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Parametrized leptogenesis from linear seesaw

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Abstract We present a purely linear seesaw mechanism in a mirror left-right symmetric framework and then realize a novel leptogenesis scenario for parametrizing the cosmic baryon asymmetry by the charged lepton masses, the light Majorana neutrino mass matrix and the heavy mirror charged lepton masses. Through the same Yukawa couplings, the lepton-number-conserving decays of the mirror charged leptons can generate three individual lepton asymmetries stored in the ordinary lepton flavors, while the lepton-number-violating processes for the Majorana neutrino mass generation can wash out part of these lepton asymmetries. The remnant lepton asymmetries then can be partially converted to a baryon asymmetry by the sphaleron processes.

1 Introduction

The precise measurements on the atmospheric, solar, accelerator and reactor neutrinos have established the phenomenon of neutrino oscillations. This fact implies three flavors of neutrinos should be massive and mixed [1]. Meanwhile, the cosmological observations indicate the neutrinos should be extremely light [1]. The tiny but nonzero neutrino masses call for new physics beyond the $SU(3)_c \times SU(2)_L \times U(1)_Y$ standard model (SM). Furthermore, the SM is challenged by other puzzles such as the cosmic baryon asymmetry [1]. Currently a seesaw [2–5] extension of the SM has become very attractive since it can simultaneously explain the small neutrino masses and the cosmic baryon asymmetry [6]. In this popular scenario [6–22], we do not know much about the masses and couplings involving the non-SM fields. Consequently, we cannot get an exact relation between the cosmic baryon asymmetry and the neutrino mass matrix. For example, we can expect a successful leptogenesis in the canonical seesaw model even if the neutrino mass matrix does not contain any CP phases [23].

In a previous paper [28], we showed that the cosmic baryon asymmetry could be well described by *the Dirac neutrino mass matrix, the charged lepton masses and the heavy mirror charged lepton masses* within a mirror left-right symmetric framework based on the gauge groups $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_X$ [29–32]. However, we had to introduce several scalars irrelevant to the neutrino mass generation so that we could eliminate some harmful lepton asymmetries. The model thus seemed complicated and uneconomical.

In the present paper, we shall develop a novel leptogenesis [6] scenario where the neutrinos can obtain their tiny Majorana masses in the so-called linear [25–27] seesaw way while the cosmic baryon asymmetry can be parameterized by the neutrino and charged lepton mass matrices as well as the heavy mirror charged lepton masses. This scenario shares some features of the work [28]. For example, it is also in the mirror left-right symmetric framework. However, this model does not contain any particles irrelevant to the neutrino mass generation. Specifically, some Yukawa interactions can accommodate the lepton-number-conserving decays of the mirror electron-positron pairs to produce three individual lepton asymmetries in the ordinary lepton flavors although the lepton number is exactly conserved. The same Yukawa interactions can participate in the Majorana neutrino mass generation and then can lead to some lepton-number-violating processes to wash out part of the produced lepton asymmetries. The $SU(2)_L$ sphaleron [33] processes then can partially convert the remnant lepton asymmetries to a baryon asymmetry.

2 The model

We extend the SM fermions and Higgs scalar, i.e.

$$q_L \left(2, 1, +\frac{1}{3}\right)(+1), \ d_R \left(1, 1, -\frac{2}{3}\right)(+1),$$
$$u_R \left(1, 1, +\frac{4}{3}\right)(+1), l_L (2, 1, -1)(-1), \ l_L(2, 1, -1)(-1),$$

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$$e_R(1, 1, -2)(-1); \phi_L(2, 1, -1)(0), \tag{1}$$

by a mirror copy [28,34–45],

$$Q_R\left(1, 2, +\frac{1}{3}\right)(-1), \quad D_L\left(1, 1, -\frac{2}{3}\right)(-1),$$

$$U_L\left(1, 1, +\frac{4}{3}\right)(-1), \quad L_R(1, 2, -1)(+1),$$

$$E_L(1, 1, -2)(+1); \phi_R(1, 2, -1)(0). \tag{2}$$

Here and thereafter the first brackets following the fields describe the transformations under the $SU(2)_L \times SU(2)_R \times U(1)_X$ gauge groups, while the second brackets denote an imposed $U(1)_{3B-L}$ global symmetry. Note the SM fields have a $U(1)_{3B-L}$ charge opposite to their mirror partners. Our model also contains the Higgs scalars $\chi_{L,R}$, Σ, ξ and the fermions N_L as below,

