

Neutrinoless double beta decay in left-right symmetric models with a universal seesaw mechanism

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We discuss a class of left-right symmetric theories with a universal seesaw mechanism for fermion masses and mixing and the implications for neutrinoless double beta ($0\nu\beta\beta$) decay where neutrino masses are governed by natural type-II seesaw dominance. The scalar sector consists of left- and right-handed Higgs doublets and triplets, while the conventional Higgs bidoublet is absent in this scenario. We use the Higgs doublets to implement the left-right and the electroweak symmetry breaking. On the other hand, the Higgs triplets with induced vacuum expectation values can give Majorana masses to light and heavy neutrinos and mediate $0\nu\beta\beta$ decay. In the absence of the Dirac mass terms for the neutrinos, this framework can naturally realize type-II seesaw dominance even if the right-handed neutrinos have masses of a few TeV. We study the implications of this framework in the context of $0\nu\beta\beta$ decay.

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I. INTRODUCTION

The neutrino oscillation experiments have established that the neutrinos have nonzero masses. However, the question regarding the fundamental nature of the neutrinos; whether they are Dirac [1] or Majorana [2], is yet to be answered. From a theoretical point of view several frameworks predict Majorana neutrinos, while the experimental searches are still inconclusive in this regard. To this end the detection of neutrinoless double beta ($0\nu\beta\beta$) decay, which requires neutrinos to be Majorana particles regardless of the underlying mechanism, plays the crucial role in confirming the nature of the neutrinos. This rare process, the conversion of two neutrons into two protons, two electrons and nothing else, if observed, would pave a path towards the search for new physics beyond the standard model (SM). Currently, the KamLAND-Zen experiment [3] (using

^{136}Xe) claims the best limit on the half-life $T_{1/2}$ to be less than 1.07×10^{26} yr corresponding to an upper bound on effective Majorana mass $m_{\text{eff}} \lesssim 0.06\text{--}0.17$ eV, depending on the nuclear matrix element calculation used.

Models naturally accommodating neutrino masses are the need of the time and the left-right symmetric model (LRSM) [4–9] is one of the most popular candidates for this purpose. The right-handed neutrinos and mirror gauge bosons present make it quintessential to explain the $V - A$ nature of the weak interactions, tiny but nonzero masses of neutrinos [10–14] and a lot more. Moreover, the spontaneous breaking of the left-right symmetry at TeV scale offers a plethora of possibilities in collider phenomenology. The LHC gives a lower bound on the right-handed charged gauge boson mass M_{W_R} of $\mathcal{O}(3 \text{ TeV})$ [15], while the kaon $K_L - K_S$ mass difference results in a lower bound on M_{W_R} of 2.5 TeV [16–19].¹ Such a low scale W_R gauge boson associated with right-handed charged currents can give rise to new contributions to $0\nu\beta\beta$ decay and can be accessible at the LHC. The current limits on left-right symmetry and neutrinoless double beta decay using the LHC dilepton channel data can be found in Refs. [23,24].

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¹Lower values of M_{W_R} in the range 1.9–2.5 TeV are allowed within TeV scale LRSMs with spontaneous D-parity breaking [20–22].

In this work we study a LRSM framework with vectorlike fermions realizing a universal seesaw mechanism for fermion masses, except for the neutrinos. The scalar sector consists of left- and right-handed Higgs doublets and triplets, while the conventional Higgs bidoublet is absent in this scenario. We use the Higgs doublets to implement the left-right and the electroweak symmetry breaking. On the other hand the Higgs triplets with induced vacuum expectation values give Majorana masses to the light and heavy neutrinos in the absence of the Dirac mass terms. Consequently, one can naturally realize a type-II seesaw dominance in this framework even if the right-handed neutrinos have masses around a few TeV. We study its implications for the $0\nu\beta\beta$ decay in this framework.

