Letter

Gauged $U(1)_X$ breaking as origin of neutrino masses, dark matter and leptogenesis at TeV scale

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We propose a new mechanism which simultaneously explains tiny neutrino masses, stability of dark matter and baryon asymmetry of the Universe via leptogenesis due to the common origin; a spontaneous breaking of a $U(1)_X$ gauge symmetry at TeV scale. The $U(1)_X$ breaking provides small Majorana masses of vectorlike leptons which generate small mass differences among them, and enhance their CP-violating decays via the resonant effect. Such CP-violation and lepton-number violation turns out to be a sufficient amount of the observed baryon asymmetry through leptogenesis. The Majorana masses from the $U(1)_X$ breaking also induce radiative generation of masses for active neutrinos at one-loop level. Furthermore, a Z_2 symmetry appears as a remnant of the $U(1)_X$ breaking, which guarantees the stability of dark matter. We construct a simple renormalizable model to realize the above mechanism, and show a benchmark point which can explain observed neutrino oscillations, dark matter data and the baryon asymmetry at the same time.

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I. INTRODUCTION

Neutrino oscillations, existence of dark matter, and baryon asymmetry of the Universe have been well-known and established phenomena which cannot be explained in the standard model (SM) of particle physics. Thus, there is no doubt about the necessity for new physics beyond the SM. So far, a plethora of models have been proposed to explain these phenomena, some of which can simultaneously explain all of them.

One of the simplest such new physics models is that with right-handed neutrinos. Masses and mixings of active neutrinos can be explained by the type-I seesaw mechanism [1–4]. Assuming one of the right-handed neutrinos to be odd under a Z_2 symmetry, it can be a candidate of dark matter. In addition, decays of the Z_2-even right-handed neutrinos can generate CP-violation and the lepton number, which can be converted into the baryon asymmetry of the Universe via the sphaleron process [5], that is, the leptogenesis scenario [6]. In this scenario, right-handed neutrinos "unify" the explanation of three phenomena at the same time. Although this simple model works well, its experimental probe is generally quite challenging because masses of the right-handed neutrinos typically have to be of order 10^{10} GeV or larger, see e.g., [7]. Furthermore, the *ad hoc* Z_2 symmetry does not originate from dynamics.

In this paper, we propose a new mechanism which simultaneously explains tiny neutrino masses, stability of dark matter and the baryon asymmetry via leptogenesis, in which all of them originate from a spontaneous breaking of a $U(1)_X$ gauge symmetry at the TeV scale. In our scenario, vectorlike leptons with nonzero $U(1)_X$ charges are introduced, which have Dirac masses at tree level. After the $U(1)_X$ breaking, Majorana masses for these vectorlike leptons appear, by which small mass differences are generated in their mass eigenstates. Such a mass difference can enhance CP-violating (CPV) decays of the heavy neutral leptons due to the resonant effect [8,9], and then the sufficient amount of the baryon asymmetry is explained via the leptogenesis. The $U(1)_X$ breaking also provides a Z_2 symmetry as a remnant, by which stability of the lightest Z_2-odd particle is guaranteed, and it can be a candidate of dark matter. Furthermore, the Majorana masses for the vectorlike leptons and the Z_2 symmetry realize the so-called scotogenic mechanism [10], where tiny masses for active neutrinos are generated at one-loop level. Effectively,

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our scenario is similar to the scotogenic model with the low-scale leptogenesis [11–16], but the CPV decay of heavy Majorana fermions is enhanced by not only the resonant effect but also sizable Yukawa couplings associated with a scalar field whose vacuum expectation value (VEV) breaks the $U(1)_X$ symmetry.

In the following, we construct a simple model to realize the above-mentioned mechanism, and give successful benchmark points to explain current neutrino data and the observed baryon asymmetry of the Universe.

