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Neutrinoless double beta decay process in left-right symmetric models without scalar bidoublet

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We present an alternative formulation of the left-right symmetric theory where the scalar sector consists of two Higgs doublets. This formulation differs from the standard version of the left-right model that makes use of L- and R- Higgs triplets and a Higgs bidoublet. The basic idea is to consider a few extra charged isosinglet fields; the fermion masses can be realized by integrating out these heavy isosinglet fields. We also give a detailed discussion on neutrinoless double beta decay in this particular left-right symmetric theory where the right-handed (RH) Majorana neutrino can be of MeV range. With this RH Majorana mass around the MeV scale, the contribution to neutrinoless double beta decay coming from the right-handed current can be comparable to the contributions coming from the standard left-handed sector only if the right-handed gauge boson mass is around 5 TeV. With this operative scale of W_R around a few TeV, it is possible to probe at the LHC. We have briefly commented on cosmological constraints from the big-bang nucleosynthesis and Universe cosmology on the RH neutrinos involved in this discussion.

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I. INTRODUCTION

The left-right symmetric model (LRSM) is a novel extension of the standard model of particle physics, which will treat the left-handed and right-handed particles on equal footing, and the parity violation we observe at low energies is due to the spontaneous breaking of the left-right symmetry at some high scale [1–6]. The right-handed neutrino is an automatic consequence of left-right symmetric theory; such models provide a natural explanation for the smallness of neutrino masses via a seesaw mechanism [5,7,8]. Another interesting feature of the left-right symmetric model is that the difference between the baryon number (B) and the lepton number (L) becomes a gauge symmetry, which leads to several interesting consequences.

However, until now, the fundamental fermionic representation and the Higgs sector has not been fully determined, which in turn gives a variety of possibilities for choosing these representations (of course, the representations have to be restricted by the symmetry of the known gauge group). In addition, one has to address the issue of origin of the observed fermion masses and mixing. In the standard model (SM), all of the flavor structure is determined by unknown Yukawa couplings. Hence, a new approach to address these issues has been discussed in Refs. [9–12]. In addition to the fermion sector, the basic structure of these models excludes the conventional Higgs triplets and bidoublet, but includes the new left-handed Higgs doublet ϕ_L and the right-handed Higgs doublet ϕ_R; the masses of the usual fermions can be realized by means of a universal seesaw with the aid of a few extra isosinglet fermions. In this paper, we shall follow the simplest approach, which contains a scalar sector with only two Higgs doublets and a few extra isosinglet fermions in order to realize fermion masses and mixings.

Furthermore, experimental observations on solar, atmosphere, reactor, and accelerator neutrino oscillations have revealed that neutrinos can oscillate from one flavor to another as they propagate; this is the strongest indication for nonzero neutrino masses and mixing [13–16]. Moreover, until now, there has been no information about the absolute scale of neutrino masses. One can find the bound on the absolute scale of neutrino masses via studies of lepton number (L) violating neutrinoless double beta decay ($^{A}_{Z}[\mathrm{Nucl}] \to {}^{A}_{Z+2}[\mathrm{Nucl}'] + 2e^-$), the observation of which would imply that neutrinos are Majorana fermions [17]. At present, the best limit on the half-life of this process is $T_{1/2} < 3 \times 10^{25}$ years. This value comes from the Heidelberg-Moscow [18–20] and IGEX [21] collaborations who conducted experiments with ^{76}Ge [22] and, in turn, it translated to a bound on the effective neutrino mass $m_{\mathrm{eff}} \le 0.21$–$0.53$ eV, where the maximum and minimum ranges arise due to the uncertainty in the nuclear matrix elements. Upcoming experiments are trying to improve this bound [23–25].

Along with the standard contribution to $0\nu\beta\beta$ which comes from the exchange of light neutrinos [where the effective Majorana neutrino mass is just the absolute value of the (ee) element of the low energy neutrino mass matrix in the flavor basis], there can be many other contributions to neutrinoless double beta decay in generic left-right (LR) models [26–28]. The importance of RH Majorana neutrinos for neutrinoless double beta decay has been pointed out by Mohapatra [29], while Doi and Kotani [30] gave a detailed discussion of decay rate, including terms for both left-handed and right-handed Majorana neutrinos.

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Recently, a very interesting possibility of left-right symmetry from the LHC to neutrinoless double beta decay [31,32] has been proposed, wherein the scale of left-right symmetry restoration and associated lepton number violation (the neutrinoless double beta decay) can be probed at the LHC. This idea has been discussed in great detail in Ref. [33], where the scale of new physics is at the ~TeV scale, which is phenomenologically rich for the LHC.

Since the aforementioned LR symmetric model without bidoublet offers an appealing possibility that both the light and heavy Majorana neutrino mass matrices are related to each other, it will be worthwhile to study the neutrinoless double decay process in this scenario, including the contributions coming from both left-handed and right-handed sectors. With this motivation, we shall first present the LR models with only isodoublet Higgs Φ_L and Φ_R without having a scalar bidoublet and with detailed discussion. We then extend our discussion to $0\nu\beta\beta$ with a particular emphasis on new contributions coming from the right-handed current.

