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Neutrino mass and dark matter from gauged $U(1)_{B-L}$ breaking

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We propose a new model where the Dirac mass term for neutrinos, the Majorana mass term for right-handed neutrinos, and the other new fermion masses arise via the spontaneous breakdown of the $\mathrm{U}(1)_{\mathrm{B-L}}$ gauge symmetry. The anomaly-free condition gives four sets of assignment of the B−L charge to new particles, and three of these sets have an associated global $\mathrm{U}(1)_{\mathrm{DM}}$ symmetry which stabilizes dark matter candidates. The dark matter candidates contribute to generating the Dirac mass term for neutrinos at the one-loop level. Consequently, tiny neutrino masses are generated at the two-loop level via a type-I-seesaw-like mechanism. We show that this model can satisfy current bounds from neutrino oscillation data, the lepton flavor violation, the relic abundance of the dark matter, and the direct search for the dark matter. This model would be tested at future collider experiments and dark matter experiments.

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I. INTRODUCTION

The existence of neutrino masses has been established very well by the brilliant success of neutrino oscillation measurements [1–9], in spite that neutrinos are massless in the standard model of particle physics (SM) where right-handed neutrinos ν_R are absent. If ν_R are introduced to the SM, there are two possible mass terms for neutrinos [10], the Dirac type $\overline{\nu_L}\nu_R$ and the Majorana type $\overline{(\nu_R)^c}\nu_R$.

Since fermion masses in the SM are generated via the spontaneous breakdown of the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ gauge symmetry, it seems natural that new fermion mass terms which do not exist in the SM arise from spontaneous breakdown of a new gauge symmetry. Let us take a $\mathrm{U}(1)$ as the group of the new gauge symmetry [denoted as $\mathrm{U}(1)'$]. Suppose that the $\mathrm{U}(1)'$ gauge symmetry is spontaneously broken by the vacuum expectation value (VEV) of a scalar field σ^0 which is a singlet under the SM gauge group. Then origins of the Majorana mass term of ν_R and the Dirac mass term of neutrinos can be $\sigma^0\overline{(\nu_R)^c}\nu_R$ [or $(\sigma^0)^*\overline{(\nu_R)^c}\nu_R$] and $\sigma^0\overline{\nu_R}\Phi^T\epsilon L$, respectively, where the field L is the $\mathrm{SU}(2)_L$ doublet of leptons, Φ is the Higgs doublet field in the SM, and ϵ is the complete antisymmetric tensor for the $\mathrm{SU}(2)_L$ indices. The Majorana mass term for ν_L comes from $(\sigma^0)^3\overline{L^c}\epsilon\Phi^*\Phi^T\epsilon L$ (or $\sigma^0|\sigma^0|^2\overline{L^c}\epsilon\Phi^*\Phi^T\epsilon L$).

When we decompose the dimension-5 operator $\sigma^0\overline{\nu_R}\Phi^T\epsilon L$ with renormalizable interactions, an interesting possibility is the radiative realization of the operator. A variety of models where the Dirac mass term for neutrinos is radiatively generated has been studied in Refs. [11–18] (see also Ref. [19]). In a radiative mechanism for neutrino masses, a dark matter candidate can appear by imposing an *ad hoc* unbroken Z_2 symmetry (see e.g., Refs. [15,20–28]). It would be natural that such a symmetry to stabilize the dark matter appears as a residual symmetry of a gauge symmetry which is spontaneously broken at higher energies than the electroweak scale (see e.g., Refs. [29–31]). The breaking of such a gauge symmetry can also be the origin of masses of new chiral fermions which contribute to the loop diagram. If we take a $\mathrm{U}(1)'$ symmetry as the new gauge symmetry and introduced fermions are only singlet fields under the SM gauge group, a simple choice for $\mathrm{U}(1)'$ is the $U(1)_{\mathrm{B-L}}$ because of the cancellation of the $[\mathrm{SU}(3)_C]^2 \times U(1)'$, the $[\mathrm{SU}(2)_L]^2 \times U(1)'$, the $[\mathrm{U}(1)_Y]^2 \times U(1)'$ and the $\mathrm{U}(1)_Y \times [U(1)']^2$ anomalies.[1] New physics models with the TeV-scale $\mathrm{U}(1)_{\mathrm{B-L}}$ gauge symmetry can be found in e.g., Refs. [33,34]. Collider phenomenology on the $\mathrm{U}(1)_{\mathrm{B-L}}$ gauge symmetry is discussed in e.g., Ref. [35].

Along with the scenario stated above, a model in Ref. [28] was constructed such that the breaking of the $\mathrm{U}(1)_{\mathrm{B-L}}$ gauge symmetry gives a residual symmetry for the dark matter (DM) stability and new fermion mass terms which are absent in the SM (e.g., the Majorana neutrino mass of ν_R, the one-loop generated Dirac mass term of neutrinos, and the masses of new fermions among which the lightest one can be a DM candidate). However, in order to cancel the anomalies for the $\mathrm{U}(1)_{\mathrm{B-L}}$ gauge symmetry, it is required to introduce more new fermions which do not contribute to the mechanism of generating neutrino masses.

