Neutrino Mass, Dark Matter, and Baryon Asymmetry via TeV-Scale Physics without Fine-Tuning

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We propose an extended version of the standard model, in which neutrino oscillation, dark matter, and the baryon asymmetry of the Universe can be simultaneously explained by the TeV-scale physics without assuming a large hierarchy among the mass scales. Tiny neutrino masses are generated at the three-loop level due to the exact Z_2 symmetry, by which the stability of the dark matter candidate is guaranteed. The extra Higgs doublet is required not only for the tiny neutrino masses but also for successful electroweak baryogenesis. The model provides discriminative predictions especially in Higgs phenomenology, so that it is testable at current and future collider experiments.

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Although the standard model (SM) for elementary particles has been successful for over three decades, the Higgs sector remains unknown. The discovery of a Higgs boson is the most important issue at the CERN Large Hadron Collider (LHC). On the other hand, today we have definite reasons to consider a model beyond the SM. First of all, the data indicate that neutrinos have tiny masses and mix with each other [1]. Second, cosmological observations have revealed that the energy density of dark matter (DM) in the Universe dominates that of baryonic matter [2]. Finally, the asymmetry of matter and antimatter in our Universe has been addressed as a serious problem regarding the existence of ourselves [3]. They are all beyond the scope of the SM, so that an extension of the SM is required to explain these phenomena, which would be related to the physics of electroweak symmetry breaking.

A simple scenario to generate tiny masses (m_ν) for left-handed (LH) neutrinos would be based on the seesaw mechanism with heavy right-handed (RH) neutrinos [4]; $m_\nu \simeq m_D^2/M_R$, where M_R ($\sim 10^{13-16}$ GeV) is the Majorana mass of RH neutrinos, and m_D is the Dirac mass of the electroweak scale. This scenario would be compatible with the framework with large mass scales like grand unification. However, introduction of such large scales causes a problem of hierarchy. In addition, the decoupling theorem [5] makes it far from experimental tests.

In this Letter, we propose an alternative model which would explain neutrino oscillation, origin of DM and baryon asymmetry simultaneously by an extended Higgs sector with RH neutrinos. In order to avoid large hierarchy, masses of the RH neutrinos are to be at most TeV scales. Tiny neutrino masses are then generated at the three-loop level due to an exact discrete symmetry, by which tree-level Yukawa couplings of neutrinos are prohibited. The lightest neutral odd state under the discrete symmetry is a candidate of DM. Baryon asymmetry can be generated at the electroweak phase transition (EWPT) by additional *CP* violating phases in the Higgs sector [6,7]. In this framework, a successful model can be built without contradiction of the current data.

Original idea of generating tiny neutrino masses via the radiative effect has been proposed by Zee [8]. The extension with a TeV-scale RH neutrino has been discussed in Ref. [9], where the neutrino masses are generated at the three-loop due to the exact Z_2 parity, and the Z_2-odd RH neutrino is a candidate of DM. This has been extended with two RH neutrinos to describe the neutrino data [10]. Several models with adding baryogenesis have been considered in Ref. [11]. As compared to these models, the following new advantages are in the present model: (a) all mass scales are at most at the TeV scale without large hierarchy, (b) physics for generating neutrino masses is connected with that for DM and baryogenesis, (c) the model parameters are strongly constrained by the current data, so that the model gives discriminative predictions which can be tested at future experiments.

In addition to the *known* SM fields, particle entries are two scalar isospin doublets with hypercharge $1/2$ (Φ_1 and Φ_2), charged singlets ($S^\pm$), a real scalar singlet (η) and two generation isospin-singlet RH neutrinos (N_R^α with $\alpha = 1$, 2). In order to generate tiny neutrino masses at the three-loop level, we impose an exact Z_2 symmetry as in Ref. [9], which we refer to as Z_2. We assign the Z_2 odd charge to $S^\pm$, η, and N_R^α, while ordinary gauge fields, quarks and leptons and Higgs doublets are Z_2 even. Introduction of two Higgs doublets would cause a dangerous flavor changing neutral current. To avoid this in a natural way, we impose another discrete symmetry ($\tilde{Z}_2$) that is softly broken [12]. From a phenomenological reason discussed later, we assign $\tilde{Z}_2$ charges such that only Φ_1 couples to leptons whereas Φ_2 does to quarks;

$$\mathcal{L}_Y = -y_{e_i}\bar{L}^i\Phi_1 e_R^i - y_{u_i}\bar{Q}^i\tilde{\Phi}_2 u_R^i - y_{d_i}\bar{Q}^i\Phi_2 d_R^i + \text{H.c.}, \tag{1}$$

TABLE I. Particle properties under the discrete symmetries.

