Neutrinoless double beta decay in LRSM with natural type-II seesaw dominance

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ABSTRACT: We present a detailed discussion on neutrinoless double beta decay within a class of left-right symmetric models where neutrino mass originates by natural type-II seesaw dominance. The spontaneous symmetry breaking is implemented with doublets, triplets and bidoublet scalars. The fermion sector is extended with an extra sterile neutrino per generation that helps in implementing the seesaw mechanism. The presence of extra particles in the model exactly cancels type-I seesaw and allows large value for Dirac neutrino mass matrix M_D. The key feature of this work is that all the physical masses and mixing are expressed in terms of neutrino oscillation parameters and lightest neutrino mass thereby facilitating to constrain light neutrino masses from $0\nu\beta\beta$ decay. With this large value of M_D new contributions arise due to; i) purely left-handed current via exchange of heavy right-handed neutrinos as well as sterile neutrinos, ii) the so called λ and η diagrams. New physics contributions also arise from right-handed currents with right-handed gauge boson W_R mass around $3\,\text{TeV}$. From the numerical study, we find that the new contributions to $0\nu\beta\beta$ decay not only saturate the current experimental bound but also give lower limit on absolute scale of lightest neutrino mass and favor NH pattern of light neutrino mass hierarchy.

KEYWORDS: Beyond Standard Model, Neutrino Physics, Spontaneous Symmetry Breaking

ARXIV EPRINT: 1607.07655

Contents

1 Introduction

The discovery that neutrinos have mass and they mix with each other has put before us another vital question to speculate over; whether they are Dirac or Majorana [1] particles. Even more intriguing is the theoretical origin of such a tiny mass and the mass hierarchy among them. The different seesaw mechanisms like type-I [2–5], type-II [6–10] and others that appropriately explain this tiny mass further require them to be Majorana particles.

On the contrary, Majorana nature of neutrinos violates global lepton number by 2 units that is regarded as an accidental symmetry within the Standard Model (SM). This leads to the search of a rare process called Neutrinoless Double Beta Decay ($0\nu\beta\beta$) that only can assuredly endorse the Majorana nature of neutrinos and lepton number violation in nature [11]. While new theories are trying to find new physics contributions to $0\nu\beta\beta$ decay, the experiments are looking for lower limits on the half-lives being decayed. Of yet, GERDA [12] using Ge^{76} gives lower limit on half life of $0\nu\beta\beta$ decay as $T_{1/2}^{0\nu} > 2.1 \times 10^{25}$ yrs at 90% C.L. whereas the limits provided by EXO-200 [13] and KamLAND [14] are $T_{1/2}^{0\nu} > 1.6 \times 10^{25}$ yrs and $T_{1/2}^{0\nu} > 1.9 \times 10^{26}$ yrs respectively. The combined limit from KamLAND-Zen comes to be $T_{1/2}^{0\nu} > 3.4 \times 10^{26}$ yrs at 90% C.L. This process can be mediated by the exchange of a light Majorana neutrinos or by new particles appearing in various extensions of SM [15–20, 20–28].

Within preview of BSM physics, left-right symmetric models (LRSM) [2, 29–33] are found to be best suited frameworks for explaining the origin of maximal parity violation in weak interactions and the origin of small neutrino mass. This class of models, based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_C$, when studied at TeV scale interlinks high energy collider physics to low energy phenomena like neutrinoless double beta decay and other LFV processes (see refs. [34–60]). Moreover, the left-right symmetric models can also accommodate stable dark matter candidate contributing 25% energy budget of the Universe [61–65]. In conventional left-right symmetric models where symmetry breaking is implemented with scalar triplets and bidoublet, the light neutrino mass is governed by type-I plus type-II seesaw mechanisms

$$m_\nu = -M_D M_R^{-1} M_D^T + M_L = m_\nu^{\mathrm{I}} + m_\nu^{\mathrm{II}}.$$

Here $M_L(M_R)$ is the Majorana mass term for light left-handed (heavy right-handed) Majorana neutrinos arising from respective VEVs of left-handed (right-handed) scalar triplets and M_D is the Dirac neutrino mass matrix connecting light-heavy neutrinos. The scale of M_R is decided by the vacuum expectation value of right-handed scalar triplet which spontaneously breaks LRSM to SM. Thus, the smallness of light neutrino mass is connected to high scale of parity restoration i.e, 10^{15} GeV clearly making it inaccessible to current and planned accelerator experiments. Moreover when LRSM breaks around TeV scale, the gauge bosons W_R, Z_R, right-handed neutrinos N_R and scalar triplets $\Delta_{L,R}$ get mass around that scale allowing several lepton number violating signatures at high energy as well as low energy experiments. A wide range of literature provides discussions on neutrinoless double beta decay within TeV scale LRSM assuming type-I seesaw dominance [42] or type-I plus type-II [35, 42, 43, 47, 48, 66, 67] seesaw mechanisms. Some more scenarios have been studied in [34–36, 39, 42, 54, 68, 69] where type-II seesaw dominance relates the light and heavy neutrinos with each other. Other works that discuss complementarity study of lepton number, lepton flavour violation and collider signatures in LRSM with spontaneous D-parity breaking mechanism also embed the framework in a non-SUSY SO(10) GUT [36–38, 70]. One should bear in mind that the new physics contributions to neutrinoless double beta decay mainly involves left-right mixing (or light-heavy neutrino mixing) which crucially depends on Dirac neutrino mass M_D. Necessarily M_D should be

large in order to expect LNV signatures at colliders. Contrary to this, the type-II seesaw dominance can be realized with suppressed value of M_D or with very high scale of parity restoration. Studies that assume $M_D \rightarrow 0$ therefore miss to comment on LNV, LFV and Collider aspects involving left-right mixing. We thus feel motivated to explore alternative class of left-right symmetric models which allows large value of M_D and carries light and heavy neutrinos proportional to each other.

This work considers a TeV scale LRSM where symmetry breaking is implemented with scalar bidoublet Φ, doublets $H_{L,R}$ and triplets $\Delta_{L,R}$. The scalar bidoublet carrying $B - L$ charge 0 provides Dirac masses to charged fermions as well as to neutrinos. The scalar triplets with $B - L$ charge 2 units provide Majorana masses to light and heavy neutrinos. One extra sterile fermion S_L per generation also finds place in the model that help in implementing extended type-II seesaw mechanism. The scalar doublets $H_{L,R}$ play the same role as S_L. An interesting feature of this new class of LRSM is that it provides possibility of achieving type-II seesaw dominance when parity and $\text{SU}(2)_R$ break at same scale. Moreover this framework allows large value for Dirac neutrino mass matrix M_D thereby leading to new physics contributions to neutrinoless double bea decay i.e, i) from purely left-handed currents via exchange of heavy right-handed and extra sterile neutrinos, ii) from purely right handed currents via exchange of heavy right-handed neutrinos, iii) from so called λ and η diagrams. This work aims to carefully analyze the new contributions to $0\nu\beta\beta$ in order to derive the absolute scale of light neutrino masses and mass hierarchy.

The complete work is structured as follows. In section 2, we briefly discuss the generic and TeV scale LRSMs in context of neutrino mass and associated lepton number violation. Section 3 highlights the natural realization of type-II seesaw dominance. Section 4 lays out the basic ingredients for neutrinoless double beta decay and the calculation of Feynman amplitudes. Section 5 and section 6 are devoted towards the numerical study of LNV $0\nu\beta\beta$ contributions within the present framework. In section 8 we summarize our results.

2 The left-right symmetric model and lepton number violation

The left-right symmetric model [2, 29–33] is based on the gauge group

$$\mathcal{G}_{LR} \equiv \text{SU}(2)_L \times \text{SU}(2)_R \times \text{U}(1)_{B-L} \times \text{SU}(3)_C \,. \tag{2.1}$$

In this class of models, the difference between baryon B and lepton L number is defined as a local gauge symmetry. The electric charge Q is defined as

$$Q = T_{3L} + T_{3R} + \frac{B - L}{2} = T_{3L} + Y \,. \tag{2.2}$$

Here, T_{3L} and T_{3R} are, respectively, the third component of isospin of the gauge groups $\text{SU}(2)_L$ and $\text{SU}(2)_R$, and Y is the hypercharge. The usual leptons and quarks are given by

$$\ell_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} \sim (\mathbf{2}, \mathbf{1}, -\mathbf{1}, \mathbf{1}) \,, \qquad \ell_R = \begin{pmatrix} \nu_R \\ e_R \end{pmatrix} \sim (\mathbf{1}, \mathbf{2}, -\mathbf{1}, \mathbf{1}) \,, \tag{2.3}$$

$$q_L = \begin{pmatrix} u_R \\ d_R \end{pmatrix} \sim \left(\mathbf{2}, \mathbf{1}, \frac{\mathbf{1}}{\mathbf{3}}, \mathbf{3}\right) \,, \qquad q_R = \begin{pmatrix} u_R \\ d_R \end{pmatrix} \sim \left(\mathbf{1}, \mathbf{2}, \frac{\mathbf{1}}{\mathbf{3}}, \mathbf{3}\right) \,. \tag{2.4}$$

The left-right symmetry calls for the presence of right-handed neutrinos and this makes the model suitable for explaining light neutrino masses. For generating fermion masses one needs a scalar bidoublet Φ with the following matrix representation

$$\Phi \equiv \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix} \sim (\mathbf{2}, \mathbf{2}, \mathbf{0}, \mathbf{1}). \tag{2.5}$$

The relevant Yukawa interactions are expressed as,

$$-\mathcal{L}_{\text{Yuk}} \supset \overline{q_L} \left[Y_1 \Phi + Y_2 \widetilde{\Phi} \right] q_R + \overline{\ell_L} \left[Y_3 \Phi + Y_4 \widetilde{\Phi} \right] \ell_R + \text{h.c.}, \tag{2.6}$$

where $\widetilde{\Phi} = \sigma_2 \Phi^* \sigma_2$ and σ_2 is the second Pauli matrix. The scalar bidoublet takes a non-zero VEV as,

$$\langle \Phi \rangle = \begin{pmatrix} v_1 & 0 \\ 0 & v_2 \end{pmatrix}, \tag{2.7}$$

it yields masses for quarks and charged leptons as

$$\begin{aligned} M_u &= Y_1 v_1 + Y_2 v_2^*, \qquad M_d = Y_1 v_2 + Y_2 v_1^*, \\ M_e &= Y_3 v_2 + Y_4 v_1^*. \end{aligned} \tag{2.8}$$

One can generate Dirac masses for light neutrinos using scalar bidoublet as

$$M_D^\nu \equiv M_D = Y_3 v_1 + Y_4 v_2^*. \tag{2.9}$$

However, the Majorana masses for neutrinos depend crucially on how spontaneous symmetry breaking of LRSM down to the SM i.e, $\mathcal{G}_{LR} \to \mathcal{G}_{\text{SM}}$ is implemented.

2.1 Lepton number violation and the origin of neutrino mass

The spontaneous symmetry breaking of LRSM to SM goes in favor of neutrino mass generation and associated lepton number violation. This happens in the following three ways

- with Higgs doublets $H_L(\mathbf{2}, \mathbf{1}, -\mathbf{1}, \mathbf{1}) \oplus H_R(\mathbf{1}, \mathbf{2}, -\mathbf{1}, \mathbf{1})$,

- with scalar triplets $\Delta_L(\mathbf{3}, \mathbf{1}, \mathbf{2}, \mathbf{1}) \oplus \Delta_R(\mathbf{1}, \mathbf{3}, \mathbf{2}, \mathbf{1})$,

- with the combination of doublets and triplets $H_L \oplus H_R$ and $\Delta_L \oplus \Delta_R$.

In the first case, H_R breaks the LR symmetry while the left-handed counterpart is required for left-right invariance. Though this framework holds a minimal scalar spectrum it lacks Majorana mass for neutrinos and thus forbids any signature of lepton number violation or neutrinoless double beta decay. Since the light neutrinos here owe their identity to Dirac fermions, their masses can only be explained by adjusting Yukawa couplings through the non-zero VEVs of scalar bidoublet. Other important roles that this scalar bidoublet plays are to break the SM gauge symmetry to low energy theory and provide the masses to

charged fermions. Using the Yukawa interactions given in eq. (2.6) and with $Y_3 \ll Y_4$, $v_2 \ll v_1$ and $\theta_1, \theta_2 = 0$, the masses for charged leptons and the light neutrinos are given by

$$M_e \simeq Y_4 v_1^*, \qquad M_D \simeq v_1 \left(Y_3 + M_e \frac{v_2}{v_1^2} \right). \tag{2.10}$$

However, a pleasant situation arises in the second case where Δ_R carrying $B - L$ charge 2 breaks the LR symmetry to SM. The inclusion of Δ_L and Δ_R in the framework generate Majorana masses for light as well as heavy neutrinos and thus violate lepton number by two units. This calls for a possibility of smoking-gun same-sign dilepton signatures at collider as well as neutrinoless double beta decay in low energy experiments. The interaction terms involving scalar triplets and leptons are given by

$$-\mathcal{L}_{\text{Yuk}} \supset f_{ij} \left[\overline{(\ell_{Li})^c} \ell_{Lj} \Delta_L + \overline{(\ell_{Ri})^c} \ell_{Rj} \Delta_R \right] + \text{h.c.}. \tag{2.11}$$

Using eq. (2.6) and eq. (2.11), the resulting mass matrix for neutral leptons in the basis (ν_L, N_R^c) reads as

$$\mathcal{M}_\nu = \begin{pmatrix} M_L & M_D \\ M_D^T & M_R \end{pmatrix}, \tag{2.12}$$

where, M_D is the Dirac neutrino mass matrix, $M_L(M_R)$ is the Majorana mass matrix arising from the non-zero VEV of LH (RH) scalar triplet. After diagonalization, the resulting light neutrino mass can be written as a combination of canonical type-I and type-II seesaw formula

$$m_\nu = -M_D M_R^{-1} M_D^T + M_L = m_\nu^I + m_\nu^{II}, \tag{2.13}$$

where, m_ν^I (m_ν^{II}) is denoted as the type-I (type-II) contribution to light neutrino masses,

$$m_\nu^I = -M_D M_R^{-1} M_D^T, \qquad m_\nu^{II} = f \, v_L = f \langle \Delta_L^0 \rangle.$$

In conventional left-right symmetric models, where parity and $\text{SU}(2)_R$ break at same scale, the analytic formula for induced VEV of left-handed scalar triplet Δ_L is given by,

$$v_L \simeq \gamma \frac{v^2}{v_R}.$$

In the above expression $v = \sqrt{v_1^2 + v_2^2}$ lies around electroweak scale, v_R is the VEV of right-handed scalar triplet $\langle \Delta_R \rangle$ and γ is dimensionless Higgs parameter. In order to be consistent with oscillation data $m_\nu^{II} = f v_L$ should be order of 0.1 eV and assuming natural values of f and γ, this sub-eV scale of v_L can be attained only if v_R lies around 10^{14} GeV. However such a high scale is inaccessible to LHC and thus urges to look for TeV scale LRSM. These frameworks offer numerous opportunities like low scale seesaw mechanism, LNV like neutrinoless double beta decay and its collider complementarity and have been already explored by the works mentioned in refs [28, 34–49, 54–58]. Many of the works considered either type-I seesaw dominance or type-II seesaw dominance for en extensive study of $0\nu\beta\beta$ decay. In manifest left-right symmetric model, where right-handed scale lies at TeV range,

the neutrino mass mechanism via type-I plus type-II seesaw gives negligible value to the left-right mixing. As a result of this the production cross-section of heavy neutrinos and the lepton number violating processes at LHC get suppressed. However, the extension of type-I plus type-II seesaw scheme by the inclusion of another sterile neutrino per generation changes the scenario which results large left-right mixing. Now the neutrino mass arises only from type-II seesaw dominance since type-I seesaw contribution gets exactly canceled out. We propose a new framework where type-II seesaw dominance is achieved naturally and allows large value of Dirac neutrino mass which additionally contributes to $0\nu\beta\beta$ decay from purely left-handed current via exchange of heavy neutrinos as well as from the so called λ type and η type diagrams.

