A Left-Right Symmetric Model for Neutrino Masses, Baryon Asymmetry and Dark Matter

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In the left-right symmetric models without bi-doublet Higgs scalars, the standard model fermions can obtain masses by integrating out heavy charged singlet fermions. We find the decays of heavy neutral singlet fermions, responsible for generating small neutrino masses, can simultaneously produce a left-handed lepton asymmetry for baryon asymmetry and a relic density of right-handed neutrinos for dark matter. Benefited from the left-right symmetry, the properties of the dark matter can be related to the generation of the neutrino masses and the baryon asymmetry. We also indicate that the decays of the non-thermally produced right-handed neutrinos can explain the observed fluxes of 511 keV photons from the Galactic bulge.

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I. INTRODUCTION

The $SU(3)_c \times SU(2)_L \times U(1)_Y$ standard model (SM) has been tested to a very high accuracy, but it has been suffering big challenges from particle physics and cosmology. For example, it can’t explain the existence of non-baryonic dark matter, which has been confirmed by precisely cosmological observations [1],

$$\Omega_h h^2 = 0.1099 \pm 0.0062.$$ (1)

Here $h = 0.719^{+0.026}_{-0.025}$ [1]. So far we know little on the true identity of the dark matter although there have been many interesting candidates. The cosmological observations also indicates that the present universe doesn’t contain significant amount of baryonic antimatter. This baryon asymmetry again requires supplementing new ingredients to the existing theory. The density of the baryonic matter is measured by [1],

$$\Omega_B h^2 = 0.02273 \pm 0.00062,$$ (2)

which is intriguingly comparable to that of the dark matter. The coincidence between the dark and baryonic matter implies that they may have a specific relation [2] although their creation and evolution are usually understood by unrelated mechanisms. Furthermore, observations on solar, atmospheric, reactor and accelerator neutrino oscillations have established the phenomenon of massive and mixing neutrinos [3],

$$\Delta m_{21}^2 = 7.65^{+0.23}_{-0.20} \times 10^{-5} \text{eV}^2,$$ (3a)
$$\Delta m_{31}^2 = \pm 2.4^{+0.12}_{-0.11} \times 10^{-3} \text{eV}^2,$$ (3b)
$$\sin^2 \theta_{23} = 0.50^{+0.07}_{-0.06},$$ (3c)
$$\sin^2 \theta_{12} = 0.304^{+0.022}_{-0.016},$$ (3d)
$$\sin^2 \theta_{13} = 0.01^{+0.016}_{-0.011},$$ (3e)

which is well consistent with the cosmological limits on the sum of the neutrino masses [1],

$$\Sigma m_\nu < 1.3 \text{eV} \ (95\% \CL).$$ (4)

The smallness of the neutrino masses inspires a seesaw [4–8] extension of the SM. The seesaw models [4–9] also allow the attractive leptogenesis [10–15] mechanism for generating the observed baryon asymmetry.

On the other hand, new physics beyond the SM is also motivated by some theoretical considerations. For example, the SM accommodates the left- and right-handed fermions in different ways: the left-handed fermions transform as doublets but the right-handed ones transform as singlets. The unequal treatment on the left- and right-handed fermions means parity violation. If we start with a theory with left-right symmetry [16] at high energy, where the left- and right-handed fermions are both placed in doublets and the parity is conserved, the spontaneously left-right symmetry breaking can induce the parity violation as observed at low energy. The left-right symmetric models are based on the gauge group $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B–L}$. Within this context, the electric charge ($Q$) can be well defined in terms of lepton ($L$) and baryon ($B$) numbers,

$$Q = T_L + T_R + \frac{B - L}{2},$$ (5)

where $T_L$, $T_R$ and $B - L$ are the generators of $SU(2)_L$, $SU(2)_R$ and $U(1)_{B–L}$, respectively. This provides us a natural explanation for the choice of hypercharge charge in the SM. An important prediction of the left-right symmetry is the existence of right-handed neutrinos, which are trivial and hence absent in the SM. In the most popular left-right symmetric models with bi-doublet and triplet Higgs scalars [5], the right-handed neutrinos can naturally obtain heavy Majorana masses to realize the seesaw and the leptogenesis after the left-right symmetry breaking.

