Mirror symmetry: from active and sterile neutrino masses to baryonic and dark matter asymmetries

Pei-Hong Gu
Max-Planck-Institut für Kernphysik, Saupfercheckweg 1, 69117 Heidelberg, Germany

We consider an $SU(3)_c \times SU(2)_L \times U(1)_Y$ mirror sector where the field content and dimensionless couplings are a copy of the $SU(3)_c \times SU(2)_L \times U(1)_Y$ ordinary sector. Our model also contains three gauge-singlet fermions with heavy Majorana masses and an $[SU(2)_L \times SU(2)_R]$-bidoublet Higgs scalar with seesaw-suppressed vacuum expectation value. The mirror sterile neutrino masses will have a form of canonical seesaw while the ordinary active neutrino masses will have a form of double and linear seesaw. In this canonical and double-linear seesaw scenario, we can expect one sterile neutrino at the eV scale and the other two above the MeV scale to fit the cosmological and short baseline neutrino oscillation data. Associated with the $SU(2)_L$ and $SU(2)_R$ sphaleron processes, the decays of the fermion singlets can simultaneously generate a lepton asymmetry in the $[SU(2)_L]$-doublet leptons and an equal lepton asymmetry in the $[SU(2)_R]$-doublet leptons to explain the existence of baryonic and dark matter. The lightest mirror baryon then should have a determined mass around 5 GeV to account for the dark matter relic density. The $U(1)$ kinetic mixing can open a window for dark matter direct detection.

PACS numbers: 98.80.Cq, 95.35.+d, 14.60.Pq, 12.60.Cn, 12.60.Fr

I. INTRODUCTION

Various neutrino oscillation experiments have firmly established the three active neutrino oscillation picture \cite{ref1}, however, some short baseline neutrino oscillation experiments \cite{ref2} and recent re-evaluations of reactor antineutrino fluxes \cite{ref3} hint the existence of additional sterile neutrinos with the eV-scale masses \cite{ref4,ref5}. The light active neutrinos can be elegantly understood in the seesaw \cite{ref6,ref7} or other \cite{ref8,ref9} extensions of the $SU(3)_c \times SU(2)_L \times U(1)_Y$ standard model (SM). If the light sterile neutrinos are eventually confirmed, we need explain not only the existence and small masses of the sterile neutrinos but also the mixing between the active and sterile neutrinos. Such sterile neutrinos can naturally appear in the mirror \cite{ref22,ref23} universe models discussed in the literature \cite{ref42,ref43}. There are also other ideas for the light sterile neutrinos \cite{ref52,ref53}. On the other hand, precise cosmology has indicated that in the present universe the energy density of the dark matter is comparable with that of the ordinary matter \cite{ref59}. This raises an interesting possibility that the dark and ordinary matter may have a special relation although their properties are very different. For example, the dark matter relic density may be an asymmetry between the dark matter and antimatter \cite{ref22,ref23,ref60,ref63} since the ordinary matter exists as a baryon asymmetry. In particular, the asymmetric dark matter can be well motivated in the mirror universe models \cite{ref22,ref23}.

In this paper we shall propose a novel mirror universe model to give the active and sterile neutrino masses as well as the baryonic and dark matter asymmetries. In addition to the $SU(3)_c \times SU(2)_L \times U(1)_Y$ ordinary sector and its $SU(3)_c \times SU(2)_L \times U(1)_Y$ mirror partner, our model contains three gauge-singlet Majorana fermions and an $[SU(2)_L \times SU(2)_R]$-bidoublet Higgs scalar. The mirror symmetry is softly broken to allow the symmetry breaking pattern in the mirror sector different from that in the ordinary sector. The mirror photon can become massive according to the mirror electromagnetic symmetry breaking. By integrating out the fermion singlets and the Higgs bidoublet, we can get the mirror sterile neutrino masses by canonical \cite{ref10} seesaw, as well as the ordinary active neutrino masses by double \cite{ref17} and linear \cite{ref18} seesaw. In this canonical and double-linear seesaw scenario, two sterile neutrinos are above the MeV scale and the other one is at the eV scale, so that their existence can fulfill the cosmological and short baseline neutrino oscillation data \cite{ref4}. Our model can realize a leptogenesis \cite{ref94,ref104} as a common origin of the ordinary and dark matter. Specifically, the decays of the fermion singlets can simultaneously generate a desired lepton asymmetry in the $[SU(2)_R]$-doublet leptons and an equal lepton asymmetry in the $[SU(2)_L]$-doublet leptons even if we do not resort to the fine tuned resonant effect \cite{ref96}. Through the sphaleron-induced lepton-to-baryon conversion \cite{ref103}, we can obtain an ordinary baryon asymmetry and an equal mirror baryon asymmetry to account for the number densities of ordinary and dark matter. The lightest mirror baryon then should have a determined mass around 5 GeV to explain the ratio between the ordinary and dark matter energy densities. In the presence of a $U(1)$ kinetic mixing, the dark matter particle can be verified in the ongoing and future dark matter direct detection experiments.

Our model shares some ideas of Ref. \cite{ref37}, where the authors introduced two $[SU(2)]$-triplet Higgs scalars to generate the active neutrino masses by type-II \cite{ref8} and inverse \cite{ref19} seesaw. They also assumed all of the mirror neutrinos above the MeV scale. Furthermore, they res-
onantly enhanced the CP asymmetries in the decays of the fermion singlets to produce the required lepton and baryon asymmetries.

II. THE MODEL

There are two Higgs scalars in both of the ordinary and dark sectors,

\[
\begin{align*}
\phi_d(1, 2, +1) &= \begin{bmatrix} \phi_{d}^{+} \\ \phi_{d}^{0} \end{bmatrix} \leftrightarrow \phi_d'(1, 2, +1) = \begin{bmatrix} \phi_{d}^{+} \\ \phi_{d}^{0} \end{bmatrix}, \\
\phi_u(1, 2, -1) &= \begin{bmatrix} \phi_{u}^{0} \\ \phi_{u}^{−} \end{bmatrix} \leftrightarrow \phi_u'(1, 2, -1) = \begin{bmatrix} \phi_{u}^{0} \\ \phi_{u}^{−} \end{bmatrix}.
\end{align*}
\]  

