Dominant light-heavy neutrino mixing contribution
to $0\nu\beta\beta$ in minimal left-right symmetric model with
universal seesaw

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ABSTRACT: We present a detailed discussion on neutrinoless double beta decay ($0\nu\beta\beta$) within left-right symmetric models based on the gauge symmetry of type $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ as well as $SU(3)_L \times SU(3)_R \times U(1)_X$ where fermion masses including that of neutrinos are generated through a universal seesaw mechanism. We find that one or more of the right-handed neutrinos could be as light as a few keV if left-right symmetry breaking occurs in the range of a few TeV to 100 TeV. With such light right-handed neutrinos, we perform a detailed study of new physics contributions to $0\nu\beta\beta$ and constrain the model parameters from the latest experimental bound on such a rare decay process. We find that the dominant new physics contribution to $0\nu\beta\beta$ in such a scenario comes from large light-heavy neutrino mixing and can saturate the existing experimental bounds even if the right handed sector gauge bosons are beyond the reach of present collider experiments. We also find the constraints on these models from the latest experimental bounds on charged lepton flavour violating decays like $\mu \rightarrow e\gamma$.

KEYWORDS: Neutrinoless Double Beta Decay, Neutrino masses and mixing, Universal Seesaw, Left-Right Theories
1 Introduction

The Standard Model (SM) of particle physics has been established as the most successful description of the fundamental particles and their interactions: strong, weak and electromagnetic. The model based on the local gauge symmetry $SU(2)_L \times U(1)_Y \times SU(3)_C$ describing strong, weak and electromagnetic interactions between fundamental particles gets broken down to $U(1)_Q \times SU(3)_C$ remnant gauge symmetry spontaneously due to the non-zero vacuum expectation value (vev) of the Higgs field charged under the $SU(2)_L \times U(1)_Y$ symmetry of the model. Since the discovery of the Higgs boson in 2012 at the large hadron collider (LHC), the SM has been confirmed again and again as the only theory around the electroweak scale with no signs of new physics yet. In spite of these null results for new physics beyond the standard model (BSM), there are convincing amount of evidence suggesting the presence of new physics. This need for new physics arises due to the inadequacies of the SM as it can not address several observed phenomena as well as theoretical questions. Non zero but tiny neutrino mass [1] is one such observation which the SM can not address. Due to the absence of right handed neutrinos in the SM, there is no renormalisable interaction between the neutrino and the Higgs field resulting in vanishing mass of neutrinos. The SM also can not explain the origin of parity violation seen in low energy weak interaction processes which is at sharp contrast with other interactions like electromagnetic and strong which are parity conserving. This motivates one to speculate that all fundamental interactions are parity conserving at the most fundamental level or at a very high energy scale and
the SM is a parity violating low energy manifestation of such a unified parity conserving theory. These two observations namely, non-zero neutrino mass and parity violation in weak interactions can be explained naturally within the framework of the left-right symmetric model (LRSM) [2–4]-based on the gauge symmetry $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_C$ where both left and right-handed fermions are treated on equal footing in a parity symmetric manner. The right handed fermions transform as doublets under the new $SU(2)_R$ gauge symmetry similar to the transformation of left handed fermions under the $SU(2)_L$ gauge symmetry of the SM. The inclusion of right handed neutrinos become a necessity in such a framework and hence neutrinos can naturally acquire a non-zero mass.

In the framework of the LRSM, the light neutrino masses can arise in several different ways depending on the scalar content and the way spontaneous symmetry breaking occurs from LRSM gauge symmetry to that of the SM and finally to the $U(1)_Q \times SU(3)_c$ symmetry. The minimal version of LRSM contains scalars which transform as $SU(2)$ triplets and bidoublet with the triplets playing to role of breaking the LRSM gauge symmetry to that of the SM and the bidoublet playing the same role in electroweak symmetry breaking. This promising theory of parity preserving weak interaction can have testable consequences for different experiments or observed phenomena when LRSM gauge symmetry breaking occurs at few TeV. Such tantalising consequences can be in the gauge sector in terms of additional gauge bosons \[5–35\], in the Higgs sector \[36–49\], in the context of neutrinoless double beta decay \[50–71\], in dark matter contexts \[72–77\] in low-energy charged lepton flavour violation (LFV) \[11, 55, 60, 66, 68, 70, 78–88\] and electric dipole moment (EDM) \[57, 63, 89–92\]. Apart from this most widely studied minimal LRSM, there have been alternative formulation of the left-right symmetric models as well which have different scalar content and different ways of generating fermion masses. For example, if the scalar sector of the minimal LRSM (MLRSM) is replaced by a pair of scalar doublets transforming under $SU(2)_L, SU(2)_R$ respectively then the desired symmetry breaking can be achieved in a more minimal way. However, the absence of the bidoublet prevents one from writing renormalisable mass terms for the fermions forcing one to introduce higher dimensional operators \[93\]. A renormalisable version of such LRSM with universal seesaw for all fermions can be achieved by introducing additional heavy fermions \[94–96\]. Several other realisations of fermion masses within such LRSM without scalar bidoublet can be found in \[97–99\]. The LRSM with universal seesaw (LRSM-US) for all fermions was also studied from $0\nu\beta\beta$ point of view in \[58, 100, 101\] and more recently within the 331 set up \[102\]. The extension of such a framework to the 331 models namely, the $SU(3)_L \times SU(3)_R \times U(1)_X$ models has also been studied recently \[102\]. In the 331 version of LRSM with universal seesaw, it was pointed out that due to the existence of light right handed neutrinos in the keV-MeV range, such a model can give observable contributions to $0\nu\beta\beta$ due to large light-heavy neutrino mixing. Such a version of LRSM where the scalar sector can be made more minimal than MLRSM at the price of introducing extra vector-like iso-singlet quarks and leptons was of much interest in the context of LHC anomalies \[103–105\]. Apart from having all other generic features of minimal LRSM in terms of providing an explanation to the origin neutrino mass, origin of parity violation, strong CP problem, allowing the possibility of non-supersymmetric grand unification the LRSM-US can also explain the origin of fermion
mass hierarchies through seesaw mechanism instead of arbitrarily fine tuning the Yukawa couplings.

In the present work, we intend to study the contribution of different particles in LRSM-US to lepton number violating rare decay process of neutrinoless double beta decay that have been looked for at several experiments resulting in strict upper bound (lower bound) on the amplitude (half-life) for such a process. This is an extension of the recent work [102] to do a complete scan of all parameters, though confined to $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ version of LRSM for simplicity instead of the 331 version of [102]. With the present 0νββ experiments like KamLAND-Zen [106, 107], GERDA [108, 109] probing the quasi-degenerate regime of light neutrino masses, one can expect the next generation experiments to cover the entire parameter space for 0νββ, at least in the case inverted hierarchical pattern of light neutrino masses. The current lower limit on the half-life of this rare process from these two experiments lie in the range of $10^{25} - 10^{26}$ year. The projected sensitivity of the phase III of KamLAND-Zen is $T_{1/2} > 2 \times 10^{26}$ year after two years of data taking. Similar goal is also set by the GERDA experiment to reach $T_{1/2} > 10^{26}$ year. We show that the contributions to 0νββ in LRSM-US can saturate these experimental bounds even if the scale of $SU(2)_R \times U(1)_{B-L}$ breaking is outside the reach of present collider experiments. In particular, such a model allows the right handed neutrinos to be as light as a keV without any fine-tuning even if the symmetry breaking scale is kept around a few (tens of) TeV. This also allows for the possibility of large light-heavy neutrino mixing and can contribute significantly to 0νββ through light-heavy or left-right mixing diagrams. It is interesting to note that in MLRSM such diagrams remain suppressed or at least equally dominant as other diagrams for most of the region of parameter space whereas in LRSM-US these diagrams can be the most dominant one. Such large light-heavy mixing between neutrinos can introduce non-unitary effects to leptonic mixing and can be constrained significantly from charged lepton flavour violating decay like $\mu \to e\gamma$ that have been looked for at ongoing experiments like MEG [110]. We constrain the model parameters from the requirement of satisfying both 0νββ and LFV constraints from the latest experimental data.

