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Research Article

A New Left-Right Symmetry Model

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We propose a new L-R symmetry model where the L-R symmetry transformation reverses both the L-R chirality and the local quantum number. We add to the model a global quantum number F whose value is one for fermions (minus one for antifermion) and vanishes for bosons. For each standard model (SM) particle, we have the corresponding L-R dual particle whose mass is very large and which should have decayed at the current low energy level. Due to the global quantum number F, there is no Majorana neutrino in the model but a Dirac seesaw mechanism can still occur and the usual three active neutrino oscillation can still be realized. We add two leptoquarks and their L-R duals, for generating the baryon number asymmetry and for facilitating the decay of the L-R dual particles. The decay of the L-R dual particles will produce a large entropy to the SM sector and give a mechanism for avoiding the big bang nucleosynthesis constraint.

1. Introduction

The left-right symmetry (LRS) model was originally introduced based on aesthetic reason, to explain the nonexistence of the right weak current [1–3]. Even though it is no longer the center of attention in the particle physics community, there is no physical fact/observation that ruled out the LRS model as a viable standard model extension. All fermions in the original LRS model are doublets of either $\mathrm{SU}(2)_L$ or $SU(2)_R$ and one has to add a L-R bidoublet scalar to facilitate the mass generation of the fermions [4]. Unfortunately, this bidoublet also causes the mixing between the left and right weak gauge bosons, which is unobserved until now. There is also a variant of the LRS model where for each SM particle one has its left (or right) chiral partner, thus doubling the SM particle content [5–7]. In this variant model, fermions gain their mass through seesaw mechanism [8–10], except for neutrinos that undergo a double seesaw mechanism [5–7]. There is no requirement for bidoublet in this variant model, but the seesaw mechanism for the charged fermions is not compatible with the recent result of the Higgs scalar coupling to the SM fermions, where the coupling is proportional to the fermion mass [11, 12]. There is also another variant of the LRS model where one doubles both the particle content and the gauge group of the SM. This model is also known as the mirror model [13, 14]. There is no interaction between SM particles and mirror particles, except through the U(1) kinetic mixing and the scalar mixing in the scalar potentials. This model has a very rich hidden sector and the mirror particles can be dark matter candidates. Since the mirror particles do not share any common quantum numbers with the SM particles, the chance to directly detect or generate any mirror particle at the accelerator is very small.

In this paper, we propose a new variant of the LRS model with a new LRS transformation and an additional global quantum number. The new model is still invariant under the gauge group $\mathrm{SU}(3) \otimes \mathrm{SU}(2)_L \otimes \mathrm{SU}(2)_R \otimes \mathrm{U}(1)_Y$, but with a new LRS transformation that reverses both the L-R chirality and the local quantum number of a particle. Specifically, the new LRS transformation will change the *L-R* chirality and change the gauge group representation into its charge conjugate representation, with an exception for the SU(2)

Table 1: Irreducible representation and quantum number assignment for the SM and the L-R dual particles with respect to the L-R gauge group and the global quantum number F.

SM particles	Irreps	Dual particles	Irreps
Fermion ($F = 1$)			
$L_L \equiv \begin{pmatrix} \nu \\ e \end{pmatrix}$	(**1**,**2**,**1**,-1)	$L_R \equiv \begin{pmatrix} E \\ N \end{pmatrix}$	(**1**,**1**,**2***,1)
ν_R	(**1**,**1**,**1**,0)	N_L	(**1**,**1**,**1**,0)
e_R	(**1**,**1**,**1**,-2)	E_L	(**1**,**1**,**1**,2)
$Q_L \equiv \begin{pmatrix} u \\ d \end{pmatrix}$	(**3**, **2**, **1**, 1/3)	$Q_R \equiv \begin{pmatrix} D \\ U \end{pmatrix}$	$\left(\mathbf{3}^*, \mathbf{1}, \mathbf{2}^*, \frac{-1}{3}\right)$
u_R	(**3**, **1**, **1**, 4/3)	U_L	$\left(\mathbf{3}^*, \mathbf{1}, \mathbf{1}, \frac{-4}{3}\right)$
d_R	(**3**, **1**, **1**,−2/3)	D_L	$\left(\mathbf{3}^*, \mathbf{1}, \mathbf{1}, \frac{2}{3}\right)$
Scalar ($F = 0$)			
$\chi_L = \begin{pmatrix} \chi_\nu \\ \chi_e \end{pmatrix}$	(**1**,**2**,**1**,-1)	$\chi_R = \begin{pmatrix} \chi_E \\ \chi_N \end{pmatrix}$	(**1**,**1**,**2***,1)
η	$\left(\mathbf{3}, \mathbf{1}, \mathbf{1}, \frac{4}{3}\right)$	η^*	$\left(\mathbf{3}^*, \mathbf{1}, \mathbf{1}, \frac{-4}{3}\right)$
ρ	$\left(\mathbf{3}, \mathbf{1}, \mathbf{1}, \frac{-2}{3}\right)$	ρ^*	$\left(\mathbf{3}^*, \mathbf{1}, \mathbf{1}, \frac{2}{3}\right)$

representation where the L-R chirality change is also changing $SU(2)_L$ representation into $SU(2)_R$ representation and vice versa. Since the gauge couplings are real, even though this new LRS transformation is different from the original LRS transformation, the gauge couplings of $SU(2)_L$ and $SU(2)_R$ are still equal, $g_L = g_R \equiv g$.

We add an additional global quantum number F whose value is one for all fermions (minus one for antifermions) and zero for all scalars and gauge bosons. We assume that this global quantum number F is not affected by the new LRS transformation, and therefore, the L-R dual particles are different from the SM antiparticles. The fermion and scalar content of the model are given in Table 1, where we have classified the particles into two categories: the usual SM particles (plus the right-handed singlet neutrinos) and the corresponding L-R dual particles. These two classes of particles are related by the new L-R symmetry transformation. Since both left and right chiral particles exist, this model is free from chiral anomaly.