II. MODEL

The content of new fields is shown in Table I, where all the SM fields have the same quantum number as those in the SM.[1] The relevant new terms in the following discussion are given by

$$\begin{aligned}-\mathcal{L}_{\rm rel} = {}& M_a\overline{N_L^a}N_R^a + y_\eta^{ia}\overline{L_L^i}(i\sigma_2\eta^*)N_R^a + \text{H.c.},\\ &+ y_L^{ab}\overline{N_L^{ac}}N_L^a\varphi + y_R^{ab}\overline{N_R^{ac}}N_R^a\varphi + \text{H.c.},\\ &+ \sum_{\Phi=\eta,\chi,\varphi} m_\Phi^2|\Phi|^2 + \left(\mu_1\eta^\dagger H\chi + \frac{\mu_2}{2}\varphi\chi\chi + \text{H.c.}\right),\end{aligned} \tag{1}$$

where L_L^i $(Wi = 1-3)$ and H are the ith generation of the SM lepton doublet and the Higgs doublet, respectively. The superscript c denotes the charge conjugation, and σ_2 is the second Pauli matrix. The Dirac masses M_a $(a = 1, 2)$ can be taken to be diagonal with real and positive values by the biunitarity transformation of N^a. In this basis, the Yukawa matrices y_η and $y_{L,R}$ are generally complex, where the latter are symmetric due to the $SU(2)_L$ structure. The phases of μ_1 and μ_2 can be removed by rephasing χ and η without loss of generality.

It is important to mention here that a nonzero value of $\mu_1\mu_2$ explicitly breaks the global lepton-number symmetry $U(1)_L$. In other words, if we take μ_1 and/or μ_2 to be zero, the theory recovers the $U(1)_L$ symmetry, in which the lepton number of φ can be taken to be an arbitrary value by choosing those of η, χ and N^a appropriately. This means that the $U(1)_L$ symmetry becomes exact for the case with $\mu_1\mu_2 = 0$, and thus Majorana masses for the left-handed neutrinos vanish as we will discuss below.

[1]Setup of our model is similar to that given in Ref. [16], in which an isospin triplet scalar field is introduced to close the η-loop in the one-loop diagram for neutrino masses. In our scenario, the μ_1 and μ_2 terms given in Eq. (1) play the similar role without introducing triplets. In addition, in [16], the lepton-number asymmetry is mainly produced via the resonant Z' effect in two-body to two-body scatterings, while such scatterings are negligibly small in our model, and the lepton-number asymmetry is mainly produced via the decay of heavy Majorana fermions.

TABLE I. Charge assignments under the $SU(2)_L \times U(1)_Y \times U(1)_X$ gauge symmetry, where N^a $(a = 1, 2)$ are the vectorlike leptons and all the others are scalar fields.

Fields	N^a	η	χ	φ
$SU(2)_L$	**1**	**2**	**1**	**1**
$U(1)_Y$	0	1/2	0	0
$U(1)_X$	1	1	1	−2

The $U(1)_X$ symmetry is broken down by the VEV $\langle\varphi\rangle = v_\varphi/\sqrt{2}$. Assuming the VEVs of η and χ to be zero, a Z_2 symmetry remains as the remnant of the $U(1)_X$ symmetry, where fields with an odd number of the $U(1)_X$ charge, i.e., N^a, η, and χ are Z_2-odd while all the other fields are even. Then, the lightest Z_2-odd particle can be a candidate of dark matter.

The neutral components of the Z_2-odd scalars, $\eta^0 = (\eta_H + i\eta_A)/\sqrt{2}$ and $\chi = (\chi_H + i\chi_A)/\sqrt{2}$, are mixed with each other due to the μ_1 term. We define the mass eigenstates of these scalar fields as

$$\begin{pmatrix}\eta_X\\ \chi_X\end{pmatrix} = \begin{pmatrix}\cos\theta_X & -\sin\theta_X\\ \sin\theta_X & \cos\theta_X\end{pmatrix}\begin{pmatrix}X_1\\ X_2\end{pmatrix}, \tag{2}$$

with $X = H, A$. We note that in the limit $\mu_1 \to 0$, these mixing angles become zero and $m_{H_1} = m_{A_1}$, while in the limit $\mu_2 \to 0$, these mixing angles become identical and $m_{H_i} = m_{A_i}$ $(i = 1, 2)$.

