Leptogenesis, Dark Matter and Higgs Phenomenology at TeV

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We propose an interesting model of neutrino masses to realize leptogenesis and dark matter at the TeV scale. A real scalar is introduced to naturally realize the Majorana masses of the right-handed neutrinos. We also include a new Higgs doublet that contributes to the dark matter of the universe. The neutrino masses come from the vacuum expectation value of the triplet Higgs scalar. The right-handed neutrinos are not constrained by the neutrino masses and hence they could generate leptogenesis at the TeV scale without subscribing to resonant leptogenesis. In our model, all new particles could be observable at the forthcoming Large Hardon Collider or the proposed future International Linear Collider.

Introduction: Currently many neutrino oscillation experiments [1] have confirmed that neutrinos have tiny but nonzero masses. This phenomenon is elegantly explained by the seesaw mechanism [2], in which neutrinos naturally acquire small Majorana [2] or Dirac [3] masses. At the same time, the observed matter-antimatter asymmetry [1] in the universe can be generated via leptogenesis [3, 4, 5, 6, 7, 8, 9, 10, 11, 12] in these models. In this scheme, the right-handed neutrinos or other virtual particles need have very large masses to realize a successful leptogenesis if we don't resort to the resonant effect [6, 7] by highly fine tuning.

Another big challenge to the standard model is the dark matter [1]. What is the nature of dark matter? Recently, it has been pointed out [13, 14, 15, 16] that a new Higgs doublet can be a candidate for the dark matter if it doesn't decay into the standard model particles. Although the possibility of Higgs doublet to be a dark matter candidate was proposed many years back [17], following the recent proposal [13] a thorough analysis have been carried out [14, 18] demonstrating its consistency with all the recent results. In this interesting scenario, the dark matter is expected to produce observable signals at the Large Hardon Collider (LHC) [14] and in the GLAST satellite experiment [18]. Combining this idea, the type-I seesaw and the concept [19] of generation of the cosmological matter-antimatter asymmetry along with the cold dark matter, the author of [15] successfully unified the leptogenesis and dark matter. However, this scenario need the right-handed neutrinos to be very heavy, around the order of 10^7 GeV.

In this paper, we propose a new scheme to explain neutrino masses, baryon asymmetry and dark matter at TeV scale by introducing a Higgs triplet which is responsible for the origin of neutrino masses, a new Higgs doublet that can be a candidate for the dark matter, and a real scalar which can generate the Majorana masses of the right-handed neutrinos naturally. A discrete symmetry ensures that the new Higgs doublet cannot couple to ordinary particles. This same discrete symmetry will also prevent any connection between the right-handed neutrinos and left-handed neutrino masses. This allows the right-handed neutrinos to decay at low scale generating the lepton asymmetry, which will be finally converted to the baryon asymmetry through the sphaleron processes [20]. This will then explain the observed matter-antimatter asymmetry in the universe, even if the Majorana masses of the right-handed neutrinos are not highly quasi-degenerate. In our model, all new particles could be close to the TeV scale and hence should be observable at the forthcoming LHC or the proposed future International Linear Collider (ILC).

The model: We extend the standard model with some new fields. The field content is shown in Table I, in which

$$\psi_L = \begin{pmatrix} \nu_L \\ l_L \end{pmatrix}, \quad \phi = \begin{pmatrix} \phi^0 \\ \phi^- \end{pmatrix} \tag{1}$$

are the left-handed lepton doublet and Higgs doublet of the standard model, respectively, while

$$\eta = \begin{pmatrix} \eta^0 \\ \eta^- \end{pmatrix} \tag{2}$$

is the new Higgs doublet that will be the dark matter candidate, ν_R is the right-handed neutrino, χ is the real scalar and

$$\Delta_L = \begin{pmatrix} \frac{1}{\sqrt{2}}\delta^+ & \delta^{++} \\ \delta^0 & -\frac{1}{\sqrt{2}}\delta^+ \end{pmatrix} \tag{3}$$

is the Higgs triplet. We further introduce a discrete Z_4 symmetry, under which the different fields transform as

$$\begin{aligned} \psi_L &\to \psi_L\,, \quad \phi \to \phi\,, \quad \eta \to -i\eta\,, \\ \nu_R &\to i\nu_R\,, \quad \chi \to -\chi\,, \quad \Delta_L \to \Delta_L\,. \end{aligned} \tag{4}$$

Here the other standard model fields, which are all even under the Z_4, and the family indices have been omitted for simplicity.

We write down the relevant Lagrangian for the Yukawa interactions,

$$\begin{aligned} -\mathcal{L} \supset \sum_{ij} \Big(& y_{ij}\overline{\psi_{Li}}\eta\nu_{Rj} + \frac{1}{2}g_{ij}\chi\overline{\nu^c_{Ri}}\nu_{Rj} \\ + & \frac{1}{2}f_{ij}\overline{\psi^c_{Li}}i\tau_2\Delta_L\psi_{Lj} + \text{h.c.}\Big)\,, \end{aligned} \tag{5}$$

Fields	ψ_L	ϕ	η	ν_R	χ	Δ_L
$SU(2)_L$	**2**	**2**	**2**	**1**	**1**	**3**
$U(1)_Y$	$-\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	0	0	1

TABLE I: The field content in the model. Here ψ_L, ϕ are the standard model left-handed lepton doublets and Higgs doublet, η is the new Higgs doublet, ν_R is the right-handed neutrinos, χ is the real scalar and Δ_L is the Higgs triplet. Here the other standard model fields and the family indices have been omitted for simplicity.

