Leptogenesis with TeV Scale W_R

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Successful leptogenesis within the conventional TeV-scale left-right implementation of type-I seesaw has been shown to require that the mass of the right-handed $W_R^\pm$ boson should have a lower bound much above the reach of the Large Hadron Collider. This bound arises from the necessity to suppress the washout of lepton asymmetry due to $W_R^\pm$ -mediated $\Delta L \neq 0$ processes. We show that in an alternative quark seesaw realization of left-right symmetry, the above bound can be avoided. Lepton asymmetry in this model is generated not via the usual right-handed neutrino decay but rather via the decay of new heavy scalars producing an asymmetry in the $B - L$ carrying Higgs triplets responsible for type-II seesaw, whose decay leads to the lepton asymmetry.

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I. INTRODUCTION

The discoveries on atmospheric, solar, accelerator, and reactor neutrino oscillations have established that the neutrinos are massive and are mixed among different flavors [1]. This is the first evidence for new physics beyond the standard $SU(3)_c \times SU(2)_L \times U(1)_Y$ model (SM). Meanwhile, the cosmological observations combined with oscillation data imply that the neutrino masses should be in a sub-eV range [1]. A simple beyond the standard model mechanism for understanding such tiny neutrino masses is provided by the various seesaw schemes [2–5], known as type-I, type-II, inverse, and type-III seesaw, respectively. In the seesaw scenarios, the SM is extended by the inclusion of heavy SM singlet fermions, i.e., the right-handed neutrinos (RHNs) with Majorana masses (for type-I and inverse) or heavy triplet scalars with lepton-number-violating couplings (type-II case). After the $SU(2)_L \times U(1)_Y$ electroweak symmetry is spontaneously broken down to the $U(1)_{em}$, the neutrinos acquire a Majorana mass term highly suppressed by a small ratio of the electroweak scale over the heavy Majorana fermion or scalar masses. These models have the extra advantage that they provide a way to understand the cosmic matter-antimatter asymmetry via the so-called

leptogenesis mechanism [6]. In the type-I seesaw case, the typical scenario involves the decay of heavy RHNs of the seesaw to generate a lepton asymmetry through their Yukawa interactions as long as the CP is not conserved [6]. Subsequently, this lepton asymmetry can be partially converted to a baryon asymmetry through the sphaleron processes [7].

Left-right (LR) symmetric models based on the gauge group $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [8] provide an elegant embedding of the seesaw mechanism into a more fundamental framework. It naturally provides the two fundamental ingredients of the type-I seesaw mechanism, i.e., the RHNs that arise from anomaly cancellation and the Majorana masses for RHNs that arise from the breaking of the $SU(2)_R \times U(1)_{B-L}$ symmetry. In the LR symmetric models, the SM left-handed fermions are placed in the $SU(2)_L$ doublets as they are in the SM; the SM right-handed fermions plus three RHNs are placed in the $SU(2)_R$ doublets. Implementation of type-I seesaw in the model comes from the Higgs scalars consisting of $SU(2)_{L,R}$ triplets and one $SU(2)_L \times SU(2)_R$ bidoublet. When the $[SU(2)_R]$ -triplet Higgs scalar develops its vacuum expectation value (VEV) breaking the LR symmetry down to the SM, i.e., $SU(2)_R \times U(1)_{B-L} \rightarrow U(1)_Y$ and the lepton number violation needed for seesaw emerges automatically. After the RHNs go out of equilibrium in the early Universe, their decays then can provide a realization of leptogenesis. Because the RHNs couple to the non-SM gauge bosons, their departure from equilibrium as well as the suppression of subsequent washout constrains the LR symmetry breaking [9,10], implying $M_{W_R} \geq 10\text{--}13$ TeV depending on assumptions. This implies that it would be impossible to

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simultaneously have a successful leptogenesis and LR symmetry with type-I seesaw testable at the LHC.

In addition to the neutrino mass and the baryon asymmetry, the SM faces other challenges such as the strong CP problem. It is an interesting property of the LR symmetric models that they can solve the strong CP problem without the need for an axion. This is because these models naturally contain a discrete parity symmetry that makes the tree-level quark mass matrices Hermitian, implying the vanishing of the strong CP phase at the tree level [11]. This leads to a calculable value for θ from loop effects. However in the minimal LR type-I seesaw model, additional symmetries (e.g., supersymmetry [12]) are required for this purpose. A realization of left-right symmetry, where a finite θ is generated at the two-loop level without the need for any additional symmetries [13] and is therefore naturally below the current limits, is provided by adding heavy vectorlike quarks and leptons to the model. The quark masses arise in this model from a seesaw mechanism (called quark seesaw [14] in this paper) using the vectorlike quark masses as the heavy end of the seesaw. The model has a different Higgs structure with only doublets, one for the left and another for the right $SU(2)$ instead of bidoublets as in the minimal LR type-I seesaw case.

