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Unified picture for Dirac neutrinos, dark matter, dark energy and matter–antimatter asymmetry

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Abstract

We propose a unified scenario to generate the masses of Dirac neutrinos and cold dark matter at the TeV scale, understand the origin of dark energy and explain the matter–antimatter asymmetry of the universe. This model can lead to significant impact on the Higgs searches at LHC.

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Strong evidences from neutrino oscillation experiments have confirmed the tiny neutrino masses of the order of 10^{-2} eV [1]. However, the neutrino's Majorana or Dirac nature is still unknown. The smallness of the neutrino masses can be elegantly understood via the Majorana [2] or Dirac [3,4] seesaw mechanism in various extensions of the Standard Model (SM). The nature of the dark matter, which contributes about 20% [1] to the energy density of the universe, also indicates the necessity of supplementing to the existing theory. Currently many supersymmetric or nonsupersymmetric candidates [5–11] for the dark matter have been proposed to study and search for. As for the dark energy with the energy density $\sim (3 \times 10^{-3}\ \mathrm{eV})^4$ [1], which accelerates the expansion of our universe at present, it is striking that its scale is far lower than all the known scales in particle physics except that of the neutrino masses. The intriguing coincidence between the neutrino mass scale and the dark energy scale inspires us to consider them in a unified scenario, as in the neutrino dark energy model [12]. The origin of the observed matter–antimatter asymmetry [1] of the universe poses another big challenge to the SM, but within the Majorana or Dirac seesaw scenario, it can be naturally explained through leptogenesis [13] or neutrinogenesis [14].

In this Letter, we unify the mass origin for the Dirac neutrinos and the dark matter in a nonsupersymmetric extension of the SM. After a new $U(1)$ gauge symmetry is spontaneously broken at the TeV scale, the SM neutrinos will obtain small Yukawa couplings to the new right-handed neutrinos and the SM Higgs while other new introduced fermions, which guarantee the theory free of gauge anomaly, will acquire masses of a few hundred GeV. These new fermions with the right amount of the relic density can serve as the candidate for the cold dark matter. In order to understand the origin of the dark energy, we further introduce a proper global symmetry, after which is spontaneously broken near the Planck scale, a pseudo-Nambu–Goldstone boson (pNGB) associated with the neutrino mass-generation can explain the nature of the dark energy. Meanwhile, the matter–antimatter asymmetry can be resolved via the neutrinogenesis mechanism. This model predicts new Higgs phenomenology that can be tested at LHC.

To generate the masses for the Dirac neutrinos, we can simply introduce three right-handed neutrinos to the SM. However, the Yukawa couplings of the Dirac neutrinos should be extremely small. One possibility to naturally explain this phenomena is to consider the Dirac seesaw [3,4], in which the Yukawa couplings of the neutrinos to the SM Higgs are generated by integrating out some heavy particles, meanwhile, the conventional Yukawa couplings of the neutrinos to the SM

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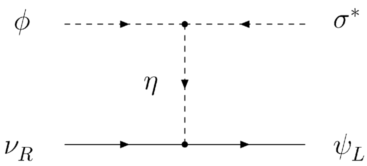

Fig. 1. The dim-5 operator for neutrino mass-generation.

Higgs should be exactly forbidden by consideration of symmetry. Here we consider a $U(1)_X$ gauge symmetry, under which the right-handed neutrinos carry the charge -1 while all SM particles transform trivially. We also introduce a new $SU(2)_L$ Higgs doublet which carries the same $U(1)_Y$ hypercharge with the SM Higgs but has the $U(1)_X$ charge $+1$. Thus the neutrinos can have the Yukawa couplings to this new Higgs,

$$\mathcal{L} \supset -y\overline{\psi_L}\eta\nu_R + \text{h.c.}, \tag{1}$$

where ψ_L, η and ν_R are the SM lepton doublets, new Higgs doublet and right-handed neutrinos, respectively. However, different from the SM Higgs, the new one has a positive quadratic term in the scalar potential so that it cannot develop a vacuum expectation value (*vev*) to generate the neutrino masses at this stage. Fortunately, we can conveniently introduce a SM Higgs singlet with the $U(1)_X$ charge $+1$ and then obtain a trilinear coupling among three types of Higgs fields,