In this paper we consider the left-right symmetric models without bi-doublet Higgs scalars [17]. If the Higgs
In the present model without the Higgs bi-doublets, we introduce some heavy charged singlets to construct the left- and right-handed fermions in the SM with the masses as below, their decays can produce the observed baryon asymmetry through the leptogenesis and can give the right-handed neutrinos a desired relic density for the dark matter. In this left-right symmetric scenario, the properties of the dark matter are related to the parameters for generating the neutrino masses and the baryon asymmetry.

II. THE MODEL

The most general scalar potential of our model is

\[ V = -\mu_1^2 \sigma^2 - \mu_2^2 \left( \phi_L^d \phi_R + \phi_R^d \phi_L \right) + \mu_3 \sigma \left( \phi_L^d \phi_R + \phi_R^d \phi_L \right) + \lambda_1 \sigma^4 + \lambda_2 \left( \phi_L^d \phi_R^2 + \phi_R^d \phi_L^2 \right) + \lambda_3 \phi_L^d \phi_R^d \phi_L^d \phi_R^d + \lambda_4 \sigma^2 \left( \phi_L^d \phi_R^2 + \phi_R^d \phi_L^2 \right), \tag{6} \]

with \( \sigma \to -\sigma \) and \( \phi_L \to \phi_R \) under the parity symmetry. Here \( \lambda_{1,2} > 0, \lambda_3 > -2\lambda_2 \) and \( \lambda_4 > -2\sqrt{\lambda_1 \lambda_2} \) so that the potential is bounded from below. For appropriate parameter choice, it is easy to give the desired vacuum expectation values (VEVs),

\[ v_R = \langle \phi_R \rangle \gg v_L = \langle \phi_L \rangle \simeq 174 \text{ GeV}. \tag{7} \]

In the present model without the Higgs bi-doublets, the charged gauge bosons \( W_L^\pm \) and \( W_R^\pm \) are mass eigenstates \[{18}\],

\[ m_{W_L} = \frac{1}{\sqrt{2}} g v_L, \quad m_{W_R} = \frac{1}{\sqrt{2}} g v_R. \tag{8} \]

The left- and right-handed fermions in the SM with the right-handed neutrinos are placed in doublets for each family. For generating the masses of the SM fermions, we introduce some heavy charged singlets to construct the Yukawa couplings to the Higgs and fermion doublets so that we can derive the SM yukawa couplings by integrating out these singlets \[{17}\],

\[ \mathcal{L} \supset -y_D \left( \bar{q}_L \phi_L D_R + \bar{q}_R \phi_R D_L \right) - M_D \tilde{D}_L D_R - y_U \left( \bar{q}_L \phi_L U_R + \bar{q}_R \phi_R U_L \right) - M_U \tilde{U}_L U_R - y_E \left( \bar{l}_L \phi_L E_R + \bar{l}_R \phi_R E_L \right) - M_E \tilde{E}_L E_R + \text{H.c.,} \tag{9} \]

where the SM Yukawa couplings are given by

\[ y_d = -y_D^L M_D^{-1} y_D^R, \tag{10a} \]

\[ y_u = -y_U^L M_U^{-1} y_U^R, \tag{10b} \]

\[ y_e = -y_E^L M_E^{-1} y_E^R. \tag{10c} \]

Here we have chosen the base where the mass matrices \( M_{D,U,E} \) are real and diagonal. Note the mass matrices \( M_{D,U,E} \) now are not hermitian as a result of the spontaneous D-parity violation, i.e.