Our goal in this paper is to focus on this class of LR symmetric models and discuss leptogenesis and neutrino masses. We extend the minimal quark seesaw LR model by adding $B - L = 2$ Higgs triplets to generate small neutrino masses using type-II seesaw mechanism. The required smallness of the triplet VEV needed for this purpose arises from a two-step suppression of the VEVs [15] with no couplings smaller than $\sim 10^{-5}$ (roughly the same order as the electron Yukawa in the SM).

We then discuss leptogenesis in this model and show that the lower bound on the W_R mass mentioned above for the minimal type-I LR model can be avoided in this case. Our strategy is to realize that in the quark seesaw realization of the left-right model, after the full gauge symmetry is broken, there is still a global lepton number symmetry left exact. We use this symmetry to generate lepton asymmetry of the Universe. The source of lepton asymmetry, however, is not the decay of RHNs but rather the decay of new heavy SM singlet scalar bosons, as for instance in [15] with nonzero lepton number that couples two sets of lepton number carrying scalars, one set coupling to leptons and the other to fields that do not couple to leptons. The former acquire vacuum expectation values at very low scale to generate neutrino masses via type-II seesaw mechanism. This heavy scalar decay leads to an asymmetry in $B - L$ carrying scalar triplets responsible for type-II seesaw as well as the others. The scalar triplets decay to leptons generating the lepton asymmetry that gets converted to the baryon asymmetry via the sphaleron interactions. This happens prior to their acquiring VEVs. The significance of this result is that if the W_R is discovered

at the LHC 14, it will not only provide evidence for the idea of LR extension of SM but also point towards possibly the quark seesaw realization of the idea as well as providing possibly an alternative solution to the strong CP problem. This has significant implications for physics beyond the SM, e.g., any possible grand unification etc. The model also leads to keV sterile neutrinos, with implications for low energy beta decay phenomenology [16].

The paper is organized as follows: in Sec. II, we review the quark seesaw realization of the left-right model with new Higgs scalar triplets for neutrino masses and the different symmetry breaking stages; in Sec. III, we discuss how small neutrino masses arise in this model; in Sec. IV, we give benchmark values for parameters of the model; in Sec. V, we present the calculation of the lepton asymmetry generation, in Sec. VI we discuss how this lepton asymmetry gets converted to the baryon symmetry, and in Sec. VII we summarize our results.

II. THE MODEL

Our model includes the chiral fermions and Higgs scalars as given below,

$$\begin{aligned}
 q_L \left(3, 2, 1, +\frac{1}{3} \right) \left(+\frac{1}{3} \right) &= \begin{bmatrix} u_L \\ d_L \end{bmatrix}, \\
 q_R \left(3, 1, 2, +\frac{1}{3} \right) \left(+\frac{1}{3} \right) &= \begin{bmatrix} u_R \\ d_R \end{bmatrix}, \\
 D_{L,R} \left(3, 1, 1, -\frac{2}{3} \right) \left(+\frac{1}{3} \right), & \quad U_{L,R} \left(3, 1, 1, +\frac{4}{3} \right) \left(+\frac{1}{3} \right); \\
 l_L (1, 2, 1, -1) (-1) &= \begin{bmatrix} \nu_L \\ e_L \end{bmatrix}, \\
 l_R (1, 1, 2, -1) (-1) &= \begin{bmatrix} \nu_R \\ e_R \end{bmatrix}, \\
 E_{L,R} (1, 1, 1, -2) (-1); \\
 \phi_L (1, 2, 1, -1) (0) &= \begin{bmatrix} \phi_L^0 \\ \phi_L^- \end{bmatrix}, \\
 \phi_R (1, 1, 2, -1) (0) &= \begin{bmatrix} \phi_R^0 \\ \phi_R^- \end{bmatrix}, \\
 \Delta_L (1, 3, 1, +2) (+2) &= \begin{bmatrix} \delta_L^+/\sqrt{2} & \delta_L^{++} \\ \delta_L^0 & \delta_L^+/\sqrt{2} \end{bmatrix}, \\
 \Delta_R (1, 1, 3, +2) (+2) &= \begin{bmatrix} \delta_R^+/\sqrt{2} & \delta_R^{++} \\ \delta_R^0 & \delta_R^+/\sqrt{2} \end{bmatrix}; \\
 \sigma (1, 1, 1, 0) (-2), & \quad \chi (1, 1, 1, 0) (+1). \tag{1}
 \end{aligned}$$

Here the first brackets following the fields describe the transformations under the $G_{LR} = SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_X$ gauge group, while the second brackets

are the global $B - L$ numbers.¹ For brevity, we do not write down the full Lagrangian. Instead we only show the following terms relevant for our discussion.