	Q^i	u_R^i	d_R^i	L^i	e_R^i	Φ_1	Φ_2	$S^\pm$	η	N_R^α
Z_2 (exact)	+	+	+	+	+	+	+	−	−	−
$\tilde{Z}_2$ (softly broken)	+	−	−	+	+	+	−	+	−	+

where Q^i (L^i) is the ordinary ith generation LH quark (lepton) doublet, and u_R^i and d_R^i (e_R^i) are RH-singlet up- and down-type quarks (charged leptons), respectively. We summarize the particle properties under Z_2 and $\tilde{Z}_2$ in Table I. Notice that the Yukawa coupling in Eq. (1) is different from that in the minimal supersymmetric SM [13]. The scalar potential is given by

$$\begin{aligned} V = &\sum_{a=1}^{2}(-\mu_a^2|\Phi_a|^2 + \lambda_a|\Phi_a|^4) - (\mu_{12}^2\Phi_1^\dagger\Phi_2 + \text{H.c.}) \\ &+ \lambda_3|\Phi_1|^2|\Phi_2|^2 + \lambda_4|\Phi_1^\dagger\Phi_2|^2 + \left\{\frac{\lambda_5}{2}(\Phi_1^\dagger\Phi_2)^2 + \text{H.c.}\right\} \\ &+ \sum_{a=1}^{2}\left(\rho_a|\Phi_a|^2|S|^2 + \sigma_a|\Phi_a|^2\frac{\eta^2}{2}\right) \\ &+ \sum_{a,b=1}^{2}\{\kappa\epsilon_{ab}(\Phi_a^c)^\dagger\Phi_b S^-\eta + \text{H.c.}\} + \mu_s^2|S|^2 \\ &+ \lambda_s|S|^4 + \mu_\eta^2\eta^2/2 + \lambda_\eta\eta^4 + \xi|S|^2\eta^2/2, \end{aligned} \qquad (2)$$

where ϵ_{ab} is the antisymmetric tensor with $\epsilon_{12} = 1$. The mass term and the interaction for N_R^α are given by

$$\mathcal{L}_Y = \sum_{\alpha=1}^{2}\left\{\frac{1}{2}m_{N_R^\alpha}\bar{N}_R^{\alpha c}N_R^\alpha - h_i^\alpha\overline{(e_R^i)^c}N_R^\alpha S^- + \text{H.c.}\right\}. \qquad (3)$$

In general, μ_{12}^2, λ_5, and κ (as well as h_i^α) can be complex. The phases of λ_5 and κ can be eliminated by rephasing $S^\pm$ and Φ_1. The remaining phase of μ_{12}^2 causes CP violation in the Higgs sector. Although the phase is crucial for successful baryogenesis at the EWPT [6], it does not much affect the following discussions. Thus, we neglect it for simplicity. We later give a comment on the case with the nonzero CP-violating phase.

As Z_2 is exact, the even and odd fields cannot mix. Mass matrices for the Z_2 even scalars are diagonalized as in the usual two Higgs doublet model (THDM) by the mixing angles α and β, where α diagonalizes the CP-even states, and $\tan\beta = \langle\Phi_2^0\rangle/\langle\Phi_1^0\rangle$ [13]. The Z_2 even physical states are two CP-even (h and H), a CP-odd (A) and charged ($H^\pm$) states. We here define h and H such that h is always the SM-like Higgs boson when $\sin(\beta - \alpha) = 1$.