In this paper, we propose a new model which is an improved version of the model in Ref. [28] from the viewpoint of the anomaly cancellation. The B−L charges of new particles are assigned such that the condition of anomaly cancellation is satisfied. Consequently, the B−L

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[1]See Ref. [32] for an anomaly-free $\mathrm{U}(1)'$ gauge symmetry when a new fermion field is introduced as an $\mathrm{SU}(2)_L$ triplet with a hypercharge $Y = 0$.

charges for some new particles turn out to be irrational numbers. Because of this charge assignment, there exists an unbroken global $U(1)$ symmetry even after the breakdown of the $U(1)_{B-L}$ symmetry. The global $U(1)$ symmetry stabilizes the dark matter, so that we hereafter call it $U(1)_{DM}$. The lightest particle with the irrational quantum number can be a dark matter candidate. In our model, the dark matter candidate is a new scalar boson with the irrational quantum number. Furthermore, the Dirac mass term of neutrinos is radiatively generated at the one-loop level due to the quantum effect of the new particles with irrational quantum numbers. Tiny neutrino masses are explained by the two-loop diagrams with a type-I-see-saw-like mechanism. We find that the model can satisfy current data from the neutrino oscillation, the lepton flavor violation (LFV), the relic abundance and the direct search for the dark matter, and the LHC experiment.

This paper is organized as follows. In Sec. II, the model is defined and the basic property is discussed. In Sec. III, the neutrino masses are induced due to the spontaneous breaking of $U(1)_{B-L}$. We find a benchmark scenario in which current experimental constraints are taken into account such as the neutrino oscillation data, the LFV, the relic abundance of the dark matter, the direct search for the dark matter, and the LHC results. Conclusions are given in Sec. IV. Some details of our calculations are shown in Appendixes.

II. THE MODEL

New particles listed in Table I are added to the SM. Assignment of $U(1)_{B-L}$ charges is different from that in the previous model in Ref. [28]. Conditions for cancellation of the $[U(1)_{B-L}] \times [\text{gravity}]^2$ and $[U(1)_{B-L}]^3$ anomalies are

$$3 - \frac{1}{3}N_{\nu_R} - \frac{2}{3}N_\psi = 0, \tag{1}$$

$$3 - \frac{1}{27}N_{\nu_R} + \left(-2x^2 - \frac{4}{3}x - \frac{8}{27}\right)N_\psi = 0, \tag{2}$$

TABLE I. Particle contents in this model. Indices i (for ψ_R and ψ_L) and a (for ν_R) run from 1 to N_ψ and from 1 to N_{ν_R}, respectively.

	s^0	η	ψ_{Ri}	ψ_{Li}	ν_{Ra}	σ^0
Spin	0	0	1/2	1/2	1/2	0
$SU(2)_L$	$\underline{\mathbf{1}}$	$\underline{\mathbf{2}}$	$\underline{\mathbf{1}}$	$\underline{\mathbf{1}}$	$\underline{\mathbf{1}}$	$\underline{\mathbf{1}}$
$U(1)_Y$	0	1/2	0	0	0	0
$U(1)_{B-L}$	$x+1$	$x+1$	x	$x+2/3$	$-1/3$	2/3

where N_ψ is the number of ψ_{Ri} (the same as the number of ψ_{Li}), and N_{ν_R} is the number of ν_{Ra}. There are four solutions as presented in Table II. Except for case III, the $U(1)_{B-L}$ charges of some new particles are irrational numbers while the $U(1)_{B-L}$ symmetry is spontaneously broken by the VEV of σ^0 whose $U(1)_{B-L}$ charge is a rational number. Therefore, the irrational charges are conserved, and the lightest particle with an irrational $U(1)_{B-L}$ charge becomes stable so that the particle can be regarded as a dark matter candidate. Notice that there is no dark matter candidate in case III. As we see later, two of three light neutrinos are massless in case I, which does not fit the neutrino oscillation data. In this paper, we take case IV as an example.[2]

The Yukawa interactions are given by

$$\begin{aligned}\mathcal{L}_{\text{Yukawa}} = \mathcal{L}_{\text{SM-Yukawa}} &- (y_R)_a \overline{(\nu_R)_a}(\nu_R)^c_a(\sigma^0)^* \\ &- (y_\Psi)_i \overline{(\psi_R)_i}(\psi_L)_i(\sigma^0)^* - h_{ia}\overline{(\psi_L)_i}(\nu_R)_a s^0 \\ &- f_{\ell i}\overline{L_\ell}(\psi_R)_i\tilde{\eta} + \text{H.c.},\end{aligned} \tag{3}$$

where $\mathcal{L}_{\text{SM-Yukawa}}$ denotes the Yukawa interactions in the SM, L_ℓ ($\ell = e, \mu, \tau$) are the $SU(2)_L$ doublet fields of the SM leptons, and $\tilde{\eta} \equiv ((\eta^0)^*, -\eta^-)^T$. Indices i and a run from 1 to N_ψ and from 1 to N_{ν_R}, respectively. Notice that a Yukawa interaction $\overline{(\nu_R)^c_a}\psi_{Ri}(s^0)^*$ which exists in the previous model is absent in this model because assignments of B−L charge to new particles are different from those in the previous paper [28].

The scalar potential in our model is the same as that in the previous model[3] [28]:

$$\begin{aligned}V(\Phi, s, \eta, \sigma) = &-\mu^2_\phi\Phi^\dagger\Phi + \mu^2_s|s^0|^2 + \mu^2_\eta\eta^\dagger\eta - \mu^2_\sigma|\sigma^0|^2 + \lambda_\phi(\Phi^\dagger\Phi)^2 + \lambda_s|s^0|^4 + \lambda_\eta(\eta^\dagger\eta)^2 + \lambda_\sigma|\sigma^0|^4 \\ &+ \lambda_{s\eta}|s^0|^2\eta^\dagger\eta + \lambda_{s\phi}|s^0|^2\Phi^\dagger\Phi + \lambda_{\phi\phi}(\eta^\dagger\eta)(\Phi^\dagger\Phi) + \lambda_{\eta\phi}(\eta^\dagger\Phi)(\Phi^\dagger\eta) \\ &+ \lambda_{s\sigma}|s^0|^2|\sigma^0|^2 + \lambda_{\sigma\eta}|\sigma^0|^2\eta^\dagger\eta + \lambda_{\sigma\phi}|\sigma^0|^2\Phi^\dagger\Phi + (\mu_3 s^0\eta^\dagger\Phi + \text{H.c.}),\end{aligned} \tag{4}$$

where μ^2_ϕ, μ^2_s, μ^2_η, and μ^2_σ are defined as positive values. Without loss of generality, we can take a real positive μ_3 by utilizing a rephasing of s^0.