$$\mathcal{L} \supset -\mu\sigma\eta^\dagger\phi + \text{h.c.}, \tag{2}$$

where σ and ϕ are the SM singlet and doublet Higgs scalars, respectively. Therefore, as shown in Fig. 1, we can obtain a dim-5 operator,

$$\mathcal{L} \supset \frac{\mu}{M_\eta^2} y\overline{\psi_L}\phi\nu_R\sigma + \text{h.c.} \tag{3}$$

by integrating out the new Higgs doublet. Once the $U(1)_X$ symmetry is spontaneously broken by the *vev*, $\langle\sigma\rangle$, the neutrinos will acquire the effective Yukawa couplings to the SM Higgs. It is straightforward to see that the effective Yukawa couplings can be highly suppressed by $\mu\langle\sigma\rangle/M_\eta^2$ and hence the neutrinos will obtain the tiny Dirac masses,

$$m_\nu = y_{\text{eff}}\langle\phi\rangle \equiv -\frac{\mu\langle\sigma\rangle}{M_\eta^2} y\langle\phi\rangle. \tag{4}$$

For instance, we find that by inputting

$$\begin{aligned} &\langle\phi\rangle \simeq 174~\text{GeV}, \qquad \langle\sigma\rangle = \mathcal{O}(\text{TeV}), \\ &\mu/M_\eta = \mathcal{O}(0.1), \quad M_\eta = \mathcal{O}\big(10^{13\text{–}15}~\text{GeV}\big), \end{aligned} \tag{5}$$

the Yukawa couplings of the neutrinos can naturally remain small, $y_{\text{eff}} \sim \mathcal{O}(10^{-13})$ for $y \sim \mathcal{O}(10^{-2}\text{–}1)$, and hence, the neutrino masses become of the order of $m_\nu \sim \mathcal{O}(10^{-2}~\text{eV})$, which is consistent with the neutrino oscillation data [1]. In fact, as shown in [4], by minimizing the full scalar potential, the new Higgs doublet will acquire a small *vev*,

$$\langle\eta\rangle \simeq -\frac{\mu\langle\sigma\rangle}{M_\eta^2}\langle\phi\rangle \tag{6}$$

with the range,

$$10^{-2}~\text{eV} \lesssim \langle\eta\rangle \lesssim 1~\text{eV} \tag{7}$$

for the parameters (5). This confirms Eq. (4) due to the mass formula,

$$m_\nu = y\langle\eta\rangle \tag{8}$$

from Eq. (1).

The requirement to ensure anomaly free indicates the necessity of supplementing the existing theory with three left-handed SM singlet fermions χ_L with the $U(1)_X$ charge $+1$. Under the present gauge symmetry, it is convenient to introduce three right-handed singlet fermions χ_R to generate the following Yukawa couplings,

$$\mathcal{L} \supset -f\sigma\overline{\chi_L}\chi_R + \text{h.c.}, \tag{9}$$

through which the singlet fermions will acquire masses,

$$m_\chi = f\langle\sigma\rangle, \tag{10}$$

after the $U(1)_X$ breaking by $\langle\sigma\rangle$. We further consider a Z_3 discrete symmetry, under which $\chi_{L,R}$ have the transformation properties $\chi_{L,R} \to \omega\chi_{L,R}$ with $\omega^3 = 1$ while all other fields are trivial. In consequence, the Yukawa coupling $\overline{\psi_L}\phi\chi_R$ and the Majorana mass term of χ_R are exactly forbidden. So, $\chi_{L,R}$ have not any decay modes and hence are inert. By diagonalizing the mass matrix (10), the inert fermions can be defined in their mass-eigenbasis $\chi_{1,2,3}$ with the corresponding masses,

$$m_{\chi_{1,2,3}} = f_{1,2,3}\langle\sigma\rangle, \tag{11}$$

where $f_1 \leqslant f_2 \leqslant f_3$ are the eigenvalues of matrix f. The inert fermions can serve as the dark matter if and only if their relic density is consistent with the cosmological observation.

