0νββ 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 which differs from the standard version of the left-right model that makes use of $L$– and $R$– Higgs triplets and a Higgs bi-doublet. The basic idea is to consider few extra charged iso-singlet fields and 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 Majorana neutrino can be of MeV range. With this right-handed Majorana mass around MeV scale, the contribution to neutrinoless double beta decay coming from the right-handed current can be comparable with the contributions coming from the standard left-handed sector only if the right-handed gauge boson mass is around 5 TeV, and with this operative scale of $W_R$ around few TeV, it is possible to probe at LHC. We have briefly commented on cosmological constraints coming from the big-bang nucleosynthesis and Universe cosmology to the right-handed neutrinos involved in this discussion.

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INTRODUCTION

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 would be due to the spontaneous breaking of the left-right symmetry at some high scale [1–3]. 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 see-saw mechanism [3, 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, till now, the fundamental fermionic representation and the Higgs sector is not fully determined which in turn gives a variety of possibilities of choosing these representation (of course the representations has 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 the flavor structure is determined by unknown Yukawa couplings. Hence, a new approach to address these issues has been discussed in Ref. [9–12]. The basic structure of these models excludes the conventional Higgs triplets and bidoublet, but includes the new left-handed Higgs doublet $\Phi_L$ and the right-handed Higgs doublet $\Phi_R$ and the masses of the usual fermions can be realized by means of a universal seesaw with the aid of few extra isosinglet fermions. In this paper, we shall follow a simplest approach which contains scalar sector with only two Higgs doublets and few extra iso-singlet fermions in order to realize fermion masses and mixings.

Furthermore, the experimental observation on solar, atmospheric, reactor and accelerator neutrino oscillations have revealed that neutrinos can oscillate from one flavor to another as they propagate is the strongest indication for nonzero neutrino masses and mixing [13–16]. Moreover, until now there is no information about the absolute scale of neutrino masses. One can find the bound on absolute scale of neutrino mass via studies of lepton number ($L$) violating neutrino less double $\beta$-decay ($^2\beta [\text{Nucl}] \rightarrow Z^0 + A [\text{Nucl}'] + 2e^-$), whose observation would imply that neutrinos are Majorana fermions [17]. At present days, the best limit on the half life of this process is $T_{1/2} \leq 3 \times 10^{25}$ years coming from the Heidelberg-Moscow [18, 19] and IGEX [20], collaborations conducted experiments with $^{76}\text{Ge}$, [21] which in turn translated to a bound on the effective neutrino mass $m_{\text{eff}} \leq 0.21 – 0.53 \text{eV}$, where the maximum and minimum range arises due to the uncertainty in the nuclear matrix elements. In addition to this bound, there are other upcoming experiments trying to improve this bound [22–24].

Along with the standard contribution to $0\nu\beta\beta$ which comes through 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 flavour basis), there can be many other contribution to neutrinoless double beta in generic left-right (LR) models [25, 26]. The importance of RH Majorana neutrinos for neutrinoless double beta decay has been pointed out by Mohapatra [27] while Doi and Kotani [28] gave a detailed discussion of decay rate including terms for both left-handed and right-handed Majorana neutrinos. Recently, a very interesting possibility of “left-right symmetry: from LHC to neutrinoless double beta” [29, 30] 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 LHC. Then this idea has been discussed in great detailed in Ref. [31] where the scale of new physics is at $\sim \text{TeV}$ scale which is phenomenologically rich for 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 with each other and hence, it will be worth to study the neutrino...
no less double decay process in this scenario including both the contributions coming from left-handed as well as right-handed sector. With this motivation, we shall first present the LR models with only isodoublets Higgs $\Phi_L$ and $\Phi_R$ without having a scalar bidoublet with detailed discussion. We then extend our discussion to $\theta\nu/\beta\beta$ with particular emphasis on new contribution coming from right-handed current.

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 which has been discussed in Ref. [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 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:

$$Q = T_{3L} + T_{3R} + \frac{B - L}{2} = T_{3L} + Y.$$  \hspace{1cm} (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:

$$q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix} \equiv [3, 2, 1, \frac{1}{3}], \quad q_R = \begin{pmatrix} u_R \\ d_R \end{pmatrix} \equiv [3, 1, 2, -\frac{1}{3}],$$

$$\ell_L = \begin{pmatrix} \nu_L \\ \ell_L \end{pmatrix} \equiv [1, 2, 1, -1], \quad \ell_R = \begin{pmatrix} \nu_R \\ \ell_R \end{pmatrix} \equiv [1, 1, 2, -1].$$

