Leptogenesis and Dark Matter from Low Scale Seesaw

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In this paper, we perform a detail analysis on leptogenesis and dark matter form low scale seesaw. In the framework of $\nu^2$HDM, we further introduce one scalar singlet $\phi$ and one Dirac fermion singlet $\chi$, which are charged under a $Z_2$ symmetry. Assuming the coupling of $\chi$ is extremely small, it serves as a FIMP dark matter. The heavy right hand neutrinos $N$ provide a common origin for tiny neutrino mass (via seesaw mechanism), leptogenesis (via $N \to \ell_L \Phi^*_\nu, \bar{\ell}_L \Phi_\nu$) and dark matter (via $N \to \chi \phi$). With hierarchical right hand neutrino masses, the explicit calculation shows that success thermal leptogenesis is viable even for TeV scale $N_1$ with $0.4 \text{ GeV} \lesssim v_\nu \lesssim 1 \text{ GeV}$ and lightest neutrino mass $m_1 \lesssim 10^{-11} \text{ eV}$. In such scenario, light FIMP dark matter in the keV to MeV range is naturally expected. The common parameter space for neutrino mass, natural leptogenesis and FIMP DM is also obtained in this paper.
I. INTRODUCTION

Besides the success of standard model (SM), there are still several open questions. In particular, tiny neutrino mass, baryon asymmetry of the Universe (BAU) and dark matter (DM) are the three outstanding evidences that require physics beyond SM. The discovery of neutrino oscillations [1, 2] indicate that the mass of neutrinos are at sub-eV scale, which is at least six order of magnitudes smaller than charged leptons. Known as type-I seesaw mechanism [3, 4], this extensively considered way to naturally incorporate neutrino masses is via introducing three right hand neutrinos $N$ together with high scale Majorana masses of $N$, 

$$
-\mathcal{L}_Y \supset y_L \Phi N + \frac{1}{2} \overline{N}^c m_N^c N + \text{h.c.},
$$

where $\Phi$ is the SM Higgs doublet. After spontaneous electroweak symmetry breaking, neutrinos achieve masses as

$$
m_\nu = -\frac{v^2}{2} y m_N^{-1} y T.
$$

Typically, $m_\nu \sim \mathcal{O}(0.1)$ eV is obtained with $y \sim \mathcal{O}(1)$ and $m_N \sim \mathcal{O}(10^{14})$ GeV. Meanwhile, the heavy neutrino can also account for BAU via leptogenesis [5]. For canonical thermal leptogenesis with hierarchal right hand neutrinos, an upper limit on the CP asymmetry exists, thus a lower limit on mass of lightest right hand neutrino $M_1$ should be satisfied [6],

$$
M_1 \gtrsim 5 \times 10^8 \text{GeV} \left(\frac{v}{246 \text{ GeV}}\right)^2.
$$

Therefore, both tiny neutrino mass and leptogenesis favor high scale $N$ in type-I seesaw. However for such high scale $N$, a naturalness problem might arise [7]. By requiring radiative corrections to the $m_\Phi^2 \Phi^\dagger \Phi$ term no larger than 1 TeV$^2$, it is found that [8]

$$
M_1 \lesssim 3 \times 10^7 \text{ GeV} \left(\frac{v}{246 \text{ GeV}}\right)^{2/3},
$$

should be satisfied. Clearly, naturalness is incompatible with leptogenesis. One viable pathway to overcome this is lowering the leptogenesis scale by imposing resonant leptogenesis [9], ARS mechanism via neutrino oscillation [10, 11], or from Higgs decays [12, 13]. All the success of these scenarios depend on the degenerate mass of right hand neutrinos [14], which seems is another sense of unnatural. An alternative scenario with hierarchal right hand neutrinos is employing intrinsic low scale neutrino mass model, e.g., $\nu$2HDM [15, 16] or Scotogenic model [17–24]. In this paper, we consider the $\nu$2HDM [25]. Based on previous brief discussion in Ref. [16, 26, 27], we perform a detailed analysis on leptogenesis, especially focus on dealing with the corresponding Boltzmann equations to obtain the viable parameter space.
On the other hand, dark matter accounts for more than five times the proportion of visible baryonic matter in our current cosmic material field. In principle, one can regard the lightest right hand neutrino $N_1$ at keV scale as sterile neutrino DM [28–31]. However, various constraints leave a quite small viable parameter space [32]. Meanwhile, leptogenesis with two hierarchal right hand neutrinos is actually still at high scale [21, 33, 34]. In this paper, we further introduce a dark sector with one scalar singlet $\phi$ and one Dirac fermion singlet $\chi$, which are charged under a $Z_2$ symmetry [35]. The stability of DM $\chi$ is protected by the $Z_2$ symmetry, therefore the tight X-ray limits can be avoided [32]. In light of the null results from DM direct detection [36] and indirect detection [37], we consider $\chi$ as a FIMP DM [38].

The structure of the paper is as follows. In Sec. II, we briefly introduce our model. Leptogenesis with hierarchal right hand neutrinos is discussed in Sec. III. The relic abundance of FIMP DM $\chi$ and constraint from free streaming length are considered in Sec. IV. Viable parameter space for leptogenesis and DM is obtained by a random scan in Sec. V. We conclude our work in Sec. VI.

II. THE MODEL

The original TeV-scale $\nu$2HDM for neutrino mass was proposed in Ref. [25]. The model is extended by one neutrinophilic scalar doublet $\Phi_\nu$ with same quantum numbers as SM Higgs doublet $\Phi$ and three right hand heavy neutrino $N$. To forbid the direct type-I seesaw interaction $L\Phi N$, a global $U(1)_L$ symmetry should be employed, under which $L_\Phi = 0$, $L_{\Phi_\nu} = -1$ and $L_N = 0$. Therefore, $\Phi_\nu$ will specifically couple to $N$, and $\Phi$ couple to quarks and charge leptons as in SM. For the dark sector, one scalar singlet $\phi$ and one Dirac fermion singlet $\chi$ are further introduced, which are charged under a $Z_2$ symmetry. Provided $m_\chi < m_\phi$, then $\chi$ serves as DM candidate.