Before calculating the relic density of the inert fermions, we need clarify the properties of the gauge and Higgs bosons in the present model since they are essential to the annihilation of the inert fermions. There exists a $U(1)_X$ gauge field C_μ in addition to the SM gauge fields B_μ and W_μ^i ($i = 1, 2, 3$). Since the new Higgs doublet with *vev* carries both $U(1)_Y$ and $U(1)_X$ charge, the $U(1)_X$ gauge field should mix with the SM ones. By diagonalizing the vector boson mass matrix, we obtain the charged bosons $W_\mu^\pm = \frac{1}{\sqrt{2}}(W_\mu^1 \mp iW_\mu^2)$ with the mass $m_W^2 = \frac{1}{2}g^2(\langle\phi\rangle^2 + \langle\eta\rangle^2)$, the photon $A_\mu = B_\mu\cos\theta + W_\mu^3\sin\theta$ with the mixing angle $\tan\theta = g'/g$ as well as the two massive neutral vector bosons Z_μ and Z'_μ,

$$Z_\mu = Z_\mu^0\cos\xi - C_\mu\sin\xi, \qquad Z'_\mu = Z_\mu^0\sin\xi + C_\mu\cos\xi \tag{12}$$

with the masses,

$$\begin{aligned} m_Z^2 = \big(g^2 + g'^2\big)\bigg\{&\big(\langle\sigma\rangle^2 + \langle\eta\rangle^2\big)\sin^2\theta + \frac{1}{4}\big(\langle\phi\rangle^2 + \langle\eta\rangle^2\big) \\ &- \bigg\{\bigg[\big(\langle\sigma\rangle^2 + \langle\eta\rangle^2\big)\sin^2\theta - \frac{1}{4}\big(\langle\phi\rangle^2 + \langle\eta\rangle^2\big)\bigg]^2 \\ &+ \langle\eta\rangle^4\sin^2\theta\bigg\}^{\frac{1}{2}}\bigg\}, \end{aligned} \tag{13}$$

$$m_{Z'}^2 = \left(g^2 + g'^2\right)\left\{\left(\langle\sigma\rangle^2 + \langle\eta\rangle^2\right)\sin^2\theta + \frac{1}{4}\left(\langle\phi\rangle^2 + \langle\eta\rangle^2\right) + \left\{\left[\left(\langle\sigma\rangle^2 + \langle\eta\rangle^2\right)\sin^2\theta - \frac{1}{4}\left(\langle\phi\rangle^2 + \langle\eta\rangle^2\right)\right]^2 + \langle\eta\rangle^4\sin^2\theta\right\}^{\frac{1}{2}}\right\} \tag{14}$$

and the mixing angle,

$$\sin 2\xi = \langle\eta\rangle^2 \sin\theta\left\{\left[\left(\langle\sigma\rangle^2 + \langle\eta\rangle^2\right)\sin^2\theta - \frac{1}{4}\left(\langle\phi\rangle^2 + \langle\eta\rangle^2\right)\right]^2 + \langle\eta\rangle^4\sin^2\theta\right\}^{-\frac{1}{2}}. \tag{15}$$

Here $Z_\mu^0 = -B_\mu \sin\theta + W_\mu^3 \cos\theta$ corresponds to the neutral vector boson of the SM. As shown in Eq. (7), $\langle\eta\rangle$ is much smaller than $\langle\phi\rangle$ and $\langle\sigma\rangle$, thus we obtain

$$m_Z^2 \simeq \frac{1}{2}\left(g^2 + g'^2\right)\langle\phi\rangle^2, \qquad m_{Z'}^2 \simeq 2g'^2\langle\sigma\rangle^2. \tag{16}$$

Meanwhile, the mixing angle (15) is tiny and hence free of the constraint from the precise measurement. In consequence, C_μ and Z_μ^0 can be approximately identified with Z'_μ and Z_μ, respectively.