Two scalar fields ϕ^0 and σ^0 obtain VEVs v_ϕ [$= \sqrt{2}\langle\phi^0\rangle = 246$ GeV] and v_σ [$= \sqrt{2}\langle\sigma^0\rangle$]. Then $SU(2)_L \times U(1)_Y$ and $U(1)_{B-L}$ gauge symmetries are spontaneously broken by v_ϕ and v_σ, respectively. These VEVs are given by

[2]If the B − L charge of σ^0 is 2 as in the model in Ref. [28], the B − L charges for $\{s^0, \eta, \psi_R, \psi_L, \nu_R\}$ will be assigned as $\{x+1, x+1, x, x+2, -1\}$. There is only an anomaly-free solution $x = -1$. We do not take this possibility because there is no residual symmetry to stabilize the dark matter.

[3]For case III in Table II, there are additional terms e.g., $(s^0)^*(\sigma^0)^2$. See also Ref. [36].

TABLE II. Sets of N_ψ, N_{ν_R} and x, for which the $\mathrm{U}(1)_{\mathrm{B-L}}$ gauge symmetry is free from anomaly. Here, N_ψ is the number of ψ_{Ri} (the same as the number of ψ_{Li}), N_{ν_R} is the number of ν_{Ra}, and x is the B − L charge of ψ_{Ri}.

	Case I	Case II	Case III	Case IV
N_ψ	1	2	3	4
N_{ν_R}	7	5	3	1
x	$\frac{2\sqrt{3}-1}{3}$	$\frac{\sqrt{6}-1}{3}$	$\frac{1}{3}$	$\frac{\sqrt{3}-1}{3}$

$$\begin{pmatrix} v_\phi^2 \\ v_\sigma^2 \end{pmatrix} = \frac{1}{\lambda_\sigma\lambda_\phi - \lambda_{\sigma\phi}^2/4}\begin{pmatrix} \lambda_\sigma & -\lambda_{\sigma\phi}/2 \\ -\lambda_{\sigma\phi}/2 & \lambda_\phi \end{pmatrix}\begin{pmatrix} \mu_\phi^2 \\ \mu_\sigma^2 \end{pmatrix}. \quad (5)$$

The VEV v_σ provides a mass of the $\mathrm{U}(1)_{\mathrm{B-L}}$ gauge boson Z' as $m_{Z'} = (2/3)g_{\mathrm{B-L}}v_\sigma$, where $g_{\mathrm{B-L}}$ is the $\mathrm{U}(1)_{\mathrm{B-L}}$ gauge coupling constant. After the gauge symmetry breaking with v_ϕ and v_σ, we can confirm in Eqs. (3) and (4) that there is a residual global $\mathrm{U}(1)_{\mathrm{DM}}$ symmetry, for which irrational $\mathrm{U}(1)_{\mathrm{B-L}}$-charged particles ($\eta$, s^0, ψ_{Li}, and ψ_{Ri}) have the same $\mathrm{U}(1)_{\mathrm{DM}}$ charge while the other particles are neutral.

We have two CP-even scalar particles h^0 and H^0 as

$$\begin{pmatrix} h^0 \\ H^0 \end{pmatrix} = \begin{pmatrix} \cos\theta_0 & -\sin\theta_0 \\ \sin\theta_0 & \cos\theta_0 \end{pmatrix}\begin{pmatrix} \phi_r^0 \\ \sigma_r^0 \end{pmatrix},$$
$$\sin 2\theta_0 = \frac{2\lambda_{\sigma\phi}v_\phi v_\sigma}{m_{H^0}^2 - m_{h^0}^2}, \quad (6)$$

where $\phi^0 = (v_\phi + \phi_r^0 + iz_\phi)/\sqrt{2}$ and $\sigma^0 = (v_\sigma + \sigma_r^0 + iz_\sigma)/\sqrt{2}$. Nambu-Goldstone bosons z_ϕ and z_σ are absorbed by Z and Z' bosons, respectively. Masses of h^0 and H^0 are given by

$$m_{h^0}^2 = \lambda_\phi v_\phi^2 + \lambda_\sigma v_\sigma^2 - \sqrt{(\lambda_\phi v_\phi^2 - \lambda_\sigma v_\sigma^2)^2 + \lambda_{\sigma\phi}^2 v_\phi^2 v_\sigma^2},$$
$$m_{H^0}^2 = \lambda_\phi v_\phi^2 + \lambda_\sigma v_\sigma^2 + \sqrt{(\lambda_\phi v_\phi^2 - \lambda_\sigma v_\sigma^2)^2 + \lambda_{\sigma\phi}^2 v_\phi^2 v_\sigma^2}. \quad (7)$$

On the other hand, η^0 and s^0 do not mix with ϕ^0 and σ^0 even though the $\mathrm{U}(1)_{\mathrm{B-L}}$ symmetry is broken by v_σ. Two neutral complex scalars $\mathcal{H}_1^0$ and $\mathcal{H}_2^0$ are obtained by

$$\begin{pmatrix} \mathcal{H}_1^0 \\ \mathcal{H}_2^0 \end{pmatrix} = \begin{pmatrix} \cos\theta_0' & -\sin\theta_0' \\ \sin\theta_0' & \cos\theta_0' \end{pmatrix}\begin{pmatrix} \eta^0 \\ s^0 \end{pmatrix},$$
$$\sin 2\theta_0' = \frac{\sqrt{2}\mu_3 v_\phi}{m_{\mathcal{H}_2^0}^2 - m_{\mathcal{H}_1^0}^2}. \quad (8)$$