The scalar doublets could be denoted as

$$\Phi = \begin{pmatrix} \phi^+ \\ v + \phi^0_r + i \phi^0_i \end{pmatrix}, \quad \Phi_\nu = \begin{pmatrix} \phi^+_\nu \\ v_\nu + \phi^0_{\nu_r} + i \phi^0_{\nu_i} \end{pmatrix}.$$ (5)

The corresponding Higgs potential is then

$$V = m_\Phi^2 \Phi^\dagger \Phi + m_{\Phi_\nu}^2 \Phi_\nu^\dagger \Phi_\nu + m_\phi^2 \phi^\dagger \phi + \frac{\lambda_1}{2} (\Phi^\dagger \Phi)^2 + \frac{\lambda_2}{2} (\Phi_\nu^\dagger \Phi_\nu)^2$$

$$+ \lambda_3 (\Phi^\dagger \Phi_\nu^\dagger \Phi_\nu) + \lambda_4 (\Phi^\dagger \Phi_\nu)^2 (\Phi^\dagger_\nu \Phi) - (\mu^2 \Phi^\dagger \Phi_\nu + h.c.)$$

$$+ \frac{\lambda_5}{2} (\phi^\dagger \phi)^2 + \lambda_6 (\phi^\dagger \phi) (\Phi^\dagger \Phi) + \lambda_7 (\phi^\dagger \phi) (\Phi^\dagger_\nu \Phi_\nu),$$ (6)

where the $U(1)_L$ symmetry is broken explicitly but softly by the $\mu^2$ term. For the unbroken $Z_2$ symmetry, $\langle \phi \rangle = 0$ should be satisfied. Meanwhile, VEVs of Higgs doublets in terms of parameters of the Higgs
potential can be found by deriving the minimization condition
\[ \begin{align*}
  v \left[ m_\Phi^2 + \frac{\lambda_1}{2} v^2 + \frac{\lambda_3 + \lambda_4}{2} v^2 \right] - \mu^2 v = 0, \\
v_\nu \left[ m_{\Phi,\nu}^2 + \frac{\lambda_2}{2} v_\nu^2 + \frac{\lambda_3 + \lambda_4}{2} v^2 \right] - \mu^2 v = 0.
\end{align*} \tag{8, 9} \]

Taking the parameter set
\[ m_\Phi^2 < 0, m_{\Phi,\nu}^2 > 0, |\mu^2| \ll m_{\Phi,\nu}^2, \tag{10} \]
we can obtain the relations of VEVs as
\[ v \simeq \sqrt{-\frac{2m_\Phi^2}{\lambda_1}}, v_\nu \simeq \frac{\mu^2 v}{m_{\Phi,\nu}^2 + (\lambda_3 + \lambda_4)v^2/2}. \tag{11} \]

Typically, \( v_\nu \sim 1 \text{ GeV} \) is obtained with \( \mu \sim 10 \text{ GeV} \) and \( m_{\Phi,\nu} \sim 100 \text{ GeV} \). Since \( \mu^2 \) term is the only source of \( U(1)_L \) breaking, radiative corrections to \( \mu^2 \) are proportional to \( \mu^2 \) itself and are only logarithmically sensitive to the cutoff \[39\]. Thus, the VEV hierarchy \( v_\nu \ll v \) is stable against radiative corrections \[40, 41\].

After SSB, the physical Higgs bosons are given by \[42\]
\[ \begin{align*}
  H^+ &= \phi^0_\nu \cos \beta - \phi^+ \sin \beta, \\
  A &= \phi^0_\nu \cos \beta - \phi^0 \sin \beta, \\
  H &= \phi^0_r \cos \alpha - \phi^0 \sin \alpha, \\
  h &= \phi^0_r \cos \alpha + \phi^0 \sin \alpha, \tag{12, 13} \end{align*} \]
where the mixing angles \( \beta \) and \( \alpha \) are determined by
\[ \tan \beta = \frac{v_\nu}{v}, \quad \tan 2\alpha \simeq \frac{2v_\nu}{v} - \frac{\mu^2 + (\lambda_3 + \lambda_4)v^2_v}{-\mu^2 + \lambda_1 v^2_v}. \tag{14} \]

Neglecting terms of \( O(v^2_\nu) \) and \( O(\mu^2) \), masses of the physical Higgs bosons are
\[ m_{H^+}^2 \simeq m_{\Phi,\nu}^2 + \frac{1}{2} \lambda_3 v^2, \\
m_{H}^2 \simeq m_{H^+}^2 + \frac{1}{2} \lambda_4 v^2, \\
m_{A}^2 \simeq m_{H}^2 + \frac{1}{2} \lambda_1 v^2. \tag{15} \]

Since the mixing angles are suppressed by the small value of \( v_\nu \), \( h \) is almost identically to the 125 GeV SM Higgs boson \[43, 44\]. A degenerate mass spectrum of \( \Phi_\nu \) as \( m_{H^+} = m_H = m_A = m_{\Phi,\nu} \) is adopted in our following discussion for simplicity, which is certainly allowed by various constraints \[45\]. Due to the unbroken \( Z_2 \) symmetry, the dark scalar singlet \( \phi \) do not mix with the Higgs doublets.

The new Yukawa interaction and mass terms are
\[- \mathcal{L}_Y \supset \bar{y}L \tilde{\Phi}_\nu N + \lambda \chi \phi N + \frac{1}{2} \overline{\chi m_N N + m_\chi \chi \chi + h.c.,} \tag{16} \]
where \( \tilde{\Phi}_\nu = i\sigma_2 \Phi^*_\nu \). Similar to the canonical Type-I seesaw \[3\], the mass matrix for light neutrinos can be derived from Eq. \[16\] as:
\[ m_\nu = -\frac{v^2_\nu}{2} y m_{\nu}^{-1} y^T = U_{\text{PMNS}} \hat{m}_\nu U_{\text{PMNS}}^T, \tag{17} \]
where \( \hat{m}_\nu = \text{diag}(m_1, m_2, m_3) \) is the diagonalized neutrino mass matrix, and \( U_{\text{PMNS}} \) is the PMNS (Pontecorvo-Maki-Nakagawa-Sakata) matrix:

\[
U_{\text{PMNS}} = \begin{pmatrix}
    c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\
    -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{-i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{-i\delta} & s_{23}c_{13} \\
    s_{12}s_{23} - c_{12}c_{23}s_{13}e^{-i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{-i\delta} & c_{23}c_{13}
\end{pmatrix} \times \text{diag}(e^{i\varphi_1/2}, 1, e^{i\varphi_2/2})
\]