The outline for the rest of this paper is as follows. In Sec. II we start with a brief overview of the model, followed by discussions on the generation of fermion masses via universal seesaw in Sec. III, neutrino masses via type-II seesaw dominance in Sec. IV and the gauge bosons in Sec. V. In Sec. VI we discuss the implications for $0\nu\beta\beta$ decay. Finally, in Sec. VII we summarize our results and conclude.

II. LEFT-RIGHT SYMMETRY WITH VECTORLIKE FERMIONS

The gauge group for LRSMs is $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_C$. The charge equation relating the electric charge to the third component of isospin in the $SU(2)_{L,R}$ gauge groups and the difference between baryon and lepton number $B-L$ is given by

$$Q = T_{3L} + T_{3R} + \frac{(B-L)}{2}. \quad (1)$$

The usual fermion content of the model is

$$\begin{aligned} Q_L &= \begin{pmatrix} u_L \\ d_L \end{pmatrix} \equiv \left[2, 1, \frac{1}{3}, 3 \right], \\ Q_R &= \begin{pmatrix} u_R \\ d_R \end{pmatrix} \equiv \left[1, 2, \frac{1}{3}, 3 \right], \\ \ell_L &= \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} \equiv [2, 1, -1, 1], \\ \ell_R &= \begin{pmatrix} \nu_R \\ e_R \end{pmatrix} \equiv [1, 2, -1, 1], \end{aligned} \quad (2)$$

where the numbers in brackets correspond to the transformations under $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_C$. We also consider additional vectorlike quarks and charged leptons [25–28],

$$\begin{aligned} U_{L,R} &\equiv [1, 1, 4/3, 3], \\ D_{L,R} &\equiv [1, 1, -2/3, 3], \\ E_{L,R} &\equiv [1, 1, -2, 1]. \end{aligned} \quad (3)$$

TABLE I. Field content of the LRSM with universal seesaw.

Field	$SU(2)_L$	$SU(2)_R$	$B-L$	$SU(3)_C$
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
$U_{L,R}$	1	1	4/3	3
$D_{L,R}$	1	1	-2/3	3
$E_{L,R}$	1	1	-2	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

In Table I, we present the field content of the left-right symmetric model with universal seesaw and their transformations under the LRSM gauge group.

We implement a scalar sector consisting of $SU(2)_{L,R}$ doublets and triplets, however the conventional scalar bidoublet is absent. We use the Higgs doublets to implement the left-right and the electroweak symmetry breaking: $H_R \equiv (h_R^0, h_R^-)^T \equiv [1, 2, -1, 1]$ breaks the left-right symmetry, while $H_L \equiv (h_L^0, h_L^-)^T \equiv [2, 1, -1, 1]$ breaks the electroweak symmetry once they acquire vacuum expectation values (VEVs),

$$\langle H_R \rangle = \begin{pmatrix} \frac{v_R}{\sqrt{2}} \\ 0 \end{pmatrix}, \quad \langle H_L \rangle = \begin{pmatrix} \frac{v_L}{\sqrt{2}} \\ 0 \end{pmatrix}. \quad (4)$$

Note that the present framework requires only doublet Higgs fields for spontaneous symmetry breaking. However, in the absence of a Higgs bidoublet, we use the vectorlike new fermions to generate correct charged fermion masses through a universal seesaw mechanism. For the neutrinos we note that in the absence of a scalar bidoublet there is no Dirac mass term for light neutrinos and without scalar triplets no Majorana masses are generated either. To remedy this fact we introduce additional scalar triplets Δ_L and Δ_R ,

$$\Delta_{L,R} = \begin{pmatrix} \delta_{L,R}^+/\sqrt{2} & \delta_{L,R}^{++} \\ \delta_{L,R}^0 & -\delta_{L,R}^+/\sqrt{2} \end{pmatrix}, \quad (5)$$