$$\chi_L(2, 1, -1)(-2), \chi_R(1, 2, -1)(+2), \Sigma(2, 2, 0)(-2),$$

$$\xi(1, 1, 0)(+1); N_L(1, 1, 0)(-1). \tag{3}$$

The Higgs bidoublet Σ can be expressed in terms of two $SU(2)_L$ doublets, i.e. $\Sigma = [\sigma_1 \, \tilde{\sigma}_2]$.

The usual electric charge Q and hypercharge Y are related to the $SU(2)_L \times SU(2)_R \times U(1)_X$ quantum numbers as [29–32],

$$Y = T_{3R} + \frac{X}{2}, \quad Q = T_{3L} + T_{3R} + \frac{X}{2} = T_{3L} + Y, \tag{4}$$

where $T_{3L(R)} = 0$ for the $SU(2)_{L(R)}$ singlets, $T_{3L(R)} = +1/2$ for the up components of the $SU(2)_{L(R)}$ doublets, while $T_{3L(R)} = -1/2$ for the down components. So, we can easily read the hypercharges of the fermions and scalars after the $SU(2)_L \times SU(2)_R \times U(1)_X$ symmetry is spontaneously broken down to the SM $SU(2)_L \times U(1)_Y$ symmetry. For example, the SM fermions and Higgs scalar $(q_L, d_R, u_R, l_L, e_R, \phi_L)$ have their hypercharges $(+\frac{1}{6}, -\frac{1}{3}, +\frac{2}{3}, -\frac{1}{2}, -1, -\frac{1}{2})$ as they are in the SM.

In this work, the discrete left-right symmetry is taken to be a mirror CP (charge conjugation and parity) symmetry:

$$\phi_L \xleftrightarrow{CP} \phi_R^*, \quad \chi_L \xleftrightarrow{CP} \chi_R^*, \quad \Sigma \xleftrightarrow{CP} \Sigma^T, \quad \xi \xleftrightarrow{CP} \xi,$$

$$q_L \xleftrightarrow{CP} Q_R^c, \quad d_R \xleftrightarrow{CP} D_L^c, \quad u_R \xleftrightarrow{CP} U_L^c,$$

$$l_L \xleftrightarrow{CP} L_R^c, \quad e_R \xleftrightarrow{CP} E_L^c, \quad N_L \xleftrightarrow{CP} N_L. \tag{5}$$

The global $U(1)_{3B-L}$ symmetry and the discrete mirror CP symmetry are both conserved exactly.

For simplicity we do not write down the full scalar potential. Instead we show the cubic terms and some quartic terms as below,

$$V \supset \kappa \xi^2 \left(\phi_L^\dagger \chi_L + \phi_R^T \chi_R^*\right) + \rho \left(\phi_L^\dagger \Sigma \chi_R\right.$$

$$\left. + \phi_R^T \Sigma^T \chi_L^*\right) + \text{H.c.}. \tag{6}$$

As for the Yukawa interactions, they are given by

$$\mathcal{L}_Y = -y_d \left(\bar{q}_L \tilde{\phi}_L d_R + \bar{Q}_R^c \tilde{\phi}_R^* D_L^c\right)$$

$$\quad - y_u \left(\bar{q}_L \phi_L u_R + \bar{Q}_R^c \phi_R^* U_L^c\right) - y_Q \bar{q}_L \tilde{\Sigma} Q_R$$

$$\quad - y_e \left(\bar{l}_L \tilde{\phi}_L e_R + \bar{L}_R^c \tilde{\phi}_R^* E_L^c\right)$$

$$\quad - y_N \left(\bar{l}_L \chi_L N_L^c + \bar{L}_R^c \chi_R^* N_L^c\right) - y_L \bar{l}_L \Sigma L_R + \text{H.c.}$$

$$\quad \text{with} \quad y_Q = y_Q^T, \quad y_L = y_L^T. \tag{7}$$

Note the $U(1)_{3B-L}$ global symmetry has forbidden the gauge-invariant mass terms of the $[SU(2)]$-singlet fermions.