II. THE MODEL

We now recapitulate the important features of the minimal left-right symmetric model without any scalar bidoublet, where spontaneous parity breaking occurs through only Higgs doublets. This model has been discussed in Refs. [9–12]. At this stage, we shall write the particle content and corresponding Lagrangian for the aforementioned minimal model without invoking any horizontal symmetry, although inclusion of horizontal symmetry is a more complete one. The gauge group of this particular model is $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, where the electric charge is related to the generators of the group as

$$Q = T_{3L} + T_{3R} + \frac{B-L}{2} = T_{3L} + Y. \tag{1}$$

The fermion content of the minimal $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge model is well known, i.e., quarks and leptons transform under the left-right symmetric gauge group as

$$q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix} \equiv \left[3, 2, 1, \frac{1}{3}\right],$$

$$q_R = \begin{pmatrix} u_R \\ d_R \end{pmatrix} \equiv \left[3, 1, 2, \frac{1}{3}\right],$$

$$\ell_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} \equiv [1, 2, 1, -1],$$

$$\ell_R = \begin{pmatrix} N_R \\ e_R \end{pmatrix} \equiv [1, 1, 2, -1].$$

In the generic left-right models, a scalar bidoublet transforming as (1, 2, 2, 0) is introduced for the obvious reason that we want masses for quarks and leptons. Also, a few attempts have been made to explain fermion masses in minimal left-right symmetric models without adding a scalar bidoublet and, in this case, scalar doublets were added to do the job.

The simplest way to achieve this symmetry breaking is to introduce two Higgs doublets which are given below:

$$\Phi_L = (\phi_L^+, \phi_L^0), \qquad \Phi_R = (\phi_R^+, \phi_R^0). \tag{2}$$

Thus, the complete Lagrangian density could be read as

$$\begin{aligned} L = & -\frac{1}{4} W_{\mu\nu L}.W^{\mu\nu L} - \frac{1}{4} W_{\mu\nu R}.W^{\mu\nu R} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} \\ & + \bar{\psi}_L \gamma^\mu \Big(i\partial_\mu - g\frac{1}{2}\tau.W_{\mu L} - g'\frac{Y}{2} B_\mu \Big) \psi_L \\ & + \bar{\psi}_R \gamma^\mu \Big(i\partial_\mu - g\frac{1}{2}\tau.W_{\mu R} - g'\frac{Y}{2} B_\mu \Big) \psi_R \\ & + \left| \Big(i\partial_\mu - g\frac{1}{2}\tau.W_{\mu L} - g'\frac{Y}{2} B_\mu \Big) \Phi_L \right|^2 \\ & + \left| \Big(i\partial_\mu - g\frac{1}{2}\tau.W_{\mu R} - g'\frac{Y}{2} B_\mu \Big) \Phi_R \right|^2 - V(\Phi_L, \Phi_R), \end{aligned} \tag{3}$$

where $g_L = g_R = g$ are the $SU(2)$ couplings, g' is the $U(1)$ coupling, γ^μ are the Dirac matrices, τ's are the Pauli spin matrices, $V(\Phi_L, \Phi_R)$ is the Higgs potential, and Y is the hypercharge ($Y = B - L$). Also, ψ is a fermionic spinner valid for both quarks (q) and leptons (ℓ). Here, the vacuum expectation values of two doublets (v_L and v_R with the relation $v_R \gg v_L$) could contribute to the gauge bosons' masses.

The Higgs sector only consists of a pair of left-right symmetric isodoublets $\Phi_L(2,1,1) \oplus \Phi_R(1,2,1)$ with the following Higgs potential:

$$\begin{aligned} \mathcal{V} = & -(\mu_L^2 \Phi_L^\dagger \Phi_L + \mu_R^2 \Phi_R^\dagger \Phi_R) + \frac{\rho_1}{2}[(\Phi_L^\dagger \Phi_L)^2 \\ & + (\Phi_R^\dagger \Phi_R)^2] + \rho_2^2 (\Phi_L^\dagger \Phi_L)(\Phi_R^\dagger \Phi_R). \end{aligned} \tag{4}$$

The minimum of the potential corresponds to $\langle \Phi_L \rangle = v_L/\sqrt{2}$, $\langle \Phi_R \rangle = v_R/\sqrt{2}$. Choosing $\mu_R \geq \mu_L$ guarantees $v_R \geq v_L$. In the unitary gauge, there are two physical Higgs bosons: $h_L \equiv \mathrm{Re}\Phi_L^0$ and $h_R \equiv \mathrm{Re}\Phi_R$. These two states, in principle, could mix with each other with a mixing angle $\theta_\phi \simeq (\frac{\rho_2}{\rho_1})(\frac{v_L}{v_R})$ for $v_R \gg v_L$. Their masses are given by

$$M_{h_R}^2 \simeq \rho_1 v_R^2, \qquad M_{h_L}^2 \simeq \rho_1 \Big(1 - \frac{\rho_2^2}{\rho_1^2} \Big) v_R^2.$$

A. Gauge boson mass

From Eq. (3), we can see that the relevant gauge boson mass terms are as follows:

$$\begin{aligned} L_{\text{boson}} = & \left| \Big(-g\frac{1}{2}\tau.W_{\mu L} - g'\frac{Y}{2} B_\mu \Big) \Phi_L \right|^2 \\ & + \left| \Big(-g\frac{1}{2}\tau.W_{\mu R} - g'\frac{Y}{2} B_\mu \Big) \Phi_R \right|^2. \end{aligned} \tag{5}$$

