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Leptogenesis with testable Dirac neutrino mass generation

Pei-Hong Gu [a,b,*]

[a] School of Physics, Jiulonghu Campus, Southeast University, Nanjing 211189, China
[b] School of Physics and Astronomy, Shanghai Jiao Tong University, 800 Dongchuan Road, Shanghai 200240, China

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ABSTRACT

A TeV-scale Higgs doublet can acquire a tiny vacuum expectation value via its small mixing with the standard model Higgs doublet. This new Higgs doublet then can offer a testable Dirac neutrino mass generation through its sizable Yukawa couplings with several right-handed neutrinos and the standard model lepton doublets. We show the small mixing between the two Higgs doublets can be naturally induced by a seesaw mechanism after an additional symmetry is spontaneously broken. In association with the second Higgs doublet decays, this seesaw mechanism can also accommodate a leptogenesis mechanism to generate the baryon asymmetry in the universe.

1. Introduction

The phenomena of neutrino oscillations have been established by the atmospheric, solar, accelerator and reactor neutrino experiments [1]. This means three flavors of neutrinos should be massive and mixed. Meanwhile, the cosmological observations have indicated that the neutrinos should be extremely light [1]. The tiny neutrino masses can be naturally induced at tree level by the so-called seesaw [2–4] extensions of the standard model (SM). Alternatively, the neutrino masses can be achieved in some radiative seesaw models [5]. In these popular seesaw scenarios, the neutrino masses originate from some lepton-number-violating interactions and hence the neutrinos have a Majorana nature. However, we should keep in mind that the theoretical assumption of the lepton number violation and then the Majorana neutrinos has not been confirmed by any experiments. So it is worth studying the possibility of Dirac neutrinos [6–31]. In analogy to the conventional seesaw mechanisms for the Majorana neutrinos, we can construct the tree-level Dirac seesaw [6,8,12,13] as well as the radiative Dirac seesaw [14–16] for the Dirac neutrinos. In the Majorana or Dirac seesaw models, the cosmic baryon asymmetry, which is another big challenge to the SM, can be understood in a natural way. This is the so-called leptogenesis mechanism [32] and has been widely studied [5,7–9,12–14,18,19,32–41].

In this paper we shall present a double Dirac seesaw scenario to simultaneously generate the tiny neutrino masses and the cosmic baryon asymmetry. Besides the SM gauge symmetries, our model respects an additionally global or gauge symmetry, under which the right-handed neutrinos have no Yukawa couplings with the SM. After spontaneously breaking this additional symmetry, two or more heavy Higgs singlets can acquire their suppressed vacuum expectation values (VEVs) and then a second Higgs doublet, which realizes the Yukawa couplings of the right-handed neutrinos to the SM lepton doublets, can obtain a small mixing with the SM Higgs doublet. Therefore, the second Higgs doublet can also acquire a suppressed VEV even if it is set at the TeV scale. This means a testable Dirac neutrino mass generation [10,17]. Furthermore, the heavy Higgs singlet decays can produce an asymmetry stored in the second Higgs doublet. This asymmetry can result in a desired baryon asymmetry in association with the sphaleron processes [42]. So, unlike the previous works, our model can simultaneously realize a testable neutrino mass generation and accommodate a successful leptogenesis.

2. The models

We denote the SM fermions and Higgs by

$$q_L(3,2,+\tfrac{1}{6}) = \begin{bmatrix} u_L \\ d_L \end{bmatrix}, \quad d_R(3,1,-\tfrac{1}{3}), \quad u_R(3,1,+\tfrac{2}{3}),$$

$$l_L(1,2,-\tfrac{1}{2}) = \begin{bmatrix} \nu_L \\ e_L \end{bmatrix}, \quad e_R(1,1,-1),$$

$$\phi(1,2,-\tfrac{1}{2}) = \begin{bmatrix} \phi^0 \\ \phi^- \end{bmatrix}. \quad (1)$$

* Correspondence to: School of Physics, Jiulonghu Campus, Southeast University, Nanjing 211189, China.
E-mail address: phgu@seu.edu.cn.

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Here and thereafter the brackets following the fields describe the transformations under the $SU(3)_c \times SU(2)_L \times U(1)_Y$ gauge groups. The SM charged fermions can obtain their masses through the Yukawa interactions as follows,

$$\mathcal{L} \supset -y_d \bar{q}_L \tilde{\phi} d_R - y_u \bar{q}_L \phi u_R - y_e \bar{l}_L \tilde{\phi} e_R + \text{H.c.}\,. \tag{2}$$

Similarly, we can introduce two or more right-handed neutrinos,

$$\nu_R(1,1,0)\,, \tag{3}$$

to construct the Yukawa interactions for generating a neutrino mass matrix with at least two nonzero eigenvalues,

$$\mathcal{L} \supset -y_\nu \bar{l}_L \phi \nu_R + \text{H.c.}\,. \tag{4}$$

In this Dirac neutrino scenario, the tiny neutrino masses would enforce the Yukawa couplings y_ν to be extremely small. This smallness could be naturally explained by certain Dirac seesaw mechanisms. In the Dirac seesaw models, the Yukawa interactions (4) will not appear before a new symmetry is spontaneously broken.