3 Extended seesaw mechanism and natural type-II seesaw dominance

3.1 Extended seesaw mass matrix

In order to implement the extended seesaw mechanism[1] within left-right symmetric models, one has to add a complete left-right gauge symmetry singlet neutral fermion S_L per generation to the usual quarks and leptons. Along with this the Higgs sector includes scalar bidoublet Φ with $B - L = 0$, scalar triplets $\Delta_L \oplus \Delta_R$ with $B - L = 2$ and scalar doublets $H_L \oplus H_R$ with $B - L = -1$. The complete particle spectrum is given in table 1.

The relevant leptonic Yukawa interaction terms for extended seesaw mechanism are given by

$$
\begin{aligned}
-\mathcal{L}_{\text{Yuk}} = &\, \overline{\ell_L} \left[Y_3 \Phi + Y_4 \widetilde{\Phi} \right] \ell_R + f \left[\overline{(\ell_L)^c} \ell_L \Delta_L + \overline{(\ell_R)^c} \ell_R \Delta_R \right] \\
&+ F \, \overline{(\ell_R)} H_R S_L^c + F' \, \overline{(\ell_L)} H_L S_L + \mu_S \overline{S_L^c} S_L + \text{h.c.}.
\end{aligned}
\tag{3.1}
$$

$$
\begin{aligned}
\supset &\, M_D \overline{\nu_L} N_R + M_L \overline{\nu_L^c} \nu_L + M_R \overline{N_R^c} N_R \\
&+ M \overline{N_R} S_L + \mu_L \overline{\nu_L^c} S_L + \mu_S \overline{S_L^c} S_L
\end{aligned}
\tag{3.2}
$$

After spontaneous symmetry breaking, the resulting neutral lepton mass matrix for extended seesaw mechanism in the basis (ν_L, N_R^c, S_L) is given by

$$
\mathbb{M}_\nu = \begin{pmatrix} M_L & M_D & \mu_L \\ M_D^T & M_R & M^T \\ \mu_L^T & M & \mu_S \end{pmatrix},
\tag{3.3}
$$

where $M_D = Y \langle \Phi \rangle$ is the Dirac neutrino mass matrix connecting left-handed light neutrinos with right-handed heavy neutrinos, $M_N = f \, v_R = f \, \langle \Delta_R \rangle$ $(M_L = f \, v_L = f \, \langle \Delta_L \rangle)$ is the Majorana mass term for heavy (light) neutrinos, $M = F \langle H_R \rangle$ is the $N - S$ mixing matrix, $\mu_L = F' \langle H_L \rangle$ is the small mass term connecting $\nu - S$ and μ_S is the bare Majorana mass term for extra singlet fermion.

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[1]The discussion of extended seesaw mechanism can be found in refs. [71, 72].

	Fields	SU(2)$_L$	SU(2)$_R$	$B-L$	SU(3)$_C$
Fermions	q_L	2	1	1/3	3
	q_R	1	2	1/3	3
	ℓ_L	2	1	-1	1
	ℓ_R	1	2	-1	1
	S_L	1	1	0	1
Scalars	Φ	2	2	0	1
	H_L	2	1	-1	1
	H_R	1	2	-1	1
	Δ_L	3	1	2	1
	Δ_R	1	3	2	1

Table 1. LRSM representations of extended field content.

Inverse seesaw. In eq. (3.3), if we assume $M_L, M_R, \mu_L \to 0$ and the mass hierarchy $M \gg M_D \gg \mu_S$, we will arrive at the inverse seesaw mass formula for light neutrinos [73]

$$m_\nu = \left(\frac{M_D}{M} \right) \mu \left(\frac{M_D}{M} \right)^T .$$

The light neutrino mass can be parametrized in terms of model parameters of inverse seesaw framework as,

$$\left(\frac{m_\nu}{0.1\,\text{eV}} \right) = \left(\frac{M_D}{100\,\text{GeV}} \right)^2 \left(\frac{\mu}{\text{keV}} \right) \left(\frac{M}{10^4\,\text{GeV}} \right)^{-2} .$$

This expression bears M of few TeV which allows large left-right mixing and thus leads to interesting testable collider phenomenology. Extension of such a scenario has been discussed in the context of allowing large LNV and LFV in the work [38].

Linear seesaw. Similarly in eq. (3.3), if we assume $M_L, M_R, \mu_S \to 0$, the linear seesaw mass formula for light neutrinos is given by [55]

$$m_\nu = M_D^T M^{-1} \mu_L + \text{transpose} , \tag{3.4}$$

whereas the heavy neutrinos form pair of pseudo-Dirac states with masses

$$M_\pm \approx \pm M + m_\nu . \tag{3.5}$$

The following discussion considers the same eq. (3.3) with the assumption that $\mu_L, \mu_S \to 0$ which leads to natural realization of type-II seesaw dominance allowing large left-right mixing.

3.2 Natural realization of type-II seesaw

The natural realization of type-II seesaw dominance is considered here within a class of left-right symmetric models where both discrete left-right parity symmetry and SU(2)$_R$

gauge symmetry break at same scale. The scalar sector is comprising of SU(2) doublets $H_{L,R}$, triplets $\Delta_{L,R}$ and bidoublet Φ whereas the fermion sector is extended with one neutral fermion S_L per generation which is complete singlet under both LRSM as well as SM gauge group. We denote this class of LR model as *Extended LR models* and thus, the corresponding seesaw formula which is type-II dominance in this case is termed as *Extended type-II seesaw mechanism*. In principle, there could be a gauge singlet mass term in the Lagrangian for extra fermion singlet, i.e, $\mu_S \overline{S^c} S$ which can take any value. But we have taken this mass parameter to be either zero or very small so that the generic inverse seesaw contribution involving μ_S is very much suppressed. In addition, we have assumed the induced VEV for H_L is taken to be zero, i.e, $\langle H_L \rangle \to 0$.

The relevant interaction terms necessary for realizing natural type-II seesaw dominance is given by

$$-\mathcal{L}_{\text{Yuk}} = \overline{\ell_L}\left[Y_3\Phi + Y_4\widetilde{\Phi}\right]\ell_R + f\left[\overline{(\ell_L)^c}\ell_L\Delta_L + \overline{(\ell_R)^c}\ell_R\Delta_R\right] + F\overline{(\ell_R)}H_R S_L^c + \text{h.c.} \qquad (3.6)$$

$$\supset M_D\overline{\nu_L}N_R + M_L\overline{\nu_L^c}\nu_L + M_R\overline{N_R^c}N_R + M\overline{N_R}S_L + \text{h.c.}. \qquad (3.7)$$

With $\langle H_L \rangle \to 0$ and $\mu_S \to 0$, the complete 9×9 neutral fermion mass matrix in the flavor basis of (ν_L, S_L, N_R^c) is read as

$$\mathbb{M} = \begin{pmatrix} & \nu_L & S_L & N_R^c \\ \hline \nu_L & M_L & 0 & M_D \\ S_L & 0 & 0 & M \\ N_R^c & M_D^T & M^T & M_R \end{pmatrix}. \qquad (3.8)$$

Using standard formalism of seesaw mechanism and using mass hierarchy $M_R > M > M_D \gg M_L$, we can integrate out the heaviest right-handed neutrinos as follows

$$\mathbb{M}' = \begin{pmatrix} M_L & 0 \\ 0 & 0 \end{pmatrix} - \begin{pmatrix} M_D \\ M \end{pmatrix} M_R^{-1}\begin{pmatrix} M_D^T & M^T \end{pmatrix}$$

$$= \begin{pmatrix} M_L - M_D M_R^{-1} M_D^T & -M_D M_R^{-1} M^T \\ M M_R^{-1} M_D^T & -M M_R^{-1} M^T \end{pmatrix} \qquad (3.9)$$

where the intermediate block diagonalised neutrino states modified as

$$\begin{aligned} \nu' &= \nu_L - M_D M_R^{-1} N_R^c, \\ S' &= S_L - M_D M_R^{-1} N_R^c, \\ N' &= N_R^c + (M_R^{-1} M_D^T)^* \nu_L + (M_R^{-1} M^T)^* S_L. \end{aligned} \qquad (3.10)$$

Thus, the intermediate block diagonalised neutrino states are related to flavor eigenstates in the following transformation,

$$\begin{pmatrix} \nu' \\ S' \\ N' \end{pmatrix} = \begin{pmatrix} \mathbb{I} & \mathbb{O} & -M_D M_R^{-1} \\ \mathbb{O} & \mathbb{I} & -M M_R^{-1} \\ (M_D M_R^{-1})^\dagger & (M M_R^{-1})^\dagger & \mathbb{I} \end{pmatrix} \begin{pmatrix} \nu_L \\ S_L \\ N_R^c \end{pmatrix} \qquad (3.11)$$

It is found that the $(2,2)$ entries of mass matrix $\mathbb{M}'$ is larger than other entries in the limit $M_R > M > M_D \gg M_L$. As a result of this, we can repeat the same procedure in eq. (3.9) to integrate out S'. Thus, the light neutrino mass formula becomes

$$
\begin{aligned}
m_\nu &= \left[M_L - M_D M_R^{-1} M_D^T\right] - \left(-M_D M_R^{-1} M^T\right)\left(-M M_R^{-1} M^T\right)^{-1}\left(-M M_R^{-1} M_D^T\right) \\
&= \left[M_L - M_D M_R^{-1} M_D^T\right] + M_D M_R^{-1} M_D^T \\
&= M_L = m_\nu^{\mathrm{II}}\,,
\end{aligned}
\tag{3.12}
$$

and the physical block diagonalised states are

$$
\begin{aligned}
\hat{\nu} &= \nu_L - M_D M^{-1} S_L \\
\hat{S} &= S_L - M M_R^{-1} N_R^c + (M_D M^{-1})^\dagger S_L
\end{aligned}
\tag{3.13}
$$

with the corresponding block diagonalised transformation as

$$
\begin{pmatrix} \hat{\nu} \\ \hat{S} \end{pmatrix} = \begin{pmatrix} \mathbb{I} & -M_D M^{-1} \\ (M M^{-1})^\dagger & \mathbb{I} \end{pmatrix} \begin{pmatrix} \nu' \\ S' \end{pmatrix}
\tag{3.14}
$$

With this block diagonalization procedure and after few simple algebra, the flavor eigenstates are related to mass eigenstates in the following transformation,

$$
\begin{pmatrix} \nu_L \\ S_L \\ N_R^c \end{pmatrix} = \begin{pmatrix} \mathbb{I} & M_D M^{-1} & M_D M_R^{-1} \\ (M_D M^{-1})^\dagger & \mathbb{I} & M M_R^{-1} \\ \mathbb{O} & -(M M_R^{-1})^\dagger & \mathbb{I} \end{pmatrix} \begin{pmatrix} \nu' \\ S' \\ N' \end{pmatrix}
\tag{3.15}
$$

Subsequently, the final block diagonalised mass matrices can be diagonalised in order to give physical masses by a 9×9 unitary matrix $V_{9\times9}$. The transformation of the block diagonalised neutrino states in terms of mass eigenstates are given by

$$
\hat{\nu}_\alpha = U_{\nu\alpha i}\nu_i\,, \quad \hat{S}_\alpha = U_{S\alpha i}S_i\,, \quad \hat{N}_\alpha = U_{N\alpha i}N_i\,.
\tag{3.16}
$$

while the block diagonalised mass matrices for light left-handed neutrinos, heavy right-handed neutrinos and extra sterile neutrinos are

$$
\begin{aligned}
m_\nu &= M_L\,, \\
M_N &\equiv M_R = \frac{v_R}{v_L} M_L\,, \\
M_S &= -M M_R^{-1} M^T\,.
\end{aligned}
\tag{3.17}
$$

These block diagonalised mass matrices can be further diagonalised by respective 3×3 unitarity matrices as follows

$$
\begin{aligned}
m_\nu^{\mathrm{diag}} &= U_\nu^\dagger m_\nu U_\nu^* = \mathrm{diag}.\{m_1, m_2, m_3\}\,, \\
M_S^{\mathrm{diag}} &= U_S^\dagger M_S U_S^* = \mathrm{diag}.\{M_{S_1}, M_{S_2}, M_{S_3}\}\,, \\
M_N^{\mathrm{diag}} &= U_N^\dagger M_N U_N^* = \mathrm{diag}.\{M_{N_1}, M_{N_2}, M_{N_3}\}\,.
\end{aligned}
\tag{3.18}
$$

Finally, the complete block diagonalization yields

$$\widehat{\mathbb{M}} = \mathrm{V}_{9\times9}^{\dagger} \mathbb{M} \mathrm{V}_{9\times9}^{*} = (\mathbb{W} \cdot \mathbb{U})^{\dagger} \mathbb{M} (\mathbb{W} \cdot \mathbb{U})$$
$$= \mathrm{diag}.\{m_1, m_2, m_3; M_{S_1}, M_{S_2}, M_{S_3}; M_{N_1}, M_{N_2}, M_{N_3}\} \qquad (3.19)$$

Here the block diagonalised mixing matrix $\mathbb{W}$ and the unitarity matrix $\mathbb{U}$ are given by

$$\mathbb{W} = \begin{pmatrix} \mathbb{I} & M_D M^{-1} & M_D M_R^{-1} \\ (M_D M^{-1})^{\dagger} & \mathbb{I} & M M_R^{-1} \\ \mathbb{O} & -(M M_R^{-1})^{\dagger} & \mathbb{I} \end{pmatrix}, \quad \mathbb{U} = \begin{pmatrix} U_\nu & \mathbb{O} & \mathbb{O} \\ \mathbb{O} & U_S & \mathbb{O} \\ \mathbb{O} & \mathbb{O} & U_N \end{pmatrix}. \qquad (3.20)$$

Thus, the complete 9×9 unitary mixing matrix diagonalizing the neutral leptons is as follows

$$\mathrm{V} = \mathbb{W} \cdot \mathbb{U} = \begin{pmatrix} U_\nu & M_D M^{-1} U_S & M_D M_R^{-1} U_N \\ (M_D M^{-1})^{\dagger} U_\nu & U_S & M M_R^{-1} U_N \\ \mathbb{O} & -(M M_R^{-1})^{\dagger} U_S & U_N \end{pmatrix} \qquad (3.21)$$

3.3 Expressing masses and mixing in terms of U_{PMNS} and light neutrino masses

The light neutrinos are generally diagonalised by standard PMNS mixing matrix U_{PMNS} in the basis where charged leptons are already diagonal i.e, $m_\nu^{\mathrm{diag}} = U_{\mathrm{PMNS}}^{\dagger} m_\nu U_{\mathrm{PMNS}}^{*}$. The Dirac neutrino mass matrix M_D in general is a complex matrix. The structure of M_D in LRSM can be approximately taken to be up-quark type mass matrix whose origin can be motivated from high scale Pati-Salam symmetry or SO(10) GUT. If we consider M to be diagonal and degenerate i.e, $M = m_S \mathrm{diag}\{1,1,1\}$, then the mass formulas for neutral leptons are given by

$$m_\nu = M_L = f v_L = U_{\mathrm{PMNS}} m_\nu^{\mathrm{diag}} U_{\mathrm{PMNS}}^{T},$$

$$M_N \equiv M_R = f v_R = \frac{v_R}{v_L} M_L = \frac{v_R}{v_L} U_{\mathrm{PMNS}} m_\nu^{\mathrm{diag}} U_{\mathrm{PMNS}}^{T},$$

$$M_S = -M M_R^{-1} M^T = -m_S^2 \left[\frac{v_R}{v_L} U_{\mathrm{PMNS}} m_\nu^{\mathrm{diag}} U_{\mathrm{PMNS}}^{T} \right]^{-1}, \qquad (3.22)$$

