Generation of a radiative neutrino mass in the linear seesaw framework, charged lepton flavor violation, and dark matter

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We investigate a model with local $\mathrm{U}(1)_{B-L}$ and discrete Z_2 symmetries where two types of weak isospin singlet neutrinos, vectorlike charged leptons, and exotic scalar fields are introduced. The linear seesaw mechanism is induced at the one-loop level through Yukawa interactions associated with the standard model leptons and exotic fields. We also discuss lepton flavor violation and a muon anomalous dipole magnetic moment induced by the new Yukawa interaction. In addition, our model has dark matter candidate which is the lightest Z_2 odd neutral particle. We calculate the relic density and constraints from direct detection.

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I. INTRODUCTION

Different experiments [1–5] on neutrino oscillation phenomena [6] are consistently giving firm indications of the existence of tiny neutrino mass and flavor mixing. The existence of neutrino mass allows us to extend the Standard Model (SM) which is an essential window for searching for new physics. The simplest idea to extend the SM with an SM singlet right-handed heavy Majorana neutrino was introduced in [7–10]. The heavy right-handed Majorana neutrinos create a lepton number violating mass term "for the light neutrinos" through a dimension five operator which can naturally explain the tiny neutrino masses. This procedure is called the seesaw mechanism. The seesaw scale (the mass scale of the heavy Majorana neutrinos) varies from the electroweak scale to the intermediate scale ($\sim 10^{15}$ GeV) as the neutrino Dirac Yukawa coupling (Y_D) varies from the scale of electron Yukawa coupling ($Y_e \sim 10^{-6}$) up to that of the top quark ($Y_t \sim 1$). If we consider the scale of the seesaw mechanism at the TeV scale or lower, the Dirac Yukawa coupling (Y_D) becomes very small [$\mathcal{O}(10^{-6})$] to produce appropriate light neutrino masses as suggested by neutrino oscillation experiments and cosmological observations.

Apart from the seesaw mechanism there is another type of mechanism where a small lepton number violating term plays a key role in generating the tiny neutrino mass. Such a mechanism is commonly called the canonical inverse seesaw mechanism [11,12]. In this scenario, unlike the seesaw mechanism, the light neutrino mass is not obtained by the suppression of the heavy neutrino mass. Due to the smallness of the lepton number violating parameter, the heavy right-handed neutrinos are pseudo-Dirac in nature. Their Dirac Yukawa couplings with the SM lepton doublets and the SM Higgs doublet could be order one to produce the light neutrino mass.

There is another type of TeV scale seesaw model which is called the linear seesaw [13–20]. This is a simple variation of the canonical inverse seesaw model. In a linear seesaw model we introduce two heavy right-handed SM singlet neutrinos with opposite lepton numbers where four right-handed SM singlet Majorana heavy neutrinos are used as in the canonical inverse seesaw. It has been shown in [18] that from the vacuum metastability bounds the unknown Dirac Yukawa coupling can be constrained. The vacuum stability bounds on the Dirac Yukawa coupling for the canonical type-I frame-work has been studied in [21–23]. In our paper we consider the linear seesaw model where we show the (13) and (31) elements of the neutrino mass matrix are nonzero but (22) and (33) elements are zero; here the elements of neutrino mass matrix are considered in the basis of (ν_L, N_R^C, S_L) where N_R and S_L are SM singlet fermions. Whereas in the inverse seesaw model (13) and (31) elements are zero, the (33) element is nonzero, and (22) may or may not be zero. In our model we generate the (13) and (31) elements of the neutrino mass matrix at the one-loop level and study various features of this model.

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In our model, we apply an extended gauged $B-L$ framework with an additional Z_2 parity where we also introduce a vectorlike charged lepton, two types of weak isospin [which are equivalent to $SU(2)_L$] singlet neutrinos, and new scalar fields. The one-loop induced linear seesaw mechanism is realized by Yukawa couplings associated with SM leptons and new fields. These Yukawa couplings also induce a muon anomalous magnetic dipole moment (muon $g-2$) where current measurement indicates $\Delta a_\mu = a_\mu^{\rm exp} - a_\mu^{\rm SM} = (28.8 \pm 8.0) \times 10^{-10}$ [24], and lepton flavor violating (LFV) processes such as $\ell_i \to \ell_j \gamma$, which are taken as constraints [25]. In addition, the lightest Z_2 odd particle is stable which can be a good candidate of dark matter (DM) if it is neutral [26–28]. Then we discuss relic density and constraint from direct detection for our DM candidate.

The paper is organized as follows. In Sec. II, we introduce our model representing particle contents, new interactions, and a neutrino mass matrix where we have studied neutrino masses and mixing in the light of neutrino experimental data. In Sec. III, we study lepton flavor violation and a muon anomalous dipole magnetic moment. In Sec. IV we analyze dark matter physics in the model. In Sec. V, we give a conclusion.

II. MODEL

In this model we extend the SM with a $U(1)_{B-L}$ gauge group and a discrete Z_2 parity. The relevant part of the particle content has been displayed in Table I. The N_{R_i} is the heavy right-handed Majorana neutrino with three generations to keep the model free from $U(1)_{B-L}$ anomalies. The fermion S_L is also a left-handed Majorana heavy neutrino which has three generations and is neutral under the $U(1)_{B-L}$ gauge group. The iso-singlet charged fermion E is vectorlike with odd Z_2 parity. We also consider that E also has three generations in our model. Notice that the lightest Z_2 odd particle is stable and can be a good DM candidate if it is electrically neutral.

We can write the Lagrangian which is relevant for the neutrino mass matrix at tree level as follows:

TABLE I. The relevant part of the particle content.

