physical Higgses h_1 and h_2:

$$h_1 = h \cos\gamma + \tilde{\phi} \sin\gamma$$
$$h_2 = -h \sin\gamma + \tilde{\phi} \cos\gamma. \qquad (11)$$

The mixing angle γ can be quantified as

$$\sin\gamma \approx \frac{\sqrt{2}\mu_\phi v}{2\lambda_H v^2 - M_\phi^2 - \frac{\lambda_{H\phi} v^2}{2}}. \qquad (12)$$

We identify h_1 to be the SM-like Higgs with mass $M_{h_1} = 125.18$ GeV, while h_2 is the second Higgs whose mass is going to be determined from the relic abundance requirement. In fact, in Sec. VI, we obtain the light scalar mass, from the requirement of depletion of the symmetric component of the DM, to be $M_{h_2} \approx 2M_\phi' \approx 2M_\chi \approx 2.32826$ GeV. In Fig. 2, we show the contours of $M_{h_1} = 125.18$ GeV (dashed lines), $M_{h_2} = 2.32826$ GeV (solid lines), and $\sin\gamma = 0.14, 0.9$ (dot-dashed lines) in the plane of λ_H versus μ_ϕ for $\lambda_{H\phi} = 0.1$ (meeting at point A), 0.01 (meeting at point B). We see that the large range of mixing is allowed to explain simultaneously the masses of h_1, h_2. Later, we will see that the large mixing angles are strongly constrained by other phenomenological requirements. When the mixing goes to zero (i.e., $\mu_\phi \to 0$, which implies $\sin\gamma \to 0$), we recover the SM Higgs mass $M_{h_1} = 2\lambda_H v^2 = 125.18$ GeV for $\lambda_H = 0.13$. As the mixing angle increases (i.e., $\mu_\phi \neq 0$), we still satisfy the required masses of h_1 and h_2 with small λ_H. In what follows, we will take $\sin\gamma$ as the measure of mixing.

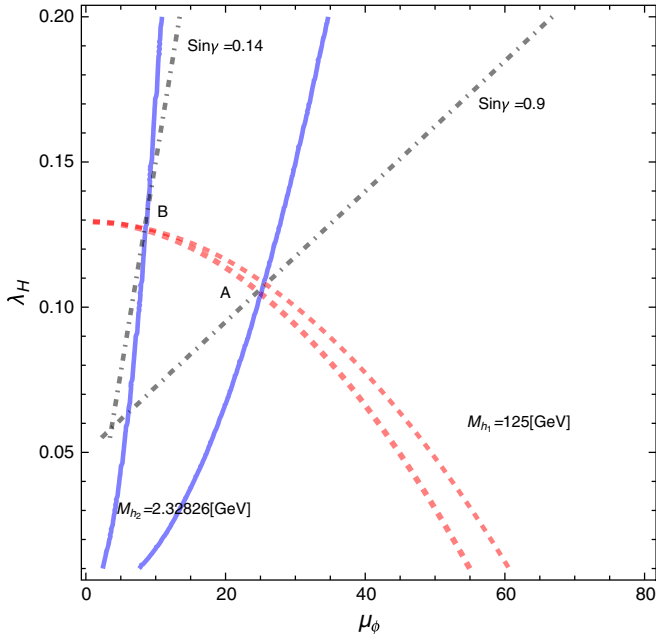

FIG. 2. Contours of $M_{h_1} = 125.18$ GeV (dashed lines), $M_{h_2} = 2.32826$ GeV (solid lines), $\sin\gamma = 0.14, 0.9$ (dot dashed lines) in the plane of λ_H versus μ_ϕ. We set $M_\phi = 0$.
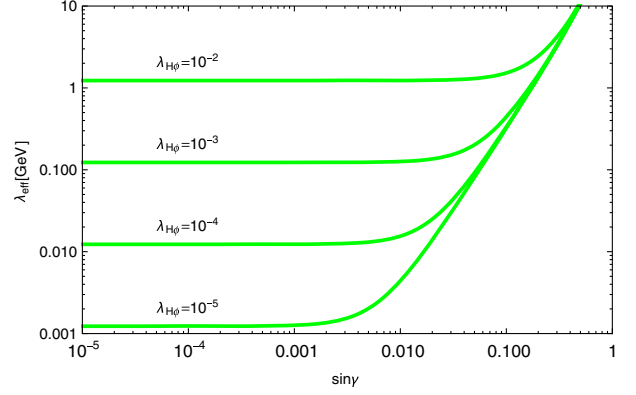

FIG. 3. Effective coupling of h_1 to $h_2 h_2$ as a function of $\sin\gamma$ for $\lambda_H = 0.13$.

The effective coupling of $h_1 h_2 h_2$ from Eq. (4) can be given as

$$\lambda_{\text{eff}} = 3\lambda_H v \cos\gamma \sin^2\gamma + \frac{\lambda_{H\phi}}{2} v \cos^3\gamma + \frac{\mu_\phi}{2} \sin^3\gamma$$
$$- \mu_\phi \sin\gamma \cos^2\gamma - \lambda_{H\phi} v \sin^2\gamma \cos\gamma. \qquad (13)$$

In Fig. 3, we have shown the effective coupling of SM-like Higgs to $h_2 h_2$ as a function of $\sin\gamma$ for various values of $\lambda_{H\phi}$. We see that λ_{eff} is almost independent of $\lambda_{H\phi}$ for $\sin\gamma \sim 0.1$. We will come back to this issue while calculating the invisible decay width of SM-like Higgs in Sec. VI.

