Masses of dark matter and neutrino from TeV-scale spontaneous $U(1)_{B\text{-}L}$ breaking

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We propose a simple testable model with mass generation mechanisms for dark matter and neutrino based on the gauged $U(1)_{B\text{-}L}$ symmetry and an exact Z_2 parity. The $U(1)_{B\text{-}L}$ symmetry is spontaneously broken at the TeV scale, by which Z_2-odd right-handed neutrinos receive Majorana masses of the electroweak scale. The lightest one is a dark matter candidate, whose stability is guaranteed by the Z_2 parity. Resulting lepton number violation is transmitted to the left-handed neutrinos ν_L^i via the loop-induced dimension-six operator. Consequently, the tiny masses of ν_L^i can be generated without excessive fine-tuning. The observed dark matter abundance can be reproduced by the pair annihilation via the s-channel scalar exchange due to mixing of neutral components of Φ and S, where Φ and S, respectively, represent the Higgs doublet and the additional scalar singlet with the B-L charge. The model can be tested at collider experiments as well as flavor experiments through the discriminative predictions such as two light neutral Higgs bosons with large mixing, invisible decays of the Higgs bosons as well as the B-L gauge boson, and lepton flavor violation.

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Physics of electroweak symmetry breaking (EWSB), which is responsible for generating masses of weak bosons as well as quarks and charged leptons, is the last unknown part in the standard model (SM) for elementary particles. Its exploration is one of the top priorities at the Fermilab Tevatron and the CERN Large Hadron Collider (LHC).

Despite its great success, the SM cannot explain several phenomena established experimentally. First, previous neutrino oscillation experiments have clarified that neutrinos have tiny masses (less than 1 eV), which are much lower than the electroweak scale, 100 GeV. Such a difference in mass scales may indicate that neutrino masses are of Majorana type. Second, the WMAP data have shown that more than one fifth of the energy density of the Universe is occupied by dark matter. If the nature of dark matter is a weakly interacting massive particle (WIMP), the observed thermal relic abundance naturally suggests that the dark matter mass is around the EW scale. Such a mass would be generated by the physics just above the scale of the EWSB, the TeV scale.

We here address the following questions:

(i) What is the origin of the mass scale of neutrinos and the WIMP dark matter?
(ii) What is the relation between these scales and that of the EWSB?

If tiny masses of the left-handed neutrinos are of Majorana type, they would be generated as higher dimensional operators $LL\Phi\Phi/\Lambda$ at low energy where L, Φ, and Λ are, respectively, the lepton doublet field, the Higgs doublet field, and a dimensionful parameter. Such operators can be realized at tree-level in three different ways; 1) the exchange of $SU(2)_W$ singlet right-handed (RH) neutrinos [1], 2) that of $SU(2)_W$ triplet scalar fields [2], and 3) $SU(2)_W$ triplet fermions [3]. In these tree-level mechanisms, very large masses of RH neutrinos or triplet scalars/fermions as compared to the scale of EWSB are required if the coupling constants are not too small. The lowered masses are allowed in so-called radiative seesaw models, where the neutrino masses are generated at one-loop [4,5], two-loop [6], and three-loop level [7,8], or in seesaw models with higher order (dimension >5) operators [9].

In a class of the radiative seesaw models [5,7,8], a Z_2 parity is imposed to RH neutrinos to forbid the Yukawa coupling for neutrinos at tree-level. The Z_2 parity also plays a role to stabilize the dark matter candidate. The scale of Majorana masses of the RH neutrinos is that of lepton number violation. If it comes from spontaneous breaking of an additional gauge symmetry such as $U(1)_{B\text{-}L}$, its breaking would be at the TeV scale in these models with TeV-scale RH neutrinos.

