Realistic Neutrinogenesis with Radiative Vertex Correction

Pei-Hong Gu[1],* Hong-Jian He[2],† and Utpal Sarkar[3‡]
[1] *The Abdus Salam International Centre for Theoretical Physics, Strada Costiera 11, 34014 Trieste, Italy*
[2] *Center for High Energy Physics, Tsinghua University, Beijing 100084, China*
[3] *Physical Research Laboratory, Ahmedabad 380009, India*

We propose a new model for naturally realizing light Dirac neutrinos and explaining the baryon asymmetry of the universe through neutrinogenesis. To achieve these, we present a minimal construction which extends the standard model with a real singlet scalar, a heavy singlet Dirac fermion and a heavy doublet scalar besides three right-handed neutrinos, respecting lepton number conservation and a Z_2 symmetry. The neutrinos acquire small Dirac masses due to the suppression of weak scale over a heavy mass scale. As a key feature of our construction, once the heavy Dirac fermion and doublet scalar go out of equilibrium, their decays induce the CP asymmetry from the interference of tree-level processes with the *radiative vertex corrections* (rather than the self-energy corrections). Although there is no lepton number violation, an equal and opposite amount of CP asymmetry is generated in the left-handed and the right-handed neutrinos. The left-handed lepton asymmetry would then be converted to the baryon asymmetry in the presence of the sphalerons, while the right-handed lepton asymmetry remains unaffected.

PACS numbers: 98.80.-k, 11.30.Er, 14.60.Pq
arXiv: 0709.1019 [hep-ph]

Strong evidences from neutrino oscillation experiments [1] so far have pointed to tiny but nonzero masses for active neutrinos. The smallness of the neutrino masses can be elegantly understood via seesaw mechanism [2] in various extensions of the standard model (SM). The origin of the observed baryon asymmetry [1] in the universe poses a real challenge to the SM, but within the seesaw scenario, it can be naturally explained through leptogenesis [3, 4, 5, 6, 7, 8].

In the conventional leptogenesis scenario, the lepton number violation is essential as it is always associated with the mass-generation of Majorana neutrinos. However, the Majorana or Dirac nature of the neutrinos is unknown a priori and is awaiting for the upcoming experimental determination. It is important to note [9, 10] that even with lepton number conservation, it is possible to generate the observed baryon asymmetry in the universe. Since the sphaleron processes [11] have no direct effect on the right-handed fields, a nonzero lepton asymmetry stored in the left-handed fields, which is equal but opposite to that stored in the right-handed fields, can be partially converted to the baryon asymmetry as long as the interactions between the left-handed lepton number and the right-handed lepton number are too weak to realize an equilibrium before the electroweak phase transition, the sphalerons convert the lepton asymmetry in the left-handed fields, leaving the asymmetry in the right-handed fields unaffected [10, 12, 13, 14, 15].

For all the SM species, the Yukawa interactions are sufficiently strong to rapidly cancel the stored left- and right-handed lepton asymmetry. However, the effective Yukawa interactions of the ultralight Dirac neutrinos are exceedingly weak [16, 17] and thus will not reach equilibrium until the temperatures fall well below the weak scale. In some realistic models [12, 14, 15], the effective Yukawa couplings of the Dirac neutrinos are naturally suppressed by the ratio of the weak scale over the heavy mass scale. Simultaneously, the heavy particles can decay with the CP asymmetry to generate the expected left-handed lepton asymmetry after they are out of equilibrium. This new type of leptogenesis mechanism is called neutrinogenesis [10].

In this paper, we propose a new model to generate the small Dirac neutrino masses and explain the origin of cosmological baryon asymmetry, by extending the SM with a real scalar, a heavy Dirac fermion singlet and a heavy doublet scalar besides three right-handed neutrinos. In comparison with all previous realistic neutrinogenesis models [12, 14, 15], the Dirac neutrino masses in our new model are also suppressed by the ratio of the weak scale over the heavy mass scale, but the crucial difference is that in the decays of the heavy particles, the *radiative vertex corrections* (instead of the self-energy corrections) interfere with the tree-level diagrams to generate the required CP asymmetry and naturally realize neutrinogenesis.

We summarize the field content in Table I, in which ψ_L, ϕ, ν_R, $D_{L,R}$, η and χ denote the left-handed lepton doublets, the SM Higgs doublet, the right-handed neutrinos, the heavy singlet Dirac fermion, the heavy doublet scalar and the real scalar, respectively. Here ψ_L, ν_R, D_L and D_R carry lepton number 1 while ϕ, η and χ have zero lepton number. For simplicity, we have omitted the family indices as well as other SM fields, which carry even parity under the discrete symmetry Z_2. It should be noted that the conventional dimension-4 Yukawa inter-

*Electronic address: pgu@ictp.it
†Electronic address: hjhe@tsinghua.edu.cn
‡Electronic address: utpal@prl.res.in

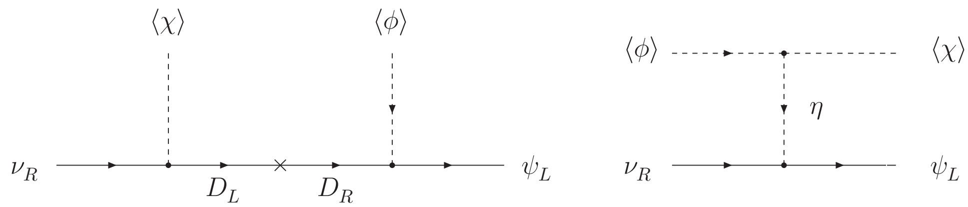

FIG. 1: The neutrino mass-generation. The left diagram is the type-I Dirac seesaw while the right one is the type-II Dirac seesaw.

