Neutrino masses, dark matter and leptogenesis with $U(1)_{B-L}$
gauge symmetry

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Abstract

We propose a model with an $U(1)_{B-L}$ gauge symmetry, in which small neutrino masses, dark matter and the matter-antimatter asymmetry in the Universe can be simultaneously explained. In particular, the neutrino masses are generated radiatively, while the matter-antimatter asymmetry is led by the leptogenesis mechanism, at TeV scale. We also explore allowed regions of the model parameters and discuss some phenomenological effects including lepton flavor violating processes.

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I. INTRODUCTION

Radiative neutrino mass generation is one of the most promising candidates to naturally explain the small mass scales of active neutrinos. Some models to realize this type of the approach can also accommodate dark matter (DM). Normally, when one considers the DM candidate in a theory, an additional symmetry such as $Z_2$ is imposed in order to stabilize it. The representative model has been shown in ref. [1]. In recent years, a lot of applications have been presented in the literature. In particular, a model in ref. [2] with introducing $Z_2 \times Z_2$ symmetry and two inert isospin doublet bosons was proposed to understand the cosmic ray anomaly [3] by a decaying fermionic DM. In this model, the Baryon Asymmetry of the Universe (BAU) is understood via the leptogenesis mechanism within the TeV scale.

Our paper extends the study in ref. [2] by having a gauged $U(1)_{B-L}$ symmetry instead of $Z_2 \times Z_2$, in which three right-handed neutrinos are naturally introduced as a usual model with the $U(1)_{B-L}$ symmetry. However, their charge assignments are taken with very unique manner, i.e. $-4$, $-4$ and $5$ for three right-handed neutrinos [4, 5], respectively.\footnote{Several applications along this ideas can be found in refs. [6–8]} It suggests that the first two right-handed neutrinos can contribute to the active neutrinos, resulting in two massive and one massless neutrinos, while the third right-handed one can be a good DM candidate, even though two kinds of new bosons with nonzero $B - L$ charges have to be added to give the masses of right-handed neutrinos. As a result, the stabilized symmetry of DM ($Z_2$) is induced as a remnant symmetry after the spontaneous symmetry breaking (SSB) of $U(1)_{B-L}$.

This paper is organized as follows. In Sec. II, we first set up our model. We then discuss the active neutrinos, lepton flavor violating (LFV) processes, leptogenesis and dark matter. In Sec. III, we give the numerical analysis to explore the allowed parameter space of the model. We conclude in Sec. IV.
II. MODEL SETUP AND PHENOMENOLOGIES

First of all, we impose an additional $U(1)_{B-L}$ gauge symmetry with three right-handed neutral fermions $N_{R_i} (i = 1, 2, 3)$ in the standard model (SM), where the right-handed neutrinos have $U(1)_{B-L}$ charges of $-4, -4$ and $5$, respectively. Consequently, all the anomalies to be considered are the triangular $U(1)^3_{B-L}$ and mixed gravity-gauged $U(1)_{B-L}$ ones, which are found to be zero [4, 5], due to the uniqueness of the charge assignments in the SM [9]. On the other hand, even when we include two types of isospin singlet bosons $\varphi_1$ and $\varphi_2$ to acquire nonzero Majorana masses after the SSB of $U(1)_{B-L}$, one cannot find active neutrino masses due to the absence of the Yukawa term $\bar{L}_L H N_R$. Thus, we introduce an isospin singlet and doublet inert bosons $s$ and $\eta$ with nonzero $U(1)_{B-L}$ charges, so that neutrino masses are radiatively generated at one-loop level. Also the stability of DM is assured by a remnant $Z_2$ symmetry after the SSB of $U(1)_{B-L}$. Field contents and their assignments for fermions and bosons are given in Table I, respectively. The renormalizable Lagrangian for lepton sector and Higgs potential are given by

$$-\mathcal{L}_L = y_{\epsilon_L} \bar{L}_L \epsilon_R H + (y_{\zeta})_{ai} \bar{L}_L \tilde{\zeta} N_{R_i} + \frac{y_{N_i}}{2} \bar{N}_{R_i}^C N_{R_i} \varphi_8 + \frac{y_{N_3}}{2} \bar{N}_{R_3}^C N_{R_3} \varphi_{10}^* + \text{h.c.},$$