C. Neutrino masses

The lepton number is violated by the Majorana mass term of the heavy right-handed neutrinos. Note that the term $\overline{N_R} \tilde{H}^\dagger L$ is not allowed, as N_R is odd under the Z_2 symmetry. However, the dimension-eight operator $\mathcal{O}_\nu = \frac{(\overline{N_R} \tilde{H}^\dagger L)^2}{\Lambda^4}$ is allowed, where Λ is the scale of symmetry breaking. The relevant diagram generating neutrino masses radiatively is shown in Fig. 4.[1]

By taking Λ as the cutoff scale, the neutrino mass can be calculated from Fig. 4 as

$$M_\nu = \frac{1}{8\pi^2} \frac{v^2 M_N}{\Lambda^2} \left[1 + \frac{M_N^2}{\Lambda^2} Log\left(\frac{M_N^2}{M_N^2 + \Lambda^2} \right) \right], \qquad (14)$$

where the v is the vacuum expectation value of the SM Higgs and M_N is the mass scale of the heavy right-handed neutrino. Inverting the above formula, we get the symmetry breaking scale:

$$\Lambda \approx 7.66 \times 10^{11} \text{ GeV} \left(\frac{0.1 \text{ eV}}{M_\nu} \right) \left(\frac{M_N/\Lambda}{0.1} \right). \qquad (15)$$
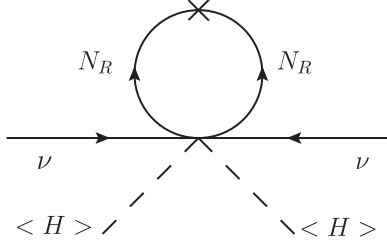
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[1]See for a recent review [18].

FIG. 4. Radiative neutrino mass at one-loop level, generated by the operator $\mathcal{O}_\nu$.

In Sec. III, we take the Majorana mass of heavy right-handed neutrinos to be $M_N \approx 10^{10}$ GeV.

III. GENERATION OF ASYMMETRY IN DARK MATTER SECTOR

In the early Universe, the right-handed neutrinos at a temperature above their mass scales are assumed to be in thermal equilibrium. As the Universe expands, the temperature falls. As a result, the right-handed neutrinos, below their mass scales, go out of equilibrium and decay through the process: $y_i \bar{N}_{Ri} \chi \phi' + \text{H.c.}$. Without loss of generality we choose the mass basis of right-handed neutrinos to be diagonal. In this basis, the heavy Majorana neutrinos are defined by $N_i = \frac{1}{\sqrt{2}}[N_{iR} + (N_{iR})^c]$ and hence their decay violate $B - L$ by two units. In the mass basis of N_1, the decay rate of N_1 is given by: $\Gamma_1 = \frac{(y^\dagger y)_{11}}{16\pi} M_1$. Comparing it with the Hubble expansion parameter $H = 1.67 g_*^{1/2} T^2 / M_{\text{Pl}}$ at $T \sim M_1$, we get the out-of-equilibrium condition: $y \lesssim \mathcal{O}(10^{-3})\sqrt{M_1/10^{10} \text{ GeV}}$. Thus depending on the mass of right-handed neutrinos the decoupling epoch can be different. We assume a normal hierarchy among the heavy Majorana neutrinos. As a result, the CP-violating decay of lightest heavy neutrino (N_1) to ϕ' and χ, generates a net asymmetry in χ and ϕ'. Since both χ and ϕ' are stable, the asymmetry in χ and ϕ' together represents the DM abundance.

The CP asymmetry in the decay of N_1 arises via the interference of a tree-level diagram with one-loop self energy and vertex diagrams as shown in Fig. 5. The asymmetry ϵ_χ is estimated to be [19]

$$\epsilon_\chi = \frac{\Gamma(N_1 \to \chi\phi') - \Gamma(N_1 \to \bar{\chi}\phi')}{\Gamma_{N_1}}$$
$$\simeq -\frac{1}{8\pi}\left(\frac{M_1}{M_2}\right)\frac{\text{Im}[(y^\dagger y)^2]_{12}}{(y^\dagger y)_{11}}. \tag{16}$$

where we assume $M_1 \ll M_2 \ll M_3$, and M_i, $i = 1, 2, 3$ are the masses of heavy right-handed neutrinos. Thus, below the mass scale of N_1, we get a net $B - L$ asymmetry: [20,21]
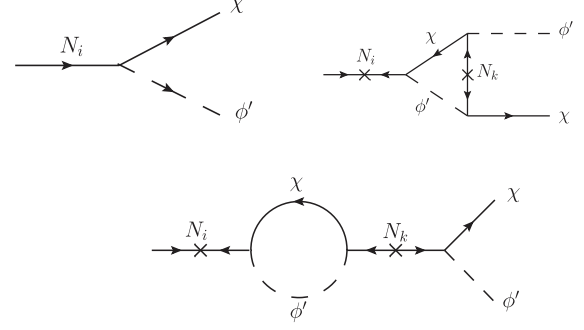

FIG. 5. CP-violation arising from the interference of tree-level diagram with vertex and self energy correction diagrams in the decay of N_1.

$$(n_{B-L})_{\text{total}} = \epsilon_\chi \kappa s \times \frac{n_{N_1}^{eq}(T \to \infty)}{s} \tag{17}$$

where $(n_{N_1}^{eq}/s)(T \to \infty) = 135\zeta(3)/(4\pi^4 g_*)$ is the relativistic equilibrium abundance of N_1. κ is the washout factor, arises via inverse decay and scattering processes and $s = (2\pi^2/45)g_* T^3$ is the entropy density. Depending on the strength of Yukawa coupling, the value of κ can vary between 0 to 1. However, for definiteness we choose $\kappa = 0.01$. The details of $B - L$ asymmetry generated in the dark sector can be obtained by solving the required Boltzmann equations [20], which is beyond the scope of this paper. The generated $B - L$ asymmetry will be distributed between visible and dark sectors via a higher-dimension operator as we introduce in Sec. IV.