After substituting the vacuum expectation values of the Higgs fields,

$$\langle\Phi_L\rangle = \begin{pmatrix} 0 \\ v_L \end{pmatrix}, \qquad \langle\Phi_R\rangle = \begin{pmatrix} 0 \\ v_R \end{pmatrix}, \tag{6}$$

into relation (5), we obtain

$$L_{\rm boson} = \frac{g^2 v_L^2}{4}\{(W^1_{\mu L})^2 + (W^2_{\mu L})^2\} + \frac{v_L^2}{4}(gW^3_{\mu L} - g'B_\mu)^2 + \frac{g^2 v_R^2}{4}\{(W^1_{\mu R})^2 + (W^2_{\mu R})^2\} + \frac{v_R^2}{4}(gW^3_{\mu R} - g'B_\mu)^2. \tag{7}$$

Let us define

$$W^\pm_\alpha = \frac{1}{\sqrt{2}}(W^1_{\mu\alpha} \mp iW^2_{\mu\alpha}), \qquad Z_{\mu\alpha} = \frac{gW^3_{\mu\alpha} - g'B_{\mu\alpha}}{\sqrt{g^2+g'^2}}, \tag{8}$$

$$A_{\mu\alpha} = \frac{g'W^3_{\mu\alpha} + gB_{\mu\alpha}}{\sqrt{g^2+g'^2}}, \qquad Z'_{\mu\alpha} = W^3_{\mu\alpha}, \tag{9}$$

where $\alpha = L, R$. With this definition, the gauge boson mass can read from Eq. (7) as

$$L_{\rm boson} = M^2_{W_L} W^+_L W^-_L + M^2_{W_R} W^+_R W^-_R + M^2_{Z_L} Z_{\mu L} Z^\mu_L + M^2_{Z_R} Z_{\mu R} Z^\mu_R + M^2_A A_\mu A^\mu, \tag{10}$$

where the respective masses that appear in the above Lagrangian are given below:

$$M_{W_L} = \frac{gv_L}{2}, \qquad M_{W_R} = \frac{gv_R}{2}, \qquad M_A = 0, \qquad M_{Z_L} = \frac{v_L\sqrt{g^2+g'^2}}{2}, \qquad M_{Z_R} = \frac{v_R\sqrt{g^2+g'^2}}{2}. \tag{11}$$

B. Fermion mass

We shall discuss here how fermion masses arise in this particular approach. The key idea of the model is to suppose the existence of weak isosinglet heavy fermions in one-to-one correspondence with the light ones. In order to generate the masses of the usual SM fermions, we introduce some heavy charged singlets to construct the Yukawa couplings to the Higgs and fermion doublets so that we can derive the SM Yukawa couplings by integrating out these singlets (see Refs. [9–12]). These heavy isosinglet vector-like fermions include color triplets with electric charge $+2/3$ as $U_{L,R}$, color triplets with electric charge $-1/3$ as $D_{L,R}$, and color singlets with electric charge -1 as $E_{L,R}$. With these extra fields, the Yukawa terms can be written as

$$\begin{aligned}\mathcal{L} \supset & -y_D(\bar q_L\Phi_L D_R + \bar q_R\Phi_R D_L) - M_D\bar D_L D_R \\ & - y_U(\bar q_L\tilde\Phi_L U_R + \bar q_R\tilde\Phi_R U_L) - M_U\bar U_L U_R \\ & - y_E(\bar l_L\Phi_L E_R + \bar l_R\Phi_R E_L) - M_E\bar E_L E_R + \text{H.c.} \\ & \Rightarrow -y_d\bar q_L\Phi_L d_R - y_u\bar q_L\tilde\Phi_L u_R - y_e\bar l_L\Phi_L e_R + \text{H.c.},\end{aligned} \tag{12}$$

where the SM Yukawa couplings are given by

$$y_d = -y^L_D\frac{v_R}{M_D}y^{R\dagger}_D, \tag{13a}$$

$$y_u = -y^L_U\frac{v_R}{M_U}y^{R\dagger}_U, \tag{13b}$$

$$y_e = -y^L_E\frac{v_R}{M_E}y^{R\dagger}_E. \tag{13c}$$

Here we have chosen the base where the mass matrices $M_{D,U,E}$ are real and diagonal.

In the neutrino sector, we consider the left- and right-handed neutral singlets $S_{L,R}$ with the Yukawa couplings and the masses as

$$\mathcal{L} \supset -y_S(\bar\ell_L\tilde\Phi_L S_R + \bar\ell_R\tilde\Phi_R S_L) - M^D_S\bar S_L S_R - \frac{1}{2}M^M_S(\bar S^c_L S_L + \bar S^c_R S_R) + \text{H.c.} \tag{14}$$

At this stage, we do not want the Yukawa couplings $\bar\ell_L\tilde\Phi_L S^c_L$, $\bar\ell_R\tilde\Phi_R S^c_R$ and their CP conjugates. This can be achieved by imposing a discrete symmetry as well as global and local symmetries. For example, let us consider a $U(1)_X$ local symmetry under which $D_{L,R}$, $U^c_{L,R}$, $E_{L,R}$, $S^c_{L,R}$, $\Phi^*_{L,R}$ carry a quantum number $X = 1$. Clearly, this $U(1)_X$ is free of a gauge anomaly. In this context, the Yukawa couplings and the Dirac mass terms in Eqs. (12) and (14) are allowed while the Majorana mass terms in Eq. (14) are forbidden. To break this $U(1)_X$, we can introduce a singlet scalar η with Yukawa couplings to the neutral singlets $S_{L,R}$,

$$\mathcal{L} \supset -\frac{1}{2}f_S(\eta\bar S^c_L S_L + \eta^*\bar S^c_R S_R) + \text{H.c.} \tag{15}$$