Let us subsequently consider the Higgs sector,

$$V(\phi,\sigma,\eta) = -m_1^2\phi^\dagger\phi - m_2^2\sigma^\dagger\sigma + M_\eta^2\eta^\dagger\eta + \lambda_1\left(\phi^\dagger\phi\right)^2 + \lambda_2\left(\sigma^\dagger\sigma\right)^2 + \lambda_3\left(\eta^\dagger\eta\right)^2 + \frac{1}{2}\lambda_4\left(\phi^\dagger\phi\right)\left(\sigma^\dagger\sigma\right) + \frac{1}{2}\lambda_5\left(\phi^\dagger\phi\right)\left(\eta^\dagger\eta\right) + \frac{1}{2}\lambda_6\left(\sigma^\dagger\sigma\right)\left(\eta^\dagger\eta\right) + \mu\sigma\eta^\dagger\phi + \text{h.c.} \tag{17}$$

Similar to [4], we can deduce the *vevs*, $\langle\phi\rangle$, $\langle\sigma\rangle$ and $\langle\eta\rangle$ by minimizing the above scalar potential. For $\langle\eta\rangle \ll \langle\phi\rangle, \langle\sigma\rangle$, the contribution from η to σ and ϕ can be neglected, we thus have the two neutral bosons,

$$h = \frac{1}{\sqrt{2}}\phi - \langle\phi\rangle, \qquad h' = \frac{1}{\sqrt{2}}\sigma - \langle\sigma\rangle, \tag{18}$$

which are the linear combinations of the mass eigenstates h_1 and h_2,

$$h_1 = h\cos\beta - h'\sin\beta, \qquad h_2 = h\sin\beta + h'\cos\beta \tag{19}$$

with the masses,

$$m_{h_1}^2 = \lambda_1\langle\phi\rangle^2 + \lambda_2\langle\sigma\rangle^2 - \left[\left(\lambda_2\langle\sigma\rangle^2 - \lambda_1\langle\phi\rangle^2\right)^2 + 4\lambda_4^2\langle\sigma\rangle^2\langle\phi\rangle^2\right]^{\frac{1}{2}}, \tag{20}$$

$$m_{h_2}^2 = \lambda_1\langle\phi\rangle^2 + \lambda_2\langle\sigma\rangle^2 + \left[\left(\lambda_2\langle\sigma\rangle^2 - \lambda_1\langle\phi\rangle^2\right)^2 + 4\lambda_4^2\langle\sigma\rangle^2\langle\phi\rangle^2\right]^{\frac{1}{2}} \tag{21}$$

and the mixing angle,

$$\tan 2\beta = \frac{2\lambda_4\langle\sigma\rangle\langle\phi\rangle}{\lambda_2\langle\sigma\rangle^2 - \lambda_1\langle\phi\rangle^2}. \tag{22}$$

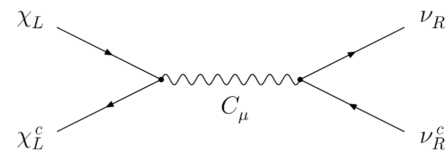

Fig. 2. The inert fermions annihilate into the right-handed neutrinos through the gauge couplings.

Similar to [15], here the couplings of h_1 and h_2 to the SM gauge bosons, quarks and charged leptons have the same structure as the corresponding Higgs couplings in the SM, however, their size is reduced by $\cos\beta$ and $\sin\beta$, respectively. For $\langle\sigma\rangle \simeq \mathcal{O}(\text{TeV})$, the mixing angle β and the mass splitting between h_1 and h_2 may be large. In consequence, there could be significant impact on the Higgs searches at LHC [16]. For example, the couplings of the lighter h_1 to the quarks and leptons would even vanish in the extreme case $\beta = \frac{\pi}{2}$.