	$SU(2)$	$U(1)$	$U(1)_{B-L}$	Z_2
$L_L(\equiv[\nu_L, \ell_L]^T)$	2	$-\frac{1}{2}$	-1	$+$
e_R	1	-1	-1	$+$
$N_{R_{i=1,2,3}}$	1	0	-1	$+$
$S_{L_{j=1,2,3}}$	1	0	0	$+$
$E_{L,R_{a=1,2,3}}$	1	-1	-1	$-$
Φ	2	$\frac{1}{2}$	0	$+$
η	2	$\frac{1}{2}$	0	$-$
χ^-	1	-1	-1	$-$
ϕ	1	0	-1	$+$

$$\mathcal{L}_{\rm int} \supset y_\ell \overline{L_L}\Phi e_R + y_N \overline{L_L}\,\tilde{\Phi} N_R + y_{\rm NS}\overline{N_R} S_L \phi + M_S \overline{S_L^C} S_L + {\rm H.c.}, \tag{2.1}$$

where the first three terms induce the Dirac mass terms after Φ and ϕ getting vacuum expectation value (VEV), and the fourth term with M_S is the lepton number violating Majorana mass term. We use the SM Higgs field Φ as

$$\Phi = \begin{pmatrix} \Phi^+ \\ \Phi_0 \end{pmatrix}, \qquad \Phi^* = \begin{pmatrix} \Phi^{+^*} \\ \Phi_0^* \end{pmatrix},$$
$$\tilde{\Phi} = i\sigma_2 \Phi^* = \begin{pmatrix} \Phi_0^* \\ -\Phi^{+^*} \end{pmatrix}, \quad \Phi^- = \Phi^{+^*} \text{ and } \quad \eta = \begin{pmatrix} \eta^+ \\ \eta_0 \end{pmatrix} \tag{2.2}$$

where neutral components are written by $\Phi_0 \equiv (v+h)/\sqrt{2}$, and $\phi \equiv (v_\phi + \varphi)/\sqrt{2}$, $\eta_0 \equiv \frac{\eta^{\rm Re} + i\eta^{\rm Im}}{\sqrt{2}}$ and $\tilde{\eta} \equiv i\sigma_2 \eta^*$.

After the symmetry breaking, one can write the neutrino mass matrix in Eq. (2.3)

$$\mathcal{L}_{\rm mass} = \begin{pmatrix} \overline{\nu_L^C} & \overline{N_R} & \overline{S_L^C} \end{pmatrix} \begin{pmatrix} 0 & m_D^* & 0 \\ m_D^\dagger & 0 & m_{\rm NS} \\ 0 & m_{\rm NS}^T & M_S \end{pmatrix} \begin{pmatrix} \nu_L \\ N_R^C \\ S_L \end{pmatrix}, \tag{2.3}$$

where Dirac masses can be written by $m_D = y_N \frac{v}{\sqrt{2}}$ and $m_{\rm NS} = y_{\rm NS} v_\phi$. The $B-L$ symmetry forbids the (22) term in the neutrino mass matrix of Eq. (2.3).

At this point it must be pointed out that ϕ is a $B-L$ charged scalar whose VEV is denoted by v_ϕ. The breaking of the electroweak and $B-L$ symmetry is induced spontaneously through the potential:

$$V_1 = m_H^2 \Phi^\dagger\Phi + \frac{\lambda_1}{2}(\Phi^\dagger\Phi)(\Phi^\dagger\Phi) + m_\phi^2 \phi^\dagger\phi + \frac{\lambda_2}{2}(\phi^\dagger\phi)(\phi^\dagger\phi) + \frac{\lambda_{12}}{2}(\Phi^\dagger\Phi)(\phi^\dagger\phi). \tag{2.4}$$

After $\mathrm{U}(1)_{B-L}$ breaking, we have Z' boson whose mass is given by v_ϕ. In our analysis, we assume the Z' boson is sufficiently heavy evading collider constraints. Then the mixing between Z and Z' is essentially given in terms of their masses as

$$\tan\theta_{Z-Z'} \approx \frac{m_Z^2}{m_{Z'}^2} \lesssim \mathcal{O}(10^{-4}) \tag{2.5}$$

that is negligible tiny, where we take $m_{Z'} = 3.5$ TeV. Since Z' does not contribute to neutrino mass and DM physics, we just assume the gauge coupling for Z' is sufficiently small satisfying the current constraint. Thus we will not discuss phenomenology of Z'. There are other two scalars η

and χ^- with odd Z_2 parity. The potential containing η and χ^- can be written as

$$\begin{aligned} V_2 = {} & m_\eta \eta^\dagger \eta + \frac{\lambda_\eta}{2}(\eta^\dagger\eta)(\eta^\dagger\eta) + m_{\chi^-}\chi^{-\dagger}\chi^- \\ & + \frac{\lambda_\chi}{2}(\chi^{-\dagger}\chi^-)(\chi^{-\dagger}\chi^-) + \frac{\lambda_{\eta\chi^-}}{2}(\eta^\dagger\eta)(\chi^{-\dagger}\chi^-) \\ & + \lambda_{\Phi\eta}(\Phi^\dagger\Phi)(\eta^\dagger\eta) + \lambda'_{\Phi\eta}(\Phi^\dagger\eta)(\Phi^\dagger\eta) \\ & + \frac{\lambda''_{\Phi\eta}}{2}[(\Phi^\dagger\eta)^2 + \text{c.c.}] + \lambda_{\eta\phi}(\eta^\dagger\eta)(\phi^\dagger\phi) \\ & + \mu(\Phi^T.\eta)\chi^-\phi, \end{aligned} \tag{2.6}$$

which is invariant under $SU(2)\times U(1)\times U(1)_{B-L}$ and the Z_2 symmetry.[1] Therefore the complete potential of our system will be given as

$$V_{\rm sys} = V_1 + V_2. \tag{2.7}$$

Using seesaw approximation, from Eq. (2.3) we can write the effective 3×3 light neutrino mass matrix as

$$m_\nu = (m_D^* m_{\rm NS}^{-1}) M_s (m_D^* m_{\rm NS}^{-1})^T. \tag{2.8}$$

Note that the light neutrino mass is directly proportional to the M_s. Therefore the degree of smallness regulates the smallness of the light neutrino mass and if $M_s \to 0$, the light neutrino becomes massless, which is the inverse seesaw scenario [11,12], and the Feynman diagram of the inverse seesaw operator is given in Fig. 1. If there is a nonzero (22) term in [29–31] in the neutrino mass matrix which provides a nontrivial contribution to light neutrino masses at the one-loop level which does not vanish in the limit (33) term (M_S) going to zero. However, at tree level the light neutrino masses go to zero in the limit $M_S \to 0$, even if $M_R \neq 0$.