We first discuss neutrinoless double beta decay in the framework of left-right symmetric model with universal seesaw in two regimes: i) firstly considering the $SU(2)_R$ gauge boson masses $M_{W_R} \approx 3.5$ TeV and equivalently the right handed neutrino masses $M_R \approx O(\text{keV})$, ii) secondly, with $M_{W_R} \approx 100$ TeV (or equivalently, $M_R \approx O(\text{MeV})$). In the former case, the diagrams mediated by $W_L - W_R$ as well as $W_L - W_L$ where light-heavy neutrino mixing can play the major role. On the other hand, the purely $W_R$ mediated diagram remains suppressed in this case. However, with $M_{W_R} \approx 100$ TeV, the $W_R - W_R$ mediated as well as $W_L - W_R$ mediated diagrams give negligible contribution to 0νββ. The only dominant contribution arises from $W_L - W_L$ mediated diagrams with the exchange of heavy neutrinos having masses in the MeV range where large light-heavy neutrino mixing plays an important role, once again. The importance of light-heavy neutrino mixing in the study of 0νββ within generic LRSM with type I/II seesaw was pointed out by [63]. As shown in other 0νββ studies of generic LRSM, there are other sizeable 0νββ contributions also that can saturate the present experimental bounds. However, in the present model with universal seesaw, we find that the light-heavy neutrino mixing contribution is the most dominant one, a feature
which distinguishes it from the generic LRSM. After showing the new physics contribution to $0 \nu \beta \beta$ for some benchmark values of parameters, we also scan the entire parameter space and put the constraints on $W_R$ mass, light-heavy neutrino mixing parameter as well as the lightest neutrino mass from the requirement of satisfying the current experimental bounds on $0 \nu \beta \beta$. We also check that for the entire region of our interest with $W_R$ mass being varied all the way up to 100 TeV, the gauge boson mediated diagrams contributing to $\mu \rightarrow e\gamma$ remain very much suppressed compared to the latest experimental bound. We also identify another possible new physics contribution to this LFV decay process mediated by neutral scalars and charged vector like leptons. We constrain the corresponding Yukawa couplings and vector fermion mass from the requirement of satisfying the MEG 2016 bound [110] on this LFV decay process using a benchmark value of scalar mass. Apart from using the latest experimental constraints on light neutrino parameters that have appeared in global fit work [111], we also use the limit on the sum of light neutrino masses from the Planck mission data as $\sum_i m_i < 0.17$ eV [112].

The paper is organised as follows. In section 2, we briefly discuss the LRSM with universal seesaw for all fermions. In section 3 we discuss different contributions to $0 \nu \beta \beta$ in the model followed by a brief discussion on possible new physics contribution to the charged lepton flavour violating decay $\mu \rightarrow e\gamma$ in section 4. We discuss our numerical calculations in section 5 and finally conclude in section 6.

## 2 Left-Right Symmetric Model with Universal Seesaw

In this section, we briefly recapitulate the left-right symmetric model without scalar bidoublet where all fermion masses are generated by a common universal seesaw mechanism. The particle content of the model transforms non-trivially under the gauge symmetry of the model given by

$$G_{LR} \equiv SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_c,$$

(2.1)

which gets broken down to the $U(1)_Q \times SU(3)_c$ of electromagnetism and colour spontaneously at two stages such that the electromagnetic charge $Q$ is defined as

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

(2.2)

We denote $T_{3L}$ and $T_{3R}$ are, respectively, the third component of isospin corresponding to the gauge groups $SU(2)_L$ and $SU(2)_R$, and $Y$ is the hypercharge. Here difference between baryon and lepton number is promoted to local gauge symmetry. It should be noted that the model also has an in built discrete $Z_2$ symmetry or left-right symmetry (D parity) under which forces the couplings in the left and right sectors equal, making the theory left-right symmetric. The fermion content of the model is

$$Q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}, \quad Q_R = \begin{pmatrix} u_R \\ d_R \end{pmatrix},$$

$$\ell_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}, \quad \ell_R = \begin{pmatrix} \nu_R \\ e_R \end{pmatrix},$$  

(2.3)
plus additional vector-like quarks and charged leptons,

\[ U_{L,R}, \ D_{L,R}, \ E_{L,R}, \ N_{L,R}. \]  

The spontaneous symmetry breaking is implemented with a scalar sector consisting of \( SU(2)_{L,R} \) doublets \( H_L \oplus H_R \) and the conventional scalar bidoublet of MLRSM is absent. All these fields with their transformations under the gauge symmetry are shown in table 1. The scalar potential of the model can be written as

\[
V = \mu_H^2 (H_L^\dagger H_L + H_R^\dagger H_R) + \lambda_H \left((H_L^\dagger H_L)^2 + (H_R^\dagger H_R)^2\right) + \lambda'_H (H_L^\dagger H_L)(H_R^\dagger H_R) \\
+ \lambda''_H (H_L^\dagger H_L)(H_R^\dagger H_R) + \text{h.c.} \tag{2.5}
\]

The scalar fields can acquire non-zero vev as

\[
\langle H_R \rangle = \left( \frac{v_R}{\sqrt{2}} \right), \quad \langle H_L \rangle = \left( \frac{v_L}{\sqrt{2}} \right). \tag{2.6}
\]

The vev of the neutral component of \( H_R \) spontaneously breaks the symmetry of the LRSM to that of the SM while the vev of the neutral component of \( H_L \) gives rise to the usual electroweak symmetry breaking. In other words, the desired symmetry breaking chain is

\[
SU(2)_L \times SU(2)_R \times U(1)_{B-L} \xrightarrow{\langle H_R^0 \rangle} SU(2)_L \times U(1)_Y \xrightarrow{\langle H_L^0 \rangle} U(1)_Q
\]

In the scalar potential written above, the discrete left-right symmetry is assumed which ensures the equality of left and right sector couplings. However, as shown in earlier works \[93–96\] the scalar potential of such a model with exact discrete left-right symmetry is too restrictive and gives to either parity preserving \( v_L = v_R \) solution or a solution with \( (v_R \neq 0, v_L = 0) \) at tree level. While the first one is not phenomenologically acceptable the latter solution can be acceptable if a non-zero vev \( v_L \neq 0 \) can be generated through radiative

\[
\begin{array}{|c|c|c|c|c|}
\hline
\text{Field} & SU(2)_L & SU(2)_R & U(1)_{B-L} & SU(3)_c \\
\hline
Q_L & 2 & 1 & 1/3 & 3 \\
Q_R & 1 & 2 & 1/3 & 3 \\
\ell_L & 2 & 1 & -1 & 1 \\
\ell_R & 1 & 2 & -1 & 1 \\
U_{L,R} & 1 & 1 & 4/3 & 3 \\
D_{L,R} & 1 & 1 & -2/3 & 3 \\
E_{L,R} & 1 & 1 & -2 & 1 \\
N_{L,R} & 1 & 1 & 0 & 1 \\
H_L & 2 & 1 & -1 & 1 \\
H_R & 1 & 2 & -1 & 1 \\
\hline
\end{array}
\]

Table 1. Field content and their transformations under the gauge symmetry of LRSM with universal seesaw.