$$\begin{aligned} \mathcal{L} \supset & -y_D(\bar{q}_L\tilde{\phi}_L D_R + \bar{q}_R\tilde{\phi}_R D_L) - M_D\bar{D}_R D_L \\ & -y_U(\bar{q}_L\phi_L U_R + \bar{q}_R\phi_R U_L) - M_U\bar{U}_R U_L \\ & -y_E(\bar{l}_L\tilde{\phi}_L E_R + \bar{l}_R\tilde{\phi}_R E_L) - M_E\bar{E}_R E_L \\ & -\frac{1}{2}f(\bar{l}_L^c i\tau_2 \Delta_L l_L + \bar{l}_R^c i\tau_2 \Delta_R l_R) \\ & -\kappa\sigma(\phi_L^T i\tau_2 \Delta_L \phi_L + \phi_R^T i\tau_2 \Delta_R \phi_R) - \mu\sigma\chi^2 \\ & + \text{H.c.} + M_{\Delta_L}^2 \text{Tr}(\Delta_L^\dagger \Delta_L) + M_{\Delta_R}^2 \text{Tr}(\Delta_R^\dagger \Delta_R) + M_\sigma^2 \sigma^\dagger \sigma. \end{aligned} \quad (2)$$

The $[SU(2)]$ -doublet Higgs scalars $\phi_{L,R}$ are responsible for the left-right and electroweak symmetry breaking, i.e.,

$$\begin{aligned} & SU(2)_L \times SU(2)_R \times U(1)_X \\ & \xrightarrow{\langle\phi_R\rangle} SU(2)_L \times U(1)_Y \xrightarrow{\langle\phi_L\rangle} U(1)_{em}. \end{aligned} \quad (3)$$

This can be achieved by softly breaking the parity symmetry in Eq. (2) in the Higgs doublet masses. This keeps the strong CP solution unaffected.

Note that at this stage the gauge symmetry is fully broken down to $SU(3)_c \times U(1)_{em}$ as desired but it leaves a global $B - L$ symmetry unbroken. Thus the neutrinos only have a tiny Dirac mass from two-loop diagram [17]. When the $[G_{LR}]$ -singlet Higgs scalar χ develops its VEV $\langle\chi\rangle$ it breaks the global $U(1)_{B-L}$ symmetry. We assume the scale of this breaking to be ~ 100 TeV. The corresponding massless Goldstone boson decouples from the thermal bath at a high temperature and does not affect the subsequent evolution of the Universe. The other $[G_{LR}]$ -singlet Higgs scalars σ then acquire a seesaw-suppressed VEV,

$$\langle\sigma\rangle \simeq -\frac{\mu\langle\chi\rangle^2}{M_\sigma^2} \ll \langle\chi\rangle \quad \text{for } M_\sigma^2 \gg \mu\langle\chi\rangle. \quad (4)$$

As for the $[SU(2)]$ -triplet Higgs scalars $\Delta_{L,R}$, their VEVs arise from the σ VEV and are therefore highly suppressed by the small VEVs $\langle\sigma\rangle$ as well as by the small value of the coupling parameter κ that is required for generation of matter-antimatter asymmetry (see later), i.e.,

$$\langle\Delta_{L,R}\rangle \simeq -\frac{\kappa\langle\sigma\rangle\langle\phi_{L,R}\rangle^2}{M_{\Delta_{L,R}}^2} \ll \langle\phi_{L,R}\rangle. \quad (5)$$

This for example leads to a VEV of Δ_L in the eV range, which explains the tiny neutrino masses.

¹The present $U(1)_{B-L}$ global symmetry can be a gauged symmetry. But we do not discuss it here.

We emphasize that the phase in the cubic coupling $\kappa\langle\sigma\rangle$ can contribute an overall phase to the mass matrix M_E by redefining the fields (Δ_L, l_L, E_R) and (Δ_R, l_R, E_L) . This overall phase cannot be removed because the parity symmetry now is softly broken but its effect on the strong CP phase arises in high loop level and the model still solves the strong CP problem.

III. FERMION MASSES

After the left-right and electroweak symmetry breaking, we obtain the charged fermion masses,

$$\begin{aligned} \mathcal{L} \supset & -[\bar{u}_L \quad \bar{U}_L] \begin{bmatrix} 0 & y_U\langle\phi_L\rangle \\ y_U^\dagger\langle\phi_R\rangle & M_U^\dagger \end{bmatrix} \begin{bmatrix} u_R \\ U_R \end{bmatrix} \\ & -[\bar{d}_L \quad \bar{D}_L] \begin{bmatrix} 0 & y_D\langle\phi_L\rangle \\ y_D^\dagger\langle\phi_R\rangle & M_D^\dagger \end{bmatrix} \begin{bmatrix} d_R \\ D_R \end{bmatrix} \\ & -[\bar{e}_L \quad \bar{E}_L] \begin{bmatrix} 0 & y_E\langle\phi_L\rangle \\ y_E^\dagger\langle\phi_R\rangle & M_E^\dagger \end{bmatrix} \begin{bmatrix} e_R \\ E_R \end{bmatrix} + \text{H.c.} \end{aligned} \quad (6)$$