There is another possibility to obtain the light neutrino mass through switching on the 31 term in the mass matrix in Eq. (2.3). This can restore the small neutrino mass even if we have a vanishing M_s. Here vanishing M_s can be justified by assigning a charge of some global symmetry to S_L, ϕ and χ^- as -1, 1, and -1, for example, where only the M_s term explicitly breaks the charge conservation in our model. In that case we can interpret that the M_s term softly breaks the symmetry and it is natural to take a small value for the M_s. However in our model it is not possible to generate the mass term at tree level because $U(1)_{B-L}$ symmetry forbids us from writing the terms like $\overline{N_R^C} N_R \phi^*$, $\overline{L_L^C}\Phi^* S_L$, and $\overline{e_R^C} S_L \chi^+$ where the first and the second terms respectively induce 22 and 13(31) terms of the neutrino mass matrix while the third term would contribute to a LFV

[1]$\mu(\Phi^T.\eta)\chi^-\phi = \mu(\Phi^T\epsilon\eta)\chi^-\phi = \mu(\Phi^+\,\Phi^0)\begin{pmatrix}0 & 1\\ -1 & 0\end{pmatrix}\begin{pmatrix}\eta^+\\ \eta^0\end{pmatrix}\times(\chi^-\phi) = \mu(\Phi^+\eta^0\chi^-\phi - \Phi^0\eta^+\chi^-\phi) \sim -\mu\Phi^0\eta^+\chi^-\phi$. After symmetry breaking, $-\mu\Phi^0\eta^+\chi^-\phi \supset -\mu\frac{v_\phi v}{\sqrt{2}}\eta^+\chi^-$.

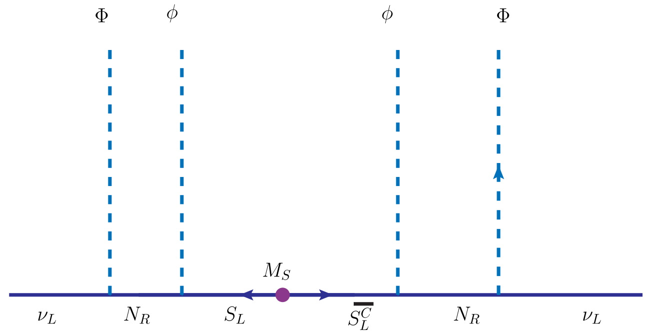

FIG. 1. Feynman diagram for the inverse seesaw.

process. Although some terms in the neutrino mass matrix are forbidden at tree level, our particle content in Table I allows us to write the Dirac mass term of E and the gauge invariant Yukawa terms which can generate the 13 (or 31) term of the neutrino mass matrix through a one-loop diagram;

$$\mathcal{L} \supset (y_1)_{i\alpha}\overline{L_{L_i}}\eta E_{R_\alpha} + (y_2)_{\alpha j}\overline{E_{L_\alpha}} S^C_{L_j}\chi^- + M_E \overline{E_{L_\alpha}} E_{R_\alpha}, \tag{2.9}$$

where α and j are the generation index of the fermions E and S respectively. The third term of Eq. (2.9) is a Dirac mass term of E and will contribute in the neutrino mass generation at one-loop level. After generating the 31 (or 13) term radiatively we can write the neutrino mass matrix[2]

$$m_\nu^{\rm tree+1-loop} = \begin{pmatrix} 0 & m_D^* & \delta_1^* \\ m_D^\dagger & 0 & m_{\rm NS} \\ \delta_1^\dagger & m_{\rm NS}^T & M_S \end{pmatrix}. \tag{2.10}$$

Using the seesaw approximation, from Eq. (2.10) we can write the effective 3×3 light neutrino mass matrix as

$$\begin{aligned} (m_\nu^{\rm light})^{\rm tree+1-loop} & \\ = {} & (\,m_D^* \quad \delta_1^*\,)\begin{pmatrix} 0 & m_{\rm NS} \\ m_{\rm NS}^T & M_s \end{pmatrix}^{-1}\begin{pmatrix} m_D^\dagger \\ \delta_1^\dagger \end{pmatrix} \\ = {} & (\,m_D^* \quad \delta_1^*\,)\left(\frac{-1}{m_{\rm NS} m_{\rm NS}^T}\right)\begin{pmatrix} M_s & -m_{\rm NS} \\ -m_{\rm NS}^T & 0 \end{pmatrix}\begin{pmatrix} m_D^\dagger \\ \delta_1^\dagger \end{pmatrix} \\ = {} & -(m_D^* m_{\rm NS}^{-1}) M_s (m_D^* m_{\rm NS}^{-1})^T \\ & + (m_{\rm NS}^T)^{-1} m_D^* \delta_1^\dagger + \delta_1^* m_D^\dagger (m_{\rm NS}^{-1}). \end{aligned} \tag{2.11}$$

Therefore the vanishing limit of M_s in Eq. (2.11) will switch off the tree-level mass term and the light neutrino

[2]It must be mentioned that in Eq. (2.3) we have three generations of ν_L, three generations of N_R, and S_L which makes the Dirac mass matrix, m_D as a 3×3 matrix as Y_N is carrying the flavors. The same structure is for Eq. (2.10) where δ_1 is a 3×3 matrix keeping the other matrices the same and the total mass matrix has a 9×9 structure.

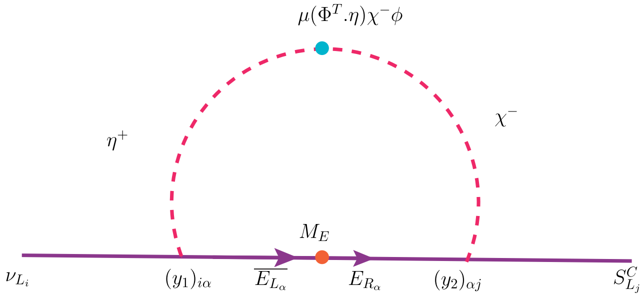

FIG. 2. Feynman diagram of the radiative loop to generate the 13 (or 31) mass term in the neutrino mass matrix in Eq. (2.10) to induce the linear seesaw.

mass term will be generated only from the one-loop term leading to

$$m_\nu^{\mathrm{light}^{1-\mathrm{loop}}} = (m_{\mathrm{NS}}^T)^{-1} m_D^* \delta_1^\dagger + \delta_1^* m_D^\dagger (m_{\mathrm{NS}}^{-1}). \tag{2.12}$$

Therefore, we can resolve the light neutrino mass through the radiative one-loop process in the linear seesaw mechanism. The one-loop diagram in Fig. 2 shows the radiative mass term for the (13) and (31) elements in the neutrino mass matrix. Now solving the diagram, we can calculate the value of δ_1. To do this we first rotate the charged scalar sector using an arbitrary orthogonal matrix

$$\begin{pmatrix} \chi^- \\ \eta^- \end{pmatrix} = \begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} H_1^- \\ H_2^- \end{pmatrix}. \tag{2.13}$$