The present work will be based on the models as below,

$$\begin{aligned}\mathcal{L} \supset & -\mu_\xi^2 \xi^\dagger \xi - \lambda_\xi \left(\xi^\dagger \xi\right)^2 - \mu_\phi^2 \phi^\dagger \phi - \lambda_\phi \left(\phi^\dagger \phi\right)^2 \\ & -\lambda_{\xi\phi} \xi^\dagger \xi \phi^\dagger \phi - M_a^2 \sigma_a^\dagger \sigma_a - (\mu_\eta^2 + \lambda_{\xi\eta} \xi^\dagger \xi + \lambda_{\phi\eta} \phi^\dagger \phi \\ & + \lambda_{a\eta} \sigma_a^\dagger \sigma_a) \eta^\dagger \eta - \lambda'_{\phi\eta} \eta^\dagger \phi \phi^\dagger \eta - \kappa_a \sigma_a \xi^\dagger \xi^\dagger \xi^\dagger \\ & - \rho_a \sigma_a \eta^\dagger \phi - f \bar{l}_L \eta \nu_R + \text{H.c.}\,, \quad (a = 1, \ldots, n \geq 2)\,,\end{aligned} \tag{5}$$

where σ and ξ are the SM singlets while η is a new Higgs doublet,

$$\sigma(1,1,0)\,, \quad \xi(1,1,0)\,, \quad \eta(1,2,-\tfrac{1}{2}) = \begin{bmatrix} \eta^0 \\ \eta^- \end{bmatrix}. \tag{6}$$

Note that one of the four parameters $\kappa_{a,b\neq a}$ and $\rho_{a,b\neq a}$ can always keep to be complex.

In order to forbid the Yukawa interactions (4) and then construct the model (5), we can introduce a $U(1)_{B-L}$ gauge symmetry under which three right-handed neutrinos $\nu_{R1,2,3}$ carry the $B-L$ numbers [43],

$$(-4, -4, +5) \quad \text{for} \quad (\nu_{R1}, \nu_{R2}, \nu_{R3})\,, \tag{7}$$

while the new scalars η, ξ and σ carry the $B-L$ numbers,

$$(+3, +1, +3) \quad \text{for} \quad (\eta, \xi, \sigma)\,. \tag{8}$$

In this $U(1)_{B-L}$ scenario, although the third right-handed neutrino ν_{R3} has no Yukawa couplings and hence keeps massless, it can decouple at a temperature above the QCD scale and hence can escape from the BBN constraint when the $U(1)_{B-L}$ symmetry is broken above the TeV scale. Alternatively, we can consider a $U(1)_X$ global symmetry under which only the non-SM fields are nontrivial, i.e.

$$(+3, -3, -1, -3) \quad \text{for} \quad (\nu_R, \eta, \xi, \sigma)\,. \tag{9}$$

Remarkably, this $U(1)_X$ global symmetry can be always allowed even if we have introduced the previous $U(1)_{B-L}$ gauge symmetry.

It is easy to see that in Lagrangian (5), we can introduce an arbitrary definition of the global lepton numbers of the right-handed neutrinos ν_R and the new scalars (η, ξ, σ) to conserve the global lepton number, i.e.

$$\nu_R(x)\,, \quad \eta(1-x)\,, \quad \xi(\tfrac{1-x}{3})\,, \quad \sigma(1-x)\,. \tag{10}$$

We will show later our leptogenesis mechanism does not depend on the parameter x.

3. Dirac neutrino mass

The $U(1)_{B-L}$ gauge symmetry or the $U(1)_X$ global symmetry is expected to be spontaneously broken before the electroweak symmetry breaking. The Higgs singlet ξ and the Higgs doublet ϕ are responsible for spontaneously breaking the $U(1)_{B-L}$ or $U(1)_X$ symmetry and the electroweak symmetry respectively. When the $U(1)_{B-L}$ or $U(1)_X$ symmetry is spontaneously broken while the electroweak symmetry is still conserved, we can minimalize the scalar potential, i.e.

$$\begin{aligned}\frac{\partial V}{\partial \langle\xi\rangle} &= \langle\xi\rangle(2\mu_\xi^2 + 4\lambda_\xi \langle\xi\rangle^2 + 3\kappa_a \langle\sigma_a\rangle\langle\xi\rangle + 3\kappa_a^* \langle\sigma_a^*\rangle\langle\xi\rangle) \\ &= 0\,,\end{aligned} \tag{11}$$

$$\frac{\partial V}{\partial \langle\sigma_a^*\rangle} = M_a^2 \langle\sigma_a\rangle + \kappa_a^* \langle\xi\rangle^3 = 0\,. \tag{12}$$

In the limiting case of $\langle\sigma_a\rangle, \langle\sigma_a^*\rangle \ll \langle\xi\rangle$, we can obtain

$$\langle\xi\rangle \simeq \sqrt{-\frac{\mu_\xi^2}{2\lambda_\xi}} \quad \text{for} \;\; \mu_\xi^2 < 0\,, \; \lambda_\xi > 0\,, \tag{13}$$

as well as

$$\langle\sigma_a\rangle \simeq -\frac{\kappa_a^* \langle\xi\rangle^3}{M_a^2} \ll \langle\xi\rangle \;\; \text{for} \;\; M_a^2 \gg \langle\xi\rangle^2\,. \tag{14}$$