After some simple algebra, the active LH neutrinos ν_L, active RH neutrinos N_R and heavy sterile neutrinos S_L in the flavor basis are related to their mass basis as

$$\begin{pmatrix} \nu_L \\ S_L \\ N_R^c \end{pmatrix}_\alpha = \begin{pmatrix} \mathrm{V}^{\nu\nu} & \mathrm{V}^{\nu S} & \mathrm{V}^{\nu N} \\ \mathrm{V}^{S\nu} & \mathrm{V}^{SS} & \mathrm{V}^{SN} \\ \mathrm{V}^{N\nu} & \mathrm{V}^{NS} & \mathrm{V}^{NN} \end{pmatrix}_{\alpha i} \begin{pmatrix} \nu_i \\ S_i \\ N_i \end{pmatrix} \qquad (3.23)$$

$$= \begin{pmatrix} U_{\mathrm{PMNS}} & \frac{1}{m_S} M_D U_{\mathrm{PMNS}}^{*} & \frac{v_L}{v_R} M_D U_{\mathrm{PMNS}}^{-1} m_\nu^{\mathrm{diag}.\,-1} \\ \frac{1}{m_S} M_D^{\dagger} U_{\mathrm{PMNS}} & U_{\mathrm{PMNS}}^{*} & \frac{v_L}{v_R} m_S U_{\mathrm{PMNS}}^{-1} m_\nu^{\mathrm{diag}.\,-1} \\ \mathbb{O} & \frac{v_L}{v_R} m_S U_{\mathrm{PMNS}}^{-1} m_\nu^{\mathrm{diag}.\,-1} & U_{\mathrm{PMNS}} \end{pmatrix}_{\alpha i} \begin{pmatrix} \nu_i \\ S_i \\ N_i \end{pmatrix}$$

4 Neutrinoless double beta decay in LRSM

In this section, we shall present a detailed discussion on Feynman amplitudes for neutrinoless double beta decay within TeV scale LRSM where light neutrino mass mechanism is governed by natural type-II seesaw dominance. The basic charge current interaction Lagrangian for leptons as well quarks are given by

$$\mathcal{L}^{\ell}_{\text{CC}} = \sum_{\alpha=e,\mu,\tau} \left[\frac{g_L}{\sqrt{2}} \overline{\ell}_{\alpha L} \gamma_\mu \nu_{\alpha L} W_L^\mu + \frac{g_R}{\sqrt{2}} \overline{\ell}_{\alpha R} \gamma_\mu N_{\alpha R} W_R^\mu \right] + \text{h.c.}$$

$$= \frac{g_L}{\sqrt{2}} \overline{e}_L \gamma_\mu \nu_{eL} W_L^\mu + \frac{g_R}{\sqrt{2}} \overline{e}_R \gamma_\mu N_{eR} W_R^\mu + \text{h.c.} + \cdots \tag{4.1}$$

$$\mathcal{L}^{\text{q}}_{\text{CC}} = \left[\frac{g_L}{\sqrt{2}} \overline{u}_L \gamma_\mu d_L W_L^\mu + \frac{g_R}{\sqrt{2}} \overline{u}_R \gamma_\mu d_R W_R^\mu \right] + \text{h.c.} \tag{4.2}$$

Using eq. (3.23) of section 3, the flavor eigenstates (ν_L) and N_R^c are expressed in terms of admixture of mass eigenstates (ν_i, S_i, N_i) in the following way,

$$\nu_{eL} = \text{V}_{ei}^{\nu\nu} \nu_i + \text{V}_{ei}^{\nu S} S_i + \text{V}_{ei}^{\nu N} N_i,$$
$$N_{eR} = \text{V}_{ei}^{NS} S_i + \text{V}_{ei}^{NN} N_i. \tag{4.3}$$

This modifies the charged current interaction for leptons as

$$\mathcal{L}^{\text{mass}}_{\text{CC}} = \frac{g_L}{\sqrt{2}} \left[\overline{e}_L \gamma_\mu \{ \text{V}_{ei}^{\nu\nu} \nu_i + \text{V}_{ei}^{\nu S} S_i + \text{V}_{ei}^{\nu N} N_i \} W_L^\mu \right] + \text{h.c.}$$

$$+ \frac{g_R}{\sqrt{2}} \left[\overline{e}_R \gamma_\mu \{ \text{V}_{ei}^{NS} S_i + \text{V}_{ei}^{NN} N_i \} W_R^\mu \right] + \text{h.c.} \tag{4.4}$$

In the above charged-current interaction, there is a possibility that both left-handed W_L and right-handed W_R gauge bosons can mix with each other which can eventually contribute to $0\nu\beta\beta$ transition amplitude. In the present framework, the resulting mass matrix for LH (RH) charged gauge bosons (W_L, W_R) is given by

$$\mathbb{M}_W = \frac{1}{4} \left(\begin{array}{c|cc} & W_L^+ & W_R^+ \\ \hline W_L^- & g_L^2 \left(v_1^2 + v_2^2 + 2v_L^2 + u_L^2 \right) & -2 g_L g_R v_1^* v_2 \\ W_R^- & -2 g_L g_R v_1 v_2^* & g_R^2 \left(2v_R^2 + u_R^2 + v_1^2 + v_2^2 \right) \end{array} \right) \tag{4.5}$$

The physical masses of the charged gauge bosons derived with $g_L = g_R$ after diagonalization are given by

$$M_{W_1}^2 \approx \frac{1}{4} g_L^2 \left[\left(v_1^2 + v_2^2 \right) - \frac{4v_1^2 v_2^2}{u_R^2 + 2v_R^2} \right],$$

$$M_{W_2}^2 \approx \frac{1}{4} g_R^2 \left[u_R^2 + 2v_R^2 + v_1^2 + v_2^2 \right]. \tag{4.6}$$

The physical gauge boson states W_1 and W_2 are related to the mixture of weak eigenstates W_L and W_R as

$$\mathbb{R}_W \equiv \begin{pmatrix} \cos\xi & \sin\xi \\ -\sin\xi & \cos\xi \end{pmatrix}, \tag{4.7}$$

where,

$$|\tan 2\xi| \sim \frac{2v_1 v_2}{u_R^2 + 2v_R^2 - u_L^2 - 2v_L^2} \,. \tag{4.8}$$

Thus, one can express physical states in terms of W_L and W_R as follows

$$\begin{cases} W_1 = \quad \cos\xi \; W_L + \sin\xi \; W_R \\ W_2 = -\sin\xi \; W_L + \cos\xi \; W_R \end{cases} \tag{4.9}$$

We classify all contributions to neutrinoless double beta decay in the present TeV scale LRSM as:

- due to standard mechanism mediated by purely left-handed currents ($W_L - W_L$ mediation) via exchange of light neutrinos ν_i,

- due to purely left-handed currents via $W_L^- - W_L^-$ mediation through the exchange of the heavy RH Majorana neutrino N_i and heavy sterile neutrinos S_i,

- due to purely right-handed currents ($W_R - W_R$ mediation) via exchange of heavy right-handed Majorana neutrinos N_i,

- due to purely right-handed currents via $W_R^- - W_R^-$ mediation through the exchange of the light neutrinos ν_i and extra sterile neutrinos S_i,

- due to mixed helicity so called λ and η diagrams through mediation of ν_i, S_i, N_i neutrinos.

Before deducing Feynman amplitudes for various contributions to neutrinoless double beta decay, it is desirable to discuss few points regarding the chiral structure of the matrix element with the neutrino propagator as [21]

$$\begin{aligned} P_L \frac{\slashed{p} + m_i}{p^2 - m_i^2} P_L &= \frac{m_i}{p^2 - m_i^2} \,, & P_R \frac{\slashed{p} + m_i}{p^2 - m_i^2} P_R &= \frac{m_i}{p^2 - m_i^2} \,, \\ P_L \frac{\slashed{p} + m_i}{q^2 - m_i^2} P_R &= \frac{\slashed{p}}{p^2 - m_i^2} \,, & P_R \frac{\slashed{p} + m_i}{p^2 - m_i^2} P_L &= \frac{\slashed{p}}{p^2 - m_i^2} \,, \end{aligned} \tag{4.10}$$

$$\frac{m_i}{p^2 - m_i^2} \simeq \begin{cases} \frac{m_i}{p^2} \,, & m_i^2 \ll p^2 \\ -\frac{1}{m_i} & m_i^2 \gg p^2 \end{cases} \tag{4.11}$$

and

$$\frac{\slashed{p}}{p^2 - m_i^2} \propto \begin{cases} \frac{1}{|p|} \,, & m_i^2 \ll p^2 \\ -\frac{|p|}{m_i^2} & m_i^2 \gg p^2 \,. \end{cases} \tag{4.12}$$

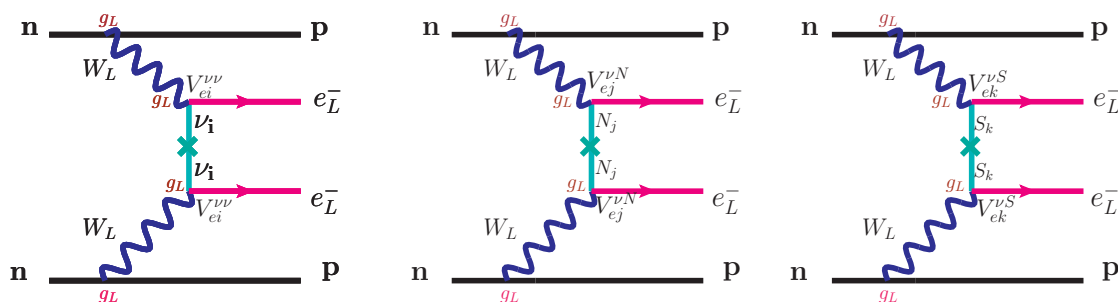**Figure 1**. Feynman diagrams for neutrinoless double beta decay via $W_L^- - W_L^-$ mediation with the exchange of virtual Majorana neutrinos ν_i, N_j and S_k.

4.1 Feynman amplitudes for $0\nu\beta\beta$ decay due to purely left-handed currents

The Feynman amplitudes for $W_L^- - W_L^-$ mediated diagrams shown in figure 1 with the exchange of Majorana neutrinos ν_i, N_j and S_k, respectively, are given by

$$\mathcal{A}_{LL}^{\nu} \propto G_F^2 \sum_{i=1,2,3} \frac{\mathrm{V}_{ei}^{\nu\nu 2} \, m_{\nu_i}}{p^2} \, ,$$

$$\mathcal{A}_{LL}^{N} \propto G_F^2 \sum_{j=1,2,3} \left(-\frac{\mathrm{V}_{ej}^{\nu N 2}}{M_{N_j}} \right) \, ,$$

$$\mathcal{A}_{LL}^{S} \propto G_F^2 \sum_{k=1,2,3} \left(-\frac{\mathrm{V}_{ek}^{\nu S 2}}{M_{S_k}} \right) \, , \tag{4.13}$$

where p is the typical momentum exchange of the $0\nu\beta\beta$ decay process and $G_F = 1.2 \times 10^{-5}\,\mathrm{GeV}^{-2}$ is the Fermi coupling constant. The analytic expressions for suitably normalized dimensionless lepton number violating particle physics parameters for these contributions are as follows

$$\mid \eta_{LL}^{\nu} \mid = \sum_{i=1,2,3} \frac{\mathrm{V}_{ei}^{\nu\nu 2} \, m_{\nu_i}}{m_e} \, , \; \mid \eta_{LL}^{N} \mid = m_p \sum_{j=1,2,3} \frac{\mathrm{V}_{ej}^{\nu N 2}}{M_{N_j}} \, , \; \mid \eta_{LL}^{S} \mid = m_p \sum_{k=1,2,3} \frac{\mathrm{V}_{ek}^{\nu S 2}}{M_{S_k}} \, . \tag{4.14}$$

Though we shall discuss in detail about the lepton number violating effective mass parameters and half-life in the following section, it will be better if one can express normalized effective mass parameters representing LNV due to these above mentioned Feynman diagrams and are given below

$$|\langle m_{ee} \rangle_L^{\nu}| = \sum_{i=1,2,3} \mathrm{V}_{ei}^{\nu\nu 2} \, m_{\nu_i} \, , |\langle m_{ee} \rangle_L^{N}| = \langle p^2 \rangle \sum_{j=1,2,3} \frac{\mathrm{V}_{ej}^{\nu N 2}}{M_{N_j}} \, , |\langle m_{ee} \rangle_L^{S}| = \langle p^2 \rangle \sum_{k=1,2,3} \frac{\mathrm{V}_{ek}^{\nu S 2}}{M_{S_k}} \, .$$
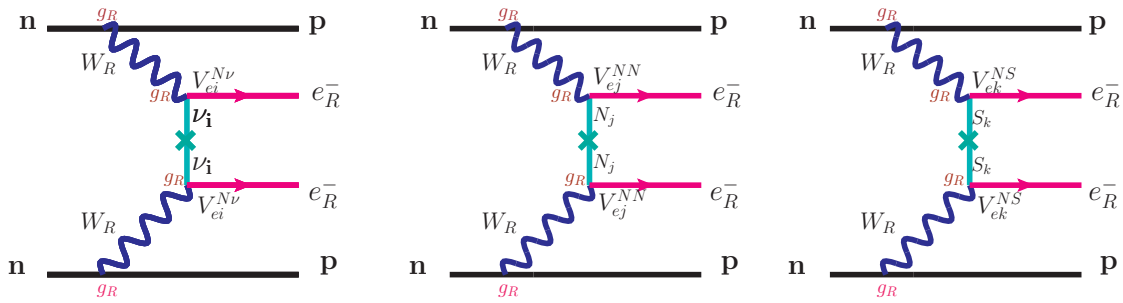
Figure 2. Feynman diagrams for neutrinoless double beta decay ($0\nu\beta\beta$) via $W_R^- - W_R^-$ mediation with the exchange of virtual Majorana neutrinos ν_i, N_j and S_k.

4.2 Feynman amplitudes for $0\nu\beta\beta$ decay due to purely right-handed currents

One of our major contribution in this work is that with $W_R^- - W_R^-$ mediation as shown in first one of figure 2 by the exchange of mainly heavy right-handed neutrinos within the type-II seesaw dominance can yield significantly large contribution to $0\nu\beta\beta$ decay rate than the standard one. The Feynman amplitudes for these diagrams displayed in figure 2 normalized in terms of G_F are given by

$$\mathcal{A}_{RR}^\nu \propto G_F^2 \sum_{i=1,2,3} \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{g_R}{g_L}\right)^4 \frac{\mathrm{V}_{ei}^{N\nu\,2} m_{\nu_i}}{p^2},$$

$$\mathcal{A}_{RR}^N \propto G_F^2 \sum_{j=1,2,3} \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{g_R}{g_L}\right)^4 \left(-\frac{\mathrm{V}_{ej}^{NN\,2}}{M_{N_j}}\right),$$

$$\mathcal{A}_{RR}^S \propto G_F^2 \sum_{k=1,2,3} \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{g_R}{g_L}\right)^4 \left(-\frac{\mathrm{V}_{ek}^{NS\,2}}{M_{S_k}}\right). \tag{4.15}$$

The resulting dimensionless LNV particle physics parameters due to $W_R^- - W_R^-$ mediated diagrams are as follows

$$|\eta_R^\nu| = \sum_{i=1,2,3} \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{g_R}{g_L}\right)^4 \frac{\mathrm{V}_{ei}^{N\nu\,2} m_{\nu_i}}{m_e},$$

$$|\eta_R^N| = \sum_{j=1,2,3} m_p \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{g_R}{g_L}\right)^4 \frac{\mathrm{V}_{ej}^{NN\,2}}{M_{N_j}},$$

$$|\eta_R^S| = \sum_{k=1,2,3} m_p \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{g_R}{g_L}\right)^4 \frac{\mathrm{V}_{ek}^{NS\,2}}{M_{S_k}}. \tag{4.16}$$
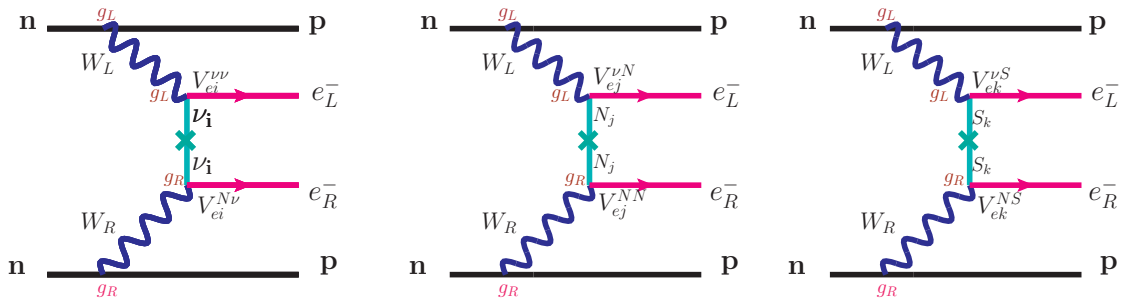
Figure 3. The λ diagram for $0\nu\beta\beta$ decay within LRSM via $W_L^- - W_R^-$ mediation and by the exchange of virtual Majorana neutrinos ν_i, N_j and S_k.
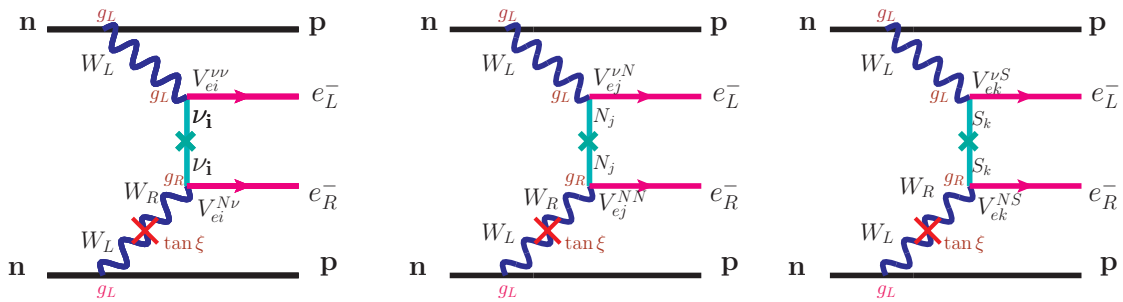

Figure 4. Feynman diagrams for $0\nu\beta\beta$ decay for standard η-contributions which involve mixing between W_L and W_R, i.e, $\tan\xi$ and $W_L^- - W_R^-$ mediation.