$$V = \frac{\mu_H^2}{2} |H|^2 + \frac{\mu_n^2}{2} |\eta|^2 + \frac{\mu_\zeta^2}{2} |\zeta|^2 + \frac{\mu_{\varphi_8}^2}{2} |\varphi_8|^2 + \frac{\mu_{\varphi_{10}}^2}{2} |\varphi_{10}|^2 + \frac{\lambda_n}{2} [(H^\dagger \zeta)(\eta^\dagger \zeta) + \text{h.c.}]$$

$$+ \frac{\lambda_H}{4} |H|^4 + \frac{\lambda_n}{4} |\eta|^4 + \frac{\lambda_\zeta}{4} |\zeta|^4 + \frac{\lambda_{\varphi_8}}{4} |\varphi_8|^4 + \frac{\lambda_{\varphi_{10}}}{4} |\varphi_{10}|^4 + \lambda_{H\eta} |H|^2 |\eta|^2 + \lambda_{H\zeta} |H|^2 |\zeta|^2 + \lambda_{H\varphi_8} |H|^2 |\varphi_8|^2 + \lambda_{H\varphi_{10}} |H|^2 |\varphi_{10}|^2$$

$$+ \lambda_{n\eta} |\eta|^2 |\varphi_8|^2 + \lambda_{n\varphi_{10}} |\eta|^2 |\varphi_{10}|^2 + \lambda_{\zeta\varphi_8} |\zeta|^2 |\varphi_8|^2 + \lambda_{\zeta\varphi_{10}} |\zeta|^2 |\varphi_{10}|^2,$$

respectively, where $\tilde{\zeta} \equiv (i\sigma_2)\zeta^*$ with $\sigma_2$ being the second Pauli matrix, and $a(i)$ runs over 1 to 3(2).

| TABLE I: Field contents of fermions and their charge assignments under $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$, where $i = 1, 2$. |
|-----------------|-----|-----|-----|-----|---|---|
|                 | $Q_L$ | $u_R$ | $d_R$ | $L_L$ | $e_R$ | $N_{R_i}$ | $N_{R_3}$ |
| $SU(3)_C$       | 3    | 3    | 3    | 1    | 1   | 1   | 1          |
| $SU(2)_L$       | 2    | 1    | 1    | 2    | 1   | 1   | 1          |
| $U(1)_Y$        | $\frac{1}{6}$ | $\frac{2}{3}$ | $-\frac{1}{3}$ | $-\frac{1}{2}$ | -1 | 0 | 0 |
| $U(1)_{B-L}$    | $\frac{1}{3}$ | $\frac{1}{3}$ | $\frac{1}{3}$ | $-1$ | $-1$ | -4 | 5 |
In the scalar sector, the scalar fields are parameterized as

\[ H = \begin{pmatrix} w^+ \\
\frac{v+h+iz}{\sqrt{2}} \end{pmatrix}, \quad \eta = \begin{pmatrix} \eta^+ \\
\frac{\nu_0+\eta_R+iv}{\sqrt{2}} \end{pmatrix}, \quad \zeta = \begin{pmatrix} \zeta^+ \\
\frac{\zeta_R+i\zeta_I}{\sqrt{2}} \end{pmatrix}, \quad \varphi_i = \frac{v_{\varphi_j} + \varphi_{Rj} + iz_{\varphi_j}}{\sqrt{2}}, \quad (j = 8, 10), \]