Their masses and the mass of the charged scalar $\eta^\pm$ are given by

$$m_{\mathcal{H}_1^0}^2 = \frac{1}{2}\left(m_\eta^2 + m_s^2 - \sqrt{(m_\eta^2 - m_s^2)^2 + 2\mu_3^2 v_\phi^2}\right), \quad (9)$$

$$m_{\mathcal{H}_2^0}^2 = \frac{1}{2}\left(m_\eta^2 + m_s^2 + \sqrt{(m_\eta^2 - m_s^2)^2 + 2\mu_3^2 v_\phi^2}\right), \quad (10)$$

$$m_{\eta^\pm}^2 = m_\eta^2 - \lambda_{\eta\phi}\frac{v_\phi^2}{2}, \quad (11)$$

where $m_s^2 \equiv \mu_s^2 + \lambda_{s\phi}v_\phi^2/2 + \lambda_{s\sigma}v_\sigma^2/2$ and $m_\eta^2 \equiv \mu_\eta^2 + (\lambda_{\phi\phi} + \lambda_{\eta\phi})v_\phi^2/2 + \lambda_{\sigma\eta}v_\sigma^2/2$.

III. NEUTRINO MASS AND DARK MATTER

A. Neutrino mass

Tiny neutrino masses are generated by two-loop diagrams in Fig. 1 [28]. The mass matrix m_ν is expressed in the flavor basis as

$$(m_\nu)_{\ell\ell'} = \frac{1}{(16\pi^2)^2}\sum_{i,j,a} f_{\ell i}h_{ia}(m_R)_a(h^T)_{aj}(f^T)_{j\ell'}\{(I_1)_{ija} + (I_2)_{ija}\}. \quad (12)$$

Loop functions $(I_1)_{ija}$ and $(I_2)_{ija}$ correspond to contributions of diagram (a) and (b) in Fig. 1, respectively. Explicit forms of these loop functions are shown in Appendix A.

Let us define the following matrix:

$$A_{ij} \equiv \sum_a h_{ia}(m_R)_a(h^T)_{aj}\{(I_1)_{ija} + (I_2)_{ija}\}. \quad (13)$$

If $N_\psi = 1$, the matrix A_{ij} becomes just a number and then $(m_\nu)_{\ell\ell'}$ becomes a rank-1 matrix which is not consistent with neutrino oscillation data. Therefore, case I in Table II is not acceptable. On the other hand, $N_{\nu_R} = 1$ does not mean that $(m_\nu)_{\ell\ell'}$ is a rank-1 matrix because of the existence of $(I_2)_{ija}$. We will see later that our benchmark point for case IV in Table II does not include massless neutrinos even though $N_{\nu_R} = 1$.

The neutrino mass matrix $(m_\nu)_{\ell\ell'}$ is diagonalized by a unitary matrix U_{MNS}, the so-called Maki-Nakagawa-Sakata (MNS) matrix [37], as $U_{\mathrm{MNS}}^\dagger m_\nu U_{\mathrm{MNS}}^* = \mathrm{diag}(m_1 e^{i\alpha_1}, m_2 e^{i\alpha_2}, m_3 e^{i\alpha_3})$. We take m_i ($i = 1$–3) to be real and positive values. Two differences of three phases α_i are physical Majorana phases [38]. The MNS matrix can be parametrized as

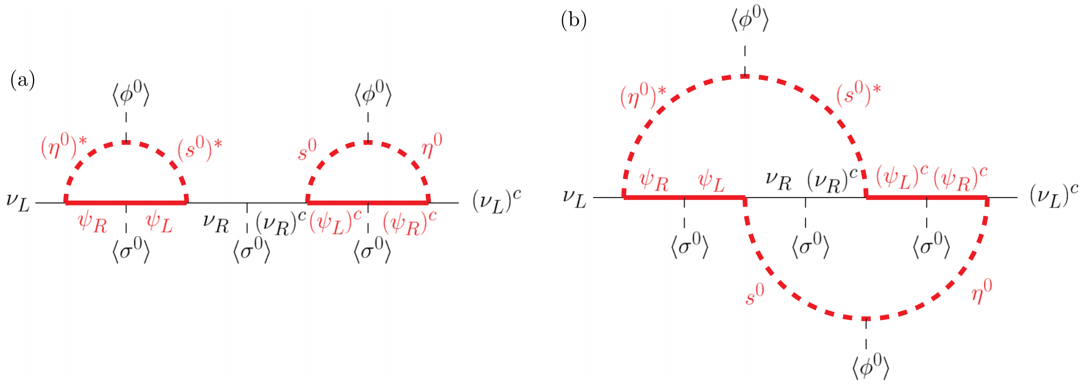

FIG. 1 (color online). Two-loop diagrams for tiny neutrino masses in this model. Bold (red) lines are propagators of particles of irrational $\mathrm{U}(1)_{\mathrm{B-L}}$ charges.