Fields	$SU(2)_L$	$U(1)_Y$	Z_2
ψ_L	**2**	$-1/2$	$+$
ϕ	**2**	$-1/2$	$+$
ν_R	**1**	0	$-$
$D_{L,R}$	**1**	0	$+$
η	**2**	$-1/2$	$-$
χ	**1**	0	$-$

TABLE I: The field content of our model, where ψ_L is the left-handed lepton doublet, ϕ is the SM Higgs doublet, ν_R is the right-handed neutrino, $D_{L,R}$ is the heavy singlet Dirac fermion, η is the heavy doublet scalar, and χ is the real scalar. For simplicity the family indices and the other SM fields (carrying even Z_2 parity) are omitted from the Table.

actions among the left-handed lepton doublets, the SM Higgs doublet and the right-handed neutrinos are forbidden under the Z_2 symmetry. Our model also exactly conserves the lepton number, so we can write down the relevant Lagrangian as below,

$$\begin{aligned} -\mathcal{L} \supset \; & \{f_i\overline{\psi_{Li}}\phi D_R + g_i\chi\overline{D_L}\nu_{Ri} + y_{ij}\overline{\psi_{Li}}\eta\nu_{Ri} - \mu\chi\eta^\dagger\phi \\ & +M_D\overline{D_L}D_R + \text{h.c}\} + M_\eta^2\eta^\dagger\eta\,, \end{aligned} \tag{1}$$

where f_i, g_i and y_{ij} are the Yukawa couplings, while the cubic scalar coupling μ has mass-dimension equal one. The parameters M_D and M_η in (1) are the masses of the heavy singlet fermion D and the heavy Higgs doublet η, respectively. Note that in the Higgs potential the scalar doublet η has a positive mass-term as shown in the above Eq. (1), while the Higgs doublet ϕ and singlet χ both have negative mass-terms [1].

The lepton number conservation ensures that there is no Majorana mass term for all fermions. As we will discuss below, the vacuum expectation value (vev) of η comes out to be much less than the vev of the other fields. Thus the first two terms generate mixings of the light Dirac neutrinos with the heavy Dirac fermion, while the third term gives the light Dirac neutrino mass term. The complete mass matrix can now be written in the basis $\{\nu_L,\ D_L,\ \nu_R,\ D_R\}$ as

$$M = \begin{bmatrix} 0 & 0 & a & b \\ 0 & 0 & c & d \\ a^\dagger & c^\dagger & 0 & 0 \\ b^\dagger & d^\dagger & 0 & 0 \end{bmatrix}, \tag{2}$$

where $a \equiv y\langle\eta\rangle$, $b \equiv f\langle\phi\rangle$, $c \equiv g\langle\chi\rangle$ and $d \equiv M_D$. As will be shown below, $d \gg a, b, c$. So, the diagonalization of the mass matrix (2) generates the light Dirac neutrino masses of order $a - bc/d$ and a heavy Dirac fermion mass of order d.

As shown in Fig. 1, at low energy we can integrate out the heavy singlet fermion as well as the heavy doublet scalar. Then we obtain the following effective dimension-5 operators,

$$\mathcal{O}_5 = \frac{f_i g_j}{M_D}\overline{\psi_{Li}}\phi\nu_{Rj}\chi - \frac{\mu y_{ij}}{M_\eta^2}\overline{\psi_{Li}}\phi\nu_{Rj}\chi + \text{h.c.}\,. \tag{3}$$

Therefore, once the SM Higgs doublet ϕ and the real scalar χ both acquire their vevs, the neutrinos naturally acquire small Dirac masses,

$$\mathcal{L}_m \;=\; -\,(m_\nu)_{ij}\,\overline{\nu_{Li}}\nu_{Rj} + \text{h.c.}\,, \tag{4}$$

where

$$m_\nu \;\equiv\; m_\nu^I + m_\nu^{II}\,, \tag{5}$$

with [16]

$$\left(m_\nu^I\right)_{ij} \;=\; -\,f_i g_j\frac{\langle\phi\rangle\langle\chi\rangle}{M_D} \;=\; -\,\frac{(bc)_{ij}}{d}\,, \tag{6}$$

and [15]

$$\left(m_\nu^{II}\right)_{ij} \;=\; y_{ij}\frac{\mu\langle\phi\rangle\langle\chi\rangle}{M_\eta^2} \;=\; a_{ij}\,. \tag{7}$$

[1] The general Higgs potential $V(\phi,\eta,\chi)$ was given in the Appendix of our first paper in Ref. [15].

To quantify the second equality in (7), we note that different from the SM Higgs doublet, the heavy scalar doublet η has a positive mass-term in the Higgs potential, so it will develop a tiny nonzero *vev* until ϕ and χ both acquire their *vevs* [15],

$$\langle\eta\rangle \;\simeq\; \frac{\mu\langle\phi\rangle\langle\chi\rangle}{M_\eta^2}\,. \tag{8}$$

With this we can derive the neutrino mass formula $m_\nu^{II} = y\langle\eta\rangle \equiv a$ from the Lagrangian (1), which confirms the Eq. (7) above. In the reasonable parameter space of $M_D \sim M_\eta \sim \mu \gg \langle\chi\rangle\,,\langle\phi\rangle$ and $(f,\,g,\,y) = O(1)\,,$ we can naturally realize $a \ll b, c \ll d\,.$ Furthermore, using the second relations in (6) and (7) we can re-express the summed neutrino mass matrix as

$$m_\nu \;\equiv\; m_\nu^I + m_\nu^{II} = -bc/d + a\,. \tag{9}$$

This is consistent with the direct diagonalization of the original Dirac mass matrix (2), which we have mentioned below (2).