which transform as $\Delta_L \equiv [3, 1, 2, 1]$ and $\Delta_R \equiv [1, 3, 2, 1]$, respectively. They generate Majorana masses for the light and heavy neutrinos although they are not essential in spontaneous symmetry breaking here. In the presence of the Higgs triplets, the manifestly left-right symmetric scalar potential has the form

$$\begin{aligned} \mathcal{L} &= (D_\mu H_L)^\dagger D^\mu H_L + (D_\mu H_R)^\dagger D^\mu H_R \\ &+ (D_\mu \Delta_L)^\dagger D^\mu \Delta_L + (D_\mu \Delta_R)^\dagger D^\mu \Delta_R \\ &- \mathcal{V}(H_L, H_R, \Delta_L, \Delta_R), \end{aligned} \quad (6)$$

where the scalar potential is given by

$$\begin{aligned}
\mathcal{V}(H_L, H_R, \Delta_L, \Delta_R) = & -\mu^2(H_L^\dagger H_L + H_R^\dagger H_R) + \lambda[(H_L^\dagger H_L)^2 + (H_R^\dagger H_R)^2] + \beta(H_L^\dagger H_L)(H_R^\dagger H_R) \\
& -\mu'^2[\text{Tr}(\Delta_L^\dagger \Delta_L) + \text{Tr}(\Delta_R^\dagger \Delta_R)] + \lambda_1[(\text{Tr}(\Delta_L^\dagger \Delta_L))^2 + (\text{Tr}(\Delta_R^\dagger \Delta_R))^2] \\
& + \lambda_2[\text{Tr}(\Delta_L^\dagger \Delta_L \Delta_L^\dagger \Delta_L) + \text{Tr}(\Delta_R^\dagger \Delta_R \Delta_R^\dagger \Delta_R)] + \lambda_3[\text{Tr}(\Delta_L^\dagger \Delta_L^\dagger) \text{Tr}(\Delta_L \Delta_L) + \text{Tr}(\Delta_R^\dagger \Delta_R^\dagger) \text{Tr}(\Delta_R \Delta_R)] \\
& + \beta_1 \text{Tr}(\Delta_L^\dagger \Delta_L) \text{Tr}(\Delta_R^\dagger \Delta_R) + \beta_2 \text{Tr}(\Delta_L \Delta_L) \text{Tr}(\Delta_R^\dagger \Delta_R^\dagger) + \rho_1 (\text{Tr}(\Delta_L^\dagger \Delta_L)(H_R^\dagger H_R) + \text{Tr}(\Delta_R^\dagger \Delta_R)(H_L^\dagger H_L)) \\
& + \rho_2 (\text{Tr}(\Delta_L^\dagger \Delta_L)(H_L^\dagger H_L) + \text{Tr}(\Delta_R^\dagger \Delta_R)(H_R^\dagger H_R)) + \rho_3 (H_L^\dagger \Delta_L \Delta_L^\dagger H_L + H_R^\dagger \Delta_R \Delta_R^\dagger H_R) \\
& + \rho_4 (H_L^\dagger \Delta_R \Delta_R^\dagger H_L + H_R^\dagger \Delta_L \Delta_L^\dagger H_R) + \mu(H_L^T i\sigma_2 \Delta_L H_L + H_R^T i\sigma_2 \Delta_R H_R) + \text{H.c.} \quad (7)
\end{aligned}$$

We will assign nonzero VEVs to the Higgs doublets H_R and H_L and triplets Δ_R and Δ_L as follows

$$\begin{aligned}
\langle H_R \rangle &= \begin{pmatrix} \frac{v_R}{\sqrt{2}} \\ 0 \end{pmatrix}, & \langle H_L \rangle &= \begin{pmatrix} \frac{v_L}{\sqrt{2}} \\ 0 \end{pmatrix}, \\
\langle \Delta_R \rangle &= \begin{pmatrix} 0 & 0 \\ \frac{u_R}{\sqrt{2}} & 0 \end{pmatrix}, & \langle \Delta_L \rangle &= \begin{pmatrix} 0 & 0 \\ \frac{u_L}{\sqrt{2}} & 0 \end{pmatrix}. \quad (8)
\end{aligned}$$