$$U_{\mathrm{MNS}} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_{23} & s_{23} \\ 0 & -s_{23} & c_{23} \end{pmatrix} \begin{pmatrix} c_{13} & 0 & s_{13}e^{-i\delta} \\ 0 & 1 & 0 \\ -s_{13}e^{i\delta} & 0 & c_{13} \end{pmatrix} \begin{pmatrix} c_{12} & s_{12} & 0 \\ -s_{12} & c_{12} & 0 \\ 0 & 0 & 1 \end{pmatrix}, \tag{14}$$

where $c_{ij} \equiv \cos\theta_{ij}$ and $s_{ij} \equiv \sin\theta_{ij}$. In our analysis, the following values [2,5,8] obtained by neutrino oscillation measurements are used in order to search for a benchmark point of model parameters:

$$m_1 = 10^{-4}\ \mathrm{eV}, \tag{15}$$

$$\Delta m_{21}^2 = 7.46\times 10^{-5}\ \mathrm{eV}^2, \quad \Delta m_{32}^2 = +2.51\times 10^{-3}\ \mathrm{eV}^2, \tag{16}$$

$$\sin^2 2\theta_{23} = 1, \qquad \sin^2 2\theta_{13} = 0.09, \qquad \tan^2\theta_{12} = 0.427, \tag{17}$$

$$\delta = 0, \qquad \{\alpha_1, \alpha_2, \alpha_3\} = \{0, 0, 0\}, \tag{18}$$

where $\Delta m_{ij}^2 \equiv m_i^2 - m_j^2$. By using an ansatz presented in Appendix B for the structure of Yukawa matrix $f_{\ell i}$, we found a benchmark point as

$$f = \begin{pmatrix} 1.79 & -2.49 & -1.97 & 2.56 \\ -1.82 & 1.10 & 1.30 & -0.818 \\ 1.40 & -0.598 & -0.905 & 0.222 \end{pmatrix} \times 10^{-2}, \tag{19}$$

$$h = (\,0.7 \quad 0.8 \quad 0.9 \quad 1\,)^T, \tag{20}$$

$$(m_R)_1 = 250\ \mathrm{GeV}, \tag{21}$$

$$\{m_{\psi_1}, m_{\psi_2}, m_{\psi_3}, m_{\psi_4}\} = \{650, 750, 850, 950\ \mathrm{GeV}\}, \tag{22}$$

$$\{m_{h^0}, m_{H^0}, \cos\theta_0\} = \{125\ \mathrm{GeV}, 1000\ \mathrm{GeV}, 1\}, \tag{23}$$

$$\{m_{\mathcal{H}_1^0}, m_{\mathcal{H}_2^0}, \cos\theta_0'\} = \{60\ \mathrm{GeV}, 450\ \mathrm{GeV}, 0.05\}, \tag{24}$$

$$m_{\eta^\pm} = 420\ \mathrm{GeV}, \tag{25}$$

$$\{g_{\mathrm{B-L}}, m_{Z'}\} = \{0.1, 4000\ \mathrm{GeV}\}. \tag{26}$$

The values of $\{g_{\mathrm{B-L}}, m_{Z'}\}$ mean $v_\sigma = 60$ TeV. The values of $\{m_{h^0}, m_{H^0}, \cos\theta_0\}$ correspond to $\lambda_\phi \simeq 0.13$, $\lambda_\sigma \simeq 2.8\times 10^{-4}$ and $\lambda_{\sigma\phi} = 0$. The values of $\{m_{\mathcal{H}_1^0}, m_{\mathcal{H}_2^0}, \cos\theta_0'\}$ and $m_{\eta^\pm}$ can be produced by $m_s \simeq 60$ GeV, $m_\eta \simeq 450$ GeV, $\mu_3 \simeq 57$ GeV and $\lambda_{\eta\phi} \simeq 0.86$.

B. Lepton flavor violation

The charged scalar $\eta^\pm$ contributes to the LFV decays of charged leptons. The formula for the branching ratio (BR) of $\mu \to e\gamma$ can be calculated [39] as

$$\mathrm{BR}(\mu \to e\gamma) = \frac{3\alpha_{\mathrm{EM}}}{64\pi G_F^2}\left|\frac{1}{m_{\eta^\pm}^2} f_{\mu i} F\left(\frac{m_{\psi_i}^2}{m_{\eta^\pm}^2}\right)(f^\dagger)_{ie}\right|^2, \tag{27}$$

where

$$F(x) \equiv \frac{1 - 6x + 3x^2 + 2x^3 - 6x^2\ln(x)}{6(1-x)^4}. \tag{28}$$

At the benchmark point, we have $\mathrm{BR}(\mu \to e\gamma) = 6.1\times 10^{-14}$ which satisfies the current constraint $\mathrm{BR}(\mu \to e\gamma) < 5.7\times 10^{-13}$ (90% C.L.) [40].

C. Dark matter

In principle, ψ_1 or $\mathcal{H}_1^0$ can be a dark matter candidate. However, due to the following reason, the scalar $\mathcal{H}_1^0$ turns out to be the dark matter candidate. If the dark matter is the

TABLE III. Branching ratios of Z' decays.

$q\overline{q}$	$\ell\overline{\ell}$	$\nu_L\overline{\nu_L}$	$\nu_R\overline{\nu_R}$	$\psi_1\overline{\psi_1}$	$\psi_2\overline{\psi_2}$	$\psi_3\overline{\psi_3}$	$\psi_4\overline{\psi_4}$	$\mathcal{H}_1^0\mathcal{H}_1^{0*}$	$\mathcal{H}_2^0\mathcal{H}_2^{0*}$	$\eta^+\eta^-$
0.21	0.32	0.16	0.0059	0.046	0.045	0.044	0.043	0.041	0.038	0.039

fermion ψ_1, it annihilates into a pair of SM particles via the s-channel process mediated by h^0 and H^0. The cross section of the process is proportional to $\sin^2 2\theta_0$. In order to obtain a sufficient annihilation cross section of ψ_1, a large mixing $\cos\theta_0 \simeq 1/\sqrt{2}$ is preferred [34]. Even for a maximal mixing $\cos\theta_0 = 1/\sqrt{2}$, the observed abundance of the dark matter [41] requires $v_\sigma \lesssim 10$ TeV. The current constraint from direct searches of the dark matter [42] requires larger v_σ in order to suppress the Z' contribution.[4]