From Fig. 2 and using Eq. (2.13) we write

$$\begin{aligned} (y_1)_{i\alpha}\overline{\nu_{L_i}}\eta^+ E_{R_\alpha} &= (y_1)_{i\alpha}\overline{\nu_{L_i}} E_{R_\alpha}(\sin\theta H_1^+ + \cos\theta H_2^+), \\ (y_2)_{\alpha j}\overline{E_{L_\alpha}} S^C_{L_j}\chi^- &= (y_2)_{\alpha j} E_{L_\alpha} S^C_{L_\alpha}(\cos\theta H_1^- - \sin\theta H_2^-), \end{aligned} \tag{2.14}$$

$$\begin{aligned} -i(\delta_1)_{ij} &= \int \frac{d^4k}{(2\pi)^4}(-iy_2)_{\alpha j} P_R \frac{i(\not{k} + M_{E_\alpha})}{(k^2 - M_{E\alpha}^2)}(-iy_1)_{i\alpha} P_R \sin\theta\cos\theta\left(\frac{i}{k^2 - m_{H_1}^2} - \frac{i}{k^2 - m_{H_2}^2}\right) \\ &= \frac{-i\sin\theta\cos\theta}{(4\pi)^2}(m_{H_1}^2 - m_{H_2}^2)\sum_\alpha (y_1)_{i\alpha} M_{E_\alpha}(y_2)_{\alpha j} \\ &\quad \times \int_0^1 dx \int_0^{1-x} dy \left[\frac{1}{xM_{E_\alpha}^2 + ym_{H_1}^2 + (1-x-y)m_{H_2}^2}\right] \\ &= \frac{-i\sin\theta\cos\theta}{(4\pi)^2}(m_{H_1}^2 - m_{H_2}^2)\sum_\alpha (y_1)_{i\alpha} M_{E_\alpha}(y_2)_{\alpha j} \\ &\quad \times \left[\frac{2\left(M_{E_\alpha}^2 m_{H_1}^2 \ln\left[\frac{M_{E_\alpha}}{m_{H_1}}\right] + m_{H_2}^2\left(m_{H_1}^2 \ln\left[\frac{m_{H_1}}{m_{H_2}}\right] + M_{E_\alpha}^2 \ln\left[\frac{m_{H_2}}{M_{E_\alpha}}\right]\right)\right)}{(M_{E_\alpha} - m_{H_1})(M_{E_\alpha} + m_{H_1})(M_{E_\alpha} - m_{H_2})(M_{E_\alpha} + m_{H_2})(m_{H_1}^2 - m_{H_2}^2)}\right] \\ &= -i\frac{\sin 2\theta}{16\pi^2}\sum_\alpha (y_1)_{i\alpha}(y_2)_{\alpha j} M_{E_\alpha}\left[\frac{\left(M_{E_\alpha}^2 m_{H_1}^2 \ln\left[\frac{M_{E_\alpha}}{m_{H_1}}\right] + m_{H_2}^2\left(m_{H_1}^2 \ln\left[\frac{m_{H_1}}{m_{H_2}}\right] + M_{E_\alpha}^2 \ln\left[\frac{m_{H_2}}{M_{E_\alpha}}\right]\right)\right)}{(M_{E_\alpha}^2 - m_{H_1}^2)(M_{E_\alpha}^2 - m_{H_2}^2)}\right], \end{aligned} \tag{2.15}$$

where we have assumed $M_{E_\alpha} \neq m_{H_1} \neq m_{H_2}$, and $m_{H_{1(2)}}$ is defined as the mass of the singly charged boson of $H^\pm_{1(2)}$. When we take $M_{E_\alpha} \gg m_{H_k}$ $(k = 1, 2)$ the typical size of δ_1 is approximately given by

$$(\delta_1)_{ij} \sim \frac{\sin 2\theta}{16\pi^2}\sum_\alpha\sum_k (y_1)_{i\alpha}(y_2)_{\alpha j}\frac{m_{H_i}}{M_{E_\alpha}} m_{H_k}, \tag{2.16}$$

where the $\ln[m_{H_k}/M_{E_\alpha}]$ factor is omitted here.

Depending on the mass scales and the scales of the Yukawa couplings, one can justify the degree of smallness of the mass term δ_1 so as to reproduce the light neutrino masses at the correct scale.

A. Neutrino data

In this analysis we assume that $(m_{\mathrm{NS}}^T)^{-1} m_D^* \ll 1$ which allows us to express the flavor eigenstates (ν) of the light Majorana neutrinos in terms of the mass eigenstates of the light (ν_m) and heavy (N_m) Majorana neutrinos where

$$\nu \sim \mathcal{N}\nu_m + \mathcal{R}N_m. \tag{2.17}$$

For simplicity we may consider δ, m_D, and m_{NS} are real quantities. Here

$$\mathcal{R} = m_D m_{\mathrm{NS}}^{-1}, \qquad \mathcal{N} = \left(1 - \frac{1}{2}\epsilon\right)U_{\mathrm{PMNS}}, \qquad \epsilon = \mathcal{R}^*\mathcal{R}^T \tag{2.18}$$

and $U_{\rm PMNS}$ is the usual neutrino mixing matrices which can diagonalize m_ν in the following way:

$$U^T_{\rm PMNS} m_\nu U_{\rm PMNS} = {\rm diag}(m_1, m_2, m_3). \tag{2.19}$$

Due to the presence of ϵ, the mixing matrix $\mathcal{N}$ is nonunitary. For simplicity we consider that there are three degenerate heavy neutrinos.