Now considering the above VEVs, the part of the scalar potential involving only triplet Higgs takes the form

$$\begin{aligned}
\mathcal{V}(\langle \Delta_L^0 \rangle, \langle \Delta_R^0 \rangle) = & -\frac{1}{2}\mu'^2(u_L^2 + u_R^2) + \frac{1}{4}(\lambda_1 + \lambda_2)(u_L^4 + u_R^4) \\
& + \frac{1}{4}\beta_1 u_L^2 u_R^2. \quad (9)
\end{aligned}$$

Next, let us parametrize $u_L = u \sin \alpha$ and $u_R = u \cos \alpha$ to obtain the extremum condition after differentiating $\mathcal{V}(\langle \Delta_L^0 \rangle, \langle \Delta_R^0 \rangle)$ with respect to α ,

$$[\beta_1 - 2(\lambda_1 + \lambda_2)]u^4 \sin(2\alpha) \cos(2\alpha) = 0. \quad (10)$$

The solution $\alpha = \frac{\pi}{4}$ of Eq. (10) corresponds to the minimum for the choice $\beta_1 < 2(\lambda_1 + \lambda_2)$, as can be readily verified by taking a double derivative of the potential with respect to α , whereas $\alpha = 0$ and $\pi/2$ correspond to the maximum for $\beta_1 < 2(\lambda_1 + \lambda_2)$. The solution $\alpha = \frac{\pi}{4}$ of Eq. (10) implies $u_L = u_R = 0$. Once we consider the full potential now to include the triplet interactions with H_L and H_R such that a nonzero VEV $\langle H_R^0 \rangle = v_R$ breaks the LRSM to the SM at a high scale and $\langle H_L^0 \rangle = v_L$ breaks the SM at the electroweak scale we can choose the induced VEVs for scalar triplets much smaller than VEVs of Higgs doublets for a range of parameters in the scalar potential, i.e. $u_L, u_R \ll v_L, v_R$. Consequently, one is well justified write down the induced VEVs in the Higgs triplets as

$$u_L = \frac{\mu v_L^2}{M_{\delta_L^0}^2}, \quad u_R = \frac{\mu v_R^2}{M_{\delta_R^0}^2}, \quad (11)$$

assuming that the term in the scalar potential containing the masses of the scalar triplets dominate over other competing terms.

III. FERMION MASSES VIA UNIVERSAL SEESAW

As discussed earlier, in this scheme normal Dirac mass terms for the SM fermions are not allowed due to the absence of a bidoublet Higgs scalar. However, in the presence of vectorlike copies of quark and charged lepton gauge isosinglets, the charged fermion mass matrices can assume a seesaw structure. The Yukawa interaction Lagrangian in this model is given by

$$\begin{aligned}
\mathcal{L} = & -Y_U^L H_L \bar{Q}_L U_R + Y_U^R H_R \bar{Q}_R U_L + Y_D^L \tilde{H}_L \bar{Q}_L D_R \\
& + Y_D^R \tilde{H}_R \bar{Q}_R D_L + Y_E^L \tilde{H}_L \bar{\ell}_L E_R + Y_E^R \tilde{H}_R \bar{\ell}_R E_L \\
& + \frac{1}{2} f (\bar{\ell}_L^c i\tau_2 \Delta_L \ell_L + \bar{\ell}_R^c i\tau_2 \Delta_R \ell_R) \\
& - M_U \bar{U}_L U_R - M_D \bar{D}_L D_R - M_E \bar{E}_L E_R + \text{H.c.}, \quad (12)
\end{aligned}$$