where y_{ij}, g_{ij}, f_{ij} are all dimensionless. We also display the general scalar potential of ϕ, η, χ and Δ_L,

$$\begin{aligned} & V(\chi,\phi,\eta,\Delta_L) \\ = & \ \frac{1}{2}\mu_1^2\chi^2+\frac{1}{4}\lambda_1\chi^4+\mu_2^2\left(\phi^\dagger\phi\right)+\lambda_2(\phi^\dagger\phi)^2 \\ + & \ \mu_3^2\left(\eta^\dagger\eta\right)+\lambda_3(\eta^\dagger\eta)^2+M_\Delta^2\mathrm{Tr}\left(\Delta_L^\dagger\Delta_L\right) \\ + & \ \lambda_4\mathrm{Tr}\left[\left(\Delta_L^\dagger\Delta_L\right)^2\right]+\lambda_5\left[\mathrm{Tr}\left(\Delta_L^\dagger\Delta_L\right)\right]^2 \\ + & \ \alpha_1\chi^2\left(\phi^\dagger\phi\right)+\alpha_2\chi^2\left(\eta^\dagger\eta\right)+\alpha_3\chi^2\mathrm{Tr}\left(\Delta_L^\dagger\Delta_L\right) \\ + & \ 2\beta_1\left(\phi^\dagger\phi\right)\left(\eta^\dagger\eta\right)+2\beta_2\left(\phi^\dagger\eta\right)\left(\eta^\dagger\phi\right) \\ + & \ 2\beta_3\left(\phi^\dagger\phi\right)\mathrm{Tr}\left(\Delta_L^\dagger\Delta_L\right)+2\beta_4\phi^\dagger\Delta_L^\dagger\Delta_L\phi \\ + & \ 2\beta_5\left(\eta^\dagger\eta\right)\mathrm{Tr}\left(\Delta_L^\dagger\Delta_L\right)+2\beta_6\eta^\dagger\Delta_L^\dagger\Delta_L\eta \\ + & \ \left(\mu\phi^Ti\tau_2\Delta_L\phi+\kappa\chi\eta^Ti\tau_2\Delta_L\eta+\mathrm{h.c.}\right), \end{aligned} \tag{6}$$

where $\mu_{1,2,3}$ and μ have the mass dimension-1, while $\lambda_{1,\ldots,5}$, $\alpha_{1,2,3}$, $\beta_{1,\ldots,6}$ and κ are all dimensionless, M_Δ^2 is the positive mass-square of the Higgs triplet. Without loss of generality, μ and κ will be conveniently set as real after proper phase rotations.

The vacuum expectation values: For $\lambda_1>0$ and $\mu_1^2<0$, we can guarantee that before the electroweak phase transition, the real scalar χ acquires a nonzero vacuum expectation value (VEV),

$$\langle\chi\rangle\equiv u=\sqrt{-\frac{\mu_1^2}{\lambda_1}}\,. \tag{7}$$

We can then write the field χ in terms of the real physical field σ as

$$\chi\equiv\sigma+u\,, \tag{8}$$

so that the explicit form of the Yukawa couplings become

$$\begin{aligned} -\mathcal{L} \supset & \ y_{ij}\overline{\psi_{Li}}\eta\nu_{Rj}+\frac{1}{2}M_{ij}\overline{\nu_{Ri}^c}\nu_{Rj}+\frac{1}{2}f_{ij}\overline{\psi_{Li}^c}i\tau_2\Delta_L\psi_{Lj} \\ + & \ \mu\phi^Ti\tau_2\Delta_L\phi+\tilde{\mu}\eta^Ti\tau_2\Delta_L\eta+\frac{1}{2}g_{ij}\sigma\overline{\nu_{Ri}^c}\nu_{Rj} \\ + & \ \kappa\sigma\eta^Ti\tau_2\Delta_L\eta+\mathrm{h.c.}+M_\Delta^2\mathrm{Tr}\left(\Delta_L^\dagger\Delta_L\right), \end{aligned} \tag{9}$$

where we defined,

$$M_{ij}\equiv g_{ij}u \quad \text{and} \quad \tilde{\mu}\equiv\kappa u. \tag{10}$$

For convenience, we diagonalize $g_{ij}\to g_i$ as well as $M_{ij}\to M_i$ by redefining ν_{Ri} and then simplify the Lagrangian (9) as

$$\begin{aligned} -\mathcal{L} \supset & \ y_{ij}\overline{\psi_{Li}}\eta N_j+\frac{1}{2}f_{ij}\overline{\psi_{Li}^c}i\tau_2\Delta_L\psi_{Lj}+\mu\phi^Ti\tau_2\Delta_L\phi \\ + & \ \tilde{\mu}\eta^Ti\tau_2\Delta_L\eta+\kappa\sigma\eta^Ti\tau_2\Delta_L\eta+\mathrm{h.c.} \\ + & \ \frac{1}{2}g_i\sigma\overline{N_i}N_i+\frac{1}{2}M_i\overline{N_i}N_i+M_\Delta^2\mathrm{Tr}\left(\Delta_L^\dagger\Delta_L\right) \end{aligned} \tag{11}$$

with

$$N_i\equiv\nu_{Ri}+\nu_{Ri}^c \tag{12}$$

being the heavy Majorana neutrinos.