We would like to give some comments on the original models for linear seesaw [25–27]. In these models, three gauge-singlet fermions, i.e. the N_L fermions in our model, can always have a Majorana mass term and hence the left-handed neutrinos can always obtain their Majorana masses through an inverse/double seesaw besides the linear seesaw [25–27]. There are no symmetries constraining the size of this Majorana mass term. Accordingly, although the linear seesaw is allowed to dominate over the inverse/double seesaw, such choice can not be explained by certain symmetries. Moreover, these models can accommodate a conventional leptogenesis scenario. As we will demonstrate in the following, our model provides a purely linear seesaw and a new leptogenesis scenario.

3 Fermion masses

From the full potential which are not shown for simplicity, we can expect the VEVs to be

$$\langle\xi\rangle, \langle\chi_R^0\rangle, \langle\phi_R^0\rangle \gg \langle\phi_L^0\rangle, \langle\chi_L^0\rangle, \langle\sigma_{1,2}^0\rangle. \tag{8}$$

For example, after ξ and χ_R develop their VEVs, ϕ_R can obtain a VEV due to the second term in Eq. (6). Subsequently, σ_1 develops its VEV and then ϕ_L and χ_L can acquire their VEVs through the third and forth terms in Eq. (6). The hierarchical VEVs (8) can only give a negligible effect on the precision measurements such as the $W - W'$ and $Z - Z'$ mixing.

The Yukawa interactions (7) then can reasonably yield

$$y_Q\langle\sigma_2^0\rangle \ll y_d\langle\phi_R^0\rangle, \quad y_Q\langle\sigma_1^0\rangle \ll y_u\langle\phi_R^0\rangle,$$

$$y_L\langle\sigma_2^0\rangle \ll y_e\langle\phi_R^0\rangle, \quad y_L\langle\sigma_1^0\rangle \ll y_N\langle\chi_R^0\rangle. \tag{9}$$

This means the mixing between the ordinary charged fermions $(f = d, u, e)$ and their mirror partners $(F = D, U, E)$ should be of the order of $\mathcal{O}(y_{Q,L}\langle\sigma_{1,2}^0\rangle/M_F) \ll 1$ and hence can be safely ignored. Thus the mass eigenstates of the charged fermions can come from

$$\mathcal{L} \supset -m_f \bar{f}_L f_R - M_F \bar{F}_R^c F_L^c + \text{H.c.} \quad \text{with}$$

$$m_f = y_f\langle\phi_L^0\rangle \ll M_F = y_f\langle\phi_R^0\rangle. \tag{10}$$

Meanwhile, we can apply the linear seesaw mechanism [25–27] to the neutral fermions, i.e.

$$
\mathcal{L} \supset - \begin{bmatrix} \bar{\nu}_L & \bar{N}_L & \bar{N}_R^c \end{bmatrix} \begin{bmatrix} 0 & y_N \langle \chi_L^0 \rangle & y_L \langle \sigma_1^0 \rangle \\ y_N^T \langle \chi_L^0 \rangle & 0 & y_N^T \langle \chi_R^0 \rangle \\ Y_L^T \langle \sigma_1^0 \rangle & y_N \langle \chi_R^0 \rangle & 0 \end{bmatrix} \begin{bmatrix} \nu_L^c \\ N_L^c \\ N_R \end{bmatrix} + \text{H.c.}
$$

$$
\simeq -M_N \bar{N}_R^c N_L^c - \frac{1}{2} m_\nu \bar{\nu}_L \nu_L^c + \text{H.c. with}
$$

$$
M_N = y_N \langle \chi_R^0 \rangle, \quad m_\nu = m_\nu^T = -2 y_L \langle \sigma_1^0 \rangle \frac{\langle \chi_L^0 \rangle}{\langle \chi_R^0 \rangle}. \quad (11)
$$

Here N_R denotes the neutral components of the $SU(2)_R$ doublets L_R.