It is well known the above quark mass matrices can solve the strong CP problem without introducing the axion. The details can be found in [13]. Furthermore, since the $[SU(2)]$ -singlet fermions are heavy, they can be integrated out. The SM fermion masses thus should be

$$\begin{aligned} m_d &= -y_D \frac{\langle\phi_L\rangle\langle\phi_R\rangle}{M_D} y_D^\dagger, \\ m_u &= -y_U \frac{\langle\phi_L\rangle\langle\phi_R\rangle}{M_U} y_U^\dagger, \\ m_e &= -y_E \frac{\langle\phi_L\rangle\langle\phi_R\rangle}{M_E} y_E^\dagger. \end{aligned} \quad (7)$$

We now study the neutrino mass. Clearly, through the Yukawa couplings of the $[SU(2)]$ -triplet Higgs scalars $\Delta_{L,R}$ to the $[SU(2)]$ -doublet leptons $l_{L,R}$, the left- and right-handed neutrinos can obtain their Majorana masses,

$$\begin{aligned} \mathcal{L} \supset & -\frac{1}{2}m_L\bar{\nu}_L^c\nu_L - \frac{1}{2}m_R\bar{\nu}_R^c\nu_R + \text{H.c.} \\ & \text{with } m_{L,R} = f\langle\Delta_{L,R}\rangle. \end{aligned} \quad (8)$$

Thanks to the small VEV $\langle\Delta_L\rangle$, the left-handed neutrino masses m_L can be made tiny even if the Yukawa couplings f are sizable. This is the essence of the type-II seesaw. As shown in Eq. (4), the present model has a first stage seesaw suppression of the parameter $\kappa\langle\sigma\rangle$ in a natural way. Therefore, this two-step seesaw scheme can be called a double type-II seesaw mechanism [15]. Clearly, the double type-II seesaw mechanism now is also applied to the generation of the right-handed neutrino masses m_R . As a result, the right-handed neutrino masses are much below the left-right symmetry breaking scale and for our benchmark choice of parameters (Sec. VI) are in few keV range.

As noted, at the two-loop level, the charged current interactions and the charged fermion masses can always mediate a Dirac mass term between the left- and right-handed neutrinos [17],²

$$\mathcal{L} \supset -m_D \bar{\nu}_L \nu_R + \text{H.c.} \quad \text{with} \quad m_D \simeq \frac{3g^4}{4(16\pi^2)^2} \frac{m_t m_b}{M_{W_R}^2} \hat{m}_e. \quad (9)$$

The $\nu_L - \nu_R$ mixing is small enough to escape from the big bang nucleosynthesis (BBN) constraint [18].

IV. NUMERICAL EXAMPLE

In this section, we give a benchmark scenario to show how the model works for right-handed W_R mass in the low TeV range so that it is accessible at the LHC. For example, we take

$$\langle \phi_R \rangle = 5 \text{ TeV} \Rightarrow M_{W_R} = \frac{1}{\sqrt{2}} g \langle \phi_R \rangle \simeq 2.3 \text{ TeV}. \quad (10)$$

In addition, we assume the global lepton number breaking scale to be³

$$\langle \chi \rangle = 100 \text{ TeV}. \quad (11)$$

By inserting

$$\begin{aligned} M_{\Delta_L} &= 3 \text{ TeV}, & M_{\Delta_R} &= 1 \text{ TeV}, & M_{\sigma_{1,2}} &= 333 \text{ TeV}, \\ \mu_{1,2} &= 10^{-6} M_{\sigma_{1,2}}, & \kappa_{1,2} &= 10^{-5}, \end{aligned} \quad (12)$$

in Eqs. (4) and (5), we then read

$$\langle \sigma \rangle \simeq 10 \text{ MeV}, \quad \langle \Delta_R \rangle \simeq 2.5 \text{ keV}, \quad \langle \Delta_L \rangle \simeq 0.3 \text{ eV}. \quad (13)$$

Eventually, the neutrino masses should be

$$\begin{aligned} m_L &= 0.3 \text{ eV} \cdot f_\Delta, & m_R &= 2.5 \text{ keV} \cdot f_\Delta, \\ m_D &\simeq 1.3 \text{ eV}. \end{aligned} \quad (14)$$

Note, the right-handed neutrinos ν_R can be quasidegenerate even if the left-handed neutrinos ν_L have a hierarchical spectrum. This is because the parity symmetry is softly

broken and hence the right-handed Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix is allowed to be different from the left-handed PMNS matrix. Accordingly, the three ν_R can be all at the keV scale while the $\nu_L - \nu_R$ mixing $\theta_{LR} \sim m_D/m_R$ can be very small when the Yukawa couplings f_Δ are sizable, i.e.,

$$\theta_{LR} \sim m_D/m_R \lesssim 10^{-3} \quad \text{for} \quad f_\Delta = 0.5. \quad (15)$$