After the electroweak symmetry breaking, we denote the different VEVs as $\langle\phi\rangle\equiv\frac{1}{\sqrt{2}}v$, $\langle\eta\rangle\equiv\frac{1}{\sqrt{2}}v'$, $\langle\Delta_L\rangle\equiv\frac{1}{\sqrt{2}}v_L$ and $\langle\chi\rangle\equiv u'$ and then analyze the potential as a function of these VEVs,

$$\begin{aligned} & V(u',v,v',v_L) \\ = & \ \frac{1}{2}\mu_1^2u'^2+\frac{1}{4}\lambda_1u'^4+\frac{1}{2}\mu_2^2v^2+\frac{1}{4}\lambda_2v^4 \\ + & \ \frac{1}{2}\mu_3^2v'^2+\frac{1}{4}\lambda_3v'^2+\frac{1}{2}M_\Delta^2v_L^2+\frac{1}{4}(\lambda_4+\lambda_5)v_L^4 \\ + & \ \frac{1}{2}\alpha_1u'^2v^2+\frac{1}{2}\alpha_2u'^2v'^2+\frac{1}{2}\alpha_3u'^2v_L^2 \\ + & \ \frac{1}{2}\left(\beta_1+\beta_2\right)v^2v'^2+\frac{1}{2}\left(\beta_3+\beta_4\right)v^2v_L^2 \\ + & \ \frac{1}{2}\left(\beta_5+\beta_6\right)v'^2v_L^2+\frac{1}{\sqrt{2}}\mu v^2v_L+\frac{1}{\sqrt{2}}\tilde{\mu}'v'^2v_L \end{aligned} \tag{13}$$

with $\tilde{\mu}'\equiv\kappa u'$. Using the extremum conditions, $0=\partial V/\partial u'=\partial V/\partial v=\partial V/\partial v'=\partial V/\partial v_L$, we obtain,

$$\begin{aligned} 0 = & \ \lambda_1u'^3+\mu_1^2u'+\alpha_1v^2u'+\alpha_2v'^2u'+\alpha_3v_L^2u' \\ + & \ \frac{1}{\sqrt{2}}\kappa v'^2v_L\,, \end{aligned} \tag{14}$$

$$\begin{aligned} 0 = & \ \mu_2^2+\alpha_1u'^2+\left(\beta_1+\beta_2\right)v'^2+\left(\beta_3+\beta_4\right)v_L^2 \\ + & \ 2\sqrt{2}\mu v_L+\lambda_2v^2\,, \end{aligned} \tag{15}$$

$$\begin{aligned} 0 = & \ \mu_3^2+\alpha_2u'^2+\left(\beta_1+\beta_2\right)v^2+\left(\beta_5+\beta_6\right)v_L^2 \\ + & \ 2\sqrt{2}\tilde{\mu}'v_L+\lambda_3v'^2\,, \end{aligned} \tag{16}$$

$$\begin{aligned} 0 = & \ \frac{1}{\sqrt{2}}\mu v^2+\frac{1}{\sqrt{2}}\tilde{\mu}'v'^2+\left[M_\Delta^2+\alpha_3u'^2+\left(\beta_3+\beta_4\right)v^2\right. \\ + & \ \left.\left(\beta_5+\beta_6\right)v'^2\right]v_L+\left(\lambda_4+\lambda_5\right)v_L^3\,. \end{aligned} \tag{17}$$

For

$$\begin{cases} \lambda_3>0\,, \\ \mu_3^2+\alpha_2u'^2+\left(\beta_1+\beta_2\right)v^2+\left(\beta_5+\beta_6\right)v_L^2 \\ +2\sqrt{2}\tilde{\mu}'v_L>0\,, \end{cases} \tag{18}$$

the new Higgs doublet η gets a zero VEV, i.e., $v' = 0$. We assume $\mu < M_\Delta$ and $v^2 \ll M_\Delta^2$, u'^2, and then deduce

$$\begin{aligned} v_L &\simeq \frac{1}{\sqrt{2}} \frac{\mu v^2}{M_\Delta^2 + \alpha_3 u'^2 + (\beta_3 + \beta_4) v^2} \\ &\simeq \frac{1}{\sqrt{2}} \frac{\mu v^2}{M_\Delta^2 + \alpha_3 u'^2} \\ &\simeq \frac{1}{\sqrt{2}} \frac{\mu v^2}{M_\Delta^2} \quad \text{for} \quad M_\Delta^2 \gg \alpha_3 u'^2 \,. \end{aligned} \tag{19}$$

Subsequently, u' and v can be solved,

$$\begin{aligned} u' &= \sqrt{-\frac{\mu_1^2 + \alpha_1 v^2 + \alpha_3 v_L^2}{\lambda_1}} \\ &\simeq \sqrt{-\frac{\mu_1^2 + \alpha_1 v^2}{\lambda_1}} \,, \end{aligned} \tag{20}$$

$$\begin{aligned} v &= \sqrt{-\frac{\mu_2^2 + \alpha_1 u'^2 + (\beta_3 + \beta_4) v_L^2 + 2\sqrt{2}\mu v_L}{\lambda_2}} \\ &\simeq \sqrt{-\frac{\mu_2^2 + \alpha_1 u'^2}{\lambda_2}} \,, \end{aligned} \tag{21}$$

for

$$\begin{cases} \lambda_1 > 0 \,, \\ \mu_1^2 + \alpha_1 v^2 + \alpha_3 v_L^2 < 0 \,, \end{cases} \tag{22}$$

$$\begin{cases} \lambda_2 > 0 \,, \\ \mu_2^2 + \alpha_1 u^{'2} + (\beta_3 + \beta_4) v_L^2 + 2\sqrt{2}\mu v_L < 0 \,, \end{cases} \tag{23}$$