IV. ASYMMETRY TRANSFER FROM DARK SECTOR TO VISIBLE SECTOR

The asymmetry generated via the decay of lightest heavy Majorana neutrino N_1 can be transferred to the visible sector by a higher-dimension operator [14]:

$$\mathcal{O}_8 = \frac{1}{M_{\text{asy}}^4} \bar{\chi}^2 (LH)^2. \tag{18}$$

Depending on the value of M_{asy}, the transfer operator will decouple from thermal plasma at different temperatures. We can find the decoupling temperature by comparing the interaction rate of the transfer operator with the Hubble expansion rate of the universe at the decoupling epoch T_D. For the operator (18), the rate of interaction between visible and dark sector at the decoupling epoch T_D is given as

$$\Gamma_D \simeq \left(\frac{T_D^4}{M_{\text{asy}}^4}\right)^2 T_D, \tag{19}$$

where M_{Pl} is the Planck mass. By comparing the above interaction rate with the Hubble expansion parameter $H = 1.67 g_*^{1/2} T_D^2 / M_{\text{Pl}}$ we get

$$M_{\mathrm{asy}}^8 > M_{\mathrm{Pl}} T_D^7. \qquad (20)$$

We assume that $T_D \gtrsim T_{\mathrm{sph}}$, where T_{sph} is the sphaleron decoupling temperature. For Higgs mass $M_{h_1} = 125.18 \, \mathrm{GeV}$, the sphaleron decoupling temperature is $T_{\mathrm{sph}} \geq M_W$. As a result, from Eq. (20), we get the constraint on M_{asy} to be $M_{\mathrm{asy}} = 0.9 \times 10^4 \, \mathrm{GeV} = M_{\mathrm{asy}}^*$ for $T_D = M_W$. In other words, Eq. (20) indicates that, if $M_{\mathrm{asy}} > M_{\mathrm{asy}}^*$, then the interaction rate of transfer operator will be in thermal equilibrium for $T > T_D$. The same condition also implies that the processes allowed by the operator will remain out-of-equilibrium below electroweak phase transition. Notice that the estimation of Eq. (20) holds only for the case where χ mass is much smaller than T_D. However, if one were to study heavier χ, chemical decoupling can take place when the number density of χ becomes Boltzmann suppressed. See for instance [22].

The asymmetry in the equilibrium number densities of particle n_i and antiparticle $\bar{n}_i$ can be given as

$$n_i - \bar{n}_i = \frac{g_i}{2\pi^2} \int_0^\infty dq q^2 \left[\frac{1}{e^{\frac{E_i(q) - \mu_i}{T}} \pm 1} - \frac{1}{e^{\frac{E_i(q) + \mu_i}{T}} \pm 1} \right] \qquad (21)$$

where the g_i is the internal degrees of freedom of the particle species i. In the above equation, E_i and q_i are the energy and momentum of the corresponding particle species i. In the approximation of a weakly interacting plasma, where $\beta\mu_i \ll 1$, $\beta \equiv 1/T$, we get [4]

$$n_i - \bar{n}_i \sim \frac{g_i T^3}{6} \times [2\beta\mu_i + \mathcal{O}((\beta\mu_i)^3) \quad \text{bosons}$$

$$\sim \frac{g_i T^3}{6} \times [\beta\mu_i + \mathcal{O}((\beta\mu_i)^3) \quad \text{fermions.} \qquad (22)$$

By comparing Eq. (22) with Eq. (17), we see that $\beta\mu \sim k\epsilon_\chi \ll 1$. This justifies the weak interaction of thermal plasma. We will comeback to this issue at the end of this section.

Now we will estimate the B asymmetry in the visible sector at a temperature above the sphaleron decoupling temperature. To find that we will use the chemical equilibration [23] between different fermions until sphaleron decoupling temperature as discussed below. All the left-handed charged lepton e_{iL}, $\forall \, i$, right-handed charged lepton e_{iR}, $\forall \, i$, left-handed neutrino ν_{iL}, $\forall \, i$, left-handed up-type quark u_{iL}, $\forall \, i$, right-handed up-type quark u_{iR}, $\forall \, i$, left-handed down-type quark d_{iL}, $\forall \, i$, right-handed down-type quark d_{iR}, $\forall \, i$, $W^\pm$, Z-boson, photon (γ), Higgs boson(h) are in thermal equilibrium until sphaleron decoupling temperature. Here the index $i = 1$, 2, 3 is written for three generations. All three generations up-type quark have the same chemical potential, all the three generations down-type quark have the same chemical potential. Similarly, all three left-handed neutrinos have the

same chemical potential. But the three different charge leptons may have different chemical potential. So we omit index i from chemical potential of quarks and neutrinos. The chemical potential of physical Higgs boson, Z boson and photon are set to zero.

Below the electroweak phase transition, the Yukawa interactions can be given as

$$\mathcal{L}_{\mathrm{Yukawa}} = g_{e_i} \bar{e}_{iL} h e_{iR} + g_{u_i} \bar{u}_{iL} h u_{iR}$$
$$+ g_{d_i} \bar{d}_{iL} h d_{iR} + \mathrm{H.c.}, \qquad (23)$$

which gives the following chemical potential condition,

$$0 = \mu_h = \mu_{u_L} - \mu_{u_R} = \mu_{d_L} - \mu_{d_R} = \mu_{e_{iL}} - \mu_{e_{iR}}. \qquad (24)$$