In the generic left-right models, a scalar bidoublet transforming $(1, 2, 2, 0)$ is introduced for obvious reason that we want masses for quarks and leptons. Also there are few attempts has been made in order to explain fermion masses in minimal left-right symmetric models without adding 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

$$\Phi_L = (\phi^+_L, \phi^0_L), \quad \Phi_R = (\phi^+_R, \phi^0_R) \hspace{1cm} (2)$$

Thus, the complete Lagrangian density could be read as:

$$L = -\frac{1}{4} W_{\mu
u L} W^{\mu
u L} - \frac{1}{4} W_{\mu
u R} W^{\mu
u R} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} + \bar{\psi}_L \gamma^\mu \left( i\partial_\mu - g_2/2 \tau^i W_{\mu L} - g Y/2 B_\mu \right) \psi_L$$

$$+ \bar{\psi}_R \gamma^\mu \left( i\partial_\mu - g_2/2 \tau^i W_{\mu R} - g Y/2 B_\mu \right) \psi_R$$

$$+ \left( i\partial_\mu - g_2/2 \tau^i W_{\mu L} - g Y/2 B_\mu \right) \Phi_L^2$$

$$+ \left( i\partial_\mu - g_2/2 \tau^i W_{\mu R} - g Y/2 B_\mu \right) \Phi_R^2$$

$$- V(\Phi_L, \Phi_R) \hspace{1cm} (3)$$

where $g_L = g_R = g$ are the $SU(2)$ couplings, $g'$ is the $U(1)$ coupling, $\gamma^\mu$ are the Dirac matrices, $\tau^i$ are the Pauli spin matrices, $V(\Phi_L, \Phi_R)$ is the Higgs potential, $Y$ is the hypercharge ($Y = B - L$). Also $\psi$ is a fermionic spinor valid for both quarks ($q$) and leptons ($\ell$). Here, the vacuum expectation values of two doublets ($\nu_L$ and $\nu_R$ with the relation $v_R >> v_L$) could contribute to the gauge bosons masses.

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

$$V = -\left( \mu_L^2 \Phi^+_L \Phi_L + \mu_R^2 \Phi^+_R \Phi_R \right)$$

$$+ \rho_1 \left( (\Phi^+_L \Phi_L)^2 + (\Phi^+_R \Phi_R)^2 \right)$$

$$+ \rho_2 \left( \Phi^+_L \Phi_L \right) \left( \Phi^+_R \Phi_R \right) \hspace{1cm} (4)$$

The minimum of the potential corresponds to $\langle \Phi_L \rangle = \nu_L/\sqrt{2}$, $\langle \Phi_R \rangle = \nu_R/\sqrt{2}$. Choosing $\mu_R \geq \mu_L$ guarantees $v_R \geq v_L$. In the unitary gauge, there are two physical Higgs bosons: $h_L \equiv \text{Re}\Phi^0_L$ and $h_R \equiv \text{Re}\Phi^0_R$; these two states, in principle could mix with each other with mixing angle $\theta_{h} \simeq \left( \frac{\rho_2}{\rho_1} \right) \left( \frac{\nu_L}{\nu_R} \right)$ for $v_R >> v_L$. Their masses are given by

$$M^2_{h_R} \simeq \rho_1 v_R^2, \quad M^2_{h_L} \simeq \rho_1 \left( 1 - \frac{\rho_2^2}{\rho_1^2} \right) v_R^2$$

Gauge boson mass

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

$$L_{\text{boson}} = \left| -\frac{g}{2} \tau^i W_{\mu L} - g_2 Y/2 B_\mu \right| \Phi_L^2$$

$$+ \left| -\frac{g}{2} \tau^i W_{\mu R} - g_2 Y/2 B_\mu \right| \Phi_R^2 \hspace{1cm} (5)$$
After substituting the vacuum expectation values of the Higgs fields:

\[
\langle \Phi_L \rangle = \left( \begin{array}{c} 0 \\ v_L \end{array} \right), \quad \langle \Phi_R \rangle = \left( \begin{array}{c} 0 \\ v_R \end{array} \right),
\]

(6)

into the relation \(\Phi\), we obtain

\[
L_{\text{boson}} = \frac{g^2 v_L^2}{4} \left\{ (W^1_{\mu L})^2 + (W^2_{\mu L})^2 \right\} + \frac{2 v^2}{4} \left( g W^3_{\mu L} - g' B_{\mu} \right)^2 + \frac{g^2 v_R^2}{4} \left\{ (W^1_{\mu R})^2 + (W^2_{\mu R})^2 \right\} + \frac{v^2}{4} \left( g W^3_{\mu R} - g' B_{\mu} \right)^2
\]

(7)