The LH neutrino mass matrix M_{ij} is generated by the three-loop diagrams in Fig. 1. The absence of lower order loop contributions is guaranteed by Z_2. $H^\pm$ and e_R^i play a crucial role to connect LH neutrinos with the one-loop subdiagram by the Z_2-odd states. We obtain

$$M_{ij} = \sum_{\alpha=1}^{2} C_{ij}^\alpha F(m_{H^\pm}, m_{S^\pm}, m_{N_R^\alpha}, m_\eta), \qquad (4)$$

where $C_{ij}^\alpha = 4\kappa^2\tan^2\beta(y_{e_i}^{\text{SM}}h_i^\alpha)(y_{e_j}^{\text{SM}}h_j^\alpha)$ and

$$\begin{aligned} F(m_{H^\pm}, m_{S^\pm}, m_{N_R}, m_\eta) = &\left(\frac{1}{16\pi^2}\right)^3\frac{(-m_{N_R}v^2)}{m_{N_R}^2 - m_\eta^2}\int_0^\infty dx\left[x\left\{\frac{B_1(-x, m_{H^\pm}, m_{S^\pm}) - B_1(-x, 0, m_{S^\pm})}{m_{H^\pm}^2}\right\}^2 \right. \\ &\left. \times\left(\frac{m_{N_R}^2}{x + m_{N_R}^2} - \frac{m_\eta^2}{x + m_\eta^2}\right)\right], \qquad (m_{S^\pm}^2 \gg m_{e_i}^2), \end{aligned} \qquad (5)$$

with m_f representing the mass of the field f, $y_{e_i}^{\text{SM}} = \sqrt{2}m_{e_i}/v$, $v \simeq 246$ GeV and B_1 being the tensor coefficient function in Ref. [14]. Magnitudes of $\kappa\tan\beta$ as well as F determine the universal scale of M_{ij}, whereas variation of h_i^α ($i = e, \mu, \tau$) reproduces the mixing pattern indicated by the neutrino data [1]. M_{ij} is related to the data by $M_{ij} = U_{is}(M_\nu^{\text{diag}})_{st}(U^T)_{tj}$, where U_{is} is the unitary matrix and $M_\nu^{\text{diag}} = \text{diag}(m_1, m_2, m_3)$. Under the *natural* requirement $h_e^\alpha \sim \mathcal{O}(1)$, and taking the $\mu \to e\gamma$ search results into account [15], we find that $m_{N_R^\alpha} \sim \mathcal{O}(1)$ TeV, $m_{H^\pm} \lesssim \mathcal{O}(100)$ GeV, $\kappa\tan\beta \gtrsim \mathcal{O}(10)$, and $m_{S^\pm}$ being several times 100 GeV. On the other hand, the LEP direct search results indicate $m_{H^\pm}$ (and $m_{S^\pm}$) $\gtrsim 100$ GeV [1]. In addition, with the LEP precision measurement for the ρ parameter, possible values uniquely turn out to be $m_{H^\pm} \simeq m_H$ (or m_A) $\simeq 100$ GeV for $\sin(\beta - \alpha) \simeq 1$. Thanks to the Yukawa coupling in Eq. (1), such a light $H^\pm$ is not excluded by the $b \to s\gamma$ data [16]. Since we cannot avoid to include the hierarchy among y_i^{SM}, we only require $h_i^\alpha y_i \sim \mathcal{O}(y_e) \sim 10^{-5}$ for values of h_i^α. Several sets for h_i^α are shown in Table II with the predictions on the branching ratio of $\mu \to e\gamma$ assuming the normal hierarchy, $m_1 \simeq m_2 \ll m_3$ with $m_1 = 0$. For the inverted hierarchy ($m_3 \ll m_1 \simeq m_2$ with $m_3 = 0$), $\kappa\tan\beta$ is required to be larger. Our

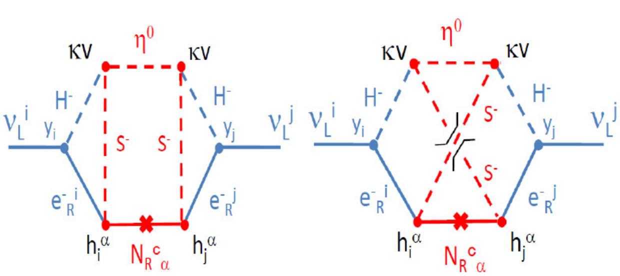

FIG. 1 (color online). The diagrams for generating tiny neutrino masses.