Thus we see that for quark and charge leptons the left-handed and right-handed fields have the same chemical potential. Sphaleron transitions are efficient down to the decoupling temperature T_{sph} and, hence, we get

$$\mu_{u_L} + 2\mu_{d_L} + \mu_\nu = 0. \qquad (25)$$

At a temperature below electroweak phase transition the electric charge neutrality of the Universe holds. However, at this epoch, the top quark is already decoupled from the thermal plasma and, hence, does not take part in the charge neutrality condition. Therefore, we get

$$Q = 4(\mu_{u_L} + \mu_{u_R}) + 6\mu_W - 3(\mu_{d_L} + \mu_{d_R}) - \sum_{i=1}^3 (\mu_{e_{iL}} + \mu_{e_{iR}})$$

$$= 0. \qquad (26)$$

The charge current interactions:

$$\mathcal{L}_{\mathrm{int}}^{(W)} = g W_\mu^+ \bar{u}_L \gamma^\mu d_L + g W_\mu^+ \bar{e}_{iL} \gamma^\mu \bar{\nu}_{e_{iL}}. \qquad (27)$$

are also in thermal equilibrium below electroweak phase transition down to sphaleron decoupling temperature and hence satisfies the following chemical equilibrium condition:

$$\mu_W = \mu_{u_L} - \mu_{d_L}, \qquad (28)$$

$$\mu_W = \mu_\nu - \mu_{e_{iL}}, \quad \forall \, i. \qquad (29)$$

Thus, Eq. (29) ensures that three generations of charge leptons also have the same chemical potential.

Thus solving the Eqs. (24)–(29), we get the total baryon and lepton number densities in the visible sector:

$$n_B = -\frac{90}{29}\mu_\nu \quad \text{and} \quad n_L = \frac{201}{29}\mu_\nu. \qquad (30)$$

In Eq. (30), we have dropped the common factor $\frac{g_i T^3}{6} \times \beta$ and follow the same notation throughout the draft as we are

interested in the ratio of densities, rather than their individual values. From the above Eq. (30), we get the total $B - L$ asymmetry in the visible sector n_{B-L}:

$$(n_{B-L})_{\text{vis}} = -\frac{291}{29}\mu_\nu.$$

(31)

Moreover, from Eqs. (30) and (31), we get the total baryon asymmetry:

$$n_B = \frac{30}{97}(n_{B-L})_{\text{vis}}.$$

(32)

We assume that, the dark matter χ is also in thermal equilibrium with the visible sector via the dimension-eight operator $\mathcal{O}_8$ until the sphaleron decoupling temperature $T_{\text{sph}} > M_W$. This gives chemical equilibrium condition:

$$-\mu_\chi + \mu_\nu = 0$$

(33)

Thus, from Eqs. (31) and (33), we get the number density of χ asymmetry, which is also the $B - L$ number density in the dark sector:

$$n_\chi = (n_{B-L})_{\text{dark}} = -2\mu_\chi = \frac{58}{291}(n_{B-L})_{\text{vis}}.$$

(34)

The total n_{B-L} of the Universe, generated by the CP-violating out-of-equilibrium decay of the lightest right-handed neutrino (N_1), is the sum of n_{B-L} in the visible and dark sectors. Therefore, we get

$$\begin{aligned}(n_{B-L})_{\text{total}} &= (n_{B-L})_{\text{vis}} + (n_{B-L})_{\text{dark}} \\ &= (n_{B-L})_{\text{vis}} + \frac{58}{291}(n_{B-L})_{\text{vis}} \\ &= \frac{349}{291}(n_{B-L})_{\text{vis}}.\end{aligned}$$

(35)

Comparing Eq. (35) with Eq. (17) and using Eq. (32), we get the required asymmetry for observed DM abundance $\epsilon_\chi = 141.23(\eta/\kappa)(s/n_{N_1}^{eq}(T \to \infty))$. Thus for $\kappa \sim 0.01$ we get $\epsilon_\chi \sim 10^{-6}$. This is in accordance with the weakly interacting plasma with $\beta\mu \sim \epsilon_\chi \approx 10^{-6}$. Using Eq. (35) in Eqs. (32) and (34), we can get,

$$n_B = \frac{90}{349}(n_{B-L})_{\text{total}}, \qquad n_\chi = \frac{58}{349}(n_{B-L})_{\text{total}}$$

(36)

The Asymmetry generated in ϕ' can be written as,

$$n_{\phi'} = (n_{B-L})_{\text{total}}$$

(37)

The ratio of DM to baryon abundance, given by WMAP and the PLANCK data, to be $\Omega_{\text{DM}}/\Omega_B \approx 5$. This implies from Eqs. (36) and (37),
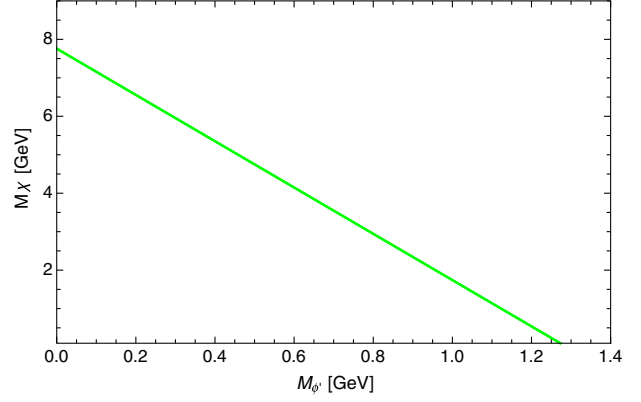

FIG. 6.   Allowed mass range for both DM candidates.

$$M_\chi = \frac{450M_p - 349M_{\phi'}}{58},$$

(38)

where M_p is the proton mass. We have shown the allowed range of masses of χ and ϕ' in Fig. 6. In what follows, we take $M_\chi = M_{\phi'} \approx 1$ GeV.

V. ANNIHILATION OF THE SYMMETRIC COMPONENT OF THE DARK MATTER

The symmetric component of χ and ϕ' can be efficiently depleted through the ϕ mediated interactions. In particular, $\phi - H$ mixing provides a portal for annihilation of χ and ϕ' to the SM particles. We show that when the extra scalar mass (M_{h_2}) is twice of the DM mass we get Breit-Wigner enhancement in the cross section which actually annihilates the symmetric component of the DM candidates, as shown in Figs. 7 and 8.