Let us define:

\[
W^\pm_\alpha = \frac{1}{\sqrt{2}} (W^{\mu\alpha}_1 \pm i W^{\mu\alpha}_2), \quad Z_{\mu\alpha} = \frac{g W_{\mu\alpha}^3 - g' B_{\mu\alpha}}{\sqrt{g^2 + g'^2}}, \quad A_{\mu\alpha} = \frac{g' W_{\mu\alpha}^3 + g B_{\mu\alpha}}{\sqrt{g^2 + g'^2}}, \quad Z''_{\mu\alpha} = W_{\mu\alpha}^3,
\]

(8)

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

\[
L_{\text{boson}} = M_{W_L}^2 W^+_L W^-_L + M_{W_R}^2 W^+_R W^-_R + M_{Z_L}^2 Z_{\mu L} Z''_{\mu L} + M_{Z_R}^2 Z_{\mu R} Z''_{\mu R} + M_A^2 A_{\mu} A^\mu
\]

(10)

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

\[
M_{W_L} = \frac{g v_L}{2}, \quad M_{W_R} = \frac{g v_R}{2}, \quad M_A = 0,
\]

\[
M_{Z_L} = \frac{v\sqrt{g^2 + g'^2}}{2}, \quad M_{Z_R} = \frac{v R \sqrt{g^2 + g'^2}}{2}
\]

(11)

**Fermion mass**

We shall discuss here how fermion masses arise in this particular approach. The key idea of the model is to impose the existence of weak iso-singlets 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 the Ref. [9, 12, 32]). These heavy isosinglet vector like fermions include: color triplet with electric charge +2/3 as \(U_{L,R}\), color triplet with electric charge −1/3 as \(D_{L,R}\) and color singlet with electric charge −1 as \(E_{L,R}\) and with these extra fields, the Yukawa terms can be written as

\[
\mathcal{L} \supset -y_D (\bar{q}_L \Phi_L D_R + \bar{q}_R \Phi_R D_L) - M_D \bar{D}_L D_R - y_U (\bar{q}_L \Phi_L U_R + \bar{q}_R \Phi_R U_L) - M_U \bar{U}_L U_R - y_E (\bar{l}_L \Phi_L E_R + \bar{\nu}_R \Phi_R E_L) - M_E \bar{E}_L E_R + \text{H.c.}
\]

\[
\Rightarrow -y_d \bar{q}_L \Phi_L d_R - y_u \bar{q}_R \Phi_R u_R - y_e \bar{l}_L \Phi_L e_R + \text{H.c.}, \quad (12)
\]

where the SM Yukawa couplings are given by

\[
y_d = -y_d D_\mu \bar{d}_R, \quad (13a)
\]

\[
y_u = -y_u D_\mu \bar{u}_R, \quad (13b)
\]

\[
y_e = -y_e D_\mu \bar{e}_R. \quad (13c)
\]

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 \(\tilde{S}_{L,R}\) with the Yukawa couplings and the masses as below,

\[
\mathcal{L} \supset -y_S (\bar{\nu}_L \tilde{\Phi}_L S_R + \bar{\nu}_R \tilde{\Phi}_R S_L) - M_S \bar{S}_L S_R - \frac{1}{2} M^M_S \tilde{S}_L S_R + \tilde{S}_R S_R + \text{H.c.}. \quad (14)
\]

At this stage, we do not want the Yukawa couplings \(\bar{\nu}_L \tilde{\Phi}_L S'_C, \bar{\nu}_R \tilde{\Phi}_R S'_C\) 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 \(U(1)_X\) symmetry under which \(D_{L,R}, U_{L,R}, E_{L,R}, \tilde{S'}_{L,R}\) carry a quantum number \(X = 1\). Clearly, this \(U(1)_X\) is free of 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 \(\eta\) with Yukawa couplings to the neutral singlets \(S_{L,R}\),

\[
\mathcal{L} \supset -\frac{1}{2} f_S (\eta \tilde{S}_L S_L + \eta^* \tilde{S}_R S_R) + \text{H.c.}. \quad (15)
\]

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

\[
M^M_S = f_S(\eta). \quad (16)
\]

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

\[
\mathcal{L} \supset -\frac{1}{2} \overline{\nu}_L M_{L,L} \nu_L - \frac{1}{2} \overline{N}_R M_{R,R} N_R - \overline{\nu}_L M_{D,D} N_R + \text{H.c.}(17)
\]
with
\[ M_L = -g_S \frac{1}{M_S^2} y_S^T v_L, \]  
\[ M_R = -g_S \frac{1}{M_S^2} y_S^T v_R^2, \]  
\[ M_D = y_S \frac{1}{M_S} (M_S^D)^T \frac{1}{M_S^2} y_S^T v_L v_R. \]  