The particle content in this new model is a variant of the one proposed around two decades ago by Coutinho et al. [5]. However, due to the new L-R symmetry transformation and the global quantum number F, unlike in [5], there is no need of a seesaw mechanism for generating charged fermion masses (no need of double seesaw for neutrinos) in this model. But due to the global quantum number F, there is no Majorana neutrinos and therefore no leptogenesis in this model. The L-R dual particles are very massive but they cannot decay into the SM particles via any gauge interactions. Therefore, we introduce two leptoquarks ρ and η with their corresponding L-R duals to the model as facilitators for the decay of the massive L-R dual particles into the SM particles and the neutrinos. These two leptoquarks and their L-R duals will also produce a baryon asymmetric universe (BAU) for the SM and the L-R dual sectors. Thus, the existence of the leptoquarks is necessary in this model, not only for facilitating the decay of L-R dual particle but also for producing BAU.

2. Scalar Sector

The scalar potential that is invariant under the gauge group and the new L-R transformation is given by the following:

$$\begin{aligned} v = & -\mu_L^2 |\chi_L|^2 - \mu_R^2 |\chi_R|^2 + \mu_1^2 |\eta|^2 + \mu_2^2 |\rho|^2 \\ & + \lambda_1 \left(|\chi_L|^4 + |\chi_R|^4 \right) + \lambda_2 |\eta|^4 + \lambda_3 |\rho|^4 \\ & + \epsilon_1 |\eta|^2 \left(|\chi_L|^2 + |\chi_R|^2 \right) + \epsilon_2 |\rho|^2 \left(|\chi_L|^2 + |\chi_R|^2 \right) \\ & + \epsilon_3 |\eta|^2 |\rho|^2 + \epsilon_4 \left(|\chi_L|^2 |\chi_R|^2 \right) + \epsilon_5 \rho^\dagger \eta \eta^\dagger \rho \cdot \end{aligned} \quad (1)$$

The parameter $\mu_L \neq \mu_R$ corresponds to a soft L-R symmetry breaking term. The parameters μ_i's, λ_i's, and ϵ_i's above can be chosen so that the scalar potential can lead to nonzero vacuum expectation values (VEVs) for certain scalars. The leptoquarks η and ρ should not have any nonzero VEV, since otherwise the gluons will be massive. χ_L and χ_R can have different VEVs, written as follows:

$$\begin{aligned} \langle \chi_L \rangle &= \begin{pmatrix} v_L \\ 0 \end{pmatrix}, \\ \langle \chi_R \rangle &= \begin{pmatrix} 0 \\ v_R \end{pmatrix}, \end{aligned} \quad (2)$$

where in general $v_L \neq v_R$. These nonzero VEVs for χ_L and χ_R will give masses to the left and right weak gauge bosons. Note that because there is no bidoublet scalar in this model, then there is no direct mixing between W_L and W_R.

The relevant terms for the mass of the weak gauge bosons are as follows:

$$\begin{aligned} & \left| \sqrt{\frac{1}{2}} \left(-\frac{i}{2} g \tau \cdot \mathbf{W}_L + \frac{i}{2} g' B_\mu \right) \begin{pmatrix} v_L \\ 0 \end{pmatrix} \right|^2 \\ & + \left| \sqrt{\frac{1}{2}} \left(-\frac{i}{2} g \tau \cdot \mathbf{W}_R - g' \frac{i}{2} B_\mu \right) \begin{pmatrix} 0 \\ v_R \end{pmatrix} \right|^2 \\ & = \left(\frac{1}{2} g v_L \right)^2 W_{\mu L}^+ W_L^{\mu -} + \left(\frac{1}{2} g v_R \right)^2 W_{\mu R}^+ W_R^{\mu -} \\ & + \frac{1}{2} \left(W_{\mu L}^3 W_{\mu R}^3 B_\mu \right) \mathbf{M}_{\mathrm{WB}} \begin{pmatrix} W_L^{3\mu} \\ W_R^{3\mu} \\ B_\mu \end{pmatrix}, \end{aligned} \quad (3)$$

where $\mathbf{M}_{WB}$ is the mass matrix of $W_L^{3\mu}$, $W_R^{3\mu}$, and B^μ fields which is given by

$$\mathbf{M}_{WB} = \frac{v_R^2}{4}\begin{pmatrix} g^2\omega^2 & 0 & -gg'\omega^2 \\ 0 & g^2 & -gg' \\ -gg'\omega^2 & -gg' & g'^2\left(1+\omega^2\right) \end{pmatrix}, \quad (4)$$

where $\omega = v_L/v_R$. The first and second terms in equation (3) directly give the mass of charged gauge bosons, W_L and W_R, namely,

$$\begin{aligned} M_{W_L}^2 &= \frac{1}{4}g^2 v_L^2, \\ M_{W_R}^2 &= \frac{1}{4}g^2 v_R^2. \end{aligned} \quad (5)$$