We now show that with our choice of parameters, the keV right-handed neutrinos do affect the BBN discussion. First, the right-handed neutrinos decouple from the weak interactions at the temperature around $T \sim 3 \text{ MeV} \theta_{LR}^{-2/3} \simeq 300 \text{ MeV} (\theta_{LR}/10^{-3})^{-2/3}$ where 3 MeV is the decoupling temperature for the left-handed neutrinos. Secondly, the annihilating and scattering processes mediated by the other gauge bosons decouple at the temperature around $T \sim 3 \text{ MeV} (M_{W_R}/M_{W_L})^{4/3} \simeq 260 \text{ MeV} (M_{W_R}/2.3 \text{ TeV})^{4/3}$. Thirdly, we consider the annihilating and scattering processes mediated by the singly charged component of the Δ_R scalar. The decoupling temperature should be about $T \sim 3 \text{ MeV} (g/f_\Delta)^{4/3} (M_{\Delta_R}/M_{W_L})^{4/3} \simeq 120 \text{ MeV} (0.5/f_\Delta)^{4/3} (M_{\Delta_R}/1 \text{ TeV})^{4/3}$. In conclusion, the existence of the keV right-handed neutrinos does not conflict with the BBN. As already noted, the $\nu_L - \nu_R$ mixing is small enough so that oscillations do not affect the BBN considerations [18].

V. LEPTON ASYMMETRY

We now discuss the generation of lepton asymmetry in our model. As noted, it is different from the mechanism in type-I seesaw models where decay of the right-handed neutrinos was the main source. In our model, it is the decay of heavy singlet scalar σ out of thermal equilibrium with the cosmic plasma when $T \sim M_\sigma$ for our choice of parameters and its decay plays this role. Its decay takes place at temperatures above the $\langle \chi \rangle$ VEV so that χ and χ^* are still different particles. We also assume the couplings involving σ fields and other scalars of type $\sigma^* \sigma \phi^\dagger \phi$ etc. are $\sim 10^{-6}$ so that these scattering processes do not keep σ fields in equilibrium at $T \sim M_\sigma$. Understanding the smallness of these parameters is beyond the scope of this paper.

Before the left-right symmetry breaking, the $[G_{LR}]$ -singlet scalars σ_a can have three decay modes, $\sigma_a \rightarrow \phi_L^* \phi_L^* \Delta_L^*$, $\sigma_a \rightarrow \phi_R^* \phi_R^* \Delta_R^*$, and $\sigma_a \rightarrow \chi^* \chi^*$. We calculate the decay width at tree level,

$$\begin{aligned} \Gamma_{\sigma_a} &= \Gamma(\sigma_a \rightarrow \phi_L^* \phi_L^* \Delta_L^*) + \Gamma(\sigma_a \rightarrow \phi_R^* \phi_R^* \Delta_R^*) \\ &\quad + \Gamma(\sigma_a \rightarrow \chi^* \chi^*) \\ &\equiv \Gamma(\sigma_a^* \rightarrow \phi_L \phi_L \Delta_L) + \Gamma(\sigma_a^* \rightarrow \phi_R \phi_R \Delta_R) \\ &\quad + \Gamma(\sigma_a^* \rightarrow \chi \chi), \\ &= \frac{1}{8\pi} \left(\frac{|\mu_a|^2}{M_{\sigma_a}^2} + \frac{3|\kappa_a|^2}{16\pi^2} \right) M_{\sigma_a}. \end{aligned} \quad (16)$$

²Because of some quartic couplings in the scalar potential, the singly charged components of the $[SU(2)]$ -triplet Higgs scalars Δ_L and Δ_R can mix with each other at tree level after the left-right and electroweak symmetry breaking. These singly charged scalars with the charged leptons can mediate a one-loop diagram to generate a Dirac mass term between the left- and right-handed neutrinos. This Dirac mass term could be much smaller than the two-loop contribution since the VEVs $\langle \Delta_{L,R} \rangle$ are assumed to be much smaller than the quark masses $m_{t,b}$.

³If the present $U(1)_{B-L}$ global symmetry is a gauged symmetry, we need to consider the departure from equilibrium of the related gauge interactions. Furthermore, we need to check other experimental constraints [19]. For this paper, we keep it as a global symmetry.