where each of the lightest states (=massless states) of \( w^\pm, \eta^\pm, \eta_I \), and \( \varphi_{8,10} \) is absorbed by the SM gauge bosons of \( W^\pm \) and \( Z \), and the \( B-L \) gauge boson of \( Z' \), induced after the SSB. Here, the structure of \( w^\pm - \eta^\pm \) is same as one of the two Higgs doublet models \(^{10}\). Inserting tadpole conditions, the CP even matrix with 5 by 5 in the basis of \((h, \eta_R, \varphi_{R8}, \varphi_{R10})^T\) is defined by \( M_R \), which is diagonalized by the orthogonal matrix \( O_R \) as \( m_{hi}^2 = O_RM_R^TO_R^T, (i = 1 - 5) \), where \( h_i \) is the mass eigenstate and \( m_{hi} \) its mass eigenvalue. Furthermore, the SM Higgs is defined by \( h_{SM} \equiv h_1 \) with \( m_{h_{SM}} \equiv m_{h_1} = 125 \text{ GeV} \). In the inert sector, the mass eigenstates of \( \eta_{R(i)} \) are given by \(^{11}\)

\[ m_{\zeta_R}^2 = M_\zeta^2 + \lambda_0 v_\eta, \quad m_{\zeta_i}^2 = M_\zeta^2 - \lambda_0 v_\eta, \quad M_\zeta^2 = \mu_\zeta^2 + \lambda_\varphi^2 v_{\varphi_8}^2 + \lambda_{\varphi_{10}} v_{\varphi_{10}}^2 + (\lambda_\eta + \lambda_\eta') v_\eta^2. \]

Here, we will briefly discuss the breaking scale of \( U(1)_{B-L} \). Because of our large number of \( B-L \) charge assignments for bosons \( \varphi_8 \) and \( \varphi_{10} \) in Table\(^{11}\) that cause the SSB of \( U(1)_{B-L} \) in order to give masses of the right-handed neutrinos with \( B-L \) charges \((-4, -4, 5)\), our theory can really be within the TeV scale. The breaking scale could be evaluated by the mass of the \( B-L \) gauge boson. Once we fix the gauge coupling of \( B-L \) to be \( O(1) \), its typical mass is greater than 7 TeV from the LEP constraint \(^{12}\). On the other hand, its theoretical mass in our model can be given by \( m_{Z'} = \sqrt{(8v_{\varphi_8})^2 + (10v_{\varphi_{10}})^2} \). Even assuming \( v_{\varphi_8} \gg v_{\varphi_{10}} \), the typical breaking scale of \( B-L \) can be less than 1 TeV!

**TABLE II:** Field contents of bosons and their charge assignments under \( SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L} \)

| Bosons | \( H \) | \( \zeta \) | \( \eta \) | \( \varphi_8 \) | \( \varphi_{10} \) |
|--------|---------|---------|---------|---------|---------|
| \( SU(3)_C \) | 1 | 1 | 1 | 1 | 1 |
| \( SU(2)_L \) | 2 | 2 | 2 | 1 | 1 |
| \( U(1)_Y \) | \( \frac{1}{2} \) | \( \frac{1}{2} \) | \( \frac{1}{2} \) | 0 | 0 |
| \( U(1)_{B-L} \) | 0 | -3 | -6 | 8 | 10 |
A. Neutrino masses

Before discussing the neutrino sector, several remarks are as follows: Since $y_{\ell}$ and $y_N$ can be diagonal without loss of generality in Eq. (2), we define $m_{\ell_a} = y_{\ell_a} v / \sqrt{2}$ ($a = e, \mu, \tau$) $M_{N_i} = y_{N_i} v_{\varphi_b} / \sqrt{2}$ ($i = 1, 2$) after the electroweak symmetry breaking and the $U(1)_{B-L}$ symmetry breaking, where $m_\ell$ and $M_N$ are the masses.