(18)

Here, we use abbreviations \( c_{ij} = \cos \theta_{ij} \) and \( s_{ij} = \sin \theta_{ij} \), \( \delta \) is the Dirac phase and \( \varphi_1, \varphi_2 \) are the two Majorana phases. Due to smallness of \( v_\nu \), TeV scale \( m_N \) could be viable to realise \( 0.1 \text{ eV} \) scale light neutrino masses. Using the Casas-Ibarra parametrization \[46, 47\], the Yukawa matrix \( y \) can be expressed in terms of neutrino oscillation parameters

\[
y = \sqrt{2} \frac{v_\nu}{v_\nu} U_{\text{PMNS}} \hat{m}_\nu^{1/2} R (\hat{m}_N)^{1/2},
\]

(19)

where \( R \) is an orthogonal matrix in general and \( \hat{m}_N = \text{diag}(M_1, M_2, M_3) \) is the diagonalized heavy neutrino mass matrix. In this work, we parameterize matrix \( R \) as

\[
R = \begin{pmatrix}
    \cos \omega_{12} & -\sin \omega_{12} & 0 \\
    \sin \omega_{12} & \cos \omega_{12} & 0 \\
    0 & 0 & 1
\end{pmatrix} \begin{pmatrix}
    \cos \omega_{13} & 0 & -\sin \omega_{13} \\
    0 & 1 & 0 \\
    \sin \omega_{13} & 0 & \cos \omega_{13}
\end{pmatrix} \begin{pmatrix}
    1 & 0 & 0 \\
    0 & \cos \omega_{23} & -\sin \omega_{23} \\
    0 & \sin \omega_{23} & \cos \omega_{23}
\end{pmatrix},
\]

(20)

where \( \omega_{12,13,23} \) are arbitrary complex angles.

**III. LEPTOGENESIS**

Now we consider the leptogenesis in this model. The lepton asymmetry is generated by the out-of-equilibrium CP-violating decays of right hand neutrino \( N \rightarrow \ell_L \Phi^*_\nu, \bar{\ell}_L \Phi_\nu \). Neglecting the flavor effect \[48\], the CP asymmetry is given by

\[
\epsilon_i = \frac{1}{8\pi (y_i \bar{y}_i)_{ii}} \sum_{j \neq i} \text{Im}[(y_i \bar{y}_j)_{ij}] G \left( \frac{M^2_{ij}}{M^2_i}, \frac{m^2_{\Phi}}{M^2_i} \right),
\]

(21)

where the function \( G(x, y) \) is defined as \[54\]

\[
G(x, y) = \sqrt{x} \left[ \frac{(1 - y)^2}{1 - x} + 1 + \frac{1 - 2y + x}{(1 - y^2)^2} \ln \left( \frac{x - y^2}{1 - 2y + x} \right) \right].
\]

(22)

Using the parametrization of Yukawa coupling \( y \) in Eq. (19), it is easy to verify

\[
y^\dagger y = \frac{2}{v_\nu^2} \hat{m}_N^{1/2} R^\dagger \hat{m}_\nu R \hat{m}_N^{1/2}.
\]

(23)
Hence, the matrix $y^\dagger y$ does not depend on the PMNS matrix, which means that the complex matrix $R$ is actually the source of CP asymmetry $\epsilon_i$. The asymmetry is dominantly generated by the decay of $N_1$. Further considering the hierarchal mass spectrum $m_{3\nu}^2 \ll M_1^2 \ll M_{2,3}^2$, the asymmetry $\epsilon_1$ is simplified to

$$\epsilon_1 \simeq -\frac{3}{16\pi(y^\dagger y)_{11}} \sum_{j=2,3} \text{Im}[(y^\dagger y)_{1j}^2] \frac{M_1}{M_j}$$  \hspace{1cm} (24)

Similar to the Davidson-Ibarra bound [6], an upper limit on $\epsilon_1$ can be derived

$$|\epsilon_1| \lesssim \frac{3}{16\pi} \frac{M_1 m_3}{v_\nu^2}.$$  \hspace{1cm} (25)

Comparing with the bound in type-I seesaw, the asymmetry could be enhanced due to the smallness of VEV $v_\nu$. Therefore, low scale leptogenesis seems to be viable in the $\nu2$HDM [16, 26]. Meanwhile, the washout effect is quantified by the decay parameter

$$K = \frac{\Gamma_1}{H(z = 1)}.$$  \hspace{1cm} (26)

where $\Gamma_1$ is the decay width of $N_1$, $H$ is the Hubble parameter and $z \equiv M_1/T$ with $T$ being the temperature of the thermal bath. The decay width is given by

$$\Gamma_1 = \frac{M_1}{8\pi} (y^\dagger y)_{11} \left(1 - \frac{m_{3\nu}^2}{M_1^2}\right)^2,$$  \hspace{1cm} (27)

and the Hubble parameter is

$$H = \sqrt{\frac{8\pi^3 g_*}{90}} \frac{T^2}{M_{pl}} = H(z = 1) \frac{1}{z^2},$$  \hspace{1cm} (28)

with $g_*$ the effective number of relativistic degrees of freedom and $M_{pl} = 1.2 \times 10^{19}$ GeV. Using Eq. (23), one can verify

$$K \simeq 897 \left(\frac{v}{v_\nu}\right)^2 \frac{(\hat{m}_\nu^R)_{11}}{\text{eV}},$$  \hspace{1cm} (29)

where $\hat{m}_\nu^R \equiv R^\dagger \hat{m}_\nu R$, and thus

$$(\hat{m}_\nu^R)_{11} = m_1 |\cos \omega_{12}|^2 |\cos \omega_{13}|^2 + m_2 |\sin \omega_{12}|^2 |\cos \omega_{13}|^2 + m_3 |\sin \omega_{13}|^2.$$  \hspace{1cm} (30)

It is obvious that the decay parameter $K$ does not depend on $\omega_{23}$, and it is also enhanced by smallness of $v_\nu$. Since $(\hat{m}_\nu^R)_{11}$ is typically of the order of $m_3 \sim 0.1$ eV, the decay parameter $K \simeq 5.4 \times 10^6$ when $v_\nu = 1$ GeV. So even with maximum asymmetry $\epsilon_1^{\text{max}} \sim -6.0 \times 10^{-7}$ for $M_1 = 10^5$ GeV obtained from Eq. (25), a rough estimation of final baryon asymmetry gives $Y_{\Delta B} \sim -10^{-3} \frac{\epsilon_1^{\text{max}}}{K} \sim 1.1 \times 10^{-16}$ for strong washout [49], which is far below current observed value $Y_{\Delta B}^{\text{obs}} = (8.72 \pm 0.04) \times 10^{-11}$ [50]. Hence,
only obtaining an enhanced CP asymmetry \( \epsilon_1 \) is not enough, one has to deal with the washout effect more carefully.