After the $U(1)_X$ breaking, N^a obtain the Majorana masses and their mass term is expressed as

$$\mathcal{L}_{\rm mass} = -\frac{1}{2}\overline{\tilde{\Psi}_R^c}\tilde{M}_\Psi\tilde{\Psi}_R + \text{H.c.}, \tag{3}$$

where $\tilde{\Psi}_R \equiv (N_L^{1c}, N_R^1, N_L^{2c}, N_R^2)^T$, and $\tilde{M}_\Psi$ is the 4×4 mass matrix given as

$$\tilde{M}_\Psi = \begin{pmatrix}(\delta m_L^*)_{11} & M_1 & (\delta m_L^*)_{12} & 0\\ M_1 & (\delta m_R)_{11} & 0 & (\delta m_R)_{12}\\ (\delta m_L^*)_{12} & 0 & (\delta m_L^*)_{22} & M_2\\ 0 & (\delta m_R)_{12} & M_2 & (\delta m_R)_{22}\end{pmatrix}. \tag{4}$$

In the mass matrix, we introduce $\delta m_{L,R} \equiv \sqrt{2}v_\varphi y_{L,R}$. We can diagonalize the mass matrix by introducing the 4×4 unitary matrix V as $M_\Psi \equiv V^T\tilde{M}_\Psi V = \text{diag}(m_{\psi_1}, m_{\psi_2}, m_{\psi_3}, m_{\psi_4})$ with $m_{\psi_4} \geq m_{\psi_3} \geq m_{\psi_2} \geq m_{\psi_1}$. The mass eigenstates are then given by

$$\Psi \equiv (\psi_1, \psi_2, \psi_3, \psi_4)^T = V\tilde{\Psi}. \tag{5}$$

III. NEUTRINO MASS

Majorana masses for active neutrinos are generated from the one-loop diagram shown in Fig. 1. The mass matrix is calculated as

$$(m_\nu)_{ij} = \sum_{I=1}^{4} \frac{Y_\eta^{*iI} Y_\eta^{*jI}}{32\pi^2} m_{\psi_I} \sum_{X=H,A} p^X \times \left(\frac{m_{X_1}^2 \cos^2\theta_X}{m_{X_1}^2 - m_{\psi_I}^2} \ln \frac{m_{X_1}^2}{m_{\psi_I}^2} + \frac{m_{X_2}^2 \sin^2\theta_X}{m_{X_2}^2 - m_{\psi_I}^2} \ln \frac{m_{X_2}^2}{m_{\psi_I}^2} \right), \tag{6}$$

where $p^H(p^A) = 1(-1)$, and $Y_\eta^{iI} = y_\eta^{i1} V^{2I} + y_\eta^{i2} V^{4I}$. We note that the above matrix vanishes for $\mu_1\mu_2 = 0$, because the $U(1)_L$ symmetry is recovered in this limit. This can explicitly be shown by using the properties mentioned just below Eq. (2). Therefore, both μ_1 and μ_2 should be nonzero in order to obtain finite neutrino masses. In addition, the mass matrix also vanishes in the limit of $v_\varphi \to 0$. Although this can be seen by looking at Fig. 1, but we can explicitly show it as follows. In this limit, the mass matrix for $\tilde{\Psi}_R$ given in Eq. (4) becomes a block diagonal form, and we obtain $m_{\psi_1} = m_{\psi_2}$ and $m_{\psi_3} = m_{\psi_4}$. At the same time, the unitary matrix V becomes a simple form of

$$V \to \frac{1}{\sqrt{2}} \begin{pmatrix} -i & 1 & 0 & 0 \\ i & 1 & 0 & 0 \\ 0 & 0 & -i & 1 \\ 0 & 0 & i & 1 \end{pmatrix}. \tag{7}$$

Using the above matrix and the mass degeneracy, we can show that the contributions from ψ_1 and ψ_2 (ψ_3 and ψ_4) are exactly canceled. This, however, does not mean that active neutrinos become massless, because the $U(1)_L$ symmetry is explicitly broken at Lagrangian level as mentioned above. In fact, we can find higher-loop contributions to the Majorana mass for active neutrinos, and one of such examples is shown in Fig. 2. Throughout this paper, we do not take into account such higher-loop contributions, and suppose that the one-loop contribution given in Eq. (6) is dominant. We also note that our mass matrix has rank 2, so that the lightest neutrino becomes massless.