We consider a situation where the Dirac mass term carries the flavor, where as the δ term is proportional to unity. Therefore

$$m_\nu = \frac{m_D}{m_{\rm NS}}\delta + \delta\frac{m_D}{m_{\rm NS}} = 2\delta\frac{m_D}{m_{\rm NS}} = 2\delta\mathcal{R} = U^*_{\rm PMNS} D_{\rm NH/IH} U^\dagger_{\rm PMNS}, \tag{2.20}$$

$$\mathcal{R} = \frac{1}{2\delta} U^*_{\rm PMNS} D_{\rm NH/IH} U^\dagger_{\rm PMNS}, \tag{2.21}$$

where NH(IH) represents the shorthand symbol for "normal (inverted) hierarchy." Using the neutrino oscillation data [24,32] $\sin^2 2\theta_{13} = 0.092$, $\sin^2 2\theta_{12} = 0.87$, $\sin^2 2\theta_{23} = 1.0$, $\Delta m^2_{\rm sol} = 7.6\times 10^{-5}$ eV^2 and $\Delta m^2_{\rm atm} = 2.4\times 10^{-3}$ eV^2 we can write

$$D_{NH} = \begin{pmatrix} \sqrt{0.1*\Delta m^2_{12}} & 0 & 0 \\ 0 & \sqrt{\Delta m^2_{12}} & 0 \\ 0 & 0 & \sqrt{\Delta m^2_{12} + \Delta m^2_{23}} \end{pmatrix} \tag{2.22}$$

and

$$D_{IH} = \begin{pmatrix} \sqrt{\Delta m^2_{23} - \Delta m^2_{12}} & 0 & 0 \\ 0 & \sqrt{\Delta m^2_{23}} & 0 \\ 0 & 0 & \sqrt{0.1*\Delta m^2_{23}} \end{pmatrix} \tag{2.23}$$

respectively. We have expressed D_{NH} in Eq. (2.22) in terms of $m_2^2 - m_1^2 = \Delta m^2_{12} = 0.9*\Delta m^2_{\rm sol}$ and $m_2^2 - m_3^2 = -\Delta m^2_{23} = -\Delta m^2_{\rm atm}$ whereas D_{IH} in Eq. (2.23) has been expressed in terms of $m_2^2 - m_1^2 = \Delta m^2_{12} = \Delta m^2_{\rm sol}$ and $m_2^2 - m_3^2 = \Delta m^2_{23} = 0.9*\Delta m^2_{\rm atm}$. Without the loss of generality we can also replace the least eigenvalues by zero for the NH and IH cases, however, the choices of the smallness of these values do not affect the smallness of δ_1.

Therefore

$$\mathcal{R}^*\mathcal{R}^T = \frac{1}{4\delta^2} U_{\rm PMNS} D_{\rm NH/IH} U^T_{\rm PMNS} U^*_{\rm PMNS} D_{\rm NH/IH} U^\dagger_{\rm PMNS}. \tag{2.24}$$

Using the updated result of the nonunitarity matrix from the LFV bounds we can write $\mathcal{N}\mathcal{N}^\dagger \sim \mathbf{1} - \epsilon$. Due to its nonunitarity, the elements of the mixing matrix $\mathcal{N}$ are severely constrained by the combined data from the neutrino oscillation experiments, the precision measurements of weak boson decays, and the lepton-flavor-violating decays of charged leptons [33–37]. We update the results by using more recent data on the lepton-favor-violating decays [38–40]:

$$|\mathcal{N}\mathcal{N}^\dagger| = \begin{pmatrix} 0.994 \pm 0.00625 & 1.288\times 10^{-5} & 8.76356\times 10^{-3} \\ 1.288\times 10^{-5} & 0.995\pm 0.00625 & 1.046\times 10^{-2} \\ 8.76356\times 10^{-3} & 1.046\times 10^{-2} & 0.995\pm 0.00625 \end{pmatrix}. \tag{2.25}$$

Since $\mathcal{N}\mathcal{N}^\dagger \simeq \mathbf{1} - \epsilon$, we have the constraints on ϵ such that

$$|\epsilon| = \begin{pmatrix} 0.006 \pm 0.00625 & < 1.288\times 10^{-5} & < 8.76356\times 10^{-3} \\ < 1.288\times 10^{-5} & 0.005\pm 0.00625 & < 1.046\times 10^{-2} \\ < 8.76356\times 10^{-3} & < 1.046\times 10^{-2} & 0.005\pm 0.00625 \end{pmatrix}. \tag{2.26}$$

The most stringent bound is given by the (12)-element which is from the constraint on the lepton-flavor-violating muon decay $\mu \to e\gamma$. Using these bounds we can find the minimum value of δ_1 as $\delta_{1_{\rm min}} \sim \mathcal{O}(10\ {\rm eV})$.

III. CHARGED LEPTON FLAVOR VIOLATION

In our model the fermion E_α and the scalar η is involved in the charged lepton flavor violation (cLFV) processes through the interaction

$$\mathcal{L}_{\rm int} \supset \overline{E_{R_\alpha}}(y_1)^\dagger_{\alpha i} L_{L_i}\tilde{\eta} \supset \frac{(y_{1_{i\alpha}})^\dagger}{\sqrt{2}} \overline{E_{R_\alpha}} \ell_{L_i}(\eta^{\rm Re} - i\eta^{\rm Im}), \tag{3.1}$$

where $\tilde{\eta} \equiv i\sigma_2\eta^*$. The Feynman diagram for the corresponding $\ell_i \to \ell_j\gamma$ process(es) are given in Fig. 3. The scattering amplitude for Fig. 3 is given as[3]

$$\begin{aligned} iM &= \int \frac{d^4k}{(2\pi)^4}\overline{u(p_2)}\left[-\frac{-i(y_1^\dagger)_{j\alpha}P_R}{\sqrt{2}}\right]\frac{i(\not{k}+\not{p}1)-\not{p}3+M_{E_\alpha})}{(k+p_1+p_3)^2-M_{E_\alpha}{}^2}(ie\gamma_\mu)\frac{i(\not{k}+\not{p}1)+M_{E_\alpha})}{(k+p_1)^2-M_{E_\alpha}{}^2} \\ &\quad\times\left[-\frac{-i(y_1^\dagger)_{\alpha i}P_L}{\sqrt{2}}\right]u(p_1)\frac{1}{k^2-m_\eta^2}\epsilon(p_3)^\mu \\ &= i(2ep_1.\epsilon^*)\overline{u(p_2)}[a_RP_R+a_LP_L]u(p_1), \end{aligned} \tag{3.2}$$