We then obtain the masses of resulting physical scalar bosons after the electroweak symmetry breaking,

$$M_{\delta^{++}}^2 \simeq M_\Delta^2 + \alpha_3 u'^2 + (\beta_3 + \beta_4) v^2 \,, \tag{24}$$

$$M_{\delta^+}^2 \simeq M_\Delta^2 + \alpha_3 u'^2 + \left(\beta_3 + \frac{1}{2}\beta_4\right) v^2 \,, \tag{25}$$

$$M_{\delta^0}^2 \simeq M_\Delta^2 + \alpha_3 u'^2 + \beta_3 v \,, \tag{26}$$

$$m_{\eta^\pm}^2 \simeq \mu_3^2 + \alpha_2 u'^2 + \beta_1 v^2 \,, \tag{27}$$

$$m_{\eta_R^0}^2 \simeq \overline{m}_\eta^2 + \delta m_\eta^2 \,, \tag{28}$$

$$m_{\eta_I^0}^2 \simeq \overline{m}_\eta^2 - \delta m_\eta^2 \,, \tag{29}$$

$$m_{h_1}^2 \simeq \overline{m}_h^2 - \delta m_h^2 \,, \tag{30}$$

$$m_{h_2}^2 \simeq \overline{m}_h^2 + \delta m_h^2 \,, \tag{31}$$

with

$$\overline{m}_\eta^2 \equiv \mu_3^2 + \alpha_2 u'^2 + (\beta_1 + \beta_2) v^2 \,, \tag{32}$$

$$\delta m_\eta^2 \equiv \frac{\tilde{\mu}' \mu}{M_\Delta^2 + \alpha_3 u'^2} v^2 \simeq \frac{\tilde{\mu}' \mu}{M_\Delta^2} v^2 \,, \tag{33}$$

$$\overline{m}_h^2 \equiv \lambda_1 u'^2 + \lambda_2 v^2 \,, \tag{34}$$

$$\delta m_h^2 \equiv \left[\left(\lambda_1 u'^2 - \lambda_2 v^2\right)^2 + 4\alpha_1^2 u'^2 v^2 \right]^{\frac{1}{2}} \,. \tag{35}$$

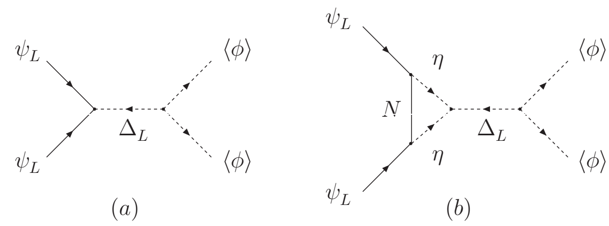

FIG. 1: The neutrino mass-generation. (a) is the type-II seesaw. (b) is the radiative contribution.

Here η^+ and $\eta_{R,I}^0$ are defined by

$$\eta^+ \equiv \left(\eta^-\right)^* \,, \tag{36}$$

$$\eta^0 \equiv \frac{1}{\sqrt{2}} \left(\eta_R^0 + i\eta_I^0\right) \,. \tag{37}$$

In addition, the mass eigenstates $h_{1,2}$ are the linear combinations of h and σ', i.e.,

$$h_1 \equiv \sigma' \sin\vartheta + h \cos\vartheta \,, \tag{38}$$

$$h_2 \equiv \sigma' \cos\vartheta - h \sin\vartheta \,, \tag{39}$$

where h, σ' are defined by

$$\phi \equiv \frac{1}{\sqrt{2}} \begin{pmatrix} v + h \\ 0 \end{pmatrix} \,, \quad \chi \equiv u' + \sigma' \,, \tag{40}$$

and the mixing angle is given by

$$\tan 2\vartheta \simeq \frac{2\alpha_1 u' v}{\lambda_2 v^2 - \lambda_1 u'^2} \,. \tag{41}$$

Neutrino masses: The first diagram of Fig. 1 shows the type-II seesaw approach to the generation of the neutrino masses. It is reasonable to take the scalar cubic coupling μ less than the triplet mass M_Δ in (19). In consequence, the triplet VEV in (19) is seesaw-suppressed by the ratio of the electroweak scale v over the heavy mass M_Δ. Substantially, the neutrinos naturally obtain the small Majorana masses,

$$(m_\nu^{II})_{ij} \equiv \frac{1}{\sqrt{2}} f_{ij} v_L \simeq -f_{ij} \frac{\mu v^2}{2M_\Delta^2} \,. \tag{42}$$

For the zero VEV of new Higgs doublet η, we can not realize the neutrino masses via the type-I seesaw. However, similar to [15], it is possible to generate the radiative neutrino masses at one-loop order due to the trilinear scalar interactions in (11). As shown in the second diagram of Fig. 1, the one-loop process will induce a contribution to the neutrino masses,

$$\begin{aligned} (\widetilde{m}_\nu^I)_{ij} &= \frac{1}{16\pi^2} \sum_k y_{ik} y_{jk} M'_k \left[\frac{m_{\eta_R^0}^2}{m_{\eta_R^0}^2 - M'^2_k} \ln\left(\frac{m_{\eta_R^0}^2}{M'^2_k}\right) \right. \\ &\quad \left. - \frac{m_{\eta_I^0}^2}{m_{\eta_I^0}^2 - M'^2_k} \ln\left(\frac{m_{\eta_I^0}^2}{M'^2_k}\right) \right] \,. \end{aligned} \tag{43}$$