Here and thereafter the mirror fields are denoted by a prime on a symbol, and hence the parentheses following the mirror fields are the quantum numbers under the \(G' = SU(3)_{c}' \times SU(2)_{L}' \times U(1)'_Y\) gauge group, while the parentheses following the mirror fields are the quantum numbers under the \(G = SU(3)_{c} \times SU(2)_{L} \times U(1)_{Y}\) gauge group. Our model also contains three families of ordinary and mirror fermions,

\[
\begin{align*}
q_L(3, 2, +1) = \begin{bmatrix} u_L \\ d_L \end{bmatrix}, & \leftrightarrow q_l'(3, 2, +1) = \begin{bmatrix} u_L' \\ d_L' \end{bmatrix}, \\
d_R(3, 1, -2) = \begin{bmatrix} u_R \\ d_R \end{bmatrix}, & \leftrightarrow d_R'(3, 1, -2), \\
u_R(3, 1, +4) = \begin{bmatrix} u_R' \\ d_R' \end{bmatrix}, & \leftrightarrow u_R'(3, 1, +4), \\
\ell_L(1, 2, -1) = \begin{bmatrix} \nu_L \\ e_L \end{bmatrix}, & \leftrightarrow \ell_L'(1, 2, -1) = \begin{bmatrix} \nu_L' \\ e_L' \end{bmatrix}, \\
e_R(1, 1, -2) = \begin{bmatrix} \nu_R \\ e_R \end{bmatrix}, & \leftrightarrow e_R'(1, 1, -2),
\end{align*}
\]

with the family indices being omitted for simplicity. We further introduce three gauge-singlet fermions \[12, 43,

\[
N_{Ri}(1, 1, 0)(1, 1, 0) \leftrightarrow N_{Ri} \ (i = 1, 2, 3),
\]

and an \([SU(2)_L \times SU(2)_{L}']\)-doublet scalar \[40, 42,

\[
\Sigma(1, 2, +1)(1, 2, +1) = \begin{bmatrix} \sigma^0 & \sigma^{(−, 0)} \\ \sigma^{(0, −)} & \sigma^{(−, −)} \end{bmatrix} \leftrightarrow \Sigma,
\]

where the first and second parentheses are the quantum numbers under the \(G\) and \(G'\) gauge groups, respectively. Besides the gauge symmetries and the mirror discrete symmetry, we impose a \(U(1)_{G} \times U(1)_{G} \) global symmetry under which only the \(SU(2)\) doublets and the \([SU(2)_L \times SU(2)_{L}']\)-doublet are nontrivial: \((1, 0)\) for the \(SU(2)_L\) doublets, \((0, 1)\) for the \([SU(2)_L]\)-doublets and \((1, 1)\) for the \([SU(2)_L \times SU(2)_{L}']\)-doublet.

A. Interactions

For simplicity, we only write down the following terms of the full Lagrangian,

\[
\mathcal{L} \supset -y_u(\bar{U}_L \phi_d d_R + \bar{U}_L \phi_d d_R') - y_u(\bar{U}_L \phi_d u_R + \bar{U}_L \phi_d u_R')
- y_u(\bar{U}_L \phi_d e_R + \bar{U}_L \phi_d e_R') - h(\bar{L}_L \phi_u N_R + \bar{L}_L \phi_u N_R')
- \frac{1}{2} M_N N_{R'} N_R - f_{\ell} L_y c Y_L - \rho \delta_u Y_u \sigma_u^* + \text{h.c.}
- \frac{1}{2} M_{\Sigma}^2 \text{Tr}(\Sigma^\dagger \Sigma) - \frac{\epsilon}{2} B_{\mu \nu} B^{\mu \nu},
\]

where \(B_{\mu \nu}\) and \(B_{\mu \nu}^*\) are the \(U(1)_Y\) and \(U(1)'_Y\) field strength tensors. Note the other gauge-invariant trilinear couplings have been forbidden by the \(U(1)_G \times U(1)'_G\) global symmetry.

As a result of the mirror symmetry, the Yukawa couplings of the \([SU(2)_L \times SU(2)_{L}']\)-doublet scalar to the \([SU(2)_L]\)-doublets should have a symmetric structure,

\[
f = f^T.
\]

Furthermore, the mass term of the gauge-singlet fermions and the trilinear coupling of the \([SU(2)_L \times SU(2)_{L}']\)-doublet scalar to the \([SU(2)_L]\)-doublets softly break both of the ordinary and dark lepton numbers. Without loss of generality and for convenience, we will choose the base where the gauge-singlet fermions have a diagonal and real mass matrix,

\[
M_N = \text{diag}\{M_{N_1}, M_{N_2}, M_{N_3}\},
\]

to define three Majorana fermions,

\[
N_i = N_{Ri} + N_{Ri}^*.
\]

Meanwhile, the \([SU(2)_L \times SU(2)_{L}']\)-doublet scalar can have a real cubic coupling to the \([SU(2)_L]\)-doublet scalars by a proper phase rotation, i.e.

\[
\rho = |\rho|.
\]

B. Vacuum expectation values

We allow the discrete mirror symmetry and the global \(U(1)_G \times U(1)'_G\) symmetry to be softly broken by the quadratic terms in the full scalar potential, which is not shown for simplicity. So, the mirror Higgs scalars can develop the vacuum expectation values (VEVs) different from those of the ordinary Higgs scalars. In particular, the charged components of the mirror Higgs scalars can have the nonzero VEVs \[106, 107\]. In this case, the symmetry breaking pattern should be

\[
SU(3)_c \times SU(2)_L \times U(1)_Y \rightarrow SU(3)_c \times U(1)_e, \\
SU(3)_c \times SU(2)_{L'} \times U(1)'_Y \rightarrow SU(3)'_c \times U(1)'_e, \\
SU(3)_c \times SU(2)_{L} \times U(1)'_Y \rightarrow SU(3)'_c \times U(1)'_e.
\]
Accordingly, the mirror photon can become massive although the ordinary photon keeps massless.