where we suppress the flavor and color indices on the fields and couplings. $\tilde{H}_{L,R}$ denotes $\tau_2 H_{L,R}^*$, where τ_2 is the usual second Pauli matrix. Note that there is an ambiguity regarding the breaking of parity, which can either be broken spontaneously with the left-right symmetry at around the TeV scale or at a much higher scale independent of the left-right symmetry breaking. In the latter case, the Yukawa couplings corresponding to the right-type and left-type Yukawa terms can be different because of the renormalization group running below the parity breaking scale, $Y_X^R \neq Y_X^L$. Thus, while writing the Yukawa terms above we distinguish the left- and right-handed couplings explicitly with the subscripts L and R .

After spontaneous symmetry breaking we can write the mass matrices for the charged fermions as [26]

$$\begin{aligned}
M_{uU} &= \begin{pmatrix} 0 & Y_U^L v_L \\ Y_U^R v_R & M_U \end{pmatrix}, & M_{dD} &= \begin{pmatrix} 0 & Y_D^L v_L \\ Y_D^R v_R & M_D \end{pmatrix}, \\
M_{eE} &= \begin{pmatrix} 0 & Y_E^L v_L \\ Y_E^R v_R & M_E \end{pmatrix}. \quad (13)
\end{aligned}$$

The corresponding generation of fermion masses is diagrammatically depicted in Fig. 1. Note that we are interested in a scenario where the VEVs of the Higgs doublets are much larger than the VEVs of the Higgs triplets i.e. $u_L \ll v_L, u_R \ll v_R$.

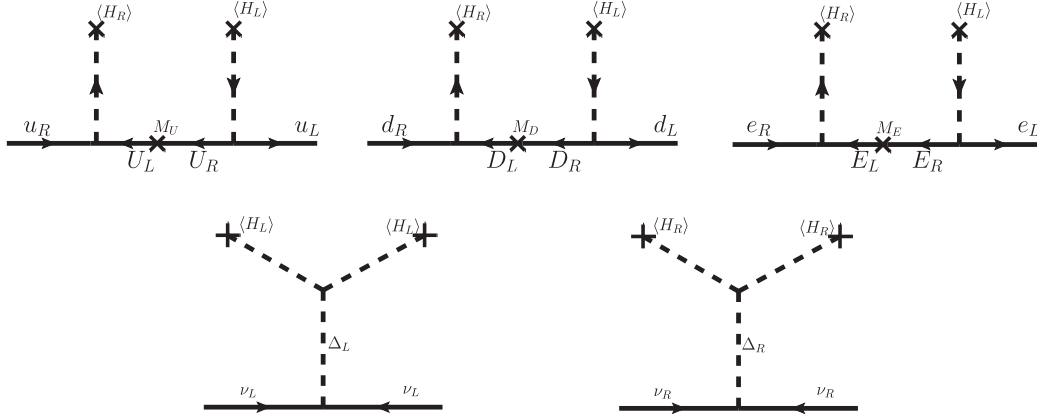


FIG. 1. Generation of fermion masses through universal seesaw and induced triplet VEVs.

Assuming all parameters to be real one can obtain the mass eigenstates by rotating the mass matrices via left and right orthogonal transformations $\mathcal{O}^{L,R}$. For example, up to leading order in $Y_U^L v_L$, the SM and heavy vector partner up-quark masses are

$$m_u \approx Y_U^L Y_U^R \frac{v_L v_R}{\hat{M}_U}, \quad \hat{M}_U \approx \sqrt{M_U^2 + (Y_U^R v_R)^2}, \quad (14)$$

and the mixing angles $\theta_U^{L,R}$ in $\mathcal{O}^{L,R}$ are determined as

$$\tan(2\theta_U^{L,R}) \approx 2Y_U^{L,R} \frac{v_{L,R} M_U}{M_U^2 \pm (Y_U^R v_R)^2}. \quad (15)$$

The other fermion masses and mixing are obtained in an analogous manner. Note that here we have neglected the flavor structure of the Yukawa couplings $Y_X^{L,R}$ which will determine the observed quark and charged lepton mixings. The hierarchy of SM fermion masses can be explained by assuming either a hierarchical structure of the Yukawa couplings or a hierarchical structure of the vectorlike fermion masses.