As a result, the second Higgs doublet η can have a small mixing with the SM Higgs doublet ϕ, i.e.

$$\mathcal{L} \supset -\mu_{\eta\phi}^2 \eta^\dagger \phi + \text{H.c. with } \mu_{\eta\phi}^2 = \rho_a \langle\sigma_a\rangle\,. \tag{15}$$

This means the second Higgs doublet η can also pick up a suppressed VEV after the SM Higgs doublet ϕ drives the electroweak symmetry breaking, i.e.

$$\langle\eta\rangle \simeq -\frac{\rho_a \langle\sigma_a\rangle\langle\phi\rangle}{m_{\eta^0}^2} \ll \langle\phi\rangle\,. \tag{16}$$

Here m_{η^0} is the mass of the neutral component η^0 of the second Higgs doublet η, i.e.

$$m_{\eta^0}^2 = \mu_\eta^2 + \lambda_{\xi\eta}\langle\xi\rangle^2 + (\lambda_{\phi\eta} + \lambda'_{\phi\eta})\langle\phi\rangle^2 + \lambda_{a\eta}\langle\sigma_a\rangle^2\,, \tag{17}$$

which could have a split with the mass $m_{\eta^\pm}$ of the charged component $\eta^\pm$, i.e.

$$m_{\eta^\pm}^2 = \mu_\eta^2 + \lambda_{\xi\eta}\langle\xi\rangle^2 + \lambda_{\phi\eta}\langle\phi\rangle^2 + \lambda_{a\eta}\langle\sigma_a\rangle^2\,. \tag{18}$$

Through the Yukawa interactions of the second Higgs doublet η to the right-handed neutrinos ν_R and the SM lepton doublets l_L, we now can obtain a tiny neutrino mass term,

$$\mathcal{L} \supset -m_{\alpha i} \bar{\nu}_{L\alpha} \nu_{Ri} + \text{H.c. with } m_{\alpha i} = f_{\alpha i}\langle\eta\rangle\,. \tag{19}$$

This Dirac neutrino mass generation depends on two-step suppression of the VEVs $\langle\sigma\rangle$ and $\langle\eta\rangle$ so that it may be titled as a double type-II Dirac seesaw mechanism, in analogy to our double type-II seesaw [44] for the Majorana neutrinos. We can conveniently understand this double Dirac seesaw in Fig. 1.

Note the heavy Higgs singlets σ_a in the above Dirac neutrino mass generation will give a huge radiative correction to the quadratic terms of the Higgs doublets ϕ and η. The one-loop contributions are,

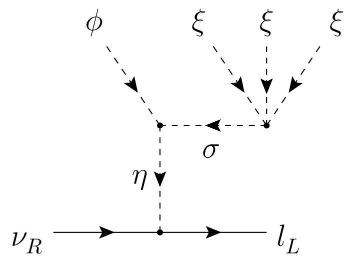

Fig. 1. The Dirac neutrino mass generation.

$$\delta\mu_\eta^2 = -\frac{1}{16\pi^2}\left[2\mathrm{Tr}\left(f^\dagger f\right)\Lambda^2 + |\rho_a|^2 \ln\left(\frac{\Lambda^2}{M_a^2}\right)\right],$$
$$\delta\mu_\phi^2 = -\frac{1}{16\pi^2}|\rho_a|^2 \ln\left(\frac{\Lambda^2}{M_a^2}\right), \quad (20)$$

with Λ being the regulator cutoff. This means a large fine-tuning would be necessary to obtain the SM Higgs boson and set the second Higgs doublet at the TeV scale. Actually, such fine-tunings also exist in the conventional seesaw models for the Majorana neutrino mass generation [45].

4. Heavy Higgs singlet decays

As shown in Fig. 2, the heavy Higgs singlets σ_a have two decay modes,

$$\sigma_a \to \xi\xi\xi\,, \quad \sigma_a \to \eta\phi^*\,. \quad (21)$$

As long as the CP is not conserved, we can expect a CP asymmetry in the above decays,

$$\varepsilon_a = \frac{\Gamma(\sigma_a \to \eta\phi^*) - \Gamma(\sigma_a^* \to \eta^*\phi)}{\Gamma_a}$$
$$= -\frac{\Gamma(\sigma_a \to \xi\xi\xi) - \Gamma(\sigma_a^* \to \xi^*\xi^*\xi^*)}{\Gamma_a} \neq 0\,, \quad (22)$$

where Γ_a is the total decay width,

$$\Gamma_a = \Gamma(\sigma_a \to \xi\xi\xi) + \Gamma(\sigma_a \to \eta\phi^*)$$
$$= \Gamma(\sigma_a^* \to \xi^*\xi^*\xi^*) + \Gamma(\sigma_a^* \to \eta^*\phi)\,. \quad (23)$$

We can calculate the decay width at tree level and the CP asymmetry at one-loop level,

$$\Gamma_a = \frac{1}{8\pi}\left(\frac{|\rho_a|^2}{M_a^2} + \frac{3|\kappa_a|^2}{32\pi^2}\right) M_a\,, \quad (24)$$

$$\varepsilon_a = -\frac{3}{64\pi^3}\sum_{b\neq a} \frac{\mathrm{Im}\left(\kappa_a^*\kappa_b\rho_a\rho_b^*\right)}{\frac{|\rho_a|^2}{M_a^2} + \frac{3|\kappa_a|^2}{32\pi^2}} \frac{1}{M_b^2 - M_a^2}$$
$$= -\frac{3}{64\pi^3}\sum_{b\neq a} \frac{|\kappa_a\kappa_b\rho_a\rho_b|\sin\delta_{ab}}{\frac{|\rho_a|^2}{M_a^2} + \frac{3|\kappa_a|^2}{32\pi^2}} \frac{1}{M_b^2 - M_a^2}\,. \quad (25)$$

Here δ_{ab} is the relative phase among the parameters $\rho_{a,b}$ and $\kappa_{a,b}$.