The VEVs of the \([SU(2)]_L\)-doublet Higgs scalars \(\phi_u\) and \(\phi_d\) should be fixed by

\[
\langle \phi_u \rangle = \begin{bmatrix} \langle \phi^0_u \rangle \\ 0 \end{bmatrix}, \quad \langle \phi_d \rangle = \begin{bmatrix} 0 \\ \langle \phi^0_d \rangle \end{bmatrix}
\]

with

\[
\langle \phi \rangle^2 = \sqrt{\langle \phi_u \rangle^2 + \langle \phi_d \rangle^2} \approx 174 \text{ GeV}, \tag{11}
\]

while the VEVs of the \([SU(2)]_L\)-doublet Higgs scalars \(\phi'_u\) and \(\phi'_d\) can be described by

\[
\langle \phi'_u \rangle = \begin{bmatrix} \langle \phi^0_u \rangle \\ 0 \end{bmatrix}, \quad \langle \phi'_d \rangle = \begin{bmatrix} \langle \phi^0_d \rangle \\ \langle \phi^0_d \rangle \end{bmatrix}.
\]

\[
\tag{12}
\]

The \([SU(2)]_L \times SU(2)]_L\)-bidoublet Higgs scalar \(\Sigma\) can pick up the VEV through its cubic coupling to the \([SU(2)]_L\)-doublet Higgs scalars \(\phi_u\) and \(\phi'_u\),

\[
\langle \Sigma \rangle = \begin{bmatrix} \langle \phi^0_u \rangle \\ 0 \\ 0 \end{bmatrix}
\]

with \(\langle \sigma \rangle \approx -\frac{\rho \langle \phi^0_u \rangle \langle \phi^0_u \rangle}{M^2_\Sigma}\). \tag{13}

Clearly, the VEV \(\langle \Sigma \rangle\) can be much smaller than the VEVs \(\langle \phi^0_u \rangle\) and \(\langle \phi^0_u \rangle\) in the seesaw scenario, i.e.

\[
\langle \Sigma \rangle \ll \langle \phi^0_u \rangle, \quad \langle \phi^0_u \rangle \ll \rho \approx M_\Sigma. \tag{14}
\]

**C. Mirror photon**

We can make a non-unitary transformation

\[
\tilde{B}_\mu = B_\mu + \epsilon B'_\mu, \quad \tilde{B}'_\mu = \sqrt{1 - \epsilon^2} B'_\mu, \tag{15}
\]

to remove the \(U(1)_Y \times U(1)'_Y\) kinetic mixing and then define the orthogonal fields,

\[
A_\mu = W^3_\mu \sin \theta_W + \tilde{B}_\mu \cos \theta_W, \\
Z_\mu = W^3_\mu \cos \theta_W - \sin \tilde{B}_\mu \theta_W, \\
A'_\mu = W'_{3\mu} \sin \theta_W + \tilde{B}'_\mu \cos \theta_W, \\
Z'_\mu = W'_{3\mu} \cos \theta_W - \tilde{B}'_\mu \sin \theta_W. \tag{16}
\]

Here \(\theta_W\) with \(\sin^2 \theta_W \approx 0.231\) is the Weinberg angle while \(W^3_\mu\) and \(W'_{3\mu}\) are the \([SU(2)]_L\) and \([SU(2)]_L'\) gauge fields. In the above orthogonal base, the field \(A_\mu\) is exactly massless and is the physical mass-eigenstate field, the ordinary photon, according to the unbroken electromagnetic symmetry \(U(1)_{em}\) in the ordinary sector, while the others \(Z_\mu, Z'_\mu\) and \(A'_\mu\) will mix together. The mirror electromagnetic symmetry \(U(1)'_{em}\) is broken by the charged VEV \(\langle \phi^0_u \rangle\) given in Eq. \([12]\). Consequently, the \(W^{\pm}\) boson will also mix with the \(Z'\) boson and the mirror photon \(A'_\mu\), which is massive now.

The mirror photon can couple to the ordinary fermions besides the mirror fermions,

\[
\mathcal{L} \supset e A'_\mu \left\{ \frac{e}{4} (\bar{e} \gamma^\mu (3 + \gamma_5) e + \bar{\nu} \gamma^\mu (1 - \gamma_5) \nu) + \bar{d} \gamma^\mu \left( \frac{1}{3} + \gamma_5 \right) d - \bar{u} \gamma^\mu \left( \frac{5}{3} + \gamma_5 \right) u + \left( -\frac{1}{3} \bar{d} \gamma^\mu d' + \frac{2}{3} \bar{u} \gamma^\mu u' - \bar{e} \gamma^\mu e' \right) \right\} \quad \text{for} \quad \epsilon \ll 1. \tag{17}
\]

In the case with \(\langle \phi^0_u \rangle = \mathcal{O}(100 \text{ MeV})\), the mirror photon can have a mass

\[
m_{A'} \approx \sqrt{8 \pi \alpha \langle \phi^0_u \rangle} = \mathcal{O}(100 \text{ MeV}), \tag{18}
\]

and its decay width will not be smaller than

\[
\Gamma_{A'} = \Gamma_{A' \rightarrow \nu \bar{\nu}} + \Gamma_{A' \rightarrow e\bar{e}} + \Gamma_{A' \rightarrow uu} + \Gamma_{A' \rightarrow dd} \approx \frac{5}{9} \epsilon^2 m_{A'}. \tag{19}
\]

Here \(\alpha = e^2/(4\pi) \approx 1/137\) is the fine structure constant.

**III. SEESAW FOR ACTIVE AND STERILE NEUTRINO MASSES**

From Eq. \([5]\), the ordinary active neutrinos, the mirror sterile neutrinos and the gauge-singlet fermions will have the mass matrix as below,

\[
\mathcal{L} \supset -\frac{1}{2} \begin{bmatrix} \bar{\nu}_L & \bar{\nu}'_L & \bar{N}_R \end{bmatrix} \begin{bmatrix} 0 & f \langle \Sigma \rangle & h \langle \phi_u \rangle \\ f^T \langle \Sigma \rangle & 0 & h \langle \phi'_u \rangle \\ h^T \langle \phi_u \rangle & h^T \langle \phi'_u \rangle & M_N \end{bmatrix} \begin{bmatrix} \nu^c_L \\ \nu^c_L' \\ N_R \end{bmatrix} + \text{H.c.}, \tag{20}
\]

after the ordinary and mirror electroweak symmetry breaking.

**A. Active and sterile neutrino masses and mixing**

As long as the gauge-singlet fermions are heavy enough, i.e.

\[
M_N \gg h \langle \phi_u \rangle, \quad h \langle \phi'_u \rangle, \quad f \langle \Sigma \rangle, \tag{21}
\]
we can make use of the seesaw mechanism to get the mass matrix of the active and sterile neutrinos,
Here the mass eigenvalues have been introduced,

\[
\hat{m}_\nu = \text{diag}\{m_{\nu_1}, m_{\nu_2}, m_{\nu_3}\}, \quad \hat{m}_{\nu'} = \text{diag}\{m_{\nu'_1}, m_{\nu'_2}, m_{\nu'_3}\}.
\]