4.3 Feynman amplitudes for λ-diagram due to $W_L^- - W_R^-$ mediation

There are Feynman diagrams for neutrinoless double beta decay as shown in figure 3 due to mixed helicity of emitted electrons in the final state via $W_L^- - W_R^-$ mediation and the Feynman amplitudes for these diagrams with the exchange of virtual Majorana neutrinos ν_i, N_j and S_k are

$$\mathcal{A}_\lambda^\nu \propto \mathbf{G_F}^2 \left(\frac{M_{W_L}}{M_{W_R}}\right)^2 \left(\frac{g_R}{g_L}\right)^2 \sum_{i=1,2,3} \mathrm{V}_{ei}^{\nu\nu} \mathrm{V}_{ei}^{N\nu} \frac{1}{|p|},$$

$$\mathcal{A}_\lambda^N \propto \mathbf{G_F}^2 \sum_{j=1,2,3} \left(\frac{M_{W_L}}{M_{W_R}}\right)^2 \left(\frac{g_R}{g_L}\right)^2 \mathrm{V}_{ej}^{\nu N} \mathrm{V}_{ej}^{NN} \frac{|p|}{M_{N_j}^2},$$

$$\mathcal{A}_\lambda^S \propto \mathbf{G_F}^2 \sum_{k=1,2,3} \left(\frac{M_{W_L}}{M_{W_R}}\right)^2 \left(\frac{g_R}{g_L}\right)^2 \mathrm{V}_{ek}^{\nu S} \mathrm{V}_{ek}^{NS} \frac{|p|}{M_{S_k}^2}. \tag{4.17}$$

4.4 Feynman amplitudes for η-diagram with $W_L - W_R$ mixing

There are Feynman diagrams for neutrinoless double beta decay as shown in figure 4 due to mixed helicity of emitted electrons in the final state via $W_L^- - W_R^-$ mediation as well as involves mixing between W_L and W_R gauge boson. The Feynman amplitudes for these
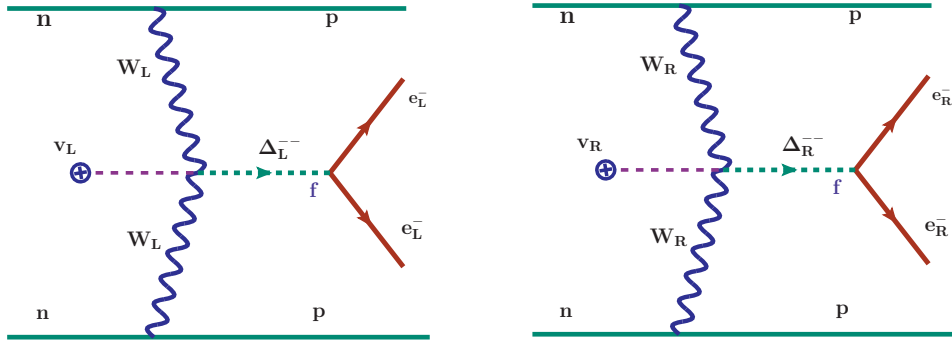
Figure 5. Feynman diagrams for $0\nu\beta\beta$ decay due to doubly charged scalar triplets.

diagrams with the exchange of virtual Majorana neutrinos ν_i, N_j and S_k are given by

$$\mathcal{A}_\eta^\nu \propto \mathbf{G_F}^2 \sum_{i=1,2,3} \left(\frac{g_R}{g_L}\right) \tan\xi \mathrm{V}_{ei}^{\nu\nu} \mathrm{V}_{ei}^{N\nu} \frac{1}{|p|},$$

$$\mathcal{A}_\eta^N \propto \mathbf{G_F}^2 \sum_{j=1,2,3} \left(\frac{g_R}{g_L}\right) \tan\xi \mathrm{V}_{ej}^{\nu N} \mathrm{V}_{ej}^{NN} \frac{|p|}{M_{N_j}^2},$$

$$\mathcal{A}_\eta^S \propto \mathbf{G_F}^2 \sum_{k=1,2,3} \left(\frac{g_R}{g_L}\right) \tan\xi \mathrm{V}_{ek}^{\nu S} \mathrm{V}_{ek}^{NS} \frac{|p|}{M_{S_k}^2}. \tag{4.18}$$

4.5 Feynman amplitudes for $0\nu\beta\beta$ decay due to doubly charged scalar

The Feynman amplitudes due to doubly charged Higgs scalars Δ_L^{--} (Δ_R^{--}) exchanges as displayed in figure 5 are given by

$$\mathcal{A}_{LL}^{\Delta_L} \propto G_F^2 \sum_{i=1,2,3} \frac{1}{M_{\Delta_L}^2} \mathrm{V}_{ei}^{\nu\nu 2} m_{\nu_i},$$

$$\mathcal{A}_{RR}^{\Delta_R} \propto G_F^2 \sum_{j=1,2,3} \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{g_R}{g_L}\right)^4 \frac{1}{M_{\Delta_R}^2} \mathrm{V}_{ej}^{NN 2} M_{N_j},$$

result in lepton number violating dimensionless particle physics parameters.

5 Half-life and normalized LNV effective mass parameters

From the earlier discussion, we found that there are various contributions to neutrino-less double beta decay arising from purely left-handed currents, from purely right-handed currents, from mixed diagrams with left-handed as well as right-handed currents and possible interference effects. In this regard we closely follow the refs. [35, 74, 75] where the QRPA calculations of the matrix elements for the mixed diagrams leads to life-time of $0\nu\beta\beta$ transition as

$$\left[T_{1/2}^{0\nu}\right]^{-1} = G_{01}^{0\nu} |\mathcal{M}_{\mathrm{GT}}^{0\nu}|^2 \left\{ |X_L|^2 + |X_R|^2 + \tilde{C}_2 |\eta_\lambda||X_L|\cos\psi_1 + \tilde{C}_3|\eta_\eta||X_L|\cos\psi_2 \right.$$
$$\left. + \tilde{C}_4|\eta_\lambda|^2 + \tilde{C}_5|\eta_\eta|^2 + \tilde{C}_6|\eta_\lambda||\eta_\eta|\cos(\psi_1-\psi_2) + \mathrm{Re}\left[\tilde{C}_2 X_R \eta_\lambda + \tilde{C}_3 X_R \eta_\eta\right]\right\}, \tag{5.1}$$

where $G_{01}^{0\nu}$ is the phase space factor, $\mathcal{M}_{\mathrm{GT}}^{0\nu}$ is the matrix element for $0\nu\beta\beta$ transition. Here X_L and X_R represent the relevant contributions arising from left-handed and right-handed currents respectively. The coefficients $\tilde{C}_i$ stand for combination of matrix elements and integrated kinematically factors and ψ_i represents complex phases. The LNV dimensionless particle physics parameters are denoted by $\eta's$. In above eq. (5.1), we omitted the interference terms between left-handed and right-handed currents as they are suppressed due to different electron helicities. However the interference terms arising from mixed helicity λ and η diagrams are included.

The neutrino virtual momentum $|p^2| \simeq (100 \text{ MeV})^2$ plays a crucial role as the formula for $0\nu\beta\beta$ transition could be different for $M_i^2 \ll p^2$ or $M_i^2 \gg p^2$ where M_i denoted as mass of any type of neutrinos. It is observed that the light neutrinos having $m_i \ll p^2$ contributing to nuclear matrix element different than the mediating particles having masses $M_i^2 \gg p^2$. To demonstrate this, we assume that the $0\nu\beta\beta$ transition is only mediated by light neutrinos and heavy neutrinos while neglecting the right-handed current effects. In this scenario, the analytic formula for inverse half-life for a given isotope from purely left-handed currents due to exchange of light ν and heavy N neutrinos is given by

$$
\begin{aligned}
\left[T_{1/2}^{0\nu}\right]^{-1} &= G_{01}^{0\nu}\left[|\mathcal{M}_\nu^{0\nu} \cdot \eta_\nu|^2 + |\mathcal{M}_N^{0\nu} \cdot \eta_N|^2\right] \\
&= G_{01}^{0\nu}\left[|\mathcal{M}_\nu^{0\nu}|^2 \frac{\langle m_{ee}^\nu\rangle^2}{m_e^2} + |\mathcal{M}_N^{0\nu}|^2 \left|\left(\frac{m_p}{\langle M_N\rangle}\right)\right|^2\right]
\end{aligned}
\tag{5.2}
$$

Here

$$
\langle m_{ee}^\nu\rangle = \sum_i U_{ei}^2 m_i \quad \text{and} \quad \frac{1}{\langle M_N\rangle} = -\sum_i \frac{V_{ei}^2}{M_i}.
\tag{5.3}
$$

Actually, we normalized here the inverse half-life for standard mechanism due to exchange of light neutrino mechanism as

$$
\left[T_{1/2}^{0\nu}\right]^{-1} = G_{01}^{0\nu}\left|\frac{\mathcal{M}_\nu^{0\nu}}{m_e}\right|^2 |m_{ee}^\nu|^2.
\tag{5.4}
$$

Now we take $G_{01}^{0\nu}\left|\frac{\mathcal{M}_\nu^{0\nu}}{m_e}\right|^2$ as common factor and normalized others with respect to this common factor. Then using eq. (5.2), one can express

$$
\begin{aligned}
\left[T_{1/2}^{0\nu}\right]^{-1} &= G_{01}^{0\nu}\left|\frac{\mathcal{M}_\nu^{0\nu}}{m_e}\right|^2 \left[|\sum_i U_{ei}^2 m_i|^2 + |\left(-m_p m_e \frac{\mathcal{M}_N^{0\nu}}{\mathcal{M}_\nu^{0\nu}}\right)\sum_i \frac{V_{ei}^2}{M_i}|^2\right] \\
&= G_{01}^{0\nu}\left|\frac{\mathcal{M}_\nu^{0\nu}}{m_e}\right|^2 |m_{ee}^{\nu+N}|^2
\end{aligned}
\tag{5.5}
$$

where

$$
\begin{aligned}
\left|m_{ee}^{\nu+N}\right|^2 &\equiv \left|m_{ee}^{\mathrm{eff}}\right|^2 = |m_{ee}^\nu|^2 + \left|m_{ee}^N\right|^2 \\
m_{ee}^N &= \left(-m_p m_e \frac{\mathcal{M}_N^{0\nu}}{\mathcal{M}_\nu^{0\nu}}\right)\sum_i \frac{V_{ei}^2}{M_i} \equiv \langle p\rangle^2 \sum_i \frac{V_{ei}^2}{M_i}.
\end{aligned}
\tag{5.6}
$$

It is clear now that the virtual momentum can be expressed in terms of known masses and nuclear matrix elements.

$$\langle p \rangle^2 = -m_e\, m_p\, \frac{\mathcal{M}_N^{0\nu}}{\mathcal{M}_\nu^{0\nu}} \simeq (100\,\text{MeV})^2\,. \tag{5.7}$$

We discuss here another situation to get a clear idea about how heavy Majorana neutrinos contribute to neutrinoless double beta decay mediated by purely left-handed currents and purely right-handed currents and difference between them. Since we have already discussed heavy neutrino contributions to $0\nu\beta\beta$ transition, one can express inverse half-life formula arising from right-handed currents due to exchange of heavy neutrinos N_R as

$$\frac{1}{T_{1/2}^{0\nu}} = G_{01}^{0\nu}\left[\, |\mathcal{M}_N^{0\nu} \cdot \eta_N^R|^2\, \right]$$
$$= G_{01}^{0\nu}\left[|\mathcal{M}_N^{0\nu}|^2\, |\left(\frac{g_R}{g_L}\right)^4 \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{m_p}{\langle M_N \rangle}\right)|^2\, \right]\,. \tag{5.8}$$

Again following eq. (5.2)–eq. (5.6) and normalized with respect to standard factor $G_{01}^{0\nu}\left|\frac{\mathcal{M}_\nu^{0\nu}}{m_e}\right|^2$, one can express

$$\left[T_{1/2}^{0\nu}\right]^{-1} = G_{01}^{0\nu}\left|\frac{\mathcal{M}_\nu^{0\nu}}{m_e}\right|^2 \left[\, |\left(-m_p m_e \frac{\mathcal{M}_N^{0\nu}}{\mathcal{M}_\nu^{0\nu}}\right)\left(\frac{g_R}{g_L}\right)^4 \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \sum_i \frac{V_{ei}^2}{M_i}|^2\, \right]$$
$$= G_{01}^{0\nu}\left|\frac{\mathcal{M}_\nu^{0\nu}}{m_e}\right|^2 \left|m_{ee,R}^N\right|^2 \tag{5.9}$$

where

$$m_{ee,R}^N = \left(-m_p m_e \frac{\mathcal{M}_N^{0\nu}}{\mathcal{M}_\nu^{0\nu}}\right)\left(\frac{g_R}{g_L}\right)^4 \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \sum_i \frac{V_{ei}^2}{M_i} \equiv \langle p \rangle^2 \left(\frac{g_R}{g_L}\right)^4 \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \sum_i \frac{V_{ei}^2}{M_i}\,.$$

It is seen that the proton mass m_p appears whenever neutrinoless double beta decay is mediated by heavy particles regardless of left-handed or right-handed currents. However, with right-handed current an additional factor of $\left(\frac{g_R}{g_L}\right)^4 \left(\frac{M_{W_L}}{M_{W_R}}\right)^4$ appears. Similarly, one can express half-life for mixed helicity λ and η diagrams and their interference terms in terms of effective Majorana mass parameters.

In order to arrive at a common normalization factor for all types of contributions, at first we use the expression for inverse half-life for $0\nu 2\beta$ decay process due to only light active Majorana neutrinos, $\left[T_{1/2}^{0\nu}\right]^{-1} = G_{01}^{0\nu} |\mathcal{M}_\nu^{0\nu}|^2 |\eta_\nu|^2$. Using the numerical values given in table 7, we rewrite the inverse half-life in terms of effective mass parameter

$$\left[T_{1/2}^{0\nu}\right]^{-1} = G_{01}^{0\nu}\left|\frac{\mathcal{M}_\nu^{0\nu}}{m_e}\right|^2 |\mathbf{m}_\nu^{ee}|^2 = 1.57 \times 10^{-25}\,\text{yrs}^{-1}\,\text{eV}^{-2}|\mathbf{m}_\nu^{ee}|^2 = \mathcal{K}_{0\nu}\,|\mathbf{m}_\nu^{ee}|^2$$

where $\mathbf{m}_\nu^{ee} = \sum_i \left(V_{ei}^{\nu\nu}\right)^2 m_{\nu_i}$ and $\mathcal{K}_{0\nu} \simeq 1.57 \times 10^{-25}\,\text{yrs}^{-1}\,\text{eV}^{-2}$.