After the heavy Higgs singlets σ_a go out of equilibrium, their decays can generate an X asymmetry X_η and a lepton asymmetry L_η stored in the second Higgs doublet η. Meanwhile, these decays can produce an X asymmetry X_ξ and a lepton asymmetry L_ξ stored in the Higgs singlet ξ. For demonstration, we simply assume a hierarchical spectrum of the heavy Higgs singlets σ_a, i.e. $M_{\sigma_1}^2 \ll M_{\sigma_{2,...}}^2$. In this case, the decays of the lightest σ_1 should dominate the final asymmetries in the second Higgs doublet η, i.e.

$$X_\eta = -3\frac{\varepsilon_1}{g_*}\kappa\,, \quad L_\eta = (1-x)\frac{\varepsilon_1}{g_*}\kappa\,, \quad (26)$$

where the factors -3 and $1-x$ respectively are the X number and the lepton number of the second Higgs doublet η, the character $\kappa < 1$ is a washout factor and the symbol $g_* = 112.75 + 1.75n$ is the relativistic degrees of freedom (the SM fields plus the right-handed neutrinos ν_R, the second Higgs doublet η and the Higgs singlet ξ) [46]. Similarly, the final asymmetries in the Higgs singlet ξ are

$$X_\xi = 3\times(-1)\times\frac{-\varepsilon_1}{g_*}\kappa = 3\frac{\varepsilon_1}{g_*}\kappa = -X_\eta\,,$$
$$L_\xi = 3\times\left(\frac{1-x}{3}\right)\frac{-\varepsilon_1}{g_*}\kappa = -(1-x)\frac{\varepsilon_1}{g_*}\kappa = -L_\eta\,. \quad (27)$$

Here we have taken Eqs. (9), (10) and (22) into account. Note the asymmetries in the Higgs singlet ξ and the asymmetries in the second Higgs doublet η will decouple from each other once they are produced from the out-of-equilibrium decays of the heavy Higgs singlets σ_1. Furthermore, the $U(1)_{B-L}$ or $U(1)_X$ symmetry has forbidden the terms only including one Higgs singlet ξ. This means the Higgs singlet ξ can not have any decay channels before the $U(1)_{B-L}$ or $U(1)_X$ symmetry breaking. The asymmetries in the Higgs singlet ξ eventually will not exist after the Higgs singlet ξ develops its VEV.

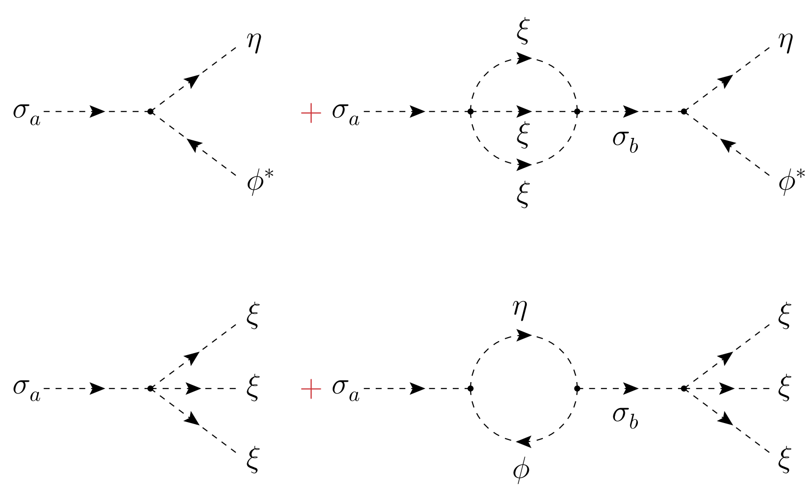

Fig. 2. The heavy Higgs singlet decays.

5. Baryon asymmetry

The X asymmetry X_η and the lepton asymmetry L_η stored in the second Higgs doublet η will lead to a lepton asymmetry stored in the SM lepton doublets l_L because of the related Yukawa interactions in Eq. (5). The sphaleron processes then can partially transfer this SM lepton asymmetry to a baryon asymmetry.

We now analyze the chemical potentials [47] to discuss the details of these conversions. For this purpose, we denote μ_q, μ_d, μ_u, μ_l, μ_e, μ_ν, μ_ϕ and μ_η for the chemical potentials of the fields q_L, d_R, u_R, l_L, e_R, ν_R, ϕ and η. We then can consider the chemical potentials in two phases,

- phase-I: during the inert Higgs singlet decays and the second Higgs doublet decays,
- phase-II: during the second Higgs doublet decays and the electroweak symmetry breaking.

