Realistic Neutrinogenesis with Radiative Vertex Correction

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We propose a new model for naturally realizing light Dirac neutrinos and explaining the baryon asymmetry of the universe through neutrinogenesis. To achieve this, 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.

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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\(^3\) in the universe poses a real challenge to the SM, but within the seesaw scenario, it can be naturally explained through leptogenesis\(^4, 5, 6, 7, 8\).

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 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\(^\text{I}\) in which $\psi_L, \phi, \nu_R, D_{L,R}, \eta$ and $\chi$ 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 $\psi_L, \nu_R, D_L$ and $D_R$ carry lepton number 1 while $\phi$, $\eta$ and $\chi$ 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-

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The parameters $M$ negative mass-terms and the Higgs doublet.

$\psi_L$ is the left-handed lepton doublet, $\phi$ is the SM Higgs doublet, $\nu_R$ is the right-handed neutrino, $D_{L,R}$ is the heavy singlet Dirac fermion, $\eta$ is the heavy doublet scalar, and $\chi$ is the real scalar. For simplicity the family indices and the other SM fields (carrying even $Z_2$ parity) are omitted from the Table.

TABLE I: The field content of our model, where $\psi_L$ is the left-handed lepton doublet, $\phi$ is the SM Higgs doublet, $\nu_R$ is the right-handed neutrino, $D_{L,R}$ is the heavy singlet Dirac fermion, $\eta$ is the heavy doublet scalar, and $\chi$ is the real scalar.

| Fields | $SU(2)_L$ | $U(1)_Y$ | $Z_2$ |
|--------|-----------|-----------|-------|
| $\psi_L$ | 2 | $-1/2$ | + |
| $\phi$ | 2 | $-1/2$ | + |
| $\nu_R$ | 1 | 0 | $-1/2$ |
| $D_{L,R}$ | 1 | 0 | + |
| $\eta$ | 2 | $-1/2$ | $-1/2$ |
| $\chi$ | 1 | 0 | $-1/2$ |

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,

$$- \mathcal{L} \supset \begin{cases} f_i \overline{\psi}_L i D_R + g_i \overline{D_L} \nu R_i + y_{ij} \overline{\psi}_L \eta \nu R_i - \mu \chi \eta^\dagger \phi \\ + M_D \overline{D_L} D_R + \text{h.c.} \end{cases} + M_\eta^2 \eta \eta$$

(1)

where $f_i$, $g_i$ and $y_{ij}$ are the Yukawa couplings, while the cubic scalar coupling $\mu$ has mass-dimension equal one. The parameters $M_D$ and $M_\eta$ in (1) are the masses of the heavy singlet fermion $D$ and the heavy Higgs doublet $\eta$, respectively. Note that in the Higgs potential the scalar doublet $\eta$ has a positive mass-term as shown in the above Eq. (1), while the Higgs doublet $\phi$ and singlet $\chi$ 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 $\eta$ 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

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

(2)

where $a \equiv y(\eta)$, $b \equiv f(\phi)$, $c \equiv g(\chi)$ 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,

$$\mathcal{O}_5 = \frac{f_i g_j}{M_D} \overline{\psi}_{Li} \overline{\nu} R_j \chi - \frac{\mu y_{ij}}{M_\eta^2} \overline{\psi}_{Li} \overline{\nu} R_j \chi + \text{h.c.}.$$ (3)

Therefore, once the SM Higgs doublet $\phi$ and the real scalar $\chi$ both acquire their vevs, the neutrinos naturally acquire small Dirac masses,

$$\mathcal{L}_m = - (m_\nu)_{ij} \overline{\nu}_i R_j + \text{h.c.},$$ (4)

where

$$m_\nu \equiv m_\nu' + m_\nu'',$$ (5)

with [16]

$$(m_\nu')_{ij} = - f_i g_j \frac{\langle \phi \rangle \langle \chi \rangle}{M_D} = - \frac{(bc)_{ij}}{d},$$ (6)

and [15]

$$(m_\nu'')_{ij} = y_{ij} \frac{\mu \langle \phi \rangle \langle \chi \rangle}{M_\eta^2} = a_{ij}.$$ (7)