It is clear that this mechanism of the neutrino mass generation has two essential features: (i) it generates Dirac masses for neutrinos, and (ii) it retains the essence of the conventional seesaw [2] by making the neutrino masses tiny via the small ratio of the weak scale over the heavy mass scale. It is thus called Dirac Seesaw [15]. In particular, compared to the classification of the conventional type-I and type-II seesaw, we may refer to Eqs. (6) and (7) as the type-I and type-II Dirac seesaw, respectively.

From Eq. (9) we see that both type-I and type-II seesaws can contribute to the 3×3 mass-matrix m_ν for the light neutrinos. There are three possibilities in general: (i) $m_\nu^I \gg m_\nu^{II}$, or (ii) $m_\nu^I \sim m_\nu^{II}$, or (iii) $m_\nu^I \ll m_\nu^{II}$. We note that for case-(iii), the type-II contribution alone can accommodate the neutrino oscillation data even if type-I is fully negligible; while for case-(i) and -(ii), the type-II contribution should still play a nontrivial role for ν-mass generation because m_ν^I is rank-1 and additional contribution from m_ν^{II} is necessary. The rank-1 nature of $m_\nu^I = bc/d$ is due to that there is only one singlet heavy fermion in our current minimal construction, which means that m_ν^I has two vanishing mass-eigenvalues. Hence, to accommodate the neutrino oscillation data [1] in the case-(i) and -(ii) of our minimal construction always requires nonzero contribution m_ν^{II} from the type-II Dirac seesaw [2]. Let us explicitly analyze how this can be realized for the case-(i) and -(ii). As m_ν^I is rank-1, we can consider a basis for m_ν^I where one of the two massless states is manifest, i.e., $b_1 = c_1 = 0\,.$ For the remaining components of b and c, we choose a generic parameter set [3], $b_2 \approx b_3 \approx -c_2 \approx -c_3\,,$ which naturally realizes the maximal mixing angle $\theta_{23} = 45^\circ$ for explaining the atmospheric neutrino mixing. Including the type-II Dirac-seesaw matrix $m_\nu^{II} = a$ will then account for the other mixing angles $(\theta_{12},\,\theta_{13})$ and the two other neutrino masses. Thus, we can naturally realize the light neutrino mass-spectrum via both normal hierarchy (NH) and inverted hierarchy (IH) schemes. To be concrete, the NH-scheme is realized in our case-(i) where the type-II Dirac-seesaw matrix $m_\nu^{II} = a \equiv \delta \ll m_0 \sim m_\nu^I\,,$ with m_0 the neutrino mass scale (fixed by the atmospheric neutrino mass-squared-difference $\Delta_{\rm a}$ with $m_0 \equiv \sqrt{\Delta_{\rm a}}\,$) and its relations to the nonzero $(b_j,\,c_j)$ are defined via $b_j = \sqrt{m_0 d/2} + O(\delta)$ and $c_j = -\sqrt{m_0 d/2} + O(\delta)$ for $j = 2, 3$. Thus we have

$$m_\nu = -\frac{bc}{d} + a = m_0 \begin{pmatrix} 0 & 0 & 0 \\ 0 & \frac{1}{2} & \frac{1}{2} \\ 0 & \frac{1}{2} & \frac{1}{2} \end{pmatrix} + O(\delta)\,. \tag{10}$$

It is clear that Eq. (10) predicts the neutrino masses, $(m_1,\,m_2,\,m_3) = m_0(0,\,0,\,1) + O(\delta)\,,$ consistent with the NH mass-spectrum. Next, the IH-scheme can be realized in our case-(ii) where the type-II Dirac-seesaw matrix $m_\nu^{II} = a \equiv m_0{\rm diag}(1,0,0) + \delta \sim m_\nu^I$ with $\delta \ll m_0\,,$ while the structure of the type-I Dirac seesaw matrix m_ν^I remains the same,

$$m_\nu = -\frac{bc}{d} + a = m_0 \begin{pmatrix} 1 & 0 & 0 \\ 0 & \frac{1}{2} & \frac{1}{2} \\ 0 & \frac{1}{2} & \frac{1}{2} \end{pmatrix} + O(\delta)\,. \tag{11}$$

From this equation we deduce the neutrino masses, $(m_1,\,m_2,\,m_3) = m_0(1,\,1,\,0) + O(\delta)\,,$ consistent with the IH mass-spectrum. So far we have discussed all three possibilities, the case-(i), -(ii) and -(iii), regarding the relative contributions of the type-I versus type-II seesaw to the neutrino mass matrix m_ν for accommodating the oscillation data. The question of which one among these three possibilities is realized in nature should be answered by a more fundamental theory which can precisely predict the Yuakawa couplings and masses for D and η as well as the *vev* of χ. Finally, we also note that as the neutrinos being Dirac particles, our Dirac-seesaw construction will be consistent with the possible non-observation of the neutrinoless double beta decays $(0\nu\beta\beta)$ which will be tested in the upcoming $0\nu\beta\beta$-experiments [19].