Because of the tiny mixing $\cos\theta_0' = 0.05$, the scalar dark matter $\mathcal{H}_1^0$ at the benchmark point is dominantly made from s^0 which is a gauge-singlet field under the SM gauge group. The annihilation of $\mathcal{H}_1^0$ into a pair of the SM particles is dominantly caused by the s-channel scalar mediation via h^0 [43] because H^0 is assumed to be heavy. The coupling constant $\lambda_{\mathcal{H}_1^0\mathcal{H}_1^0 h^0}$ for the $\lambda_{\mathcal{H}_1^0\mathcal{H}_1^0 h^0} v_\phi \mathcal{H}_1^0\mathcal{H}_1^{0*} h^0$ interaction controls the annihilation cross section, the invisible decay $h^0 \to \mathcal{H}_1^0\mathcal{H}_1^{0*}$ in the case of kinematically accessible, and the h^0 contribution to the spin-independent scattering cross section $\sigma_{\rm SI}$ on a nucleon. In Ref. [44], for example, we see that $\mathcal{H}_1^0$ with $m_{\mathcal{H}_1^0} = 60$ GeV and $\lambda_{\mathcal{H}_1^0\mathcal{H}_1^0 h^0} \sim 10^{-3}$ can satisfy constraints from the relic abundance of the dark matter and the invisible decay of h^0. We see also that the h^0 contribution to $\sigma_{\rm SI}$ is small enough to satisfy the current constraint $\sigma_{\rm SI} < 9.2 \times 10^{-46}$ cm^2 for $m_{\rm DM} = 60$ GeV [42]. Although the scattering of $\mathcal{H}_1^0$ on a nucleon is mediated also by the Z' boson in this model, the contribution can be suppressed by taking a large v_σ. The benchmark point corresponds to $v_\sigma = 60$ TeV and gives about 6.6×10^{-47} cm^2 for the scattering cross section via Z', which is smaller than the current constraint [42] by an order of magnitude. Thus, the constraint from the direct search of the dark matter is also satisfied at the benchmark point.

D. Collider phenomenology

The light CP-even neutral scalar h^0 is made from an SU$(2)_L$-doublet field Φ because we take $\cos\theta_0 = 0$. The mass $m_{h^0} = 125$ GeV at the benchmark point is consistent with $m_{h^0_{\rm SM}} = 125.5 \pm 0.2({\rm stat})^{+0.5}_{-0.6}({\rm sys})$ GeV in the ATLAS experiment [45] and $m_{h^0_{\rm SM}} = 125.7 \pm 0.3({\rm stat}) \pm 0.3({\rm sys})$ GeV in the CMS experiment [46]. The branching ratio of the invisible decay $h^0 \to \mathcal{H}_1^0\mathcal{H}_1^{0*}$ is about 7×10^{-4} for $\lambda_{\mathcal{H}_1^0\mathcal{H}_1^0 h^0} = 0.001$, where the recommended value 4.07 MeV [47] for the total width of $h^0_{\rm SM}$ is used.

For the Z' boson, the LEP-II bound $m_{Z'}/g_{\rm B-L} \gtrsim 7$ TeV [48] is satisfied at the benchmark point because of

[4]This is because $m_{Z'}/g_{\rm B-L}$ is not $2v_\sigma$ as usual but $2v_\sigma/3$ in this model.

TABLE IV. Branching ratios of ν_R decays.

$W^+\ell^- + W^-\ell^+$	$Z\nu_L + Z\overline{\nu_L}$	$h^0\nu_L + h^0\overline{\nu_L}$	$H^0\nu_L + H^0\overline{\nu_L}$
0.56	0.28	0.16	0

$m_{Z'}/g_{\rm B-L} = 40$ TeV which we take for a sufficient suppression of $\sigma_{\rm SI}$ for the direct search of the dark matter. The production cross section of Z' with $g_{\rm B-L} = 0.1$ and $m_{Z'} = 4000$ GeV is about 0.3 fb at the LHC with $\sqrt{s} = 14$ TeV [35].[5] Decay branching ratios of Z' are shown at the benchmark point in Table III.

Decays of ψ_i are dominated by $\psi_i \to \nu_R\mathcal{H}_1^0$ with the Yukawa coupling constants h_{i1} because $y_{\ell i}$ for $\psi_i \to \ell^\pm\eta^\mp$ are small in order to satisfy the $\mu \to e\gamma$ constraint. The $\mathcal{H}_2^0$ ($\simeq\eta^0$) decays into $h^0\mathcal{H}_1^0$ via the trilinear coupling constant μ_3. The main decay mode of the charged scalar is $\eta^\pm \to W^\pm\mathcal{H}_1^0$ through the mixing θ_0' between η^0 and s^0.

In this model, ν_R is not the dark matter and can decay into the SM particles. Decay branching ratios for ν_R are shown in Table IV. The decay into H^0 is forbidden because it is heavier than ν_R at the benchmark point. Since the B−L charge of ν_R is rather small, ν_R is not produced directly from Z'. However, ν_R can be produced through the decays of ψ_i. As a result, about 18% of Z' produces ν_R. For $\nu_R \to W\ell$ (56%) followed by the hadronic decay of W (68%), the ν_R would be reconstructed. In this model, an invariant mass of a pair of the reconstructed ν_R is not at $m_{Z'}$ in contrast with a naive model where only three ν_R with ${\rm B-L} = -1$ are introduced to the SM.[6] This feature of ν_R also enables us to distinguish this model from the previous model in Ref. [28] where ν_R with ${\rm B-L} = 1$ can be directly produced by the Z' decay.

