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Loop suppression of Dirac neutrino mass in the neutrinophilic two-Higgs-doublet model

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ABSTRACT

We extend the scalar sector of the neutrinophilic two-Higgs-doublet model, where small masses of Dirac neutrinos are obtained via a small vacuum expectation value v_ν of the neutrinophilic $SU(2)_L$-doublet scalar field which has a Yukawa interaction with only right-handed neutrinos. A global $U(1)_X$ symmetry is used for the neutrinophilic nature of the second $SU(2)_L$-doublet scalar field and also for eliminating Majorana mass terms of neutrinos. By virtue of an appropriate assignment of the $U(1)_X$-charges to new particles, our model has an unbroken Z_2 symmetry, under which the lightest Z_2-odd scalar boson can be a dark matter candidate. In our model, v_ν is generated by the one-loop diagram to which Z_2-odd particles contribute. We briefly discuss a possible signature of our model at the LHC.

1. Introduction

It has been well established that neutrinos have nonzero masses as shown in the neutrino oscillation measurements [1–6] although they are massless particles in the standard model (SM) of particle physics. Since the scale of neutrino masses is much different from that of the other fermion masses, they might be generated by a different mechanism from the one for the other fermions. Usually, the possibility that neutrinos are Majorana fermions is utilized as a characteristic feature of the neutrino masses. The most popular example is the seesaw mechanism [7] where very heavy right-handed Majorana neutrinos are introduced. However, lepton number violation which is caused by masses of the Majorana neutrinos has not been discovered. Thus it is worth considering the possibility that neutrinos are not Majorana fermions but Dirac fermions similarly to charged fermions.

The neutrinophilic two-Higgs-doublet model (νTHDM) is a new physics model where neutrinos are regarded as Dirac fermions. The second $SU(2)_L$-doublet scalar field which couples only with right-handed neutrinos ν_R was first introduced in Ref. [8] for Majorana neutrinos. Phenomenology in the model of Majorana neutrinos is discussed in Refs. [9,10]. The neutrinophilic doublet field is also utilized for Dirac neutrinos [11] where a spontaneously broken Z_2 parity is introduced in order to achieve the neutrinophilic property. Smallness of neutrino masses are explained by a tiny vacuum expectation value (VEV) of the neutrinophilic scalar without extremely small Yukawa coupling constant for neutrinos. Instead of the Z_2 parity, the model in Ref. [12] uses a global $U(1)_X$ symmetry that is softly broken in the scalar potential. The $U(1)_X$ symmetry forbids Majorana mass terms of ν_R, and then neutrinos are Dirac fermions.[1] We refer to the model in Ref. [12] as the νTHDM.

The new particle which was discovered at the LHC [13,14] is likely to be the SM Higgs boson [15–18]. It opens the new era of probing the origin of particle masses. Then it would be a natural desire to expect that the origin of neutrino masses are also uncovered. If the neutrinophilic scalars in the νTHDM exist within the experimentally accessible energy scale (namely the TeV-scale), decays of the neutrinophilic charged scalar into leptons can provide direct information on the neutrino mass matrix because it is proportional to the matrix of new Yukawa coupling constants for the neutrinophilic scalar field [12,19]. In such a case, the smallness of a new VEV which is relevant to Dirac neutrino masses is interpreted by the smallness of a soft-breaking parameter of the global $U(1)_X$ symmetry. It seems then better to have a suppression mechanism for the soft-breaking parameter by extending the νTHDM with TeV scale particles including a dark matter candidate. The existence of dark matter has also been established in cosmological observations [20,21], and it is an important guideline for constructing new physics models.

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[1] Since the Majorana mass terms of ν_R can also be acceptable as soft breaking terms of the $U(1)_X$, the lepton number conservation may be imposed to the Lagrangian.

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The reason why the neutrino masses are tiny can be explained by a mechanism that the interaction of neutrinos with the SM Higgs boson is generated via a loop diagram involving a dark matter candidate in the loop while the interaction is forbidden at the tree level [22–32]. Notice that smallness of neutrino masses in such radiative mechanisms does not require new particles to be very heavy. Similarly, if neutrino masses arise from a new VEV, smallness of neutrino masses can be explained by assuming that the VEV is generated at the loop level by utilizing a dark matter candidate [33]. In this Letter, we extend the νTHDM such that the new VEV is generated at the one-loop level (see also Ref. [34]) where a dark matter candidate is involved in the loop.

This Letter is organized as follows. We briefly introduce the νTHDM in Section 2. The νTHDM is extended in Section 3 such that a small VEV is generated via the one-loop diagram which involving a dark matter candidate in the loop. Section 4 is devoted to discussion on phenomenology in the extended νTHDM. We conclude in Section 5.