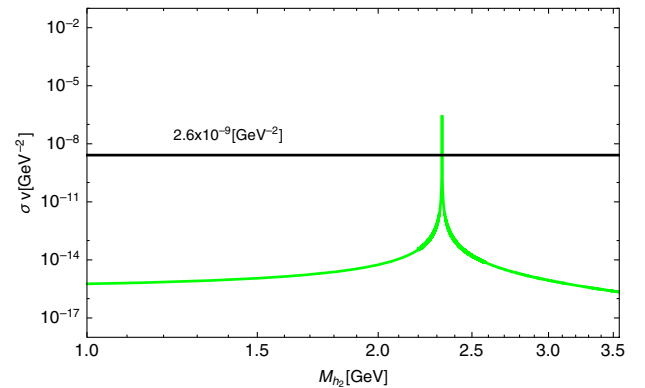The annihilation cross section for the process $\bar{\chi}\chi \to \bar{f}f$ is given by

FIG. 7.   The annihilation cross section of $\bar{\chi}\chi \to \bar{f}f$ as a function of M_{h_2} for a typical value of $\lambda_{\text{DM}} = 1 \times 10^{-2}$, $\lambda_{\text{H}\phi'} = 1 \times 10^{-3}$, $\mu_{\phi'} = 1 \times 10^{-3}$ and $\sin\gamma = 0.1$.
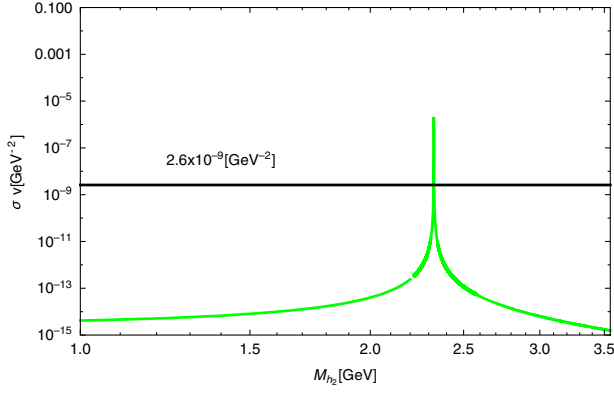
FIG. 8. The annihilation cross section of $\bar{\phi}'^{\dagger}\phi' \to \bar{f}f$ as a function of M_{h_2} for a typical value of $\lambda_{DM} = 1 \times 10^{-2}$, $\lambda_{H\phi'} = 1 \times 10^{-3}$, $\mu_{\phi'} = 1 \times 10^{-3}$ and $\sin\gamma = 0.1$.

$$\sigma_\chi v = \frac{\sqrt{s - 4M_f^2}}{16\pi s \sqrt{s}}$$
$$\times \frac{\lambda_{DM}^2 \lambda_f^2 \cos^2\gamma \sin^2\gamma}{[(s - M_{h_1}^2)^2 + \Gamma_{h_1}^2 M_{h_1}^2][(s - M_{h_2}^2)^2 + \Gamma_{h_2}^2 M_{h_2}^2]}$$
$$\times \{[2s - (M_{h_1}^2 + M_{h_2}^2)]^2 + [\Gamma_{h_1}M_{h_1} + \Gamma_{h_2}M_{h_2}]^2\}$$
$$\times \{(s - 2M_\chi^2)(s - 2M_f^2) - 2M_f^2(s - 2M_\chi^2)$$
$$- 2M_\chi^2(s - 2M_f^2) + 4M_\chi^2 M_f^2\} \quad (39)$$

where M_f represents the mass of SM fermions and $\lambda_f = M_f/v$. The decay width of h_1 is given by:

$$\Gamma_{h_1} = \cos^2\gamma \Gamma_{h_1}^{SM} + \sin^2\gamma \Gamma_{h_1}^{\bar{\chi}\chi} + \Gamma_{h_1}^{h_2 h_2} + \Gamma_{h_1}^{\phi'^{\dagger}\phi'}, \quad (40)$$

where $\Gamma_{h_1}^{SM} = 4.2$ MeV,

$$\Gamma_{h_1}^{\bar{\chi}\chi} = M_{h_1} \frac{\lambda_{DM}^2}{8\pi} \left[1 - \frac{4M_\chi^2}{M_{h_1}^2}\right]^{\frac{3}{2}}, \quad (41)$$

$$\Gamma_{h1}^{h_2 h_2} = \frac{\lambda_{eff}^2}{32\pi M_{h_1}} \left[1 - \frac{4M_{h_2}^2}{M_{h_1}^2}\right]^{\frac{1}{2}} \quad (42)$$

and

$$\Gamma_{h_1}^{\phi'^{\dagger}\phi'} = \frac{(\mu_\phi' \sin\gamma + \lambda_{H\phi'} v \cos\gamma)^2}{32\pi M_{h_1}} \left[1 - \frac{4M_{\phi'}^2}{M_{h_1}^2}\right]^{\frac{1}{2}}. \quad (43)$$

The decay width of h_2 is given by:

$$\Gamma_{h_2} = \sum_f \frac{C_f M_{h_2} \sin^2\gamma}{8\pi} \left(\frac{M_f}{v}\right)^2 \left[1 - \frac{4M_f^2}{M_{h_2}^2}\right]^{3/2}$$
$$+ \frac{M_{h_2} \lambda_{DM}^2 \cos^2\gamma}{8\pi} \left[1 - \frac{4M_\chi^2}{M_{h_2}^2}\right]^{3/2}, \quad (44)$$

where C_f accounts the color factor of SM fermions.