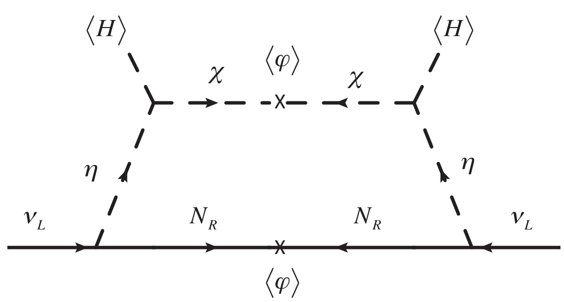

FIG. 1. One-loop generation of neutrino masses.

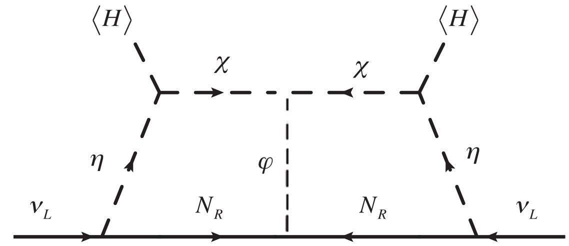

FIG. 2. Example of higher-loop contributions to neutrino masses which do not vanish in the limit $v_\varphi \to 0$.

IV. LEPTOGENESIS

In our scenario, the lepton-number density n_L or the $B-L$ number density n_{B-L} can be generated through the decay of the Majorana fermions ψ_I shown in Fig. 3, if we take nonzero CPV phases of the Yukawa couplings, and if the decay occurs in the out-of-thermal equilibrium. The produced lepton number is then converted into the baryon number density n_B via the sphaleron process [5] according to the following equation [17]:

$$n_B = \frac{8}{23} n_{B-L}, \tag{8}$$

which is derived by using the relation given by the chemical equilibrium for the sphaleron process, Yukawa interactions (except for that with the vectorlike leptons), and conservation of the hypercharge. For the discussion of leptogenesis, we neglect the mixing effect shown in Eq. (2) for simplicity, which does not essentially change the conclusion.

We first consider the out-of-equilibrium decay of ψ_I whose amount can be described by introducing the following efficiency parameter K_I [18]:

$$K_I \equiv \frac{\langle \Gamma(\psi_I \to L\eta) \rangle_{T=m_{\psi_I}}}{2H(m_{\psi_I})} \sim \frac{|\bar{Y}_\eta|^2}{16\pi\sqrt{g_*}} \frac{M_{\rm pl}}{m_{\psi_I}}, \tag{9}$$

where $H(m_{\psi_I})$ is the Hubble parameter at the temperature T to be m_{ψ_I}, $\bar{Y}_\eta$ is the averaged value of the Yukawa couplings Y_η^{Ii}, $M_{\rm pl} = 1.22 \times 10^{19}$ GeV is the Planck mass, and $g_* \simeq 124$ is the effective massless degrees of freedom assuming all the particles in our model being massless. In Eq. (9), we introduced the thermally averaged decay rate defined as

$$\langle \Gamma(\psi_I \to L\eta) \rangle \equiv \sum_{\rm fin} [\Gamma(\psi_I \to L\eta) + \Gamma(\psi_I \to L^c\eta^c)] \frac{\mathcal{K}_1(m_{\psi_I}/T)}{\mathcal{K}_2(m_{\psi_I}/T)}, \tag{10}$$

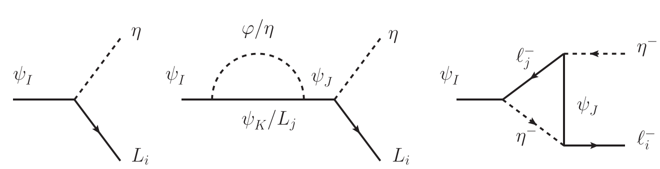

FIG. 3. CP-violating decays of ψ_I.