TABLE II. Values of h_i^α for $m_{H^\pm}(m_{S^\pm}) = 100(400)$ GeV $m_\eta = 50$ GeV, $m_{N_R^1} = m_{N_R^2} = 3.0$ TeV for the normal hierarchy. For Set A (B), $\kappa \tan\beta = 28(32)$ and $U_{e3} = 0(0.18)$. Predictions on the branching ratio of $\mu \to e\gamma$ are also shown.

Set	h_e^1	h_e^2	h_μ^1	h_μ^2	h_τ^1	h_τ^2	$B(\mu \to e\gamma)$
A	2.0	2.0	−0.019	0.042	−0.0025	0.0012	6.9×10^{-12}
B	2.2	2.2	0.0085	0.038	−0.0012	0.0021	6.1×10^{-12}

model turns out to prefer the normal hierarchy scenario [17].

The lightest Z_2-odd particle is stable and can be a candidate of DM if it is neutral. In our model, N_R^α must be heavy, so that the DM candidate is identified as η. When η is lighter than the W boson, η dominantly annihilates into $b\bar{b}$ and $\tau^+\tau^-$ via tree-level s-channel Higgs (h and H) exchange diagrams, and into $\gamma\gamma$ via one-loop diagrams. From their summed thermal averaged annihilation rate $\langle\sigma v\rangle$, the relic mass density $\Omega_\eta h^2$ is evaluated as

$$\Omega_\eta h^2 = 1.1 \times 10^9 \frac{(m_\eta/T_d)}{\sqrt{g_*} M_P \langle\sigma v\rangle} \text{ GeV}^{-1}, \qquad (6)$$

where M_P is the Planck scale, g_* is the total number of relativistic degrees of freedom in the thermal bath, and T_d is the decoupling temperature [18]. Fig. 2 shows $\Omega_\eta h^2$ as a function of m_η. Strong annihilation can be seen near 50 GeV $\simeq m_H/2$ (60 GeV $\simeq m_h/2$) due to the resonance of H (h) mediation. The data ($\Omega_{\rm DM}h^2 \simeq 0.11$ [2]) indicate that m_η is around 40–65 GeV.

The model satisfies the necessary conditions for baryogenesis [3]. Especially, departure from thermal equilibrium can be realized by the strong first order EWPT. The free energy is given at a high temperature T by [19]

$$V_{\rm eff}[\varphi, T] = D(T^2 - T_0^2)\varphi^2 - ET\varphi^3 + \frac{\lambda_T}{4}\varphi^4 + \ldots, \qquad (7)$$

where φ is the order parameter, and

$$E \simeq \frac{1}{12\pi v^3}(6m_W^3 + 3m_Z^3 + m_A^3 + 2m_{S^\pm}^3), \qquad (8)$$

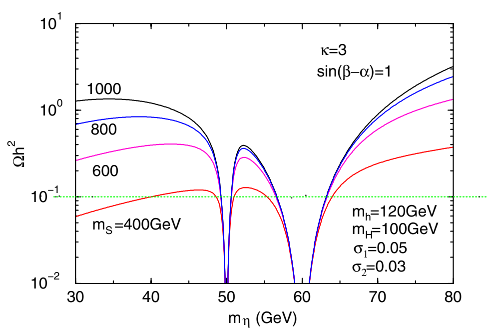

FIG. 2 (color online). The relic abundance of η.

with $D \simeq (6m_W^2 + 3m_Z^2 + 6m_t^2 + m_A^2 + 2m_{S^\pm}^2)/(24v^2)$, $T_0^2 \sim m_h^2/(4D)$ and $\lambda_T \sim m_h^2/(2v^2)$. A large value of the coefficient E is crucial for the strong first order EWPT [7]. In Eq. (8), quantum effects by h, H and $H^\pm$ are neglected since they are unimportant for $\sin(\beta - \alpha) \simeq 1$ and $m_{H^\pm} \simeq m_H \simeq M(\equiv \sqrt{2\mu_{12}^2/\sin 2\beta})$ (the soft $\tilde{Z}_2$ breaking scale [20]). For sufficient sphaleron decoupling in the broken phase, it is required that [21]

$$\frac{\varphi_c}{T_c}\left(\simeq \frac{2E}{\lambda_{T_c}}\right) \gtrsim 1, \qquad (9)$$

where φ_c ($\neq 0$) and T_c are the critical values of φ and T at the EWPT. In Fig. 3, the allowed region under the condition of Eq. (9) is shown. The condition is satisfied when $m_{S^\pm} \gtrsim 350$ GeV for $m_A \gtrsim 100$ GeV, $m_h \simeq 120$ GeV, $m_H \simeq m_{H^\pm}(\simeq M) \simeq 100$ GeV, $\mu_S \simeq 200$ GeV and $\sin(\beta - \alpha) \simeq 1$. Unitarity bounds are also satisfied unless m_A (m_S) is too larger than M (μ_S) [7,22].