If the entries in the mass matrix have the following hierarchy,

\[
-h \frac{(\phi_u')}{M_N} h^T \gg f(\Sigma) - h \frac{(\phi_u)}{M_N} h^T, \quad -h \frac{(\phi_u')}{M_N} h^T
\]

the sterile neutrino masses should have a form of the canonical seesaw,

\[
\mathcal{L} \supset -\nu_L m_{\nu'} \nu' L + \text{H.c.\ with\ } m_{\nu'} = -h \frac{(\phi_u')}{M_N} h^T, \quad \text{(25)}
\]

while the active neutrino masses should have a form of the double and linear seesaw,

\[
\mathcal{L} \supset f \frac{(\Sigma)^2}{h \frac{(\phi_u)}{M_N} h^T} - 2 f \frac{(\phi_u)}{h \frac{(\phi_u)}{M_N} h^T} \cdot \text{(26)}
\]

Under the seesaw condition, the active mixing matrix \(U_\nu\) and the sterile mixing matrix \(U_{\nu'}\) can approximate to the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrices,

\[
U_\nu^\dagger U_\nu = U_{\nu'}^\dagger U_{\nu'} = 1, \quad U_{\nu'}^\dagger U_{\nu'} = U_{\nu'}^\dagger U_{\nu'} = 1, \quad \text{(27)}
\]

while the active-sterile mixing matrices \(U_{\nu}\) and \(U_{\nu'}\) can be simplified by

\[
U_{\nu\nu'} = f U_{\nu'} \frac{(\Sigma)}{m_{\nu'}} + \frac{(\phi_u)}{h (\phi_u)} U_{\nu' U_{\nu'}}, \quad U_{\nu\nu'} = -U_{\nu'} U_{\nu'} U_{\nu'} \cdot \text{(28)}
\]

**B. Sterile neutrino decays**

Due to their mixing with the ordinary neutrinos, the sterile neutrinos can decay into the ordinary fermions

\[
\Gamma_{\nu' \rightarrow \nu \nu \nu} = \frac{G_F^2 m_{\nu'}^5}{96 \pi^3} (U_{\nu'}^\dagger U_{\nu'} \nu_{\nu} \nu_{\nu}) ii, \quad \text{(29a)}
\]

\[
\Gamma_{\nu' \rightarrow \nu e^+ e^-} = \frac{5 G_F^2 m_{\nu'}^5}{768 \pi^3} (U_{\nu'}^\dagger U_{\nu'} \nu_{\nu} \nu_{\nu}) ii, \quad \text{(29b)}
\]

\[
\Gamma_{\nu' \rightarrow \nu \nu \bar{\nu}} = \frac{G_F^2 m_{\nu'}^5}{32 \pi^3} (1 - \frac{8}{3} W^2 + \frac{32}{9} s_W^2) (U_{\nu'}^\dagger U_{\nu'} \nu_{\nu} \nu_{\nu}) ii, \quad \text{(29c)}
\]

\[
\Gamma_{\nu' \rightarrow e^- \bar{\nu} d} = \Gamma_{\nu' \rightarrow e^- d^*} = \frac{|V_{ud}|^2 G_F^2 m_{\nu'}^5}{32 \pi^3} (U_{\nu'}^\dagger U_{\nu'} \nu_{\nu} \nu_{\nu}) ii, \quad \text{(29d)}
\]

if the kinematics is allowed. Here \(G_F = 1.16637 \times 10^{-5} \text{GeV}^{-2}\) is the Fermi constant, \(s_W^2 = \sin^2 \theta_W\) is the Weinberg angle, while \(V_{ud} \simeq 0.97419\) is an element of the Cabibbo-Kobayashi-Maskawa matrix.

**IV. LEPTOGENESIS FOR ORDINARY AND MIRROR BARYON ASYMMETRIES**

If CP is not conserved, the decays of the heavy Majorana fermions \(N_i\) can simultaneously generate two types of lepton asymmetries: one is stored in the \([SU(2)_L]_d\)-doublet leptons \(l_L\), i.e.

\[
\eta_L = \frac{n_{l_L}}{s}, \quad \text{(30)}
\]

and the other is stored in the mirror leptons \(l'_L\), i.e.

\[
\eta_L' = \frac{n_{l'_L}}{s}, \quad \text{(31)}
\]

Here \(n_{l_L}\) and \(n_{l'_L}\) are the number densities while \(s\) is the entropy density. The relevant diagrams are shown in Fig. 1.

The \([SU(2)_L]_d\) sphaleron processes then will partially transfer the ordinary lepton asymmetry to an ordinary baryon asymmetry,

\[
\eta_B = -\frac{28}{79} \eta_L. \quad \text{(32)}
\]
FIG. 1: The heavy Majorana fermions ($N_i = N_{Ri} + N_{ci}$, with $N_{Ri}$ being three gauge singlets) decays into the $[SU(2)_L]$-doublet leptons and Higgs scalar ($l_L, \phi_u$) as well as into the $[SU(2)_{L}']$-doublet dark leptons and Higgs scalar ($l'_{L}, \phi'_{u}$). Here $\Sigma$ is a heavy $[SU(2)_L \times SU(2)'_L]$-bidoublet scalar.

Similarly, the mirror lepton asymmetry will be partially converted to a mirror baryon asymmetry,

$$\eta'_B = -\frac{28}{79} \eta'_{L'},$$

through the $SU(2)'_L$ sphaleron processes. Due to the Yukawa couplings in Eq. (5), the ordinary lepton asymmetry and then the ordinary baryon asymmetry should equal to the mirror ones, i.e.

$$\eta_L = \eta'_{L'} \propto \varepsilon_{N_i} \Rightarrow \eta_B = \eta'_B \propto \varepsilon_{N_i}.$$

Here $\varepsilon_{N_i}$ is the CP asymmetry in the decays of the heavy Majorana fermions $N_i$.