Then, the neutrino mass matrix is induced at one-loop level as shown in Fig. 1, and its mass-insertion-approximation form is given by

$$\langle M_\nu \rangle_{\alpha\beta} = \frac{1}{16\pi^2} \sum_{i=1,2} (y_\zeta)_{\alpha i} (y_\zeta^T)_{i\beta} M_{N_i} \left( \frac{m_{\zeta R}^2}{m_{\zeta R}^2 - M_{N_i}^2} \ln \left[ \frac{m_{\zeta R}^2}{M_{N_i}^2} \right] - \frac{m_{\zeta I}^2}{m_{\zeta I}^2 - M_{N_i}^2} \ln \left[ \frac{m_{\zeta I}^2}{M_{N_i}^2} \right] \right).$$

We note that the Casas-Ibarra parametrization is a convenient method to achieve the numerical analysis [13]. Once we define $D_\nu \equiv U_{MNS} M_\nu U_{MNS}^T$, $y_\zeta$ can be replaced by observables with several arbitral parameters given by

$$y_\zeta = U_{MNS}^\dagger D_\nu^{1/2} O R^{-1/2},$$

$$R \equiv \frac{1}{16\pi^2} \sum_{i=1,2} M_{N_i} \left( \frac{m_{\zeta R}^2}{m_{\zeta R}^2 - M_{N_i}^2} \ln \left[ \frac{m_{\zeta R}^2}{M_{N_i}^2} \right] - \frac{m_{\zeta I}^2}{m_{\zeta I}^2 - M_{N_i}^2} \ln \left[ \frac{m_{\zeta I}^2}{M_{N_i}^2} \right] \right),$$

where $U_{MNS}$ and $D_\nu$ are measured by neutrino oscillation experiments, and $O$ is an arbitral complex 3 by 2 rotation matrix with $OO^T = \text{Diag}(0, 1, 1)$ ($O^T O = 1_{2 \times 2}$), which can be parametrized by the following matrices for the normal hierarchy (NH) and inverted hierarchy

\[ \text{footnote}{\text{Notice here that our one-loop function is different from the Ma’s model [1], since we apply a mass insertion approximation method.}} \]
respectively, where $z$ can be complex. In our numerical analysis, we will use the global fit of the current neutrino oscillation data as the best fit values for NH and IH:

\[ s_{12}^2 = 0.304, \quad s_{23}^2 = 0.452, \quad s_{13}^2 = 0.0218, \quad \delta_{CP} = \frac{306}{180}\pi, \]

\[ (m_{\nu_1}, m_{\nu_2}, m_{\nu_3}) \approx (0, 8.66, 49.6) \text{ meV}, \]  

\[ (m_{\nu_1}, m_{\nu_2}, m_{\nu_3}) \approx (49.5, 50.2, 0) \text{ meV}, \]

where $s_{12,13,23}$ are the short-hand notations of $\sin \theta_{12,13,23}$ for three mixing angles of $U_{\text{MNS}}$, while two Majorana phases are taken to be zero.

**B. Lepton flavor violating processes**

The processes $\ell \to \ell' \gamma$ violate the lepton number, which are induced from the neutrino Yukawa couplings at one-loop level, and their forms are given by

\[
\text{BR}(\ell_\alpha \to \ell_\beta \gamma) = \frac{\alpha_{em} C_{\alpha \beta}}{768\pi G_F^2} \left| \sum_{i=1,2} (y_\zeta)^{\alpha_i} \frac{(y_\zeta^\dagger)^{\alpha_i}}{M_{N_i}^2} 2 + 3r_{N_i} - 6r_{N_i}^2 + r_{N_i}^3 + 6r_{N_i} \ln r_{N_i} \right|^2, \]

where $r_{N_i} \equiv m_{\zeta_i}^2/M_{N_i}^2$, $\alpha_{em} \approx 1/134$ is the fine-structure constant, $G_F \approx 1.17 \times 10^{-5} \text{ GeV}^{-2}$ is the Fermi constant, $C_{\mu e} \approx 1$, $C_{\tau e} \approx 0.1784$, and $C_{\tau \mu} \approx 0.1736$. Experimental upper bounds for the branching ratios of the LFV decays are found to be:

\[
\text{BR}(\ell_\mu \to \ell_e \gamma) \lesssim 4.2 \times 10^{-13}, \quad \text{BR}(\ell_\tau \to \ell_e \gamma) \lesssim 3.3 \times 10^{-8}, \quad \text{BR}(\ell_\tau \to \ell_\mu \gamma) \lesssim 4.4 \times 10^{-13}. \]

The stringent constraint comes from $\mu \to e \gamma$, and it roughly gives the upper limit $y_\zeta \lesssim \mathcal{O}(10^{-2})$ for $M_{N_i} = \mathcal{O}(100) \text{ GeV}$.