Through the above Yukawa interactions, the Majorana masses in Eq. (14) can be given by

$$M^M_S = f_S\langle\eta\rangle. \tag{16}$$

By integrating out the neutral singlets, the full neutrino masses would contain a Dirac mass term and two Majorana ones:

$$\mathcal{L} \supset -\frac{1}{2}\bar\nu_L M_L\nu^c_L - \frac{1}{2}\bar N_R M_R N^c_R - \bar\nu_L M_D N_R + \text{H.c.}, \tag{17}$$

with

$$M_L = -y_S\frac{1}{M^M_S}y^T_S v_L^2, \tag{18a}$$

$$M_R = -y_S\frac{1}{M^M_S}y^T_S v_R^2, \tag{18b}$$

$$M_D = y_S\frac{1}{M^M_S}(M^D_S)^T\frac{1}{M^M_S}y^\dagger_S v_L v_R. \tag{18c}$$

Here we have assumed

$$M^M_S \gg M^D_S, y_S v_R, y_S v_L, \tag{19}$$

by choosing the base where the Majorana mass matrix M^M_N is real and diagonal:

$$M_S^M = \mathrm{diag}\{M_1, M_2, M_3\} \simeq M. \tag{20}$$

Clearly, the right-handed neutrinos will give their left-handed partners an additional Majorana mass term through the seesaw, since their Dirac masses are not vanishing. This contribution is indeed negligible:

$$\delta M_L = -M_D \frac{1}{M_R^\dagger} M_D^T = \mathcal{O}\left[\left(\frac{M_S^D}{M_S^M}\right)^2\right] M_L \ll M_L. \tag{21}$$

Therefore, we can well define the left- and right-handed Majorana neutrinos:

$$\nu = \nu_L + \nu_L^c, \tag{22a}$$

$$N = N_R + N_R^c. \tag{22b}$$

Diagonalization of the light neutrino mass matrix $m_\nu = M_L$, through the lepton flavor mixing matrix U_{PMNS} [34] gives us three light Majorana neutrinos $m_{\mathrm{light}}^{\mathrm{diag}} = U_{\mathrm{PMNS}} M_L U_{\mathrm{PMNS}}^T = \mathrm{diag}\{m_1, m_2, m_3\}$. If we look at the structure of the light neutrino mass matrix M_L and heavy neutrino mass matrix M_N, then it is clear that both matrices can be simultaneously diagonalized by the same unitary matrix U_{PMNS}, i.e., $M_{\mathrm{heavy}}^{\mathrm{diag}} = U_{\mathrm{PMNS}} M_N U_{\mathrm{PMNS}}^T \frac{v_R^2}{v_L^2} = \mathrm{diag}\{M_1, M_2, M_3\}$. Hence, one can correlate the eigenvalues of the light and heavy Majorana neutrinos, which in turn gives $m_\nu \propto M_N$.

In other words, one can write the light left-handed and heavy right-handed Majorana neutrino mass matrices in terms of the diagonal eigenvalues of light neutrinos as

$$m_\nu = M_L = U_{\mathrm{PMNS}}^\dagger \mathrm{diag}\{m_1, m_2, m_3\} U_{\mathrm{PMNS}}^*,$$

$$M_N = M_R = U_{\mathrm{PMNS}}^\dagger \mathrm{diag}\{m_1, m_2, m_3\} U_{\mathrm{PMNS}}^* \frac{v_R^2}{v_L^2},$$

where m_1, m_2, and m_3 are the absolute masses of light Majorana neutrinos and are chosen to be real.

III. NEUTRINOLESS DOUBLE BETA DECAY

In this section, we shall present the lepton number violating processes such as neutrinoless double beta decay in the left-right symmetric model without having a scalar bidoublet. We shall examine how the $0\nu\beta\beta$ is controlled by heavy Majorana neutrinos having mass around 1–10 MeV. If left-right symmetry exists at high energy, then the contribution of the right-handed current is expected at low energy from the exchange of right-handed weak W_R bosons. The Feynman diagrams that give rise to neutrinoless double beta decay are depicted in Fig. 1.

The corresponding Feynman amplitude for these above diagrams is depicted in Table I.

In this table, $G_F = 1.2 \times 10^{-5}\ \mathrm{GeV}^{-2}$ is the Fermi constant, M_{W_R} is the right-handed charged gauge boson mass, ζ_{L-R} is the $W_L - W_R$ mixing, and p^2 is the neutrino virtuality. In order to estimate the relative contributions of different terms, it is worthwhile to note here that we shall analyze the effect of neutrinoless double beta decay

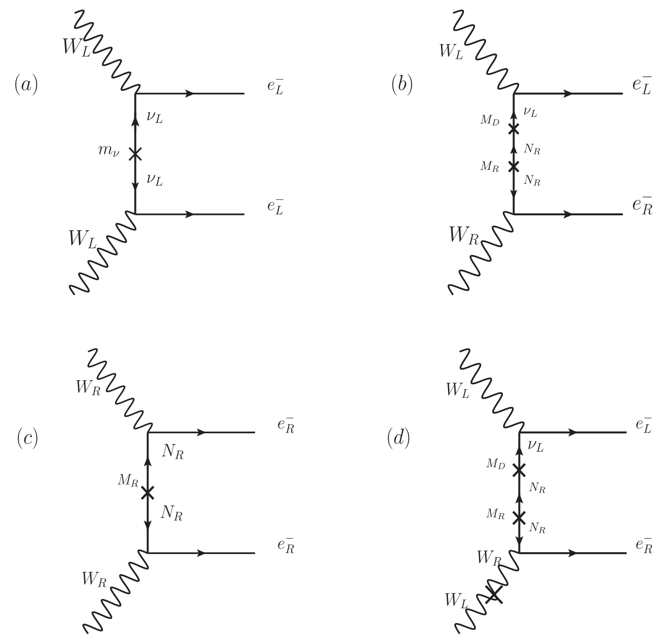

FIG. 1. The Feynman diagrams leading to neutrinoless double beta decay in the presence of a right-handed current. The nucleon part that couples to W bosons is omitted here. (a) The standard process via light left-handed neutrino exchange. (b)–(d) Those involving a right-handed current.