where $\Gamma(\psi_I \to XY)$ denotes the decay rate of $\psi_I \to XY$ with $\sum_{\rm fin}$ denoting the summation for all the possible final states, i.e., isospin components and lepton flavors, and $\mathcal{K}_n(z)$ are the modified Bessel functions of the nth kind. This K_I parameter can be of order one, i.e., the decay rate is compatible with the expansion rate of the Universe and provide sizable amount of out-of thermal equilibrium, for $m_{\psi_I} = \mathcal{O}(10)$ TeV and $\bar{Y}_\eta = \mathcal{O}(10^{-6})$. However, to reproduce the active neutrino masses to be $\mathcal{O}(0.1)$ eV, $\bar{Y}_\eta$ has to be of order 10^{-4} or larger, so that typically we obtain $K_I > 10^4$, which corresponds to the so-called strong washout regime. The yield for the baryon number $Y_B \equiv n_B/s$ with s being the entropy density is roughly estimated as $Y_B \sim -8/23 \times 0.3\epsilon/[g_* K(\ln K)^{0.6}]$ [18] with ϵ and K to be $\max(|\epsilon_I|)$ describing the amount of CP-violation defined in Eq. (11) and the corresponding K_I value, respectively.[2] Thus, we need a larger value of ϵ parameter, typically $\mathcal{O}(10^{-2})$, to compensate the suppression factor of $1/K$. For the actual calculation of Y_B, we numerically solve the Boltzmann equations, as discussed below.

Next, we discuss the CPV decay of ψ_I. As in the ordinal leptogenesis scenario, the CPV effect appears from the interference between the tree diagram and the one-loop diagrams shown in Fig. 3 at leading order. The amount of the CP-violation is expressed by introducing the following asymmetric parameter:

$$\epsilon_I \equiv \frac{\sum_{\rm fin}[\Gamma(\psi_I \to L\eta) - \Gamma(\psi_I \to L^c\eta^c)]}{\sum_{\rm fin}[\Gamma(\psi_I \to L\eta) + \Gamma(\psi_I \to L^c\eta^c)]}. \tag{11}$$

The magnitude of these ϵ_I parameters is determined by two types of the Yukawa couplings, i.e., y_η and $y_{L,R}$. As aforementioned, the magnitude of y_η has to be of order 10^{-4} to reproduce the active neutrino masses and to avoid a too strong wash-out of the generated lepton number, while $y_{L,R}$ can be of order one. Thus, the contribution from the φ loop shown in Fig. 3 is dominated with respect to the η loop one, and hence we can safely ignore the vertex correction shown as the third diagram. The self-energy diagram (the second one in Fig. 3) can also be enhanced by using the resonant effect of the intermediate ψ_J if a small mass difference between ψ_I and ψ_J is taken, because the amplitude is proportional to $(m_{\psi_I}^2 - m_{\psi_J}^2)^{-1}$. We note that in order to make the contribution from the self-energy diagram with φ-loop nonzero, the sum of the masses of ψ_K and φ must be smaller than that of ψ_I, because a nonzero value of ϵ_I requires both "weak phase" coming from the imaginary part of the Yukawa coupling and the "strong phase" coming from loop functions. The latter becomes nonzero when the particles in the loop are on shell.

[2]This expression gives a good approximation particularly for $K = K_I \gg K_{I'}$ with $m_{\psi_I} > m_{\psi_{I'}}$.

V. NUMERICAL RESULTS

To evaluate Y_B, we numerically solve the following set of Boltzmann equations:

$$\frac{dY_I}{dz} = -\frac{z}{H(m_{\psi_1})}\left[\langle\Gamma(\psi_I \to L\eta)\rangle(Y_I - Y_I^{\rm eq}) + \sum_{J\neq I}\langle\Gamma(\psi_I \to \psi_J\varphi)\rangle\left(Y_I - \frac{Y_I^{\rm eq}}{Y_J^{\rm eq}}Y_J\right)\right], \tag{12}$$

$$\frac{dY_L}{dz} = \frac{z}{H(m_{\psi_1})}\sum_I\langle\Gamma(\psi_I \to L\eta)\rangle \times \left[(Y_I - Y_I^{\rm eq})\epsilon_I - \frac{Y_L}{2Y_{\rm rel}^{\rm eq}}Y_I^{\rm eq}\right], \tag{13}$$