The real scalar χ is expected to acquire its *vev* near the weak scale, so we will set $\langle\chi\rangle$ around $O(\text{TeV})$ [4]. Under this setup, it is straightforward to see that m_ν^I will

[2] Note that the type-I Dirac seesaw alone can accommodate the oscillation data once we extend the current minimal construction to include a second heavy fermion D' which makes m_ν^I rank 2; this is similar to the minimal (Majorana) neutrino seesaw studied before [18].

[3] Here we consider the difference between any two of the four components $|b_j|$ and $|c_j|$ $(j = 2, 3)$ to be much smaller themselves.

[4] Here we comment on the cosmological domain wall problem as-

be efficiently suppressed by the ratio of the weak scale over the heavy mass. For instance, we find that $m_\nu^I = O(0.1)\,\text{eV}$ for $M_D = O(10^{13-15})\,\text{GeV}$ and $(f,\,g,\,y) = O(0.1-1)$, where $\langle\phi\rangle \simeq 174\,\text{GeV}$. It also is reasonable to set the trilinear scalar coupling $|\mu|$ to be around the scale of the η mass M_η. In consequence, the neutrino mass m_ν^{II} in (7) will be highly suppressed, similar to m_ν^I. For example, we derive $m_\nu^{II} = O(0.1)\,\text{eV}$ for $M_\eta = O(10^{13-15})\,\text{GeV}$ and $(y,\,\mu/M_\eta) = O(0.1-1)$. So, we can naturally realize the Dirac neutrino masses around $O(0.1)\,\text{eV}$.

We now demonstrate how to generate the observed baryon asymmetry in our model by invoking the neutrinogenesis [10] mechanism. Since the sphaleron processes [11] have no direct effect on the right-handed neutrinos, and the effective Yukawa interactions of the Dirac neutrinos are too weak to reach the equilibrium until temperatures fall well below the weak scale, the lepton asymmetry stored in the left-handed leptons, which is equal but opposite to that stored in the right-handed neutrinos, can be partially converted to the baryon asymmetry by sphalerons. In particular, the final baryon asymmetry should be

$$B \;=\; \frac{28}{79}\,(B-L_{SM}) \;=\; -\frac{28}{79}L_{SM}\,, \tag{12}$$

for the SM with three generation fermions and one Higgs doublet.

In the pure type-I Dirac seesaw scenario [14], we can generate the CP asymmetry through the interferences between the tree-level decay and the self-energy loops if there exist at least two heavy fermion singlets. Similarly, the pure type-II Dirac seesaw model [15] also needs two heavy scalar doublets to obtain the self-energy loops in the decays. In the following, we shall focus on the minimal construction with only one heavy singlet fermion and one heavy doublet scalar to realize the radiative vertex corrections for the CP asymmetry, although further extensions are allowed in our current scenario.

In this framework, depending on the values of the masses and couplings, the leptogenesis can be realized either from the decay of the heavy singlet fermion or from the decay of the heavy doublet scalar. From the decay of the heavy singlet fermion to the left-handed leptons and the SM Higgs doublet, as shown in Fig. 2, the CP asymmetry is given by

$$\begin{aligned}\varepsilon^I \;&\equiv\; \frac{\Gamma\left(D_R\to\psi_L\phi^*\right)-\Gamma\left(D_R^c\to\psi_L^c\phi\right)}{\Gamma_D}\\ &=\; \frac{1}{4\pi}\frac{\text{Im}\left[\text{Tr}\left(f^\dagger y g^\dagger\right)\mu\right]M_\eta^2}{\left[\text{Tr}\left(f^\dagger f\right)+\frac{1}{2}\text{Tr}\left(g^\dagger g\right)\right]M_D^3}\\ &\qquad\times\ln\left(1+\frac{M_D^2}{M_\eta^2}\right)\,,\end{aligned} \tag{13}$$

where

$$\Gamma_D \;=\; \frac{1}{16\pi}\left[\text{Tr}\left(f^\dagger f\right)+\frac{1}{2}\text{Tr}\left(g^\dagger g\right)\right]M_D \tag{14}$$

is the total decay width of D or D^c. Here we have taken M_D to be real after proper phase rotation. Furthermore, from the decay of the heavy doublet scalar to the left-handed leptons and the right-handed neutrinos, a CP asymmetry can also be produced. It is given by the interference of the tree-level process with the one-loop vertex diagram as shown in Fig. 3,

$$\begin{aligned}\varepsilon^{II} \;&\equiv\; \frac{\Gamma\left(\eta\to\psi_L\nu_R^c\right)-\Gamma\left(\eta^*\to\psi_L^c\nu_R\right)}{\Gamma_\eta}\\ &=\; \frac{1}{4\pi}\frac{\text{Im}\left[\text{Tr}\left(f^\dagger y g^\dagger\right)\mu\right]M_D}{\left[\text{Tr}\left(y^\dagger y\right)M_\eta^2+|\mu|^2\right]}\ln\left(1+\frac{M_\eta^2}{M_D^2}\right)\end{aligned} \tag{15}$$

where

$$\Gamma_\eta \;=\; \frac{1}{16\pi}\left[\text{Tr}\left(y^\dagger y\right)+\frac{|\mu|^2}{M_\eta^2}\right]M_\eta \tag{16}$$

is the total decay width of η or η^*.