while the representative set of parameters in this model is $M_{W_R} \sim 10$ TeV and the heaviest right-handed neutrino mass is ~1–10 MeV. With this set of parameters, the relevant dominant contributions are found to be

$$\mathcal{A}_a \propto \frac{G_F^2}{p^2}(U_{ei}^2 m_{\nu i}) \sim \frac{G_F^2}{p^2} \times 10^{-2}\ \mathrm{eV},$$

$$\mathcal{A}_c \propto \frac{G_F^2}{p^2} 10^{-8} 10^7\ \mathrm{eV} \sim \frac{G_F^2}{p^2} \times 10^{-1}\ \mathrm{eV}.$$

A. The standard contribution from the left-handed current

The generic contribution to total decay width for neutrinoless double beta decay ($0\nu\beta\beta$), which comes from the left-handed light neutrinos as an exchange particle, is given as

TABLE I. Analytic formulas for amplitudes for different Feynman diagrams in neutrinoless double beta decay process as described in the text.

Feynman diagrams	Amplitude
Fig. 1(a)	$\mathcal{A}_a \propto G_F^2 \frac{U_{ei}^2 m_{\nu i}}{p^2}$
Fig. 1(b)	$\mathcal{A}_b \propto G_F^2 (\frac{M_{W_L}^2}{M_{W_R}^2}) U_{ei}^2 (\frac{M_D}{M_R}) \frac{1}{\lvert p \rvert}$
Fig. 1(c)	$\mathcal{A}_c \propto G_F^2 (\frac{M_{W_L}^4}{M_{W_R}^4}) U_{ei}^2 \frac{M_{Ri}}{p^2}$
Fig. 1(d)	$\mathcal{A}_d \propto G_F^2 U_{ei}^2 (\frac{M_D}{M_R}) \zeta_{L-R} \frac{1}{\lvert p \rvert}$

$$\Gamma_{0\nu} = G^{0\nu} \left| \frac{\mathcal{M}^{0\nu}}{m_e} \right|^2 |M_\nu^{ee}|^2, \tag{23}$$

where $G^{0\nu}$ is a phase space factor, m_e is the electron mass, $\mathcal{M}^{0\nu}$ is the nuclear matrix element, and the effective Majorana mass is given by

$$|M_\nu^{ee}| = |U_{ej}^2 m_j|. \tag{24}$$

Here U_{ej} are the elements of the lepton mixing matrix $U_{\rm PMNS}$ given in Ref. [34] which contains three mixing angles and three phases (one Dirac and two Majorana phases). It is worthwhile to emphasize here that the neutrinoless double beta decay experiment can probe the phases which crucially depend on the pattern of the neutrino masses, i.e., whether neutrinos are normal, inverted, or quasidegenerate and on the magnitude of the neutrino masses. One can parametrize the effective Majorana mass in terms of the elements of $U_{\rm PMNS}$ and mass eigenvalues as

$$|M_\nu^{ee}| = |\cos^2_{12}\cos^2_{13} m_1 + e^{2i\alpha_2}\sin^2_{12}\cos^2_{13} m_2 + e^{2i\alpha_3}\sin^2_{13} m_3|. \tag{25}$$

This contribution of effective Majorana mass is depicted in Fig. 2, which gives the value of the effective Majorana mass as a function of lightest neutrino mass. To generate the required plot, we have used the 3σ ranges, the best-fit values of the mass squared differences, the mixing angles $\sin^2\theta_{12}$, $\sin^2\theta_{23}$ from a global analysis of oscillation data [35], and the value of $\sin^2\theta_{13}$ from the recent measurement of a Daya Bay experiment [36]. In particular, the representative values of the parameters which have been taken in this model, in order to give the result shown in Fig. 2, are as follows:

$\Delta m^2_{\rm sol}[10^{-5}\ {\rm eV}^2]$	7.58[best-fit]	6.99–8.18[3-sigma]
$\|\Delta m^2_{\rm atm}\|[10^{-3}\ {\rm eV}^2]\|$	2.35[best-fit]	2.06–2.67[3-sigma]
$\sin^2\theta_{12}$	0.306[best-fit]	0.259–0.359[3-sigma]
$\sin^2\theta_{23}$	0.42[best-fit]	0.34–0.64[3-sigma]
$\sin^2\theta_{13}$	0.023[best-fit]	0.009–0.037[3-sigma].