Note the VEVs $\langle \phi_L^0 \rangle$, $\langle \chi_L^0 \rangle$ and $\langle \sigma_{1,2}^0 \rangle$, should be constrained by

$$
\sqrt{\langle \phi_L^0 \rangle^2 + \langle \chi_L^0 \rangle^2 + \langle \sigma_1^0 \rangle^2 + \langle \sigma_2^0 \rangle^2} \equiv v = 174 \, \text{GeV},
$$

$$
\langle \phi_L^0 \rangle = \frac{m_t}{y_t} > \frac{m_t}{\sqrt{4\pi}} \simeq 48.5 \, \text{GeV}, \quad (12)
$$

which implies

$$
\mu_2^2 \equiv 2 \langle \chi_L^0 \rangle \langle \sigma_1^0 \rangle \le \langle \chi_L^0 \rangle^2 + \langle \sigma_1^0 \rangle^2 < \left(174^2 - 48.5^2 \right) \text{GeV}^2
$$

$$
= 2.79 \times 10^4 \, \text{GeV}^2. \quad (13)
$$

In addition, the tiny but nonzero neutrino masses require

$$
\mu_1 \equiv \frac{2 \langle \chi_L^0 \rangle \langle \sigma_1^0 \rangle}{\langle \chi_R^0 \rangle} = \frac{\hat{m}_\nu}{\hat{y}_L} > \frac{m_{\text{max}}}{\sqrt{4\pi}}
$$

$$
= 0.014 \, \text{eV} \left(\frac{m_{\text{max}}}{0.05 \, \text{eV}} \right), \quad (14)
$$

with m_{max} being the largest eigenvalue of the neutrino mass matrix,

$$
m_\nu = U \hat{m}_\nu U^T = U \text{diag}\{m_1, m_2, m_3\} U^T. \quad (15)
$$

Here the PMNS matrix U contains three mixing angles, one Dirac phase and two Majorana phases, i.e.

$$
U = \begin{bmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23}-c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23}-s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23}-c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23}-s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{bmatrix}
$$
$$
\times \text{diag}\{e^{i\alpha_1/2}, e^{i\alpha_2/2}, 1\}. \quad (16)
$$

The linear seesaw can be also understood at the electroweak level. From the Yukawa interactions (7), we can read

$$
\mathcal{L} \supset -\hat{y}_e \bar{l}_L \tilde{\phi}_L e_R - \hat{M}_E \bar{E}_R^c E_L^c - X_R \hat{y}_N \bar{l}_L \chi_L N_L^c
$$
$$
-\hat{M}_N \bar{N}_R^c N_L^c - y_L X_R^* \bar{l}_L \sigma_1 N_R - y_L \bar{l}_L \tilde{\sigma}_2 E_R + \text{H.c.}
$$
$$
\text{with } y_N = X_R \hat{y}_N X_L^\dagger, \quad X_{L,R}^\dagger X_{L,R} = X_{L,R} X_{L,R}^\dagger = 1. \quad (17)
$$

Here the yukawa couplings y_e and the mass matrices $M_{E,N}$ have been chosen to be diagonal and real without loss of generality. As shown in Fig. 1, the left-handed neutrinos ν_L

can acquire their Majorana masses (11) by integrating out the heavy Dirac pairs $N = N_L + N_R$, i.e.

$$
\mathcal{L} \supset -\frac{1}{2} y_L X_R^* \frac{1}{\hat{M}_N} \hat{y}_N X_R^T \bar{l}_L \sigma_1 \chi_L^T l_L^c
$$

$$
-\frac{1}{2} X_R \hat{y}_N \frac{1}{\hat{M}_N} X_R^\dagger y_L^T \bar{l}_L \chi_L \sigma_1^T l_L^c + \text{H.c.}
$$

$$
= -\frac{1}{2} \frac{y_L}{\langle \chi_R^0 \rangle} \left(\bar{l}_L \sigma_1 \chi_L^T l_L^c + \bar{l}_L \chi_L \sigma_1^T l_L^c \right) + \text{H.c..} \quad (18)
$$