where $z = m_{\psi_1}/T$ and Y_I (Y_L) is the yield for ψ_I (lepton number). The values of $Y_I^{\rm eq}$ and $Y_{\rm rel}^{\rm eq}$ respectively denote the yields for ψ_I and a relativistic SM lepton given in the thermal equilibrium,

$$Y_I^{\rm eq} = \frac{45}{2\pi^4}\frac{z^2}{g_*}\mathcal{K}_2\left(z\frac{m_{\psi_1}}{m_{\psi_I}}\right), \qquad Y_{\rm rel}^{\rm eq} = \frac{45}{2\pi^4}\frac{3}{2g_*}\zeta(3), \tag{14}$$

where $\zeta(3) \simeq 1.2$ is the zeta function.

In order to show the typical behavior of ϵ_I and Y_B, we consider the following simplified input parameters

$$(M_1, M_2, v_\varphi, m_\eta, m_\chi, m_\varphi) = (10, 15, 1, 0.3, 0.3, 0.3)\ \text{TeV},$$
$$y_L^{11} = y_R^{11} = e^{-i\pi/4}, \qquad y_L^{12} = -y_R^{12} =: y_{12},$$
$$y_L^{22} = y_R^{22} =: y_{22}, \qquad y_\eta^{i1} = r y_\eta, \qquad y_\eta^{i2} = y_\eta. \tag{15}$$

As mentioned above, the magnitude of y_η should be of order 10^{-4}–10^{-5} to reproduce the active neutrino masses and to avoid too strong washout. For $y_{22} \ll 1$, we obtain $m_{\psi_3} \simeq m_{\psi_4} \simeq M_2$ with $(m_{\psi_4} - m_{\psi_3})/M_2 \ll 1$. In this case, $\epsilon_{3,4}$ are significantly enhanced from the second diagram in Fig. 3 with $\psi_{I,J} = \psi_{3,4}$ and $\psi_K = \psi_{1,2}$ due to the resonance between ψ_3 and ψ_4 as well as the large CPV effect coming from the larger Yukawa coupling $y_{L,R}^{11}$. On the other hand, the φ-loop contribution to $\epsilon_{1,2}$ are kinematically suppressed, so that the η-loop contribution is dominated. Therefore, we obtain $|\epsilon_{3,4}| \gg |\epsilon_{1,2}|$.

In Fig. 4, we show the contour plot for the value of $|\epsilon_3|$ as a function of y_{22} and y_{12}. We note that $\epsilon_{1,2} \sim 0$ and $\epsilon_4 \sim \epsilon_3$. As expected, larger $|\epsilon_3|$ is realized for smaller y_{22}, because the resonant effect of ψ_3-ψ_4 becomes stronger. This enhancement is, however, terminated at some values of y_{22} depending on the value of y_{12}, because of the effect of the finite width of $\psi_{3,4}$. We also see that the dependence on y_{12} is also significant, which determines the size of the connection between the $\psi_{1,2}$ sector and the $\psi_{3,4}$ sector. We find that the $\epsilon_{3,4}$ parameters can be of order 1 at, for instance, $(y_{12}, y_{22}) \sim (10^{-1.5}, 10^{-5})$.