In the plot, we need to explain how an effective Majorana mass probes which kind of mass pattern of neutrinos. As shown in Fig. 2, the cyan band for the normal hierarchy (NH) corresponds to varying the parameters in their 3σ range, whereas the red band corresponds to the best-fit parameters where the $\sin^2\theta_{12}$ values are taken from a recent Daya Bay result. In both figures, the Majorana phases are varied between 0 to 2π. In the same manner, the green band for the inverted hierarchy (IH) corresponds to varying the parameters in their 3σ range, whereas the blue band corresponds to the best-fit parameters. We will not present the detailed analysis of this figure since this has already been discussed elaborately in Ref. [33]. We shall now move on to the next subsection, where the dominant

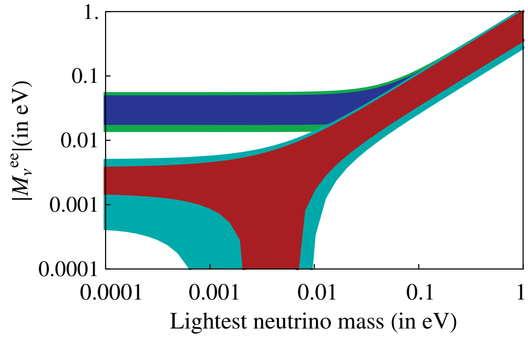

FIG. 2 (color online). The generic contribution from a light neutrino mass with θ_{13} from Ref. [36] to the neutrinoless double beta decay. Here the top (blue and green) horizontal bands are for inverted neutrino mass hierarchy, and the bottom (red and cyan) bands are for hierarchical neutrino masses.

contribution comes from the right-handed current, and present an analysis for the result obtained with a MeV mass range of RH Majorana neutrinos.

B. New contribution from the right-handed current

From the discussion of the light and heavy Majorana neutrino masses which is stated in the end of Sec. II, it is found that they are related to each other as $m_j \propto M_j$, where the proportionality factor is v_R^2/v_L^2. Before relating heavy RH neutrinos in terms of light neutrino masses, we will first present the different hierarchy patterns of the light neutrinos as follows:

(i) In the case of the NH, the light neutrino masses m_2 and m_3 can be expressed in terms of the lightest light neutrino mass m_1 as

$$m_2 = \sqrt{m_1^2 + \Delta m^2_{\rm sol}},$$
$$m_3 = \sqrt{m_1^2 + \Delta m^2_{\rm atm} + \Delta m^2_{\rm sol}},$$

and their mass hierarchy is $m_1 < m_2 \ll m_3$.

(ii) The IH implies $m_3 \ll m_1 \sim m_2$ and the light neutrino masses m_1 and m_2 can be written in terms of the lightest light neutrino mass, which is m_3 in this case, as

$$m_1 = \sqrt{m_3^2 + \Delta m^2_{\rm atm}},$$
$$m_2 = \sqrt{m_3^2 + \Delta m^2_{\rm sol} + \Delta m^2_{\rm atm}}.$$

(iii) The quasidegenerate limit corresponds to $m_1 \approx m_2 \approx m_3 \gg \sqrt{\Delta m^2_{\rm atm}}$.

In the following, we will present the relation between heavy right-handed neutrino masses in terms of light left-handed neutrinos for various mass spectra and try to analyze the behavior of effective Majorana mass M_N^{ee} as a function of the lightest light left-handed neutrinos.

1. Hierarchical pattern of the neutrino masses

It is important to mention here that the value of M_{W_R} has to be at least 10 TeV in order to get a MeV scale of the heaviest RH neutrino mass so that the new contributions to neutrinoless double beta decay coming from the right-handed current can be comparable. In presenting the analytical behavior of the neutrinoless double beta decay contribution from the right-handed current, one should first give the heavy right-handed (RH) neutrino mass ratios to those of light neutrinos, which are given below:

$$\frac{M_1}{M_3} = \frac{m_1}{m_3}, \quad \text{and} \quad \frac{M_2}{M_3} = \frac{m_2}{m_3},$$

where the value of the heaviest RH neutrino mass M_3 is fixed around the MeV range. With this input, the expression for M_N^{ee} is given by

$$|M_N^{ee}|_{NH} = \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \sum_j U_{ej}^2 M_j = \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 M_3 \left| \cos^2\theta_{12}\cos^2\theta_{13}\frac{m_1}{m_3} + \sin^2\theta_{12}\cos^2\theta_{13}e^{2i\alpha_2}\frac{m_2}{m_3} + \sin^2\theta_{13}e^{2i\alpha_3} \right|. \quad (26)$$

In the purely hierarchical case, 10^{-5} eV $< m_1 <$ 10^{-3} eV, one can write $m_2 \simeq \sqrt{\Delta m_{\rm sol}^2}$, $m_3 \simeq \sqrt{\Delta m_{\rm atm}^2}$. Given the input parameters in our model, the ratio between left- and right-handed charged gauge boson masses is found to be 10^{-8}, and the ratio between the solar and atmospheric mass squared differences is $m_2/m_3 \simeq \sqrt{\Delta m_{\rm sol}^2/\Delta m_{\rm atm}^2} = \{0.16, 0.2\}$, corresponding to minimum and maximum values, respectively. Since m_1 is very small, the first term in Eq. (26) gives a negligible contribution and hence can be neglected. With the choice made for M_3 at the 5 MeV scale, the effective Majorana mass is

$$|M_N^{ee}|_{NH} = 0.05 \left| \sin^2\theta_{12}\cos^2\theta_{13}\sqrt{\frac{\Delta m_{\rm sol}^2}{\Delta m_{\rm atm}^2}}e^{2i\alpha_2} + \sin^2\theta_{13}e^{2i\alpha_3} \right|. \quad (27)$$

The maximum and minimum values of $|M_N^{ee}|_{NH}$ correspond to the phase values α_2, $\alpha_3 = 0$, π; $\alpha = 0$, π; and $\alpha_3 = \pi/2$, respectively. As shown in Fig. 3, the blue band in the regime 10^{-5} eV $< m_1 < 10^{-3}$ eV corresponds to the minimum and maximum values as follows:

$$\begin{cases} |M_N^{ee}|_{NH}(\max) = 0.0075 \\ |M_N^{ee}|_{NH}(\min) = 0.0055 \end{cases}.$$

For intermediate hierarchical values of m_1 (say, 10^{-3} eV $< m_1 < 10^{-2}$ eV), the first term in Eq. (26) can still be neglected. For illustration, one can see that the first term of Eq. (27) is small because of the smallness of $\sqrt{\Delta m_{\rm sol}^2/\Delta m_{\rm atm}^2}$; at the same time, the second term is also suppressed due to the factor $\sin^2\theta_{13}$. As a result, cancellation occurs in this regime due to the relative phase cancellation of α_2 and α_3.