A. CP violation in decays

The total decay width in the decays of the heavy Majorana fermions $N_i$ can be calculated at tree level,

$$\Gamma_{N_i} = \Gamma_{N_i \rightarrow l_L \phi_u} + \Gamma_{N_i \rightarrow l'_{L'} \phi'_{u}} + \Gamma_{N_i \rightarrow l'_{L'_{L}} \phi'_{u}} + \Gamma_{N_i \rightarrow l_{L_{L}} \phi_u} = \frac{1}{4\pi} (h^\dagger h)_{ii} M_{N_i}.$$  

We then can compute the CP asymmetry $\varepsilon_{N_i}$ appeared in Eq. (34) at one-loop level,
The BNN constraint enforces the self-energy correction, $V(x)$ and $\tilde{V}(x,y)$ being the vertex corrections, and $\tilde{\Sigma}$ the bidoublet Higgs, can have an upper bound, $\varepsilon_{N_i} \leq M_{N_i}^2$. Similar to the Davidson-Ibarra bound [98] in the processes such as $2\tilde{\Sigma}L_1\phi^c_3 \rightarrow 2\tilde{\Sigma}L^\prime_1\phi^c_3$, $M_{N_i}$ is the Hubble constant with $\rho_{Pl} \equiv 1.22 \times 10^{19}$ GeV being the Planck mass and $g_* = 2 \times (106.75 + 2) = 217.5$ being the relativistic degrees of freedom. Below the masses of the mediators $N_i$ and $\Sigma$, the interaction rate can be given by

$$\Gamma_s = \frac{2}{\pi^3} \left(\frac{T^3}{\langle \phi_u^i \rangle^4}\right) \text{Tr} \left[ m_{\nu_i}^2 m_{\nu_i} + \left(\frac{\langle \phi_u^i \rangle}{\langle \phi_u^i \rangle^4}\right) m_{\nu_i}^{\text{Linear}} \right]$$

\[ \times \left( m_{\nu_i} + \frac{\langle \phi_u^i \rangle}{2\langle \phi_u^i \rangle^4} m_{\nu_i}^{\text{Linear}} \right) \text{ for } T \ll M_N, M_\Sigma, \] (37b)

for the constraint [99]. Alternatively, the scattering processes can decouple at a temperature above or near the mediator’s mass if the interactions are weak enough to satisfy [112]

$$K_{N_i} = \frac{\Gamma_{N_i}}{2H(T)} \big|_{T=M_{N_i}} \ll 1, \quad K_{\Sigma} = \frac{\Gamma_{\Sigma}}{2H(T)} \big|_{T=M_{\Sigma}} \ll 1.$$ (45)

\[ K_{N_i} \]

\[ K_{\Sigma} \]

C. Final baryon asymmetries

In the case the lightest Majorana fermion $N_i$ has a mass smaller than the decouple temperature of the scattering processes mediated by the other Majorana fermions $N_{2,3}$ and the Higgs bidoublet $\Sigma$, i.e.

$$M_{N_i} < T_{D},$$

the final baryon asymmetries can be approximately solved by [112]

$$\eta_B = \eta_B^i \simeq \frac{-28}{79} \times \frac{\varepsilon_{N_i}}{g_*} \times \kappa$$

\[ = 0.888 \times 10^{-10} \left(\frac{\varepsilon_{N_i}}{5.45 \times 10^{-8}}\right) \left(\frac{K_{N_i}}{1}\right) \text{ with } \kappa \simeq \begin{cases} 1 & K_{N_i} \ll 1, \\ \frac{\phi_{N_i}^3}{\ln \phi_{N_i}^{1/4}} & K_{N_i} \gg 1. \end{cases} \] (47)

\[ \eta_B \]

\[ \eta_B^i \]

\[ \kappa \]
V. IMPLICATIONS AND CONSTRAINTS

Before giving the concrete parameter choice, we shall demonstrate some general implications and constraints on the model.

A. Dark matter mass

From the Yukawa couplings in Eq. (5), we can easily read the relation between the ordinary and mirror charged fermion masses,

\[ \frac{\langle \phi_d' \rangle}{\langle \phi_d \rangle} = \frac{m_{d'}}{m_d} = \frac{m_{d'}}{m_s} = \frac{m_{b'}}{m_b} = \frac{m_{e'}}{m_e} = \frac{m_{\tau'}}{m_\tau}, \]

\[ \frac{\langle \phi_u' \rangle}{\langle \phi_u \rangle} = \frac{m_{u'}}{m_u} = \frac{m_{u'}}{m_c} = \frac{m_{t'}}{m_t}. \] (48)

As a result of the mirror symmetry, the ordinary and mirror strong coupling constants should become equal at sufficiently high scales. Therefore, the beta functions of the ordinary and mirror QCD can govern the dependence of the mirror hadronic scale on the ordinary one \[37\],

\[ \Lambda_{\text{QCD'}} \left( \frac{\langle \phi_u \rangle}{\langle \phi_u' \rangle} \right) = \left( \frac{\tan \beta}{\tan \beta'} \right) \left( m_u m_d m_s \right) \Lambda_{\text{QCD}} \]

for \[\Lambda_{\text{QCD'}} < m_{u'}, m_{d'}\],

where we have defined

\[ \tan \beta = \frac{\langle \phi_u \rangle}{\langle \phi_d \rangle}, \quad \tan \beta' = \frac{\langle \phi_u' \rangle}{\langle \phi_d' \rangle}. \] (50)

In the ordinary sector, the current quark masses \[m_u\] and \[m_d\] are much smaller than the hadronic scale \[\Lambda_{\text{QCD}}\] so that they can only have a negligible contribution to the nucleon masses,

\[ m_p \simeq m_n \simeq 940 \text{ MeV} = m_N. \] (51)

In addition, the \[\Delta\] baryons and the neutron has a mass split from the hyperfine interaction among the constituent quarks \[37\],

\[ m_{\Delta} - m_{n} \simeq 300 \text{ MeV} \ll \frac{\Lambda_{\text{QCD}}^3}{m_q^2}, \] (52)

with \[m_q \simeq 300 \text{ MeV}\] being the constituent quark mass. In the mirror sector, the quark masses \[m_{u'}\] and \[m_{d'}\] may be larger than the hadronic scale \[\Lambda_{\text{QCD'}}\]. The mirror proton and neutron masses then can approximately equal to the sum of the mirror quark masses,

\[ m_{p'} = 2m_{u'} + m_{d'}, \quad m_{n'} = 2m_{d'} + m_{u'}, \] (53)

which implies

\[ m_{p'} < m_{n'} \text{ for } m_{u'} < m_{d'}, \] (54a)

\[ m_{p'} > m_{n'} \text{ for } m_{u'} > m_{d'}. \] (54b)

In the case \[54b\], the mirror \[\Delta^{(-)}\] baryon can be lighter than the mirror neutron if the mirror hyperfine interaction doesn’t compensate the mass difference \[m_{u'} - m_{d'}\] \[37\],

\[ m_{\Delta^{-}} = 3m_{d'} + (m_{\Delta} - m_{n}) \frac{\Lambda_{\text{QCD'}}}{\Lambda_{\text{QCD}}} \left( \frac{m_q^2}{m_{u'} m_{d'}} \right)^3. \] (55)

As the lightest mirror baryon is expected to serve as the dark matter particle, its mass should be

\[ m_{\text{DM}} \simeq 5m_N \simeq 5 \text{ GeV}, \] (56)

to explain the cosmological observations,

\[ \Omega_B h^2 : \Omega_{\text{DM}} h^2 = \frac{m_N \eta_B}{m_{\text{DM}}} \eta_B = m_N : m_{\text{DM}} \simeq 1 : 5. \] (57)