The annihilation cross section for the process $\phi'^{\dagger}\phi' \to \bar{f}f$ is given by

$$\sigma_{\phi'} v = \frac{(s - 4M_f^2)^{\frac{3}{2}}}{8\pi s \sqrt{s}}$$
$$\times \left[\frac{\lambda_1'^2 \lambda_2'^2}{(s - M_{h_2})^2 + \Gamma_{h_2}^2 M_{h_2}^2} + \frac{\lambda_1''^2 \lambda_2''^2}{(s - M_{h_1})^2 + \Gamma_{h_1}^2 M_{h_1}^2} \right.$$
$$+ 2\lambda_1' \lambda_2' \lambda_1'' \lambda_2''$$
$$\left. \times \left(\frac{(s - M_{h_2}^2)(s - M_{h_1}^2) + \Gamma_{h_1} M_{h_1} \Gamma_{h_2} M_{h_2}}{[(s - M_{h_2}^2)^2 + \Gamma_{h_2}^2 M_{h_2}^2][(s - M_{h_1}^2)^2 + \Gamma_{h_1}^2 M_{h_1}^2]} \right) \right]$$
$$\quad (45)$$

where $\lambda_1' = \mu_\phi' \cos\gamma - \lambda_{H\phi'} v \sin\gamma$, $\lambda_2' = -(M_f/v) \sin\gamma$, $\lambda_1'' = \mu_\phi' \sin\gamma + \lambda_{H\phi'} v \cos\gamma$ and $\lambda_2'' = (M_f/v) \cos\gamma$. The Γ_{h_1} and Γ_{h_2} are given in Eqs. (40) and (44).

In our case, $\bar{\chi}\chi$ is annihilating dominantly to a pair of muons. In Eqs. (39) and (45), the unknown parameters which dominantly contribute to the annihilation cross section are the mass of h_2, i.e., M_{h_2}, and the singlet-doublet Higgs mixing, i.e., $\sin\gamma$, the coupling of h_2 with χ, i.e., λ_{DM}, and the coupling of h_2 with ϕ', i.e., $\lambda_{H\phi'}$. However, these parameters are strongly constrained by invisible Higgs decay [24], relic abundance of DM measured by PLANCK [3] and WMAP [2], and spin-independent direct detection cross sections at XENON100 [25], LUX [26], XENON1T [27] and CRESST-II [28] and the Higgs signal strength measured at LHC [29,30]. For a typical value of the parameters $\lambda_{DM} = 1 \times 10^{-2}$, $\lambda_{H\phi'} = 1 \times 10^{-3}$, $\mu_{\phi'} = 1 \times 10^{-3}$ and $\sin\gamma = 0.1$, we have plotted σv as a function of M_{h_2} in Figs. 7, 8, respectively.

As shown in Figs. 7, 8, the value of the χ and ϕ' annihilation cross section $\sigma v < \langle \sigma|v| \rangle_F = 2.6 \times 10^{-9}/\text{GeV}^2$ in most of the parameter space except at the resonance, where $\sigma v > \langle \sigma|v| \rangle_F$. A crucial observation here is that mass of h_2 has to be twice the DM mass in order to get a large cross section via resonance. Note that a large cross section is required to deplete the symmetric component of the DM.

VI. PHENOMENOLOGICAL CONSTRAINTS

A. Higgs signal strength

The signal strength of SM-like Higgs in a particular channel $h_1 \to xx$ can be measured at LHC and can be defined as

$$\mu_{h_1 \to xx} = \frac{\sigma_{h_1}}{\sigma_{h_1}^{\text{SM}}} \frac{\text{Br}_{h_1 \to xx}}{\text{Br}_{h_1 \to xx}^{\text{SM}}} = \frac{\cos^4 \gamma \Gamma_{h_1}^{\text{SM}}}{\Gamma_{h_1}}, \qquad (46)$$

where Γ_{h_1} is given by Eq. (40). In the absence of any new physics, $\mu = 1$. However, in our case, the mixing between the two Higgses can reduce the signal strength of SM-like Higgs. Therefore, the mixing cannot be arbitrarily large and can be strongly constrained from the observation. The combined signal strength is measured to be $\mu = 1.17 \pm 0.1$ [29,30]. In Figs. 9 and 10, we have shown the contours of different values of μ in the planes of λ_{DM} and $\lambda_{\text{H}\phi'}$ with $\sin\gamma$, respectively. From the Figs. 9 and 10 we see that as the mixing increases the signal strength reduces accordingly. For an optimistic low value $\mu = 0.80$, the allowed mixing angle can be as large as 0.28. Thus in the rest of the draft we restrict the range of $\sin\gamma$ up to 0.3.

B. Constraints from invisible Higgs decay

The singlet-doublet Higgs mixing in this model allows the SM-like Higgs h_1 to decay via invisible channels:
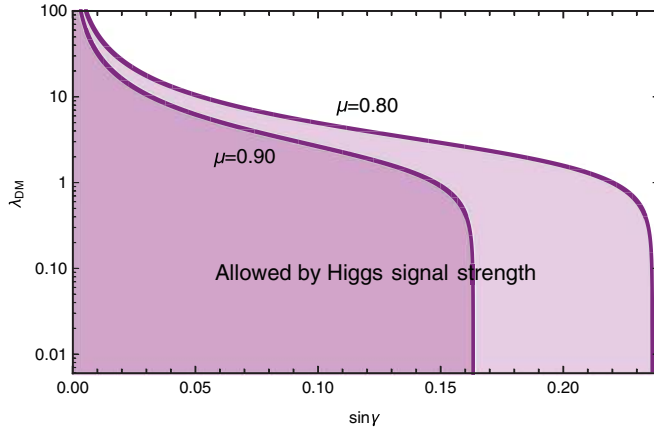

FIG. 9. Contours of signal strength μ of the SM-like Higgs h_1 in the plane of λ_{DM} versus $\sin\gamma$.
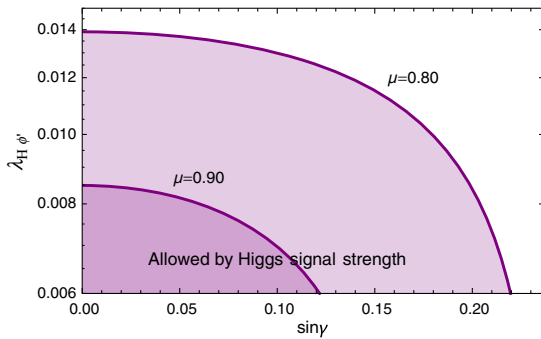

FIG. 10. Contours of signal strength μ of the SM-like Higgs h_1 in the plane of $\lambda_{\text{H}\phi'}$ versus $\sin\gamma$.