While diagonalizing the matrix $\mathbf{M}_{WB}$ in equation (3), we obtain the masses of the neutral right and left weak gauge bosons (Z_R, Z_L) and the photon (A)

$$\begin{aligned} M_{Z_R}^2 &= \frac{v_R^2}{8}\left(g^2+g'^2\right)\left(\left(1+\omega^2\right)+\left(\left(1-\omega^2\right)^2+\frac{4\omega^2 g'^4}{\left(g^2+g'^2\right)^2}\right)^{1/2}\right) \\ &\simeq v_R^2 g^2 \frac{\tan^2\theta_W}{\sin^2 2\beta}\left(1+\omega^2\sin^4\beta\right), \end{aligned} \quad (6)$$

$$\begin{aligned} M_{Z_L}^2 &= \frac{v_R^2}{8}\left(g^2+g'^2\right)\left(\left(1+\omega^2\right)-\left(\left(1-\omega^2\right)^2+\frac{4\omega^2 g'^4}{\left(g^2+g'^2\right)^2}\right)^{1/2}\right) \\ &\simeq \frac{1}{4}\frac{v_L^2 g^2}{\cos^2\theta_W}\left(1-\omega^2\sin^4\beta\right), \end{aligned} \quad (7)$$

$$M_A^2 = 0, \quad (8)$$

where the mixing angles θ_W and β are given by

$$\begin{aligned} \sin^2\theta_W &= \frac{g^2 g'^2}{g^4 + 2g^2 g'^2}, \\ \sin^2\beta &= \frac{g'^2}{g^2+g'^2}. \end{aligned} \quad (9)$$

The phenomenology of our model is similar to many LRS models. In particular, it is the same as in [5] where we have reproduced some of their results here. The mass basis weak gauge bosons are related to the original gauge bosons as

$$A_\mu = \sin\theta_W\, W_{\mu L}^3 + \sin\theta_W\, W_{\mu R}^3 + \cos\beta\cos\theta_W\, B_\mu, \quad (10)$$

$$Z_{\mu L} \simeq -\cos\theta_W\, W_{\mu L}^3 + \sin\theta_W\sin\beta\, W_{\mu R}^3 + \sin\theta_W\cos\beta\, B_\mu, \quad (11)$$

$$Z_{\mu R} \simeq -\omega^2\sin^2\beta\cos\beta\, W_{\mu L}^3 - \cos\beta\, W_{\mu L}^3 + \sin\beta\, B_\mu. \quad (12)$$

From equation (10), we have the usual relation for the electromagnetic charge operator $Q = T_L^3 + T_R^3 + Y/2$. The neutral currents coupled to the massive vector bosons Z_L and Z_R are given by

$$\begin{aligned} J_{Z_L}^\mu = \frac{g}{\cos\theta_W}\gamma^\mu\Bigg\{&\left(1-\omega^2\sin^4\beta\right)T_L^3 - \omega^2\sin^2\beta\cos^2\beta\, T_R^3 \\ &- Q\sin^2\theta_W\left(1-\frac{\omega^2\sin^4\beta}{\sin^2\theta_W}\right)\Bigg\}, \end{aligned} \quad (13)$$

$$\begin{aligned} J_{Z_R}^\mu = g\tan\theta_W\tan\beta\,\gamma^\mu\Bigg\{&\left(1+\frac{\omega^2\sin^2\beta\cos^2\beta}{\sin^2\theta_W}\right)T_L^3 \\ &+ \frac{T_R^3}{\sin^2\beta} - Q\left(1+\omega^2\cos^2\beta\sin^2\beta\right)\Bigg\}. \end{aligned} \quad (14)$$

From equation (14), it is clear that the SM fermion's coupling to Z_R is not suppressed by factor ω. This has been used in [5] to calculate the corrections of the SM fermion's coupling Z_L. By fitting this correction to the experimental data, Coutinho et al. have obtained a lower bound for v_R, i.e., $v_R > 30 v_L$ [5].

3. Fermion Sector

The Yukawa terms in the Lagrangian that are invariant under the gauge and the new LRS transformation are

$$\begin{aligned} &-G_e\left(\bar{l}_L\chi_L e_R + \bar{L}_R\tilde{\chi}_R E_L\right) - G_\nu\left(\bar{l}_L\tilde{\chi}_L \nu_R + \bar{L}_R\chi_R N_L\right) \\ &\quad - G_d\left(\bar{q}_L\chi_L d_R + \bar{Q}_R\tilde{\chi}_R D_L\right) - G_u\left(\bar{q}_L\tilde{\chi}_L u_R + \bar{Q}_R\chi_R U_L\right) \\ &\quad - G_{d\nu}\left(\bar{\nu}_R\rho D_L + \bar{N}_L\rho^\dagger d_R\right) - G_{ue}\left(\bar{U}_L\rho^\dagger e_R + \bar{u}_R\rho E_L\right) \\ &\quad - G_{du}\left(\bar{U}_L\rho d_R + \bar{u}_R\rho^\dagger D_L\right) - G_{u\nu}\left(\bar{U}_L\eta^\dagger \nu_R + \bar{u}_R\eta N_L\right) \\ &\quad - G_{dd}\bar{D}_L\eta d_R - M\bar{\nu}_R N_L + \text{h.c.} \end{aligned} \quad (15)$$

Note that the Yukawa couplings G's and the mass M are 3×3 matrices to account for the three generations. After χ_R and χ_L gain nonzero VEV, the above Yukawa terms will generate mass for the charged fermions and mixing mass terms for the neutrinos

$$\begin{aligned} &-G_d\left(v_L\bar{d}_L d_R + v_R\bar{D}_R D_L\right) - G_u\left(v_L\bar{u}_L u_R + v_R\bar{U}_R U_L\right) \\ &\quad - G_e\left(v_L\bar{e}_L e_R + v_R\bar{E}_R E_L\right) - G_\nu\left(v_L\bar{\nu}_L \nu_R + v_R\bar{N}_R N_L\right) \\ &\quad - M\bar{\nu}_R N_L + \text{h.c.} \end{aligned} \quad (16)$$