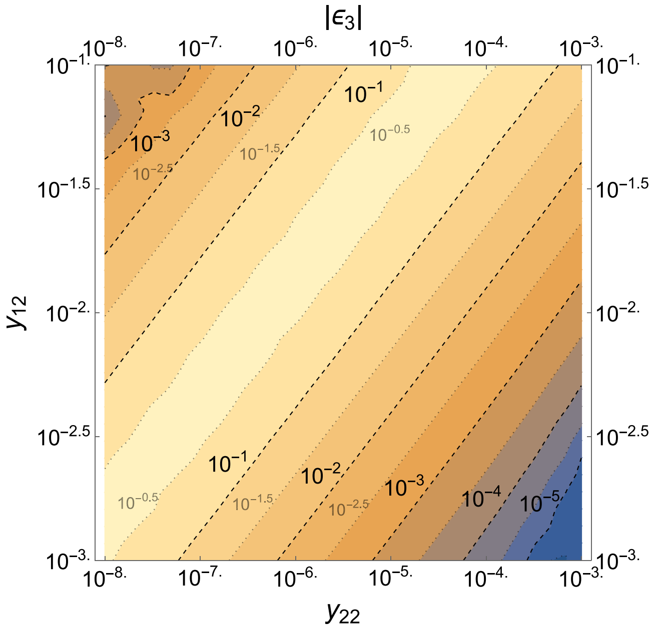

FIG. 4. Contour plot for $|\epsilon_3|$ with $y_\eta = 10^{-4}$ and $r = 1$.

In Fig. 5, we show the contour plot for $|Y_B|/Y_B^{\rm obs}$ as a function of r and y_η, where $Y_B^{\rm obs} = 8.7 \times 10^{-11}$ [19] is the observed value of the present baryon number of the Universe. For a fixed value of r, we see that $|Y_B|$ significantly becomes larger for smaller y_η, because the sizable out-of-equilibrium decay of $\psi_{3,4}$ is realized. We also see that smaller r gives larger $|Y_B|$. This is because for $r > 1$ the decays of $\psi_{1,2}$ can be more active than those of $\psi_{3,4}$, i.e., $K_{1,2} > K_{3,4}$, so that the produced lepton number from the former decay is washed out by the latter with negligibly small $\epsilon_{1,2}$. For $r \ll 1$, such washout does not happen as the decays of $\psi_{1,2}$ are already decoupled from the thermal equilibrium, and thus the produced lepton number from the decays of $\psi_{3,4}$ is kept.

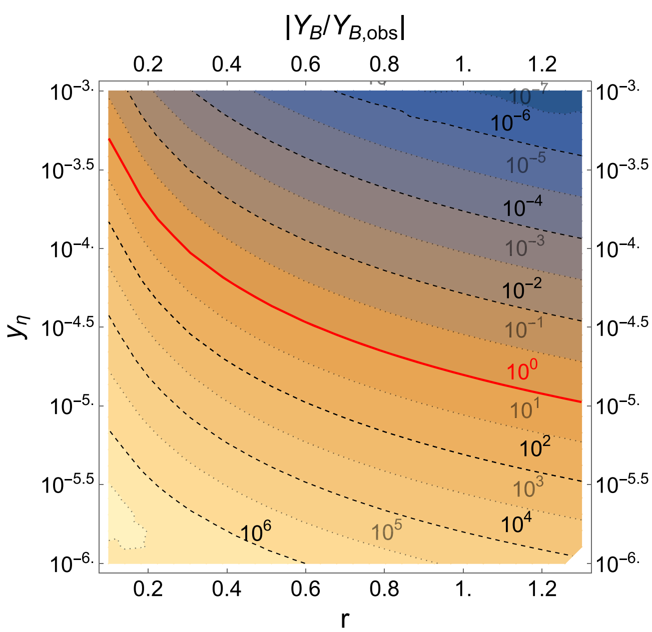

FIG. 5. Contour plot for $|Y_B|/Y_B^{\rm obs}$ with $y_{12} = 10^{-1.5}$ and $y_{22} = 10^{-5}$.

Finally, we would like to show a concrete benchmark point which satisfies the observed neutrino oscillation data and the baryon asymmetry as follows:

$$y_\eta = \begin{bmatrix} (3.89 + 1.84i)10^{-6} & (-7.38 + 1.10i)10^{-4} \\ (-8.58 + 4.79i)10^{-5} & (-3.58 + 1.44i)10^{-5} \\ (1.63 + 3.02i)10^{-5} & (5.43 - 12.8i)10^{-4} \end{bmatrix},$$
$$\mu_1 = 19.3~\text{GeV}, \qquad \mu_2 = 22.1~\text{GeV}, \tag{16}$$

while all the other inputs are taken to be the same way as in Eq. (15). We then obtain $Y_B = 8.6 \times 10^{-11}$ and