We now discuss the possibility of the inert fermions as the dark matter. The pairs of the inert fermions have the gauge couplings to C_μ and the Yukawa couplings to h'. For $\langle\eta\rangle \ll \langle\phi\rangle, \langle\sigma\rangle$, C_μ can be looked on as the mass eigenstate Z'_μ. Furthermore, since the systematic analyses of the implication from the quartic interaction $\lambda_4(\sigma^\dagger\sigma)(\phi^\dagger\phi)$ on the relic density of the inert fermions will be presented elsewhere, we for simplicity take $\lambda_4 = 0$ and hence h and h' are exactly identified with h_1 and h_2 with the masses,

$$m_{h_1}^2 = 2\lambda_1\langle\phi\rangle^2, \qquad m_{h_2}^2 = 2\lambda_2\langle\sigma\rangle^2, \tag{23}$$

respectively. For the purpose of calculating the relic density of the inert fermions, we take $\langle\sigma\rangle = 720$ GeV and then give $m_{Z'} \simeq 360$ GeV, $m_{h'} \simeq 600$ GeV with $\lambda_2 = 0.35$, $m_{\chi_{1,2,3}} \simeq 200$ GeV with $f_{1,2,3} = 0.28$ by using Eqs. (16), (23) and (11), respectively. With this mass spectrum, as shown in Fig. 2, the channel of a pair of the inert fermions to a pair of the right-handed neutrinos should dominate the annihilation process of the inert fermions. We calculate

$$\sigma_i v = \frac{g'^4}{96\pi}\frac{s - m_{\chi_i}^2}{(s - m_{Z'}^2)^2 + m_{Z'}^2\Gamma_{Z'}^2}, \tag{24}$$

where σ_i is the annihilation cross section of a pair of χ_i to a pair of ν_R, v is the relative speed between the two χ_i's in their center-of-mass system (cms), s is the usual Mandelstam variable, and

$$\Gamma_{Z'} = \frac{g'^2}{24\pi}m_{Z'} \tag{25}$$

is the decay width of Z'. Comparing the annihilation rate $\Gamma_i = n_{\chi_i}^{eq}\langle\sigma_i v\rangle$ to the Hubble constant, we find that the freeze out should happen at the temperature $T_F \simeq 10$ GeV. Here $\langle\sigma_i v\rangle \simeq 2.6$ pb is the thermal-average cross section. The relic density of the inert fermions is then approximately given by

$$\Omega_\chi h^2 \simeq \sum_{i=1}^{3}\frac{0.1\ \text{pb}}{\langle\sigma_i v\rangle} \simeq 0.1, \tag{26}$$

which is equal to the right amount [1] of the relic density for the cold dark matter. Thus the inert fermions in the present model can serve as the candidate for the cold dark matter.

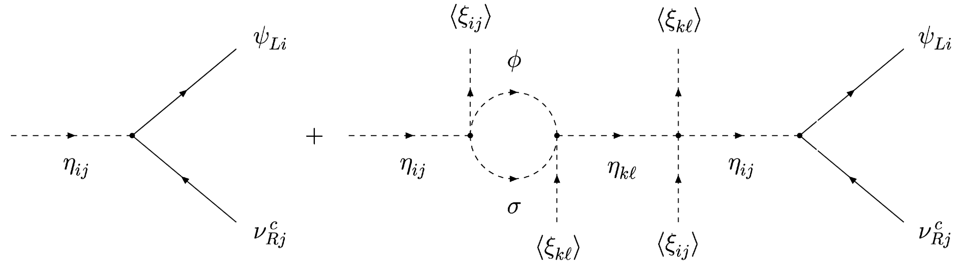

Fig. 3. The new Higgs doublets decay into the leptons at one-loop order. Here $i \neq j$, $k \neq \ell$ and $ij \neq k\ell$.

Note that for the above parameters, the right-handed neutrinos will also decouple at T_F. So the ratio of the relic density of the right-handed neutrinos over that of the left-handed neutrinos is about [17]

$$\frac{n_{\nu_R}}{n_{\nu_L}} \simeq \frac{g_{*S}(\text{MeV})}{g_{*S}(10\ \text{GeV})} = \mathcal{O}(0.1), \tag{27}$$

which is consistent with the current cosmological observation.