In the case where the masses of the heavy singlet fermion and heavy doublet scalar locate around the same scale, and also their couplings are of the same order of magnitude, the two types of asymmetry of Eqs. (13) and (15) can be both important for the neutrinogenesis. For illustration below, we will analyze two typical scenarios where one process dominates over the other.

Scheme-1 is defined for $M_D \ll M_\eta$ and $f\sim g\sim y$, under which the final left- or right-handed lepton asymmetry mainly comes from the pair decays of $(D,\,D^c)$. We can simplify the CP asymmetry (13) as

$$\begin{aligned}\varepsilon^I \;&\simeq\; \frac{1}{64\pi^2}\frac{M_D M_\eta^2\text{Im}\left[\text{Tr}\left(m_\nu^{I\dagger}m_\nu^{II}\right)\right]}{\langle\phi\rangle^2\langle\chi\rangle^2\Gamma_D}\\ &=\; \left[\frac{45}{(4\pi)^7 g_*}\right]^{\frac{1}{2}}\frac{1}{K_D}\frac{M_\eta^2}{M_D^2}\\ &\qquad\times\frac{M_{\text{Pl}}M_D\text{Im}\left[\text{Tr}\left(m_\nu^{I\dagger}m_\nu^{II}\right)\right]}{\langle\phi\rangle^2\langle\chi\rangle^2}\end{aligned} \tag{17}$$

sociated with spontaneous breaking of a discrete Z_2 symmetry. This problem arises during the phase transition (when the broken discrete symmetry gets restored at the transition temperature) because of the production of topological defects – domain walls which carry too much energy and trouble the standard big-bang cosmology [20]. This can be avoided by inflation as long as phase transition temperature is above the inflation scale [21, 22]. Another resolution [23] to the domain wall problem is realized by the possibility of symmetry non-restoration at high temperature [24]. It is also very possible that a discrete symmetry like Z_2 is not a basic symmetry but appears as a remnant of a continuous symmetry such as $U(1)$ which is free from the domain wall problem [25]. Finally, the Z_2 symmetry in our model can also be replaced by a global $U(1)_D$ as in the pure type-I Dirac seesaw model [16], and the phenomenology of the Goldstone boson associated with this $U(1)_D$ breaking was discussed in [14].

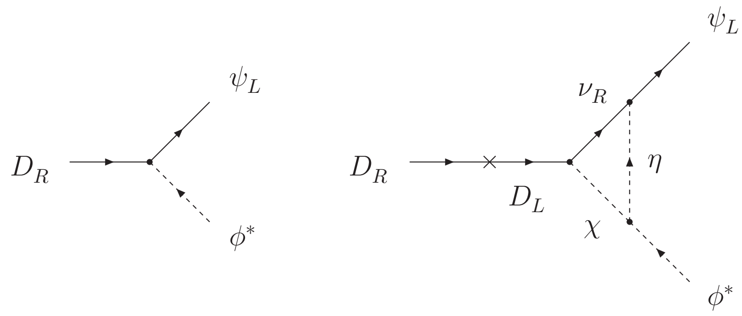

FIG. 2: The heavy singlet Dirac fermion decays to the left-handed leptons and the SM Higgs boson at one-loop order.

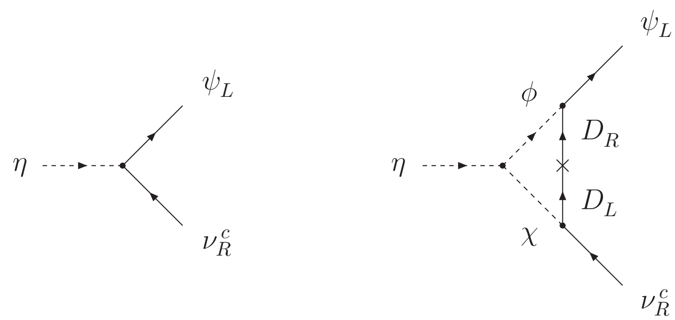

FIG. 3: The heavy doublet scalar decays to the leptons at one-loop order.

with

$$K_D \equiv \left.\frac{\Gamma_D}{H}\right|_{T=M_D} \qquad (18)$$

as a measurement of the deviation from equilibrium for D. Here H is the Hubble constant,

$$H(T) = \left(\frac{4\pi^3 g_*}{45}\right)^{\frac{1}{2}} \frac{T^2}{M_{\rm Pl}}, \qquad (19)$$

with $g_* = O(100)$ and $M_{\rm Pl} \simeq 1.2 \times 10^{19}\,\text{GeV}$. Note that there is a correlation between K_D and m_ν^I,

$$\begin{aligned}
\overline{m}_I^2 &\equiv \text{Tr}\left(m_\nu^{I\dagger} m_\nu^I\right) \\
&= \text{Tr}\left(f^\dagger f g g^\dagger\right) \frac{\langle\phi\rangle^2\langle\chi\rangle^2}{M_D^2} \\
&= \sum_i |f_i|^2 |g_i|^2 \frac{\langle\phi\rangle^2\langle\chi\rangle^2}{M_D^2} \\
&< \sum_i |f_i|^2 \sum_j |g_j|^2 \frac{\langle\phi\rangle^2\langle\chi\rangle^2}{M_D^2} \\
&= 2\,(16\pi)^2\, B_L B_R \Gamma_D^2 \frac{\langle\phi\rangle^2\langle\chi\rangle^2}{M_D^4} \\
&= \frac{2\,(4\pi)^5\, g_*}{45} B_L B_R K_D^2 \frac{\langle\phi\rangle^2\langle\chi\rangle^2}{M_{\rm Pl}^2}\,,
\end{aligned} \qquad (20)$$

and hence

$$K_D > \left[\frac{45}{2\,(4\pi)^5\, g_* B_L B_R}\right]^{\frac{1}{2}} \frac{M_{\rm Pl}\overline{m}_I}{\langle\phi\rangle\langle\chi\rangle}\,. \qquad (21)$$