$$\Delta m_{21}^2 = 6.94 \times 10^{-5}~\text{eV}^2, \quad \Delta m_{31}^2 = 2.51 \times 10^{-3}~\text{eV}^2,$$
$$s_{\{12,23,13\}} = \{0.524, 0.783, 0.143\}, \quad \delta_{\rm CP} = 117°, \tag{17}$$

where s_{ij} indicates $\sin\theta_{ij}$, and all the values given in Eq. (17) are within the 3σ range of the global-fit results [20]. We check that the prediction of the lepton-flavor violating decays given in the above benchmark is much smaller than the current upper limit. For instance, the branching ratio of the $\mu \to e\gamma$ decay is given to be $\mathcal{O}(10^{-30})$ due to the small y_η values. We also find Z_2 odd neutral scalar masses and mixing as $\{m_{H_1}, m_{H_2}, m_{A_1}, m_{A_2}\} = \{326, 299, 301, 271\}$ GeV and $\sin\theta_A = -\sin\theta_H \simeq 0.1$.

VI. DISCUSSIONS AND CONCLUSIONS

We briefly discuss dark matter physics in the model. The dark matter candidate in our model is the lightest Z_2 odd scalar boson since new fermions ψ_I are heavier to realize the successful leptogenesis scenario discussed above. For example, in our benchmark, the lightest one is A_2 that dominantly comes from the imaginary component of χ. The scalar boson A_2 interacts with the SM gauge bosons similar to scalar dark matter given in the scotogenic model, but the coupling is suppressed by the factor $\sin\theta_A$. The annihilation cross section via electroweak processes, $A_2A_2 \to W^+W^-/ZZ$, is typically given by $\langle\sigma v\rangle \sim 10^{-2} \times (\sin\theta_A/0.1)^4 \times (100~\text{GeV}/m_{A_2})^2$ pb [21]. Thus, the cross section is too small for $\sin\theta_A = 0.1$ to explain the observed dark matter relic density, i.e., $\Omega h^2 \sim 0.12$ [22], which corresponds to $\langle\sigma v\rangle \sim 0.1$ pb [18]. We can, however, accommodate the observed relic density from the annihilation process via $A_2A_2 \to Z'Z'$ with Z' being the $U(1)_X$ gauge boson. The annihilation cross section is roughly given by $\langle\sigma v\rangle \sim 3(g_X\cos\theta_A)^4/(64\pi m_{A_2}^2) \sim 0.1 \times (g_X\cos\theta_A/0.2)^4 \times (300~\text{GeV}/m_{A_2})^2$ pb with g_X being the $U(1)_X$ gauge coupling, so that $\Omega h^2 \sim 0.12$ can be reproduced by taking $g_X \simeq 0.2$ with $\cos\theta_A \simeq 1$. Regarding the constraint from dark matter direct detections, gauge interactions only induce inelastic scatterings between dark matter and nucleus at tree level, whose cross section is negligibly small in our benchmark point with a mass difference of $\mathcal{O}(10)$ GeV between dark matter and the other Z_2-odd scalars. Although the process via the Higgs portal interaction can be important, such a coupling can be taken to be appropriately small to avoid the current upper limit on the cross section.

Finally, let us mention the collider phenomenology to test our scenario. One of the promising signatures would be $pp \to \varphi \to Z'Z' \to 4\ell$ at LHC with ℓ being $e^{\pm}$, $\mu^{\pm}$, which can be realized when the mass of φ is larger than twice the Z' mass. The Z' boson can decay into a pair of SM fermions via the kinetic mixing term in the Lagrangian, and the branching ratio of $Z' \to \ell^+\ell^-$ can be sizable if the Z' mass is a few 100 MeV. We leave more detailed phenomenological studies including dark matter physics and collider analyses as future projects [23].

In conclusion, we have proposed a simple model at TeV scale which can explain neutrino oscillations, stability of dark matter and baryon asymmetry of the Universe via leptogenesis from the common origin: the spontaneous breaking of the $U(1)_X$ symmetry. We have shown in Figs. 4 and 5 the typical orders of the magnitude for Yukawa couplings that are required for the successful leptogenesis scenario and generation of neutrino masses, and then we have presented a concrete benchmark point satisfying the neutrino oscillation data, dark matter data and the observed baryon asymmetry of the Universe.

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