The mass of the L-R dual fermions (except neutrinos) will be ω^{-1} times the mass of its corresponding SM particles. The current limit on the charged lepton mass is around 100 GeV [15]. This means we should have $v_R > 10^5 v_L$. This bound surpasses the bound obtained by Coutinho et al. [5] from the correction of Z_L coupling to the SM fermion that has been mentioned in the previous section. For the case of neutrinos, in general, the value of M in equation (16) is not restricted by any gauge symmetry mechanism since ν_R, N_L are gauge singlet; thus, we assume that $M \gg \nu_L, \nu_R$. This will lead to a seesaw mechanism [8–10]. The terms in equation (16) related to the neutrino masses can be written as $\bar{\psi}\mathcal{M}\psi$ where $\psi = (\nu_L, N_R, N_L, \nu_R)^T$ and

$$\mathcal{M} = \begin{pmatrix} 0 & 0 & 0 & G_\nu v_L \\ 0 & 0 & G_\nu v_R & 0 \\ 0 & G_\nu v_R & 0 & M \\ G_\nu v_L & 0 & M & 0 \end{pmatrix}, \tag{17}$$

which is similar in form to a Dirac seesaw mechanism case in [16]. From this, it is clear that in our model the neutrinos are Dirac particles. Therefore, any positive result from the neutrinoless double beta decay will rule out this model. The mass matrix in equation (17) can be diagonalized and written in terms of mass basis neutrino state $\Psi = (\nu, \nu', N', N)$, where $\psi = \mathcal{U}\Psi$, with the mixing matrix $\mathcal{U}$ diagonalizing $\mathcal{M}$, i.e., $\mathcal{M} = \mathcal{U}\mathcal{M}_\delta \mathcal{U}^T$. Using the seesaw mechanism, we have

$$\mathcal{U} \simeq \begin{pmatrix} \frac{I}{\sqrt{2}} & \frac{I}{\sqrt{2}} & v_L G_\nu M^{-1} & -v_L G_\nu M^{-1} \\ \frac{-I}{\sqrt{2}} & \frac{I}{\sqrt{2}} & v_R G_\nu M^{-1} & v_R G_\nu M^{-1} \\ -M^{-1} v_L G_\nu^T & -M^{-1} v_R G_\nu^T & \frac{I}{\sqrt{2}} & \frac{-I}{\sqrt{2}} \\ M^{-1} v_L G_\nu^T & -M^{-1} v_R G_\nu^T & \frac{I}{\sqrt{2}} & \frac{I}{\sqrt{2}} \end{pmatrix} \mathcal{V}, \tag{18}$$

where I is a 3×3 identity matrix, and $\mathcal{V}$ is the following block diagonal matrix

$$\mathcal{V} \simeq \begin{pmatrix} V_1 & 0 & 0 & 0 \\ 0 & V_2 & 0 & 0 \\ 0 & 0 & V_3 & 0 \\ 0 & 0 & 0 & V_4 \end{pmatrix}, \tag{19}$$

where the 3×3 matrices V_i $(i = 1, \cdots, 4)$ are the matrix that will diagonalize each block submatrix in the following diagonal matrix $\mathcal{M}_\delta$:

$$\mathcal{M}_\delta \simeq \mathcal{V}^T \begin{pmatrix} v_R v_L G_\nu M^{-1} G_\nu^T & 0 & 0 & 0 \\ 0 & -v_R v_L G_\nu M^{-1} G_\nu^T & 0 & 0 \\ 0 & 0 & M & 0 \\ 0 & 0 & 0 & -M \end{pmatrix} \mathcal{V}. \tag{20}$$

For example, V_1 will diagonalize $v_R v_L G_\nu M^{-1} G_\nu^T$. Because the first two block matrices are the same, then $V_1 = V_2$. Similarly, for the next two block matrices, we have $V_3 = V_4$. From equation (20), we have two different orders for the mass eigenvalues

$$\begin{aligned} m_\nu = m_{\nu'} &\simeq v_L v_R V_1^T G_\nu M^{-1} G_\nu^T V_1, \\ m_N = m_{N'} &\simeq V_3^T M V_3. \end{aligned} \tag{21}$$

The doublet neutrinos (ν_L and N_R) are dominated by lighter neutrinos ν and ν', while the singlet neutrinos (ν_R and N_L) are dominated by heavier neutrinos N and N'. But ν_L, N_R still contain a small portion of heavier component neutrinos, while ν_R, N_L still contain a small portion of lighter component neutrinos. We can have an estimation for the mass order of M. Assuming that the light neutrinos have masses in the range of $\sqrt{\Delta m_{12}^2}$ to $\sqrt{\Delta m_{23}^2}$ as in the neutrino oscillation results [15], then the order of the light neutrino is $m_\nu \simeq 10^{-12} \sim 10^{-11}$ GeV. If we assume that the value of G_ν is of the same order as the Yukawa coupling for charged lepton, i.e., $G_e \simeq 10^{-6} \sim 10^{-3}$, then M should be in the order of $(10^5 \sim 10^{12})\omega^{-1}$ GeV.