$h_1 \to h_2 h_2$, $h_1 \to \bar{\chi}\chi$ and $h_1 \to \bar{\phi}'^\dagger \phi'$. The branching ratio for the invisible Higgs decay can be defined as

$$\text{Br}_{\text{inv}} = \frac{\sin^2\gamma \Gamma_{h_1}^{\bar{\chi}\chi} + [\text{Br}(h_2 \to \bar{\chi}\chi + h_2 \to \phi'^\dagger \phi')]\Gamma_{h_1}^{h_2 h_2} + \Gamma_{h_1}^{\phi'^\dagger \phi'}}{\cos^2\gamma \Gamma_{h_1}^{\text{SM}} + \sin^2\gamma \Gamma_{h_1}^{\bar{\chi}\chi} + \Gamma_{h_1}^{h_2 h_2} + \Gamma_{h_1}^{\phi'^\dagger \phi'}},$$

(47)

where $\Gamma_{h_1}^{\bar{\chi}\chi}$, $\Gamma_{h_1}^{h_2 h_2}$ and $\Gamma_{h_1}^{\phi'^\dagger \phi'}$ are given by Eqs. (41)–(43), respectively. Note that LHC give an upper bound to the invisible Higgs decay to be $\text{Br}_{\text{inv}} \leq 24\%$ [24]. For a given h_2 mass, the allowed invisible Higgs decay width will constraint λ_{DM}, $\lambda_{\text{H}\phi'}$ and $\sin\gamma$ as we discuss below.

C. Constraints from direct detection of dark matter

The singlet-doublet scalar mixing also allows the DM χ and ϕ' to scatter off the nucleus at terrestrial laboratories. The spin independent DM-nucleon scattering cross section can be written as [31–34]

$$\sigma^{SI} = \frac{\mu_r^2}{\pi A^2} [Z f_p + (A - Z) f_n]^2 \qquad (48)$$

Where the Z and A are the atomic and mass numbers of the target nucleus. In Eq. (48), the reduced mass $\mu_r = M_\chi m_n / (M_\chi + m_n)$, where m_n is the mass of the nucleon (proton or neutron) and f_p and f_n are the effective interaction strengths of DM with proton and neutron of the target nucleus and are given by:

$$f_{p,n} = \sum_{q=u,d,s} f_{T_q}^{p,n} \alpha_q \frac{m_{p,n}}{m_q} + \frac{2}{27} f_{TG}^{p,n} \sum_{q=c,t,b} \alpha_q \frac{m_{p,n}}{m_q}, \quad (49)$$

where in case of χ DM can be written as

$$\alpha_q = \lambda_{\text{DM}} \left(\frac{m_q}{v}\right) \left[\frac{1}{M_{h_2}^2} - \frac{1}{M_{h_1}^2}\right] \sin\gamma \cos\gamma. \qquad (50)$$

where in case of ϕ' DM can be written as

$$\alpha_q = \left(\frac{\mu_{\phi'} \sin\gamma + \lambda_{\text{H}\phi'} v \cos\gamma}{M_{h_2}^2} - \frac{\mu_{\phi'} \cos\gamma - \lambda_{\text{H}\phi'} v \sin\gamma}{M_{h_1}^2}\right) \quad (51)$$

In Eq. (49) above, the $f_{T_q}^{p,n}$ are given by $f_{Tu}^{(p)} = 0.020 \pm 0.004$, $f_{Td}^{(p)} = 0.026 \pm 0.005$, $f_{Ts}^{(p)} = 0.118 \pm 0.062$, $f_{Tu}^{(n)} = 0.014 \pm 0.003$, $f_{Td}^{(n)} = 0.036 \pm 0.008$, $f_{Ts}^{(n)} = 0.118 \pm 0.062$ [35]. The coupling of DM with the gluons in target nuclei is parametrized by

$$f_{TG}^{p,n} = 1 - \sum_{q=u,d,s} f_{T_q}^{p,n}. \qquad (52)$$

We summarize all the constraints from invisible Higgs decay, relic abundance of DM, and null detection of DM at CRESST-II [28] as the allowed regions in the plane of λ_{DM} versus $\sin\gamma$ in Fig. 11 and $\lambda_{\mathrm{H}\phi'}$ versus $\sin\gamma$ in Fig. 12. We see from Fig. 11 that the region above the top purple line is not allowed by the invisible Higgs decay, where $\mathrm{Br}_{\mathrm{inv}} \geq 24\%$ [24]. The region below the bottom red line gives large relic abundance of DM since in this
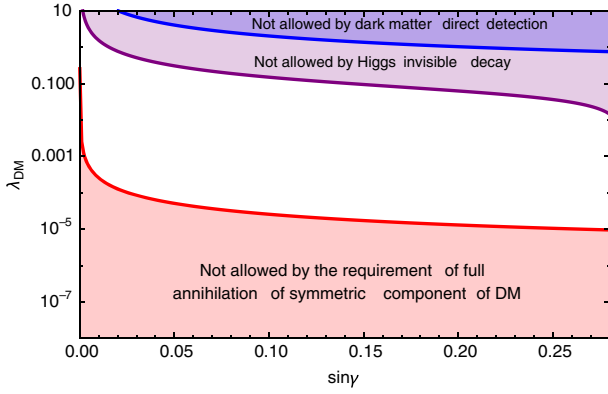