(18a, 18b, 18c)

Here we have assumed
\[ M_S^M \gg M_S^D, y_S v_R, y_S v_L, \]
by choosing the base where the Majorana mass matrix \( M_N^M \) is real and diagonal,
\[ M_S^M = \text{diag}(M_1, M_2, M_3) \simeq M. \]

(19)

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,
\[ \delta M_L = -M_D \frac{1}{M_R^2} M_D^T = O \left( \left( \frac{M_D^2}{M_R^2} \right)^2 \right), \]
\[ \delta M_L \ll M_L. \]

(20)

Therefore, we can well define the left- and right-handed Majorana neutrinos,
\[ \nu = \nu_L + \nu_R^c, \]  
\[ N = N_R + N_R^c, \]

(22a, 22b)

Diagonalization of the light neutrino mass matrix \( m_{\nu}^{} \) through lepton flavour mixing matrix \( U_{\text{PMNS}} \) gives us three light Majorana neutrinos \( m_{\nu}^{\text{light}} = U_{\text{PMNS}} M_L U_{\text{PMNS}}^T = \text{diag}(m_1, m_2, m_3) \).

If we look the structure of light neutrino mass matrix \( M_L \) and heavy neutrino mass matrix \( M_N \), then it is clear that both the matrix can be simultaneously diagonalized by the same unitary matrix \( U_{\text{PMNS}} \), i.e \( M_{\text{heavy}}^{\text{diag}} = U_{\text{PMNS}} M_N U_{\text{PMNS}}^T \).

Hence, one can correlate the eigenvalues of the light and heavy Majorana neutrino 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
\[ m_{\nu} = M_L = U_{\text{PMNS}}^{\text{diag}} \{ m_1, m_2, m_3 \} U_{\text{PMNS}}^*, \]  
\[ M_N = M_R = U_{\text{PMNS}}^{\text{diag}} \{ m_1, m_2, m_3 \} U_{\text{PMNS}}^* \frac{v_L^2}{v_R^2}. \]

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

**NEUTRINOLESS DOUBLE BETA DECAY**

In this section, we shall present the lepton number violating processes such as neutrinoless double beta decay in left-right symmetric model without having a scalar bidoublet. We shall examine how the \( 0^{\nu}B\beta \) 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 \( W_R \) boson. The Feynman diagrams that give rise to neutrinoless double beta decay are depicted in Fig. (1a), (1b), (1c), and (1d).

The corresponding Feynman amplitude for these above diagrams is depicted in following table as

| Feynman diagrams | Amplitude |
|------------------|-----------|
| Fig. 1(a)        | \( A_0 \propto G_F^2 \frac{U_{e1}^2 m_{\nu}}{p^2} \) |
| Fig. 1(b)        | \( A_0 \propto G_F^2 \left( \frac{M_R}{M_W} \right)^2 U_{e1}^2 \left( \frac{M_D}{M_R} \right) \frac{1}{|p|} \) |
| Fig. 1(c)        | \( A_c \propto G_F^2 \left( \frac{M_R}{M_W} \right)^2 U_{e1}^2 \frac{M_D}{p} \) |
| Fig. 1(d)        | \( A_d \propto G_F^2 \left( \frac{M_R}{M_W} \right)^2 U_{e1}^2 \frac{M_D}{p} \frac{1}{|p|} \) |

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

In this table, \( G_F \approx 1.2 \times 10^{-5}\text{GeV}^{-2} \) is the Fermi con-
stant, $M_{W_R}$ is the right-handed charged gauge boson mass, $\bar{\zeta}_{L-R}$ is the $W_L - W_R$ mixing and $p^2$ is the neutrino virtuality. In order to evaluate the relative contributions of different terms, it is worth to note here that we shall analyze the effect of neutrinoless double beta decay while the representative set of parameters in this model are: $M_{W_R} \sim 10$ TeV and the heaviest right-handed neutrino mass around $\sim (1 - 10)$ MeV. WeV. With this set of parameters, the relevant dominant contributions are found to be

$$A_e \propto \frac{G_F^2}{p^2} \left( U_{e1}^2 m_{\nu e} \right) \sim \frac{G_F^2}{p^2} \times 10^{-2} \text{eV}$$
$$A_e \propto \frac{G_F^2}{p^2} \times 10^{-8} \text{eV} \sim \frac{G_F^2}{p^2} \times 10^{-1} \text{eV}$$