Here $M'_k \equiv \frac{u'}{u} M_k$. For $|\mu_1^2| \gg |\alpha_1| v^2$, we have $u' \simeq u$ and then $M'_k \simeq M_k$, so the above formula can be simplified as

$$(\widetilde{m}_\nu^I)_{ij} \simeq \frac{1}{16\pi^2} \sum_k y_{ik} y_{jk} \frac{1}{M_k} \times \left[m^2_{\eta_R^0} \ln\left(\frac{M_k^2}{m^2_{\eta_R^0}}\right) - m^2_{\eta_I^0} \ln\left(\frac{M_k^2}{m^2_{\eta_I^0}}\right)\right] . \quad (44)$$

by taking $m^2_{\eta^0_{R,I}} \ll M_k^2$. Moreover, from (32) and (33), if $|\tilde{\mu}'\mu| \ll M_\Delta^2$, we have $\delta m_\eta^2 \ll \overline{m}_\eta^2$ and then obtain

$$\begin{aligned}(\widetilde{m}_\nu^I)_{ij} &\simeq -\frac{1}{8\pi^2} \sum_k y_{ik} y_{jk} \frac{\delta m_\eta^2}{M_k} \left[1 - \ln\left(\frac{M_k^2}{\overline{m}_\eta^2}\right)\right] \\ &= -\xi \sum_k y_{ik} y_{jk} \frac{v^2}{2M_k} \end{aligned} \quad (45)$$

for

$$\begin{aligned}\xi &= \mathcal{O}\left(\frac{1}{4\pi^2} \frac{\delta m_\eta^2}{v^2} \left[1 - \ln\left(\frac{M_k^2}{\overline{m}_\eta^2}\right)\right]\right) \\ &= \mathcal{O}\left(\frac{1}{4\pi^2} \frac{\tilde{\mu}'\mu}{M_\Delta^2} \left[1 - \ln\left(\frac{M_k^2}{\overline{m}_\eta^2}\right)\right]\right) . \end{aligned} \quad (46)$$

Note that the above loop-contribution will be absent once the values of κ and then $\tilde{\mu}'$ are taken to be zero.

Baryon asymmetry: We now demonstrate how the observed baryon asymmetry is generated in this model. In the Lagrangian (11), the lepton number of the left-handed lepton doublets and the Higgs triplet are 1 and -2, respectively, while those of the heavy Majorana neutrinos, the Higgs doublets and the real scalar are all zero. There are two sources of lepton number violation, one is the trilinear interaction between the Higgs triplet and the Higgs doublets, the other is the Yukawa couplings of the heavy Majorana neutrinos to the left-handed lepton doublet and the new Higgs doublet. Therefore, both the Higgs triplet and the heavy Majorana neutrinos could decay to produce the lepton asymmetry if their decays are CP-violation and out-of-equilibrium[1].

[1] Note that there is an equivalent choice of lepton number: $L = 1$ for η and $L = 0$ for ν_R, which makes only the $\mu\phi^T i\tau_2 \Delta_L \phi$ term to be lepton number violating. So, the CP asymmetry in the decays of N_i and Δ_L can only create an asymmetry in the numbers of ψ_L and an equal and opposite amount of asymmetry in the numbers of η. Thus there is no net lepton number asymmetry at this stage. However, since only the left-handed fields take part in the sphaleron transitions, only the ψ_L asymmetry gets converted to a $B-L$ asymmetry before the electroweak phase transition. After the electroweak phase transition, we are thus left with a baryon asymmetry equivalent to the $B-L$ asymmetry generated from the ψ_L asymmetry and an equivalent amount of η asymmetry or lepton number asymmetry, which does not affect the baryon asymmetry of the universe. In the rest of the article we shall not discuss this possibility, since the final amount of baryon asymmetry comes out to be the same.

We can obtain the CP asymmetry in the decay of N_i through the interference between the tree-level process and three one-loop diagrams of Fig. 3, in which the first two one-loop diagrams are the ordinary self-energy and vertex correction involving another heavy Majorana neutrinos, while the third one-loop diagram is mediated by the Higgs triplet [21]. So it is convenient to divide the total CP asymmetry into two independent parts,

$$\begin{aligned}\varepsilon_i &\equiv \frac{\sum_j \left[\Gamma\left(N_i \to \psi_{Lj}\eta^*\right) - \Gamma\left(N_i \to \psi^c_{Lj}\eta\right)\right]}{\Gamma_i} \\ &= \varepsilon_i^N + \varepsilon_i^\Delta ,\end{aligned} \quad (47)$$

where

$$\begin{aligned}\Gamma_i &\equiv \sum_j \left[\Gamma\left(N_i \to \psi_{Lj}\eta^*\right) + \Gamma\left(N_i \to \psi^c_{Lj}\eta\right)\right] \\ &= \frac{1}{8\pi} \left(y^\dagger y\right)_{ii} M_i \end{aligned} \quad (48)$$

is the total decay width of N_i, while

$$\begin{aligned}\varepsilon_i^N &= \frac{1}{8\pi} \frac{1}{(y^\dagger y)_{ii}} \sum_{k\neq i} \mathrm{Im}\left[\left(y^\dagger y\right)^2_{ik}\right] \\ &\times \sqrt{\frac{a_k}{a_i}} \left[1 - \left(1 + \frac{a_k}{a_i}\right) \ln\left(1 + \frac{a_i}{a_k}\right) + \frac{a_i}{a_i - a_k}\right] ,\end{aligned} \quad (49)$$

$$\begin{aligned}\varepsilon_i^\Delta &= \frac{3}{2\pi} \frac{1}{(y^\dagger y)_{ii}} \sum_{jm} \mathrm{Im}\left(f^\dagger_{jm} y^\dagger_{ij} y^\dagger_{im}\right) \frac{\tilde{\mu}}{M_i} \\ &\times \left[1 - \frac{a_\Delta}{a_i} \ln\left(1 + \frac{a_i}{a_\Delta}\right)\right] \end{aligned} \quad (50)$$

are the contributions of the first two one-loop diagrams and the third one, respectively. Here the definitions

$$a_i \equiv \frac{M_i^2}{M_1^2} , \quad a_\Delta \equiv \frac{M_\Delta^2}{M_1^2} \quad (51)$$

have been adopted.