In the scalar potential written above, the discrete left-right symmetry is assumed which ensures the equality of left and right sector couplings. However, as shown in earlier works \[93–96\] the scalar potential of such a model with exact discrete left-right symmetry is too restrictive and gives to either parity preserving \( v_L = v_R \) solution or a solution with \( (v_R \neq 0, v_L = 0) \) at tree level. While the first one is not phenomenologically acceptable the latter solution can be acceptable if a non-zero vev \( v_L \neq 0 \) can be generated through radiative
corrections [113]. While it may naturally explain the smallness of \( v_L \) compared to \( v_R \), it will constrain the parameter space significantly [113]. Another way of achieving a parity breaking vacuum is to consider softly broken discrete left-right symmetry by considering different mass terms for the left and right sector scalars [3, 93–96]. As it was pointed out by the authors of [3], such a model which respects the discrete left-right symmetry everywhere except in the scalar mass terms, preserve the naturalness of the left-right symmetry in spite of radiative corrections. Another interesting way is to achieve parity breaking vacuum is to decouple the scale of parity breaking and gauge symmetry breaking by introducing a parity odd singlet scalar [114]. In this work, we simply assume that the desired symmetry breaking can be achieved by considering different mass terms for left and right sector scalars [3] without incorporating any new field content. Such a minimal assumption is not going to affect our discussion of 0νββ and LFV or even the origin of fermion masses.

After the neutral components of the scalar fields acquire non-zero vev’s, the resulting gauge boson masses can be derived as

\[
M_{W_L} = \frac{g}{2}v_L, \quad M_{W_R} = \frac{g}{2}v_R, \quad M_{Z_L} = \frac{g}{2}v_L \sqrt{1 + \frac{g_1^2}{g^2 + g_1^2}}, \quad M_{Z_R} = \frac{v_R}{2} \sqrt{(g^2 + g_1^2)}
\]

Here \( g_L = g_R = g \) is the SU(2)_L,R gauge coupling whereas \( g_1 \) is the corresponding gauge coupling for \( U(1)_{B-L} \) symmetry. Unlike in the minimal LRSM, here there is no tree level mixing between the left and right gauge bosons \( W_L, W_R \). However, they can mix at one-loop level with fermions going in the loop. The mixing angle \( \xi \) can be estimated as

\[
\xi \approx \frac{\alpha}{4\pi \sin^2 \theta_W} \frac{m_b m_t}{M_{W_R}^2}
\]

Using \( \alpha = 1/137, \sin^2 \theta_W \approx 0.23, m_b \approx 4.2 \text{ GeV}, m_t \approx 174 \text{ GeV}, M_{W_R} \approx 3 \text{ TeV} \), we find \( \xi \approx 2 \times 10^{-7} \).

In the absence of scalar bidoublet one can not write down a Dirac mass term for fermions including quarks and lepton. Thus, we introduce vector-like fermions so that both left-handed and right handed fermion doublets of minimal left-right symmetric model can couple to each other with the following interaction Lagrangian,

\[
L \supset Y_U(Q_L H_L U_L + Q_R H_R U_R) + Y_D(Q_L H_L^* D_L + Q_R H_R^* D_R) + M_U U_L U_R + M_D D_L D_R \\
+ Y_E(\bar{e}_L H_L^* E_L + \bar{e}_R H_R^* E_R) + Y_L(\bar{L}_L H_L N_L + \bar{L}_R H_R N_R) + M_E E_L E_R + M_N^D N_L N_R \\
+ \frac{1}{2} M_N^M (N_L N_L + N_R N_R) + h.c. \quad (2.8)
\]

After spontaneous symmetry breaking, the charge fermion mass matrices are given by

\[
M_{uU} = \begin{pmatrix} 0 & Y_U v_L \\ Y_U^T v_R & M_U \end{pmatrix}, \quad M_{dD} = \begin{pmatrix} 0 & Y_D v_L \\ Y_D^T v_R & M_D \end{pmatrix}, \\
M_{eE} = \begin{pmatrix} 0 & Y_E v_L \\ Y_E^T v_R & M_E \end{pmatrix}.
\]

(2.9)
The usual quarks get their Dirac masses via universal seesaw as follows,

\[ M_u \approx Y_U^T \frac{1}{M_U} Y_U v_L v_R, \quad M_d \approx Y_D^T \frac{1}{M_D} Y_D v_L v_R \]  

(2.10)

and the mixing angles \( \theta_{UL} \) are found to be

\[ \tan(2\theta_U) \approx 2Y_U \frac{v_L R M_U}{M_U^2 \pm (Y_U v_R)^2}. \]  

(2.11)

The charged leptons get their mass as

\[ M_l \approx Y_E^T \frac{1}{M_E} Y_E v_L v_R \]  

(2.12)

The neutrino mass matrix in the basis \((\nu_L, \nu_R \equiv N)\) can have three independent terms

\[ M_L = -Y_\nu^T \frac{1}{M_N^L} Y_\nu v_L^2 \]
\[ M_R = -Y_\nu^T \frac{1}{M_N^R} Y_\nu v_R^2 \]
\[ M_D = Y_\nu^T \frac{1}{M_N^D} M_N^D Y_\nu v_L v_R \]

after integrating out the heavy neutral fermions \(N_{L,R}\). Considering \(M_N^D = M_N^M\), the neutral lepton mass matrix in the basis \((\nu_L, \nu_R \equiv N)\) can be written as

\[ M_\nu = \begin{pmatrix} Y_\nu^T \frac{1}{M_N^L} Y_\nu v_L^2 & Y_\nu^T \frac{1}{M_N^M} Y_\nu v_L v_R \\ Y_\nu^T \frac{1}{M_N^R} Y_\nu v_R^2 & Y_\nu^T \frac{1}{M_N^M} Y_\nu v_R^2 \end{pmatrix} = \begin{pmatrix} M_L & M_D \\ M_D & M_R \end{pmatrix} \]  

(2.13)

In the limit \(M_L \ll M_D \ll M_R\), the type-I seesaw contribution to light neutrino mass is given by

\[ M_\nu = M_L - M_D \frac{1}{M_R} M_D \]  

(2.14)

Using the above definitions of \(M_D, M_R\), we find that

\[ M_D \frac{1}{M_R} M_D = Y_\nu^T \frac{1}{M_N^M} Y_\nu \left( Y_\nu^T \frac{1}{M_N^L} Y_\nu \right)^{-1} Y_\nu^T \frac{1}{M_N^M} Y_\nu v_L^2 = Y_\nu^T \frac{1}{M_N^M} Y_\nu v_L^2 = M_L \]

resulting in a vanishing light neutrino mass matrix. One can avoid such a scenario by discarding the assumption \(M_N^D = M_N^M\) considered earlier. To have a realistic scenario of non-vanishing light neutrino mass, we will consider \(M_N^D \neq M_N^M\) case in our calculations.

The vector-like fermions, crucial for the implementation of the universal seesaw mechanism are tightly constrained from direct searches. For example, the vector-like quark masses have a lower limit \(m_q \geq 750 - 920\) GeV depending on the particular channel of decay \([115, 116]\) whereas this bound gets relaxed to \(m_q \geq 400\) GeV \([117]\) for long lived vector-like quarks. These limits however, get more uplifted by the latest analysis of the
13 TeV centre of mass energy data from the LHC. For example, the recent analysis [118]
constrains the vector-like top quark mass to be $m_{T} > 870 - 1170$ GeV depending on the
weak isospin properties of it. Further constraints on vector-like quarks can be found in
[119]. The constraints on vector-like leptons are much weaker $m_{l} \geq 114 - 176$ GeV [120].
The experimental constraints put these lower bounds not only on the vector-like fermions,
but also on the new gauge bosons of the model. The right handed gauge boson masses are
primarily constrained from $K - \bar{K}$ mixing and direct searches at the LHC. While
$K - \bar{K}$ mixing puts a constraint $M_{W_{R}} > 2.5$ TeV [121], direct search bounds depend on the particular channel under study. For example, the dijet resonance search in ATLAS experiment puts a bound $M_{W_{R}} > 2.45$ TeV at 95% CL [122] in the $g_{L} = g_{R}$ limit. On the other hand, the CMS search for same sign dilepton plus dijet $pp \rightarrow l^{\pm}l^{\pm}jj$ mediated by heavy right handed neutrinos at 8 TeV centre of mass energy excludes some parameter space in the $M_{i}^{\text{lightest}} - M_{W_{R}}$ plane [33] where $M_{i}^{\text{lightest}}$ is the mass of the lightest neutral fermion from right handed lepton doublets. More recently, the results on dijet searches at ATLAS experiment at 13 TeV centre of mass energy and 37 fb$^{-1}$ of $pp$ collision data have put even stronger limits on such heavy charged gauge bosons [123].