We also compute the one-loop CP asymmetries in the above decays, i.e.,

$$\begin{aligned}\varepsilon_{\sigma_a}^{\Delta_L} &= 2 \frac{\Gamma(\sigma_a \rightarrow \phi_L^* \phi_L^* \Delta_L^*) - \Gamma(\sigma_a^* \rightarrow \phi_L \phi_L \Delta_L)}{\Gamma_{\sigma_a}} \\ &= \frac{1}{4\pi} \frac{\text{Im}(\mu_b^* \mu_a \kappa_b \kappa_a^*)}{M_{\sigma_a}^2 - M_{\sigma_b}^2} \frac{1}{\frac{|\mu_a|^2}{M_{\sigma_a}^2} + \frac{3|\kappa_a|^2}{16\pi^2}}, \\ \varepsilon_{\sigma_a}^{\Delta_R} &= 2 \frac{\Gamma(\sigma_a \rightarrow \phi_R^* \phi_R^* \Delta_R^*) - \Gamma(\sigma_a^* \rightarrow \phi_R \phi_R \Delta_R)}{\Gamma_{\sigma_a}} \\ &= \frac{1}{4\pi} \frac{\text{Im}(\mu_b^* \mu_a \kappa_b \kappa_a^*)}{M_{\sigma_a}^2 - M_{\sigma_b}^2} \frac{1}{\frac{|\mu_a|^2}{M_{\sigma_a}^2} + \frac{3|\kappa_a|^2}{16\pi^2}}, \\ \varepsilon_{\sigma_a}^{\chi} &= \frac{\Gamma(\sigma_a \rightarrow \chi^* \chi^*) - \Gamma(\sigma_a^* \rightarrow \chi \chi)}{\Gamma_{\sigma_a}} \\ &= -\frac{1}{4\pi} \frac{\text{Im}(\mu_b^* \mu_a \kappa_b \kappa_a^*)}{M_{\sigma_a}^2 - M_{\sigma_b}^2} \frac{1}{\frac{|\mu_a|^2}{M_{\sigma_a}^2} + \frac{3|\kappa_a|^2}{16\pi^2}}.\end{aligned}\quad (17)$$

It is easy to check that

$$\varepsilon_{\sigma_a}^{\Delta_L} = \varepsilon_{\sigma_a}^{\Delta_R} \equiv \frac{1}{2} \varepsilon_{\sigma_a}^{\Delta}, \quad (18)$$

due to the parity symmetry, meanwhile,

$$\varepsilon_{\sigma_a}^{\Delta_L} + \varepsilon_{\sigma_a}^{\Delta_R} + \varepsilon_{\sigma_a}^{\chi} \equiv 0, \quad (19)$$

due to the lepton number conservation.

When the $[G_{LR}]$ -singlet scalars σ go out of equilibrium, their decays can produce a lepton asymmetry stored in the $[SU(2)]$ -triplet Higgs scalars $\Delta_{L,R}$ and an opposite lepton asymmetry stored in the $[G_{LR}]$ -singlet Higgs scalar χ . Clearly, the $\Delta_{L,R}$ asymmetry and the χ asymmetry decouple from each other once they are induced. This is clearly satisfied in our model since the communication between the χ field and the $\Delta_{L,R}$ field only occurs via the σ mixing with χ 's from μ coupling and κ coupling after $\langle \phi_R \rangle$ switches on. This mixing is given by the product of two small parameters μ and κ in our choice of parameters and is of the order of $\sim 10^{-11}$. Thus this mixing is not cosmologically relevant at the epoch of our leptogenesis. The $[SU(2)]$ -triplet Higgs scalars $\Delta_{L,R}$ can transfer their lepton asymmetry to the $[SU(2)]$ -doublet leptons $l_{L,R}$ when they decay and this happens before the sphaleron processes become inactive at temperature $T_{\text{sph}} \sim 100$ GeV.

The spontaneous breaking of the global $U(1)_{B-L}$ symmetry can happen after the electroweak symmetry breaking. In this case, no lepton-number-violating processes would emerge before the sphalerons had partially converted the lepton asymmetry to a baryon asymmetry. In the alternative case where the lepton number is spontaneously broken before the electroweak symmetry breaking, the cubic couplings $\kappa \langle \sigma \rangle$ can be highly suppressed by the small

VEVs $\langle \sigma \rangle$ and hence the lepton-number-violating processes $l_{L(R)} l_{L(R)} \leftrightarrow \phi_{L(R)} \phi_{L(R)}$, $l_{L(R)}^c l_{L(R)}^c \leftrightarrow \phi_{L(R)}^* \phi_{L(R)}^*$, $l_{L(R)} \phi_{L(R)}^* \leftrightarrow l_{L(R)}^c \phi_{L(R)}$, will not go into equilibrium before the sphalerons stop working, i.e.,

$$\Gamma_{\Delta L=2} < H(T) \quad \text{for } T > T_{\text{sph}}. \quad (20)$$

Here the Hubble constant $H(T)$ is given by

$$H(T) = \left(\frac{8\pi^3 g_*}{90} \right)^{\frac{1}{2}} \frac{T^2}{M_{\text{Pl}}}, \quad (21)$$

with $M_{\text{Pl}} \simeq 1.22 \times 10^{19}$ GeV being the Planck mass and $g_* = \mathcal{O}(100)$ being the relativistic degrees of freedom. As for the lepton-number-violating interaction rates $\Gamma_{\Delta L=2}$, they can be estimated by

$$\begin{aligned}\Gamma_{\Delta L=2} &\sim \frac{\text{Tr}(f^\dagger f) |\kappa_a \langle \sigma_a \rangle|^2}{T} \quad \text{for } T > M_{\Delta_{L,R}}, \\ \Gamma_{\Delta L=2} &\sim \frac{\text{Tr}(f^\dagger f) |\kappa_a \langle \sigma_a \rangle|^2 T^3}{M_{\Delta_{L,R}}^4} \quad \text{for } T < M_{\Delta_{L,R}}.\end{aligned}\quad (22)$$

Both these conditions are satisfied for our choice of Δ_L mass of 3 TeV.