The standard contribution from left-handed current

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

$$\Gamma_{0\nu} = G^{0\nu} \left| \frac{M^{0\nu}}{m_e} \right|^2 |M^{e\nu}|^2$$

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

$$|M^{e\nu}| = |U_{ej}^2 m_j|, \quad (24)$$

Here $U_{ej}$ are the elements of the lepton mixing matrix $U_{PMNS}$ given in [13] which contains three mixing angles and three phases (one Dirac and two Majorana phases). It is worth to emphasize here that the neutrinoless double beta decay experiment can probe the phases which crucially depends on the pattern of the neutrino masses i.e whether neutrinos are Normal, or, Inverted, or, quasi-degenerate and on the magnitude of the neutrino masses. One can parametrize the effective Majorana mass in terms of the elements of $U_{PMNS}$ and mass eigenvalues as

$$|M^{e\nu}| = |U_{e1}^2 m_1 + e^{2i\alpha_2} \sin^2 \theta_{12} m_2 + e^{2i\alpha_3} \sin^2 \theta_{13} m_3|,$$  

(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 3-$\sigma$ ranges and the best-fit values of the mass squared differences and mixing angles $\sin^2 \theta_{12}$, $\sin^2 \theta_{23}$ from global analysis of oscillation data [33] and value of $\sin^2 \theta_{13}$ from the recent measurement of DayaBay experiment [38]. In particular, the representative values of the parameters which has been taken in this model, in order to give the result shown in Fig. 2 are as follows

$$\Delta m^4 = \begin{align*}
\Delta m^{0\text{sol}} [10^{-5} \text{eV}^2] & = 7.58 \text{[best-fit]} \quad 6.99 - 8.18 \text{[3-sigma]} \\
|\Delta m^{0\text{atm}} [10^{-3} \text{eV}^2]| & = 2.35 \text{[best-fit]} \quad 2.06 - 2.67 \text{[3-sigma]} \\
\sin^2 \theta_{12} & = 0.306 \text{[best-fit]} \quad 0.259 - 0.359 \text{[3-sigma]} \\
\sin^2 \theta_{23} & = 0.42 \text{[best-fit]} \quad 0.34 - 0.64 \text{[3-sigma]} \\
\sin^2 \theta_{13} & = 0.023 \text{[best-fit]} \quad 0.009 - 0.037 \text{[3-sigma]} 
\end{align*}$$

In the plot, we need to explain how effective Majorana mass probes which kind of mass pattern of neutrinos. As shown in Fig. 2 the cyan band for NH corresponds to varying the parameters in their $3\sigma$ range whereas the red band corresponds to the best-fit parameters where the $\sin^2 \theta_{13}$ values are taken from recent Daya-Bay result. In both figures the Majorana phases are varied between 0 to 2$\pi$. In the same manner, The green band for IH corresponds to varying the parameters in their $3\sigma$ 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. [31]. We shall now move to next subsection where the dominant contribution comes from the right-handed current and present an analysis for the result obtained with MeV mass range of RH Majorana neutrinos.

New contribution from right-handed current

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

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

$$m_2 = \sqrt{m_1^2 + \Delta m_{sol}^2}, \quad m_3 = \sqrt{m_1^2 + \Delta m_{atm}^2 + \Delta m_{sol}^2}$$

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

- Inverted hierarchy (IH) implies $m_3 << m_1 \sim m_2$ 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

$$m_1 = \sqrt{m_3^2 + \Delta m_{atm}^2}, \quad m_2 = \sqrt{m_3^2 + \Delta m_{sol}^2 + \Delta m_{atm}^2}$$

- The quasi-degenerate limit correspond to $m_1 \approx m_2 \approx m_3 >> \sqrt{\Delta m_{atm}^2}$.

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 $M_N^{\nu}$ as a function of lightest light-handed neutrinos.

**Hierarchical pattern of the neutrino masses**

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

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

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

$$|M_N^{\nu}|_{NH} = \left( \frac{M_{W_R}}{M_{W_L}} \right)^4 \sum_j U_{e j}^2 M_j$$

$$= \left( \frac{M_{W_R}}{M_{W_L}} \right)^4 M_3 \cos^2 \theta_{12} \cos^2 \theta_{13} \frac{m_1}{m_3} + \sin^2 \theta_{12} \cos^2 \theta_{13} e^{2i\alpha_3} \frac{m_2}{m_3} + \sin^2 \theta_{13} e^{2i\alpha_3} \right) \quad (26)$$