In this paper, we consider a simple scenario to explain the mass scales of the WIMP dark matter and neutrino simultaneously, in which $U(1)_{B\text{-}L}$ is spontaneously broken at the TeV scale by developing the vacuum expectation value (VEV) of the scalar field S. The Z_2-odd RH neutrinos then receive the mass $m_{N_R} \sim y_R\langle S\rangle$ of the electroweak scale [10,11], and the lightest one is the dark matter whose stability is guaranteed by the Z_2 parity. The tiny neutrino masses are generated via the one-loop induced dimension-six operator of $LL\Phi\Phi S/\Lambda^2$. The generated mass can be written as

$$m_{\nu_L} \sim c\left(\frac{1}{16\pi^2}\right)\left(\frac{v}{M}\right)^2\langle S\rangle, \tag{1}$$

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where $v \simeq 246$ GeV, M represents the mass scale of the heaviest particle in the loop diagram and the coefficient c is a dimensionless parameter. As compared to the simplest seesaw scenario, generated neutrino masses have the suppression factors $1/(16\pi^2)$ and (m_{N_R}/M) (if $m_{N_R} \ll M$), so that for M to be the TeV scale the smallness of neutrino masses can be more naturally explained. Consequently, spontaneous breaking of the $U(1)_{B\text{-}L}$ is Mother of mass for both dark matter and neutrino.

We here propose the minimal model where the scenario described above is realized without contradicting the current data. The model is invariant under the gauge symmetry $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_{B\text{-}L}$. Although the gauge group in this model is the same as that in Ref. [12], we introduce neither additional fermions nor a nonrenormalizable term in Lagrangian. Instead, the unbroken Z_2 parity is imposed in our model. In addition to the B-L gauge boson Z', we introduce the second $SU(2)_W$ scalar doublet η which is Z_2 odd, the singlet scalar boson S with a B-L charge and three Z_2-odd RH neutrinos N_R^α ($\alpha = 1$–3). The particle properties under these symmetries are summarized in Table I.

The interaction part in the model is described as

$$\mathcal{L}_{\rm int} = \mathcal{L}_{\rm Yukawa}^{\rm SM} + \mathcal{L}_N - V(\Phi, \eta, S), \tag{2}$$

where $\mathcal{L}_{\rm Yukawa}^{\rm SM}$ is the SM Yukawa interaction, and

$$\mathcal{L}_N = \sum_{\alpha=1}^{3}\left(\sum_{i=1}^{3} g_{i\alpha}\bar{L}^i\tilde{\eta}N_R^\alpha - \frac{y_R^\alpha}{2}\bar{N}_R^\alpha S N_R^\alpha + \text{H.c.}\right), \tag{3}$$

with $\tilde{\eta} = i\tau_2\eta^*$. Without loss of generality, y_R can be taken to be flavor diagonal. Under the Z_2 parity, neutrino Yukawa couplings among L, Φ and N_R^α are forbidden. The scalar potential $V(\Phi, \eta, S)$ is given by

$$\begin{aligned} V(\Phi, \eta, S) = {} & +\mu_1^2|\Phi|^2 + \mu_2^2|\eta|^2 + \mu_S^2|S|^2 + \lambda_1|\Phi|^4 \\ & + \lambda_2|\eta|^4 + \lambda_3|\Phi|^2|\eta|^2 + \lambda_4|\Phi^\dagger\eta|^2 \\ & + \frac{\lambda_5}{2}[(\Phi^\dagger\eta)^2 + \text{H.c.}] + \lambda_S|S|^4 + \tilde{\lambda}|\Phi|^2|S|^2 \\ & + \lambda|\eta|^2|S|^2, \end{aligned} \tag{4}$$

where λ_5 can be taken as real. We assume that μ_1^2 and μ_S^2 are negative while μ_2^2 is positive.

The $U(1)_{B\text{-}L}$ symmetry is spontaneously broken when the S develops the VEV $\langle S\rangle = v_S/\sqrt{2}$ at the TeV scale [10,11]. We assume for simplicity that kinetic mixing [13] between $U(1)_Y$ and $U(1)_{B\text{-}L}$ gauge bosons is small such that it satisfies the EW precision measurement. The Z' boson then acquires the mass $M_{Z'}^2 = 4g_{B\text{-}L}^2 v_S^2$ where $g_{B\text{-}L}$ is the gauge coupling constant for $U(1)_{B\text{-}L}$. From the LEP experiment, the lower bound on the Z' boson mass has been found to be $v_S \gtrsim 3$–3.5 TeV [14,15]. Recent bound from Tevatron is comparable to the LEP bound [16]. The Z_2-odd RH neutrinos N_R^α also obtain masses from the $U(1)_{B\text{-}L}$ breaking as