Note from the linear seesaw (11),

$$
m_\nu = -2 y_L \langle \sigma_1^0 \rangle \frac{\langle \chi_L^0 \rangle}{\langle \chi_R^0 \rangle} = -y_L \frac{\mu_2^2}{\hat{M}_N} \hat{y}_N, \quad (19)
$$

the conditions $y_{N,L} < \sqrt{4\pi}$, $\mu_2^2 < 2.79 \times 10^4 \, \text{GeV}^2$ and $m_{\text{max}} \ge 0.05 \, \text{eV}$ should put an upper bound on the heavy Dirac fermion masses,

$$
M_N < 7 \times 10^{15} \, \text{GeV}. \quad (20)
$$

4 Lepton and baryon asymmetries

As shown in Fig. 2, the mirror charged leptons E_β can decay into the ordinary lepton doublets $l_{L\alpha}$ and the Higgs doublet σ_2. These decays can generate three individual lepton asymmetries $L_{e,\mu,\tau}$ stored in the ordinary lepton flavors $l_{Le, L\mu, L\tau}$ if the CP is not conserved, i.e.

$$
L_{\beta\alpha} \propto \Gamma(E_\beta \to L_{L\alpha} \sigma_2) - \Gamma(E_\beta^c \to L_{L\alpha}^c \sigma_2^*) \neq 0,
$$

$$
L_\alpha = \sum_\beta L_{\beta\alpha}. \quad (21)
$$

We calculate the decay width at tree level,

$$
\Gamma_\beta \equiv \Gamma(E_\beta \to l_{L\alpha} + \sigma_2) = \Gamma(E_\beta^c \to l_{L\alpha}^c + \sigma_2^*)
$$

$$
= \frac{1}{16\pi} \left(y_L^\dagger y_L \right)_{\beta\beta} M_{E_\beta}
$$

$$
= \frac{1}{16\pi} \frac{(m_\nu^\dagger m_\nu)_{\beta\beta}}{\mu_1^2} M_{E_\beta}, \quad (22)
$$

and then the CP asymmetry at one-loop level,

$$
\varepsilon_{\beta\alpha} = \frac{\Gamma(E_\beta \to l_{L\alpha} + \sigma_2) - \Gamma(E_\beta^c \to l_{L\alpha}^c + \sigma_2^*)}{\Gamma_\beta}
$$

$$
= \frac{1}{4\pi} \sum_\rho \frac{\text{Im}\left[\left(y_L^\dagger y_L \right)_{\rho\beta} y_{L\alpha\beta}^* y_{L\alpha\rho} \right]}{\left(y_L^\dagger y_L \right)_{\beta\beta}} \frac{M_{E_\beta}^2}{M_{E_\beta}^2 - M_{E_\rho}^2}
$$

$$
= \frac{1}{4\pi} \sum_\rho \frac{\text{Im}\left[(m_\nu^\dagger m_\nu)_{\rho\beta} m_{\alpha\beta}^* m_{\alpha\rho} \right]}{\mu_1^2 \left(m_\nu^\dagger m_\nu \right)_{\beta\beta}} \frac{m_\beta^2}{m_\beta^2 - m_\rho^2}. \quad (23)
$$

It is easy to check

$$\varepsilon_{\beta e} + \varepsilon_{\beta \mu} + \varepsilon_{\beta \tau} \equiv 0, \tag{24}$$

as a result of the lepton number conservation.

The decays of the mirror electron-positron pairs should dominate the individual lepton asymmetries $L_{e,\mu,\tau}$ since the mirror electron is much lighter than the mirror muon and tau. Actually, the out-of-equilibrium decays of the mirror muons and taus in principle could also produce the individual lepton asymmetries $L_{e,\mu,\tau}$. However, at that moment the Yukawa interactions involving the mirror electrons were still in equilibrium so that they would wash out the produced individual lepton asymmetries $L_{e,\mu,\tau}$ from the mirror muon and tau decays. When the mirror electrons and positrons go out of equilibrium at a temperature T_D, the individual lepton asymmetries $L_{e,\mu,\tau}$ can be produced, i.e.

$$L_\alpha \simeq \varepsilon_{e\alpha} \left(\frac{n_{E_e}}{s} \right) \Bigg|_{T=T_D} . \tag{25}$$