IV. NEUTRINO MASSES AND TYPE II SEESAW DOMINANCE

In the model under consideration there is no tree level Dirac mass term for the neutrinos due to the absence of a Higgs bidoublet. The scalar triplets acquire induced VEVs $\langle \Delta_L \rangle = u_L$ and $\langle \Delta_R \rangle = u_R$ giving the neutral lepton mass matrix in the basis (ν_L, ν_R) given by

$$M_\nu = \begin{pmatrix} f u_L & 0 \\ 0 & f u_R \end{pmatrix}. \quad (16)$$

Thus the light and heavy neutrino masses are simply $m_\nu = f u_L \propto M_N = f u_R$. A Dirac mass term is generated at the two-loop level via the one-loop W boson mixing θ_W (see the next section) and the exchange of a charged lepton.

It is of the order $m_D \lesssim g_L^4 / (16\pi^2)^2 m_\tau m_b m_t / M_{W_R}^2 \approx 0.1$ eV for $M_{W_R} \approx 5$ TeV. This is intriguingly of the order of the observed neutrino masses; as long as the right-handed neutrinos are much heavier than the left-handed neutrinos, the type-II seesaw dominance is preserved and the induced mixing m_D/M_N is negligible. The mixing between charged gauge bosons $\theta_W \approx g_L^2 / (16\pi^2) m_b m_t / M_{W_R}^2$ is generated through the exchange of bottom and top quarks, and their vector-like partners. This yields a very small mixing of the order $\theta_W \approx 10^{-7}$ for TeV scale W_R bosons.

Incorporating three fermion generations leads to the mixing matrices for the left- and right-handed matrices which we take to be equal

$$V_N = V_\nu \equiv U, \quad (17)$$

where U is the phenomenological Pontecorvo-Maki-Nakagawa-Sakata mixing matrix. Thus the unmeasured mixing matrix for the right-handed neutrinos is fully determined by the left-handed counterpart. The present framework gives a natural realization of type-II seesaw providing a direct relation between light and heavy neutrinos, $M_i \propto m_i$, i.e. the heavy neutrino masses M_i can be expressed in terms of the light neutrino masses m_i as $M_i = m_i(M_3/m_3)$, for a normal and $M_i = m_i(M_2/m_2)$ for an inverse hierarchy of light and heavy neutrino masses.

V. GAUGE BOSONS

As discussed in the previous section, we consider a scenario where the VEVs of the Higgs doublets are much larger than the VEVs of the Higgs triplets i.e. $u_L \ll v_L$, $u_R \ll v_R$.² Thus, the masses for the gauge bosons get small corrections from scalar triplets. The mass matrix for charged gauge bosons is given by

²An interesting situation arises if one assumes $u_L, u_R \ll v_L \ll v_R$. Then one can allow right-handed Majorana neutrinos with masses below GeV which can play an important role in $0\nu\beta\beta$ decay.

$$M_{W'}^2 = \frac{1}{4} \begin{pmatrix} g_L^2(v_L^2 + 2u_L^2) & 0 \\ 0 & g_R^2(v_R^2 + 2u_R^2) \end{pmatrix}, \quad (18)$$

with the gauge couplings g_L and g_R associated with $SU(2)_L$ and $SU(2)_R$, respectively. The tree-level mixing between charged gauge bosons W_L and W_R is zero due to absence of a scalar bidoublet. The physical masses for charged gauge bosons are thus easily found,

$$M_{W_1}^2 \approx \frac{g_L^2}{4} v_L^2, \quad M_{W_2}^2 \approx \frac{g_R^2}{4} (v_R^2 + 2u_R^2), \quad (19)$$

where we neglect u_L . At the one-loop level, a mixing of the order $\theta_W \approx g_L^2/(16\pi^2)m_b m_t/M_{W_R}^2$ is generated through the exchange of bottom and top quarks, and their vectorlike partners. This yields a very small mixing of the order $\theta_W \approx 10^{-7}$ for $M_{W_R} \approx 5$ TeV.