$$m_{N_R^\alpha} = -y_R^\alpha\frac{v_S}{\sqrt{2}}. \tag{5}$$

After the EWSB, the neutral components of the Z_2 even scalar fields are parameterized as

$$\Phi^0 = \frac{1}{\sqrt{2}}(v + \phi + iz), \qquad S = \frac{1}{\sqrt{2}}(v_S + \phi_S + iz_S), \tag{6}$$

where z and z_S are the Nambu-Goldstone bosons absorbed by the longitudinal modes of the weak and $U(1)_{B\text{-}L}$ gauge bosons Z and Z', respectively. We define the ratio of the two VEVs as $\tan\beta = v_S/v$. The mass matrix of ϕ and ϕ_S is diagonalized by a mixing angle α;

$$\begin{pmatrix} h \\ H \end{pmatrix} = \begin{pmatrix} \cos\alpha & -\sin\alpha \\ \sin\alpha & \cos\alpha \end{pmatrix}\begin{pmatrix} \phi \\ \phi_S \end{pmatrix}, \tag{7}$$

where h and H represent the eigenstates corresponding to the mass eigenvalues

$$m_h^2 = 2(\lambda_1 c_\alpha^2 + \lambda_S s_\alpha^2\tan^2\beta - \tilde{\lambda}s_\alpha c_\alpha\tan\beta)v^2, \tag{8}$$

$$m_H^2 = 2(\lambda_1 s_\alpha^2 + \lambda_S c_\alpha^2\tan^2\beta + \tilde{\lambda}s_\alpha c_\alpha\tan\beta)v^2, \tag{9}$$

where $c_\alpha = \cos\alpha$ and $s_\alpha = \sin\alpha$. The mixing angle α is a free parameter. The constraints on the Higgs mixing from the precision measurements have been derived in Ref. [17]. The Yukawa interactions for h and H with N_R^α are then given by

$$\mathcal{L}_{\rm yukawa} = -\frac{1}{2}y_R^\alpha\bar{N}_R^\alpha\frac{(-h\sin\alpha + H\cos\alpha)}{\sqrt{2}}N_R^\alpha + \cdots. \tag{10}$$

The component fields of the Z_2-odd isospin doublet η are parametrized as

TABLE I. Particle properties.

	Q^i	d_R^i	u_R^i	L^i	e_R^i	Φ	η	S	N_R^α
$SU(3)_C$	3	3	3	1	1	1	1	1	1
$SU(2)_W$	2	1	1	2	1	2	2	1	1
$U(1)_Y$	1/6	−1/3	+2/3	1/2	−1	1/2	1/2	0	0
$U(1)_{B\text{-}L}$	1/3	1/3	1/3	−1	−1	0	0	+2	−1
Z_2	+	+	+	+	+	+	−	+	−

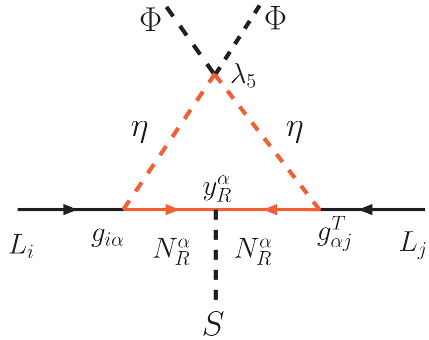

FIG. 1 (color online). The one-loop diagram of the dimension-six operator which generates neutrino masses.

$$\eta = \begin{pmatrix} H'^{+} \\ \frac{1}{\sqrt{2}}(H' + iA') \end{pmatrix}, \qquad (11)$$

whose masses are mainly determined by the invariant mass parameter μ_2 as long as $\mu_2^2 \gg \lambda_i v^2 \sim \lambda v_S^2$.