Furthermore, as shown in Fig. 2, in the decay of Δ_L, the tree-level diagram interferes with the one-loop correction to generate the CP asymmetry,

$$\begin{aligned}\varepsilon_\Delta &\equiv 2 \frac{\sum_{ij} \left[\Gamma\left(\Delta_L^* \to \psi_{Li}\psi_{Lj}\right) - \Gamma\left(\Delta_L \to \psi^c_{Li}\psi^c_{Lj}\right)\right]}{\Gamma_\Delta} \\ &= \frac{2}{\pi} \frac{\sum_{ijk} \left(y_{ki} y_{kj} f_{ij}\right) \tilde{\mu} M_k \ln\left(1 + M_\Delta^2/M_k^2\right)}{\mathrm{Tr}\left(f^\dagger f\right) M_\Delta^2 + 4\tilde{\mu}^2 + 4\mu^2} \end{aligned} \quad (52)$$

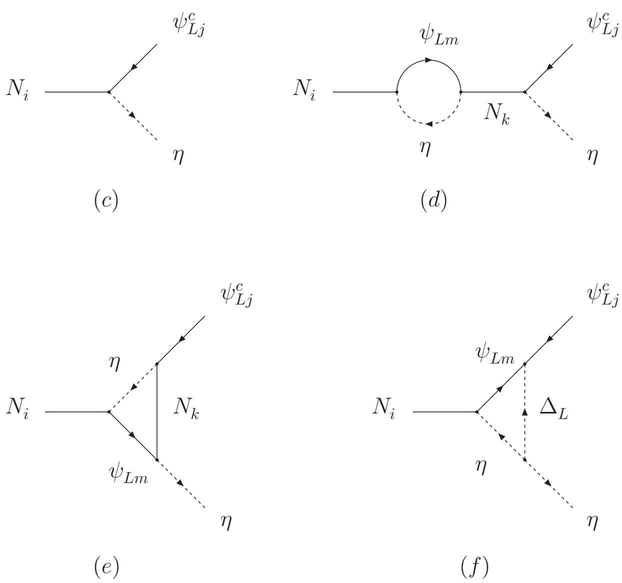

FIG. 2: The heavy Majorana neutrinos decay at one-loop order.

with

$$\begin{aligned}
\Gamma_\Delta &\equiv \sum_{ij}\Gamma\left(\Delta_L\to\psi^c_{Li}\psi^c_{Lj}\right)\\
&+ \Gamma\left(\Delta_L\to\eta\eta\right)+\Gamma\left(\Delta_L\to\phi\phi\right)\\
&\equiv \sum_{ij}\Gamma\left(\Delta^*_L\to\psi_{Li}\psi_{Lj}\right)\\
&+ \Gamma\left(\Delta^*_L\to\eta^*\eta^*\right)+\Gamma\left(\Delta^*_L\to\phi^*\phi^*\right)\\
&= \frac{1}{8\pi}\left[\frac{1}{4}\mathrm{Tr}\left(f^\dagger f\right)+\frac{\tilde{\mu}^2+\mu^2}{M^2_\Delta}\right]M_\Delta \qquad (53)
\end{aligned}$$

being the total decay width of Δ_L or Δ^*_L.

Note that we have not considered the cases where σ directly decay to produce the leptons and anti-leptons through the imaginary N_i or Δ_L if $m_\sigma > 2M_i$, $M_\Delta + 2m_\eta$ with m_σ and m_η being the masses of σ and η, respectively. For simplicity, here we will not discuss these cases.

It is straightforward to see that ε_Δ and ε^Δ_i will both be zero for $\kappa = 0$ and then $\tilde{\mu} = 0$. In the following, to illustrate how to realize non-resonant TeV leptogenesis, we first focus on the simple case where ε^N_i is the unique source of the CP asymmetry. Note that $\tilde{\mu}' = 0$ for $\kappa = 0$, accordingly, the one-loop diagram of Fig. 1 is absent and N_i have no possibility for the neutrino masses, we thus

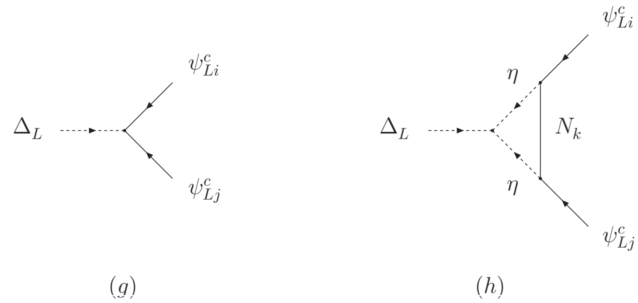

FIG. 3: The Higgs triplets decay to the leptons at one-loop order.