In the purely hierarchical case, $10^{-5} \text{eV} < m_1 < 10^{-3} \text{eV}$, one can write $m_2 \sim \sqrt{\Delta m_{sol}^2}, m_3 \sim \sqrt{\Delta m_{atm}^2}$. Given the input parameters in our model, the ratio between left- and right-handed charged gauge boson masses is found to be $10^{-8}$, the ratio between solar and atmospheric mass square difference is $m_2/m_3 = \sqrt{\Delta m_{sol}^2/\Delta m_{atm}^2} = \{0.16, 0.2\}$ corresponds to minimum and maximum value respectively. Since

$$m_1$$ is very small, the first term in eqn. (26) gives negligible contribution and hence can be neglected. With the choice made for $M_3$ at 5 MeV scale, the effective Majorana mass is

$$|M_N^{\nu}|_{NH} = 0.05 \sin^2 \theta_{12} \cos^2 \theta_{13} \sqrt{\frac{\Delta m_{sol}^2}{\Delta m_{atm}^2}} e^{2i\alpha_2}$$

$$+ \sin^2 \theta_{13} e^{2i\alpha_3} \quad (27)$$

The maximum and minimum values of $|M_N^{\nu}|_{NH}$ corresponds to the phase values $\alpha_2$, $\alpha_3 = 0, \pi$ and $\alpha = 0, \pi; \alpha_3 = \pi/2$ respectively. As shown in Fig. 3, the blue band in the regime $10^{-5} \text{eV} < m_1 < 10^{-3} \text{eV}$ corresponds to minimum and maximum values as follows

$$\left\{ \begin{array}{ll}
|M_N^{\nu}|_{NH}(\text{max}) = 0.0075 \\
|M_N^{\nu}|_{NH}(\text{min}) = 0.0055 
\end{array} \right.$$

For intermediate hierarchical values of $m_3$, say $10^{-3} \text{eV} < m_3 < 10^{-2} \text{eV}$, still the first term in eqn. (26) can be neglected. For illustration, one can see that the first term of eqn. (27) is small because the smallness of $\sqrt{\Delta m_{sol}^2/\Delta m_{atm}^2}$, at the same time, the second term is also suppressed due to the factor $\sin^2 \theta_{13}$. As a result, there is cancellation occurs in this regime due to relative phase cancellation of $\alpha_2$ and $\alpha_3$.

**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

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

![FIG. 3: The new dominant contribution to neutrinoless double beta decay coming from the right-handed current having $M_1$ 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.](image)
Now the expression for $M_N^\nu$ becomes

$$|M_N^\nu|_{IH} = \left( \frac{M_W/\sin^2\theta_W}{M_W} \right)^4 M_1 \cos^2\theta_{12} \cos^2\theta_{13} \frac{m_3}{m_2} + \sin^2\theta_{12} \cos^2\theta_{13} e^{2i\alpha_2} + \sin^2\theta_{13} e^{2i\alpha_3} \frac{m_1}{m_2} \right).$$

Before illustrating the analytical behavior of this contribution, it should be noted here that the value of $m_3$ in the case inverted hierarchy is such that $m_3 \ll \sqrt{\Delta m^2_{\text{sol}}}$, $m_1 \simeq \sqrt{\Delta m^2_{\text{atm}}}$ and $m_2 \simeq \sqrt{\Delta m^2_{\text{sol}}}$. 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 the eqn. (28) can be safely neglected. Now the effective Majorana mass in this inverted hierarchical scheme is given below

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

Similarly, the same arguments discussed in above subsection will gives the maximum and minimum values of the $|M_N^\nu|_{IH}$ as

$$\begin{cases} |M_N^\nu|_{IH} \text{ (max)} = 0.1 \\ |M_N^\nu|_{IH} \text{ (min)} = 0.05 \end{cases}$$

**Quasi-degenerate pattern of the neutrino masses**

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

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

Since quasi-degeneracy pattern of light neutrino masses also implies quasi-degeneracy in heavy right-handed neutrino sector which implies

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

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

$$|M_N^\nu|_{OD} = \left( \frac{M_W/\sin^2\theta_W}{M_W} \right)^4 M_0 \cos^2\theta_{12} \cos^2\theta_{13} + \sin^2\theta_{12} \cos^2\theta_{13} e^{2i\alpha_2} + \sin^2\theta_{13} e^{2i\alpha_3}.$$  (30)

From this relation, we can conclude that effective neutrino mass from RH current is independent of lightest neutrino mass in the quasi-degenerate limit. In other words, the value of $|M_N^\nu|$ remains constant with increasing $m_1$.