The tiny neutrino masses are generated via the one-loop induced dimension-six operator $LLS\Phi\Phi$ shown in Fig. 1. The induced mass matrix is evaluated as

$$m_{\nu_L}^{ij} \simeq \frac{\lambda_5}{8\pi^2}\left(\sum_{\alpha=1}^{3} g_{i\alpha} y_R^\alpha g_{\alpha j}^T\right)\left(\frac{v}{m_{\phi'}}\right)^2 v_S, \qquad (12)$$

for $m_{\phi'}^2 \gg m_{N_R^\alpha}^2$, where $m_{\phi'}$ denotes the scale of the quasi degenerate masses of H' and A'. For $v/m_{\phi'} \sim 10^{-1}$ and $\lambda_5 \sim g_{i\alpha} \sim y_R^\alpha \sim 10^{-2}$, the correct mass scale (~ 0.1 eV) can be realized from the TeV scale v_S. It is easy to see that the observed neutrino oscillation data can be reproduced by the result in Eq. (12), because the flavor structure is the same as that of the standard seesaw mechanism [18].

The dark matter candidates are the lightest RH neutrino N_R^1 and the lightest neutral component of the Z_2-odd doublet (H' or A'). In the model by Ma [5], the thermal relic density has been estimated for $H'(A')$ [19] as well as N_R^1 [20]. However, the scenario for the N_R^1 dark matter suffers from the constraint from lepton flavor violation (LFV) such as $\mu \to e\gamma$. The sufficient annihilation of dark matter requires large couplings $g_{i\alpha}$ whereas the null result for LFV gives the upper bound on these couplings [20]. To avoid the difficulty, additional mechanism such as coannihilation has to be introduced [21]. On the contrary, in our model with the spontaneously broken $U(1)_{B\text{-}L}$ symmetry, the neutral component of S can largely mix with the SM Higgs boson ϕ. Consequently, N_R^1 can annihilate via $N_R^1 N_R^1 \to h(H) \to f\bar{f}$ [22], so that the required relic abundance $\Omega_{N_R^1} h^2 \simeq 0.11$ can easily be attained when $m_{N_R^1} \sim m_h/2$ or $m_H/2$ without contradicting the LFV constraints. Therefore, N_R^1 can be the dark matter. In Fig. 2, the thermal relic abundance of N_R^1 is shown as a function of m_{N_R} for $m_h = 100(120)$ GeV, $m_H = 140(160)$ GeV, $\sin\alpha = 1/\sqrt{2}$ and $\tan\beta = 15$. It is seen that the relic abundance of N_R^1 is significantly reduced around $m_{N_R} = 50(60)$ and $70(80)$ GeV, respectively, and becomes consistent with the observed abundance of the dark matter. The RH neutrinos annihilate mainly via the S-channel exchange of the Higgs scalars. The annihilation of the RH neutrinos into the SM particles is resonantly enhanced through S-channel exchange of h or H when $m_h \simeq 2m_{N_R}$ and $m_h \simeq 2m_{N_R}$. Note that RH neutrinos can also annihilate via the B-L gauge boson exchange. However this process is less important because the cross section of the Z' exchange, $\langle\sigma v\rangle$, is proportional to $1/v_S^4$ and is much smaller than that of the Higgs exchange.

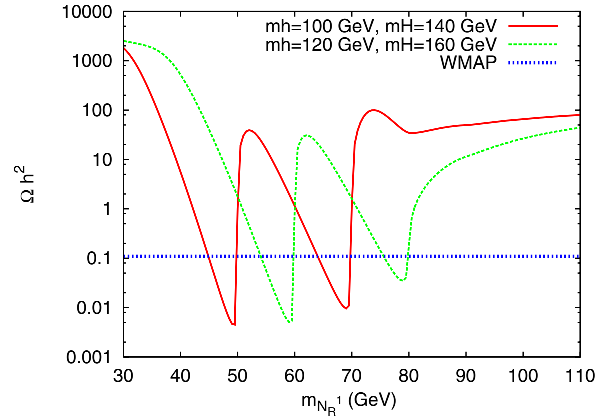

FIG. 2 (color online). The abundance of N_R^1 as a function of the mass $m_{N_R^1}$. The masses of h and H are taken to be $m_h =$ 100 GeV and $m_H = 140$ GeV (red curve), and $m_h = 120$ GeV and $m_H = 160$ GeV (green curve) for $\sin\alpha = 1/\sqrt{2}$ and $\tan\beta = 15$.