\[ M_{D,U,E} = M_{D,U,E}^0 + h_{D,U,E} \langle \sigma \rangle \tag{11} \]

with \( M_{D,U,E}^0 = M_{D,U,E} \) while \( h_{D,U,E} = -h_{D,U,E} \). If the Yukawa couplings \( h_{D,U,E} \) vanish, the left-handed yukawa couplings \( y_{D,U,E}^L \) should be equal to the right-handed ones \( y_{D,U,E}^R \) as \( M_{D,U,E} \) are hermitian matrices. Since the Yukawa coupling of top quark is very close to 1, Eq. \[10b\] will constrain \( M_T \lesssim v_R \) for \( y_T < \sqrt{4\pi} \). For simplicity, we shall conveniently assume \( M_D = M_U = M_E = v_R \) to make the heavy charged singlets decoupled from the following discussions.

In the neutrino sector, we consider the left- and right-handed neutral singlets with the Yukawa couplings and the masses as below,

\[ \mathcal{L} \supset -y_N \left( \bar{l}_L \phi_L N_R + \bar{l}_R \phi_R N_L \right) - M_N^0 \tilde{N}_L N_R - \frac{1}{2} M_N^M \left( \tilde{N}_L^c N_L + \tilde{N}_R^c N_R \right) + \text{H.c.} \tag{12} \]

In the presence of the spontaneous D-parity violation, the Dirac mass matrix \( M_N^D \) is not hermitian, i.e.

\[ M_N^D = M_N^0 + h_N \langle \sigma \rangle \tag{13} \]

with \( M_N^0 = M_N^0 \) while \( h_N = -h_N \). Furthermore, we have forbidden the Yukawa couplings \( \bar{l}_L \phi_L N_R, \bar{l}_R \phi_R N_L \) and their CP conjugates. This can be achieved by imposing a discrete, global or local symmetry. For example, we consider a \( U(1)_X \) local symmetry under which \( D_{L,R}, U_{L,R}, E_{L,R}, N_{L,R}, \phi_{L,R} \) carry a quantum number \( X = 1 \). Clearly, this \( U(1)_X \) is free of gauge anomaly. In this context, the Yukawa couplings and the Dirac mass terms in Eqs. \[9\] and \[12\] are allowed while the Majorana mass terms in Eq. \[12\] are forbidden. To break this
masses in Eq. (12) can be given by

\[
\mathcal{L} \supset - \frac{1}{2} f_N \left( \delta \tilde{N}_L N_L + \bar{\nu} \tilde{N}_R N_R \right) + \text{H.c.} \tag{14}
\]

Through the above Yukawa interactions, the Majorana masses would contain a Dirac mass term and two Majorana ones,

\[
\mathcal{L} \supset - \frac{1}{2} \bar{\nu}_L m_L \nu_L^c - \frac{1}{2} \bar{\nu}_R m_R \nu_R^c - \bar{\nu}_L m_D \nu_R + \text{H.c.} \tag{16}
\]

with

\[
\begin{align*}
m_L &= -y_N \frac{1}{M_N^2} U v_L^c, \\
m_R &= -y_N \frac{1}{M_N^2} U v_R^c, \\
m_D &= y_N \frac{1}{M_N^2} (M_D^T)^T \nu_R v_L.
\end{align*}
\tag{17}
\]

Here we have assumed

\[
M_N^M \gg M_N^D, y_N v_R, y_N v_L,
\tag{18}
\]

and then have defined the left- and right-handed Majorana fermions,

\[
\begin{align*}
L_i &= N_{L_i} + N_{L_i}^c, \\
R_i &= N_{R_i} + N_{R_i}^c,
\end{align*}
\tag{19}
\]

by choosing the base where the Majorana mass matrix \( M_N^M \) is real and diagonal,

\[
M_N^M = \text{diag} \{ M_1, M_2, M_3 \} \approx M. \tag{20}
\]

Clearly, the right-handed neutrinos will give their left-handed partners an additional Majorana mass term through the seesaw since their Dirac masses are not vanishing. This gift is negligible,

\[
\delta m_L = -m_D \frac{1}{m_R} m_D^T = O \left( \frac{M_D^2}{M_N^2} \right)^2 m_L \ll m_L. \tag{21}
\]