**C. Leptogenesis**

Here, we discuss the resonant leptogenesis mechanism, followed by those in ref. [2]. First of all, we expect the source of the CP asymmetry is induced from $N_1$ via the two-body decay,
with its decay width given by

$$\Gamma(N_1 \to e^\pm \zeta^{\mp}) = \frac{(y_\zeta^1 y_\zeta)_{11}}{16\pi} M_{N_1} \left(1 - \frac{m_{\zeta^\pm}^2}{M_{N_1}^2}\right)^2,$$

(13)

where $1 - m_{\zeta^\pm}^2/M_{N_1}^2 \sim 10^{-4}$ [17] is required in order to satisfy the out-of-equilibrium condition. When $3M_{N_1} < M_{N_2}$, the CP violating parameter $\epsilon$ is given by

$$\epsilon \simeq -\frac{3}{16\pi} \text{Im} \left[\frac{(y_\zeta^1 y_\zeta)_{12}}{(y_\zeta^1 y_\zeta)_{11}}\right] M_{N_1}/M_{N_2}.$$

(14)

Then, the resulting BAU is found as

$$\frac{n_B}{s} \simeq -\frac{1}{15} \frac{\epsilon}{g_s} = (5.8 \sim 6.6) \times 10^{-10},$$

(15)

where $g_s \approx 100$ is the relativistic degrees of freedom, and the last value is the current bound on the BAU [18].

**D. dark matter**

In our model, the possible DM candidates are $\eta_{R/I}$ and the lightest one in $N_{R,2,3}$. For the scalar DM of $\eta_{R/I}$, its nature is similar to the isospin doublet inert boson [19] except Yukawa and additional (gauged) boson interactions. However, since our typical scale of the Yukawa coupling $y_\zeta$ is around $10^{-3}$, its modes cannot be dominant to explain the relic density of DM. Clearly, $\eta_{R/I}$ cannot help to rely on the interactions of the Higgs potential and/or kinetic term. Nevertheless, it is known that there exist a lot of solutions to satisfy the present exclusion limits from the direct DM experimental detections, even when a model is minimal. To escape the limits from the spin independent direct detection searches reported by LUX [20], XENON1T [21], and PandaX-II [22], we have to consider two dominant modes from $Z^{(\prime)}$ and CP-even Higgs portals. The former one with the $Z$ boson mediation can easily be evaded by giving the mass difference between $\eta_R$ and $\eta_I$ to be greater than $O(100)$ keV. Note that the constraint with the $Z'$ mediation is always weaker than the $Z$ one, since its cross section is proportional to $(g'/m_{Z'})^1 < (1/(7\text{TeV}))^2$ given by the LEP experiment. The latter one can also be used to avoid the limits by taking the corresponding quartic Higgs couplings, which can be written in terms of linear combinations $\lambda_{H\eta}^{(\prime)}, \lambda_{H\phi_{8(10)}}$, to be less than 0.01, when the SM Higgs is mediating and the other masses of CP-even Higgses are assumed to be heavier than the mass of the SM Higgs.
In case of the fermion DM candidate such as the lightest state of \( N_{R1,2,3} \), its main modes to the relic density can be found in the kinetic term with the additional gauge boson, and/or Higgs potential. As for both the modes, its solution tends to be at around the pole with half masses of mediating fields. A comprehensive analysis has recently been done in refs. \[23, 24\] \(^3\) To satisfy the direct detection bounds, we have the similar modes from \( Z' \) and CP-even Higgses portals. With the same reason as the \( Z' \) mediation, there is almost no constraint. As a result, we should consider the CP-even Higgs only. The exclusion limits can be evaded by taking the corresponding quartic Higgs couplings of \( y_{N_{i,3}} \) depending on the lightest fermion of \( N_{R1,2,3} \), to be less than 0.01, when the SM Higgs is mediating and the other masses of CP-even Higgses are assumed to be heavier than the mass of the SM Higgs. Once we identify either of \( N_{R_{1,2}} \) as DM, the allowed parameter space might be restricted a little \(^4\). On the other hand, if \( N_{R_3} \) is DM, we can discuss the DM issue independently.