3  $0\nu\beta\beta$ in LRSM with Universal Seesaw

In minimal left-right symmetric model with universal seesaw (MLRSM-US) with additional
vector-like leptons, the Dirac as well as Majorana masses for neutral leptons arise from
Higgs fields $H_{L}$ and $H_{R}$. The resulting seesaw contributions to neutrino masses and their
Majorana nature leads to rare process like neutrinoless double beta decay. The charge
current interaction Lagrangian for leptons and quarks is given by

$$
\mathcal{L}_{CC}^{\text{lep}} = \frac{g_{L}}{\sqrt{2}} \left[ \sum_{\alpha = e, \mu, \tau} \bar{\ell}_{\alpha} \gamma^{\mu} P_{L} \nu_{\alpha} W_{L_{\mu}} + \text{h.c.} \right] + \frac{g_{R}}{\sqrt{2}} \left[ \sum_{\alpha = e, \mu, \tau} \bar{\ell}_{\alpha} \gamma^{\mu} P_{R} N_{\alpha} W_{L_{\mu}} + \text{h.c.} \right],
$$

$$
\mathcal{L}_{CC}^{\text{q}} = \left[ \frac{g_{L}}{\sqrt{2}} \bar{d} \gamma^{\mu} P_{L} u W_{L_{\mu}} + \frac{g_{R}}{\sqrt{2}} \bar{d} \gamma^{\mu} P_{R} u W_{R_{\mu}} + \text{h.c.} \right].
$$

In the present type-I plus type-II seesaw mechanism, the flavor neutrino eigenstates $\nu_{\alpha} \equiv \nu_{L_{\alpha}}$ and $N_{\beta} \equiv \nu_{R_{\beta}}$ are related to mass eigenstates $\nu_{i}$ and $N_{i}$ as,

$$
\nu_{\alpha} = U_{\alpha i} \nu_{i} + S_{\alpha i} N_{i},
$$

$$
N_{\beta} = T_{\beta i} \nu_{i} + V_{\beta i} N_{i}.
$$

The mixing matrices $U, V, S, T$ are given by

$$
\begin{pmatrix}
U & S \\
T & V
\end{pmatrix} = \begin{pmatrix}
1 - \frac{1}{2} RR^{\dagger} & R \\
-R^{\dagger} & 1 - \frac{1}{2} R^{\dagger} R
\end{pmatrix} \begin{pmatrix}
U_{L} & 0 \\
0 & U_{R}
\end{pmatrix}
$$

(3.1)
such that $U_L, U_R$ are the diagonalising matrices of light and heavy neutrino mass matrices $M_{\nu L}, M_{\nu R}$ respectively. Here $R = M_DM_R^{-1}$. Simplifying the above equation gives rise to

$$U = U_L - \frac{1}{2} R R^* U_L, \quad S = R U_R$$

$$T = -R^* U_L, \quad V = U_R - \frac{1}{2} R^* R U_R.$$  

Within MLRSM-US with neutral leptons $\nu_\alpha$ and $N_\beta$ and with negligible $W_L - W_R$ mixing, the relevant contributions to neutrinoless double beta decay are as follows:

- due to exchange of light left-handed and keV scale right-handed neutrinos via purely left-handed currents ($W_L - W_L$ mediation),
- due to exchange of light left-handed and keV scale right-handed neutrinos via purely right-handed currents ($W_R - W_R$ mediation),
- due to mixed helicity so called $\lambda$ diagrams which involves left-right neutrino mixing through mediation of $\nu_i, N_i$ neutrinos,
- due to mixed helicity $\eta$ diagrams through mediation of $\nu_i, N_i$ neutrinos involving $W_L - W_R$ gauge boson mixing as well as left-right neutrino mixing.

Before estimating Feynman amplitude and corresponding LNV effective Majorana mass parameters, we should have knowledge about the chiral structure of the matrix element with the neutrino propagator as follows,

$$P_L \frac{p + m_i}{p^2 - m_i^2} P_L = \frac{m_i}{p^2 - m_i^2}, \quad P_R \frac{p + m_i}{p^2 - m_i^2} P_R = \frac{m_i}{p^2 - m_i^2},$$

$$P_L \frac{\phi}{p^2 - m_i^2} P_R = \frac{\phi}{p^2 - m_i^2}, \quad P_R \frac{\phi + m_i}{p^2 - m_i^2} P_L = \frac{\phi}{p^2 - m_i^2},$$  

(3.2)

\[ m_i \frac{p^2 - m_i^2}{p^2} \simeq \begin{cases} \frac{m_i}{p^2}, & m_i^2 \ll p^2 \\ -\frac{1}{m_i} m_i^2 \gg p^2 \end{cases} \]  

(3.3)

and

\[ \frac{\phi}{p^2 - m_i^2} \propto \begin{cases} \frac{1}{|p|}, & m_i^2 \ll p^2 \\ -\frac{|p|}{m_i^2} m_i^2 \gg p^2 \end{cases} \]  

(3.4)
Figure 1. Feynman diagram for $0\nu\beta\beta$ transition due to exchange of $\nu_i$ and $N_j$ via purely left-handed currents.

3.1 Feynman amplitudes for different contributions to $0\nu\beta\beta$ transition

The amplitudes for $0\nu\beta\beta$ transition as shown in Fig. 1 due to exchange of light left-handed neutrinos and keV scale right-handed neutrinos are given by

$$A_{\nu}^{LL} \propto G_F^2 \sum_{i=1,2,3} \frac{U_{ei}^2 m_i}{p^2},$$

$$A_{\nu}^{N} \propto G_F^2 \sum_{j=1,2,3} \left( \frac{S^{2}_{ej} M_j}{p^2} \right),$$

with $p$ being the average momentum exchange for the process. In the above expression, $m_i$ are the masses of light neutrinos for $i = 1, 2, 3$ and $M_i$ are keV scale masses for right-handed neutrinos.

Figure 2. Feynman diagram for $0\nu\beta\beta$ transition due to exchange of $\nu_i$ and $N_j$ via purely right-handed currents.
The contribution from the left-handed neutrinos, right-handed neutrinos and $W^-_R$ exchange (Feynman diagram in Fig. 2) can be written as

$$A^{\nu}_{RR} \propto G_F^2 \sum_{i=1,2,3} \left( \frac{M_{W_L}}{M_{W_R}} \right)^4 \left( \frac{g_R}{g_L} \right)^4 \frac{T_{ei}^* m_i}{p^2},$$

$$A^{N}_{RR} \propto G_F^2 \sum_{j=1,2,3} \left( \frac{M_{W_L}}{M_{W_R}} \right)^4 \left( \frac{g_R}{g_L} \right)^4 \frac{V_{ej}^* M_j}{p^2},$$

where $M_i$ are the masses of right handed neutrinos for $i = 1, 2, 3$ and $M_i \ll |p|$ is assumed.

**Figure 3.** Feynman diagram for $0\nu\beta\beta$ transition due to exchange of $\nu_i$ and $N_j$ via left-right neutrino mixing.