B. Dark matter direct detection

In the presence of the \[U(1)_Y \times U(1)_Y\] kinetic mixing, the mirror photon can mediate a scattering of the dark matter particle off the ordinary nucleons. For example, the mirror proton \[p'\] or the mirror \[\Delta^{(-)}\] baryon has a spin-independent cross section,

\[ \sigma_{XN \rightarrow XN} \simeq \epsilon^2 \frac{\alpha^2 \mu_r^2}{m_{A'}} \left[ \frac{3Z + (A - Z)}{A} \right]^2, \]

\[ \simeq 10^{-41} \text{ cm}^2 \left( \frac{\epsilon}{1.5 \times 10^{-7}} \right)^2 \times \left( \frac{\mu_r}{0.833 \text{ GeV}} \right)^2 \left( \frac{100 \text{ MeV}}{m_{A'}} \right)^4 \times \left[ \frac{3Z + (A - Z)}{A} \right]^2, \] (58)

which can be close to the XENON10 limit \[114\]. Here \[X\] denotes the mirror proton \[p'\] or the mirror \[\Delta^{(-)}\] baryon, \[Z\] and \[A - Z\] are the numbers of proton and neutron within the target nucleus, while

\[ \mu_r = \frac{m_X m_N}{m_X + m_N} \simeq 0.833 \text{ GeV} \]

for \[m_X \simeq 5m_N \simeq 5 \text{ GeV}, \] (59)

is the reduced mass. Alternatively, the mirror neutron \[n'\] can serve as the dark matter particle if it is the lightest mirror baryon. The mirror neutron as the dark matter particle can have an energy-dependent cross section. The detailed studies can be found in \[37\].

C. Constraints on sterile neutrinos and mirror photon

The Big-Bang Nucleosynthesis (BBN) stringently restricts the existence of the new relativistic degrees of
freedom. The constraint on the new degrees of freedom is conventionally quoted as the effective number of additional light neutrinos. The latest Planck 2013 results show $N_{eff} = 3.30 \pm 0.27$. So, one light sterile neutrino can be allowed at $3 \sigma$ level. A very recent analysis on the neutrino oscillation data also hint at the existence of an additional neutrino with an eV-scale mass. This means the other two sterile neutrinos should have the masses heavier than a few MeV and have a lifetime shorter than 1 second. From Eqs. (25) and (29), we hence can put

$$m_{\nu'}^\text{max} > 92 \text{ MeV} \left( \frac{\langle \phi_u' \rangle / \langle \phi_u \rangle}{\langle \phi_u' \rangle / \langle \phi_u \rangle} \right)^{3/2} \left( \frac{\langle \phi_u' \rangle / \langle \phi_u \rangle}{2000} \right)^{3/2} \times \left( \frac{1}{\tau_{\nu'}^{\text{sec}}} \right)^{3/2} \text{ for } \tau_{\nu'}^{\text{sec}} < 1 \text{ sec}. \quad (60)$$

The mirror photon should also satisfy the BBN constraint. From Eq. (14), it is easy to see

$$\tau_{A'} \simeq \left( \frac{4 \times 10^{-11}}{\epsilon} \right)^2 \left( \frac{100 \text{ MeV}}{m_{A'}} \right)^2 \text{ sec}. \quad (61)$$

Clearly, the mirror photon $A'$ with a mass $m_{A'} = 100 \text{ MeV}$ can have a lifetime shorter than 1 second if we take $\epsilon > 4 \times 10^{-11}$. Currently, the measurement on the muon magnetic moment constrains $\epsilon^2 \cos^2 \theta_W \cos^2 \theta_W < 2 \times 10^{-5}$ for $m_{A'} = 100 \text{ MeV}$ \cite{14}.

Furthermore, the active-sterile neutrino mass matrix should be constrained by the neutrinoless double beta decay experiments. In the regime of $m_{\nu'}^\text{max} \lesssim 100 \text{ MeV}$, we can perform \cite{17}

$$|m_{\beta\beta}| = \left| \left( -h \frac{\langle \phi_u' \rangle}{M_N} h^T \right)_{11} \right| < 0.2 \text{ eV}$$

$$\Rightarrow |(m_{\nu'})_{11}| = \left| \sum_i \left[ (U_{\nu'})^{*}_{11} \right]^2 m_{\nu'} \right| < 0.8 \text{ MeV} \left( \frac{\langle \phi_u' \rangle / \langle \phi_u \rangle}{2000} \right)^2. \quad (62)$$

D. Leptogenesis scale

From Eqs. (40) and (47), the CP asymmetry $|\varepsilon_{N_1}|$ should be bigger than

$$\varepsilon_{N_1}^\text{max} > 5.45 \times 10^{-8}, \quad (63)$$

to explain the observed baryon asymmetry. Accordingly, we can have a low limit on the leptogenesis scale,

$$M_{N_1} > 1.3 \times 10^6 \text{ GeV} \left( \frac{100 \text{ MeV}}{m_{\nu'}^\text{max}} \right)^2 \left( \frac{\langle \phi_u' \rangle / \langle \phi_u \rangle}{2000} \right)^2, \quad (64)$$

which is expected to be below the critical temperature \cite{49},

$$T_D > 1.1 \times 10^7 \text{ GeV} \left[ \frac{2(m_{\nu'}^\text{max})^2}{\sum_i m_{\nu'_i}^2} \right] \left( \frac{100 \text{ MeV}}{m_{\nu'}^\text{max}} \right)^2 \times \left[ \left( \frac{\langle \phi_u' \rangle / \langle \phi_u \rangle}{2000} \right)^4 \right]. \quad (65)$$

VI. AN EXAMPLE OF PARAMETER CHOICE

As an example, let us set

$$\langle \phi_u' \rangle = 2000 \langle \phi_u \rangle, \quad \tan \beta' = 380, \quad \tan \beta = 50, \quad (66)$$

to give the mirror charged fermion masses \cite{109},

$$m_{\nu'} = 1.3 \text{ GeV for } m_{\nu} = 4.8 \text{ MeV},$$
$$m_{\nu'} = 4.6 \text{ GeV for } m_{\nu} = 2.3 \text{ MeV},$$
$$m_{\nu'} = 25 \text{ GeV for } m_{\nu} = 95 \text{ MeV},$$
$$m_{\nu'} = 2.55 \times 10^4 \text{ TeV for } m_{\nu} = 1.27 \text{ GeV},$$
$$m_{\nu'} = 1.10 \times 10^4 \text{ TeV for } m_{\nu} = 4.18 \text{ GeV},$$
$$m_{\nu'} = 3.47 \times 10^5 \text{ TeV for } m_{\nu} = 173.5 \text{ GeV},$$
$$m_{\nu'} = 134 \text{ MeV for } m_{\nu} = 0.511 \text{ MeV},$$
$$m_{\nu'} = 27.82 \text{ GeV for } m_{\nu} = 105.7 \text{ MeV},$$
$$m_{\nu'} = 467.6 \text{ GeV for } m_{\nu} = 1.777 \text{ GeV}. \quad (67)$$