Here B_L and B_R are the branching ratios of the heavy fermion singlet decaying into the left-handed lepton doublets and the right-handed neutrinos, respectively. They satisfy the following relationship,

$$B_L + B_R \equiv 1\,, \quad \Rightarrow \quad B_L B_R \leqslant \frac{1}{4}\,. \tag{22}$$

For instance, we may choose the sample inputs, $M_\eta = 10\,M_D = 1.8\times 10^{12}\,\mathrm{GeV} \gg M_D$, $\langle\phi\rangle = 174\,\mathrm{GeV}$, $\langle\chi\rangle = 400\,\mathrm{GeV}$ and $(y,\,f,\,g,\,\mu/M_\eta) \simeq (0.02,\,0.033,\,0.005,\,0.01) = O(0.01)$. Thus, we can estimate the light neutrino mass scale, $\overline{m}_I = O(m_\nu^I) = O(10\,m_\nu^{II}) \simeq 0.06\,\mathrm{eV}$. In consequence, we can estimate $B_L B_R \simeq 0.99\times 0.011 \simeq 0.011$, and thus $K_D \simeq 88$. This leads to $\varepsilon^I \simeq -2.4\times 10^{-5}$ for the maximal CP phase. We then use the approximate relation [21, 26] to deduce the final baryon asymmetry,

$$\begin{aligned} Y_B \equiv \frac{n_B}{s} \;&\simeq\; -\frac{28}{79}\times\frac{0.3\left(\varepsilon^I/g_*\right)}{K_D\left(\ln K_D\right)^{0.6}} \\ &\simeq\; 10^{-10}\,. \end{aligned} \tag{23}$$

This is consistent with the current observations [1]. Furthermore, the relationship $m_\nu^I = O(0.1\,\mathrm{eV}) \gg m_\nu^{II}$, shows the dominance of type-I Dirac seesaw. As we mentioned earlier, this can be realized via the NH mass-spectrum of the light neutrinos.

We also note that even if the Dirac fermion singlet is at a fairly low mass scale, such as TeV, it is still feasible to efficiently enhance the CP asymmetry as long as the ratio M_η/M_D is large enough. In other words, we can realize the low-scale neutrinogenesis without invoking the conventional resonant effect to enhance the CP asymmetry (which requires at least two heavy Dirac fermion singlets).

Scheme-2 is defined for the other possibility with $M_\eta \ll M_D$ and $f\sim g\sim y$. Hence the final left- or right-handed lepton asymmetry is dominated by the pair decays of $(\eta,\,\eta^*)$. We derive the following CP asymmetry from (15),

$$\begin{aligned} \varepsilon^{II} \;&\simeq\; \frac{1}{64\pi^2}\frac{M_\eta^3 \mathrm{Im}\left[\mathrm{Tr}\left(m_\nu^{I\dagger} m_\nu^{II}\right)\right]}{\langle\phi\rangle^2\langle\chi\rangle^2\Gamma_\eta} \\ &=\; \left[\frac{45}{(4\pi)^7\,g_*}\right]^{\frac{1}{2}}\frac{1}{K_\eta} \\ &\quad\times\frac{M_{\mathrm{Pl}} M_\eta \mathrm{Im}\left[\mathrm{Tr}\left(m_\nu^{I\dagger} m_\nu^{II}\right)\right]}{\langle\phi\rangle^2\langle\chi\rangle^2}\,, \end{aligned} \tag{24}$$

where K_η is given by

$$K_\eta \;\equiv\; \left.\frac{\Gamma_\eta}{H}\right|_{T=M_\eta}\,, \tag{25}$$

with the Hubble constant $H(T)$ expressed in Eq. (19). Here the parameter K_η measures the deviation from the equilibrium for η. We deduce the correlation between K_η and m_ν^{II},

$$\begin{aligned} \overline{m}_{II}^2 \;&\equiv\; \mathrm{Tr}\left(m_\nu^{II\dagger} m_\nu^{II}\right) \\ &=\; \mathrm{Tr}\left(y^\dagger y\right)\frac{|\mu|^2\langle\phi\rangle^2\langle\chi\rangle^2}{M_\eta^4} \\ &=\; (16\pi)^2 B_f B_s \Gamma_\eta^2\frac{\langle\phi\rangle^2\langle\chi\rangle^2}{M_\eta^4} \\ &=\; \frac{(4\pi)^5 g_*}{45} B_f B_s K_\eta^2 \frac{\langle\phi\rangle^2\langle\chi\rangle^2}{M_{\mathrm{Pl}}^2}\,, \end{aligned} \tag{26}$$

and also

$$K_\eta = \left[\frac{45}{(4\pi)^5\,g_* B_f B_s}\right]^{\frac{1}{2}}\frac{M_{\mathrm{Pl}}\overline{m}_{II}}{\langle\phi\rangle\langle\chi\rangle}\,, \tag{27}$$

where B_f and B_s are the branching ratios of the heavy scalar doublet decaying into the light fermions and the scalars, respectively. Similar to Eq. (22), they satisfy

$$B_f + B_s \equiv 1\,, \quad \Rightarrow \quad B_f B_s \leqslant \frac{1}{4}\,. \tag{28}$$