A successful scenario which can simultaneously solve the above three issues under the data [1,15,16] would be

$$\begin{gathered} \sin(\beta-\alpha) \simeq 1, \quad \kappa\tan\beta \simeq 30, \quad m_h = 120 \text{ GeV}, \\ m_H \simeq m_{H^\pm}(\simeq M) \simeq 100 \text{ GeV}, \quad m_A \gtrsim 100 \text{ GeV}, \\ m_{S^\pm} \sim 400 \text{ GeV}, \quad m_\eta \simeq 40\text{–}65 \text{ GeV}, \\ m_{N_R^1} \simeq m_{N_R^2} \simeq 3 \text{ TeV}. \end{gathered} \qquad (10)$$

This is realized without assuming unnatural hierarchy among the couplings. All the masses are between $\mathcal{O}(100)$ GeV and $\mathcal{O}(1)$ TeV. As they are required by the data, the model has a predictive power. We note that the masses of A and H can be exchanged with each other.

We outline phenomenological predictions in the scenario in (10) in order. The detailed analysis is shown elsewhere [22]. (i) h is the SM-like Higgs boson, but decays into $\eta\eta$ when $m_\eta < m_h/2$. The branching ratio is about 36% (25%) for $m_\eta \simeq 45(55)$ GeV. This is related to the

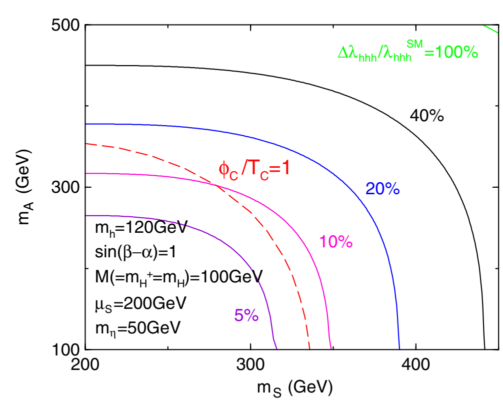

FIG. 3 (color online). The region of strong first order EWPT. Deviations from the SM value in the hhh coupling are also shown.

DM abundance, so that our DM scenario is testable at the LHC. (ii) η is potentially detectable by direct DM searches [23], because η can scatter with nuclei via the scalar exchange [24]. (iii) For successful baryogenesis, the hhh coupling has to deviate from the SM value by more than 10%–20% [7] (see Fig. 3), which can be tested at the International Linear Collider (ILC) [25]. (iv) H (or A) can predominantly decay into $\tau^+\tau^-$ instead of $b\bar{b}$ for $\tan\beta \gtrsim 3$. When A (or H) is relatively heavy it can decay into $H^\pm W^\mp$ and HZ (or AZ). (v) the scenario with light $H^\pm$ and H (or A) can be directly tested at the LHC via $pp \to W^* \to HH^\pm$ and $AH^\pm$ [26]. (vi) $S^\pm$ can be produced in pair at the LHC (the ILC) [27], and decay into $\tau^\pm \nu \eta$. The signal would be a hard hadron pair [28] with a large missing energy. (vii) The couplings h_i^α cause lepton flavor violation such as $\mu \to e\gamma$ which would provide information on $m_{N_R^\alpha}$ at future experiments.

Finally, we comment on the case with the CP violating phases. Our model includes the THDM, so that the same discussion can be applied in evaluation of baryon number at the EWPT [6]. The mass spectrum would be changed to some extent, but most of the features discussed above should be conserved with a little modification.

We have discussed the model solving neutrino oscillation, DM and baryon asymmetry by the TeV-scale physics without fine-tuning. It gives specific predictions in Higgs phenomenology, DM physics and flavor physics, so that it is testable at current and future experiments.

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