One promising pathway is to reduce the decay parameter \( K \). For instance, if weak washout condition \( K \lesssim 1 \) is realised, then \( Y_{\Delta B} \sim -10^{-3} \epsilon_1^{\text{max}} \sim 6.0 \times 10^{-10} > Y_{\Delta B}^{\text{obs}} \). Thus, correct baryon asymmetry can be obtained by slightly tuning \( \epsilon_1 \). As pointed out in Ref. [21], small value of \( K \) can be realised by choosing small \( \omega_{12,13} \). In Fig. 1 we illustrate the dependence of \( K \) on lightest neutrino mass \( m_1 \) with \( v_\nu = 10 \text{ GeV} \). The left panel shows the special case \( \omega_{13} = 0 \), where Eq. (30) is simplified to \( \langle m_{\nu} R \rangle_{11} = m_1 |\cos \omega_{12}|^2 + m_2 |\sin \omega_{12}|^2 \geq \sqrt{\Delta m_{21}^2} |\sin \omega_{12}|^2 \). It is clear that the weak washout condition \( K < 1 \) favors \( |\omega_{12}| \lesssim 10^{-2} \) and \( m_1 \lesssim 10^{-6} \text{ eV} \). The right panel shows the special case \( \omega_{12} = 0 \). Similar results are observed with left panel.

On the other hand, the \( \Delta L = 2 \) washout processes become more significant for small \( v_\nu \) [16, 26]. Notably, for low scale seesaw, the narrow width condition \( \Gamma_1/M_1 \ll 1 \) is satisfied. Therefore, the evolution of lepton asymmetry and DM abundance actually decouple from each other [51, 52]. The evolution of abundance \( Y_{N_1} \) and lepton asymmetry \( Y_{\Delta L} \) are described by the Boltzmann equations

\[
\frac{dY_{N_1}}{dz} = -D(Y_{N_1} - Y_{N_1}^{\text{eq}}), \quad \frac{dY_{\Delta L}}{dz} = -\epsilon_1 D(Y_{N_1} - Y_{N_1}^{\text{eq}}) - WY_{\Delta L}.
\]

FIG. 1. Decay parameter \( K \) as a function of \( m_1 \) with \( v_\nu = 10 \text{ GeV} \). Because \( R \) must be a complex matrix, we have set \( \omega_{ijR} = \omega_{ijI} \).

\[\omega_{12} = (10^{-1} + 10^{-1}i)\]
\[\omega_{13} = (10^{-2} + 10^{-2}i)\]
\[\omega_{12} = (10^{-3} + 10^{-3}i)\]
\[\omega_{13} = (10^{-1} + 10^{-1}i)\]
\[\omega_{13} = (10^{-2} + 10^{-2}i)\]
\[\omega_{13} = (10^{-3} + 10^{-3}i)\]
The decay term is given by

\[ D = K z \frac{K_1(z)}{K_2(z)}. \]  

(33)

For the washout term, two contributions are considered, i.e., \( W = W_{ID} + W_{\Delta L=2} \), where the inverse decay term is

\[ W_{ID} = \frac{1}{4} K z^3 K_1(z), \]

(34)

and the \( \Delta L = 2 \) scattering term at low temperature is approximately \[ W_{\Delta L=2} \simeq \frac{0.186}{z^2} \left( \frac{246 \text{ GeV}}{v_{\nu}} \right)^4 \left( \frac{M_1}{10^{10} \text{ GeV}} \right)^2 \left( \frac{\bar{m}}{\text{eV}} \right)^2. \]

(35)

Here, \( \bar{m} \) is the absolute neutrino mass scale, which is calculated as

\[ \bar{m}^2 = m_1^2 + m_2^2 + m_3^2 = 3m_1^2 + \Delta m_{21}^2 + \delta m_{31}^2, \]

(36)

for normal hierarchy. According to latest global fit, we use the best fit values, i.e., \( \Delta m_{21}^2 = 7.39 \times 10^{-5} \text{ eV}^2 \) and \( \delta m_{31}^2 = 2.525 \times 10^{-3} \text{ eV}^2 \) \[53\]. For tiny lightest neutrino mass \( m_1 \ll 10^{-2} \text{ eV} \), we actually have \( \bar{m} \simeq \sqrt{\delta m_{31}^2} \sim 0.05 \text{ eV} \). Notably, the \( \Delta L = 2 \) scattering term would be greatly enhanced when \( v_{\nu} \ll v \), so this term is much more important than in vanilla leptogenesis. Then, the sphaleron processes convert the lepton asymmetry into baryon asymmetry as \[ Y_{\Delta B} = \frac{28}{79} Y_{(B-L)} = -\frac{28}{51} Y_{\Delta L}. \]

(37)