As a successful scenario, we consider the following parameters. First, v_S is supposed to be about 3.7 TeV. Second, m_h, m_H and α are taken as 100 GeV, 120 GeV and $\pi/4$, respectively. Third, the Z' mass is assumed to be 0.5–1 TeV, and the masses of H', A' and H'^{+} are commonly taken to be a few TeV. Finally, $m_{N_R^1}$ and $m_{N_R^{2,3}}$ are taken as 46 GeV and a few hundred GeV, respectively.

Phenomenological predictions are discussed in order:

(I) A characteristic feature of the Higgs bosons h and H with large mixing is that all the coupling constants of h (H) to the SM particles are multiplied by $\sin\alpha$ ($\cos\alpha$) as compared to the SM ones. This is similar to the Type-I two Higgs doublet model (THDM) without charged Higgs bosons. For maximal mixing $\sin\alpha \sim 1/\sqrt{2}$, the visible widths of h and H are about a half of that for the SM Higgs boson. The decay pattern can be discriminated from the other types of the Yukawa coupling in the THDM [23]. Production rates of h (H) at the Tevatron and the LHC are about 50% smaller than the SM value.

(II) The Higgs bosons h and H also couple to the dark matter N_R^1, so that they decay into a dark matter pair if kinematically allowed [24]. In Fig. 3, the branching ratios of the decays of h and H are shown as a function of m_h and m_H. $\tan\beta = 15$ and 12 is taken in Fig. 3(a) and 3(b), respectively. The solid curve for $N_R^1 N_R^1$ represents the decay branching ratio of h, while the dashed one does that of H. As shown in Fig. 2, the relic abundance of N_R^1 can be consistent with that of the dark matter when its mass is slightly smaller or larger than the half of the mass of h or H. We assume here that the mass of N_R^1 is slightly smaller than the half of the mass of h. The mass of N_R^1 is fixed as the half of the mass of h minus 4 GeV for the decay of h while it is fixed as 46 GeV for that of H. The invisible decay of h and H can reach to 0.7% (2.8%) for $h(H)$ in Fig. 3(a), and 1.1% (4.3%) in Fig. 3(b). For $v_S \sim 3.7$ TeV, which corresponds to $\tan\beta = 15$, the coupling constant y_R is determined as about 0.01 to obtain the correct mass $m_{N_R^1} \sim m_h/2$ or $m_H/2$. For $m_{N_R^1} = 46$ GeV, the invisible decay $h(H) \to N_R^1 N_R^1$ is then evaluated as 0.7% (2.6%). The invisible decay of the SM-like Higgs boson can be detected if it is larger than 25% at the LHC [25] and a few % at the International Linear Collider (ILC) [26], respectively. Feasibility of dark matter at colliders as well as direct searches is discussed in terms of the simple Higgs portal dark matter model in Ref. [27]. Studies on the RH neutrino production in the model without the Z_2 parity are seen in Refs. [16,28]. In our model, N_R^1 is dark matter, while heavier RH neutrinos N_R^2 and N_R^3 can be tested via the decay into N_R^1 with a lepton pair.