where m_{η^k}

$$\begin{aligned} (a_R)_{ji} &= -\sum_{\alpha=1}^{3}\sum_{k}^{\mathrm{Re,Im}}\frac{(y_1)_{j\alpha}(y_1^\dagger)_{\alpha i}}{2(4\pi)^2}m_{\ell i}\int dxdydz\frac{yz\delta(x+y+z-1)}{(x+y)M_{E_\alpha}{}^2+zm_{\eta^k}^2} \\ &= -\sum_{\alpha=1}^{3}\sum_{k}^{\mathrm{Re,Im}}\frac{(y_1)_{j\alpha}(y_1^\dagger)_{\alpha i}}{2(4\pi)^2}m_{\ell i}\left[\frac{M_{E_\alpha}{}^6-6M_{E_\alpha}{}^4m_{\eta^k}^2+3M_{E_\alpha}{}^2m_{\eta^k}^4+2m_{\eta^k}^6+12M_{E_\alpha}{}^2m_{\eta^k}^4\ln\left[\frac{M_\alpha}{m_{\eta^k}}\right]}{12(M_{E_\alpha}^2-m_{\eta^k}^2)^4}\right], \end{aligned} \tag{3.3}$$

$$\begin{aligned} (a_L)_{ji} &= -\sum_{\alpha=1}^{3}\sum_{k}^{\mathrm{Re,Im}}\frac{(y_1)_{j\alpha}(y_1^\dagger)_{\alpha i}}{2(4\pi)^2}m_{\ell j}\int dxdydz\frac{xz\delta(x+y+z-1)}{(x+y)M_{E_\alpha}{}^2+zm_{\eta^k}^2} \\ &= -\sum_{\alpha=1}^{3}\sum_{k}^{\mathrm{Re,Im}}\frac{(y_1)_{j\alpha}(y_1^\dagger)_{\alpha i}}{2(4\pi)^2}m_{\ell j}\left[\frac{M_{E_\alpha}{}^6-6M_{E_\alpha}{}^4m_{\eta^k}^2+3M_{E_\alpha}{}^2m_{\eta^k}^4+2m_{\eta^k}^6+12M_{E_\alpha}{}^2m_{\eta}^4\ln\left[\frac{M_{E_\alpha}}{m_{\eta^k}}\right]}{12(M_{E_\alpha}^2-m_{\eta^k}^2)^4}\right]. \end{aligned} \tag{3.4}$$

Now

$$\mathcal{BR}(\ell_i\to\ell_j\gamma)\sim\frac{48\pi^3\alpha_{em}C_{ij}}{G_F^2m_{\ell i}^2}(|a_L^{\eta^{\mathrm{Re}}}+a_L^{\eta^{\mathrm{Im}}}|^2+|a_R^{\eta^{\mathrm{Re}}}+a_R^{\eta^{\mathrm{Im}}}|^2)_{ji}, \tag{3.5}$$

where $\alpha_{em}\approx 1/137$ is the fine structure constant, $G_F\approx 1.17\times 10^{-5}\ \mathrm{GeV}^{-2}$ is the Fermi constant, and C_{ij} is defined by

$$\begin{aligned} C_{ij} &\approx 1 \quad \text{for } (i,j)=(\mu,e) \\ &\approx 0.1784 \quad \text{for } (i,j)=(\tau,e) \\ &\approx 0.1736 \quad \text{for } (i,j)=(\tau,\mu). \end{aligned} \tag{3.6}$$

The current experimental bound on $\mathcal{BR}(\ell_i\to\ell_j\gamma)$ is respectively given by [41,42] at 90% C.L.

$$\mathcal{BR}(\ell_\mu\to\ell_e\gamma)\lesssim 4.2\times 10^{-13},\qquad \mathcal{BR}(\ell_\tau\to\ell_e\gamma)\lesssim 3.3\times 10^{-8},\qquad \mathcal{BR}(\ell_\tau\to\ell_\mu\gamma)\lesssim 4.4\times 10^{-8}. \tag{3.7}$$

We can avoid the constraints by choosing the Yukawa coupling y_1 so that off-diagonal elements of $a_{L(R)}$ are sufficiently small.

The diagram in Fig. 3 also contributes to the muon anomalous magnetic moment Δa_μ when $i=j=2$, and it is given by

$$\Delta a_\mu = -m_\mu[a_L^{\eta^{\mathrm{Re}}}+a_R^{\eta^{\mathrm{Re}}}+a_L^{\eta^{\mathrm{Im}}}+a_R^{\eta^{\mathrm{Im}}}]_{22}, \tag{3.8}$$

including the real and imaginary parts of the neutral scalar η. The current experiments [43–45] report that its deviation is $(28.8\pm 8.0)\times 10^{-10}$. Taking $M_{E_\alpha} > m_{\eta^k}$, we roughly obtain $\Delta a_\mu \sim \sum_\alpha (y_1)_{2\alpha}(y_1^\dagger)_{\alpha 2}(m_\mu/M_{E_\alpha})^2/(96\pi^2)$. Thus we find that the product of the Yukawa coupling $\sum_\alpha (y_1)_{2\alpha}(y_1^\dagger)_{\alpha 2}$ should be order one or larger to obtain sizable Δa_μ. In addition, exotic particles are preferred not to be too heavy as $\mathcal{O}(1)$ TeV for getting sizable muon $g-2$.

[3]In our convention $Q_E=-1$.

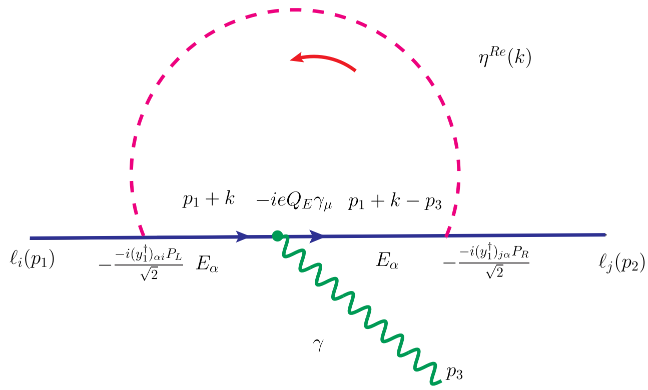

FIG. 3. Feynman diagram for the charged lepton flavor violation processes $\ell_i \to \ell_j\gamma$.