**Total contribution**

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

$$\Gamma_{0\nu} = G_{0\nu}^2 \cdot \frac{|M_{0\nu}^\nu|^2}{m_e} |m_{\text{eff}}^\nu|^2.$$  (31)

The effective neutrino mass contribution to neutrinoless double beta decay is

$$|m_{\text{eff}}^\nu|^2 = \left( \left| U_{e_j}^2 M_j \right|^2 + \frac{M_W^2 M_{W_R}^2}{M_W^4} \right).$$

where the individual contribution are $M_{\nu}^\nu = U_{e_j}^2 M_j$ and $M_N^\nu = \frac{M_W^2}{M_W^4} U_{e_j}^2 M_N$. This combine contribution is illustrated in Fig. 4.

![FIG. 4: 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 lightest neutrino masses respectively.](image)

**COMMENTS ON COSMOLOGICAL CONSTRAINTS**

We shall discuss in this section whether the MeV scale RH neutrinos for $M_W$ lying in 1-10 TeV region is consistent with the big-bang nucleosynthesis (BBN) bound and from the over closing of the Universe. We are in a problematic situation when $M_W$ lies around TeV scale, which in turn gives over-abundance of RH neutrinos $N$, because the $SU(2)_R$ gauge interaction keep 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 turns gives $\tau_N \lesssim 0.1$ 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
\[
\Gamma_{N \to \ell \pi} = \frac{G_F^2 |V_{ud}^R|^2 |V_{\ell N}^R|^2 f_\pi^2 M_N^3}{8\pi M_{WR}^4} \left(1 - x_\pi^2\right)^2 - x_\pi^2 \left(1 + x_\pi^2\right) \left[1 - (x_\pi + x_\ell^2)\right] \left(1 - (x_\pi - x_\ell^2)\right) \frac{1}{2},
\]
where $x_\pi, x_\ell = m_\pi, m_\ell / M_N$, $V^R$ is the right-handed lepton mixing matrix, $V_{ud}^R$ is the analog quark one and $f_\pi = 130$ MeV is the pion decay constant. We recall that $V_{ud}^R \simeq V_{ud}^L \simeq 0.97$; 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 $\tau_N$ is safely shorter than a second. Hence the constraints coming from the cosmology gives $M_N > 140$ MeV. This range of RH neutrino mass will push up the $M_{WR}$ scale beyond TeV scale which spoils the possible probe of our scenario in near future like at LHC.

The prescribed scenario discussed above suffers from serious problem when $M_N$ lies below $140$ MeV, the life-time becomes longer than a second, a decaying $N$ would pump too much entropy into the universe. The point is that they decouple relativistically at the temperature
\[
T_D^N = T_D^\nu \left(\frac{M_{WR}}{M_W}\right) \frac{1}{2},
\]
where $T_D^\nu \simeq 1$ MeV is the neutrino decoupling temperature. Therefore, for a representative value of $M_{WR} \sim 5$ TeV,
\[
T_D^N \simeq 250$ MeV.
\]

Then, since between $T_D^N$ and 1 MeV, only muons and pions decouple, at BBN $N$’s are almost equally abundant as light neutrinos. The only way out would be to make $N$ stable and to avoid the over-closure of the universe, lighter than about eV [36, 37]. 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 [38, 39] that four light neutrinos give the best fit to cosmological data, while five is disfavored and six is basically excluded.

CONCLUSION

We have discussed neutrinoless double beta decay in the context of left-right symmetric models with the minimal Higgs content, which is different from standard version of $L-R$ model that make 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 $\Phi_L$ and $\Phi_R$ without invoking triplets and bi-doublet, and the fermion masses are generated by integrating out the extra vector-like heavy quarks and leptons. In gauge sector, there is no mixing between left- and right-handed weak gauge bosons at tree level, but can induced at one loop level. In this particular scenario where the light neutrino and heavy Majorana neutrino are related with each other and diagonalized by the same Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix, the neutrinoless double beta decay receives important contributions from the right handed current. In fact, the choice we made for the right handed Majorana neutrino mass around MeV range in the context of neutrinoless double beta decay, the last formula of Sec-II predicts the right-handed gauge boson mass $M_{WR}$ to be at least order of 5 TeV.