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$^1$ 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 $\eta$ has a positive mass-term in the Higgs potential, so it will develop a tiny nonzero $\langle \eta \rangle$ until $\phi$ and $\chi$ both acquire their vevs $\langle \phi \rangle$ and $\langle \chi \rangle$.

$$\langle \eta \rangle \simeq \frac{\mu(\phi \langle \chi \rangle)}{M^2}.$$  \hspace{1cm} (8)

With this we can derive the neutrino mass formula $m_{\nu J} = y_\eta \langle \eta \rangle \equiv \alpha$ from the Lagrangian (1), which confirms the Eq. (1) above. In the reasonable parameter space of $M_D \sim M_\eta \sim \mu \gg \langle \chi \rangle, \langle \phi \rangle$ and $(f, \eta, y) = O(1)$, we can naturally realize $a \ll b, c \ll d$. Furthermore, using the second relations in (10) and (11) we can re-express the summed neutrino mass matrix as

$$m_{\nu} \equiv m_{\nu I} + m_{\nu II} = -bc/d + a.$$ \hspace{1cm} (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 (13). 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 $3 \times 3$ mass-matrix $m_{\nu}$ 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 $\nu$-mass generation because $m_{\nu I}$ is rank-1 allows additional contribution from $m_{\nu II}$ 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 $m_{\nu I}$ 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 $m_{\nu II}$ from the type-II Dirac seesaw (2). Let us explicitly analyze how this can be realized for the case-(i) and -(ii). As $m_{\nu I}$ is rank-1, we can consider a basis for $m_{\nu I}$ 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 $m_{\nu II} = a \equiv \delta \ll m_0 \sim m_I$, with $m_0$ the neutrino mass scale (fixed by the atmospheric neutrino mass-squared-difference $\Delta m^2$ with $m_0 = \sqrt{\Delta m^2}$) and its relations to the nonzero $(b_j, c_j)$ are defined via $b_j = \sqrt{m_0 d_j^2 + O(\delta)}$ and $c_j = -\sqrt{m_0 d_j^2 + O(\delta)}$ for $j = 2, 3$. Thus we have

$$m_{\nu I} = -\frac{bc}{d} + a = m_0 \begin{pmatrix} 0 & 0 & 0 \\ 0 & \frac{1}{2} & \frac{1}{2} \\ 0 & \frac{1}{2} & \frac{1}{2} \end{pmatrix} + O(\delta).$$ \hspace{1cm} (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 $m_{\nu II} = a \equiv \delta \ll m_0 \sim m_I$, with $m_0$ the neutrino mass scale (fixed by the atmospheric neutrino mass-squared-difference $\Delta m^2$ with $m_0 = \sqrt{\Delta m^2}$) and its relations to the nonzero $(b_j, c_j)$ are defined via $b_j = \sqrt{m_0 d_j^2 + O(\delta)}$ and $c_j = -\sqrt{m_0 d_j^2 + O(\delta)}$ for $j = 2, 3$. Thus we have

$$m_{\nu I} = -\frac{bc}{d} + a = m_0 \begin{pmatrix} 0 & 0 & 0 \\ 0 & \frac{1}{2} & \frac{1}{2} \\ 0 & \frac{1}{2} & \frac{1}{2} \end{pmatrix} + O(\delta).$$ \hspace{1cm} (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 $m_{\nu I}$ 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 Yukawa couplings and masses for $D$ and $\eta$ as well as the vev of $\chi$. 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 $\chi$ 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 $m_{\nu I}$ will

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$^2$ 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 $m_{\nu I}$ rank 2; this is similar to the minimal (Majorana) neutrino seesaw studied before (13).

$^3$ 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 $m_{ll}^I = O(0.1)\,eV$ for $M_D = O(10^{13-15})\,GeV$ and $(f, g, y) = O(0.1 - 1)$, where $\langle \phi \rangle \approx 174\,GeV$. It also is reasonable to set the trilinear scalar coupling $|\mu|_D$ to be around the scale of the $\eta$ mass $M_\eta$. In consequence, the neutrino mass $m_{\nu}^{II}$ in \[1\] will be highly suppressed, similar to $m_{\nu}^I$. For example, we derive $m_{\nu}^{II} = O(0.1)\,eV$ for $M_\eta = O(10^{13-15})\,GeV$ and $(y, \mu/M_\eta) = O(0.1 - 1)$. So, we can naturally realize the Dirac neutrino masses around $O(0.1)\,eV$.