For instance, given the sample inputs, $M_\eta = 26\mu = 0.1\,M_D = 2\times 10^{13}\,\mathrm{GeV} \ll M_D$, $\langle\phi\rangle = 174\,\mathrm{GeV}$, $\langle\chi\rangle = 400\,\mathrm{GeV}$ and $(f,\,g,\,y) = (0.16,\,0.16,\,0.34) = O(0.1)$, we can estimate the light neutrino mass scale, $\overline{m}_{II} = O(m_\nu^{II}) = O(10\,m_\nu^I) \simeq 0.05\,\mathrm{eV}$. Subsequently, we derive, $B_f B_s \simeq 0.99\times 0.013 \simeq 0.012$, and $K_\eta \simeq 84$. This leads to, $\varepsilon^{II} \simeq -2.3\times 10^{-5}$, for the maximal CP phase. Using the approximate analytical formula [21, 26] for the baryon asymmetry, we arrive at

$$Y_B \;\simeq\; -\frac{28}{79}\times\frac{0.3\left(\varepsilon^{II}/g_*\right)}{K_\eta\left(\ln K_\eta\right)^{0.6}} \;\simeq\; 10^{-10}\,, \tag{29}$$

consistent with the present observation [1]. Furthermore, we note that in the Scheme-2, the active neutrino masses are dominated by the type-II Dirac seesaw, $m_\nu^{II} = O(0.1\,\mathrm{eV}) \gg m_\nu^I$, where both the NH and IH neutrino-mass-spectra can be realized.

In this paper, we have presented a new possibility to realize the neutrinogenesis in the Dirac seesaw scenario. In our minimal construction, we introduce a real scalar χ, a heavy singlet Dirac fermion D and a heavy doublet scalar η besides three right-handed singlet neutrinos to the SM. Therefore, different from previous realistic neutrinogenesis models, the *radiative vertex corrections* rather than the self-energy corrections interfere with the tree-level diagrams to generate the CP asymmetry in the decays of the heavy particles. Finally, we note that the real singlet scalar χ at the weak scale can couple to the SM Higgs doublet ϕ via the Z_2-conserving quartic interaction $\chi\chi\phi^\dagger\phi$. In consequence, the lightest neutral

Higgs boson h^0 is a mixture between ϕ^0 and χ, leading to non-SM-like anomalous couplings of h^0 with the weak gauge bosons ($W^{\pm}$, Z) and the SM-fermions. This can significantly modify the light Higgs boson (h^0) phenomenology at the Tevatron Run-2, the CERN LHC and the future International Linear Colliders (ILC) [27]. A systematical study for the collider phenomenology of ϕ^0 and χ is beyond the present scope and will be given elsewhere.

[1] Particle Data Group, W.M. Yao *et al.*, Journal of Physics G **33**, 1 (2006); R.N. Mohapatra, A.Y. Smirnov, Ann. Rev. Nucl. Part. Sci. **56**, 569 (2006); A. Strumia and F. Vissani, hep-ph/0606054; and references therein.

[2] P. Minkowski, Phys. Lett. B **67**, 421 (1977); T. Yanagida, in *Proc. of the Workshop on Unified Theory and the Baryon Number of the Universe*, ed. O. Sawada and A. Sugamoto (KEK, Tsukuba, 1979), p. 95; M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity*, ed. F. van Nieuwenhuizen and D. Freedman (North Holland, Amsterdam, 1979), p. 315; S.L. Glashow, in *Quarks and Leptons*, ed. M. Lévy *et al.* (Plenum, New York, 1980), p. 707; R.N. Mohapatra and G. Senjanović, Phys. Rev. Lett. **44**, 912 (1980); J. Schechter and J.W.F. Valle, Phys. Rev. D **22**, 2227 (1980).

[3] M. Fukugita and T. Yanagida, Phys. Lett. B **174**, 45 (1986); P. Langacker, R.D. Peccei, and T. Yanagida, Mod. Phys. Lett. A **1**, 541 (1986).

[4] M.A. Luty, Phys. Rev. D **45**, 455 (1992); R.N. Mohapatra and X. Zhang, Phys. Rev. D **46**, 5331 (1992).

[5] M. Flanz, E.A. Paschos, and U. Sarkar, Phys. Lett. B **345**, 248 (1995); M. Flanz, E.A. Paschos, U. Sarkar, and J. Weiss, Phys. Lett. B **389**, 693 (1996).

[6] E. Ma and U. Sarkar, Phys. Rev. Lett. **80**, 5716 (1998).