In phase-I, the SM Yukawa interactions are in equilibrium and hence yield,

$$-\mu_q+\mu_d-\mu_\phi=0\,, \tag{28}$$

$$-\mu_q+\mu_u+\mu_\phi=0\,, \tag{29}$$

$$-\mu_l+\mu_e-\mu_\phi=0\,, \tag{30}$$

the fast sphalerons constrain,

$$3\mu_q+\mu_l=0\,, \tag{31}$$

while the neutral hypercharge in the universe requires,

$$3\left(\mu_q-\mu_d+2\mu_u-\mu_l-\mu_e\right)-2\mu_\phi-2\mu_\eta=0\,. \tag{32}$$

In addition, the Yukawa interactions involving the right-handed neutrinos are also in equilibrium. This means

$$-\mu_l+\mu_\nu+\mu_\eta=0\,. \tag{33}$$

In the above Eqs. (28)-(33), we have identified the chemical potentials of the different-generation fermions because the Yukawa interactions establish an equilibrium between the different generations. By solving Eqs. (28)-(33), we express the chemical potentials in phase-I as below,

$$\begin{aligned}
&\mu_\phi^I=-\frac{4}{7}\mu_l^I-\frac{1}{7}\mu_\eta^I\,,\quad \mu_q^I=-\frac{1}{3}\mu_l^I\,,\\
&\mu_d^I=-\frac{19}{21}\mu_l^I-\frac{1}{7}\mu_\eta^I\,,\quad \mu_u^I=\frac{5}{21}\mu_l^I+\frac{1}{7}\mu_\eta^I\,,\\
&\mu_e^I=\frac{3}{7}\mu_l^I-\frac{1}{7}\mu_\eta^I\,,\quad \mu_\nu^I=\mu_l^I-\mu_\eta^I\,.
\end{aligned} \tag{34}$$

Now the baryon number can be given by

$$B^I=3(2\mu_q^I+\mu_d^I+\mu_u^I)=-4\mu_l^I\,. \tag{35}$$

As for the lepton number, it should be

$$L^I=L_{SM}^I+L_{\nu_R}^I+L_\eta^I\,, \tag{36}$$

with L_{SM}^I, $L_{\nu_R}^I$ and L_η^I being the lepton number in the SM leptons, the right-handed neutrinos and the second Higgs doublet, respectively,

$$\begin{aligned}
L_{SM}^I&=3(2\mu_l^I+\mu_e^I)=\frac{51}{7}\mu_l^I-\frac{3}{7}\mu_\eta^I\,,\\
L_{\nu_R}^I&=nx\mu_\nu^I=nx(\mu_l^I-\mu_\eta^I)\,,\quad (n\geq 2)\,,\\
L_\eta^I&=4(1-x)\mu_\eta^I\,.
\end{aligned} \tag{37}$$

The $B-L$ number then can be computed by

$$\begin{aligned}
(B-L)^I&=B^I-(L_{SM}^I+L_{\nu_R}^I+L_\eta^I)\\
&=-\frac{79+7nx}{7}\mu_l^I+\frac{7(n+4)x-25}{7}\mu_\eta^I\,.
\end{aligned} \tag{38}$$

According to the $U(1)_X$ global symmetry (9), we can also obtain an X number,

$$\begin{aligned}
X^I&=X_{\nu_R}^I+X_\eta^I=3n\mu_\nu^I-12\mu_\eta^I\\
&=3n\mu_l^I-3(n+4)\mu_\eta^I\,.
\end{aligned} \tag{39}$$

After the heavy Higgs singlet decays, the $B-L$ and X numbers in the SM fields, the right-handed neutrinos and the second Higgs doublet should be both conserved. Therefore, we can read

$$\begin{aligned}
(B-L)^I&=(B-L)^i=-L_\eta^i=-(1-x)\frac{\varepsilon_1}{g_*}\kappa\,,\\
X^I&=X_\eta^i=-3\frac{\varepsilon_1}{g_*}\kappa\,,
\end{aligned} \tag{40}$$

where L_η^i and X_η^i are the initial $B-L$ and X numbers from the decays of the heavy Higgs singlets into the second Higgs doublet. Clearly, we should have

$$L_\eta^i=-\frac{1-x}{3}X_\eta^i\,. \tag{41}$$

So we eventually can derive

$$\begin{aligned}
B^I&=\frac{7n+3}{3(26n+79)}X_\eta^i\,,\\
L_{SM}^I&=-\frac{4n-1}{26n+79}X_\eta^i\,,\\
L_{\nu_R}^I&=\frac{19nx}{3(26n+79)}X_\eta^i\,,\\
L_\eta^I&=-\frac{(7n+79)(1-x)}{3(26n+79)}X_\eta^i\,.
\end{aligned} \tag{42}$$