obtain

$$\begin{aligned}
\varepsilon^N_1 &\simeq -\frac{3}{16\pi}\sum_{k=2,3}\frac{\mathrm{Im}\left[\left(y^\dagger y\right)^2_{1k}\right]}{\left(y^\dagger y\right)_{11}}\frac{M_1}{M_k}\\
&\lesssim -\frac{3}{16\pi}\left(\frac{M_1}{M_2}+\frac{M_1}{M_3}\right)\sin\delta \qquad (54)\\
&\qquad (55)
\end{aligned}$$

with δ being the CP phase. Here we have assumed N_1 to be the lightest heavy Majorana neutrinos, i.e., $M_1^2 \ll M^2_{2,3}$, M^2_Δ. The final baryon asymmetry can be given by approximate relation [22]

$$Y_B \equiv \frac{n_B}{s} \simeq -\frac{28}{79}\times\begin{cases}\frac{\varepsilon_1}{g_*}\,, & (\text{for } K\ll 1)\,,\\ \frac{0.3\,\varepsilon_1}{g_* K(\ln K)^{0.6}}\,, & (\text{for } K\gg 1)\,,\end{cases} \qquad (56)$$

where the factor 28/79 is the value of $B/(B-L)$ and the parameter K is a measure of the departure from equilibrium and is defined by

$$K \equiv \left.\frac{\Gamma_1}{H(T)}\right|_{T=M_1} = \left(y^\dagger y\right)_{11}\left(\frac{45}{2^6\pi^5 g_*}\right)^{\frac{1}{2}}\frac{M_{\mathrm{Pl}}}{M_1}\,. \qquad (57)$$

Here $H(T) = (4\pi^3 g_*/45)^{\frac{1}{2}}T^2/M_{\mathrm{Pl}}$ is the Hubble constant with the Planck mass $M_{\mathrm{Pl}} \sim 10^{19}\,\mathrm{GeV}$ and the relativistic degrees of freedom $g_* \sim 100$. For example, inspecting $M_\Delta = 10\,\mathrm{TeV}$, $|\mu| = 1\,\mathrm{GeV}$ and $f \sim 10^{-6}$ to (43), we obtain $m_\nu \sim \mathcal{O}(0.1\,\mathrm{eV})$ which is consistent with the neutrino oscillation experiments. Furthermore, let $M_1 = 0.1\,M_{2,3} = 1\,\mathrm{TeV}$, $y \sim 10^{-6}$ and $\sin\delta = 10^{-3}$, we drive the sample predictions: $K \simeq 48$ and $\varepsilon_1 \simeq -1.2\times 10^{-5}$. In consequence,we arrive at $n_B/s \simeq 10^{-10}$ as desired.

For $\kappa \neq 0$ and then $\tilde{\mu}', \tilde{\mu} \neq 0$, Δ_L and N_i will both contribute to the neutrino masses and the lepton asymmetry. In the limit of $M_\Delta \ll M_i$, the final lepton asymmetry is expected to mostly produce by the decay of Δ_L. However, because the electroweak gauge scattering should be out of thermal equilibrium, it is difficult for a successful leptogenesis to lower the mass of Δ_L at TeV scale. Let us consider another possibility that N_i are much lighter than Δ_L. In this case, leptogenesis will be dominated by the decay of N_i. For $M_1^2 \ll M_{2,3}^2$, M_Δ^2 and $|\tilde{\mu}'\mu| \ll M_\Delta^2$, ε_1^N and ε_1^Δ can be simplified as [10]

$$\begin{aligned}\varepsilon_1^N &\simeq -\frac{3}{16\pi}\sum_{k=2,3}\frac{\mathrm{Im}\left[\left(y^\dagger y\right)_{1k}^2\right]}{(y^\dagger y)_{11}}\frac{M_1}{M_k}\\ &\simeq -\frac{3}{8\pi}\frac{M_1}{v^2}\sum_{jk}\frac{\mathrm{Im}\left[(\widetilde{m}_\nu^{I*})_{jk}y_{1j}^\dagger y_{1k}^\dagger\right]}{(y^\dagger y)_{11}}\frac{1}{\xi}\\ &\simeq -\frac{3}{8\pi}\frac{M_1\widetilde{m}_{\max}^I}{v^2}\frac{1}{\xi}\sin\delta'\,, \qquad (58)\\ \varepsilon_1^\Delta &\simeq -\frac{3}{8\pi}\frac{M_1}{v^2}\frac{\tilde{\mu}}{\mu}\sum_{jk}\frac{\mathrm{Im}\left[(m_\nu^{II*})_{jk}y_{1j}^\dagger y_{1k}^\dagger\right]}{(y^\dagger y)_{11}}\\ &\simeq -\frac{3}{8\pi}\frac{M_1 m_{\max}^{II}}{v^2}\left|\frac{\tilde{\mu}}{\mu}\right|\sin\delta''\,, \qquad (59)\end{aligned}$$

where δ' and δ'' are CP phases, $m_{\max}^{II}$ and $\widetilde{m}_{\max}^I$ are the maximal eigenstates of the neutrino mass matrixes (42) and (45), respectively. Inputting $y \sim 10^{-7}$, $M_1 = 1\,\mathrm{TeV}$ and $M_{2,3} = 10\,\mathrm{TeV}$, we obtain $\widetilde{m}_{\max}^I = \mathcal{O}(10^{-3}\,\mathrm{eV})$. Similarly, $m_{\max}^{II} = \mathcal{O}(0.1\,\mathrm{eV})$ for $M_\Delta = 10\,\mathrm{TeV}$, $|\mu| = 1\,\mathrm{GeV}$ and $f \sim 10^{-6}$. Under this setup, we deduce $\xi \simeq 10^{-3}$ by substituting $\overline{m}_\eta = 70\,\mathrm{GeV}$, $|\tilde{\mu}'| = 10^3\,\mathrm{TeV}$ into (46) and then have $\varepsilon_1^N \simeq -2\times 10^{-12}$ with the maximum CP phase. We also acquire $\varepsilon_1^\Delta \simeq -3\times 10^{-8}$ for $|\tilde{\mu}| \simeq |\tilde{\mu}'|$ and $\sin\delta'' = 0.15$. We thus drive the sample predictions: $K \simeq 0.5$ and $\varepsilon_1 \simeq \varepsilon_1^\Delta \simeq -3\times 10^{-8}$. In consequence, we arrive at $n_B/s \simeq 10^{-10}$ consistent with the cosmological observations.