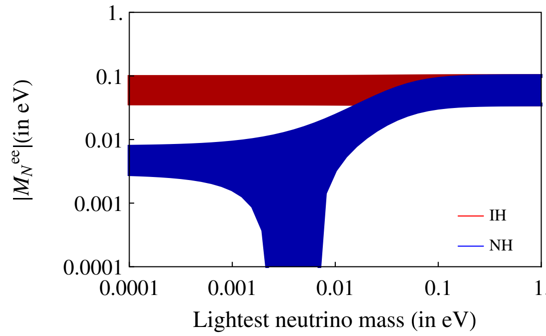

FIG. 3 (color online). The new dominant contribution to neutrinoless double beta decay coming from the right-handed current having M_j around MeV and right-handed W bosons around 10 TeV. Here the upper (red) band is for inverted hierarchical and the lower (blue) band is for hierarchical light neutrino masses.

2. Inverted hierarchy of the neutrino masses

In this case, the other heavy RH neutrino masses can be expressed in terms of light neutrino masses (keeping M_2 fixed, which is the heaviest RH neutrino mass) as

$$\frac{M_1}{M_2} = \frac{m_1}{m_2}, \quad \text{and} \quad \frac{M_3}{M_2} = \frac{m_3}{m_2}.$$

Now the expression for M_N^{ee} becomes

$$|M_N^{ee}|_{IH} = \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 M_2 \left| \cos^2\theta_{12}\cos^2\theta_{13}\frac{m_1}{m_2} + \sin^2\theta_{12}\cos^2\theta_{13}e^{2i\alpha_2} + \sin^2\theta_{13}e^{2i\alpha_3}\frac{m_1}{m_2} \right|. \quad (28)$$

Before illustrating the analytical behavior of this contribution, it should be noted here that the value of m_3 in the case of the inverted hierarchy is such that $m_3 \ll \sqrt{\Delta m_{\rm sol}^2}$, $m_1 \simeq \sqrt{\Delta m_{\rm atm}^2}$, and $m_2 \simeq \sqrt{\Delta m_{\rm sol}^2}$. Since the factor m_3/m_2 is very small in this regime and the value of $\sin^2\theta_{13}$ is also very small, the last term of Eq. (28) can be safely neglected. Now the effective Majorana mass in this inverted hierarchical scheme is

$$|M_N^{ee}|_{IH} = 0.05 \left| \cos^2\theta_{12}\cos^2\theta_{13}\sqrt{\Delta m_{\rm atm}^2/\Delta m_{\rm sol}^2} + \sin^2\theta_{12}\cos^2\theta_{13}e^{2i\alpha_2} \right|. \quad (29)$$

Similarly, the same arguments discussed in the above subsection give the maximum and minimum values of the $|M_N^{ee}|_{IH}$ as

$$\begin{cases} |M_N^{ee}|_{NH}(\max) = 0.1 \\ |M_N^{ee}|_{NH}(\min) = 0.05 \end{cases}.$$

3. Quasidegenerate pattern of the neutrino masses

In this limit, $m_1 \sim m_2 \sim m_3 \sim m_0$, which implies

$$m_0 \gg \sqrt{\Delta m^2_{\rm sol}}, \sqrt{\Delta m^2_{\rm atm}}.$$

The quasidegenerate pattern of light neutrino masses also implies quasidegeneracy in the heavy right-handed neutrino sector, which implies

$$M_1 \approx M_2 \approx M_3 = M_0,$$

where M_0 is the common absolute mass of heavy RH neutrinos at the MeV scale. In this situation, one can write the relation for the heavy neutrino contribution to the effective mass as

$$|M_N^{ee}|_{\rm QD} = \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 M_0 \left| \cos^2\theta_{12}\cos^2\theta_{13} + \sin^2\theta_{12}\cos^2\theta_{13}e^{2i\alpha_2} + \sin^2\theta_{13}e^{2i\alpha_3} \right|.$$

From this relation, we can conclude that the effective neutrino mass from the RH current is independent of the lightest neutrino mass in the quasidegenerate limit. In other words, the value of $|M_N^{ee}|$ remains constant with increasing m_1.