Effective Mass Parameters	Analytic formula
$\mathbf{m}_{ee,L}^{\nu}$	$\sum_{i=1}^{3} V_{ei}^{\nu\nu2}\, m_{\nu_i}$
$\mathbf{m}_{ee,L}^{N}$	$\sum_{i=1}^{3} V_{ei}^{\nu N2}\, \frac{|p|^2}{M_{N_i}}$
$\mathbf{m}_{ee,L}^{S}$	$\sum_{i=1}^{3} V_{ei}^{\nu S2}\, \frac{|p|^2}{M_{S_i}}$

Table 2. Effective Majorana mass parameters due to purely left-handed currents.

Effective Mass Parameters	Analytic formula
$\mathbf{m}_{ee,R}^{\nu}$	$\left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{g_R}{g_L}\right)^4 \sum_{i=1}^{3} V_{ei}^{N\nu2}\, m_{\nu_i}$
$\mathbf{m}_{ee,R}^{N}$	$\left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{g_R}{g_L}\right)^4 \sum_{i=1}^{3} V_{ei}^{NN2}\, \frac{|p|^2}{M_{N_i}}$
$\mathbf{m}_{ee,R}^{S}$	$\left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \left(\frac{g_R}{g_L}\right)^4 \sum_{i=1}^{3} V_{ei}^{NS2}\, \frac{|p|^2}{M_{S_i}}$

Table 3. Effective Majorana mass parameters due to purely right-handed currents.

We present here the analytic formula for half-life and normalized effective mass parameters for neutrinoless double beta decay for a given isotope for all relevant contributions are as follows

$$
\begin{aligned}
\left[T_{1/2}^{0\nu}\right]^{-1} &= \mathcal{K}_{0\nu}\left[|\mathbf{m}_{ee}^{\nu}|^2 + |\mathbf{m}_{ee,L}^{S,N}|^2 + |\mathbf{m}_{ee,R}^{S,N}|^2 + |\mathbf{m}_{ee}^{\lambda}|^2 + |\mathbf{m}_{ee}^{\eta}|^2\right] + \cdots \\
&= \mathcal{K}_{0\nu}\left[\left\{|\mathbf{m}_{ee}^{\nu}|^2 + |\mathbf{m}_{ee,L}^{S} + \mathbf{m}_{ee,L}^{N}|^2\right\} + \left\{|\mathbf{m}_{ee,R}^{S} + \mathbf{m}_{ee,R}^{N}|^2\right\} \right. \\
&\quad \left. + \left\{|\mathbf{m}_{ee}^{\lambda,\nu} + \mathbf{m}_{ee}^{\lambda,S} + \mathbf{m}_{ee}^{\lambda,N}|^2\right\} + \left\{|\mathbf{m}_{ee}^{\eta,\nu} + \mathbf{m}_{ee}^{\eta,S} + \mathbf{m}_{ee}^{\eta,N}|^2\right\}\right] + \text{Interference terms}
\end{aligned}
\tag{5.10}
$$

In the above expression for inverse half-life, $G_{01}^{0\nu}$ is the the phase space factor and the other nuclear matrix elements defined for different chiralities of the weak currents such as $\left(\mathcal{M}_{\nu}^{0\nu}\right)$, $\left(\mathcal{M}_{N}^{0\nu}\right)$, $\left(\mathcal{M}_{\lambda}^{0\nu}\right)$ and $\left(\mathcal{M}_{\eta}^{0\nu}\right)$ are presented in table 7. The effective Majorana mass parameters due to purely left handed currents are presented in table 2 while table 3 represents the effective Majorana mass parameters due to purely right handed currents and table 4 shows the contributions due to involvement of both left handed as well as right handed currents. However we do not take into account the interference terms in this work.

6 Numerical results within natural type-II seesaw dominance

6.1 Input model parameters

Before moving towards the numerical estimation of various contributions to neutrinoless double beta decay, it is desirable to know the model parameters and thus we list them below.

Effective Mass Parameters	Analytic formula
$\mathbf{m}^{\nu}_{\text{ee},\lambda}$	$10^{-2}\left(\frac{M_{W_L}}{M_{W_R}}\right)^2\left(\frac{g_R}{g_L}\right)^2\sum_{i=1}^{3}\mathrm{V}^{\nu\nu}_{e\,i}\mathrm{V}^{N\nu}_{e\,i}\,|p|$
$\mathbf{m}^{N}_{\text{ee},\lambda}$	$10^{-2}\left(\frac{M_{W_L}}{M_{W_R}}\right)^2\left(\frac{g_R}{g_L}\right)^2\sum_{j=1}^{3}\mathrm{V}^{\nu N}_{e\,j}\mathrm{V}^{NN}_{e\,j}\,\frac{|p|^3}{M^2_{N_j}}$
$\mathbf{m}^{S}_{\text{ee}\lambda}$	$10^{-2}\left(\frac{M_{W_L}}{M_{W_R}}\right)^2\left(\frac{g_R}{g_L}\right)^2\sum_{k=1}^{3}\mathrm{V}^{\nu S}_{e\,k}\mathrm{V}^{NS}_{e\,k}\,\frac{|p|^3}{M^2_{S_k}}$
$\mathbf{m}^{\nu}_{\text{ee},\eta}$	$\left(\frac{g_R}{g_L}\right)\sum_{i=1}^{3}\mathrm{V}^{\nu\nu}_{e\,i}\mathrm{V}^{N\nu}_{e\,i}\,\tan\zeta_{LR}\,|p|$
$\mathbf{m}^{N}_{\text{ee},\eta}$	$\left(\frac{g_R}{g_L}\right)\sum_{j=1}^{3}\mathrm{V}^{\nu N}_{e\,j}\mathrm{V}^{NN}_{e\,j}\,\tan\zeta_{LR}\,\frac{|p|^3}{M^2_{N_j}}$
$\mathbf{m}^{S}_{\text{ee},\eta}$	$\left(\frac{g_R}{g_L}\right)\sum_{k=1}^{3}\mathrm{V}^{\nu S}_{e\,k}\mathrm{V}^{NS}_{e\,k}\,\tan\zeta_{LR}\,\frac{|p|^3}{M^2_{S_k}}$

Table 4. Effective Majorana mass parameters due to so called λ and η type diagrams. It is to be noted that the suppression factor 10^{-2} arises in the $\lambda-$diagram because of normalization w.r.t. to the standard mechanism.

The method of diagonalization is given in section 3 and the resulting physical masses for all neutral fermions in terms of U_{PMNS} matrix and mass of light neutrinos are give by

$$
\begin{aligned}
m_\nu &= U_{\text{PMNS}}m_\nu^{\text{diag}}U_{\text{PMNS}}^T\,,\\
M_N \equiv M_R &= \frac{v_R}{v_L}U_{\text{PMNS}}m_\nu^{\text{diag}}U_{\text{PMNS}}^T\,,\\
M_S &= -m_S^2\frac{v_L}{v_R}U_{\text{PMNS}}^*m_\nu^{\text{diag}-1}U_{\text{PMNS}}^\dagger\,.
\end{aligned}
\tag{6.1}
$$

The flavor basis of active LH neutrinos ν_L, active RH neutrinos N_R and heavy sterile neutrinos S_L in terms of mass basis and mixing are given as follows

$$
\begin{pmatrix}\nu_L\\S_L\\N_R^c\end{pmatrix}_\alpha = \begin{pmatrix}\mathrm{V}^{\nu\nu}&\mathrm{V}^{\nu S}&\mathrm{V}^{\nu N}\\\mathrm{V}^{S\nu}&\mathrm{V}^{SS}&\mathrm{V}^{SN}\\\mathrm{V}^{N\nu}&\mathrm{V}^{NS}&\mathrm{V}^{NN}\end{pmatrix}_{\alpha i}\begin{pmatrix}\nu\\S\\N\end{pmatrix}_i
\tag{6.2}
$$

in order to express in terms of know neutrino oscillation parameters and light neutrino masses. Where

$$
\mathrm{V}^{\nu\nu}=U_{\text{PMNS}}\,,\qquad \mathrm{V}^{\nu S}=\frac{1}{m_S}M_D U_{\text{PMNS}}^*\,,\qquad \mathrm{V}^{\nu N}=\frac{v_L}{v_R}M_D U_{\text{PMNS}}^{-1}m_\nu^{\text{diag.}-1}\,,
$$

$$
\mathrm{V}^{S\nu}=\frac{1}{m_S}M_D^\dagger U_{\text{PMNS}}\,,\quad \mathrm{V}^{SS}=U_{\text{PMNS}}^*\,,\qquad \mathrm{V}^{SN}=\frac{v_L}{v_R}m_S U_{\text{PMNS}}^{-1}m_\nu^{\text{diag.}-1}\,,
$$

$$
\mathrm{V}^{N\nu}=\mathbb{0}\,,\qquad \mathrm{V}^{NS}=\frac{v_L}{v_R}m_S U_{\text{PMNS}}^{-1}m_\nu^{\text{diag.}-1}\,,\quad \mathrm{V}^{NN}=U_{\text{PMNS}}\,.
\tag{6.3}
$$

Here we consider Dirac neutrino mass matrix motivated from SO(10) GUT and assumed heavy $N-S$ mixing matrix M to be diagonal and degenerate i.e, $M=m_S\,\text{diag}\{1,1,1\}$. We fix m_S at $500\,\text{GeV}$ for all our numerical estimations. If we assume that the present TeV scale left-right symmetric model is originated from Pati-Salam symmetry [30] or SO(10) GUT [76], then the Dirac neutrino mass matrix M_D can be approximated as up-type quark

mass matrix.[2] This can be reconstructed using masses of up, charm & top quarks and the corresponding CKM mixing matrix in the quark sector [78, 79] as

$$
M_D = V_{\text{CKM}} M_u V_{\text{CKM}}^T
$$

$$
= \begin{pmatrix} 0.067 - 0.004\,i & 0.302 - 0.022\,i & 0.550 - 0.530\,i \\ 0.302 - 0.022\,i & 1.480 & 6.534 - 0.001\,i \\ 0.550 - 0.530\,i & 6.534 - 0.0009\,i & 159.72 \end{pmatrix} \text{GeV}.
$$

In the above matrix we use the PDG [78] value of up-type quark mass matrix and the corresponding CKM mixing matrix as

$$
M_u = \text{diag}\{2.3\,\text{MeV}, 1.275\,\text{GeV}, 173.210\,\text{GeV}\},
$$

$$
V_{\text{CKM}} = \begin{pmatrix} 0.97427 & 0.22534 & 0.00351 - i0.0033 \\ -0.2252 + i0.0001 & 0.97344 & 0.0412 \\ 0.00876 - i0.0032 & -0.0404 - i0.0007 & 0.99912 \end{pmatrix}. \quad (6.4)
$$

The bound derived from quark flavor changing neutral current processes is $v_R > 6\,\text{TeV}$ [48, 49, 80, 81] nonetheless we fix it at greater than $8\,\text{TeV}$. The electroweak ρ parameter gives bounds on left-handed scalar triplet VEV as $v_L < 2\,\text{GeV}$ [78] whereas we consider v_L to be $0.1\,\text{eV}$. The light neutrino masses are diagonalised by the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) mixing matrix U_{PMNS} as

$$
m_\nu^{\text{diag}} = U_{\text{PMNS}}^\dagger m_\nu U_{\text{PMNS}}^* = \text{diag.}(m_1, m_2, m_3)
$$

where

$$
U_{\text{PMNS}} = \begin{pmatrix} c_{13}c_{12} & c_{13}s_{12} & s_{13}e^{-i\delta} \\ -c_{23}s_{12} - c_{12}s_{13}s_{23}e^{i\delta} & c_{12}c_{23} - s_{12}s_{13}s_{23}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}s_{13}c_{23}e^{i\delta} & c_{13}c_{23} \end{pmatrix} \cdot P. \quad (6.5)
$$

Here, we have denoted $s_{ij} = \sin\theta_{ij}$, $c_{ij} = \cos\theta_{ij}$ and diagonal phase matrix $P = \text{diag}\left(1, e^{i\alpha}, e^{i\beta}\right)$, where δ is the Dirac CP phase and α, β are Majorana phases varied from $0 \to 2\pi$. From now onwards, we adopt the notations like $(c_\alpha, s_\alpha) \equiv (\cos\theta_\alpha, \sin\theta_\alpha)$ where the atmospheric mixing angle is defined as $\theta_a \equiv \theta_{23}$, the solar mixing angle is defined as $\theta_s \equiv \theta_{12}$, and the reactor mixing angle is defined as $\theta_r \equiv \theta_{13}$. The atmospheric, solar and reactor based neutrino oscillation experiments provide the values of mixing angles θ_{23}, θ_{12} and (θ_{13}) and mass squared differences like $(\Delta m_{\text{atm}}^2)$ and $(\Delta m_{\text{sol}}^2)$. Since the precise measurement of the sign of Δm_{atm}^2 is not confirmed, one can have different possibilities in the arrangement of light neutrino masses like

Normal hierarchy (NH): $\Delta m_{\text{atm}}^2 \equiv \Delta m_{31}^2 > 0$, which gives $m_1 < m_2 < m_3$ with

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

Inverted hierarchy (IH): $\Delta m_{\text{atm}}^2 \equiv \Delta m_{31}^2 < 0$, implying $m_3 < m_1 < m_2$ with

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

---

[2]RG effects modifies the value of Dirac neutrino mass matrix at left-right breaking scale as discussed in refs. [38, 73, 77].

Oscillation Parameters	Within 3σ range (*Schwetz et al.* [82])	within 3σ range (*Fogli et al.* [83])	within 3σ range Gonzalez-Garcia et al. ([84])
$\Delta m_{21}^2[10^{-5}\,\mathrm{eV}^2]$	7.00–8.09	6.99–8.18	7.02–8.09
$\|\Delta m_{31}^2(\mathrm{NH})\|[10^{-3}\,\mathrm{eV}^2]$	2.27–2.69	2.19–2.62	2.317–2.607
$\|\Delta m_{31}^2(\mathrm{IH})\|[10^{-3}\,\mathrm{eV}^2]$	2.24–2.65	2.17–2.61	2.307–2.590
$\sin^2\theta_s$	0.27–0.34	0.259–0.359	0.270–0.344
$\sin^2\theta_a$	0.34–0.67	0.331–0.637	0.382–0.643
$\sin^2\theta_r$	0.016–0.030	0.017–0.031	0.0186–0.0250

Table 5. The oscillation parameters like mass squared differences and mixing angles within 3σ range. However we adopt the values given in ref. [84].