Dark matter and Higgs phenomenology: Since the new Higgs doublet can not decay into the standard model particles, the neutral η_R^0 and η_I^0 can provide the attractive candidates for dark matter [13, 14, 15]. In particular, to realize dark matter, η_R^0 and η_I^0 should have the mass spectrum [14]:

$$\begin{aligned}\Delta m &\simeq (8-\ \ 9)\,\mathrm{GeV} \quad \text{for} \quad m_L = (60-73)\,\mathrm{GeV}\,, \qquad (60)\\ \Delta m &\simeq (9-12)\,\mathrm{GeV} \quad \text{for} \quad m_L = (73-75)\,\mathrm{GeV}\,. \qquad (61)\end{aligned}$$

Here $\Delta m \equiv m_{NL} - m_L$ with m_L and m_{NL} being the lightest and the next lightest masses between η_R^0 and η_I^0. Note

$$\begin{aligned}\overline{m}_\eta &\equiv \frac{1}{2}\left(m_L + m_{NL}\right)\,, \qquad (62)\\ |\delta m_\eta^2| &\equiv \frac{1}{2}\left(m_{NL}^2 - m_L^2\right)\,, \qquad (63)\end{aligned}$$

we thus deduce,

$$\begin{aligned}m_L &= \overline{m}_\eta\left(1-\frac{1}{2}\frac{|\delta m_\eta^2|}{\overline{m}_\eta^2}\right)\,, \qquad (64)\\ \Delta m &= \frac{|\delta m_\eta^2|}{\overline{m}_\eta}\,. \qquad (65)\end{aligned}$$

In the previous discussions of TeV leptogenesis with $\kappa \neq 0$, we take $M_\Delta = 10\,\mathrm{TeV}$, $|\mu| = 1\,\mathrm{GeV}$, $|\tilde{\mu}| = 10^3\,\mathrm{TeV}$ and $\overline{m}_\eta = 70\,\mathrm{GeV}$. It is straightforward to see $|\delta m_\eta| \simeq 25\,\mathrm{GeV}$ from (33). Therefore, we obtain $m_L \simeq 66\,\mathrm{GeV}$ and $\Delta m \simeq 9\,\mathrm{GeV}$, which is consistent with the mass spectrum (60).

η_R^0 and η_I^0 are expected to be produced in pairs by the standard model gauge bosons $W^\pm$, Z or γ and hence can be verified at the LHC. Once produced, $\eta^\pm$ will decay into $\eta_{R,I}^0$ and a virtual $W^\pm$, which becomes a quark-antiquark or lepton-antilepton pair. For example, if η_R^0 is lighter than η_I^0, the decay chain

$$\eta^+ \to \eta_I^0 l^+\nu\,, \quad \text{then} \quad \eta_I^0 \to \eta_R^0 l^+ l^- \qquad (66)$$

has 3 charged leptons and large missing energy, and can be compared to the direct decay

$$\eta^+ \to \eta_R^0 l^+\nu \qquad (67)$$

to extract the masses of the respective particles.

As for the phenomenology of the Higgs triplet at the LHC as well as the ILC, it has been discussed in [23]. The same-sign dileptons will be the most dominating modes of the δ^{++}. Complementary measurements of $|f_{ij}|$ at the ILC by the process $e^+e^+(\mu^+\mu^-) \to l_i^- l_j^-$ would allow us to study the structure of the neutrino mass matrix in detail.

For $\langle\chi\rangle = \mathcal{O}(\mathrm{TeV})$, which is natural to give the TeV Majorana masses of the right-handed neutrinos and then realize the TeV leptogenesis, the mixing angle ϑ and the splitting between $h_{1,2}$ may be large. Furthermore, the couplings of $h_{1,2}$ to W and Z bosons, quarks and charged leptons have essentially the same structure as the corresponding Higgs couplings in the standard model, however, their size is reduced by $\cos\vartheta$ and $\sin\vartheta$, respectively. In the extreme case $\vartheta = \frac{\pi}{2}$, the couplings of the lighter physical boson h_1 to quarks and leptons would even vanish. In other words, this mixing could lead to significant impact on the Higgs searches at the LHC [24, 25].

Summary: We propose a new model to realize leptogenesis and dark matter at the TeV scale. A real scalar is introduced to naturally realize the Majorana masses of the right-handed neutrinos. Furthermore, we also consider a new Higgs doublet to provide the attractive candidates for dark matter. Since the right-handed neutrinos have no responsibility to generate the neutrino masses, which is mostly dominated by the Higgs triplet through the type-II seesaw, they can have large CP asymmetry at a low scale, such as TeV, to produce the observed matter-antimatter asymmetry in the universe, even if their Majorana masses are not highly quasi-degenerate. It should be noticed that all new particles are close to the TeV scale and hence should be observable at the LHC or the ILC.

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