C. Total contribution

The total dominant contribution to neutrinoless double beta decay in the left-right model, in which the scalar sector consists of two isodoublets Φ_L and Φ_R without having a bidoublet, is given by

$$\Gamma_{0\nu} = G^{0\nu} \cdot \left| \frac{\mathcal{M}^{0\nu}}{m_e} \right|^2 |m_{\rm eff}^{\rm ee}|^2. \qquad (30)$$

The effective neutrino mass contribution to neutrinoless double beta decay is

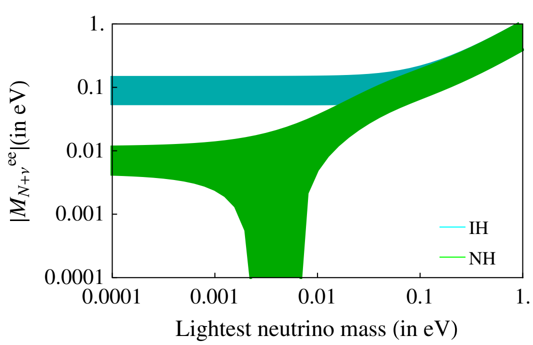

FIG. 4 (color online). The total contribution to neutrinoless double beta decay in left-right models without having a scalar bidoublet. Here the upper (cyan) band and the lower (green) band correspond to inverted hierarchy and normal hierarchy of the light neutrino masses respectively.

$$|m_{\rm eff}^{\rm ee}|^2 = \left(|U_{ej}^2 m_j|^2 + \left|\frac{M_{W_L}^4}{M_{W_R}^4} U_{ej}^2 M_{N_j}\right|^2\right) = |M_\nu^{ee}|^2 + |M_N^{ee}|^2, \qquad (31)$$

where the individual contributions are $M_\nu^{ee} = U_{ej}^2 m_j$ and $M_N^{ee} = \frac{M_{W_L}^4}{M_{W_R}^4} U_{ej}^2 M_{N_j}$. This combined contribution is illustrated in Fig. 4.

IV. COMMENTS ON COSMOLOGICAL CONSTRAINTS

We shall discuss in this section whether the MeV scale RH neutrinos for M_{W_R} lying in the 1–10 TeV region is consistent with the big-bang nucleosynthesis (BBN) bound and from the overclosing of the Universe. We are in a problematic situation when M_{W_R} lies around the TeV scale, which in turn gives an overabundance of RH neutrinos N, because the $SU(2)_R$ gauge interaction keeps them in thermal equilibrium when the temperature is high. Also, if RH neutrinos are allowed to decay later than late after the BBN era, they end up destroying the abundance of light elements, which in turn gives $\tau_N \lesssim$ sec, that translates into a lower bound on M_N.

Let us first consider the case of the heavy regime, with $M_N \gtrsim m_\pi + m_\ell$, where N decays sufficiently fast into a charged (anti)lepton and a pion with the following decay rate:

$$\Gamma_{N\to\ell\pi} = \frac{G_F^2 |V_{ud}^{qR}|^2 |V_{\ell N}^R|^2 f_\pi^2 M_N^3}{8\pi} \frac{M_W^4}{M_{W_R}^4} [(1-x_\ell^2)^2 - x_\pi^2(1+x_\ell^2)][(1-(x_\pi+x_\ell)^2)(1-(x_\pi-x_\ell)^2)]^{\frac{1}{2}}, \qquad (32)$$

where $x_{\pi,\ell} = m_{\pi,\ell}/M_N$, V^R is the right-handed lepton mixing matrix, V^{qR} is the analog quark one, and $f_\pi = 130$ MeV is the pion decay constant. We recall that $V_{ud}^{qR} \simeq V_{ud}^{qL} \simeq 0.97$; on the other hand, the leptonic mixing involved depends on the mass hierarchy and on the flavor of the charged lepton into which the RH neutrino is decaying. As one can check from (32), for $M_N > m_\pi + m_\ell$, the above process guarantees that τ_N is safely shorter than a second. Hence, the constraints coming from the cosmology give $M_N > 140$ MeV. This range of RH neutrino masses will push the M_{W_R} scale beyond the TeV scale, which spoils a possible probe of our scenario in the near future at the LHC.

The prescribed scenario discussed above suffers from a serious problem when M_N lies below <140 MeV: the lifetime becomes longer than a second and a decaying N would pump too much entropy into the Universe. The point is that they decouple relativistically at the temperature

$$T_D^N = T_D^\nu \left(\frac{M_{W_R}}{M_W}\right)^{\frac{4}{3}}, \qquad (33)$$

where $T_D^\nu \simeq 1$ MeV is the neutrino decoupling temperature. Therefore, for a representative value of $M_{W_R} \sim 5$ TeV,

$$T_D^N \simeq 250 \text{ MeV}. \tag{34}$$

Then, since between T_D^N and 1 MeV only muons and pions decouple, at BBN N's are almost equally as abundant as light neutrinos. The only way out would be to make N stable and to avoid the overclosure of the Universe, lighter than about eV [37,38]. As a result, we are in a scenario where extra species are contributing to BBN. Actually, this situation seems to be preferred and a recent study suggests [39,40] that four light neutrinos give the best fit to cosmological data, while five is disfavored and six is basically excluded.

V. CONCLUSION

We have discussed neutrinoless double beta decay in the context of left-right symmetric models with the minimal Higgs content, which are different from the standard version of the LR model that makes use of L- and R- Higgs triplets and a Higgs bidoublet for the fermion mass generation. The scalar sector model consists of two Higgs doublets Φ_L and Φ_R without invoking triplets and a bidoublet, and the fermion masses are generated by integrating out the extra vectorlike heavy quarks and leptons. In the gauge sector, there is no mixing between left- and right-handed weak gauge bosons at tree level, but it can induced at one loop level. In this particular scenario, where the light neutrino and heavy Majorana neutrino are related to each other and diagonalized by the same Pontecorvo-Maki-Nakagawa-Sakata matrix, the neutrinoless double beta decay receives important contributions from the right-handed current. In fact, given the choice we made for the right-handed Majorana neutrino mass around the MeV range in the context of neutrinoless double beta decay, the last formula of Sec. II predicts the right-handed gauge boson mass M_{W_R} to be at least of the order of 5 TeV.

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