6.2 $0\nu\beta\beta$ contributions from purely left-handed currents

The analytic expression for inverse of half-life for neutrinoless double beta decay due to purely left-handed current via W_L-W_L mediation with the exchange of Majorana neutrinos ν_L, S_L & N_R is given by

$$\left[T_{1/2}^{0\nu}\right]^{-1} = G_{01}^{0\nu}\left[|\mathcal{M}_\nu^{0\nu}|^2|\eta_L^\nu|^2 + |\mathcal{M}_N^{0\nu}|^2|\eta_L^N + \eta_L^S|^2\right]$$

$$= \mathcal{K}_{0\nu}\left[|\mathbf{m}_{\mathrm{ee}}^\nu|^2 + |\mathbf{m}_{\mathrm{ee,L}}^N|^2 + |\mathbf{m}_{\mathrm{ee,L}}^S|^2\right]$$

$$= \mathcal{K}_{0\nu}\left[\left|\sum_{i=1}^3 \mathrm{V}_{ei}^{\nu\nu 2}\, m_i\right|^2 + \left|\sum_{j=1}^3 \mathrm{V}_{ej}^{\nu N 2}\,\frac{|p|^2}{M_{N_j}}\right|^2 + \left|\sum_{k=1}^3 \mathrm{V}_{ek}^{\nu S 2}\,\frac{|p|^2}{M_{S_k}}\right|^2\right] \tag{6.6}$$

6.2.1   For standard mechanism $\mathbf{m_{ee}^\nu}$ and $\mathbf{T_{1/2}^{0\nu}\big|_\nu}$

The LNV effective Majorana mass parameter m_{ee}^ν and corresponding half-life $T_{1/2}^{0\nu}\big|_\nu$ due to standard mechanism by the exchange of light neutrinos is given by

$$|m_{ee}^\nu| = \left||U_{e1}^2|\,m_1 + |U_{e2}^2|\,m_2 e^{i\alpha} + |U_{e3}^2|\,m_3 e^{i\beta}\right| \tag{6.7a}$$

$$= \left|c_s^2 c_r^2 m_1 + s_s^2 c_r^2 m_2 e^{i\alpha} + s_r^2 m_3 e^{i\beta}\right|, \tag{6.7b}$$

$$T_{1/2}^{0\nu}\big|_\nu = \left[G_{01}^{0\nu}\left|\frac{\mathcal{M}_\nu^{0\nu}}{m_e}\right|^2\left|c_s^2 c_r^2 m_1 + s_s^2 c_r^2 m_2 e^{i\alpha} + s_r^2 m_3 e^{i\beta}\right|^2\right]^{-1} \tag{6.7c}$$

Using $m_e = 0.51$ MeV, $G_{01}^{0\nu}$ and $\mathcal{M}_\nu^{0\nu}$ from table 7 and 3σ ranges of oscillation parameters like mixing angles and mass squared differences from table 5, we examine the variation of effective mass and half-life vs. lightest neutrino mass $m_{\mathrm{lightest}} = m_1(\mathrm{NH}), m_3(\mathrm{IH})$. We plot effective Majorana mass parameter m_{ee}^ν in left-panel of figure 6 and half-life $T_{1/2}^{0\nu}\big|_\nu$ in the right-panel of figure 6 as a function of the mass of the lightest neutrino. It is observed that quasi degenerate (QD) pattern of light neutrinos i.e., $m_1 \simeq m_2 \simeq m_3$ and $m_{\mathrm{lightest}} = \sum_i m_i/3$ is disfavoured by current bound on the sum of light neutrino mass
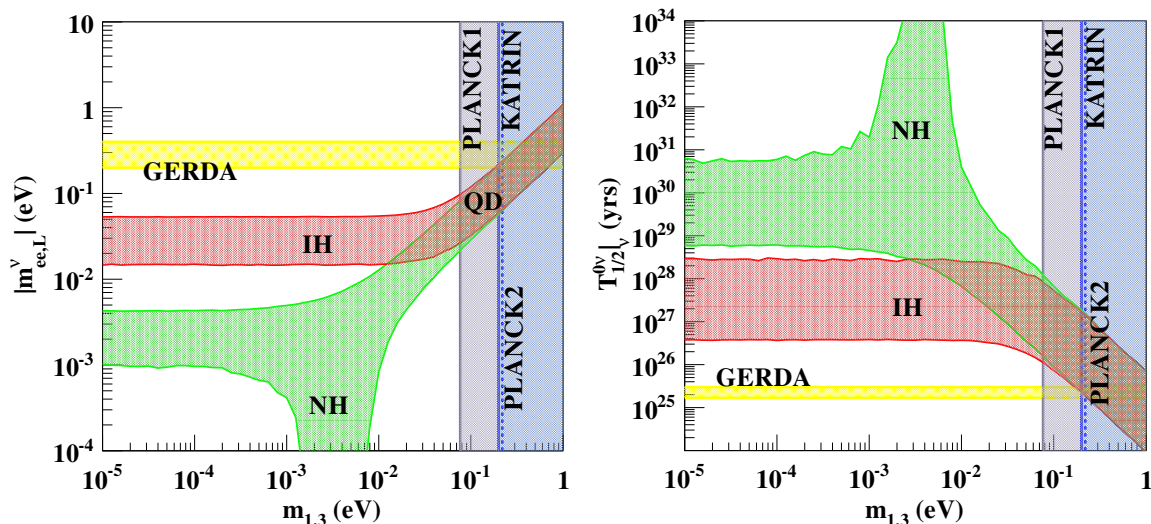
Figure 6. Left panel: the SM contribution to the plot of effective neutrino mass as a function of the lightest neutrino mass, m_1 (m_3) for NH (IH). Right panel: the SM contribution to the half life of $0\nu\beta\beta$ vs lightest neutrino mass, m_1 (m_3) for NH (IH). The NH contributions are displayed by green band while the IH contributions are given by red band. The vertical lines and the corresponding shaded areas are for constraints on the sum of light neutrino masses from recent cosmological data (PLANCK1 and PLANCK2) and KATRIN detector. The yellow band horizontal lines and the respective shaded areas are for the limits in effective Majorana mass parameter and half life by GERDA and EXO+KamLAND-Zen experiments.

Isotope	$T_{1/2}^{0\nu}$ [10^{25} yrs]	m_{ee}^{ν} [eV]	Collaboration
^{76}Ge	> 2.1	< (0.2–0.4)	GERDA [12]
^{136}Xe	> 1.6	< (0.14–0.38)	EXO [14]
^{136}Xe	> 1.9	n/a	KamLAND-Zen [13]
^{136}Xe	> 3.6	< (0.12–0.25)	EXO + KamLAND-Zen combined [13]

Table 6. The table shows the lower limits on the half life $T_{1/2}^{0\nu}$ and upper limits on the effective mass parameter $m_{\text{eff}}^{0\nu}$ for $0\nu\beta\beta$ transition for the isotopes ^{76}Ge and ^{136}Xe from different collaborations. The range for the effective mass parameter comes from the uncertainties in the nuclear matrix elements.

$m_{\Sigma} < 0.23$ derived from Planck+WP+highL+BAO data (PLANCK1) at 95% C.L. while $m_{\Sigma} < 1.08$ derived from Planck+WP+highL (PLANCK2) at 95% C.L. [85]. The experimental bound on half-lives and corresponding effective mass parameter is presented in table 6 for ^{76}Ge and ^{136}Xe isotopes. Moreover, IH as well as NH pattern will be difficult to probe within the standard mechanism even for next generation experiments. This motivates us to consider all possible new physics contributions to neutrinoless double beta decay in the present framework, that might give crucial information about lower limit on absolute scale of lightest neutrino mass as well as mass hierarchy.
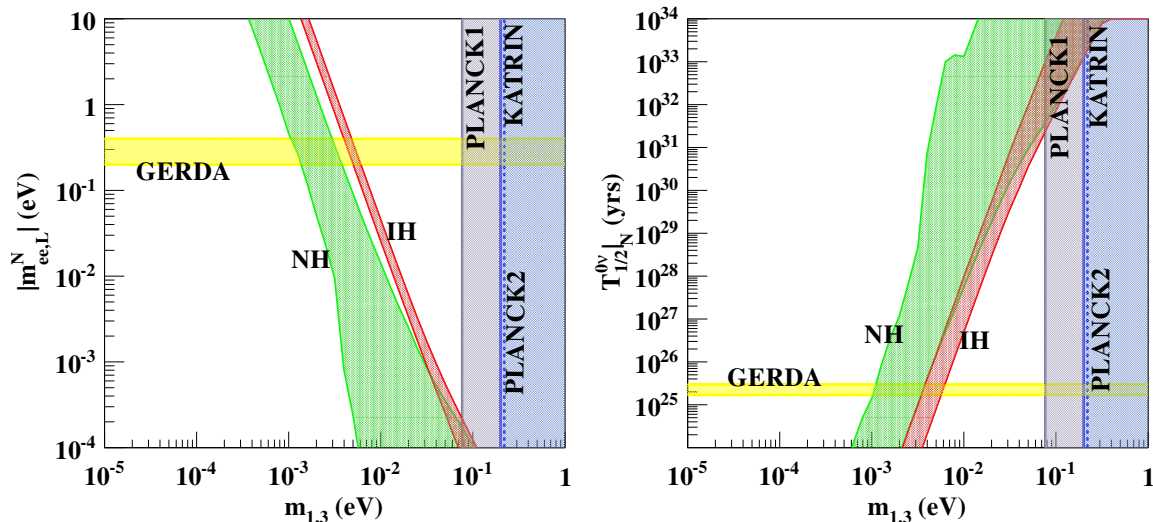
Figure 7. Left panel: the LRSM type-II seesaw dominance contribution to the plot of effective neutrino mass as a function of the lightest neutrino mass, m_1 (m_3) for NH (IH) via $W_L - W_L$ mediation with the exchange of virtual RH neutrino (N). Right panel: the corresponding half life of $0\nu\beta\beta$ vs lightest neutrino mass, m_1 (m_3) for NH (IH).

Isotope	$G_{01}^{0\nu}$ ^{76}Ge [yrs^{-1}]	$\mathcal{M}_\nu^{0\nu}$	$\mathcal{M}_N^{0\nu}$	$\mathcal{M}_\lambda^{0\nu}$	$\mathcal{M}_\eta^{0\nu}$
^{76}Ge	5.77×10^{-15}	2.58–6.64	233–412	1.75–3.76	235–637
^{136}Xe	3.56×10^{-14}	1.57–3.85	164–172	1.96–2.49	370–419

Table 7. Phase space factor $G_{01}^{0\nu}$ and Nuclear Matrix Elements taken from ref. [86].

6.2.2 Non-standard mechanism $m_{ee}^{N,S}$

The expressions for the effective Majorana mass parameter due to exchange of heavy right-handed Majorana neutrinos N_R and extra sterile neutrinos S_L are given by

$$\left| \mathbf{m}_{ee,L}^N \right| = \left| \sum_{k=1}^{3} V_{ek}^{\nu N\,2} \frac{\langle p \rangle^2}{M_{N_k}} \right| = \left| V_{e1}^{\nu N\,2} \frac{\langle p \rangle^2}{M_{N_1}} + V_{e2}^{\nu N\,2} \frac{\langle p \rangle^2}{M_{N_2}} + V_{e3}^{\nu N\,2} \frac{\langle p \rangle^2}{M_{N_3}} \right|,$$

$$\left| \mathbf{m}_{ee,L}^S \right| = \left| \sum_{k=1}^{3} V_{ek}^{\nu S\,2} \frac{\langle p \rangle^2}{M_{S_k}} \right| = \left| V_{e1}^{\nu S\,2} \frac{\langle p \rangle^2}{M_{S_1}} + V_{e2}^{\nu S\,2} \frac{\langle p \rangle^2}{M_{S_2}} + V_{e3}^{\nu S\,2} \frac{\langle p \rangle^2}{M_{S_3}} \right|. \tag{6.8}$$

The variation of LNV effective mass parameters and corresponding half life with the lightest neutrino mass are displayed in left-panel and right-panel of figure 7 due to exchange of N_R. Similarly, we have shown $\mathbf{m}_{ee,L}^S$ and $T_{1/2}^{0\nu}\big|_S$ vs. lightest neutrino mass in figure 8 and the sum of these two new physics contributions is presented in figure 9.

6.3 $0\nu\beta\beta$ from purely right-handed currents

In the present framework the right-handed gauge boson W_R and right-handed Majorana neutrinos N_R lie around few TeV thereby leading to new physics contributions to neutrinoless double beta decay due to purely right-handed currents via $W_R - W_R$ mediation
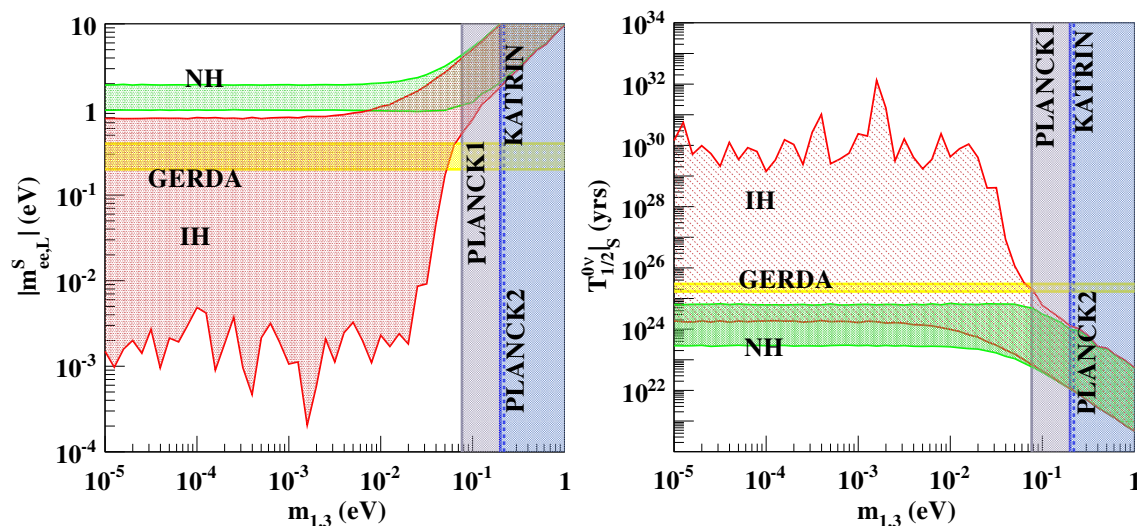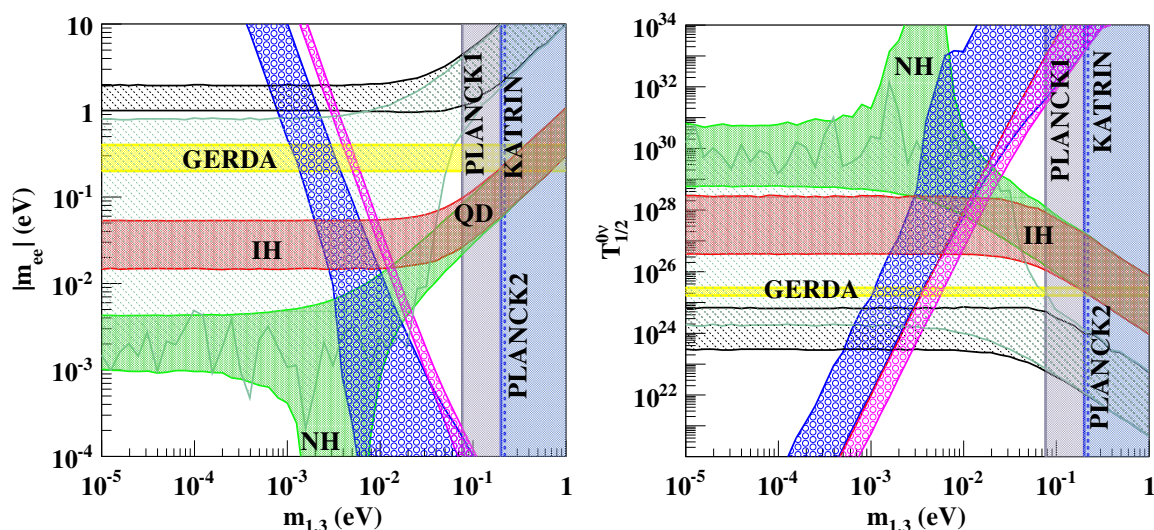
Figure 8. Left panel: the LRSM type-II seesaw dominance contribution to the plot of effective neutrino mass as a function of the lightest neutrino mass, m_1 (m_3) for NH (IH) via $W_L - W_L$ mediation with the exchange of virtual sterile neutrino (S). Right panel: the corresponding half life of $0\nu\beta\beta$ vs lightest neutrino mass, m_1 (m_3) for NH (IH).

Figure 9. Effective Majorana mass (left-panel) and half life (right-panel) as a function of the lightest neutrino mass, m_1 (m_3) for NH (IH) for combined effect of purely left-handed currents mediated by ν, N and S.

and exchange of heavy neutrinos N_R. In addition to this, the type-II seesaw dominance connects light and heavy neutrinos with each other for which one can express new physics contributions in terms of oscillation parameters.

$$m_\nu = M_L \propto M_R\,. \tag{6.9}$$

As a result of this, both light and heavy neutrino masses are diagonalised simultaneously by the U_{PMNS} and the mass eigenvalues are related as follows:

Normal hierarchy (NH):

$$m_2 = \sqrt{m_1^2 + \Delta m_{\text{sol}}^2}, \qquad m_3 = \sqrt{m_1^2 + \Delta m_{\text{atm}}^2},$$
$$M_{N_1} = \frac{m_1}{m_3} M_{N_3}, \qquad\qquad M_{N_2} = \frac{m_2}{m_3} M_{N_3}. \tag{6.10}$$

where we fixed the heaviest RH Majorana neutrino mass M_{N_3} for NH.