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[1] G. Senjanovic and R. N. Mohapatra, Phys. Rev. D12, 1502 (1975).
[2] R. N. Mohapatra and J. C. Pati, Phys. Rev. D11, 2558 (1975).
[3] R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980).
[4] R. N. Mohapatra and R. E. Marshak, Phys. Rev. Lett. 44, 1316 (1980).
[5] N. G. Deshpande, J. F. Gunion, B. Kayser, and F. I. Olness, Phys. Rev. D44, 837 (1991).
[6] R. N. Mohapatra and J. C. Pati, Phys.Rev. D11, 566 (1975).
[7] P. Minkowski, Phys. Lett. B67, 421 (1977).
[8] T. Yanagida (1979), in Proceedings of the Workshop on the Baryon Number of the Universe and Unified Theories, Tsukuba, Japan, 13-14 Feb 1979.
[9] Z. Berezhiani, Phys.Lett. B129, 99 (1983).
[10] Z. Berezhiani, Phys.Lett. B150, 177 (1985).
[11] Z. Berezhiani and R. Rattazzi, Phys.Lett. B279, 124 (1992).
[12] Z. G. Berezhiani and R. Rattazzi, Nucl.Phys. B407, 249 (1993), hep-ph/9212245.
[13] S. Fukuda et al. (Super-Kamiokande), Phys. Rev. Lett. 86, 5656 (2001), hep-ex/0103033.
[14] Q. R. Ahmad et al. (SNO), Phys. Rev. Lett. 89, 011301 (2002), nucl-ex/0204008.
[15] Q. R. Ahmad et al. (SNO), Phys. Rev. Lett. 89, 011302 (2002), nucl-ex/0204009.
[16] J. N. Bahcall and C. Pena-Garay, New J. Phys. 6, 63 (2004), hep-ph/0404061.
[17] S. Weinberg, Phys.Rev.Lett. 43, 1566 (1979).
[18] H. Klapdor-Kleingrothaus, A. Dietz, L. Baudis, G. Heusser, I. Krivosheina, et al., Eur.Phys.J. A12, 147 (2001), hep-ph/0103062.
[19] H. Klapdor-Kleingrothaus, I. Krivosheina, and I. Titkova, Mod.Phys.Lett. A21, 1257 (2006).
[20] C. Aalseth et al. (IGEX Collaboration), Phys.Rev. D65, 092007 (2002), hep-ex/0202026.
[21] H. Klapdor-Kleingrothaus, I. Krivosheina, A. Dietz, and O. Chkvorets, Phys.Lett. B586, 198 (2004), hep-ph/0404088.
[22] C. Arnaboldi et al. (CUORICINO Collaboration), Phys.Rev. C78, 035502 (2008), hep-ex/0802.3439.
[23] I. Avignone, Frank T., S. R. Elliott, and J. Engel, Rev.Mod.Phys. 80, 481 (2008), nucl-ex/0708.1033.
[24] J. Gomez-Cadenas, J. Martin-Albo, M. Mezzetto, F. Monrabal, and M. Sorel, Riv.Nuovo Cim. 35, 29 (2012), 1109.5515.
[25] S. Pascoli and S. T. Petcov, Phys.Rev. D77, 113003 (2008), hep-ph/0711.4993.
[26] H. Pas, M. Hirsch, and H. Klapdor-Kleingrothaus, Phys.Lett. B459, 450 (1999), hep-ph/9810382.
[27] R. Mohapatra, Phys.Rev. D34, 3457 (1986).
[28] M. Doi and T. Kotani, Prog.Theor.Phys. 89, 139 (1993).
[29] V. Tello, M. Nemevsek, F. Nesti, G. Senjanovic, and F. Vissani, Phys.Rev.Lett. 106, 151801 (2011), hep-ph/1011.3522.
[30] M. Nemevsek, F. Nesti, G. Senjanovic, and V. Tello (2011), hep-ph/1112.3061.
[31] J. Chakraborty, H. Z. Devi, S. Goswami, and S. Patra, JHEP 1208, 008 (2012), hep-ph/1204.2527.
[32] P.-H. Gu, Phys.Rev. D81, 095002 (2010), 1001.1341.
[33] Z. Maki, M. Nakagawa, and S. Sakata, Prog.Theor.Phys. 28, 870 (1962).
[34] G. Fogli, E. Lisi, A. Marrone, A. Palazzo, and A. Rotunno, Phys.Rev. D84, 053007 (2011), hep-ph/1106.6028.
[35] F. An et al. (DAYA-BAY Collaboration), Phys.Rev.Lett. 108, 171803 (2012), hep-ex/1203.1669.
[36] U. Seljak, A. Slosar, and P. McDonald, JCAP 0610, 014 (2006), astro-ph/0604335.
[37] G. Fogli, E. Lisi, A. Marrone, A. Melchiorri, A. Palazzo, et al., Phys.Rev. D78, 033010 (2008), hep-ph/0805.2517.
[38] J. Hamann, S. Hannestad, G. G. Raffelt, I. Tamborra, and Y. Y. Wong, Phys.Rev.Lett. 105, 181301 (2010), hep-ph/1006.5276.
[39] J. Hamann, S. Hannestad, G. G. Raffelt, and Y. Y. Wong, JCAP 1109, 034 (2011), astro-ph/1108.4136.