IV. DARK MATTER SCENARIO

A neutral component of η can be a dark matter candidate. Here we assume the real part to be DM: $\eta_R \equiv X$. A general analysis has been done by the authors of Ref. [46], where the DM mass is greater than the mass of the W boson.[4] We are interested in lower range $M_X \leq m_W$ since it is preferred to obtain a sizable muon $g-2$, and thus we focus on this range. Also we note that annihilation modes from the Higgs portal is subdominant when we are required to evade the direct detection constraint such as in the LUX experiment which is discussed below. Under this situation, the dominant mode comes from the same Yukawa coupling as Eq. (3.1), which gives a d-wave dominance in the limit of the massless final state. The interaction Lagrangian is again given by

$$\mathcal{L}_{\rm int} \supset \frac{(y_1)_{i\alpha}}{\sqrt{2}}\overline{\ell_{L_i}}E_{R_\alpha}(\eta^{\rm Re}+i\eta^{\rm Im}) \supset \frac{(y_1)_{i\alpha}}{\sqrt{2}}\overline{\ell_{L_i}}E_{R_\alpha}X. \quad (4.1)$$

The relevant Feynman diagrams for the DM annihilation are given in Fig. 4. Then the nonrelativistic cross section to explain the relic density of DM is obtained by

$$\sigma v_{\rm rel} \approx \sum_{i,j=1}^{3}\sum_{\alpha=1}^{3}\frac{|(y_1)_{i\alpha}(y_1^\dagger)_{\alpha j}|^2 M_X^6}{240\pi(M_{E_\alpha}^2+M_X^2)^4}v_{\rm rel}^4 \equiv d_{\rm eff}v_{\rm rel}^4. \quad (4.2)$$

Here we apply the relative velocity expansion approximation as follows:

$$\Omega h^2 \approx \frac{10.7\times 10^9\ [{\rm GeV}^{-1}]x_f^3}{20\sqrt{g_*}M_P d_{\rm eff}}, \quad (4.3)$$

where $M_P \approx 1.22\times 10^{19}$ GeV is the Planck mass, $g_* \approx 100$ is the total number of effective relativistic degrees of freedom at the time of freeze-out, and $x_f \approx 25$ is defined by M_X/T_f at the freeze-out temperature (T_f), and $d_{\rm eff}$ is the contribution to the d-wave. We find that $\sum|(y_1)_{i\alpha}(y_1^\dagger)_{\alpha j}|^2$ should be sizable to obtain observed relic density. Note also that even if y_1 is large, we can obtain a small scale of δ_1 in Eq. (2.16) by small values of y_2 and θ.

[4] In this case, DM mass should be greater than 500 GeV, and coannihilation should also be taken into consideration because of an oblique parameter.

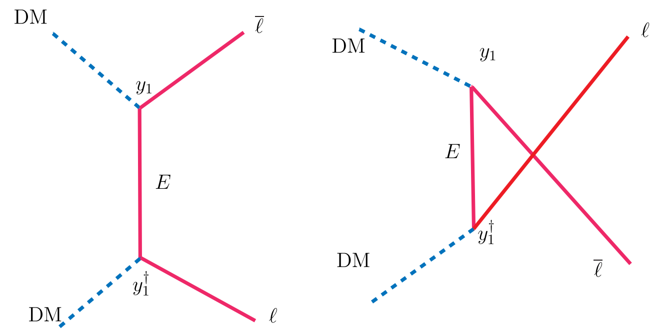

FIG. 4. Feynman diagram for DM annihilation in $s(t)$-channel (a) and u-channel (b).

A spin-independent scattering cross section can be found via the Higgs portal. The relevant terms in the Higgs potential is given in the second line of the right-hand side in Eq. (2.6). Then the CP-even Higgs mixing in the basis of (φ, h) is given by

$$\begin{pmatrix}\varphi\\ h\end{pmatrix} = \begin{pmatrix}\cos\alpha & -\sin\alpha\\ \sin\alpha & \cos\alpha\end{pmatrix}\begin{pmatrix}H_1^0\\ H_2^0\end{pmatrix}, \quad (4.4)$$

where H_2^0 is the SM Higgs and its mass $m_{H_2^0} \approx 125$ GeV, and H_1^0 is another neutral Higgs with vacuum expectation value as v'. Then its formula is given by

$$\sigma_N \approx \frac{m_N^4}{4(m_N+M_X)^2\pi}\left(\frac{C_{2XH_1^0}s_\alpha}{m_{H_1^0}^2}+\frac{C_{2XH_2}c_\alpha}{m_{H_2^0}^2}\right)^2 \times 3.29\times 10^{-29}\ {\rm cm}^2,$$
$$C_{2XH_1^0} = (\lambda_{\Phi\eta}+\lambda'_{\Phi\eta}+\lambda''_{\Phi\eta})s_\alpha + \lambda_{\eta\phi}c_\alpha\frac{v_\phi}{v},$$
$$C_{2XH_2^0} = (\lambda_{\Phi\eta}+\lambda'_{\Phi\eta}+\lambda''_{\Phi\eta})c_\alpha - \lambda_{\eta\phi}s_\alpha\frac{v_\phi}{v}, \quad (4.5)$$

where $m_N \approx 0.939$ GeV is the neutron mass. Here we give a brief estimation, where we simply fix several parameters as $\lambda \equiv \lambda_{\Phi\eta} \approx \lambda'_{\Phi\eta} \approx \lambda''_{\Phi\eta} \approx \lambda_{\eta\phi}$ and $m_{H^0} \equiv m_{H_1^0} \approx m_{H_2^0} =$ 125 GeV. Then the resulting cross section is simplified as

$$\sigma_N \approx \frac{9\lambda^2 m_N^4}{4\pi(m_N+M_X)^2 m_{H^0}^4}\times 3.29\times 10^{-29}\ {\rm cm}^2. \quad (4.6)$$

Notice here that it does not depend on s_α, v, and v'. The stringent cross section is found to be $\sigma_N \approx 2.2\times 10^{-46}$ cm^2 at $M_X \approx 50$ GeV reported by LUX experiment [47], which is supported by CoGENT [48] and CREST [49], although their results are more relaxed. Therefore, in our case, the bound on λ is found to be

$$\lambda \lesssim 0.022. \tag{4.7}$$

Here we discuss order estimation to fit the experimental values such as relic density of DM and muon $g-2$ satisfying LFVs, where notice here that the crucial parameter is y_1 and we do not need to include the neutrino sector because of a lot of independent parameters. First of all, the correct relic density can be achieved by taking $|(y_1)_{i\alpha}(y_1^\dagger)_{\alpha j}|$ to be order one, where we expect all the scales of exotic masses are of the order of 100–1000 GeV. Also sizable muon $g-2$ is achieved if we take $|(y_1)_{21}|^2 + |(y_1)_{22}|^2 + |(y_1)_{23}|^2$ to be order one, while LFVs restrict some components of y_1. For example, the most stringent constraint arises from $\mu \to e\gamma$, and its Yukawa combination $(y_1)_{11}(y_1)^*_{21} + (y_1)_{12}(y_1)^*_{22} + (y_1)_{13}(y_1)^*_{23}$ should be taken to be order $\mathcal{O}(10^{-4})$ to satisfy this bound, where we respectively take the one-loop function and the mediated fields to be order one and 500 GeV. Comparing these three combinations, one finds that there are allowed regions by controlling each component of y_1.