Therefore, we can well define the left- and right-handed Majorana neutrinos,

\[
\begin{align*}
\nu &= \nu_L + \nu_R^c, \\
\chi &= \nu_R + \nu_R^c.
\end{align*}
\tag{22}
\]

which have mass matrices with a same texture,

\[
\begin{align*}
m_{\nu} &= m_L = U^\dagger \text{diag} \{ m_1, m_2, m_3 \} U^*, \\
m_{\chi} &= m_R = U^\dagger \text{diag} \{ m_1, m_2, m_3 \} U^* v_R^2 v_L.
\end{align*}
\tag{23}
\]

Here \( U \) is the MNS lepton flavor mixing matrix \( [20] \).

### III. Baryon Asymmetry

We now discuss the physics below the scale \( M_N^M \), which is assumed much smaller than the left-right symmetry breaking scale \( v_R \), i.e.

\[
M_N^M \ll v_R. \tag{24}
\]

We thus integrate out the right-handed physical Higgs boson and then derive the following Lagrangian,

\[
\begin{align*}
\mathcal{L} \supset - y_N \bar{l}_L \phi_L L^c + y_N \frac{1}{M} M_D^T \nu_R^c L^c + \frac{\lambda_1}{4} y_N \phi_L \phi_L \nu_R^c L^c + \text{H.c.} \\
&- \frac{1}{2} \left( \bar{l}_L + \nu_R \right) + \frac{1}{2} \left( \bar{l}_L + \nu_R \right)^c.
\end{align*}
\tag{25}
\]

Here the dimension-5 operator \( \phi_L^T \phi_L \nu_R^c L^c \) and its CP conjugate have been omitted since they are sufficiently suppressed in comparison with the couplings in \( [20] \).

Clearly, Eq. (25) can accommodate the standard leptogenesis scenario, where the 2-body decays of the Majorana fermions into the SM lepton and Higgs doublets can generate a lepton asymmetry if CP is not conserved. This lepton asymmetry can be partially converted to a baryon asymmetry through sphaleron \( [21] \). So, the absence of baryonic antimatter is well explained. The final baryon asymmetry can be described by

\[
\frac{n_b}{s} = c \varepsilon_{R_j} \frac{\kappa_{R_i} \varepsilon_{R_i} + \kappa_L \varepsilon_{L_j}}{g_s}. \tag{26}
\]

Here \( c = -28/79 \) \([22]\) is the sphaleron induced lepton-to-baryon conversion coefficient, \( g_s \simeq 112 \) is the relativistic degrees of freedom (the SM fields plus three right-handed neutrinos), \( \kappa_{R_j}(y_N, M) \) and \( \kappa_{L_j}(y_N, M) \) are wash out factors, \( \varepsilon_{R_j} (\varepsilon_{L_j}) \) is the lepton asymmetry induced by the decays of a \( R_j (L_j) \) and is calculated at one-loop order,

\[
\begin{align*}
\varepsilon_{R_j}(y_N, M) &= \frac{\Gamma_{R_j \rightarrow l_L^+} \phi_L - \Gamma_{R_j \rightarrow l_L^+} \phi_L^c}{\Gamma_{R_j}}, \\
\varepsilon_{L_j}(y_N, M) &= \frac{\Gamma_{L_j \rightarrow l_L^+} \phi_L - \Gamma_{L_j \rightarrow l_L^+} \phi_L^c}{\Gamma_{L_j}}.
\end{align*}
\tag{27}
\]
where $\Gamma_{R_j}$ and $\Gamma_{L_j}$ are the decay width at tree level,
\[
\begin{align*}
\Gamma_{R_j} & = \Gamma_{R_j} \rightarrow l_j \phi_L + \Gamma_{R_j} \rightarrow \bar{l}_j \phi_L^* \\
& = \frac{1}{8\pi} \left( y_N y_N \right)_{jj} M_j, \\
\Gamma_{L_j} & = \Gamma_{L_j} \rightarrow l_j \phi_L + \Gamma_{L_j} \rightarrow \bar{l}_j \phi_L^* + \Gamma_{L_j} \rightarrow \nu \phi_L^* \\
& = \frac{1}{8\pi} \left[ M_N^2 \left( \frac{y_N}{M_N} \right)_{jj} + M_D^2 \right] M_j \\
& \approx \frac{\lambda_3^2}{2 \sqrt{2} V^2} \left( y_N y_N \right)_{jj} M_j \\
& \approx \mathcal{O} \left( \frac{M_D^2}{M_N^2} \right)^2 \frac{1}{8\pi} \left( y_N y_N \right)_{jj} M_j. 
\end{align*}
\]
(28a)