The most relevant contribution from mixed helicity $\lambda$ diagram as shown in Fig. 5 is given by

$$A^{\nu}_{\lambda} \propto G_F^2 \left( \frac{M_{W_L}}{M_{W_R}} \right)^2 \left( \frac{g_R}{g_L} \right)^2 \sum_{i=1,2,3} U_{ei} T_{ei}^* \frac{1}{|p|},$$

$$A^{N}_{\lambda} \propto G_F^2 \sum_{j=1,2,3} \left( \frac{M_{W_L}}{M_{W_R}} \right)^2 \left( \frac{g_R}{g_L} \right)^2 S_{ej} V_{ej}^* \frac{1}{|p|},$$

The Feynman amplitudes for the suppressed contributions from $\eta$ diagram as displayed in Fig. 4 are given by

$$A^{\nu}_{\eta} \propto G_F^2 \sum_{i=1,2,3} \left( \frac{g_R}{g_L} \right) \tan \xi_U \frac{1}{|p|},$$

$$A^{N}_{\eta} \propto G_F^2 \sum_{j=1,2,3} \left( \frac{g_R}{g_L} \right) \tan \xi_S \frac{1}{|p|},$$
Figure 4. Feynman diagram for $0\nu\beta\beta$ transition due to exchange of $\nu_i$ and $N_j$ via $W_L - W_R$ mixing and left-right neutrino mixing.

| Effective Mass Parameters | Analytic formula |
|---------------------------|------------------|
| $m_{\nu e, L}^\nu$       | $\sum_{i=1}^3 U_{ei}^2 m_i$ |
| $m_{\nu e, L}^N$         | $\sum_{i=1}^3 S_{ei}^2 M_i$ |

Table 2. Effective Majorana mass parameters from purely left-handed currents due to exchange of left-handed and right-handed neutrinos.

3.2 Effective Mass Parameters

| Effective Mass Parameters | Analytic formula |
|---------------------------|------------------|
| $m_{\nu e, R}^\nu$       | $\left(\frac{M_{WL}}{M_{WR}}\right)^2 \left(\frac{g_R}{g_L}\right)^2 \sum_{i=1}^3 U_{ei}^{*2} m_i$ |
| $m_{\nu e, R}^N$         | $\left(\frac{M_{WL}}{M_{WR}}\right)^2 \left(\frac{g_R}{g_L}\right)^2 \sum_{i=1}^3 V_{ei}^{*2} M_i$ |

Table 3. Effective Majorana mass parameters from purely right-handed currents due to exchange of left-handed and right-handed neutrinos.

| Effective Mass Parameters | Analytic formula |
|---------------------------|------------------|
| $m_{\nu e, \lambda}$     | $10^{-2} \left(\frac{M_{WL}}{M_{WR}}\right)^2 \left(\frac{g_R}{g_L}\right)^2 \sum_{i=1}^3 U_{ei} T_{ei}^{*} |p|$ |
| $m_{\nu e, \lambda}^N$   | $10^{-2} \left(\frac{M_{WL}}{M_{WR}}\right)^2 \left(\frac{g_R}{g_L}\right)^2 \sum_{j=1}^3 S_{ej} V_{ej}^{*} |p|$ |
| $m_{\nu e, \eta}$        | $\left(\frac{g_R}{g_L}\right) \sum_{i=1}^3 U_{ei} T_{ei}^{*} \tan \xi |p|$ |
| $m_{\nu e, \eta}^N$      | $\left(\frac{g_R}{g_L}\right) \sum_{j=1}^3 S_{ej} V_{ej}^{*} \tan \xi |p|$ |

Table 4. Effective Majorana mass parameters due $\lambda$ and $\eta$ type diagrams.
Isotope | $G_{01}^{0\nu}$ (yr$^{-1}$) | $\mathcal{M}_{\nu}^{0\nu}$ | $\mathcal{M}_{\nu}^{0\nu}$ | $\mathcal{M}_{\nu}^{0\nu}$ | $\mathcal{M}_{\eta}^{0\nu}$ | $\mathcal{M}_{\eta}^{0\nu}$ |
---|---|---|---|---|---|---|
Ge - 76 | $5.77 \times 10^{-15}$ | 2.58 - 6.64 | 233 - 412 | 1.75 - 3.76 | 235 - 637 |
Xe - 136 | $3.56 \times 10^{-14}$ | 1.57 - 3.85 | 164 - 172 | 1.92 - 2.49 | 370 - 419 |

Table 5. Values of phase space factor and nuclear matrix elements used in the analysis

Combining all the contributions, one can write down the half-life of neutrinoless double beta decay as

$$\frac{1}{T_{01}^{\nu\nu}} = G_{01}^{0\nu} \left( |\mathcal{M}_{\nu}^{0\nu} \eta_{L}^{\nu} + \mathcal{M}_{N}^{0\nu} \eta_{N}^{L}|^2 + |\mathcal{M}_{\nu}^{0\nu} \eta_{R}^{\nu} + \mathcal{M}_{\eta}^{0\nu} \eta_{R}^{\nu}|^2 \right. + |\mathcal{M}_{\nu}^{0\nu} (\eta_{L}^{\nu} + \eta_{N}^{N}) + \mathcal{M}_{\eta}^{0\nu} (\eta_{R}^{\nu} + \eta_{N}^{N})|^2 \right)$$

(3.5)

where

$$\eta_{L}^{\nu} = \sum_{i} \frac{m_{i} U_{ei}^2}{m_{e}}, \quad \eta_{R}^{\nu} = \left( \frac{M_{W_{L}}}{M_{W_{R}}} \right)^{4} \sum_{i} \frac{m_{i} T_{ei}^2}{m_{e}}$$

$$\eta_{L}^{L} = \sum_{i} \frac{S_{ei} M_{i}}{m_{e}}, \quad \eta_{R}^{R} = \left( \frac{M_{W_{L}}}{M_{W_{R}}} \right)^{4} \sum_{i} \frac{V_{ei}^2 M_{i}}{m_{e}}$$

$$\eta_{\lambda}^{\nu} = \left( \frac{M_{W_{L}}}{M_{W_{R}}} \right)^{2} \sum_{i} U_{ei} T_{ei}^{*}, \quad \eta_{N}^{\nu} = \left( \frac{M_{W_{L}}}{M_{W_{R}}} \right)^{2} \sum_{i} S_{ei} V_{ei}^{*}$$

$$\eta_{\lambda}^{\nu} = \tan \xi \sum_{i} U_{ei} T_{ei}^{*}, \quad \eta_{N}^{\nu} = \tan \xi \sum_{i} S_{ei} V_{ei}^{*}$$

Here $m_{e}, m_{p}$ are masses of electron and proton respectively. Also, the nuclear matrix elements involved are denoted by $\mathcal{M}$ the numerical values of which are shown in table 5. The numerical values of the phase space factor $G_{01}^{0\nu}$ are also shown in the table 5 for different nuclei.