Fig. 2 shows the washout effect of \( \Delta L = 2 \) processes. In Fig. 2(a), weak washout scenario is considered by fixing \( K = 10^{-2}, |\epsilon_1| = 10^{-6}, M_1 = 10^6 \text{ GeV} \) while varying \( v_{\nu} = 10, 1, 0.1 \text{ GeV} \). It shows that for \( v_{\nu} = 10 \text{ GeV} \), the \( \Delta L = 2 \) effect is not obvious, but for \( v_{\nu} = 1 \text{ GeV} \), the final baryon asymmetry \( Y_{\Delta B} \) is diluted by over three orders of magnitude. While for \( v_{\nu} = 0.1 \text{ GeV} \), the \( \Delta L = 2 \) effect is so strong that the final baryon asymmetry is negligible. The strong washout scenario with \( K = 10^2, |\epsilon_1| = 10^{-4}, M_1 = 10^6 \text{ GeV} \) and varying \( v_{\nu} = 10, 1, 0.1 \text{ GeV} \) is illustrated in Fig. 2(b), where the final baryon asymmetry \( Y_{\Delta B} \) for \( v_{\nu} = 1 \text{ GeV} \) is decreased by about six orders comparing with the case for \( v_{\nu} = 10 \text{ GeV} \). Therefore, the \( \Delta L = 2 \) washout effects set a lower bound on \( v_{\nu} \), i.e., \( v_{\nu} \gtrsim 0.3 \text{ GeV} \) as suggested by Ref. \[16\]. Furthermore, since the \( \Delta L = 2 \) washout term is also proportional to \( M_1 \), the larger \( M_1 \) is, the more obvious the washout effect is. The corresponding results are depicted in Fig. 2(c) for the weak washout and Fig. 2(d) for the strong washout. In this way, for certain value of \( v_{\nu} \), an upper bound on \( M_1 \) can be obtained. For instance, when \( v_{\nu} = 1 \text{ GeV} \), then \( M_1 \lesssim 10^5 \text{ GeV} \) should be satisfied \[26\].
IV. DARK MATTER

In our extension of the $\nu$2HDM, the right-handed heavy neutrinos $N$ also couple with fermion singlet $\chi$ and scalar singlet $\phi$ via the Yukawa interaction. The complex Yukawa coupling coefficient $\lambda$ can lead to CP violation in $N$ decays, and eventually producing asymmetric DM $\chi$ [51]. Instead, we consider another interesting scenario, i.e., the FIMP case with the real coupling $|\lambda| \ll 1$ [52]. In this way, the interaction of DM $\chi$ is so weak that it never reach thermalization. Its relic abundance is determined by the freeze-in mechanism [56], which is obtained by solving the following Boltzmann equation

\[
\frac{dY_\chi}{dz} = D Y_{N_i} \text{BR}_\chi,
\]  

(38)
FIG. 3. Evolution of dark matter abundance with parameter $z = M_i/T$. We fix $K = 10$ in the left panel and $BR_\chi = 10^{-3}$ in the right panel. The dashed horizontal lines correspond to the estimated results with Eq. (39). DM mass $m_\chi$ is obtained by setting $\Omega_\chi h^2 = 0.12$ with the numerical results of $Y_\chi(\infty)$.

where $BR_\chi$ is the branching ratio of $N_1 \rightarrow \chi \phi$. Due to the FIMP nature of $\chi$, the hierarchical condition $BR_\chi \ll BR_\ell \simeq 1$ is easily satisfied. The out of equilibrium condition for $N_1 \rightarrow \chi \phi$ decay is $\Gamma_\chi / H(z = 1) \simeq BR_\chi \Gamma_1 / H(z = 1) = BR_\chi K < 1$. In following studies, we mainly take $BR_\chi < 10^{-2}$ and $K \lesssim 10$, thus the out of equilibrium condition is always satisfied. According to the above Boltzmann equation, we can estimate the asymptotic abundances of $\chi$ as\[ Y_\chi(\infty) \simeq Y_{N_1}(0) BR_\chi \left( 1 + \frac{15\pi \zeta(5)}{16\zeta(3)} K \right). \tag{39} \]

Then, the corresponding relic abundance is \[ \Omega_\chi h^2 = \frac{m_\chi s_0 Y_\chi(\infty)}{\rho_c} h^2 \simeq 0.12 \times \left( \frac{m_\chi}{\text{keV}} \right) \left( \frac{BR_\chi}{10^{-3}} \right) \left( \frac{0.009 + \frac{K}{44}}{\text{yr}^{-1}} \right), \tag{40} \] where $s_0 = 2891.2 \text{ cm}^{-3}$, $\rho_c = 1.05371 \times 10^{-5} \text{ h}^2 \text{ GeV cm}^{-3}$\[ [57]. \] Typically, the observed relic abundance can be obtained with $m_\chi \sim 4 \text{ keV}$, $BR_\chi \sim 10^{-3}$ and $K \sim 10$. The evolution of DM abundances are shown in Fig. (3). It is clear that when the temperature goes down to $z = m_\chi / T \sim 5$, the abundances $Y_\chi$ freeze in and keep at a constant. The left panel of Fig. (3) indicates that $m_\chi$ is inverse proportional to $BR_\chi$ when the decay parameter $K$ is a constant. For instance, sub-MeV scale light DM is obtained when $BR_\chi > 10^{-6}$ with $K = 10$. Right panel of Fig. (3) shows the impact of decay parameter $K$. Affected by the constant term before $K$ in Eq. (40), we can only conclude that the smaller the $K$ is, the larger the $m_\chi$ is. Besides, we also find that the discrepancy between the numerical and analytical results of $Y_\chi(\infty)$ increases when $K$ decreases. Therefore, we adopt the numerical result of $Y_\chi(\infty)$ for a more precise calculation in the
The dominant constraint on FIMP DM $\chi$ comes from its free streaming length, which describes the average distance a particle travels without a collision $^{52}$

$$r_{FS} = \int_{a_{rh}}^{a_{eq}} \frac{\langle v \rangle}{a^2 H} da \approx \frac{a_{NR}}{H_0 \sqrt{\Omega_R}} \left( 0.62 + \ln \left( \frac{a_{eq}}{a_{NR}} \right) \right), \quad (41)$$

where $\langle v \rangle$ is the averaged velocity of DM $\chi$, $a_{eq}$ and $a_{rh}$ represent scale factors in equilibrium and reheating, respectively. We use the results $H_0 = 67.3$ km s$^{-1}$Mpc$^{-1}$, $\Omega_R = 9.3 \times 10^{-5}$ and $a_{eq} = 2.9 \times 10^{-4}$ obtained from Ref. $^{58}$. The non-relativistic scale factor for FIMP DM is

$$a_{NR} = \frac{T_0}{2m_\chi} \left( \frac{g_{*0}}{g_{*rh}} \right)^{\frac{1}{3}} K^{-\frac{1}{2}}. \quad (42)$$