In phase-II, the second Higgs doublet η has already decayed so that the condition (32) for the zero hypercharge should be modified by,

$$3\left(\mu_q-\mu_d+2\mu_u-\mu_l-\mu_e\right)-2\mu_\phi=0\,. \tag{43}$$

Ones then can solve Eqs. (28)-(31) and (43) to determine the chemical potentials,

$$\begin{aligned}
&\mu_q^{II}=-\frac{1}{3}\mu_l^{II}\,,\quad \mu_d^{II}=-\frac{19}{21}\mu_l^{II}\,,\quad \mu_u^{II}=\frac{5}{21}\mu_l^{II}\,,\\
&\mu_e^{II}=\frac{3}{7}\mu_l^{II}\,,\quad \mu_\phi^{II}=-\frac{4}{7}\mu_l^{II}\,.
\end{aligned} \tag{44}$$

At this stage, the baryon and lepton numbers in the SM should be

$$\begin{aligned}
B^{II}&=\frac{28}{79}(B-L)^{II}\,,\\
L_{SM}^{II}&=-\frac{51}{79}(B-L)^{II}\,.
\end{aligned} \tag{45}$$

Because the second Higgs doublet carries the lepton number $1-x$, its decays can produce the lepton doublets with a number $L_\eta^I/(1-x)$. The conserved $(B-L)^{II}$ number thus should be

$$(B-L)^{II}=B^I-L_{SM}^I-\frac{1}{1-x}L_\eta^I=\frac{1}{3}X_\eta^i\,. \tag{46}$$

This means the final baryon asymmetry can be given by

$$B^f = B^{II} = \frac{28}{237} X^i_\eta = -\frac{28}{79} \frac{\varepsilon_1}{g_*} \kappa\,, \tag{47}$$

which is independent on the choice of the global lepton numbers of the right-handed neutrinos and the other non-SM fields.

6. Numerical example

In order to determine the washout factor κ, we in principle should solve the Boltzmann equations. To give a numerical example, we here shall simply consider the weak washout case where the decay width Γ_1 is required to fulfil the condition as below [46],

$$K = \frac{\Gamma_1}{H(T)}\Big|_{T=M_1} \ll 1 \tag{48}$$

Here and thereafter $H(T)$ is the Hubble constant,

$$H(T) = \left(\frac{8\pi^3 g_*}{90}\right)^{\frac{1}{2}} \frac{T^2}{M_{\rm Pl}}\,, \tag{49}$$

with $M_{\rm Pl} \simeq 1.22 \times 10^{19}$ GeV being the Planck mass. The X_η and L_η asymmetries then can be approximately described by

$$X_\eta \simeq -3\varepsilon_1 \left(\frac{n^{eq}_{\sigma_1}}{s}\right)\Big|_{T=M_1}\,,\; L_\eta \simeq (1-x)\varepsilon_1 \left(\frac{n^{eq}_{\sigma_1}}{s}\right)\Big|_{T=M_1}\,,$$
$$\text{with } n^{eq}_{\sigma_1} = 2\left(\frac{M_1 T}{2\pi}\right)^{\frac{3}{2}} e^{-\frac{M_1}{T}}\,,\quad s = \frac{2\pi^2}{45} g_* T^3\,, \tag{50}$$

where the symbol $n^{eq}_{\sigma_1}$ is the equilibrium number density of the heavy Higgs singlets σ_1, while the character s is the entropy density of the universe [46]. In this weak washout case, the washout factor κ is given by

$$\kappa \simeq \frac{45}{2^{\frac{3}{2}}\pi^{\frac{7}{2}} e} = 0.11\,. \tag{51}$$

If the $U(1)_{B-L}$ gauge symmetry rather than the $U(1)_X$ global symmetry is introduced, we need also consider the annihilating and scattering processes induced by the $U(1)_{B-L}$ gauge interactions. The washout factor (51) indeed is based on an assumption that these annihilations and scatterings have already decoupled before the decays. For this purpose, we can take the heavy Higgs singlet σ_1 heavy enough and hence the processes from the gauge interactions can have a rate Γ_g smaller than the decay width Γ_1, i.e.

$$\Gamma_g \sim \frac{g^4_{B-L}}{M_1^2} n^{eq}_{\sigma_1} < \Gamma_1\,, \tag{52}$$

with g_{B-L} being the $U(1)_{B-L}$ gauge coupling.

As an example, we assume three right-handed neutrinos and then obtain

$$K = 0.052\,,$$
$$\varepsilon_1 = -7.8 \times 10^{-7} \sin\delta_{12}\,,$$
$$B^f \simeq 10^{-10}\left(\frac{\sin\delta_{12}}{0.4}\right)\,, \tag{53}$$

by inputting

$$M_1 = 10^{14}\,\text{GeV}\,,\ |\rho_1| = 10^{12}\,\text{GeV}\,,\ |\kappa_1| = 0.1\,;$$
$$M_2 = 10^{15}\,\text{GeV}\,,\ |\rho_2| = 10^{13}\,\text{GeV}\,,\ |\kappa_2| = 0.1\,. \tag{54}$$

The assumption (52) can be easily satisfied for $g_{B-L} = \mathcal{O}(0.1)$. We further take

$$\langle\xi\rangle = 10\,\text{TeV}\,,\ \ m_{\eta^0} = 1\,\text{TeV}\,, \tag{55}$$

to realize

$$\langle\eta\rangle = 1\,\text{eV}\,, \tag{56}$$

and hence

$$m_\nu = \mathcal{O}(0.01\text{-}0.1\,\text{eV})\ \text{ for }\ f = \mathcal{O}(0.01-0.1)\,. \tag{57}$$

So, the above parameter choice can explain the baryon asymmetry and the neutrino masses.