There is no oscillation between SM doublet neutrinos and L-R dual doublet neutrinos. This can be seen clearly if we write down the flavor basis neutrinos in terms of the mass basis neutrinos. For SM doublet neutrinos and its L-R duals, we have

$$\begin{aligned} \nu_{L\alpha} \simeq{}& \left(\frac{V_1}{\sqrt{2}}\right)_{\alpha i} \nu_i + \left(\frac{V_2}{\sqrt{2}}\right)_{\alpha i} \nu_i' + v_L \left(G_\nu M^{-1} V_3\right)_{\alpha i} N_i' \\ &- v_L \left(G_\nu M^{-1} V_4\right)_{\alpha i} N_i, \\ N_{R\beta} \simeq{}& -\left(\frac{V_1}{\sqrt{2}}\right)_{\beta j} \nu_j + \left(\frac{\nu_2}{\sqrt{2}}\right)_{\beta j} \nu_j' + v_R \left(G_\nu M^{-1} V_3\right)_{\beta j} N_j' \\ &+ v_R \left(G_\nu M^{-1} V_4\right)_{\beta j} N_j, \end{aligned} \tag{22}$$

where α, β are the flavor index and the sum in i and j is over 1,2,3. From this, the probability amplitude for a SM doublet neutrino $\nu_{L\alpha}$ with energy E which oscillates into an L-R dual doublet neutrino $N_{R\beta}$ after traveling a distance L is

$$\begin{aligned} \langle N_{R\beta} | \nu_{L\alpha} \rangle \simeq{}& -\frac{1}{2} (V_1)_{\beta i} (V_1)^*_{\alpha i} \exp\left(-i \frac{m_{\nu_i}^2}{2E} L\right) \\ &+ \frac{1}{2} (V_2)_{\beta i} (V_2)^*_{\alpha i} \exp\left(-i \frac{m_{\nu_i'}^2}{2E} L\right) \\ &+ v_L v_R \left(G_\nu M^{-1} V_3\right)_{\beta i} \left(G_\nu M^{-1} V_3\right)^*_{\alpha i} \exp\left(-i \frac{m_{N_i'}^2}{2E} L\right) \\ &- v_L v_R \left(G_\nu M^{-1} V_4\right)_{\beta i} \left(G_\nu M^{-1} V_4\right)^*_{\alpha i} \exp\left(-i \frac{m_{N_i}^2}{2E} L\right). \end{aligned} \tag{23}$$

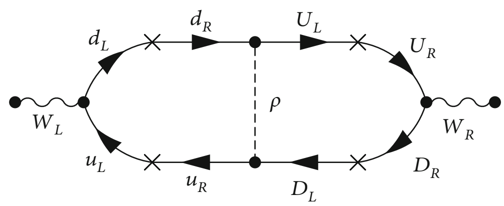

FIGURE 1: The two-loop $W_L - W_R$ mixing.

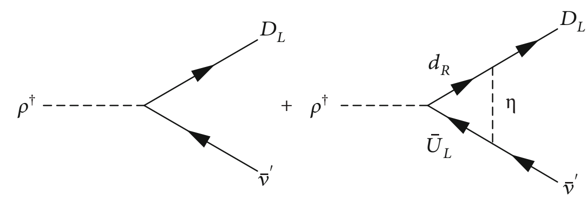

FIGURE 2: Process $\rho^\dagger \rightarrow D_L + \bar{\nu}'$.

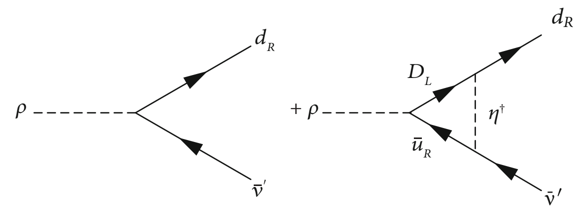

FIGURE 3: Process $\rho \rightarrow d_R + \bar{\nu}'$.

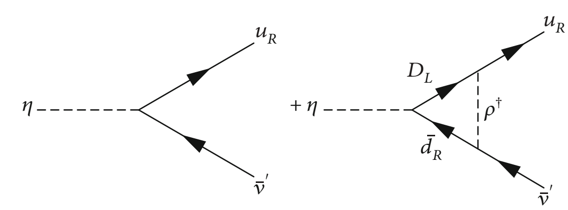

FIGURE 4: Process $\eta \rightarrow u_R + \bar{\nu}'$.

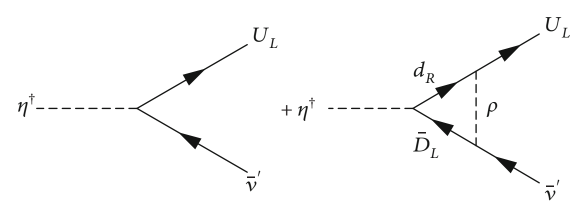

FIGURE 5: Process $\eta^\dagger \rightarrow U_L + \bar{\nu}'$.

It is clear that the above amplitude is approximately zero due to $V_1 = V_2$, $V_3 = V_4$, $m_\nu = m_{\nu'}$, and $m_N = m_{N'}$. The probability amplitude for a SM doublet neutrino $\nu_{L\alpha}$ with energy E which oscillates into a SM doublet neutrino $\nu_{L\beta}$ after traveling a distance L is

$$\begin{aligned}\langle \nu_{L\beta} | \nu_{L\alpha} \rangle &\simeq \frac{1}{2} (V_1)_{\beta i} (V_1)^*_{\alpha i} \exp\left(-i \frac{m^2_{\nu_i}}{2E} L\right) \\ &\quad + \frac{1}{2} (V_2)_{\beta i} (V_2)^*_{\alpha i} \exp\left(-i \frac{m^2_{\nu'_i}}{2E} L\right) \\ &\quad + v_L^2 \left(G_\nu M^{-1} V_3\right)_{\beta i} \left(G_\nu M^{-1} V_3\right)^*_{\alpha i} \exp\left(-i \frac{m^2_{N'_i}}{2E} L\right) \\ &\quad + v_L^2 \left(G_\nu M^{-1} V_4\right)_{\beta i} \left(G_\nu M^{-1} V_4\right)^*_{\alpha i} \exp\left(-i \frac{m^2_{N'_i}}{2E}\right) \\ &\simeq (V_1)_{\beta i} (V_1)^*_{\alpha i} \exp\left(-i \frac{m^2_{\nu_i}}{2E} L\right), \end{aligned} \tag{24}$$

where we have neglected the small terms of order $|v_L G_\nu M^{-1}|^2$. This last result will lead to the usual formulation for the three SM neutrino oscillation. The oscillation probability between doublet neutrinos (either the SM or its L-R dual) and singlet neutrinos is very small, the largest being of order $|v_R G_\nu M^{-1}|^2$.