Taking $g_{*0} = 3.91$, $g_{*rh} = 106.75$ and $T_0 = 2.35 \times 10^{-4}$ eV, finally we can get

$$r_{FS} \simeq 2.8 \times 10^{-2} \frac{\text{keV}}{m_\chi} \left( \frac{50}{K} \right)^{\frac{1}{2}} \times \left( 1 + 0.09 \ln \left( \frac{m_\chi}{\text{keV}} \frac{K}{50}^{\frac{1}{2}} \right) \right) \text{Mpc}. \quad (43)$$

The most stringent bound on $r_{FS}$ comes from small structure formation $r_{FS} < 0.1$ Mpc $^{59}$. The relationship between the mass of $\chi$ and its free streaming length is depicted in Fig. (4). Basically speaking, warm DM is obtained for $m_\chi \sim 10$ keV while $K \in [0.01, 100]$. Meanwhile, $\chi$ becomes cold DM when $\chi$ is sufficient heavy and/or the decay parameter $K$ is large enough.
V. COMBINED ANALYSIS

After studying some benchmark points, it would be better to figure out the viable parameter space for success leptogenesis and DM. We then perform a random scan over the following parameter space:

\[
m_1 \in [10^{-12}, 10^{-2}] \, \text{eV}, \quad M_1 \in [10^3, 10^8] \, \text{GeV}, \quad v_\nu \in [10^{-2}, 10^2] \, \text{GeV},
\]

\[
\text{Re}(\omega_{12,13,23}) \in [10^{-10}, 1], \quad \text{Im}(\omega_{12,13,23}) \in [10^{-10}, 1], \quad \text{BR}_\chi \in [10^{-6}, 10^{-2}].
\]

During the scan, we have fixed \(M_2/M_1 = M_3/M_2 = 10\). The final obtained baryon asymmetry \(Y_{\Delta B}\) is required to be within 3\(\sigma\) range of the observed value, i.e., \(Y_{\Delta B} \in [8.60, 8.84] \times 10^{-11}\). The results are
shown in Fig. 5 and Fig. 6 for DM and leptogenesis, respectively.

Let’s consider the DM results in Fig. 5 first. According to the dominant constraint from free streaming length $r_{FS}$, we can divide the viable samples into three scenarios in Fig. 5(a). Of course, the hot DM scenario is not favored by small structure formation. For warm DM, $m_{\chi} \in [0.3, 2 \times 10^3]$ keV is possible. Meanwhile for cold DM, $m_{\chi} \in [10, 2 \times 10^5]$ keV is allowed. And $r_{FS}$ is down to about $10^{-5}$ Mpc when $m_{\chi} \sim 10^5$ keV. From Fig. 5(b), we aware that the hot DM samples correspond to those with small DM mass $m_{\chi}$ and very weak washout effect $K \lesssim 10^{-2}$. Fig. 5(c) shows the samples in the $m_{\chi} - M_1$ plane. Three kinds of DM are all possible for certain value of $M_1$. By the way, it is interesting to obtain an upper limit on $m_{\chi}$ when $M_1 \lesssim 10^6$ GeV. This indicates that for TeV scale leptogenesis, FIMP DM should be keV
to sub-MeV. The result for $\text{BR}_\chi$ is shown in Fig. 5(d), which tells us that warm DM requires $\text{BR}_\chi \gtrsim 10^{-4}$ and cold DM requires $\text{BR}_\chi \lesssim 10^{-3}$, respectively.

Then we consider the leptogenesis results in Fig. 6. The generalised Davidson-Ibarra bound is clearly seen in Fig. 6(a). The (warm and cold DM) allowed samples show that the mass of $N_1$ for success leptogenesis could be down to about 3 TeV. The viable region in the $v_\nu - M_1$ plane is shown in Fig. 6(b), which is consistent with the theoretical bounds discussed in Ref. [16]. For completeness, the naturalness bound in Eq. (4) is also shown. Therefore, natural leptogenesis is viable for $3 \times 10^3 \text{ GeV} \lesssim M_1 \lesssim 7 \times 10^6 \text{ GeV}$ with $0.4 \text{ GeV} \lesssim v_\nu \lesssim 30 \text{ GeV}$. The result for decay parameter $K$ is given in Fig. 6(c), which shows that $K \lesssim 10$ should be satisfied when $M_1 \lesssim 10^8 \text{ GeV}$. Actually for $M_1 \lesssim 10^5 \text{ GeV}$, all the samples are within weak washout region. An upper bound on lightest neutrino mass $m_1$ is clearly seen in Fig. 6(d). Success leptogenesis in the $\nu 2\text{HDM}$ requires $m_1$ must be extremely tiny, i.e., $m_1 \lesssim 10^{-11} \text{ eV}$ for $M_1 \sim 10^4 \text{ GeV}$.

Before ending this section, we give a brief discussion on the collider signature. According to the results of leptogenesis in Fig. 6, not too small $v_\nu$ is favored. In such scenario, the branching ratios of neutrinophilic scalars are quite different from the scenario with small $v_\nu$ [42, 61–63], but are similar with type-I 2HDM [64]. Currently, if $m_{\Phi_\nu}$ is smaller than $m_t$, the most stringent constraint comes from $t \rightarrow bH^\pm (H^\pm \rightarrow \tau^\pm \nu)$ [65], which could exclude the region $v_\nu \gtrsim 18 \text{ GeV}$ [66]. Meanwhile, if $m_Z + m_h \lesssim m_{\Phi_\nu} \lesssim 2m_t$, the channel $A \rightarrow Zh(h \rightarrow b\bar{b})$ could exclude the region $v_\nu \gtrsim 24 \text{ GeV}$ [67]. For heavier additional scalars with $m_{\Phi_\nu} > 2m_t$, the signature $A/H \rightarrow tt$ is only able to probe the region $v_\nu \gtrsim 174 \text{ GeV}$ [68, 69]. Therefore, the experimental bounds on neutrinophilic scalars can be easily escaped provided $m_{\Phi_\nu}$ is large enough. At HL-LHC, the signature $A \rightarrow Zh(h \rightarrow b\bar{b})$ would reach $v_\nu \sim 10 \text{ GeV}$ [69]. Then the observation of this signature will indicate $M_1 \sim 10^6 \text{ GeV}$ and $m_1 \lesssim 10^{-7} \text{ eV}$.