Inverted hierarchy (IH):

$$m_1 = \sqrt{m_3^2 + \Delta m_{\text{atm}}^2}, \qquad m_2 = \sqrt{m_3^2 + \Delta m_{\text{atm}}^2 + \Delta m_{\text{sol}}^2}$$
$$M_{N_1} = \frac{m_1}{m_2} M_{N_2}, \qquad\qquad M_{N_3} = \frac{m_3}{m_2} M_{N_2}. \tag{6.11}$$

where we fixed the heaviest RH Majorana neutrino mass M_{N_2} for IH.

The expression for inverse half-life of $0\nu\beta\beta$ transition for a given isotope due to purely right-handed currents along with standard mechanism is given by

$$\frac{1}{T_{1/2}^{0\nu}\big|_{N,R}} = G_{01}^{0\nu} \left| \frac{\mathcal{M}_\nu^{0\nu}}{m_e} \right|^2 \cdot |m_{\text{ee,R}}^N|^2 \tag{6.12a}$$

$$\frac{1}{T_{1/2}^{0\nu}\big|_{LR}} = G_{01}^{0\nu} \left| \frac{\mathcal{M}_\nu^{0\nu}}{m_e} \right|^2 \left[|m_{\text{ee}}^\nu|^2 + |m_{\text{ee,R}}^N|^2 \right] \tag{6.12b}$$

$$= G_{01}^{0\nu} \left| \frac{\mathcal{M}_\nu^{0\nu}}{m_e} \right| |m_{\text{ee}}^{(\nu+N)}|^2, \tag{6.12c}$$

where $|m_{\text{ee}}^{(\nu+N)}|^2 = |m_{\text{ee}}^\nu|^2 + |m_{\text{ee}}^N|^2$. Under this type-II seesaw dominance, the expressions for $|m_{\text{ee}}^\nu|$ and $|m_{\text{ee,R}}^N|$ are given by

$$m_{\text{ee}}^\nu = \left| c_s^2 c_r^2 m_1 + s_s^2 c_r^2 m_2 e^{i\alpha} + s_r^2 m_3 e^{i\beta} \right|, \tag{6.13a}$$

$$m_{\text{ee}}^N \Big|_{\text{NH}} = \frac{C_N}{M_3} \left| c_s^2 c_r^2 \frac{m_3}{m_1} + s_s^2 c_r^2 \frac{m_3}{m_2} e^{i\alpha} + s_r^2 e^{i\beta} \right|, \tag{6.13b}$$

$$m_{\text{ee}}^N \Big|_{\text{IH}} = \frac{C_N}{M_2} \left[c_s^2 c_r^2 \frac{m_2}{m_1} + s_s^2 c_r^2 e^{i\alpha} + \frac{m_2}{m_3} s_r^2 e^{i\beta} \right], \tag{6.13c}$$

where $C_N = \langle p^2 \rangle (g_R/g_L)^4 (M_{W_L}/M_{W_R})^4$. We have neglected the other terms arising from purely right-handed currents. However, the right-handed scalar triplet contribution can be significant if we consider the triplet mass around $500\,\text{GeV}$.

This has been discussed in ref. [54] and the used model parameters are $M_{W_R} \simeq 2\,\text{TeV}$, $g_R \simeq 2/3 g_L$ and $M_N \simeq 1\,\text{TeV}$. In the present work, we have considered $g_L = g_R \simeq 0.65$, $M_{W_R} \simeq 3\,\text{TeV}$ and $M_N \simeq 5\,\text{TeV}$ for numerical estimation of $\mathbf{m}_{\text{ee,L}}^S$ and $T_{1/2}^{0\nu}\big|_S$ vs. lightest neutrino mass in figure 8.

The effect of right-handed currents has been estimated as displayed in figure 10 and figure 11 for effective mass parameter and half-life as a function of lightest neutrino mass.
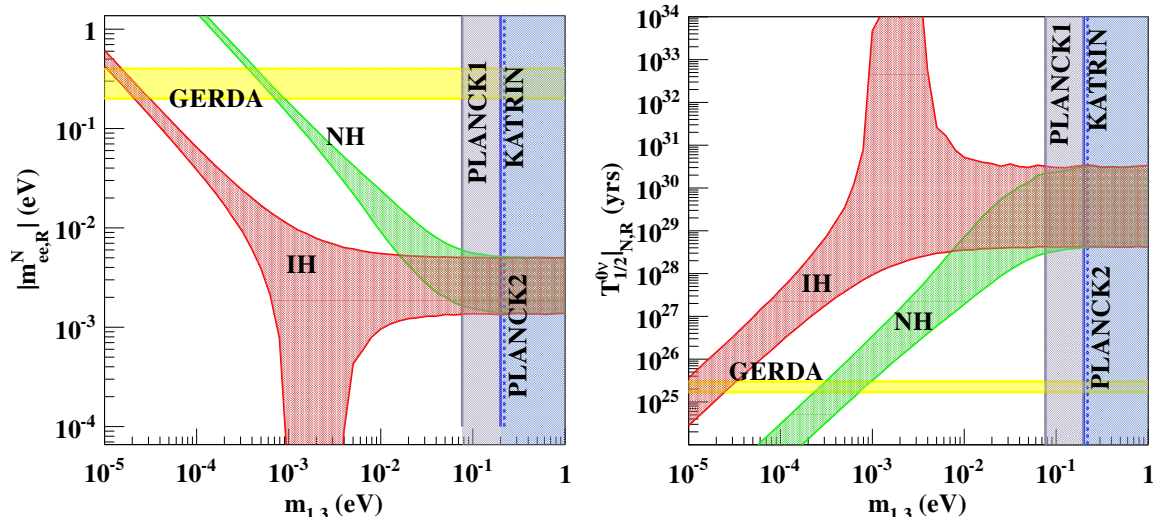
Figure 10. Left panel: the LRSM type-II seesaw dominance contribution to the plot of effective neutrino mass as a function of the lightest neutrino mass, m_1 (m_3) for NH (IH) via $W_R - W_R$ mediation with the exchange of virtual RH neutrino (N). Right panel: the corresponding half life of $0\nu\beta\beta$ vs lightest neutrino mass, m_1 (m_3) for NH (IH).
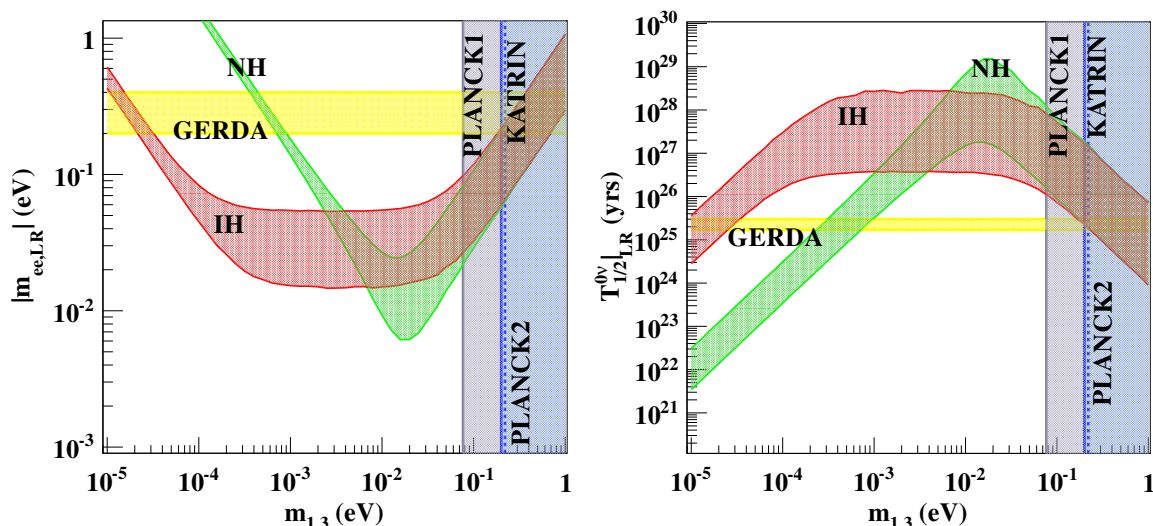

Figure 11. Effective Majorana mass (left-panel) and half life (right-panel) as a function of the lightest neutrino mass, m_1 (m_3) for NH (IH) due to combined effect of standard mechanism and right-handed currents via exchange of heavy right-handed Majorana neutrinos N.

6.4 $0\nu\beta\beta$ from λ and η-diagrams

In this framework with $g_L = g_R$ and $V^{N\nu} = 0$ from the seesaw diagonalization the relevant effective Majorana mass parameters due to so called λ and η diagrams are expressed as follows.

$$\mathbf{m}_{ee,\lambda}^N = 10^{-2} \left(\frac{M_{W_L}}{M_{W_R}}\right)^2 \sum_{j=1}^{3} V_{ej}^{\nu N} V_{ej}^{NN} \frac{|p|^3}{M_{N_j}^2} \tag{6.14}$$
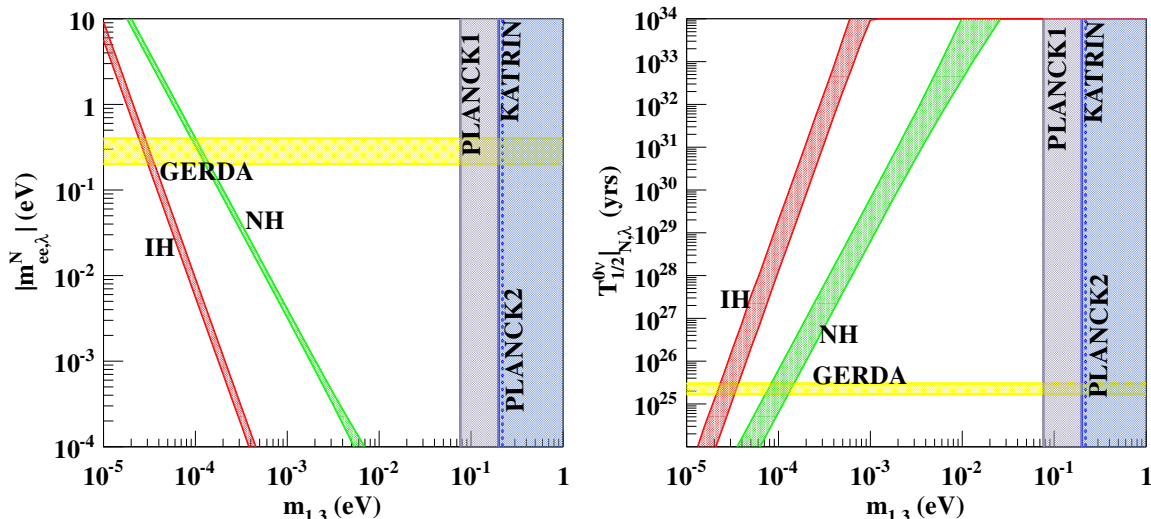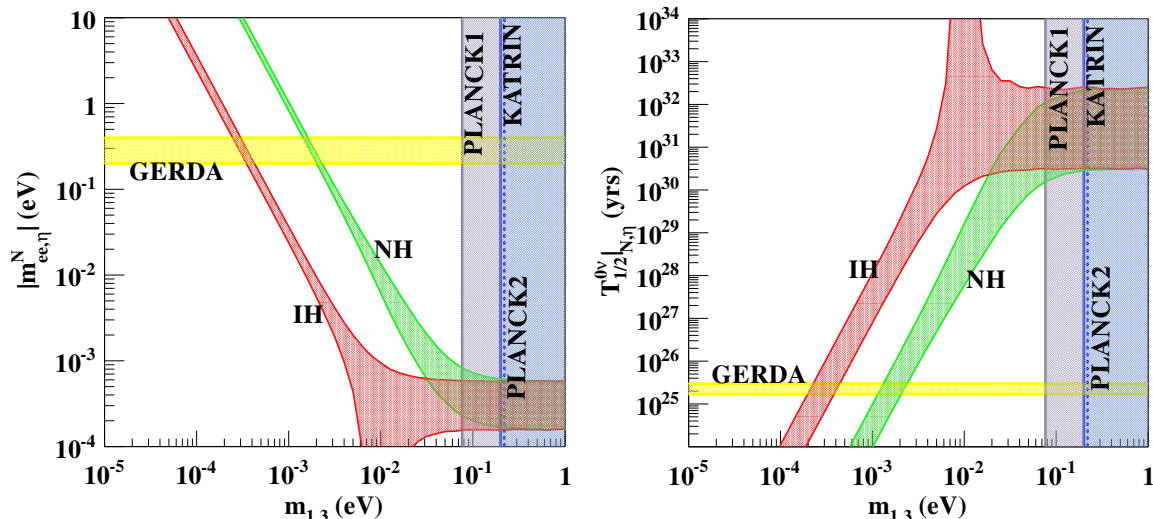
Figure 12. Effective Majorana mass (left-panel) and half life (right-panel) as a function of the lightest neutrino mass, m_1 (m_3) for NH (IH) due to λ diagram via exchange of heavy neutrinos.

$$\mathbf{m}_{ee\lambda}^{S} = 10^{-2} \left(\frac{M_{W_L}}{M_{W_R}} \right)^2 \sum_{k=1}^{3} \mathrm{V}_{ek}^{\nu S} \mathrm{V}_{ek}^{NS} \frac{|p|^3}{M_{S_k}^2} \tag{6.15}$$

$$\mathbf{m}_{ee,\eta}^{N} = \sum_{j=1}^{3} \mathrm{V}_{ej}^{\nu N} \mathrm{V}_{ej}^{NN} \tan\zeta_{LR} \frac{|p|^3}{M_{N_j}^2} \tag{6.16}$$

$$\mathbf{m}_{ee,\eta}^{S} = \sum_{k=1}^{3} \mathrm{V}_{ek}^{\nu S} \mathrm{V}_{ek}^{NS} \tan\zeta_{LR} \frac{|p|^3}{M_{S_k}^2} \tag{6.17}$$

With M_D similar to up-quark mass matrix and other input model parameters, the effective mass and corresponding half-life with the variation of lightest neutrino mass m_1(NH) and m_3(IH) due to so called λ and η diagrams to $0\nu\beta\beta$ transition are displayed in figure 12, figure 13 and figure 14.