7. Phenomenological implications

If we introduce the $U(1)_{B-L}$ gauge symmetry, the $U(1)_{B-L}$ gauge boson Z_{B-L} should have a mass,

$$M_{Z_{B-L}} = \sqrt{2} g_{B-L}\langle\xi\rangle\,. \tag{58}$$

The Z_{B-L} contribution to the cross sections for $e^+e^- \to f\bar{f}$ proceeds through an s-channel Z_{B-L} exchange (when $f = e$, there are also t- and u-channels). In the case that $M_{Z_{B-L}}$ is above 209 GeV (the maximum energy of LEP II), the bound on the $U(1)_{B-L}$ symmetry breaking scale should be [48],

$$\frac{M_{Z_{B-L}}}{g_{B-L}} \gtrsim 7\,\text{TeV} \Rightarrow \langle\xi\rangle \gtrsim 5\,\text{TeV}\,. \tag{59}$$

The Z_{B-L} boson will mediate the annihilations of the three right-handed neutrinos $\nu_{R1,2,3}$ into the SM fermions. We need to check if these annihilations can decouple above the QCD scale,

$$\sigma^g_{\nu_R} = \sum_{f=d,u,s,e,\mu,\nu_L} \sigma(\nu_R + \nu^c_R \to f + f^c)$$
$$= \frac{6g^4_{B-L}x^2}{25\pi}\frac{s}{M^4_{Z_{B-L}}} = \frac{3x^2}{50\pi}\frac{s}{\langle\xi\rangle^4}\,, \tag{60}$$

where $x = +4$ or -5 is the $B-L$ number of the right-handed neutrinos $\nu_{R1,2}$ or ν_{R3}, while s is the Mandelstam variable. The interaction rate then should be [39]

$$\Gamma^g_{\nu_R} = \frac{\frac{T}{32\pi^4}\int_0^\infty s^{3/2} K_1\left(\frac{\sqrt{s}}{T}\right)\sigma^g_{\nu_R} ds}{\frac{2}{\pi^2}T^3} = \frac{18x^2}{25\pi^3}\frac{T^5}{\langle\xi\rangle^4}\,, \tag{61}$$

with K_1 being a Bessel function. The decoupling temperature T_{ν_R} of the right-handed neutrinos $\nu_{R1,2,3}$ then can be determined by

$$\left[\Gamma^g_{\nu_R} = H(T)\right]\Big|_{T=T_{\nu_R}}\,. \tag{62}$$

The contribution of the right-handed neutrinos $\nu_{R1,2,3}$ to the effective neutrino number then should be [46]

$$\Delta N_\nu = \sum_{i=1}^{3}\left[\frac{10.75}{g_*(T_{\nu_{Ri}})}\right]^{\frac{4}{3}}\,. \tag{63}$$

We take $g_*(300\,\text{MeV}) \simeq 61.75$ [46] and then find

$$T_{\nu_{R1,2}} \simeq 348\,\text{MeV}\,,\ \ T_{\nu_{R3}} \simeq 300\,\text{MeV}\,,$$
$$\Delta N_\nu \simeq 0.29 \text{ for } \langle\xi\rangle = 11\,\text{TeV}\,. \tag{64}$$

So, the right-handed neutrinos $\nu_{R1,2,3}$ can be harmless for $\langle\xi\rangle \gtrsim 11$ TeV.

Alternatively, we can introduce the $U(1)_X$ global symmetry. In this case, a massless Goldstone boson J will emerge after the

$U(1)_X$ symmetry breaking. For $\langle\xi\rangle \gg \langle\sigma\rangle, \langle\eta\rangle$, this Goldstone J should mostly come from the Higgs singlet ξ, i.e.

$$\xi = \left(\frac{1}{\sqrt{2}}h_\xi + \langle\xi\rangle\right)\exp\left(i\frac{J}{\sqrt{2}\langle\xi\rangle}\right), \tag{65}$$

and couples to the left- and right-handed neutrinos $\nu_{L,R}$ as well as the new Higgs boson h_ξ at tree level,

$$\mathcal{L} \supset -i\frac{3}{\sqrt{2}}\frac{m_\nu}{\langle\xi\rangle}\left(J\bar{\nu}_L\nu_R + \frac{i}{\sqrt{2}\langle\xi\rangle}J^2\bar{\nu}_L\nu_R - \frac{1}{2\langle\xi\rangle^2}J^3\bar{\nu}_L\nu_R\right) + \text{H.c.} + \frac{1}{\langle\xi\rangle^2}\left(\frac{1}{4}h_\xi^2 + \frac{1}{\sqrt{2}}\langle\xi\rangle h_\xi\right)\partial_\mu J\partial^\mu J\,. \tag{66}$$

Clearly, the Goldstone J is similar to the usual majoron [49]. The Goldstone-neutrino coupling in Eq. (66) should be extremely small, i.e. $m_\nu/\langle\xi\rangle \lesssim \mathcal{O}(0.1\ \text{eV}/1-100\ \text{TeV}) = \mathcal{O}(10^{-12}-10^{-14})$ and hence can escape from the experimental constraints [49–51]. Meanwhile, the Goldstone can decouple at a temperature T_J above the QCD scale and hence it can only give a negligible contribution to the effective neutrino number [46],

$$\Delta N_\nu = \frac{8}{7}\left[\frac{10.75}{g_*\left(T_J\right)}\right]^{\frac{4}{3}}. \tag{67}$$