IV. CONCLUSIONS

We have proposed the model which is an improved version of the model in Ref. [28] by considering anomaly cancellation of the U$(1)_{\rm B-L}$ gauge symmetry. We have shown that there are four anomaly-free cases of B−L charge assignment, and three of them have an unbroken global U$(1)_{\rm DM}$ symmetry (one of the three is not acceptable because two neutrinos become massless). The U$(1)_{\rm DM}$

[5]The production cross section becomes about 6 fb if we take $g_{\rm B-L} = 0.05$ and $m_{Z'} = 2000$ GeV. Notice that the current bound $m_{Z'} \gtrsim 3$ TeV at the LHC [49] is for the case where the gauge coupling for Z' is the same as the one for Z, namely $g_{\rm B-L} \simeq 0.7$.

[6]In the naive model with $m_{Ra} = 250$ GeV (degenerate) and $m_{Z'} = 4$ TeV, the decay branching ratios of Z' into $\{q\overline{q}, \ell\overline{\ell}, \nu_L\overline{\nu_L}, \nu_R\overline{\nu_R}\}$ are $\{0.25, 0.38, 0.19, 0.19\}$.

guarantees that the lightest $\mathrm{U(1)_{DM}}$-charged particle is stable such that it can be regarded as a dark matter candidate. The spontaneous breaking of the $\mathrm{U(1)_{B-L}}$ symmetry generates new fermion mass terms which do not exist in the SM; namely, the Dirac mass term of neutrinos, the Majorana mass term of ν_R, and masses of new fermions ψ. Especially, the Dirac mass term of neutrinos is generated at the one-loop level where the dark matter candidate involved in the loop. Tiny neutrino masses are obtained at the two-loop level. The case of the fermion dark matter is excluded, and the lightest $\mathrm{U(1)_{DM}}$-charged scalar $\mathcal{H}_1^0$ should be the dark matter in this model. We have found a benchmark point of model parameters which satisfies current constraints from neutrino oscillation data, lepton flavor violation searches, the relic abundance of the dark matter, direct searches for the dark matter, and the LHC experiments.

By virtue of the radiative mechanism for the Dirac mass term of neutrinos, very heavy ν_R are not required for tiny neutrino masses. Therefore, ν_R would be produced at the LHC. In contrast to a naive model where three ν_R have $\mathrm{B-L}=-1$ and the model in Ref. [28] where ν_R have $\mathrm{B-L}=1$, the ν_R with $\mathrm{B-L}=-1/3$ in this model cannot be directly produced by the Z' decay but can be produced by the cascade decay $Z'\to\psi_i\overline{\psi}_i\to\nu_R\overline{\nu_R}\mathcal{H}_1^0\mathcal{H}_1^{0*}$. The invariant mass distribution of $\nu_R\overline{\nu_R}$ does not take a peak at $m_{Z'}$, which could be a characteristic signal of this kind of models with the unusual B−L charge of ν_R.

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APPENDIX A: LOOP INTEGRATION

A loop function $(I_1)_{ija}$ in Eq. (12) can be expressed as

$$\begin{aligned}(I_1)_{ija} &\equiv -\frac{(8\pi^2\sin 2\theta_0')^2 m_{\psi_i}m_{\psi_j}}{(m_R)_a^2}\left[\int\frac{d^4p}{(2\pi)^4}\frac{1}{p^2-m_{\psi_i}^2}\left\{\frac{1}{p^2-m_{\mathcal{H}_1^0}^2}-\frac{1}{p^2-m_{\mathcal{H}_2^0}^2}\right\}\right]\\ &\quad\times\left[\int\frac{d^4q}{(2\pi)^4}\frac{1}{q^2-m_{\psi_j}^2}\left\{\frac{1}{q^2-m_{\mathcal{H}_1^0}^2}-\frac{1}{q^2-m_{\mathcal{H}_2^0}^2}\right\}\right]\\ &=\frac{m_{\psi_i}m_{\psi_j}(m_{\mathcal{H}_1^0}^2-m_{\mathcal{H}_2^0}^2)^2\sin^2 2\theta_0'}{4(m_R)_a^2}\{C_0(0,0,m_{\psi_i},m_{\mathcal{H}_1^0}^2,m_{\mathcal{H}_2^0}^2)\\ &\quad\times C_0(0,0,m_{\psi_j},m_{\mathcal{H}_1^0}^2,m_{\mathcal{H}_2^0}^2)\},\end{aligned}\tag{A1}$$

where the C_0 function [50] is given by

$$\begin{aligned}&C_0(0,0,m_0^2,m_1^2,m_2^2)\\ &\equiv\frac{1}{(m_0^2-m_1^2)(m_1^2-m_2^2)(m_2^2-m_0^2)}\left\{m_0^2m_1^2\ln\frac{m_0^2}{m_1^2}+m_1^2m_2^2\ln\frac{m_1^2}{m_2^2}+m_2^2m_0^2\ln\frac{m_2^2}{m_0^2}\right\}.\end{aligned}\tag{A2}$$