On the other hand, the neutral gauge boson mass matrix is given by

$$M_Z^2 = \frac{1}{4} \begin{pmatrix} g_L^2 v_L^2 & 0 & -g_L g_{BL} \mu_L^2 \\ 0 & g_R^2 \mu_R^2 & -g_{BL} g_R \mu_R^2 \\ -g_{BL} g_L \mu_L^2 & -g_{BL} g_R \mu_R^2 & g_R^2 \mu_L^2 + g_{BL}^2 \mu_R^2 \end{pmatrix}, \quad (20)$$

with $\mu_{L,R}^2 = v_{L,R}^2 + 4u_{L,R}^2$ and the gauge coupling g_{BL} associated with $U(1)_{B-L}$. The diagonalization gives mass eigenvalues for the neutral gauge bosons,

$$M_{Z_1}^2 \approx \frac{g_L^2}{4c_W^2} v_L^2, \quad M_{Z_2}^2 \approx \frac{g_{BL}^2 + g_R^2}{4} (v_R^2 + 4u_R^2). \quad (21)$$

For $u_L \ll v_L$, $u_R \ll v_R$ the mixing between neutral gauge bosons Z_L and Z_R is given by $\theta_Z \approx v_L^2/v_R^2 \approx \mathcal{O}(10^{-4})$. Thus, the mixing angle for neutral sector is much larger than the mixing angle between charged gauge boson. In fact, the absence of a sizeable mixing in the charged gauge boson and neutrino sector is an important consequence of our model and searches for Z_R could constitute an important way of probing this scenario. Nevertheless, a mixing of order $\theta_Z \approx 10^{-4}$ is allowed by collider searches, electroweak precision data and searches for low-energy flavor changing neutral current processes [29–34], and we do not discuss potential avenues to probe the model in this direction but rather focus on neutrinoless double beta decay. We only comment that the mixing between Z_L and Z_R is a standard parameter being probed in collider experiments, for example at the LHC (see e.g. [35]) or especially at planned future electron-positron colliders. It could also be probed in upcoming searches for low-energy flavor changing neutral current processes.

VI. NEUTRINOLESS DOUBLE BETA DECAY

As discussed earlier, there is no tree level Dirac neutrino mass term connecting light and heavy neutrinos. Consequently, the mixing between light and heavy neutrinos is vanishing at this order. Also, the mixing between the charged gauge bosons is vanishing at the tree level due to the absence of a scalar bidoublet.

The charged current interaction in the mass basis for the leptons is given by

$$\frac{g_L}{\sqrt{2}} \sum_{i=1}^3 U_{ei} \left(\bar{\ell}_L \gamma_\mu \nu_i W_L^\mu + \frac{g_R}{g_L} \bar{\ell}_R \gamma_\mu N_i W_R^\mu \right) + \text{H.c.} \quad (22)$$

The charged current interaction for leptons leads to $0\nu\beta\beta$ decay via the exchange of light and heavy neutrinos. There are additional contributions to $0\nu\beta\beta$ decay due to doubly charged triplet scalar exchange. While the left-handed triplet exchange is suppressed because of its small induced VEV, the right-handed triplet can contribute sizeably to $0\nu\beta\beta$ decay.