Here n_{E_e} is the number density and s is the entropy density. For the following demonstration, we specify the decay width,

$$\Gamma_e = \frac{1}{16\pi} \frac{(m_\nu^\dagger m_\nu)_{ee}}{\mu_1^2} M_{E_e}, \tag{26}$$
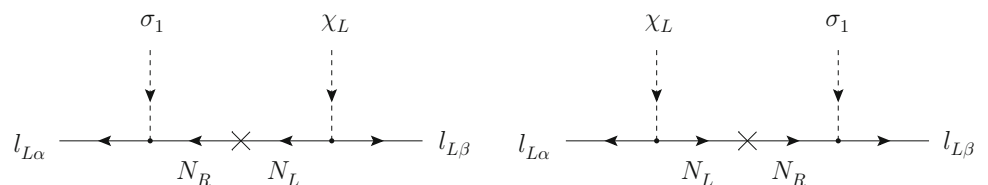
and the CP asymmetries,

$$\varepsilon_{ee} \simeq -\frac{1}{4\pi} \frac{\mathrm{Im}\left[(m_\nu^\dagger m_\nu)_{\mu e} \, m_{ee}^* m_{e\mu} \right]}{\mu_1^2 \left(m_\nu^\dagger m_\nu \right)_{ee}} \frac{m_e^2}{m_\mu^2},$$

$$\varepsilon_{e\mu} \simeq -\frac{1}{4\pi} \frac{\mathrm{Im}\left[(m_\nu^\dagger m_\nu)_{\mu e} \, m_{\mu e}^* m_{\mu\mu} \right]}{\mu_1^2 \left(m_\nu^\dagger m_\nu \right)_{ee}} \frac{m_e^2}{m_\mu^2},$$

$$\varepsilon_{e\tau} \simeq -\frac{1}{4\pi} \frac{\mathrm{Im}\left[(m_\nu^\dagger m_\nu)_{\mu e} \, m_{\tau e}^* m_{\tau\mu} \right]}{\mu_1^2 \left(m_\nu^\dagger m_\nu \right)_{ee}} \frac{m_e^2}{m_\mu^2}. \tag{27}$$

On the other hand, the model provides the lepton-number-violating interactions for generating the Majorana neutrino masses, as shown in Fig. 1. The interaction rates of the related lepton-number-violating processes such as $l_{L\alpha} l_{L\beta} \to \chi_L \sigma_1$ are computed by [46]

$$\Gamma_{\alpha\beta} = \frac{1}{\pi^3} \frac{\left| m_{\alpha\beta} \right|^2 T^3}{v^4} \quad \text{for } T < M_N. \tag{28}$$

At the temperature,

$$T_{\alpha\beta} = \left(\frac{8\pi^9 g_*}{90} \right)^{\frac{1}{2}} \frac{v^4}{M_{\mathrm{Pl}} \left| m_{\alpha\beta} \right|^2}$$

$$= 1.63 \times 10^{13} \, \mathrm{GeV} \left(\frac{0.05 \, \mathrm{eV}}{\left| m_{\alpha\beta} \right|} \right)^2, \tag{29}$$

these lepton-number-violating processes will begin to decouple, i.e.

$$\left[\Gamma_{\alpha\beta} < H(T) = \left(\frac{8\pi^3 g_*}{90} \right)^{\frac{1}{2}} \frac{T^2}{M_{\mathrm{Pl}}} \right]\Bigg|_{T<T_{\alpha\beta}}. \tag{30}$$

Here $H(T)$ is the Hubble constant with $M_{\mathrm{Pl}} \simeq 1.22 \times 10^{19}$ GeV being the Planck mass and $g_* = 110.75$ being the relativistic degrees of freedom (the SM fields plus one additional Higgs doublet σ_2).