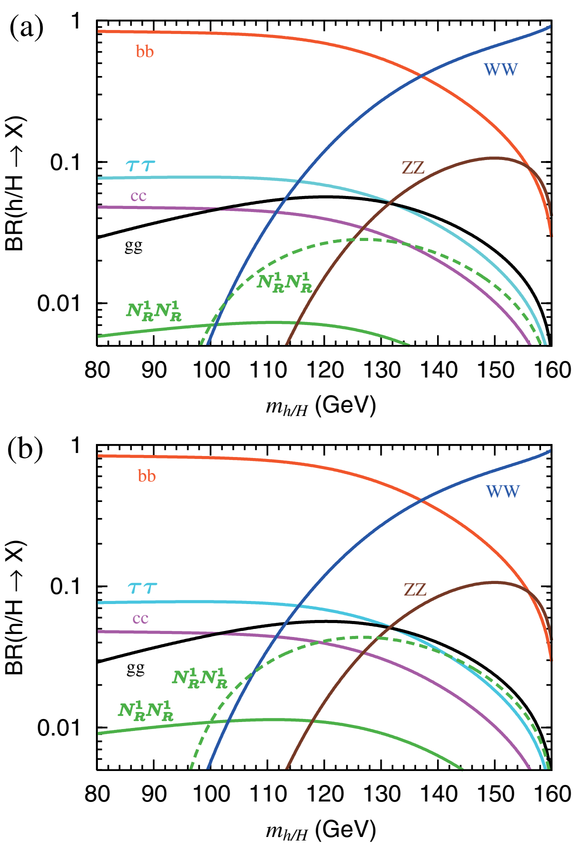

FIG. 3 (color online). Branching rations of h and H as a function of their masses. $\tan\beta$ is taken to be 15 in the figures (a) and 12 in (b), respectively. The solid curve for $N_R^1 N_R^1$ represents the branching ratio of h, while the dashed one does that of H. The mass of the lightest RH neutrino is fixed as the half of the mass of h for the decay of h, while it is fixed as 46 GeV for that of H.

(III) The existence of the Z' boson is another difference from the model by Ma [5]. Its production at the LHC in the minimal $U(1)_{B\text{-}L}$ model has been discussed [16,29]. If Z' is lighter than a few TeV, it would be detected at the LHC. The production cross section of Z' can be $\mathcal{O}(1)$ pb for $m_{Z'} \sim 1$ TeV. Decays of Z' into SM particles are proportional to the B-L charges. The branching ratios of $Z' \to q\bar{q}$, $\ell^+\ell^-$, $\nu_L\bar{\nu}_L$, $N_R N_R$ and $\phi_S\phi_S$ are given approximately by 0.2, 0.3, 0.15, 0.15 and 0.2, respectively. The branching ratio of $N_R^{2,3} \to N_R^1 \nu_L \nu_L$ is about 0.5. The invisible decay of Z' is then 0.225. This is a definite prediction in our model. It is expected that the invisible decay as well as characteristic branching ratios of visible decays can be tested at the LHC and the ILC.

(IV) An important constraint on radiative neutrino mass models comes from LFV processes such as $\mu \to e\gamma$. Since the typical scale of $g_{i\alpha}$ is of order of 10^{-2}, the present bound can be satisfied, but this would be testable by the new data from flavor experiments such as MEG.

(V) Phenomenology of the Z_2-odd scalar bosons H', A', $H'^{\pm}$ has been studied in various models [30]. In our present scenario, their masses are at the TeV scale. Although their direct detection would be difficult at the LHC, their indirect effects appear in the three-body decay of heavier right-handed neutrinos $N_R^{2(3)} \to N_R^1 \ell^- \ell^+$ that can be tested at colliders.

Therefore, the model can be tested and discriminated from the other radiative seesaw models, the minimal $U(1)_{B\text{-}L}$ model, and the other extended Higgs models such as THDMs.

Finally, we give a comment on the possibility of baryogenesis. In this minimal model, the Z_2-even part in the Higgs sector is composed of the isospin doublet Φ and the singlet S, so that there is no additional CP phase. In addition, coupling constants in the Higgs sector are so small that first-order phase transition cannot realize. Thus, for successful electroweak baryogenesis [31], the model has to be extended, for example, by introducing additional Z_2-even scalar doublets [8,32]. On the other hand, TeV-scale leptogenesis may be an alternative way. Various mechanisms have been proposed such as resonant leptogenesis [33] or three-body decay [34]. In these scenarios, however, considerable fine-tuning would be required.

We have discussed the minimal model with the mass generation mechanism for dark matter and neutrino based on the $U(1)_{B\text{-}L}$ symmetry and the Z_2 parity. Spontaneous $U(1)_{B\text{-}L}$ breaking at the TeV scale gives the Majorana masses of N_R^α, and the lightest one can be a WIMP dark matter. Its thermal relic abundance explains the WMAP result. Tiny neutrino masses are radiatively generated via the dimension-six operator without excessive fine-tuning. The model can be tested at future experiments. The detailed analyses are shown elsewhere [35].

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