We hence can determine the mirror hadronic scale \cite{49},

$$\Lambda_{\text{QCD'}} = 1.28 \text{ GeV} \text{ for } \Lambda_{\text{QCD}} = 200 \text{ MeV}. \quad (68)$$

and then the mirror baryon masses \cite{53} and \cite{55},

$$m_{\nu'} = 10.5 \text{ GeV}, \quad m_{\nu} = 7.2 \text{ GeV}, \quad m_{\Delta'} = 5 \text{ GeV} \quad (69)$$

So, the mirror $\Delta'$ baryon is the dark matter particle. For the other parameter choice, the mirror proton or neutron can act as the dark matter particle.

By further taking the masses of the sterile neutrinos,

$$m_{\nu'} = \text{diag}(0.96 \text{ eV}, 95 \text{ MeV}, 100 \text{ MeV}), \quad (70)$$

as well as the VEV and Yukawa couplings of the Higgs bidoublet,

$$\langle \Sigma \rangle = 610 \text{ eV \ for \ } M_{\Sigma} = 10 \rho = 10^{13} \text{ GeV}, \quad (71)$$

$$f = \begin{bmatrix} \frac{i 1.09 \times 10^{-4}}{1} & \frac{i 1.87 \times 10^{-4}}{1} & \frac{-2.82 \times 10^{-5}}{1} \\ \frac{i 1.87 \times 10^{-4}}{1} & \frac{-5.26 \times 10^{-2}}{1} & \frac{-3.46 \times 10^{-2}}{1} \\ \frac{-2.82 \times 10^{-5}}{1} & \frac{-3.46 \times 10^{-2}}{1} & \frac{-3.85 \times 10^{-2}}{1} \end{bmatrix}, \quad (72)$$

the active neutrino masses \cite{20} can arrive at

$$(m_{\nu})_{ij} \simeq - f_{ij} f_{11} \frac{\langle \Sigma \rangle^2}{m_{\nu'}} - 2 f_{ij} \frac{\langle \Sigma \rangle \langle \phi_u \rangle}{m_{\nu'}}$$

$$\simeq \text{ eV} \begin{bmatrix} 0.00465 & 0.00793 & 0.00172 \\ 0.00793 & 0.0321 & 0.0211 \\ 0.00172 & 0.0211 & 0.0235 \end{bmatrix}, \quad (73)$$
to give the mass eigenvalues and mixing angles,

\[ m_{\nu_1} \simeq 0.001 \text{eV}, \]
\[ m_{\nu_2}^2 - m_{\nu_1}^2 \simeq 7.6 \times 10^{-5} \text{eV}^2, \]
\[ m_{\nu_3}^2 - m_{\nu_1}^2 \simeq 2.55 \times 10^{-3} \text{eV}^2, \]
\[ \sin^2 \theta_{12} \simeq 0.32, \sin^2 \theta_{23} \simeq 0.6, \sin^2 \theta_{13} \simeq 0.025, \]

(74)

which is consistent to the neutrino oscillation data.

The lightest sterile neutrino \( \nu'_1 \) can help us to fit the short baseline neutrino oscillation data if it has a mixing with the active neutrinos as below [3],

\[ |(U_{\nu\nu'})_{11}| = 0.15, \quad |(U_{\nu\nu'})_{21}| = 0.17. \]

(75)

This can be achieved by inputting

\[ |(fU^*_{\nu\nu'})_{11}| \simeq 2.36 \times 10^{-4}, \quad |(fU^*_{\nu\nu'})_{21}| \simeq 2.68 \times 10^{-4} \]  

(76)

in Eq. (25). As shown in Eq. (62), the sterile neutrino masses (70) will also be constrained by the neutrinoless double beta decay experiments,

\[ 9.6 \times 10^{-9} |(U_{\nu\nu'})_{11}|^2 + 0.95 |(U_{\nu\nu'})_{12}|^2 + |(U_{\nu\nu'})_{13}|^2 < 0.008, \]

(77)

To fulfill the the constraints (70) and (71), we can choose the mirror PMNS matrix \( U_{\nu\nu'} \) to have the zero CP phases and the following mixing angles,

\[ \sin \theta_{13}' = 0, \quad \sin \theta_{12}' \simeq 0.0897, \quad \sin \theta_{23}' \simeq 0.817, \]

(78)

and then give the required values of the elements

\[ (U_{\nu\nu'})_{11} = \cos \theta_{12}', \quad (U_{\nu\nu'})_{12} = \sin \theta_{12}', \quad (U_{\nu\nu'})_{13} = 0, \]
\[ (U_{\nu\nu'})_{21} = -\sin \theta_{12}' \cos \theta_{23}', \quad (U_{\nu\nu'})_{31} = \sin \theta_{12}' \sin \theta_{23}'. \]

(79)

The Yukawa couplings in the sterile neutrino masses [25] can be parameterized by [118]

\[ h = -i \langle \phi_u \rangle \sqrt{m_{\nu'}} \Omega^T \sqrt{M_N}, \]

(80)

with \( \Omega \) being an arbitrary orthogonal matrix. By taking the masses of the fermion singlets,

\[ M_N = \text{diag}\{10^7 \text{GeV}, 10^8 \text{GeV}, 10^9 \text{GeV}\}, \]

(81)

we can obtain the Yukawa couplings,

\[ h = -i U_{\nu\nu'} \left[ \begin{array}{ccc} 2.82 \cdot 10^{-7} \Omega_{11} & 8.90 \cdot 10^{-7} \Omega_{21} & 2.82 \cdot 10^{-6} \Omega_{31} \\ 2.80 \cdot 10^{-3} \Omega_{12} & 8.86 \cdot 10^{-3} \Omega_{22} & 2.80 \cdot 10^{-2} \Omega_{32} \\ 2.87 \cdot 10^{-3} \Omega_{13} & 9.09 \cdot 10^{-3} \Omega_{23} & 2.87 \cdot 10^{-2} \Omega_{33} \end{array} \right]. \]