We now demonstrate how to generate the observed baryon asymmetry in our model by invoking the neutrino genesis \[10\] 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

$$B = \frac{28}{79} (B - L_{SM}) = -\frac{28}{79} L_{SM}, \quad \text{for the SM with three generation fermions and one Higgs doublet.}$$

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

$$\epsilon^I \equiv \frac{\Gamma(D_R \rightarrow \psi_L^s \phi^*) - \Gamma(D_R \rightarrow \psi_L^s \phi)}{\Gamma_D}$$

\[14\]

$$= \frac{1}{4\pi} \text{Im} \left[ \text{Tr} (f^* y g^* \mu) \right] M_\eta^2$$

\[15\]

$$\times \ln \left( 1 + \frac{M_D^2}{M_\eta^2} \right), \quad \text{for the SM with three generation fermions and one Higgs doublet.}$$

where

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

\[16\]

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.\[3\].

$$\epsilon^{II} = \frac{\Gamma (\eta \rightarrow \psi_L^s \nu_R) - \Gamma (\eta^* \rightarrow \psi_L^s \nu_R)}{\Gamma_\eta} = \frac{1}{4\pi} \text{Im} \left[ \text{Tr} (f^* y g^* \mu) \right] M_\eta^2$$

\[17\]

$$\times \ln \left( 1 + \frac{M_D^2}{M_\eta^2} \right) \cdot \Gamma_D$$

$$= \frac{1}{4\pi} \left[ \text{Tr} (y^* y) M_\eta^2 + |\mu|^2 \right] \ln \left( 1 + \frac{M_D^2}{M_\eta^2} \right)$$

$$\times \frac{1}{16\pi} \left[ \text{Tr} (y^* y) + \frac{|\mu|^2}{M_\eta^2} \right] M_\eta$$

\[18\]

is the total decay width of $\eta$ or $\eta^*$. 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. \[14\] and \[15\] can be both important for the neutrino genesis. 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- right-hand lepton asymmetry mainly comes from the pair decays of $(D, D^c)$. We can simplify the CP asymmetry \[13\] as

$$\epsilon^I \simeq \frac{1}{64\pi^2} \frac{M_D M_\eta^2 \text{Im} \left[ \text{Tr} (m_{\nu}^{II} m_{\nu}^{II}) \right]}{(\phi)^2(\chi)^2 \Gamma_D}$$

\[19\]

$$= \left[ \frac{45}{(4\pi)^7 y_\eta} \right]^\frac{1}{2} \frac{1}{K_D M_D^2} M_\eta \text{Im} \left[ \text{Tr} (m_{\nu}^{II} m_{\nu}^{II}) \right]$$

\[20\]

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

\[
K_D = \frac{\Gamma_D}{H \bigg|_{T=M_D}} \tag{18}
\]

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

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

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

\[
\begin{align*}
\overline{m}_I^2 & = \text{Tr} \left( m_{L} m_{L}^T \right) \\
& = \text{Tr} \left( f^T f g g^T \right) \frac{\langle \phi \rangle^2 \langle \chi \rangle^2}{M_D^2} \\
& = \sum_i |f_i|^2 |g_i|^2 \frac{\langle \phi \rangle^2 \langle \chi \rangle^2}{M_D^2} \\
& < \sum_i |f_i|^2 \sum_j |g_j|^2 \frac{\langle \phi \rangle^2 \langle \chi \rangle^2}{M_D^2} \\
& = 2 (16\pi)^2 B_L B_R \Gamma_D^2 \frac{\langle \phi \rangle^2 \langle \chi \rangle^2}{M_D^2} \\
& = 2 (4\pi)^5 g_* B_L B_R K_D^2 \frac{\langle \phi \rangle^2 \langle \chi \rangle^2}{M_D^2}, \tag{20}
\end{align*}
\]

and hence

\[
K_D > \left[ \frac{45}{2 (4\pi)^5 g_* B_L B_R} \right] \frac{1}{2} M_{Pl} \overline{m}_L, \tag{21}
\]

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,

$$B_L + B_R \equiv 1, \quad \Rightarrow \quad B_L B_R \leq \frac{1}{4}, \quad (22)$$