FIG. 11. Allowed regions in the plane of λ_{DM} versus $\sin\gamma$. The region above the top purple line is disallowed by invisible Higgs decay, i.e., $\mathrm{Br}_{\mathrm{inv}} \geq 24\%$. The region below the bottom red line is disallowed because $\sigma v < 2.6 \times 10^{-9}/\mathrm{GeV}^2$ and give large relic abundance. The regions above the Blue line is disallowed by the spin independent direct detection cross sections at CRESST-II 2016 for DM mass 1 GeV. we fix $M_{h_2} \approx 2M_\chi$.
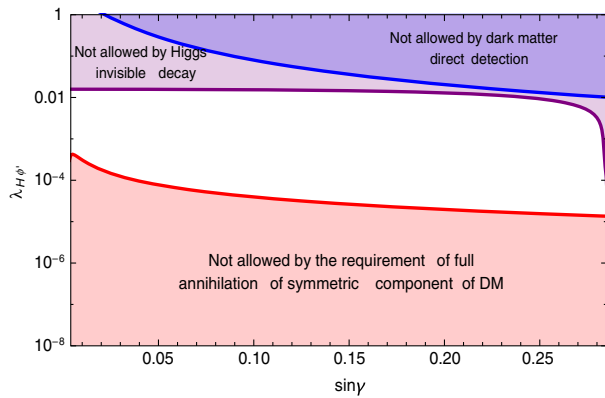

FIG. 12. Allowed regions in the plane of $\lambda_{H\phi'}$ versus $\sin\gamma$. The region above the top purple line is disallowed by invisible Higgs decay, i.e., $\mathrm{Br}_{\mathrm{inv}} \geq 24\%$. The region below the bottom red line is disallowed because $\sigma v < 2.6 \times 10^{-9}/\mathrm{GeV}^2$ and give large relic abundance. The regions above the blue line is disallowed by the spin independent direct detection cross sections at CRESST-II 2016 for DM mass 1 GeV. we fix $M_{h_2} \approx 2M'_\phi$.

region $\sigma v(\bar\chi\chi, \phi'^\dagger\phi' \to \bar{f}f) < 2.6 \times 10^{-9}/\mathrm{GeV}^2$. Since most of the annihilation occurs at the resonance, we fix $M_{h_2} \approx 2M_\chi \approx 2M'_\phi$. The blue line indicates the spin independent direct detection cross section from the CRESST-II [28] detector results $\sigma^{\mathrm{SI}} = 10^{-38}$ cm^2 corresponding to a DM mass: $M_\chi = M_{\phi'} = 1$ GeV. Therefore, the region above to that line is not allowed. Thus, are we left with a white allowed patch in the plane of λ_{DM} versus $\sin\gamma$ and $\lambda_{\mathrm{H}\phi'}$ versus $\sin\gamma$.

VII. CONCLUSION

In this paper, we extended the standard model by including a dark sector which consists of three generations of heavy right-handed neutrinos N_{iR}, $i = 1, 2, 3$, a singlet Dirac fermion χ, and a singlet scalar ϕ', where the latter two particles represent the DM. These particles are charged under an extended symmetry $U(1)_{B-L} \times U(1)_D \times Z_2$, while remaining inert with respect to the SM symmetry. An additional singlet scalar ϕ_{B-L} was introduced to break the $U(1)_{B-L}$ gauge symmetry at a high scale, say 10^{10} GeV. The breaking of $B-L$ symmetry at a high scale not only gave large Majorana masses to heavy right-handed neutrinos but also made Z_{B-L} super heavy. The global $U(1)_D$ symmetry, which was softly broken by dimension-eight operators, provides a distinction between N_R and χ since they carry the same charges under the $U(1)_{B-L} \times Z_2$.

In the early Universe, the CP-violating out-of-equilibrium decay of lightest heavy right-handed neutrino to χ and ϕ' generates a net DM asymmetry. The latter is then transferred to the visible sector via a dimension-eight operator $(\bar\chi LH)^2/M_{\mathrm{asy}}^4$ which conserves $B-L$ symmetry and is in thermal equilibrium down to the sphaleron decoupling temperature T_{sph}. $B+L$ violating sphaleron transitions are in thermal equilibrium down to a temperature T_{sph} and, hence, can convert the $B-L$ asymmetry in the visible sector to a net B asymmetry while the $B-L$ asymmetry in the dark sector remains untouched. As a result, we get a net asymmetric DM abundance (given in terms of $B-L$ asymmetry) comparable to baryon asymmetry for $M_\chi \sim M'_\phi \sim M_p$, where M_p represents the proton mass. An additional light singlet scalar ϕ was introduced, which helped in annihilating the symmetric component of the DM through its mixing with the SM Higgs. We found that the efficient annihilation of the symmetric component of DM requires the singlet scalar mass to be around twice the DM mass irrespective of all other parameters in the model. Since the observed DM abundance gives the DM mass $M_\chi = M_{\phi'} \approx 1$ GeV, we get the singlet scalar mass ≈ 2 GeV, which can be searched at the collider and via an indirect gamma ray search.

The neutrinos are massless at the tree level since the right-handed neutrinos are odd under the Z_2 symmetry and

are decoupled from the visible sector. However, at the one-loop level, the neutrinos acquired masses via a dimension-eight operator $(\bar{N}_R L H)^2/\Lambda^4$. We showed that sub-eV masses of neutrinos require the $B - L$ breaking scale to be around $\Lambda \approx 10^{11}$ GeV.

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