For a given relic density, we can determine the masses of the right-handed neutrinos by the observed baryon asymmetry,
\[
m_{\chi_i} = m_N \frac{c \Sigma_j \left( \kappa_{R_j} + \kappa_{L_j} \right) \varepsilon_j \Omega_{\chi_i}}{\sum_j \kappa_{R_j} \varepsilon_j \Omega_{\chi_i}}. 
\]
(35)

Here $m_N \sim 940$ MeV is the nucleon mass. On the other hand, as shown in Eq. (23), the mass matrix of the right-handed neutrinos has a same texture with that of the left-handed neutrinos. For example, let’s consider the case with the degenerate neutrino masses,
\[
m_1 \simeq m_2 \simeq m_3 = 0.05 \text{eV},
\]
(36)

which immediately yields
\[
m_{\chi_1} \simeq m_{\chi_2} \simeq m_{\chi_3} = 0.05 \text{eV} \frac{v_R^2}{v_L}. 
\]
(37)

With further assumptions,
\[
\kappa_{L_j} \simeq \kappa_{R_j} \simeq \kappa_{L_j} = 0.1, 
\]
(38)

which is true in the weak washout region, and
\[
\frac{\lambda_3}{\lambda_2} = 3.3 \times 10^{-4}, \quad \frac{M_D}{M_N} = 1.9 \times 10^{-7}, \quad \frac{M_j}{v_R} = 2.1 \times 10^{-4}, 
\]
(39)

we can obtain the dark matter masses
\[
m_{\chi_{1,2,3}} \simeq 1.3 \text{MeV}. 
\]
(40)

Accordingly, the left-right symmetry breaking scale can be given by
\[
v_R = v_L \sqrt{m_{\chi_1}/m_i} \simeq 5.1 \times 10^3 v_L. 
\]
(41)

The above breaking scale points to the necessity of the resonant [12, 13] leptogenesis because of the parameter choice [39], i.e. $M_j \simeq 186$ GeV. It also indicates that the gauge interactions of the right-handed neutrinos decouple at a high temperature $\sim (\frac{v_R}{v_L})^2 \mathcal{O}(\text{MeV}) = \mathcal{O}(100 \text{GeV})$ so that the non-thermally [23] produced right-handed neutrinos can successfully explain the dark matter relic density.

We should keep in mind that the right-handed neutrinos mix with their left-handed partners due to the Dirac mass term in Eq. (10) so that they will decay at tree level and loop orders. In absence of the mixing between $W_L^\pm$ and $W_R^\pm$ (cf. [38]), the decay width should be [24, 25]
\[
\begin{align*}
\Gamma_{\chi_i \rightarrow \nu \nu} & = \frac{G_F^2}{384 \pi^3} \sin^2 (2\theta_i) \frac{m_i^5}{m_\chi_i}, \\
\Gamma_{\chi_i \rightarrow \nu e^+ e^-} & = \frac{5 G_F^2}{3072 \pi^3} \sin^2 (2\theta_i) \frac{m_i^5}{m_\chi_i}, \\
\Gamma_{\chi_i \rightarrow \nu \gamma} & = \frac{9 \alpha G_F^2}{1024 \pi^4} \sin^2 (2\theta_i) \frac{m_i^5}{m_\chi_i}. 
\end{align*}
\]
(42a)