For instance, we may choose the sample inputs, $M_\eta = 10 M_D = 1.8 \times 10^{12}$ GeV $\gg M_D$, $\langle \phi \rangle = 174$ GeV, $\langle \chi \rangle = 400$ 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, $\bar{m}_I = \mathcal{O}(m^I) = \mathcal{O}(10 m^I_{\nu}) \approx 0.06$ eV. In consequence, we can estimate $B_L B_R \approx 0.99 \times 0.011 \approx 0.011$, and thus $\Delta \approx 88$. This leads to $\varepsilon^I \approx -2.4 \times 10^{-5}$ for the maximal CP phase. We then use the approximate relation [21, 26] to deduce the final baryon asymmetry,

$$Y_B = \frac{n_B}{s} \approx -\frac{28}{79} \times \frac{0.3 (\varepsilon^I/g_s)}{K_D (\ln K_D^{0.6})} \approx 10^{-10}. \quad (23)$$

This is consistent with the current observations [1]. Furthermore, the relationship $m^I_\nu = \mathcal{O}(0.1 \text{ eV}) \gg m^I_{\nu}$, 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 $M_\eta/M_D$ is large enough. In other words, we can realize the low-scale neutrinosynthesis 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),

$$\varepsilon^I \approx \frac{1}{64 \pi^2} \frac{M^2_\eta \text{Im} \left[ \text{Tr} \left( m^I_{\nu} m^I_{\nu} \right) \right]}{\langle \phi \rangle^2 \langle \chi \rangle^2 \Gamma_\eta} \approx \frac{1}{K_\eta} \frac{45}{(4\pi)^2 g_s} \frac{1}{K_\eta} \times \frac{M_{Pl} M_\eta \text{Im} \left[ \text{Tr} \left( m^I_{\nu} m^I_{\nu} \right) \right]}{\langle \phi \rangle^2 \langle \chi \rangle^2}, \quad (24)$$

where $K_\eta$ is given by

$$K_\eta \equiv \frac{\Gamma_\eta}{H_{T=M_\eta}}, \quad (25)$$

with the Hubble constant $H(T)$ expressed in Eq. (13). Here the parameter $K_\eta$ measures the deviation from the equilibrium for $\eta$. We deduce the correlation between $K_\eta$ and $m_{\nu}^I$.}

$$\bar{m}_{II}^2 \equiv \text{Tr} \left( m^I_{\nu} m^I_{\nu} \right) \approx \text{Tr} \left( y^I y^I \right) \frac{\mu^2 \langle \phi \rangle^2}{\langle \chi \rangle^2} \frac{N^2}{M^2_{Pl}}, \quad (26)$$

and also

$$K_\eta \equiv \left[ \frac{45}{(4\pi)^2 g_s B_f B_s} \right]^{1/2} \frac{M^I_{\eta}}{\langle \phi \rangle^2 \langle \chi \rangle^2}, \quad (27)$$

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

$$B_f + B_s \equiv 1, \quad \Rightarrow \quad B_f B_s \leq \frac{1}{4}. \quad (28)$$

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

$$Y_B \approx -\frac{28}{79} \times \frac{0.3 (\varepsilon^I/g_s)}{K_\eta (\ln K_\eta^{0.6})} \approx 10^{-10}, \quad (29)$$

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, $m^I_{\nu} = \mathcal{O}(0.1 \text{ eV}) \gg m^I_{\nu}$, where both the NH and IH neutrino-mass-spectra can be realized.

In this paper, we have presented a new possibility to realize the neutrinosynthesis in the Dirac seesaw scenario. In our minimal construction, we introduce a real scalar $\chi$, a heavy singlet Dirac fermion $D$ and a heavy doublet scalar $\eta$ besides three right-handed singlet neutrinos to the SM. Therefore, different from previous realistic neutrinosynthesis 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 $\chi$ at the weak scale can couple to the SM Higgs doublet $\phi$ via the $Z_2$-conserving quartic interaction $\chi \chi \phi \phi$. In consequence, the lightest neutral
Higgs boson $h^0$ is a mixture between $\phi^0$ and $\chi$, 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 $\phi^0$ and $\chi$ is beyond the present scope and will be given elsewhere.

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