Actually, the rate of the annihilation of the SM Higgs doublet into the Goldstone can be roughly estimated by $\lambda_{\xi\phi}^2 T^3\langle\xi\rangle^2/(T^2+m_{h_\xi}^2)^2$ [46]. It is easy to check for $\langle\xi\rangle = \mathcal{O}(1-100\ \text{TeV})$ and small enough $\lambda_{\xi\phi}$, this annihilation even can decouple before the electroweak symmetry breaking. In this case, we can input $g_*\left(T_J\right) = 106.75$ into Eq. (67) and then obtain $\Delta N_\nu = 0.05$.

Through the mediation of the second Higgs doublet η, the right-handed neutrinos ν_R can annihilate into the left-handed neutrinos as well as the charged leptons. These annihilations can affect the contributions of the right-handed neutrinos to the effective neutrino number. To satisfy the BBN constraint, we can require the right-handed neutrinos to decouple above the QCD scale. We calculate the annihilating cross section,

$$\begin{aligned}\sigma_{\nu_{Ri}}^\eta &= \sum_{f=e,\mu,\nu_L}\sigma(\nu_{Ri}+\nu_{Ri}^c \to f + f^c)\\ &= \frac{[(f^\dagger f)_{ii}]^2}{192\pi}\frac{s}{m_{\eta^0}^4} + \frac{|f_{ei}|^2+|f_{\mu i}|^2}{192\pi}\frac{s}{m_{\eta^\pm}^4}\\ &= \frac{s}{192\pi}\left[\frac{(\sum_{j=1}^3 m_j^2)^2}{m_{\eta^0}^4\langle\eta\rangle^4} + \frac{(|m_{ei}|^2+|m_{\mu i}|^2)^2}{m_{\eta^\pm}^4\langle\eta\rangle^4}\right]\\ &< \frac{s}{96\pi}\frac{(\sum_{j=1}^3 m_j^2)^2}{m_{\eta^0}^4\langle\eta\rangle^4}\end{aligned} \tag{68}$$

and then the interaction rate,

$$\begin{aligned}\Gamma_{\nu_{Ri}}^\eta &= \frac{\frac{T}{32\pi^4}\int_0^\infty s^{3/2}K_1\left(\frac{\sqrt{s}}{T}\right)\sigma_{\nu_{Ri}}^\eta ds}{\frac{2}{\pi^2}T^3}\\ &= \frac{T^5}{16\pi^3}\left[\frac{(\sum_{j=1}^3 m_j^2)^2}{m_{\eta^0}^4\langle\eta\rangle^4} + \frac{(|m_{ei}|^2+|m_{\mu i}|^2)^2}{m_{\eta^\pm}^4\langle\eta\rangle^4}\right]\\ &< \frac{T^5}{8\pi^3}\frac{(\sum_{j=1}^3 m_j^2)^2}{m_{\eta^0}^4\langle\eta\rangle^4}\,.\end{aligned} \tag{69}$$

Comparing the above interaction rate to the Hubble constant, we find the annihilations induced by the second Higgs doublet η can easily decouple above the QCD scale for $m_{\eta^0,\eta^\pm} = \mathcal{O}$ (TeV), $\langle\eta\rangle = \mathcal{O}$ (eV) and $m_\nu = \mathcal{O}(0.01-0.1\ \text{eV})$.

The second Higgs doublet η with tiny VEV and sizable Yukawa couplings can lead to rich observable phenomena. For example, we can get some useful information on the neutrino mass matrix by the branching ratios in the decays of the charged component $\eta^\pm$ into the SM charged leptons $l^\pm$ with the right-handed neutrinos ν_R [17,52]. The charged component $\eta^\pm$ will also have interesting implications on the muon anomalous magnetic moment a_μ, the rare process $\mu \to e\gamma$, and other precision measurements [17].

In our model, the Dirac neutrinos does not allow a neutrinoless double beta decay [1]. This is a very distinctive feature compared to the traditional Majorana seesaw models, where the origin of the cosmic baryon asymmetry is tightly connected to the lepton number violation for the Majorana neutrino mass generation. So far we have not seen the neutrinoless double beta decay in any experiments [1]. Even if no neutrinoless double beta decay is observed in the future experiments, our model can provide a common origin of the neutrino mass and the baryon asymmetry.

8. Conclusion

In this paper we have demonstrated a double Dirac seesaw mechanism. In our scenario, the heavy Higgs singlets can suppress the mixing between the second Higgs doublet and the SM Higgs doublet after the additionally gauge or global symmetry is spontaneously broken. Therefore, the second Higgs doublet can naturally pick up a tiny VEV even if it is set at the TeV scale. Through the sizable Yukawa couplings of the SM lepton doublets to the right-handed neutrinos and the second Higgs doublet, we can realize a testable Dirac neutrino mass generation. Furthermore, the interactions for the small mixing between the two Higgs doublets can also explain the observed baryon asymmetry in association with the sphaleron processes.

Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

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