We need to also check the right-handed neutrino-antineutrino oscillation. It is well known that this type of $\nu_R - \bar{\nu}_R$ oscillation should be helicity suppressed [20]. We can roughly estimate the rate by

$$\Gamma_{\nu_R - \bar{\nu}_R} \sim \frac{\text{Tr}(m_R^\dagger m_R)}{T} \quad \text{for } T \gg m_R. \quad (23)$$

If the right-handed charged current interactions,

$$\mathcal{L} \supset \left(\frac{M_{W_L}^2}{M_{W_R}^2} \right) \frac{G_F}{\sqrt{2}} \bar{u}_R \gamma^\mu d_R \bar{e}_R \gamma_\mu \nu_R + \text{H.c.}, \quad (24)$$

keep in equilibrium above the crucial temperature T_{sph} , the lepton-number-violating $\nu_R - \bar{\nu}_R$ oscillation should match the following condition,

$$\Gamma_{\nu_R - \bar{\nu}_R} < H(T) \quad \text{for } T > T_{\text{sph}}, \quad (25)$$

to avoid the elimination of the produced lepton asymmetry. It is easy to see that for keV right-handed neutrinos, the above condition is easily satisfied.

So we conclude that in any cases the $[SU(2)]$ -triplet Higgs scalars $\Delta_{L,R}$, which have become nonrelativistic above the temperature T_{sph} , can efficiently decay into the $[SU(2)]$ -doublet leptons $l_{L,R}$ and hence complete the expected lepton number conversion.

VI. BARYON ASYMMETRY

We can calculate the final baryon asymmetry in our model from the initial lepton asymmetry using the chemical potential identities [21]. For this purpose, we denote $\mu_{q,d,u,l,e,\nu,\phi}$ as the chemical potentials of the SM fermions $q_L(3,2,+1/6)$, $d_R(3,1,-1/3)$, $u_R(3,1,+2/3)$, $l_L(1,2,-1/2)$, $e_R(1,1,-1)$, the right-handed neutrinos $\nu_R(1,0,0)$, and the SM Higgs scalar $\phi(1,2,+1/2)$, respectively. Here the brackets following the fields describe the transformations under the $SU(3)_c \times SU(2)_L \times U(1)_Y$ gauge groups. Above the electroweak scale, the SM Yukawa interactions yield the following relations,

$$\begin{aligned} -\mu_q + \mu_\phi + \mu_d &= 0, & -\mu_q - \mu_\phi + \mu_u &= 0, \\ -\mu_l + \mu_\phi + \mu_e &= 0, \end{aligned} \quad (26)$$

whereas the $SU(2)_L$ sphalerons being in equilibrium lead to

$$3\mu_q + \mu_l = 0, \quad (27)$$

and the vanishing hypercharge in the Universe requires

$$3(\mu_q - \mu_d + 2\mu_u - \mu_l - \mu_e) + 2\mu_\phi = 0. \quad (28)$$

At this stage, the right-handed charged current interactions (24) may still keep in equilibrium, depending on the left-right symmetry breaking scale, and we assume that they are, in which case, we further have

$$-\mu_u + \mu_d - \mu_e + \mu_\nu = 0. \quad (29)$$

Solving for these equations, all chemical potentials can be expressed in terms of a single one (chosen here as μ_l). For example, we read

$$\begin{aligned} \mu_q &= -\frac{1}{3}\mu_l, & \mu_d &= -\frac{19}{21}\mu_l, & \mu_u &= \frac{5}{21}\mu_l, \\ \mu_e &= \frac{3}{7}\mu_l, & \mu_\nu &= \frac{11}{7}\mu_l, & \mu_\phi &= \frac{4}{7}\mu_l. \end{aligned} \quad (30)$$

The corresponding baryon and lepton asymmetries then should be

$$\begin{aligned} B &= 3(2\mu_q + \mu_u + \mu_d) = -4\mu_l, \\ L_{\text{SM}} &= 3(2\mu_l + \mu_e) = \frac{51}{7}\mu_l, \\ L_{\nu_R} &= 3\mu_\nu = \frac{33}{7}\mu_l. \end{aligned} \quad (31)$$

If the right-handed charged current interactions (24) have not decoupled before the $SU(2)_L$ sphaleron processes stop working, we can easily read the final baryon asymmetry from the initial lepton asymmetries,

$$B^f = \frac{1}{4}[B^i - (L_{\text{SM}}^i + L_{\nu_R}^i)] = -\frac{1}{4}(L_{\text{SM}}^i + L_{\nu_R}^i), \quad (32)$$

by inserting the relation (31) and assuming the initial baryon asymmetry to be 0, i.e., $B^i = 0$.