On the other hand, another loop function $(I_2)_{ija}$ in Eq. (12) is given by

$$\begin{aligned}(I_2)_{ija} &\equiv (8\pi^2\sin 2\theta_0')^2 m_{\psi_i}m_{\psi_j}\\ &\quad\times\iint\frac{d^4p}{(2\pi)^4}\frac{d^4q}{(2\pi)^4}\left\{\frac{1}{p^2-m_{\mathcal{H}_1^0}^2}-\frac{1}{p^2-m_{\mathcal{H}_2^0}^2}\right\}\frac{1}{p^2-m_{\psi_i}^2}\\ &\quad\times\frac{1}{(p+q)^2-(m_R)_a^2}\left\{\frac{1}{q^2-m_{\mathcal{H}_1^0}^2}-\frac{1}{q^2-m_{\mathcal{H}_2^0}^2}\right\}\frac{1}{q^2-m_{\psi_j}^2}\\ &=(8\pi^2\sin 2\theta_0')^2 m_{\psi_i}m_{\psi_j}\\ &\quad\times[I(m_{\mathcal{H}_1^0},m_{\psi_i}|m_{\mathcal{H}_1^0},m_{\psi_j}|(m_R)_a)-I(m_{\mathcal{H}_1^0},m_{\psi_i}|m_{\mathcal{H}_2^0},m_{\psi_j}|(m_R)_a)\\ &\quad-I(m_{\mathcal{H}_2^0},m_{\psi_i}|m_{\mathcal{H}_1^0},m_{\psi_j}|(m_R)_a)+I(m_{\mathcal{H}_2^0},m_{\psi_i}|m_{\mathcal{H}_2^0},m_{\psi_j}|(m_R)_a)],\end{aligned}\tag{A3}$$

where

$$I(m_{11}, m_{12}, \ldots, m_{1n_1}|m_{21}, m_{22}, \ldots, m_{2n_2}|m_{31}, m_{32}, \ldots, m_{3n_3}) \equiv \int \frac{d^4p_E}{(2\pi)^4}\int \frac{d^4q_E}{(2\pi)^4}\prod_{i=1}^{n_1}\prod_{j=1}^{n_2}\prod_{k=1}^{n_3}\frac{1}{p_E^2+m_{1i}^2}\frac{1}{q_E^2+m_{2j}^2}\frac{1}{(p_E+q_E)^2+m_{3k}^2}. \quad \text{(A4)}$$

We can use the following results [51]:

$$I(m_{11}, m_{12}|m_{21}, m_{22}|m_3) = \frac{I(m_{12}|m_{22}|m_3) - I(m_{11}|m_{22}|m_3) - I(m_{12}|m_{21}|m_3) + I(m_{11}|m_{21}|m_3)}{(16\pi^2)^2(m_{11}^2-m_{12}^2)(m_{21}^2-m_{22}^2)}, \quad \text{(A5)}$$

$$I(m_1|m_2|m_3) = -m_1^2 f\left(\frac{m_2^2}{m_1^2}, \frac{m_3^2}{m_1^2}\right) - m_2^2 f\left(\frac{m_1^2}{m_2^2}, \frac{m_3^2}{m_2^2}\right) - m_3^2 f\left(\frac{m_1^2}{m_3^2}, \frac{m_2^2}{m_3^2}\right), \quad \text{(A6)}$$

where

$$f(x,y) \equiv -\frac{1}{2}(\ln x)(\ln y) - \frac{1}{2}\left(\frac{x+y-1}{D}\right) \times \left\{\mathrm{Li}_2\left(\frac{-x_-}{y_+}\right) + \mathrm{Li}_2\left(\frac{-y_-}{x_+}\right) - \mathrm{Li}_2\left(\frac{-x_+}{y_-}\right) - \mathrm{Li}_2\left(\frac{-y_+}{x_-}\right) + \mathrm{Li}_2\left(\frac{y-x}{x_-}\right) + \mathrm{Li}_2\left(\frac{x-y}{y_-}\right) - \mathrm{Li}_2\left(\frac{y-x}{x_+}\right) - \mathrm{Li}_2\left(\frac{x-y}{y_+}\right)\right\}, \quad \text{(A7)}$$

and

$$D \equiv \sqrt{1-2(x+y)+(x-y)^2}, \quad \text{(A8)}$$

$$x_\pm \equiv \frac{1}{2}(1-x+y\pm D), \qquad y_\pm \equiv \frac{1}{2}(1+x-y\pm D), \quad \text{(A9)}$$

and the dilog function $\mathrm{Li}_2(x)$ is defined as

$$\mathrm{Li}_2(x) \equiv -\int_0^x dt\frac{\ln(1-t)}{t}. \quad \text{(A10)}$$

APPENDIX B: ANSATZ FOR BENCHMARK POINT

The symmetric matrix A_{ij} in Eq. (13) can be diagonalized by an orthogonal matrix X as

$$XAX^T = \mathrm{diag}(a_1, a_2, a_3, a_4). \quad \text{(B1)}$$

It is clear that a Yukawa matrix $f_{\ell i}$ of the following structure satisfies constraints from neutrino oscillation data:

$$f = 16\pi^2 U_{\mathrm{MNS}}\begin{pmatrix}\sqrt{\frac{m_1}{|a_1|}} & 0 & 0 & 0\\ 0 & \sqrt{\frac{m_2}{|a_2|}} & 0 & 0\\ 0 & 0 & \sqrt{\frac{m_3}{|a_3|}} & 0\end{pmatrix}X, \quad \text{(B2)}$$

where Majorana phases are given by $\alpha_i = \arg(a_i)$. We used

$$X = \begin{pmatrix} 0.520 & -0.520 & -0.474 & 0.484\\ -0.712 & -0.284 & 0.165 & 0.621\\ -0.425 & -0.476 & -0.522 & -0.566\\ 0.206 & -0.650 & 0.689 & -0.244\end{pmatrix}, \quad \text{(B3)}$$

where $0 < a_4 < a_1 < a_2 < a_3$. The ordering of eigenvalues a_i is preferred to suppress $y_{\ell i}$ (in order to satisfy a constraint from $\mu \to e\gamma$ search) for the normal mass ordering for neutrinos ($m_1 < m_2 < m_3$). With this ansatz, small neutrino masses are preferred to suppress $\mathrm{BR}(\mu \to e\gamma)$.

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