If all of the above lepton-number-violating processes such as $l_{L\alpha} l_{L\beta} \to \chi_L \sigma_1$ keep in equilibrium during the mirror electron decays, the individual lepton asymmetries $L_{e,\mu,\tau}$ can not be produced at all. Instead, they would be washed out once they were produced. However, we can require that the lepton-number-violating processes such as $l_{L\alpha} l_{L\beta} \to \chi_L \sigma_1$ have no effects on the lepton flavor $l_{L\gamma \neq \alpha,\beta}$. In consequence, the mirror electron decays can still contribute a lepton asymmetry $L_{\gamma \neq \alpha,\beta}$. This means we can expect the lepton asymmetry stored in certain lepton flavor(s) to survive from the lepton-number-violating processes, i.e.

$$L = \begin{cases} L_e & \text{for } T_{\mu\mu,\mu\tau,\tau\tau} < T_D < T_{ee,e\mu,e\tau}, \\ L_\mu & \text{for } T_{ee,e\tau,\tau\tau} < T_D < T_{e\mu,\mu\mu,\mu\tau}, \\ L_\tau & \text{for } T_{ee,e\mu,\mu\mu} < T_D < T_{e\tau,\mu\tau,\tau\tau}, \\ L_e + L_\mu & \text{for } T_{\tau\tau} < T_D < T_{ee,e\mu,e\tau,\mu\mu,\mu\tau}, \\ L_e + L_\tau & \text{for } T_{\mu\mu} < T_D < T_{ee,e\mu,e\tau,\mu\tau,\tau\tau}, \\ L_\mu + L_\tau & \text{for } T_{ee} < T_D < T_{e\mu,e\tau,\mu\mu,\mu\tau,\tau\tau}. \end{cases} \tag{31}$$

Here we have assumed the heavy Dirac fermions $N = N_L + N_R$ much heavier than the crucial temperature T_D. The $SU(2)_L$ sphaleron processes then can partially transfer the remnant lepton asymmetry L to a baryon asymmetry B. From Eq. (29), the lepton-number-violating processes can go out of equilibrium before the sphalerons become active, i.e.

$$T_{\alpha\beta} > T_{\mathrm{sph}} \simeq 10^{12} \, \mathrm{GeV}. \tag{32}$$

Fig. 1 The lepton-number-violating processes for generating the Majorana neutrino masses

Fig. 2 The lepton-number-conserving decays of the mirror charged leptons into the ordinary leptons
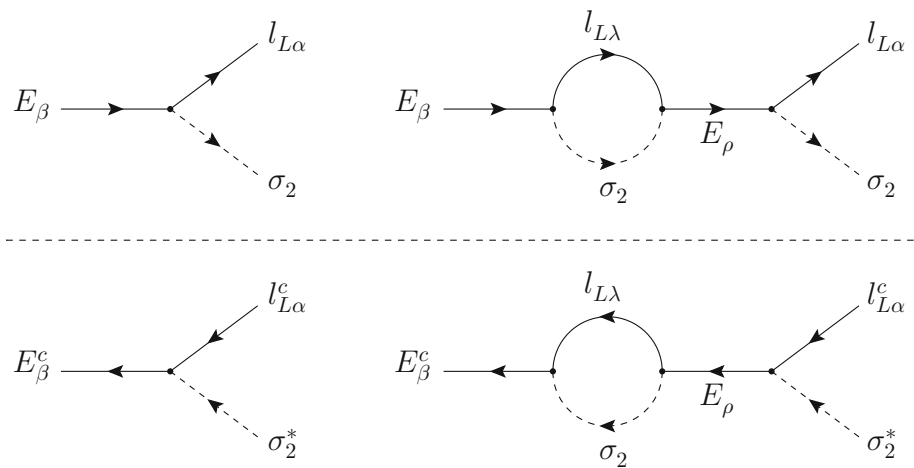
The final baryon asymmetry B then can be given by [47]

$$B = -\frac{28}{79}L. \tag{33}$$

We would like to emphasize that both the production and the washout of the individual lepton asymmetries $L_{e,\mu,\tau}$ take place above the temperatures $T > T_{\text{sph}} \simeq 10^{12}\,\text{GeV}$, at which the Yukawa interactions between the SM left- and right-handed leptons have not been in equilibrium yet. So, the flavor effect in the usual leptogenesis scenario [15,24] will not appear in the above demonstrations. Instead, the remnant lepton asymmetry L in Eq. (31) can be treated as a single-flavor lepton asymmetry when it participates in the sphaleron processes [15,24].