### 4 Lepton Flavour Violation

The new fields introduced in the model can induce LFV decays like $\mu \rightarrow e\gamma$ $^{1}$ through one-loop diagrams with heavy charged vector-like leptons and the second scalar doublets in loop. This is shown in figure 5. In the SM, such LFV decays also occur at loop level but heavily suppressed due to the smallness of neutrino masses, far beyond the current experimental sensitivity $^{[110]}$. Therefore, any experimental observation of such rare decay processes will be a clear indication of BSM physics. We calculate the new physics contribution to $\Gamma(\mu \rightarrow e\gamma)$ and check for what values of new physics parameters, it can lie close to the latest bound from the MEG collaboration is $\text{BR}(\mu \rightarrow e\gamma) < 4.2 \times 10^{-13}$ at 90% confidence level $^{[110]}$. The usual standard model contribution with $W_{L}, \nu_{L}$ in loop is given by

$^{1}$For a recent review on charged lepton flavour violation, please see $^{[87]}$
\[
\text{BR}(\mu \rightarrow e\gamma) = \frac{\sqrt{2}G_F m_\mu^5}{\Gamma_\mu} \sum_i U_{\mu i} U_{e i}^* G_\gamma \left( \frac{m_i^2}{M_{W_L}^2} \right)^2
\]

where
\[
\Gamma_\mu = \frac{G_F^2 m_\mu^5}{192\pi^3} \left( 1 - 8 \frac{m_e^2}{m_\mu^2} \right) \left( 1 + \frac{\alpha_{em}}{2\pi} \left( \frac{25}{4} - \pi^2 \right) \right)
\]
is the decay width of the muon and \( G_\gamma(x) \) is the loop function given by
\[
G_\gamma(x) = \frac{x - 6x^2 + 3x^3 + 2x^4 - 6x^3\log x}{4(1 - x)^4}
\]

The SM contribution is very suppressed due to the smallness of neutrino masses, of the order of \( 10^{-46} \). We can have three more contributions from the same diagram (right panel of figure 5) where the particles in the loop can be \( \nu_L - W_R, \nu_R - W_R, \nu_R - W_L \). In the case of \( \nu_L - W_R \), the mass of \( W_L \) will be replaced by that of \( W_R \) and the light neutrino mixing matrix elements \( U_{ai} \) will be replaced by heavy-light neutrino mixing \( T_{ai} \). Similarly, in \( \nu_R - W_L \) case, the mass of light neutrino will be replaced by that of heavy neutrino \( M_i \) and he light neutrino mixing matrix elements \( U_{ai} \) will be replaced by the light heavy mixing \( S_{ai} \). On the other hand, the diagram with \( \nu_R(\equiv N) - W_R \) in the loop, we need to do the substitutions: \( U \rightarrow V, m_i \rightarrow M_i, M_{W_L} \rightarrow M_{W_R} \). For all these cases, the contribution to \( \text{BR}(\mu \rightarrow e\gamma) \) does not come anywhere close to the latest MEG bound. If we use the heavy neutrino mass around a keV, and \( W_R \) mass at a few TeV, we can get a few order of magnitudes enhancement compared to the SM prediction, but still remains much below the experimental sensitivity.

On the other hand, the diagram on the right panel of figure 5 can give a large contribution to this LFV decay that can saturate the present experimental bound. We adopt the general prescriptions given in [124] for the calculation of the LFV decay width. From the scalar potential written above, it is clear that the scalar fields can be rotated to their mass basis as
\[
H_i^0 = O_{ij} \tilde{H}^j = L_{ij} H^j_L + R_{ij} H^j_R.
\]
Now, the Yukawa can be rewritten as
\[ \mathcal{L} \supset Y_E \sum_{i=1}^{4} (L_{ij}^{e_i} L_{j}^{h_i} E_{L} + R_{ij}^{e_i} R_{j}^{h_i} E_{R}). \] (4.3)

In order to derive processes contributing to LFV we need to use the above interaction term and bring it in the following form
\[ -i \bar{l}_\beta (\sigma_L P_L + \sigma_R P_R) \sigma^{\mu\nu} l_\alpha F_{\mu\nu} \] (4.4)

and to calculate the process \( l_\alpha \rightarrow l_\beta \gamma \) we have the following expression
\[ \Gamma(l_\alpha \rightarrow l_\beta \gamma) = \left( \frac{m_\alpha^2 - m_\beta^2}{4\pi m_\alpha^2} \right)^3 \left[ |\sigma_L|^2 + |\sigma_R|^2 \right] \] (4.5)

where \( \sigma_{L,R} \) is given as
\[
\begin{align*}
\sigma_L &= Q_F \left( (m_\alpha O_{RR} + m_\beta O_{LL}) g(t) + m_E O_{RL} h(t) \right) \\
\sigma_R &= Q_F \left( (m_\alpha O_{LL} + m_\beta O_{RR}) g(t) + m_E O_{LR} h(t) \right) \\
O_{LL} &= L_{\beta i}^* L_{\alpha i} \\
O_{RR} &= R_{\beta i}^* R_{\alpha i} \\
O_{LR} &= L_{\beta i}^* R_{\alpha i} \\
O_{RL} &= R_{\beta i}^* L_{\alpha i} \\
g(t) &= \frac{i}{192\pi^2 m_H^2} \left[ \frac{(t-1)(t^2-5t+20)+6t \ln t}{(t-1)^4} \right] \\
h(t) &= \frac{i}{32\pi^2 m_H^2} \left[ \frac{t^2-4t+3+2 \ln t}{(t-1)^3} \right] \\
\bar{g}(t) &= \frac{i}{192\pi^2 m_H^2} \left[ \frac{(t-1)(2t^2+5t-1)+6t^2(2t-1) \ln t}{(t-1)^4} \right] \\
\bar{h}(t) &= \frac{i}{32\pi^2 m_H^2} \left[ \frac{t^2-1-2t \ln t}{(t-1)^3} \right]
\end{align*}
\]

where \( Q_F \) is the electromagnetic charge of the internal fermion fields \( E \) and \( t = (M_E/m_H)^2 \) with \( M_E, m_H \) being the masses of internal fermion and scalar respectively. Although this diagram can give a large contribution to the LFV decay, the fields involved in this dominant process remain decoupled from the ones involved in the \( 0\nu\beta\beta \) process discussed above. Indirectly however, there can be a correlation between these two processes due to the structure of charged lepton mass matrix which is dictated by the structures of \( Y_E, M_E \). If the charged lepton mass matrix is non-diagonal, then it can affect the light neutrino mixing which will then affect the \( 0\nu\beta\beta \) amplitude. We however, keep these two processes decoupled in this work and constrain the relevant parameters from the latest experimental bounds.

5 Numerical Analysis

As discussed in section 2, we need to choose \( M^D_N \neq M^M_N \) in order to have non-vanishing light neutrino masses. For simplicity, we consider \( M^D_N \) and \( M^M_N \) to be equal upto a numerical
In this case, the light neutrino mass formula can be written as

\[ m_\nu = (1 - c_1)M_L \]

This also gives rise to the same diagonalising matrices for both light and heavy neutrino mass matrices \( U_L = U_R \). In such a case, the matrices \( U, S, T, V \) can be written as

\[ U = U_L - \frac{|c_1|^2}{2}U_L, \quad S = c_1 U_R, \quad T = -c_1^* U_L, \quad V = U_R - \frac{|c_1|^2}{2}U_R \]

Now, \( U_L = U_\nu \) can be parametrised as the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) leptonic mixing matrix

\[ U_{PMNS} = U_\nu^\dagger U_\nu \] (5.1)

if the charged lepton mass matrix is diagonal or equivalently, \( U_l = I \). The PMNS mixing matrix can be parametrised as

\[
U_{PMNS} = \begin{pmatrix}
  c_{12}c_{13} & s_{13}e^{-i\delta} & s_{12}s_{13}
  
  -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{13}
  
  s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}
\end{pmatrix} U_{Maj}
\] (5.2)

where \( c_{ij} = \cos \theta_{ij} \), \( s_{ij} = \sin \theta_{ij} \) and \( \delta \) is the leptonic Dirac CP phase. The diagonal matrix \( U_{Maj} = \text{diag}(1, e^{i\alpha}, e^{i(\beta + \delta)}) \) contains the Majorana CP phases \( \alpha, \beta \) which remain undetermined at neutrino oscillation experiments. The light neutrino masses that appear in the \( 0\nu\beta\beta \) half-life can be written in terms of the lightest neutrino mass and the experimentally measured mass squared differences. For normal hierarchy, the diagonal mass matrix of the light neutrinos can be written as

\[ m_\nu^{\text{diag}} = \text{diag}(m_1, \sqrt{m_1^2 + \Delta m_{21}^2}, \sqrt{m_1^2 + \Delta m_{31}^2}) \]

whereas for inverted hierarchy it can be written as

\[ m_\nu^{\text{diag}} = \text{diag}(\sqrt{m_3^2 + \Delta m_{23}^2 - \Delta m_{21}^2}, \sqrt{m_3^2 + \Delta m_{23}^2}, m_3) \]

The heavy neutrino masses are related to the light neutrinos as

\[ M_i = \frac{v_R^2}{v_L^2} \frac{1}{1 - c_1} m_i . \] (5.3)

Thus, once we choose the scale of left-right symmetry or \( v_R \), we can calculate the new physics contributions to \( 0\nu\beta\beta \) by varying the lightest neutrino mass, the CP phases for different values of \( c_1 \). The other parameters like the mixing angles and mass squared differences can be varied in their 3\( \sigma \) global fit range given in table 6.