VI. CONCLUSION

In this paper, we propose an extended $\nu 2\text{HDM}$ to interpret the neutrino mass, leptogenesis and dark matter simultaneously. This model contains one neutrinophilic scalar doublet $\Phi_\nu$, three right hand heavy neutrino $N$, which account for low scale neutrino mass generation similar to type-I seesaw. Leptogenesis is generated due to the CP-violating decays of right hand neutrino $N \rightarrow \ell_L \Phi^*_\nu, \ell_L \Phi_\nu$. The dark sector contains one scalar singlet $\phi$ and one Dirac fermion singlet $\chi$, which are charged under a $Z_2$ symmetry. Provided $m_\chi < m_\phi$ and $\lambda \ll 1$, $\chi$ is a FIMP DM candidate within this paper. The relic abundance of $\chi$ is produced by $N \rightarrow \chi \phi$. Therefore, we have a common origin, i.e., the heavy right hand neutrino $N$, for tiny neutrino mass, baryon asymmetry and dark matter.

In the frame work of $\nu 2\text{HDM}$, the asymmetry $\epsilon_1$ and decay parameter $K$ are both enhanced by the
smallness of $v_\nu$. By explicit calculation, we show that the decay parameter $K$ can be suppressed under certain circumstance. As for FIMP DM, the relic abundance mainly depends on the branching ratio $\text{BR}_\chi$ and decay parameter $K$, and $m_\chi$ is typically at the order of keV to MeV scale. Meanwhile the free streaming length sets stringent bound. The viable parameter space for success leptogenesis and DM is obtained by solving the corresponding Boltzmann equations. To keep this model natural, we find $10^3 \text{ GeV} \lesssim M_1 \lesssim 10^6 \text{ GeV}$, $0.4 \text{ GeV} \lesssim v_\nu \lesssim 30 \text{ GeV}$, $m_1 \lesssim 10^{-5} \text{ eV}$ and $K \lesssim 10$ is favored by leptogenesis. Meanwhile, the warm (cold) DM mass in the range $m_\chi \in [0.3, 2 \times 10^3] \text{ keV}$ ($m_\chi \in [10, 2 \times 10^5] \text{ keV}$) is predicted with $\text{BR}_\chi \gtrsim 10^{-4}$ ($\text{BR}_\chi \lesssim 10^{-3}$).

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[1] Y. Fukuda et al. [Super-Kamiokande Collaboration], Phys. Rev. Lett. 81, 1562 (1998) [hep-ex/9807003].
[2] Q. R. Ahmad et al. [SNO Collaboration], Phys. Rev. Lett. 89, 011301 (2002) [nucl-ex/0204008].
[3] P. Minkowski, Phys. Lett. B 67, 421 (1977).
[4] R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980).
[5] M. Fukugita and T. Yanagida, Phys. Lett. B 174, 45 (1986).
[6] S. Davidson and A. Ibarra, Phys. Lett. B 535, 25 (2002) [hep-ph/0202239].
[7] F. Vissani, Phys. Rev. D 57, 7027 (1998) [hep-ph/9709409].
[8] J. D. Clarke, R. Foot and R. R. Volkas, Phys. Rev. D 91, no. 7, 073009 (2015) [arXiv:1502.01352 [hep-ph]].
[9] A. Pilaftsis and T. E. J. Underwood, Nucl. Phys. B 692, 303 (2004) [hep-ph/0309342].
[10] E. K. Akhmedov, V. A. Rubakov and A. Y. Smirnov, Phys. Rev. Lett. 81, 1359 (1998) [hep-ph/9803255].
[11] T. Asaka and M. Shaposhnikov, Phys. Lett. B 620, 17 (2005) [hep-ph/0505013].
[12] T. Hambye and D. Teresi, Phys. Rev. Lett. 117, no. 9, 091801 (2016) [arXiv:1606.00017 [hep-ph]].
[13] T. Hambye and D. Teresi, Phys. Rev. D 96, no. 1, 015031 (2017) [arXiv:1705.00016 [hep-ph]].
[14] S. Baumholzer, V. Brdar and P. Schwaller, JHEP 1808, 067 (2018) [arXiv:1806.06864 [hep-ph]].
[15] W. Chao and M. J. Ramsey-Musolf, Phys. Rev. D 89, no. 3, 033007 (2014) [arXiv:1212.5709 [hep-ph]].
[16] J. D. Clarke, R. Foot and R. R. Volkas, Phys. Rev. D 92, no. 3, 033006 (2015) [arXiv:1505.05744 [hep-ph]].
[17] E. Ma, Mod. Phys. Lett. A 21, 1777 (2006) doi:10.1142/S0217732306021141 [hep-ph/0605180].
[18] S. Kashiwase and D. Suematsu, Phys. Rev. D 86, 053001 (2012) [arXiv:1207.2594 [hep-ph]].
[19] S. Kashiwase and D. Suematsu, Eur. Phys. J. C 73, 2484 (2013) [arXiv:1301.2087 [hep-ph]].
[20] J. Racker, JCAP 1403, 025 (2014) [arXiv:1308.1840 [hep-ph]].
[21] T. Hugle, M. Platscher and K. Schmitz, Phys. Rev. D 98, no. 2, 023020 (2018) [arXiv:1804.09660 [hep-ph]].
[22] D. Borah, A. Dasgupta and S. K. Kang. arXiv:1806.04689 [hep-ph].
[23] D. Borah, P. S. B. Dev and A. Kumar, Phys. Rev. D 99, no. 5, 055012 (2019) [arXiv:1810.03645 [hep-ph]].
[24] D. Mahanta and D. Borah, arXiv:1912.09726 [hep-ph].
[25] E. Ma, Phys. Rev. Lett. 86, 2502 (2001) [hep-ph/0011121].
[26] N. Haba and O. Seto, Prog. Theor. Phys. 125, 1155 (2011) [arXiv:1102.2889 [hep-ph]].
[27] N. Haba and O. Seto, Phys. Rev. D 84, 103524 (2011) [arXiv:1106.5354 [hep-ph]].
[28] S. Dodelson and L. M. Widrow, Phys. Rev. Lett. 72, 17 (1994) [hep-ph/9303287].
[29] M. Drewes et al., JCAP 1701, 025 (2017) [arXiv:1602.04816 [hep-ph]].
[30] A. Adulpravitchai and M. A. Schmidt, JHEP 1512, 023 (2015) [arXiv:1507.05694 [hep-ph]].
[31] Z. L. Han, B. Zhu, L. Bian and R. Ding, arXiv:1812.00637 [hep-ph].
[32] A. Boyarsky, M. Drewes, T. Lasserre, S. Mertens and O. Ruchayskiy, Prog. Part. Nucl. Phys. 104, 1 (2019) [arXiv:1807.07938 [hep-ph]].
[33] S. Antusch, P. Di Bari, D. A. Jones and S. F. King, Phys. Rev. D 86, 023516 (2012) [arXiv:1107.6002 [hep-ph]].
[34] D. Mahanta and D. Borah, JCAP 1911, no. 11, 021 (2019) [arXiv:1906.03577 [hep-ph]].
[35] M. Chianese, B. Fu and S. F. King, arXiv:1910.12916 [hep-ph].
[36] E. Aprile et al. [XENON Collaboration], Phys. Rev. Lett. 121, no. 11, 111302 (2018) [arXiv:1805.12562 [astro-ph.CO]].
[37] M. Ackermann et al. [Fermi-LAT Collaboration], Phys. Rev. Lett. 115, no. 23, 231301 (2015) [arXiv:1503.02641 [astro-ph.HE]].
[38] N. Bernal, M. Heikinheimo, T. Tenkanen, K. Tuominen and V. Vaskonen, Int. J. Mod. Phys. A 32, no. 27, 1730023 (2017) [arXiv:1706.07442 [hep-ph]].
[39] S. M. Davidson and H. E. Logan, Phys. Rev. D 80, 095008 (2009) [arXiv:0906.3335 [hep-ph]].
[40] T. Morozumi, H. Takata and K. Tamai, Phys. Rev. D 85, no. 5, 055002 (2012) Erratum: [Phys. Rev. D 89, no. 7, 079901 (2014)] [arXiv:1107.1026 [hep-ph]].
[41] N. Haba and T. Horita, Phys. Lett. B 705, 98 (2011) [arXiv:1107.3203 [hep-ph]].
[42] C. Guo, S. Y. Guo, Z. L. Han, B. Li and Y. Liao, JHEP 1704, 065 (2017) [arXiv:1701.02463 [hep-ph]].
[43] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716, 1 (2012) [arXiv:1207.7214 [hep-ex]].
[44] S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716, 30 (2012) [arXiv:1207.7235 [hep-ex]].
[45] P. A. N. Machado, Y. F. Perez, O. Sumensari, Z. Tabrizi and R. Z. Funchal, JHEP 1512, 160 (2015) [arXiv:1507.05750 [hep-ph]].
[46] J. A. Casas and A. Ibarra, Nucl. Phys. B 618, 171 (2001) [hep-ph/0103065].
[47] A. Ibarra and G. G. Ross, Phys. Lett. B 591, 285 (2004) [hep-ph/0312138].
[48] E. Nardi, Y. Nir, E. Roulet and J. Racker, JHEP 0601, 164 (2006) [hep-ph/0601084].
[49] S. Davidson, E. Nardi and Y. Nir, Phys. Rept. 466, 105 (2008) [arXiv:0802.2962 [hep-ph]].
[50] N. Aghanim et al. [Planck Collaboration], arXiv:1807.06209 [astro-ph.CO].
[51] A. Falkowski, J. T. Ruderman and T. Volansky, JHEP 1105, 106 (2011) [arXiv:1101.4936 [hep-ph]].
[52] A. Falkowski, E. Kuflik, N. Levi and T. Volansky, Phys. Rev. D 99, no. 1, 015022 (2019) [arXiv:1712.07652 [hep-ph]].