There is no mixing between the SM and L-R dual fermions in this model; thus, we do not have a $W_L - W_R$ mixing at a one-loop level which usually appears in the conventional LRS models. Instead, the SM and L-R dual particles (leptons and quarks) via the leptoquark ρ can mediate the $W_L - W_R$ mixing at two-loop level. The highest contribution for the two-loop $W_L - W_R$ mixing is shown in Figure 1, where the approximate two-loop mass term is given by

$$\delta M^2_{W_L W_R} \simeq \frac{3g^2 |V_{tb}|^2 G^2_{ud}}{2(16\pi^2)^2} \frac{M_T M_B m_t m_b}{M^2_\rho} \simeq \frac{3g^2 |V_{tb}|^2 G^2_{ud}}{2(16\pi^2)^2} m_b^2, \tag{25}$$

where V is the Cabibbo-Kobayashi-Maskawa matrix [17, 18]. M_T and M_B are the mass of top and bottom L-R dual quarks, respectively. From the above equation, it is clear that if $M_T \ll M_\rho$ then the $W_L - W_R$ mixing is a negligible quantity.

4. The Leptoquarks

The Lagrangian of the model is invariant under global quantum number F so there are no massive Majorana neutrinos whose decay can lead to leptogenesis. But baryon asymmetry can still occur because of baryogenesis from the decay of leptoquarks. The existence of leptoquarks in a model may lead to baryon number violation [19] which may lead to baryogenesis if the Sakharov's conditions [20] are satisfied. The calculation of the baryon number production from the decay of the leptoquarks follows the usual method as in [21], only here we have to consider for the leptoquarks and its L-R duals with the decay results which are the SM and L-R dual particles. The decays of leptoquarks up to one-loop order that can lead to baryogenesis are shown in Figures 2–5. We summarize the value of the baryon number production for those decays in Table 2.

Table 2: The branching for decay processes of the leptoquark particles and their baryon number violations (ΔB).

Particles	Final state	ΔB
$\rho^\dagger$	$\longrightarrow D_L + \bar{\nu}'$	−1/3
ρ	$\longrightarrow \bar{D}_L + \nu'$	+1/3
ρ	$\longrightarrow d_R + \bar{\nu}'$	+1/3
$\rho^\dagger$	$\longrightarrow \bar{d}_R + \nu'$	−1/3
η	$\longrightarrow u_R + \bar{\nu}'$	+1/3
$\eta^\dagger$	$\longrightarrow \bar{u}_R + \nu'$	−1/3
$\eta^\dagger$	$\longrightarrow U_L + \bar{\nu}'$	−1/3
η	$\longrightarrow \bar{U}_L + \nu'$	+1/3

The mean net baryon number produced in the decay of particle X is given by [21]

$$\epsilon_X = \sum_f B_f \frac{\Gamma(X \longrightarrow f) - \Gamma(\bar{X} \longrightarrow \bar{f})}{\Gamma_X}, \tag{26}$$

where the sum $\sum_f$ runs over all final state fermion f, B_f is the produced baryon number of the decay process, and Γ_X is the total decay rate of X. The contribution to ϵ comes from the interference between the lowest (tree-level) order and the one-loop order diagrams, as shown in Figures 2–5. From Figure 2, the decay rate difference between the decay $\rho^\dagger \longrightarrow D_L + \bar{\nu}'$ and $\rho \longrightarrow \bar{D}_L + \nu'$ is

$$\begin{aligned}&\Gamma\left(\rho^\dagger \longrightarrow D_L + \bar{\nu}'\right) - \Gamma\left(\rho \longrightarrow \bar{D}_L + \nu'\right)\\ &\quad = -4\left|\nu_R G_\nu M^{-1}\right|^2 \operatorname{Im}\left(I^{(1)}_{\rho\eta}\right) \operatorname{Im}\left(G_{d\nu} G^*_{du} G_{dd} G^*_{u\nu}\right),\end{aligned} \tag{27}$$

where $(I^{(1)}_{\rho\eta})$ is the kinematic factor of the internal loop in Figure 2 due to the exchange of η in the decay of $\rho^\dagger$. As for the decay rate difference between $\rho \longrightarrow d_R + \bar{\nu}'$ and $\rho^\dagger \longrightarrow \bar{d}_R + \nu'$ (Figure 3), we have

$$\begin{aligned}&\Gamma\left(\rho \longrightarrow d_R + \bar{\nu}'\right) - \Gamma\left(\rho^\dagger \longrightarrow \bar{d}_R + \nu'\right)\\ &\quad = -4\left|\nu_R G_\nu M^{-1}\right|^2 \operatorname{Im}\left(I^{(2)}_{\rho\eta}\right) \operatorname{Im}\left(G^*_{dd} G_{du} G_{u\nu} G^*_{d\nu}\right),\end{aligned} \tag{28}$$

where similar as before, $(I^{(2)}_{\rho\eta})$ is the kinematic factor of the internal loop in Figure 3 due to the exchange of $\eta^\dagger$ in the decay of ρ. Both Im $(I^{(1)}_{\rho\eta})$ and Im $(I^{(2)}_{\rho\eta})$ in equations (27) and (28) have the same value and they are given by

$$\operatorname{Im}\left(I^{(1)}_{\rho\eta}\right) = \operatorname{Im}\left(I^{(2)}_{\rho\eta}\right) = \frac{M_\rho}{128\pi^2}\left\{1 - \sigma^2 \ln\left(1 + \sigma^{-2}\right)\right\}, \tag{29}$$

where $\sigma \coloneqq M_\rho/M_\eta$, while M_ρ and M_η are the mass of the leptoquarks ρ and η, respectively. Substituting equations (27) and (28) into equation (26), we obtain

$$\epsilon_\rho = \frac{1}{3\pi} \frac{\operatorname{Im}\left(G_{d\nu} G^*_{du} G_{dd} G^*_{u\nu}\right)}{\left|G_{d\nu}\right|^2}\left\{1 - \sigma^2 \ln\left(1 + \sigma^{-2}\right)\right\}, \tag{30}$$

where we have used

$$\Gamma_\rho \simeq \left|\nu_R G_\nu M^{-1}\right|^2 \frac{M_\rho \left|G_{d\nu}\right|^2}{8\pi}. \tag{31}$$