2. Neutrinophilic two-Higgs-doublet model

In the νTHDM, the SM is extended with the second $SU(2)_L$-doublet scalar field Φ_ν which has a hypercharge $Y=1/2$ and right-handed neutrinos ν_{iR} ($i=1$–3) which are singlet fields under the SM gauge group. A global $U(1)_X$ symmetry is introduced, under which Φ_ν and ν_{iR} have the same nonzero charge while the SM particles have no charge. Then, the Yukawa interaction with Φ_ν is only the following one:

$$\mathcal{L}_{\nu\text{-Yukawa}} = -(y_\nu)_{\ell i}\overline{L_\ell}\, i\sigma_2\, \Phi_\nu^*\, \nu_{iR} + \text{h.c.}, \tag{1}$$

where $\ell(=e,\mu,\tau)$ denotes the lepton flavor and σ_i ($i=1$–3) are the Pauli matrices. Since Majorana mass terms $\overline{(\nu_{iR})^c}\nu_{iR}$ are forbidden by the $U(1)_X$ symmetry, there appears an accidental conservation of the lepton number where lepton numbers of Φ_ν and ν_{iR} are 0 and 1, respectively. When the neutral component ϕ_ν^0 of Φ_ν develops its VEV v_ν $(\equiv \sqrt{2}\langle\phi_\nu^0\rangle)$, the neutrino mass matrix arise as $(m_\nu)_{\ell i} = v_\nu (y_\nu)_{\ell i}/\sqrt{2}$. We have taken a basis where ν_{iR} are mass eigenstates. Then the mass matrix m_ν is diagonalized as $U_{\text{MNS}}^\dagger m_\nu = \text{diag}(m_1,m_2,m_3)$, where m_i ($i=1$–3) are the neutrino mass eigenvalues and a unitary matrix U_{MNS} is the so-called Maki–Nakagawa–Sakata (MNS) matrix [35]. Dirac neutrinos are constructed as $\nu_i = (\sum_\ell (U_{\text{MNS}}^\dagger)_{i\ell}\nu_{\ell L}, \nu_{iR})^T$. Smallness of neutrino masses is attributed to that v_ν is much smaller than v.

If the VEV v_ν is generated spontaneously, a CP-odd scalar $\phi_{\nu i}^0$ becomes massless as a Nambu–Goldstone boson with respect to the breaking of $U(1)_X$, where $\phi_\nu^0 = (v_\nu + \phi_{\nu r}^0 + i\phi_{\nu i}^0)/\sqrt{2}$. In addition, a CP-even neutral scalar $\phi_{\nu r}^0$ has a small mass ($\propto v_\nu \ll v$). Therefore, the scenario of the spontaneous breaking of $U(1)_X$ is not allowed by the measurement of the invisible decay of the Z boson. The scalar potential in the νTHDM is given by

$$\begin{aligned} V^{(\nu\text{THDM})} &= -\mu_{\Phi 1}^2 \Phi^\dagger\Phi + \mu_{\Phi 2}^2 \Phi_\nu^\dagger\Phi_\nu - \left(\mu_{\Phi 12}^2 \Phi_\nu^\dagger\Phi + \text{h.c.}\right) \\ &\quad + \lambda_{\Phi 1}\left(\Phi^\dagger\Phi\right)^2 + \lambda_{\Phi 2}\left(\Phi_\nu^\dagger\Phi_\nu\right)^2 + \lambda_{\Phi 12}\left(\Phi^\dagger\Phi\right)\left(\Phi_\nu^\dagger\Phi_\nu\right) \\ &\quad + \lambda'_{\Phi 12}\left(\Phi^\dagger\Phi_\nu\right)\left(\Phi_\nu^\dagger\Phi\right), \end{aligned} \tag{2}$$

where $\mu_{\Phi 12}^2$ can be real and positive by using rephasing of Φ_ν without loss of generality; we take $\mu_{\Phi 1}^2 > 0$ and $\mu_{\Phi 2}^2 > 0$. The VEV of ϕ_ν^0 is triggered by $\mu_{\Phi 12}^2$ which softly breaks the $U(1)_X$ symmetry. Since the term does not breaks the lepton number conservation, neutrinos are still Dirac particles. Taking $v_\nu/v \ll 1$ into account, the VEVs are calculated as

Table 1
New particles which are added to the SM in our model.