6.5 Mass hierarchy discrimination within natural type-II seesaw dominance

Here we discuss the comparison between the standard mechanism and the new physics contributions to $0\nu\beta\beta$ transition within the present framework with natural type-II seesaw dominance by plotting effective Majorana mass as a function of sum of light neutrino masses (m_Σ) using the cosmological limit on the light neutrino mass sum. The light neutrino mass sum is defined as $m_\Sigma = \sum_i m_i = m_1 + m_2 + m_3$. The lower limits of light neutrino mass sum m_Σ derived from cosmology as well as measurements from ongoing neutrino less double beta decay experiments at 1σ, 2σ and 3σ C.L are given by [87–89]

$$\begin{aligned} m_\Sigma &< 84 \ \text{meV} & \text{at } 1\sigma \text{ C.L.}, \\ m_\Sigma &< 146 \ \text{meV} & \text{at } 1\sigma \text{ C.L.}, \\ m_\Sigma &< 208 \ \text{meV} & \text{at } 1\sigma \text{ C.L.} \end{aligned} \tag{6.18}$$
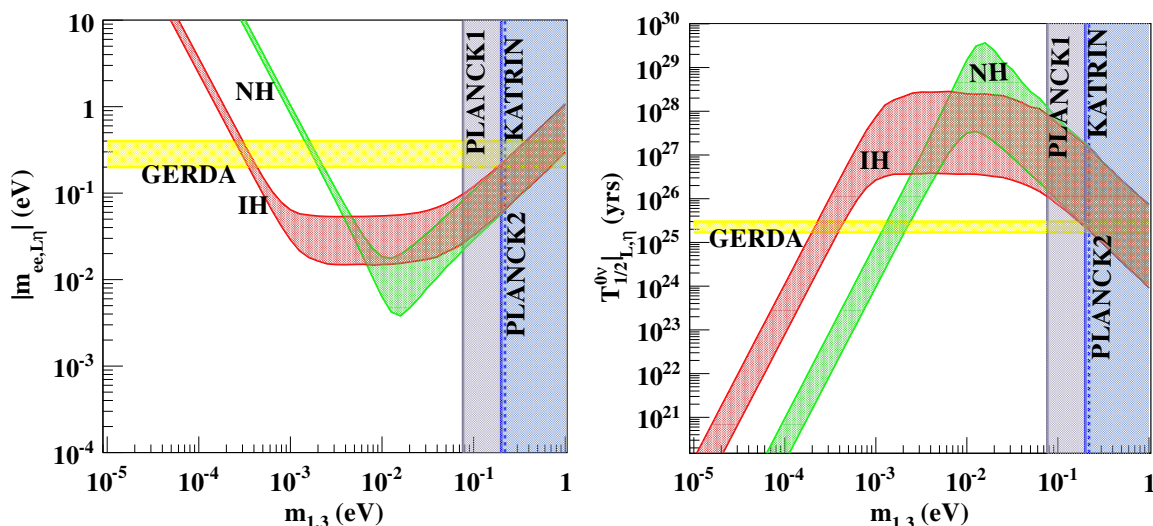
Figure 13. Effective Majorana mass (left-panel) and half life (right-panel) as a function of the lightest neutrino mass, m_1 (m_3) for NH (IH) due to η diagram via exchange of heavy right-handed Majorana neutrino N.

Figure 14. Effective Majorana mass (left-panel) and half life (right-panel) as a function of the lightest neutrino mass, m_1 (m_3) for NH (IH) due to combine effect of standard mechanism and η diagram via exchange of N and S.

We plot effective Majorana mass m_{ee} as a function of sum of light neutrino masses (m_Σ) displayed in figure 15 and figure 16 where standard mechanism is represented by the red band for NH and by the green band for IH while the new physics contributions are represented by the blue band for NH and the red hatched band for IH.

As the contribution of heavy right-handed neutrino and sterile neutrino to effective mass parameter saturate the experimental GERDA limit, figure 15 (left-panel) represents their combined effect to standard mechanism. The spectrum for IH due to the SM and others are lying within the region of cosmological bound and hence disfavoured at 1σ C.L.
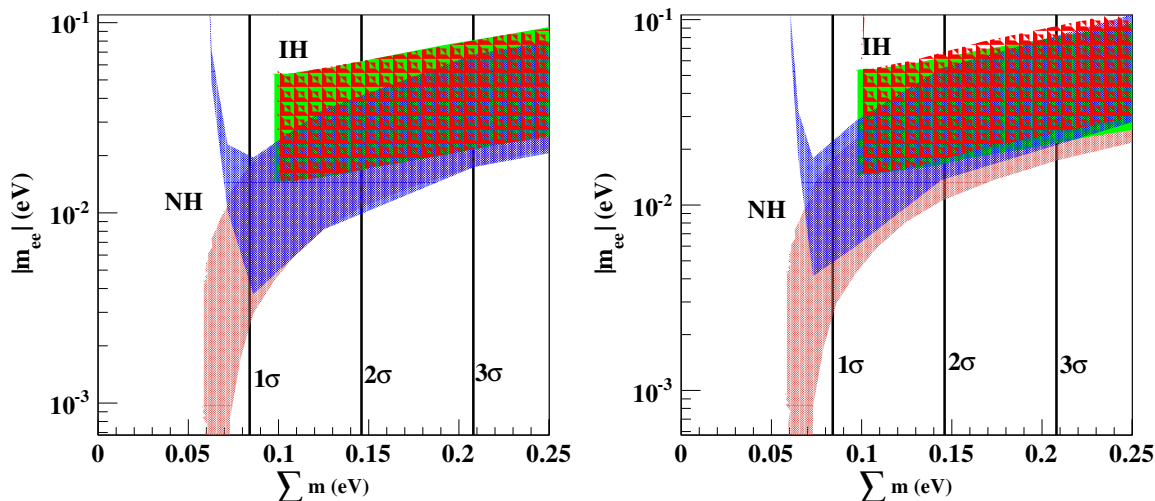
Figure 15. Left panel: preferred region of effective mass parameter $|m_{ee}|$ for standard mechanism and its addition to $W_L - W_L$ mediation with the exchange of heavy N_R and sterile neutrino as a function of sum of light neutrino masses (m_Σ). Right panel: allowed region of $|m_{ee}|$ for standard mechanism and its addition to $W_R - W_R$ mediation with the exchange of heavy N_R as a function of m_Σ.
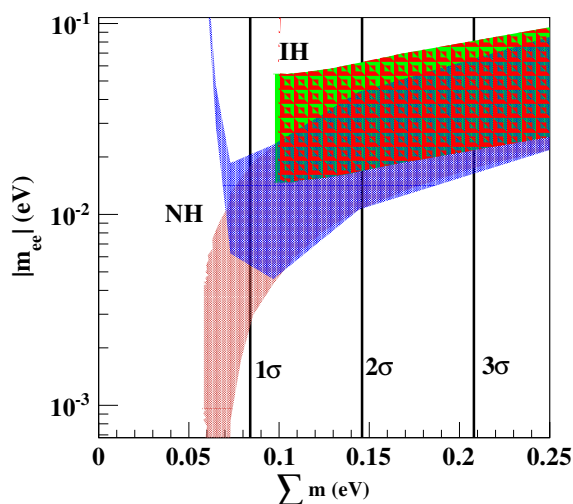

Figure 16. Allowed region of $|m_{ee}|$ for standard mechanism and its combined effect to λ diagram due to exchange of heavy RH neutrino and sterile neutrino as a function of m_Σ.

Whereas the NH spectrum are lying in the privileged region and favored for lower mass of lightest neutrino. But in case of purely right-handed currents i.e mediation via $W_R - W_R$, the contribution on effective Majorana mass parameters due to the exchange of ν_L and heavy sterile neutrinos are negligible. So the right-panel of figure 15 shows the effect of heavy N_R with standard mechanism on sum of light neutrino masses. Here also NH is favored over IH both for standard mechanism and new physics. Similarly, figure 16 indicates the effect of λ diagram due to exchange of heavy N_R and S with standard mechanism on effective mass having same characteristic as in figure 15.

Isotope	NSM (UCOM) [90]	QRPA (CCM) [91]	IBM (Jastrow) [92]
^{76}Ge	2.58	4.07–6.64	4.25–5.07
^{136}Xe	2.00	1.57–3.24	3.07

Table 8. Values of nuclear matrix elements for light neutrino exchange ($\mathcal{M}_\nu^{0\nu}$) for ^{76}Ge and ^{136}Xe.

Isotope	IBM (M-S) [93]	QRPA (CCM) [94]
^{76}Ge	48.1	233–412
^{136}Xe	35.1	164–172

Table 9. Values of nuclear matrix elements for heavy neutrino exchange ($\mathcal{M}_\nu^{0\nu}$) for ^{76}Ge and ^{136}Xe.

Isotope	$\mathcal{M}_\lambda^{0\nu}$		$\mathcal{M}_\eta^{0\nu}$	
	QRPA (CCM) [95]	QRPA (HD) [20]	QRPA (CCM) [95]	QRPA (HD) [20]
^{76}Ge	1.75–3.76	4.47	235–637	791
^{136}Xe	1.96–2.49	2.17	370–419	434

Table 10. Nuclear matrix elements for the λ- and η-diagrams with exchange of light neutrinos. However it should be noted that there are no nuclear matrix elements for lambda and eta diagrams through exchange of heavy neutrinos.

For all of them, the uncertainty on the effective Majorana mass parameter increases with increase in sum of masses in case of normal hierarchy for the contribution of new physics to standard mechanism.

7 Comparison of half-lives for $0\nu\beta\beta$ in ^{76}Ge and ^{136}Xe

We intend to make here a comparative study of half-lives for neutrinoless double beta decay in ^{76}Ge and ^{136}Xe indicating uncertainties in the nuclear matrix elements (one may refer [13] for the matrix element calculations). The half-life limits for different experiments are $T_{1/2}^{0\nu} \simeq 1.07 \times 10^{26}$yrs (for KamLAND-Zen expt. using ^{136}Xe) and $T_{1/2}^{0\nu} \simeq 5.52 \times 10^{25}$yrs (for GERDA Phase-II using ^{76}Ge). Using the values given in tables 8, 9 and 10 for nuclear matrix elements for light and heavy neutrino exchange as well as for λ- and η-diagrams we have shown the correlation plots between half-lives for ^{76}Ge and ^{136}Xe in figure 17 and figure 18 (these were first introduced in ref. [35]). The band in each plots shows the measure of uncertainties in different nuclear matrix elements while measuring half-life in ^{136}Xe to one measured in ^{76}Ge.
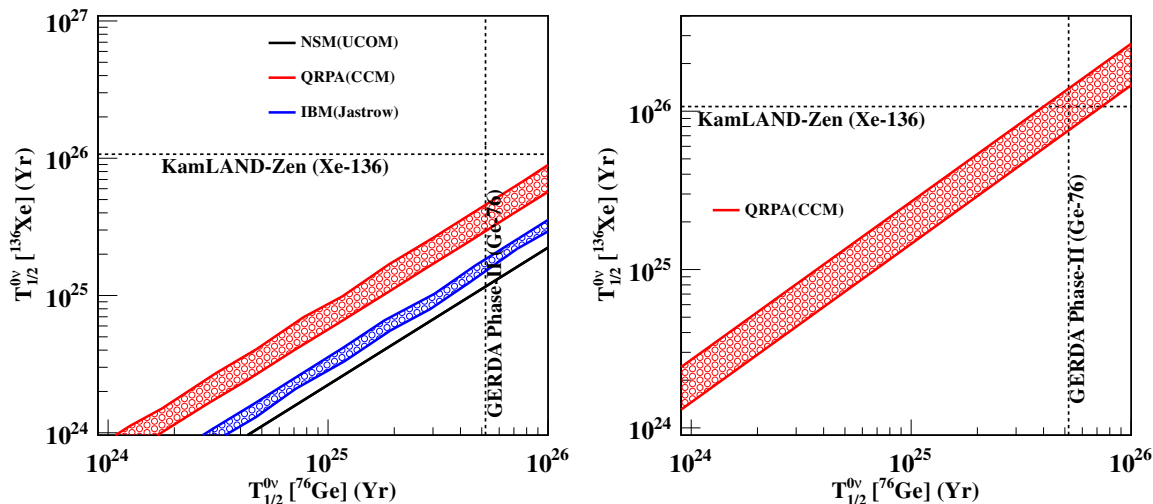
Figure 17. The left-panel shows the correlations between the $0\nu\beta\beta$ half-lives in ^{136}Xe and ^{76}Ge for different matrix element calculations and particle physics contribution due to light neutrino exchange while the right-panel is for heavy neutrino exchange.
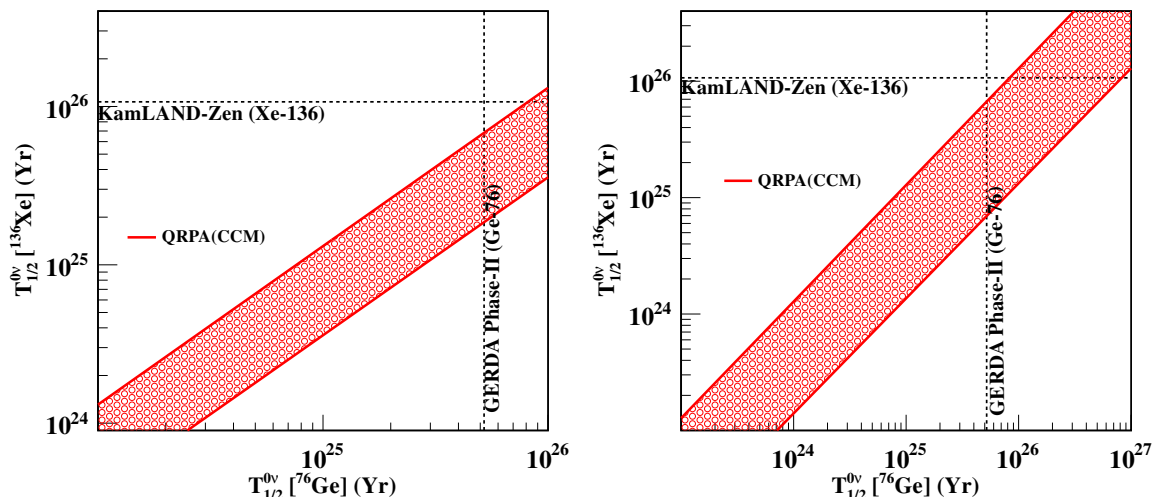

Figure 18. The left-panel shows the correlations between the $0\nu\beta\beta$ half-lives in ^{136}Xe and ^{76}Ge for different matrix element calculations and particle physics contribution due to λ-diagram while the right-panel is for η-diagram.

8 Conclusion

We have discussed natural realization of type-II seesaw dominance within a class of TeV scale left-right symmetric models where scalar sector comprises of scalar doublets $H_{L,R}$, triplets $\Delta_{L,R}$ and a bidoublet Φ, the fermion sector consists of usual quarks $q_{L,R}$, leptons $\ell_{L,R}$ plus one copy of extra sterile fermion S_L per generation. In order to achieve natural type-II seesaw dominance, we have considered negligible VEV for LH scalar doublet i.e, $\langle H_L \rangle \to 0$, negligible mass term for extra sterile neutrinos $\mu_S \to 0$ and mass hierarchy as $M_R > M > M_D \gg M_L$ where $M_L(M_R)$ is the Majorana mass term for LH (RH) neu-

trinos, M_D is the Dirac mass term connecting light-heavy neutrino and M is the $N - S$ mixing matrix. We have also discussed that the type-II seesaw dominance allows any value for M_D and thus, new physics contributions to $0\nu\beta\beta$ transition arise from the following channels; i) due to purely left-handed currents via exchange of heavy RH Majorana neutrinos N and extra sterile neutrinos S, and ii) due to so called λ and η type of diagrams. We have also demonstrated the effect of right-handed currents to $0\nu\beta\beta$ transition via $W_R - W_R$ mediation.

Most importantly we have expressed all the physical masses and mixing like ν_L, N_R and S_L which are completely Majorana in nature mediating LNV processes like neutrinoless double beta decay in terms of oscillation parameters and mass of lightest neutrino with the assumption that $N - S$ mixing matrix is diagonal and degenerate. We have demonstrated that large value of Dirac neutrino mass possibly originating from high scale Pati-Salam symmetry or SO(10) GUT plays an important role in resulting dominant contributions to these new non-standard $0\nu\beta\beta$ transition. With the model parameters like $M_{W_R} = 3\,\mathrm{TeV}$, $M_N \simeq \mathcal{O}(\mathrm{TeV})$, $M_\Delta \simeq \mathcal{O}(\mathrm{TeV})$, M_D as up-type quark mass matrix and using oscillation parameters, we numerically estimated new physics contributions to $0\nu\beta\beta$ transition and compared it with that of the standard mechanism. We have derived the lower limit on absolute scale of lightest neutrino mass by numerically estimating various new physics contributions to $0\nu\beta\beta$ transition by saturating the current experimental limit. We have shown that NH is favored over IH pattern of light neutrinos resulting from effective mass as a function of light neutrino mass sum taking into account the cosmological data. We have also presented correlation plots between half-lives for ^{76}Ge and ^{136}Xe isotopes showing the uncertainties in the nuclear matrix elements.

Acknowledgments

Note added. During the finalization of this work, another work appeared on arXiv [96] that discusses type-II seesaw dominance in LRSM. However, our work differs widely by expressing all the physical masses and mixing of heavy neutrinos in terms of oscillation parameters and lightest neutrino mass. Thus, one can get an important information about absolute scale of lightest neutrino mass and mass hierarchy from new physics contributions to neutrinoless double beta decay by saturating the current experimental limit.

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