**III. NUMERICAL ANALYSIS**

For the numerical calculations, we use the neutrino oscillation data as well as the following input parameters:

\[
\text{Re}[z] \in (0, \pi), \quad \text{Im}[z] \in (-10, -1),
\]

\[
(m_{\zeta_R}, M_{N_1}) \in (100, 1000)\text{GeV}, \quad M_{N_2} \in (3M_{N_1}, 10000)\text{GeV}
\]

where we have taken \( m_{\zeta_R} = m_{\zeta_l} = m_{\zeta_\pm} \) to avoid the constraints from the oblique parameters \([11]\). Note that we fix \( m_{\zeta_R} \approx 895M_{N_1} \) and \( \text{Min}[M_{N_2}] = 3M_{N_1} \) to guarantee the out-of-equilibrium condition and \( 1 - m_{\zeta_\pm}^2/M_{N_1}^2 = 10^{-4} \) in Eq. (14).

In our numerical analysis, we find that our results do not depend the neutrino hierarchy. In particular, all the ranges for \( \text{Re}[z] \) and \( M_{N_{1,2}} \) are allowed, whereas the allowed region of \( \text{Re}[z] \) is restricted to be \([ -9.5, -4 ]\) due to the constraint of the BAU in Eq. (15). In Figs. 2 and 3, we show the allowed regions in the planes of \( M_{N_1} - \text{BR}(\mu \to e\gamma) \), \( y_{\zeta_{11}} - y_{\zeta_{12}} \) and \( y_{\zeta_{31}} - y_{\zeta_{32}} \), to illustrate the feature of the neutrino hierarchical dependences, respectively. In

\[^3\] Notice that this kind of models always have a physical goldstone boson, it can affect the relic density as discussed in ref. \[^3\].

\[^4\] Since the related Yukawa coupling \( y_\zeta \) cannot be order one in order to lead the successful resonant leptogenesis, our final result does not change drastically.
Fig. 2: Scattering plots to satisfy all the data in the plane of BR(\(\mu \rightarrow e\gamma\)) and \(M_{N_1}\), where the left and right figures correspond to NH and IH, respectively.

Fig. 2: the constraint for NH is a little weaker than the one for IH with the fixed value of \(M_{N_1}\) by order one magnitude. This may indicate that IH needs stronger degenerated masses for active neutrinos than NH, as discussed in Sec. II.A. In Fig. 3, this feature is even clear. Since \(y_{\zeta_{11}}\) and \(y_{\zeta_{12}}\) mainly contribute to the first active neutrino mass, these values of NH are smaller than the ones of IH. While the opposite situation causes for IH, as can be expected. In sum, the typical order of all the components \(y_{\zeta_{\alpha\beta}}\) is found to be \(O(10^{-3})\), leading to all LFV processes to be unobserved in the current experimental measurements.

IV. CONCLUSION

We have proposed a radiatively generated neutrino mass model with a successful leptogenesis to produce the BAU at TeV scale, which contains a gauged \(U(1)_{B-L}\) symmetry with unusual charge assignments of (-4,-4,5) to the right-handed neutrinos. In this model, DM candidates naturally arrive without imposing any additional symmetry to stabilize DM, which is achieved by the resulting symmetry after the SSB of \(U(1)_{B-L}\). We have examined the allowed regions for the model parameters to satisfy all the experimental constraints. We have found that the Yukawa couplings \(y_{\zeta_{\alpha\beta}}\) with a typical order of \(10^{-3}\) lead to small LFV processes, which cannot be measured at the current experiments.
FIG. 3: Scattering plots to satisfy all the data in the planes of $y_{\zeta_{11}} - y_{\zeta_{12}}$ (red dots) and $y_{\zeta_{31}} - y_{\zeta_{32}}$ (blue dots), where the left (right) figures represent NH (IH).

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