	ν_{iR}	$\Phi_\nu = \begin{pmatrix}\phi_\nu^+\\ \phi_\nu^0\end{pmatrix}$	$\eta = \begin{pmatrix}\eta^+\\ \eta^0\end{pmatrix}$	s_1^0	s_2^0
$SU(2)_L$	$\underline{1}$	$\underline{2}$	$\underline{2}$	$\underline{1}$	$\underline{1}$
$U(1)_Y$	0	1/2	1/2	0	0
Global $U(1)_X$	3	3	3/2	1	1/2

$$v \simeq \frac{\mu_{\Phi 1}}{\sqrt{\lambda_{\Phi 1}}}, \qquad v_\nu \simeq \frac{2v\,\mu_{\Phi 12}^2}{2\mu_{\Phi 2}^2 + (\lambda_{\Phi 12} + \lambda'_{\Phi 12})v^2}. \tag{3}$$

If $\mu_{\Phi 2} \sim v$, we have $v_\nu \sim \mu_{\Phi 12}^2/v$. Then, $\mu_{\Phi 12}/v$ is required to be small ($\sim 10^{-6}$ for $y_\nu \sim 1$). Stability of the tiny v_ν is discussed in Refs. [10,36]. In our model presented in the next section, $\mu_{\Phi 12}/v$ becomes small because $\mu_{\Phi 12}^2$ is generated at the one-loop level.

3. An extension of the νTHDM

Since we try to generate $\mu_{\Phi 12}^2$ at the loop level, it does not appear in the Lagrangian. Then the $U(1)_X$ symmetry should be broken spontaneously. For the spontaneous breaking, we rely on an additional scalar s_1^0 which is a singlet field under the SM gauge group. Similarly to the singlet Majoron model [37] where a VEV of a singlet field spontaneously breaks the lepton number conservation by two units, the Nambu–Goldstone boson from s_1^0 is acceptable [37]; the Nambu–Goldstone boson couples first with only neutrinos among fermions. If $U(1)_X$-charges of Φ_ν and s_1^0 are 3 and 1, respectively, a dimension-5 operator $(s_1^0)^3\Phi_\nu^\dagger\Phi$ is allowed by the $U(1)_X$ symmetry although $\Phi_\nu^\dagger\Phi$ is forbidden. Then, $\mu_{\Phi 12}^2$ is generated from the dimension-5 operator with the VEV of s_1^0. In this Letter, we show the simplest realization of the dimension-5 operator at the one-loop level where dark matter candidates are involved in the loop.

Table 1 is the list of new particles added to the SM. In the table, ν_{iR} and Φ_ν are the particles which exist in the νTHDM. The $U(1)_X$ symmetry is spontaneously broken by the VEV of s_1^0. We take a scenario where η and s_2^0 do not have VEVs. Since their $U(1)_X$-charges are half-integers while the one for s_1^0 is an integer, a Z_2 symmetry remains unbroken after the $U(1)_X$ breaking. Here, η and s_2^0 are Z_2-odd particles. The Z_2 symmetry stabilizes the lightest Z_2-odd particle which can be a dark matter candidate.

The Yukawa interaction in this model is identical to those in the νTHDM (see Eq. (1)). The scalar potential in this model is expressed as

$$\begin{aligned} V &= -\mu_{s1}^2\left|s_1^0\right|^2 + \mu_{s2}^2\left|s_2^0\right|^2 - \mu_{\Phi 1}^2\Phi^\dagger\Phi + \mu_{\Phi 2}^2\Phi_\nu^\dagger\Phi_\nu + \mu_\eta^2\eta^\dagger\eta \\ &\quad - \left(\mu\, s_1^{0*}\left(s_2^0\right)^2 + \text{h.c.}\right) + \left(\lambda_{s\Phi 1\eta}\, s_1^{0*}\left(s_2^0\right)^*\Phi^\dagger\eta + \text{h.c.}\right) \\ &\quad + \left(\lambda_{s\Phi 2\eta}\, s_1^0 s_2^0 \Phi_\nu^\dagger\eta + \text{h.c.}\right) + \cdots. \end{aligned} \tag{4}$$

Only the relevant parts to our discussion are presented in Eq. (4). The other terms are shown in Appendix A. Parameters μ, $\lambda_{s\Phi 1\eta}$, and $\lambda_{s\Phi 2\eta}$ are taken to be real and positive values by rephasing of scalar fields without loss of generality. At the tree level, v_ν, v, and v_s $(=\sqrt{2}\langle s_1^0\rangle)$ are given by

$$\begin{aligned} &v_\nu = 0, \\ &\begin{pmatrix} v^2 \\ v_s^2 \end{pmatrix} = \frac{2}{4\lambda_{s1}\lambda_{\Phi 1} - \lambda_{s1\Phi 1}^2}\begin{pmatrix} 2\lambda_{s1} & -\lambda_{s1\Phi 1} \\ -\lambda_{s1\Phi 1} & 2\lambda_{\Phi 1}\end{pmatrix}\begin{pmatrix}\mu_{\Phi 1}^2 \\ \mu_{s1}^2\end{pmatrix}. \end{aligned} \tag{5}$$

The Z_2-odd scalar fields (η and s_2^0) result in the following particles: two CP-even neutral scalars ($\mathcal{H}_1^0$ and $\mathcal{H}_2^0$), two CP-odd neutral