We now demonstrate the generation of the lepton and baryon asymmetries. For the masses and couplings in Eq. (12), we can expect a weak washout condition,

$$K = \frac{\Gamma_{\sigma_{1,2}}}{H(T)} \Big|_{T=M_{\sigma_{1,2}}} \simeq 0.2, \quad (33)$$

under which the final baryon asymmetry can well approximate to [22]

$$\frac{n_B}{s} = -\frac{1}{4} \frac{n_L}{s} \sim -\frac{1}{4} \frac{\epsilon_{\sigma_1}^\Delta + \epsilon_{\sigma_2}^\Delta}{g_*}. \quad (34)$$

Here n_B , n_L , and s respectively are the baryon number density, the lepton number density, and the entropy density, while $g_* = 209.5$ is the relativistic degrees of freedom during this leptogenesis epoch. Moreover, the conditions (20) and (25) can be also satisfied. We further assume the $\sigma_{1,2}$ scalars are quasidegenerate,

$$M_{\sigma_2} - M_{\sigma_1} = 10^{-5} M_{\sigma_{1,2}} \gg \Gamma_{\sigma_{1,2}}. \quad (35)$$

The CP asymmetries $\epsilon_{\sigma_{1,2}}^\Delta$ then can arrive at

$$\begin{aligned} \epsilon_{\sigma_1}^\Delta &= \epsilon_{\sigma_2}^\Delta = 2.74 \times 10^{-7} \sin \delta \\ \text{with } \delta &= \text{Arg}(\mu_2^* \mu_1 \kappa_2 \kappa_1^*). \end{aligned} \quad (36)$$

Therefore, the final baryon asymmetry can be consistent with the observed value,

$$\frac{n_B}{s} = 10^{-10} \left(\frac{\sin \delta}{0.15} \right). \quad (37)$$

VII. DISCUSSIONS AND COMMENTS

In this section, we make a few phenomenological comments on the model.

- (i) The left-right symmetry breaking, the quark seesaw and the type-II seesaw scales are all at the TeV scale so that they can be tested at the colliders such as the LHC. For example, in this model, the dominant decay mode of W_R is not to $\ell^\pm \ell^\pm jj$ but rather to $t\bar{b}$ final states as well as to $\ell + \text{missing } E_T$. The LHC limit on W_R for this mode is 4 TeV [23]. Such right-handed keV neutrinos may also be accessible at the KATRIN experiment [16]. Also, the quark seesaw leads to flavor violating decays of t , e.g., $t \rightarrow c + g$ etc., which can be used to test these models.

- (ii) Since the masses of $\Delta_{L,R}$ are close to 3 and 1 TeV, respectively, from the flavor structure of the f -coupling from neutrino mixings, one expects flavor changing rare decays of the τ lepton. With $f \sim 0.5$, this is consistent with current limits but this model predicts that they should be observable once the search sensitivity goes down. Clearly, the predictions depend on whether the neutrino mass hierarchy is normal or inverted and we do not go into a detailed discussion of this in this paper.
- (iii) For light ν_R and low mass W_R , there are stringent supernova 1987A bound [24] on W_R mass of ~ 20 TeV or so. However, it must be pointed out that these bounds are dependent on detailed supernova properties, e.g., radial dependence of temperature as well as possible new interactions [25]. For example, in addition to the interactions in the model, if there is a singly charged particle η^+ , its interactions could trap the right-handed neutrinos [25] and help avoid the above bound.

VIII. CONCLUSION

In this paper, we show how to avoid the lower bound on W_R mass that arises from type-I leptogenesis in TeV-scale left-right models. The basic idea is to consider a quark seesaw version of the model where left-right symmetry breaking is decoupled from gauge symmetry breaking and to generate lepton number not by the decay of right-handed

neutrinos but rather via the decay of some singlet scalar bosons. The latter decays generate an asymmetry between the lepton number carrying Higgs triplets responsible for type-II seesaw, which subsequently decay to leptons via the type-II coupling and transfer their asymmetry to lepton number. The model requires a resonant enhancement of the CP asymmetries to obtain the observed baryon asymmetry in the universe, and therefore the scalars $\sigma_{1,2}$ must be quasidegenerate at the level of $\Delta M_\sigma/M_\sigma = 10^{-5}$. Our model predicts keV sterile neutrinos coupling to right-handed currents. The strong CP problem is also solved by parity symmetry. Furthermore, the left-right symmetry breaking, the quark seesaw, and the type-II seesaw are at the TeV scale so that they can be tested at the colliders such as the LHC. For example, in this model, the dominant decay mode is not to $\ell^\pm \ell^\pm jj$ but rather to $t\bar{b}$ final states as well as to $\ell + \text{missing } E_T$. Such right-handed keV neutrinos may also be accessible at the KATRIN experiment [16].

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