[7] S. Davidson and A. Ibarra, Phys. Lett. B **535**, 25 (2002); W. Buchmüller, P. Di Bari, and M. Plümacher, Nucl. Phys. B **643**, 390 (2002); P.H. Frampton, S.L. Glashow, and T. Yanagida, Phys. Lett. B **548**, 119 (2002); W. Buchmüller, P. Di Bari, and M. Plümacher, Nucl. Phys. B **665**, 445 (2003); G.F. Giudice, A. Notari, M. Raidal, A. Riotto, and A. Strumia, Nucl. Phys. B **685**, 89 (2004); V. Barger, D.A. Dicus, H.J. He, and T. Li, Phys. Lett. B **583**, 173 (2004); R.G. Felipe, F.R. Joaquim, and B.M. Nobre, Phys. Rev. D **70**, 085009 (2004); W. Rodejohann, Eur. Phys. J. C **32**, 235 (2004); P. Gu and X.J. Bi, Phys. Rev. D **70**, 063511 (2004); S. Antusch and S.F. King, Phys. Lett. B **597**, 199 (2004); R.N. Mohapatra and S. Nasri, Phys. Rev. D **71**, 033001 (2005); W. Buchmüller, P. Di Bari, and M. Plümacher, Annals Phys. **315**, 305 (2005); N. Sahu and U. Sarkar, Phys. Rev. D **74**, 093002 (2006); P.H. Gu, H. Zhang, and S. Zhou, Phys. Rev. D **74**, 076002 (2006); E.Kh. Akhmedov, M. Blennow, T. Hallgren, T. Konstandin, and T. Ohlsson, JHEP **0704**, 022 (2007); S. Antusch, Phys. Rev. D **76**, 023512 (2007); W. Buchmüller, R.D. Peccei, and T. Yanagida, Ann. Rev. Nucl. Part. Sci. **55**, 311 (2005); and references therein.

[8] For earlier related works, see M.Yu. Khlopov and A.D. Linde, Phys. Lett. B **138**, 265 (1984); F. Balestra *et al.*, Sov. J. Nucl. Phys. **39**, 626 (1984); M.Yu. Khlopov, Yu.L. Levitan, E.V. Sedelnikov, and I.M. Sobol, Phys. Atom. Nucl. **57**, 1393 (1994).

[9] E.Kh. Akhmedov, V.A. Rubakov, and A.Yu. Smirnov, Phys. Rev. Lett. **81**, 1359 (1998).

[10] K. Dick, M. Lindner, M. Ratz, and D. Wright, Phys. Rev. Lett. **84**, 4039 (2000).

[11] V.A. Kuzmin, V.A. Rubakov, and M.E. Shaposhnikov, Phys. Lett. B **155**, 36 (1985); R.N. Mohapatra and X. Zhang, Phys. Rev. D **45**, 2699 (1992).

[12] H. Murayama and A. Pierce, Phys. Rev. Lett. **89**, 271601 (2002); B. Thomas and M. Toharia, Phys. Rev. D **73**, 063512 (2006); B. Thomas and M. Toharia, Phys. Rev. D **75**, 013013 (2007).

[13] S. Abel and V. Page, JHEP **0605**, 024 (2006).

[14] D.G. Cerdeno, A. Dedes, and T.E.J. Underwood, JHEP **0609**, 067 (2006).

[15] P.H. Gu and H.J. He, JCAP **0612**, 010 (2006); P.H. Gu, H.J. He, and U. Sarkar, JCAP (2007), in press, arXiv:0705.3736 [hep-ph].

[16] M. Roncadelli and D. Wyler, Phys. Lett. B **133**, 325 (1983).

[17] P. Roy and O. Shanker, Phys. Rev. Lett. **52**, 713 (1984); S. Panda, U. Sarkar, Phys. Lett. B **139**, 42 (1984); A.S. Joshipura, P. Roy, O. Shanker, and U. Sarkar, Phys. Lett. B **150**, 270 (1985); A.S. Joshipura, A. Mukherjee, and U. Sarkar, Phys. Lett. B **156**, 353 (1985); U. Sarkar, Phys. Rev. D **35**, 1528 (1987); S. Mishra, S.P. Misra, S. Panda, and U. Sarkar, Phys. Rev. D **35**, 975 (1987); U. Sarkar, Phys. Rev. D **59**, 037302 (1999).

[18] E.g., P.H. Frampton, S.L. Glashow, and T. Yanagida, Phys. Lett. B **548**, 119 (2002); V. Barger, D.A. Dicus, H.J. He, and T. Li, Phys. Lett. B **583**, 173 (2004); and references therein.

[19] For reviews, P. Vogel, arXiv:hep-ph/0611243; S.R. Elliott and P. Vogel, Ann. Rev. Nucl. Phys. Sci. **52**, 115 (2002); and references therein.

[20] Y.B. Zeldovich, I.Y. Kobzarev, and L.B. Okun, Sov. Phys. JETP **40**, 1 (1974); T.W. Kibble, Phys. Rept. **67**, 183 (1980).

[21] E.W. Kolb and M. S. Turner, *The Early Universe* (Addison-Wesley Publishing Co, 1990).

[22] A. D. Linde, *Particle Physics and Inflationary Cosmology* (Harwood Academic, Switzerland, 1990).

[23] G. Dvali and G. Senjanovic, Phys. Rev. Lett. **74**, 5178 (1995).

[24] S. Weinberg, Phys. Rev. D **9**, 3357 (1974).

[25] G. Dvali and G. Senjanovic, Phys. Rev. Lett. **72**, 9 (1994).

[26] H.B. Nielsen and Y. Takanishi, Phys. Lett. B **507**, 241 (2001).

[27] For some related works, *e.g.* W. Buchmuller, C. Greub, and P. Minkowski, Phys. Lett. B **267**, 395 (1991); O. Bahat-Treidel, Y. Grossman, and Y. Rozen, JHEP **0705**, 022 (2007); J.R. Espinosa and M. Quirós, arXiv:hep-ph/0701145.