[53] W. Buchmuller, P. Di Bari and M. Plumacher, Annals Phys. 315, 305 (2005) [hep-ph/0401240].

[54] I. Esteban, M. C. Gonzalez-Garcia, A. Hernandez-Cabezudo, M. Maltoni and T. Schwetz, JHEP 1901, 106 (2019) [arXiv:1811.05487 [hep-ph]].

[55] J. A. Harvey and M. S. Turner, Phys. Rev. D 42, 3344 (1990).

[56] L. J. Hall, K. Jedamzik, J. March-Russell and S. M. West, JHEP 1003 (2010) 080 [arXiv:0911.1120 [hep-ph]].

[57] M. Tanabashi et al. [Particle Data Group], Phys. Rev. D 98, no. 3, 030001 (2018).

[58] P. A. R. Ade et al. [Planck Collaboration], Astron. Astrophys. 594, A13 (2016) [arXiv:1502.01589 [astro-ph.CO]].

[59] A. Berlin and N. Blinov, Phys. Rev. Lett. 120, no. 2, 021801 (2018) [arXiv:1706.07046 [hep-ph]].

[60] A. Merle, V. Niro and D. Schmidt, JCAP 1403, 028 (2014) [arXiv:1306.3996 [hep-ph]].

[61] N. Haba and K. Tsumura, JHEP 1106, 068 (2011) [arXiv:1105.1409 [hep-ph]].

[62] W. Wang and Z. L. Han, Phys. Rev. D 94, no. 5, 053015 (2016) [arXiv:1605.00239 [hep-ph]].

[63] K. Huitu, T. J. Karkkainen, S. Mondal and S. K. Rai, Phys. Rev. D 97, no. 3, 035026 (2018) [arXiv:1712.00338 [hep-ph]].

[64] G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, M. Sher and J. P. Silva, Phys. Rept. 516, 1 (2012) [arXiv:1106.0034 [hep-ph]].

[65] A. M. Sirunyan et al. [CMS Collaboration], JHEP 1907, 142 (2019) [arXiv:1903.04560 [hep-ex]].

[66] P. Sanyal, Eur. Phys. J. C 79, no. 11, 913 (2019) [arXiv:1906.02520 [hep-ph]].

[67] M. Aaboud et al. [ATLAS Collaboration], JHEP 1803, 174 (2018) Erratum: [JHEP 1811, 051 (2018)] [arXiv:1712.06518 [hep-ex]].

[68] M. Aaboud et al. [ATLAS Collaboration], Phys. Rev. Lett. 119, no. 19, 191803 (2017) [arXiv:1707.06025 [hep-ex]].

[69] N. Chen, T. Han, S. Li, S. Su, W. Su and Y. Wu, arXiv:1912.01431 [hep-ph].