(82)

By further

\[ |\Omega_{11}|, |\Omega_{21}| \ll 1, \quad |\Omega_{31}| \simeq 1, \]

we can get the out-of-equilibrium parameter [43],

\[ K_{N_1} \simeq 0.016, \]

(84)

and the CP asymmetry [38],

\[ \varepsilon_{N_1} \simeq 8.21 \times 10^{-7} \text{Re}(\Omega_{31}) \text{Im}(\Omega_{31}). \]

(85)

So, the final baryon asymmetry [47] can explain the observation [59],

\[ \eta_B = \eta'_B \simeq 0.888 \times 10^{-10} \quad \text{for} \quad \varepsilon_{N_1} \simeq -5.45 \times 10^{-8}. \]

(86)

VII. CONCLUSION

In this paper we have demonstrated a new mirror universe model, which contains three gauge-singlet Majorana fermions and an \([SU(2)_L \times SU(2)'_L]\)-bidoublet Higgs scalar in addition to the \(SU(3)_c \times SU(2)_L \times U(1)_Y\) ordinary sector and its \(SU(3)'_c \times SU(2)'_L \times U(1)'_Y\) mirror partner. In our model, the mirror sterile neutrino masses can have a form of the canonical seesaw, while the ordinary active neutrino masses can have a form of the double and linear seesaw. The mixing between the active and sterile neutrinos can also be seesaw-suppressed. Two sterile neutrinos can be above the MeV scale to avoid the BBN constraint, while the third sterile neutrino can be at the eV scale to fit the short baseline neutrino oscillation data. An ordinary lepton asymmetry and an equal mirror lepton asymmetry can be simultaneously produced from the decays of the fermion singlets. The baryonic and dark matter asymmetries then can equal each other since the ordinary and mirror sphaleron processes have a same efficiency of lepton-to-baryon conversion. Consequently, the lightest mirror baryon should have a mass around 5 GeV to serve as the dark matter particle. The \(U(1)_Y\) and \(U(1)'_Y\) kinetic mixing can mediate a testable dark matter scattering.
[91] K.Y. Choi, E.J. Chun, and C.S. Shin, arXiv:1211.5409 [hep-ph].
[92] J. Bramente and K. Fukushima, and J. Kumar, arXiv:1301.6811 [hep-ph].
[93] N.F. Bell, A. Melatos, and K. Petraki, arXiv:1301.0036 [hep-ph].
[94] M. Fukugita and T. Yanagida, Phys. Lett. B 174, 45 (1986).
[95] R.N. Mohapatra and X. Zhang, Phys. Rev. D 46, 5331 (1992); E. Ma and U. Sarkar, Phys. Rev. Lett. 80, 5716 (1998).
[96] M. Flanz, E.A. Paschos, and U. Sarkar, Phys. Lett. B 345, 248 (1995); M. Flanz, E.A. Paschos, U. Sarkar, and J. Weiss, Phys. Lett. B 389, 693 (1996); L. Covi, E. Roulet, and F. Vissani, Phys. Lett. B 384, 169 (1996); A. Pilaftsis, Phys. Rev. D 56, 5431 (1997).
[97] T. Hambye, E. Ma, and U. Sarkar, Nucl. Phys. B 602, 23 (2001).
[98] S. Davidson and A. Ibarra, Phys. Lett. B 535, 25 (2002); W. Buchmüller, P. Di Bari, and M. Plümacher, Nucl. Phys. B 665, 445 (2003).
[99] T. Hambye and G. Senjanović, Phys. Lett. B 582, 73 (2004); S. Antusch and S.F. King, Phys. Lett. B 597, 199 (2004).
[100] T. Hambye, Y. Lin, A. Notari, M. Papucci, A. Strumia, Phys. Rev. D 80, 095002 (2009).
[101] P.A.R. Ade et al., (Planck Collaboration), arXiv:1303.5076 [astro-ph.CO].
[102] J.A. Casas and A. Ibarra, Nucl. Phys. B 618, 171 (2001).

[99] T. Hambye and G. Senjanović, Phys. Lett. B 582, 73 (2004); S. Antusch and S.F. King, Phys. Lett. B 597, 199 (2004).
[100] T. Hambye, Y. Lin, A. Notari, M. Papucci, A. Strumia, Nucl. Phys. B 695, 169 (2004).
[101] T. Hambye, M. Raidal, and A. Strumia, Phys. Lett. B 632, 667 (2006).
[102] S. Davidson, E. Nardi, and Y. Nir, Phys. Rept. 466, 105 (2008).
[103] S. Blanchet and P. Di Bari, Nucl. Phys. B 807, 155 (2009).
[104] P.H. Gu, Phys. Lett. B 713, 485 (2012).
[105] V.A. Kuzmin, V.A. Rubakov, and M.E. Shaposhnikov, Phys. Lett. B 155, 36 (1985).
[106] I.F. Ginzburg and K.A. Kanishev, Phys. Rev. D 76, 095013 (2007).
[107] M. Baumgart, C. Cheung, J.T. Runderman, L.T. Wang, and I. Yavin, JHEP 0904, 014 (2009).
[108] R. Foot and X.G. He, Phys. Lett. B 267, 509 (1991).
[109] K. Nakamura et al., (Particle Data Group), J. Phys. G 37, 075021 (2010).
[110] Z. Maki, M. Nakagawa, and S. Sakata, Prog. Theor. Phys. 28, 870 (1962); B. Pontecorvo, Sov. Phys. JETP 26, 984 (1968), Zh. Eksp. Teor. Fiz. 53, 1717 (1967).
[111] F. Bezrukov, H. Hettmansperger, and M. Lindner, Phys. Rev. D 81, 085032 (2010).
[112] E.W. Kolb and M.S. Turner, The Early Universe, Addison-Wesley, 1990.
[113] M. Fukugita and T. Yanagida, Phys. Rev. D 42, 1285 (1990).
[114] J. Angle et al., (XENON10 Collaboration), Phys. Rev. Lett. 107, 051301 (2011).
[115] P.A.R. Ade et al., (Planck Collaboration), arXiv:1303.5076 [astro-ph.CO].
[116] M. Pospelov, Phys. Rev. D 80, 095002 (2009).
[117] For recent reviews, see W. Rodejohann, Int. J. Mod. Phys. E 20, 1833 (2011).
[118] J.A. Casas and A. Ibarra, Nucl. Phys. B 618, 171 (2001).