5 Numerical example

To give a numerical estimation, we consider the weak washout region [48], i.e,

$$K = \frac{\Gamma_e}{2H}\Big|_{T=M_{E_e}} \ll 1. \tag{34}$$

Accordingly, when the mirror electrons and positrons finally decay [48],

$$T_D \sim M_{E_e}\sqrt{K}, \tag{35}$$

they are very overabundant: $n_{E_e} \sim n_\gamma$, i.e. no exponential suppression in their number density. Here n_γ is the number density of photons. We hence can expect Eq. (25) to be [48]

$$L_\alpha \simeq \varepsilon_{e\alpha}\left(\frac{n_{E_e}}{s}\right)\Big|_{T=T_D} \sim \varepsilon_{e\alpha}\left(\frac{n_\gamma}{s}\right) \sim \frac{\varepsilon_{e\alpha}}{g_*}. \tag{36}$$

In this case, as shown in Eqs. (26), (27), (29), (31), (33), (34), (35) and (36), the final baryon asymmetry can be fully determined by the charged lepton masses, the Majorana neutrino mass matrix and the parameters $\mu_{1,2}$ and M_{E_e}.

We have known the charged lepton masses $m_e = 511\,\text{keV}$, $m_\mu = 107\,\text{MeV}$, and the normal(inverted) neutrino parameters $\Delta m_{21}^2 = m_2^2 - m_1^2 = 7.37 \times 10^{-5}\,\text{eV}^2$, $\Delta m_{31}^2(\Delta m_{23}^2) = m_3^2 - m_1^2(m_2^2 - m_3^2) = 2.56(2.54) \times 10^{-3}\,\text{eV}^2$, $\sin^2\theta_{12} = 0.297$, $\sin^2\theta_{23} = 0.425(0.589)$, $\sin^2\theta_{13} = 0.0215(0.0216)$ [1]. It seems difficult for the inverted hierarchical and quasi-degenerate neutrinos to fulfil the conditions in Eq. (31). So, we consider the normal hierarchical neutrinos to give a numerical estimation. As an example, we fix $m_1 = 0$, $\alpha_2 = \pi/2$, $\delta = 3\pi/2$, and then take $\mu_1 = 0.21\,\text{eV}$, $\mu_2^2 = 2.79 \times 10^4\,\text{GeV}^2$, $M_{E_e} = 3 \times 10^{16}\,\text{GeV}$. With these inputs, we obtain $K \simeq 1.9 \times 10^{-5}$, $\varepsilon_{ee} \simeq -3.4 \times 10^{-8}$, $T_D \sim 1.3 \times 10^{14}\,\text{GeV}$, $T_{\tau\tau} \simeq 4.3 \times 10^{13}\,\text{GeV}$, $T_{\mu\tau} \simeq 8.1 \times 10^{13}\,\text{GeV}$, $T_{\mu\mu} \simeq 7.5 \times 10^{13}\,\text{GeV}$, $T_{e\mu} \simeq 7.7 \times 10^{14}\,\text{GeV}$, $T_{e\tau} \simeq 2.1 \times 10^{15}\,\text{GeV}$, $T_{ee} \simeq 1.1 \times 10^{16}\,\text{GeV}$ and hence $B \sim 10^{-10}$.

6 Conclusion

In this paper we have demonstrated a novel linear seesaw scenario for paramerizing the cosmic baryon asymmetry by the charged lepton masses, the light Majorana neutrino mass matrix and the heavy mirror charged lepton masses. Through the lepton-number-conserving decays of the mirror electron-positron pairs, we can obtain three individual lepton asymmetries stored in the ordinary lepton flavors although the total lepton asymmetry is exactly zero. The lepton-number-violating processes for the neutrino mass generation can wash out the lepton asymmetry stored in certain ordinary lepton flavor(s). Remarkably, these lepton-number-conserving and lepton-number-violating interactions originate from the same Yukawa couplings. The remnant lepton asymmetry can be partially converted to a baryon asymmetry by the sphaleron processes. Our scenario seems difficult to work for an inverted hierarchical or a quasi-degenerate neutrino spectrum. Instead, it prefers to the normal hierarchical neutrino

spectrum which is hinted by the present data from neutrino oscillations and cosmological observations [1].

Data Availability Statement This manuscript has no associated data or the data will not be deposited. [Authors' comment: This is a theoretical study and no experimental data has been listed.]

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