Before closing this section we discuss Z_2 odd particle production at the LHC. The vectorlike charged leptons E_α can be produced via electroweak process $pp \to Z/\gamma \to E^+E^-$ or Z' exchange in s-channel $pp \to Z' \to E^+E^-$ where we assume Z' coupling is small and the electroweak process is dominant. Then E_α decays into charged lepton and DM via Yukawa interaction as $E^\pm \to \ell^\pm X$. We thus expect charged lepton plus missing energy signal at the LHC. Thus our E_α production signal is smiler to that of electroweak production of sleptons in supersymmetric models and we can roughly obtain mass limit as $M_E > 500$ GeV from current slepton searches [50]. Note that the mass limit for our exotic charged scalar boson will be less constrained or similar to that of $E^\pm$; the production cross section of the charged scalar $\eta^\pm$ and $\chi^\pm$ are similar to that of $E^\pm$ while they decay as $\eta^\pm \to E^{\pm(*)}\nu \to \ell^\pm X\nu$ or $\eta^\pm \to W^\pm X$ and $\chi^\pm \to E^{\pm(*)}S \to \ell^\pm XS$ ($E^{\pm*}$ is an off-shell state, and depending upon the masses, $E^\pm$ can be on-shell, too), which give more particles in final states compared to the $E^\pm$ case and the significance of finding charged scalar would be reduced. In Table II we summarize the E_α pair production cross section of $pp \to Z/\gamma \to E^+E^-$ for some benchmark values of M_E which are calculated by CalcHEP [51] with $\sqrt{s} = 13$ TeV. Therefore we expect more than 10 events for integrated luminosity 300 fb^{-1} for $M_E \lesssim 1$ TeV. More detailed analysis including simulation study is beyond the scope of this paper and will be done elsewhere.

TABLE II. The cross sections of $pp \to Z/\gamma \to E^+E^-$ for some benchmark values of M_E.

M_E [GeV]	750	1000	1250	1500
$\sigma(pp \to Z/\gamma \to E^+E^-)$ [fb]	0.3	0.064	0.016	0.0044

V. CONCLUSION

In this paper, we have proposed an extension of the SM with local $\mathrm{U}(1)_{B-L}$ symmetry and discrete Z_2 symmetry where exotic leptons and scalar particles are introduced. In particular, two types of weak isospin singlet neutrinos, N_{R_i} and S_{L_i} are introduced.

Since N_{R_i} is charged under the $\mathrm{U}(1)_{B-L}$, it has to have three generations, due to the anomaly cancelation, while S_{L_i} does not have a $B-L$ charge that suggests that the number of flavors for S_L can be arbitrary. Thus we assume to be three generations of S_L for simplicity. The model induces the linear seesaw mechanism through a one-loop diagram in which Z_2 odd particles propagate, if Majorana mass of S_{L_i} is suppressed. In addition, the lightest Z_2 odd neutral particle can be a good DM candidate.

We have shown a formula for the component of the neutrino mass matrix δ_1 which is generated by one-loop diagram. Then the neutrino mass matrix is given by δ_1 and the Dirac mass parameters in our neutrino sector through linear seesaw mechanism. To fit the neutrino oscillation data, the order of δ_1 is required to be $\delta_1 \gtrsim \mathcal{O}(10\ \mathrm{eV})$ which can easily be realized choosing the values of relevant parameters in the formula. We have also derived formulas of muon $g-2$ and lepton flavor violating decay $\ell \to \ell'\gamma$ at the one-loop level. Furthermore, relic density of DM and DM-nucleon scattering were discussed assuming a neutral component of the inert doublet scalar is a dark matter candidate. We then found that our model can accommodate with neutrino oscillation data via linear seesaw mechanism, sizable muon $g-2$, and the relic density of DM, satisfying the constraints from lepton flavor violations and the direct detection experiment of DM.

Such a model can also be tested at the collider. A small value of δ_1 ensures a sizable mixing between the SM light leptons and the beyond the Standard Model (BSM) fermions. Through such mixings the BSM fermions can be produced at a high energy collider such as Large Hadron Collider (LHC) and 100 TeV pp collider, using W boson and Z boson exchange from the charged current and neutral current interactions respectively. In fact due to the $B-L$ model framework, the pair production of such fermions can be tested through the $B-L$ gauge boson. These fermions can display the multilepton final states through the corresponding charged current and neutral current interactions [52–58] which will be interesting in the high luminosity era of the high energy collider/s. Moreover a general parameter structure can also be adopted for such models as discussed in [59,60] using the Casas-Ibarra parametrization [61].

In addition, we have several Z_2 odd scalars including DM where heavier particles decay into SM leptons and DM via Yukawa interactions. Thus the signals of charged leptons with missing transverse momentum are expected as a signature of these scalar particles. We estimated the cross section of pair production of heavy charged leptons via electroweak process. Then we found that $\mathcal{O}(0.1)$ fb cross section is obtained when heavy charged lepton mass is around 1 TeV. More detailed discussion with simulation is left as future work.

In the future a general version of this model under the $U(1)_X$ gauge group can also be considered. Recently the $U(1)_X$ extended SM has been investigated in a variety of contexts, such as the classical conformality [62,63], Z'-portal dark matter [64], and cosmological inflation scenario [65].

Finally, we also want to comment that such a model can be useful to study baryogenesis via leptogenesis [66–75] as we can do in the $B-L$, $U(1)_x$, inverse seesaw models. In this model we also have such possibilities to consider three generations of heavy fermions being couples with the SM scalar sector. Such fermions can be nondegenerate, too. Such nondegenerate heavy fermions can have sizable mixings with the SM light neutrinos which are dependent upon the neutrino oscillation data and the free model parameters such as the Dirac phase, Majorana phase, heavy fermion masses, and the Casas-Ibarra parametrization. An elaborate discussion on leptogenesis in this model is beyond the scope of this paper and will be considered as a separate work in the near future.

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