Following the same steps above, we obtain the mean net baryon number produced from the decay of η as

$$\epsilon_\eta = \frac{1}{3\pi} \frac{\operatorname{Im}\left(G^*_{dd} G_{du} G_{d\nu} G^*_{u\nu}\right)}{\left|G_{u\nu}\right|^2}\left\{1 - \sigma^{-2} \ln\left(1 + \sigma^2\right)\right\}. \tag{32}$$

Therefore, the total mean net baryon asymmetry is given by

$$\begin{aligned}\epsilon = \frac{\operatorname{Im}\left(G_{d\nu} G^*_{du} G_{dd} G^*_{u\nu}\right)}{3\pi}&\left[\frac{\left\{1 - \sigma^2 \ln\left(1 + 1/\sigma^2\right)\right\}}{\left|G_{d\nu}\right|^2}\right.\\ &\left.- \frac{\left\{1 - \sigma^{-2} \ln\left(1 + \sigma^2\right)\right\}}{\left|G_{u\nu}\right|^2}\right].\end{aligned} \tag{33}$$

For any value of σ, the value in each of the curly bracket above is of order unity. Thus, the order of the total mean net baryon asymmetry is mainly determined by the coupling constant G's which have to be complex to accommodate the C and CP violations. It is important to note that the value of ϵ will vanish if the new LRS symmetry transformation does not reverse the local quantum number of particles (check the value of ΔB in Table 2).

The leptoquark starts to become nonrelativistic at a temperature comparable to their mass $T \simeq M_{lq}$, where M_{lq} is the mass of either ρ or η which we assume to be of the same order. At this temperature, the leptoquark number density is determined primarily by their decay rate Γ_D and the cosmic expansion rate $H \simeq g_*^{1/2} T^2/m_{\mathrm{Pl}}$, where g_* is the effective relativistic degree of freedoms and m_{Pl} is the Planck mass. If at $T \simeq M_{lq}$ the leptoquark decay rate is very small compared to the cosmic expansion rate $\Gamma_D \ll H$, then the leptoquark can become overabundant. From this requirement and the result in equation (31) for the leptoquark total decay rate, we have the lower bound value for the mass of the leptoquark

$$M_{lq} \gg \alpha_{lq} g_*^{-1/2} m_{\mathrm{Pl}} \left|\nu_R G_\nu M^{-1}\right|^2, \tag{34}$$

where $\alpha_{lq} = \left|G_i\right|^2/4\pi\ (i = d\nu, u\nu, dd, du)$. When the leptoquarks start to decay, most of the particles (including the L-R duals) are still relativistic. The leptoquarks contribute $g_{lq} = 2 \times 2 \times 3$ relativistic degree of freedoms (due to ρ and η and their antiparticles, where each has colors) to g_*. We will take the value of g_* around 200 when the leptoquarks decay. The baryon asymmetry as the ratio of the

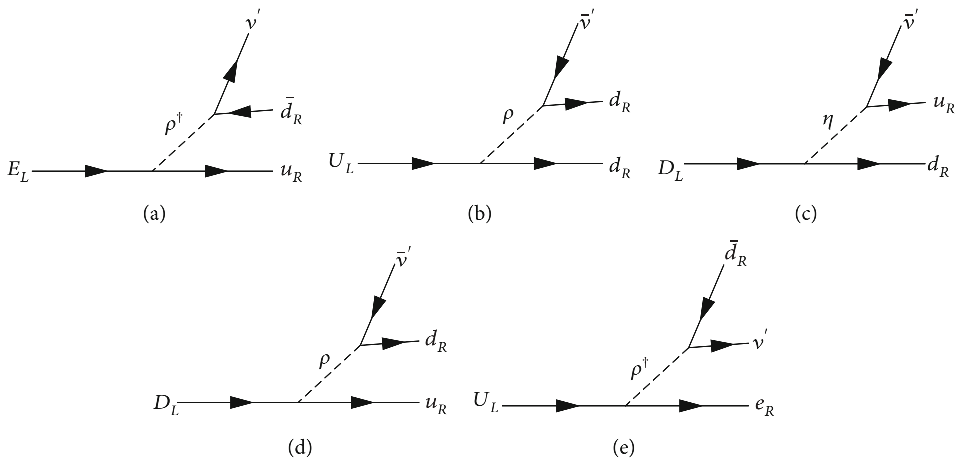

Figure 6: The decay of the L-R dual fermions into the SM fermions via the leptoquarks.