### 5.1 0\nu\beta\beta with \( M_{W_R} \) around 3.5 TeV

In the scenario where \( M_{W_R} \approx 3.5 \text{ TeV} \), all possible new physics contribution (apart from the usual light neutrino one in the SM) to \( 0\nu\beta\beta \) decay can be there except the \( \eta \)-diagrams
Parameters | Normal Hierarchy (NH) | Inverted Hierarchy (IH)  
---|---|---  
$\Delta m^2_{31}/10^{-3} \text{eV}^2$ | 2.407 – 2.643 | 2.399 – 2.635  
$\Delta m^2_{21}/10^{-5} \text{eV}^2$ | 7.03 – 8.09 | 7.02 – 8.09  
$\sin^2 \theta_{12}$ | 0.271 – 0.345 | 0.271 – 0.345  
$\sin^2 \theta_{23}$ | 0.385 – 0.635 | 0.393 – 0.640  
$\sin^2 \theta_{13}$ | 0.01934 – 0.02392 | 0.01953 – 0.02408  
$\delta$ | 0 – 2$\pi$ | 0 – 2$\pi$  

Table 6. Global fit 3$\sigma$ values of neutrino oscillation parameters [111].

Figure 6. Mixed helicity contribution (so called $\lambda$ diagram) to effective Majorana mass of neutrinoless double beta decay with the variation of lightest neutrino mass displayed as blue dots. We compare the standard light neutrino exchanged mechanism for $0\nu\beta\beta$ transition with red and green band for the 3$\sigma$ oscillation data allowed ranges in a three light neutrino scheme for normal hierarchy (NH) and for inverted hierarchy (IH), respectively. We fix $M_W$ around 3.5 TeV, $c_1 \approx 0.001$ and heaviest right-handed neutrino around keV scale. We show the limit on the sum of light neutrino masses from the Planck mission data as $\sum_i m_i < 0.17 \text{eV}$ [112] represented by vertical line and the respective shaded area.

which are proportional to $W_L – W_R$ mixing. This is due to suppressed contribution from $W_L – W_R$ mixing as this mixing is order of $\leq 10^{-7}$ and hence we ignore it from our analysis.

From right-handed currents:
With $M_i \ll |p|$, the neutrino propagator part simplifies as,

$$P_R \frac{\not{p} + M_i}{p^2 - M_i^2} P_R \approx \frac{M_i}{p^2}, \quad (5.4)$$

resulting the dimensionless particle physics parameter $\eta_N$ due to exchange of right-handed
Figure 7. Heavy neutrino contributions to the half-life of neutrinoless double beta decay with the variation of lightest neutrino mass displayed as yellow dots. Here the yellow dots are chosen such a way they saturate the experimental limit on the half-life (represented by horizontal lines for the GERDA Phase-II and KamLAND-Zen experiments). We fix $M_{W_R}$ around 4 TeV, $c_1 \approx 0.001$ and heaviest right-handed neutrino around keV scale. For comparison, we displayed standard contribution with brown and the purple areas, respectively, for the $3\sigma$ oscillation data allowed ranges in a three light neutrino scheme for normal hierarchy (NH) and for inverted hierarchy (IH), respectively. We show the limit on the sum of light neutrino masses from the Planck mission data as $\sum_i m_i < 0.17$ eV [112] represented by vertical line and the respective shaded area.

neutrinos via right-handed currents,

$$\eta_N \approx \frac{1}{m_e} \left( \frac{M_{W_L}}{M_{W_R}} \right)^4 \sum_{i=1}^{3} U_{ei}^2 M_i \propto \eta_{\nu_L}. \quad (5.5)$$

The corresponding effective mass parameter is given by

$$m_{ee}^N \approx \left( \frac{M_{W_L}}{M_{W_R}} \right)^4 \sum_{i=1}^{3} U_{ei}^2 M_i \propto m_{ee}^{\nu_L}. \quad (5.6)$$

With mass formula for charged gauged bosons as, $M_{W_R} \approx \frac{1}{2} g_R v_R$ and $M_{W} \approx \frac{1}{2} g_L v_L$, we fix $M_{W_R} \simeq 3.5$ TeV by choosing $g_R \approx g_L = 0.65$ and $v_R \approx 10$ TeV. As mentioned above, the mass relation between light and heavy neutrinos is $M_i = \frac{v_L^2}{v_R^2} \frac{1}{1-c_1} m_i$. Thus, the maximum value of heavy right-handed Majorana mass is found to be around few keV with $M_{W_R} \simeq 3.5$ TeV, unless the dimensionless parameter $c_1$ is very fine tuned to a value close to unity. Therefore, $0\nu\beta\beta$ contribution from purely right handed current is very much suppressed in this case.

From $\lambda$ diagram with large light-heavy neutrino mixing:
Figure 8. The effective Majorana mass parameter as a function of lightest neutrino mass arising from purely left-handed current effects due to exchange of right-handed Majorana neutrinos which are displayed as pink and magenta dots. We compare the standard light neutrino exchanged mechanism for $0\nu\beta\beta$ transition with red and green band for the 3$\sigma$ oscillation data allowed ranges in a three light neutrino scheme for normal hierarchy (NH) and for inverted hierarchy (IH), respectively. We fix $M_{W_R}$ around 100 TeV, $c_1 \approx 0.001$ and heaviest right-handed neutrino around MeV scale. We show the limit on the sum of light neutrino masses from the Planck mission data as $\sum_i m_i < 0.17$ eV [112] represented by vertical line and the respective shaded area.

The $\lambda$--diagram with $c_1 = 0.001$, $|p| \simeq 100$ MeV and $M_{W_R} \simeq 3.5$ TeV is estimated to be,

$$m_{ee}^\lambda \approx 10^{-2} \left( \frac{M_{W_L}}{M_{W_R}} \right)^2 \sum_{i=1}^{3} c_1 U_{ei} |p| \approx 0.25 \text{ eV} . \quad (5.7)$$

This can clearly saturate the experimental bound from the KamLAND-Zen experiment on effective neutrino mass $m_{ee} < 61 - 165$ meV [107] as well as the corresponding lower limit on half-life. This contribution is shown in figure 6 and 7 for $m_{ee}$ and $T_{1/2}$ respectively.