It is convenient to extend the present model with certain global symmetry, after which is spontaneously broken near the Planck scale, the pNGBs [18–20] are expected to arise and then explain the quintessence dark energy [21]. For example, we replace the Lagrangian (1) and (2) by

$$\begin{aligned}\mathcal{L} \supset & -\mu_0 \sigma \eta_0^\dagger \phi - \sum_{i\neq j} h_{ij}\xi_{ij}\sigma\eta_{ij}^\dagger\phi \\ & - y_{ii}\overline{\psi_{Li}}\eta_0\nu_{Ri} - \sum_{i\neq j} y_{ij}\overline{\psi_{Li}}\eta_{ij}\nu_{Rj} \\ & - \sum_{i\neq j,\, k\neq\ell} z_{ij,k\ell}\xi_{k\ell}^\dagger\xi_{ij}\eta_{ij}^\dagger\eta_{k\ell} + \text{h.c.},\end{aligned} \tag{28}$$

which is supposed to be invariant under a global $U(1)^3$ symmetry, generated by the independent phase transformations of three Higgs singlets, $\xi_{ij} \equiv \xi_{ji}^*$ $(i \neq j)$, in the limit of vanishing y_{ij} $(i \neq j)$. In other words, the $U(1)^3$ is broken down to a $U(1)^2$ due to the presence of y_{ij} $(i \neq j)$. Thus after the global symmetry is broken by the *vev*s, $\langle\xi_{ij}\rangle \simeq M$, there will be two massless Nambu–Goldstone bosons (NGBs) and one pNGB, which is associated with the neutrino mass-generation. Similar to [20], a typical term in the Coleman–Weinberg effective potential of this pNGB has the form,

$$V(Q) \simeq V_0 \cos(Q/M), \tag{29}$$

with $V_0 = \mathcal{O}(m_\nu^4)$. It is well known that if M is near the Planck scale M_{Pl}, Q will obtain a mass of the order of $\mathcal{O}(m_\nu^2/M_{\text{Pl}})$ and can be a consistent candidate for the quintessence dark energy. Therefore, the intriguing coincidence between the neutrino mass scale $\sim 10^{-2}$ eV and the dark energy scale $\sim 10^{-3}$ eV can be naturally understood. The leading phenomenology of mass varying neutrinos [12,22] is very interesting and can be tested in the present and upcoming experiments [23].

In the model described by the Lagrangian (28), the CP-violation and out-of-equilibrium decays of the new Higgs doublets, as shown in Fig. 3, can produce a lepton asymmetry stored in the left-handed leptons and an equal but opposite lepton asymmetry stored in the right-handed neutrinos. The left-handed lepton asymmetry will be partially converted to the baryon asymmetry through the sphaleron processes [24] and then explain the matter–antimatter asymmetry of the universe. This new type of leptogenesis [13] with the conserved lepton number is called neutrinogenesis [14]. For simplicity, here we will not present the detailed calculation, which is similar to that in a previous work [20].

In this Letter, the mass origin at the TeV scale for the Dirac neutrinos and the dark matter has been successfully unified in a $U(1)_X$ gauge extension of the SM. After the $U(1)_X$ breaking, the Dirac neutrinos can obtain small Yukawa couplings to the SM Higgs and then realize the tiny masses, while the inert fermions can acquire the masses of a few hundred GeV. The inert fermions can annihilate to realize the right amount of the relic density for the cold dark matter. Furthermore, the SM Higgs boson could no longer be a mass eigenstate, and its signatures at LHC could be interesting to modify. Finally, in the presence of the proper global symmetry, after which is spontaneously broken near the Planck scale, the pNGB associated with the neutrino mass-generation can provide the consistent candidate for the dark energy, meanwhile, the matter–antimatter asymmetry [25] of the universe can be generated via the out-of-equilibrium decays of the heavy Higgs doublets with the CP-violation.

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