total baryon number to the total entropy density produced by the leptoquark decay is given by

$$B = \frac{n_{lq}(T_M)\,\epsilon}{s(T_M)} = \frac{45\varsigma(3)}{2\pi^4}\left(\frac{g_{lq}}{g_*}\right)\epsilon \simeq 0.016\ \epsilon, \tag{35}$$

where $\zeta(n)$ is the Riemann-Zeta function and $T_M \simeq M_{lq}$. This baryon number will remain as long as there is no more baryon number producing reaction and there is no more entropy production. To get the current observed value of $B \simeq 10^{-10}$ [22], the value of ϵ should be around 10^{-8}. From equation (33), the order of ϵ depends on the imaginary component of $G_{d\nu}G^*_{du}G_{dd}G^*_{u\nu}$ and the absolute value of $|G_{d\nu}|^2$ or $|G_{u\nu}|^2$. These coupling constants may have a very small imaginary part and a larger real part. Taking the real part of G_i's to be of order unity, then the imaginary part should be of order 10^{-8} to get the correct value for ϵ. Using equation (34) and taking $\alpha_{lq}\sim|G_{lq}|^2\sim 1$, $|G_\nu \nu_R|\sim\omega^{-1}$GeV, and $M\sim 10^5\omega^{-1}$ (corresponding to the largest possible decay rate of leptoquark), we get the lower bound value for the leptoquark mass which is around 10^6 GeV.

The leptoquarks are also needed in this model for facilitating the decay of massive L-R dual particles into SM particles. Due to the new LRS transformation, the same CKM matrix pattern should also occur in the L-R dual quark sector; thus, any heavier generation of the L-R dual quark should decay into its lighter generation via right charged gauge bosons. In the lepton sector, any heavier generation of the L-R dual charged leptons should decay into its lighter generation and the lighter neutrinos. Thus, at lower energy, we will end up with the lighter generation of L-R dual quarks, charged leptons, and neutrinos. The lightest generation of the U, D quarks and E can decay into the SM fermions via the leptoquarks. Figure 6 shows the lowest order diagrams for the decay of L-R dual fermions.

At the lowest order, the decay rate of the L-R dual particles can be calculated easily. Assuming the decayed products are massless and the leptoquarks ρ and η are very massive and of the same order, we get the decay rates of the L-R dual particles via leptoquark which are approximately given by

$$\Gamma \simeq \frac{M_D^5|G_{ud}|^2|G_{d\nu}|^2}{M_{lq}^4}\left|\nu_R G_\nu M^{-1}\right|^2, \tag{36}$$

where M_D is the mass of the (lightest generation) L-R dual particles. The mass of the lightest quarks and charged leptons is in the order of 10^{-3} GeV; thus, M_D is in the order of 10^{-3} ω^{-1} GeV. Using the assumption for the neutrinos in the previous section, i.e., $M\sim\omega^{-1}10^5$, together with $\nu_R\sim\omega^{-1}\nu_L$, $G_\nu\nu_L\sim 1$ GeV, and $G_{ud}\sim G_{d\nu}\sim 1$, the lifetime of the L-R dual particles is approximately

$$\tau = \Gamma^{-1} \simeq \omega^5 M_{lq}^4 10^{25}\ \mathrm{GeV}^{-1} \simeq \omega^5 M_{lq}^4\ \mathrm{s}, \tag{37}$$

where M_{lq} is in GeV unit. If the leptoquarks are too massive, the decay of L-R dual particles is very slow and they should still remain in the universe. For the case of $\omega = 10^{-6}$ and $M_{lq} = 10^6$ GeV, we have $\tau \simeq 10^{-6}$ s, which means the L-R dual particles have decayed away at the current age of the universe. If the leptoquarks decayed away before the big bang nucleosynthesis (BBN), which is around several seconds after the big bang, then for the same ω the upper bound for M_{lq} is around $10^7 \sim 10^8$ GeV.

In many models that contain leptoquark, the proton can decay via a tree-level diagram. But in our model, this cannot happen since the leptoquarks always couple to fermion with massive L-R dual fermions. The existence of the L-R dual neutrinos with the same masses as the SM neutrinos clearly will add an additional relativistic degree of freedom during the BBN. This can be avoided if the L-R dual sector is somehow colder than the SM sector. In particular, the temperature ratio between L-R dual neutrino sector and the SM neutrino sector, T'/T, should be small. If the first-generation L-R dual particle (except neutrinos) energy density dominates the universe before decaying, then their decay into ν' and the SM particles will generate a large additional entropy to both sectors. Although ν' has the equal component of ν_L and N_R, their production depends on the strength of the weak interaction in left and right sectors (i.e., it depends on the ratio of M_{W_L} to M_{W_R}). Thus, the ratio of ν_L to N_R production is approximately proportional to the ratio of ω^{-4}, which is very large. Therefore, the entropy generation to the SM sector will be much greater than the one to the L-R dual neutrino sector and $T'/T \ll 1$; thus, the BBN constraint is avoided.

5. Conclusion

A new variant of L-R symmetry model has been proposed, where the L-R transformation is different from that in the original LRS model, i.e., both the L-R chirality and the local quantum number are reversed. The model also has an additional global quantum number F, which is unaffected with the new LRS transformation. The particle contents are doubled, where for each SM particle we have its L-R duals. The global quantum number F will make the L-R dual particles different from the antiparticles of the SM particles. The model does not contain a bidoublet scalar; thus, there is no mixing between left and right weak gauge bosons. The existence of the global quantum number F forbids the model from having any Majorana neutrinos. All neutrinos in the model are Dirac neutrinos, but the seesaw mechanism and the usual SM neutrino oscillation can still occur. Leptogenesis from the decay of Majorana neutrinos certainly is not available in this model, but baryon asymmetry can still be generated due to baryogenesis from the decay of the leptoquarks. Moreover, for baryogenesis, the leptoquarks are also required for the decay of L-R dual particles into SM particles. The decay of L-R dual particles will also produce a large entropy to the SM sector and give a mechanism for avoiding the BBN constraint.

Data Availability

No data were used to support this study.

Conflicts of Interest

The authors declare that they have no conflicts of interest.

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