**From $W_L - N$ mediated diagram:**

Another contribution to $0\nu\beta\beta$ due to exchange of heavy right-handed keV Majorana neutrinos via purely left-handed currents is given by

$$m_{ee,L}^N \approx \sum_{i=1}^{3} c_1^2 U_{ei}^2 M_i \approx c_1^2 \times 5 \times 10^2 \text{ eV} . \quad (5.8)$$

This can saturate the experimental bound on $0\nu\beta\beta$ transition only when $c_1 > 0.001$, but such large light-heavy neutrino mixing will give too large a contribution to $0\nu\beta\beta$ with $M_{W_R} \simeq 3.5$ TeV as derived from $\lambda$--diagram above. Thus, we fix $c_1 \approx 0.001$ such that $W_R - N, W_L - N$ contributions are suppressed whereas the $\lambda$--contribution is saturating the experimental limit on the half-life as well as effective neutrino mass (represented by horizontal bands parallel to x-axis).
Figure 9. Heavy MeV range neutrino contributions to the half-life of neutrinoless double beta decay with the variation of lightest neutrino mass displayed as cyan dots (for NH) and yellow dots (for IH). We choose these yellow and cyan dots such a way they saturate the experimental limit on the half-life (represented by horizontal lines for the GERDA Phase-II and KamLAND-Zen experiments). We fix $M_{W_R}$ around 100 TeV, $c_1 \approx 0.001$ and heaviest right-handed neutrino around keV scale. For comparison, we displayed standard contribution with brown and the purple areas, respectively, for the $3\sigma$ oscillation data allowed ranges in a three light neutrino scheme for normal hierarchy (NH) and for inverted hierarchy (IH), respectively. We show the limit on the sum of light neutrino masses from the Planck mission data as $\sum m_i < 0.17$ eV [112] represented by vertical line and the respective shaded area.

The contributions to $0\nu\beta\beta$ for this benchmark choice are presented in figure 6 and 7 with the standard light neutrino contribution as well as the $\lambda -$ diagram from new physics contribution.

5.2 $0\nu\beta\beta$ with $M_{W_R}$ around 100 TeV

With $M_{W_R} \simeq 100$ TeV which corresponds to heavy neutrino masses $M_{i>} \approx$ MeV, all the contributions to $0\nu\beta\beta$ due to purely right-handed currents, $\lambda$ and $\eta$ diagrams are negligible as they are proportional to $1/M_{W_R}^4$, $1/M_{W_R}^2$ and $1/M_{W_R}^2$, respectively. Thus, the only new physics contribution to $0\nu\beta\beta$ arises from left-handed current effects due to the exchange of right-handed Majorana neutrinos with MeV mass. The corresponding effective mass parameter is given by

$$m_{ee,L}^N \approx \sum_{i=1}^3 c_i^2 U_{ei}^2 M_i \approx 0.35 \text{ eV}$$

(5.9)

for $c_1 \approx 0.001$. This contribution is displayed in figure 8 which saturates the GERDA Phase II and KamLAND-Zen experimental bound and the corresponding half-life is presented in figure 9. This is very interesting feature of this model LRSM-US as we can have observable $0\nu\beta\beta$ even if the scale of right handed gauge bosons is beyond collider reach.
5.3 Full parameter scan for $0\nu\beta\beta$

After showing the new physics contribution to $0\nu\beta\beta$ for a few benchmark values, we perform a complete scan of parameter space and constrain the parameters from the requirement of satisfying the latest bounds from $0\nu\beta\beta$ experiment. We vary the light neutrino parameters
in their $3\sigma$ range given in table 6 and also vary the new physics parameters randomly in the following ranges

$$M_{WR} \in (3 \text{ TeV}, 100 \text{ TeV}), \quad c_1 \in (10^{-5}, 1), \quad m_{\text{lightest}} \in (10^{-5} \text{ eV}, \ 1 \text{ eV}).$$

The resulting parameter space is shown for both the hierarchies of light neutrino masses in figure 10. The lower bound on $M_{WR}$ is chosen to be near the LHC bound. It is interesting to see that the parameter $c_1$ is constrained to be below 0.01 for the entire mass range of $W_R$ between $3 - 100 \text{ TeV}$. Also, the upper bound on $c_1$ becomes weaker as the mass of $W_R$ is increased from 3 to around tens of TeV, as expected. However, it again becomes stronger and stronger as the mass of $W_R$ is increased further, all the way till 100 TeV. This is due to the increasing contribution of the heavy light neutrino mixing diagram (shown in the right panel of figure 1) which is directly proportional to the masses of right handed neutrinos. As $W_R$ mass increases, $v_R$ also increases thereby increasing the right handed neutrino masses and hence the upper bound on $c_1$ becomes more strict. It can also be seen that this scan plot in figure 10 is also compatible with benchmark scenarios shown in earlier plots. For example, in $M_{WR} \approx 3 - 4 \text{ TeV}$ as well as $M_{WR} = 100 \text{ TeV}$ cases, $c_1 = 0.001$ is not allowed from experimental bounds on $0\nu\beta\beta$, as shown by the plots in lower panels of figure 10. This is also visible from figure 6, 7, 8 and 9 respectively, where one or more of the specific contributions to $0\nu\beta\beta$ violate the experimental bounds.

### 5.4 Allowed parameter space from LFV bounds

Since the gauge boson contribution to the lepton flavour violating decay $\mu \to e\gamma$ (shown in the right panel of figure 5), is very suppressed compared to the experimental sensitivity for all the region of our interest, we numerically evaluate the scalar mediated contribution to it and constrain the relevant Yukawa couplings as well as the vector-like charged lepton mass from the requirement of satisfying the MEG 2016 bound [110]. Apart from the SM Higgs boson having mass 125 GeV, the other neutral scalar going in loop is assumed to be of mass 200 GeV. Taking all generation Yukawa couplings to be same for simplicity, we then vary the Yukawa $Y_E$ and $M_E$ and calculate the BR($\mu \to e\gamma$) numerically. The resulting parameter space satisfying the MEG 2016 bound is shown in figure 11.

### 6 Summary and Conclusion

We have studied the new physics contribution to neutrinoless double beta decay and charged lepton flavour violation in a minimal left-right symmetric model where all the fermions acquire masses through a universal seesaw mechanism where all the fermions acquire masses through a universal seesaw mechanism. Due to different ways of generating fermion mass and breaking the left-right gauge symmetry all the way down to the SM one, the leading contributions to $0\nu\beta\beta$ and LFV decay are very different in this model, in comparison to the usual LRSM with type I and type II seesaw for neutrino masses. One interesting feature of this model is the presence of light right handed neutrinos in the keV-MeV range, even if the scale of the theory is as high as 100 TeV. Since the heavy neutrinos have masses lighter than typical momentum exchange of $0\nu\beta\beta$ process, their contributions to $0\nu\beta\beta$ becomes different compared to the usual LRSM with much heavier
right handed neutrinos. We identify all possible diagrams contributing to $0\nu\beta\beta$ in this model and find that the dominant contribution comes from the diagrams with light-heavy neutrino mixing. For certain choices of parameters, the model predictions can saturate the current experimental bounds on $0\nu\beta\beta$, even if the $W_R$ mass is as heavy as 100 TeV. This offers a very interesting probe of such high scale LRSM which may not be directly accessible at collider experiments. We also check the possible sources of charged lepton flavour violating decay $\mu \rightarrow e\gamma$ which is more tightly constrained from present experimental data. We find that the charged gauge boson plus neutral fermion mediated diagrams to this process remain suppressed for all region of parameter space while the neutral scalar plus charged vector-like lepton mediated diagrams can saturate the experimental bound and hence can be probed at ongoing or near future experiments.

The model also can have very interesting cosmological signatures due to the existence of light right handed neutrinos. For example, a keV scale right handed neutrino could be long lived enough to play the role of warm dark matter [125, 126]. On the other hand, if one of the right handed neutrinos become as light as an eV, for some region of parameter space, it can affect the Planck bound on the relativistic degrees of freedom $N_{\text{eff}} = 3.15 \pm 0.23$ [112]. Since the right and left handed neutrino masses are proportional to each other, such a eV scale right handed neutrino may arise if the lightest left handed neutrino mass is vanishingly small. Such a situation will not only be interesting for neutrino oscillation experiments, but will also constrain the $W_R$ mass in order to satisfy the Planck bound on $N_{\text{eff}}$ [97, 99]. If the right handed neutrinos are in the sub GeV regime, they can also play a role in creating baryon asymmetry through neutrino oscillations [127]. We leave such studies in this particular model for future works.

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