IV. DARK MATTER

The 3-body decays of $L$ don’t have significant contributions to the total decay width and the lepton asymmetry, but it is sufficient to produce abundant right-handed neutrinos for the dark matter relic density. We now clarify this interesting scenario in details. It is easy to read the relic density,
\[
\frac{n_{\chi_i}}{s} = \sum_j \kappa'_{L_j} \frac{B_{R_{ji}}}{g_*},
\]
(32)

where $\kappa'_{L_j}(y_N, M)$ is a wash out factor and
\[
B_{R_{ji}} = \frac{\Gamma_{L_j \rightarrow \chi_i \phi_L^*}}{\Gamma_{L_j}} = \mathcal{O} \left[ \left( \frac{M_D}{M_N} \right)^2 \frac{M_j^2}{v_R^2} \frac{\lambda_3^2}{2 \sqrt{2} V^2} \left( y_N y_N \right)_{jj} \right]^2
\]
(33)

is the branching ratio. From the baryon asymmetry [20] and the relic density [22], we can perform the following ratio,
\[
\Sigma_j \kappa_{L_j} B_{R_{ji}} m_{\chi_i} : c \Sigma_j \left( \kappa_{R_j} + \kappa_{L_j} \right) \varepsilon_j m_N = \Omega_{\chi_i} : \Omega_b
\]
(34)
Here \( \alpha \) and \( G_F \), respectively, are the fine-structure constant and the Fermi constant, \( \theta_i^2 \) is the mixing angle defined by

\[
\theta_i^2 = \left( \frac{m_D m_P}{m_{\chi_i}} \right)_{ii} \propto \mathcal{O} \left[ \left( \frac{M_D}{M_N} \right)^2 \right] \frac{v^2}{v_R^2}.
\]

(43)

With the previous parameter choice, we can determine the mixing angle to be \( \theta_i^2 \approx 10^{-22} \). Therefore the decay into the electron-positron pairs (42b) can provide a natural explanation for the flux of 511 keV photons from the galactic bulge observed by INTEGRAL satellite,

\[
\Phi_{\text{gal}} / \Phi_{\text{exp}} \approx \sum_i \theta_i^2 / 10^{-22} \left( \frac{m_{\chi_i}}{1.3 \text{MeV}} \right)^4 \frac{\Omega_X}{\Omega_{\chi}}.
\]

(44)

It is easy to check that our scenario is consistent with other astrophysical and cosmological constraints. Alternatively, we may explain the observed cosmic positron/electron excess, which is probably from continuum distribution of pulsars, by fine tuning the parameters.

V. CONCLUSION

In summary we have shown the dark matter can be well determined by the neutrino masses and the baryon asymmetry in the left-right symmetric model with doublet and singlet fields. In this model, the SM fermions obtain masses by integrating out charged singlet fermions. In the neutrino sector, the right-handed neutral fermions, associated with the left-right-handed Higgs doublet, can generate the left-right-handed neutrino masses through the seesaw. The mass matrices of the left- and right-handed neutrinos have a same structure as a result of the left-right symmetry. The neutral singlets are also responsible for the baryon asymmetry and the dark matter. Specifically their 2-body decays can produce a desired lepton asymmetry in the left-handed leptons and then the observed baryon asymmetry can be realized by the sphaleron induced lepton-to-baryon conversion. At the same time, the right-handed neutrinos can serve as the dark matter as they have a right relic density from the 3-body decays of the neutral singlets. The decays of these non-thermally produced right-handed neutrinos can easily induce the observed fluxes of 511 keV photons from the Galactic bulge. The attractive feature of our scenario is that the left-right symmetry can connect the properties of the dark matter to the neutrino masses and the baryon asymmetry.

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