Before numerical estimation, let us point out the mass relations between light and heavy neutrinos under natural type-II seesaw dominance. For a hierarchical pattern of light neutrinos the mass eigenvalues are given as $m_1 < m_2 \ll m_3$. The lightest neutrino mass eigenvalue is m_1 while the other mass eigenvalues are determined using the oscillation parameters as follows, $m_2^2 = m_1^2 + \Delta m_{\text{sol}}^2$, $m_3^2 = m_1^2 + \Delta m_{\text{atm}}^2 + \Delta m_{\text{sol}}^2$. On the other hand, for the inverted hierarchical pattern of the light neutrino masses $m_3 \ll m_1 \approx m_2$ where m_3 is the lightest mass eigenvalue while other mass eigenvalues are determined by $m_1^2 = m_3^2 + \Delta m_{\text{atm}}^2$, $m_2^2 = m_3^2 + \Delta m_{\text{sol}}^2 + \Delta m_{\text{atm}}^2$. The quasi-degenerate pattern of light neutrinos is $m_1 \approx m_2 \approx m_3 \gg \sqrt{\Delta m_{\text{atm}}^2}$. In any case, the heavy neutrino masses are directly proportional to the light neutrino masses.

In the present analysis, we discuss $0\nu\beta\beta$ decay due to exchange of light neutrinos via left-handed currents, right-handed neutrinos via right-handed currents as well as a right-handed doubly charged scalar³ The half-life for a given isotope for these contributions is given by

$$[T_{1/2}^{0\nu}]^{-1} = G_{01} (|\mathcal{M}_\nu \eta_\nu|^2 + |\mathcal{M}'_N \eta_N + \mathcal{M}_N \eta_\Delta|^2), \quad (23)$$

where G_{01} corresponds to the standard $0\nu\beta\beta$ phase space factor, the \mathcal{M}_i correspond to the nuclear matrix elements for the different exchange processes and η_i are dimensionless parameters determined below.

³A detailed discussion of $0\nu\beta\beta$ decay within LRSMs can be found e.g. in Refs. [9,28,36–44] and for an early study of the effects of light and heavy Majorana neutrinos in neutrinoless double beta decay see Ref. [45].

A. Light neutrinos

The lepton number violating dimensionless particle physics parameter derived from $0\nu\beta\beta$ decay due to the standard mechanism via the exchange of light neutrinos is

$$\eta_\nu = \frac{1}{m_e} \sum_{i=1}^3 U_{ei}^2 m_i = \frac{m_{ee}^\nu}{m_e}. \quad (24)$$

Here, m_e is the electron mass and the effective $0\nu\beta\beta$ mass is explicitly given by

$$m_{ee}^\nu = |c_{12}^2 c_{13}^2 m_1 + s_{12}^2 c_{13}^2 m_2 e^{i\alpha} + s_{13}^2 m_3 e^{i\beta}|, \quad (25)$$

with the sine and cosine of the oscillation angles θ_{12} and θ_{13} , $c_{12} = \cos \theta_{12}$, etc. and the unconstrained Majorana phases $0 \leq \alpha, \beta < 2\pi$.

B. Right-handed neutrinos

The contribution to $0\nu\beta\beta$ decay arising from the purely right-handed currents via the exchange of right-handed neutrinos generally results in the lepton number violating dimensionless particle physics parameter

$$\eta_N = m_p \left(\frac{g_R}{g_L}\right)^4 \left(\frac{M_{W_L}}{M_{W_R}}\right)^4 \sum_{i=1}^3 \frac{U_{ei}^2 M_i}{|p|^2 + M_i^2}. \quad (26)$$

The virtual neutrino momentum $|p|$ is of the order of the nuclear Fermi scale, $p \approx 100$ MeV. m_p is the proton mass and for the manifest LRSM case we have $g_L = g_R$, or else the new contributions are rescaled by the ratio between these two couplings. We in general consider right-handed neutrinos that can be either heavy or light compared to nuclear Fermi scale.

If the mass of the exchanged neutrino is much higher than its momentum, $M_i \gg |p|$, the propagator simplifies as

$$\frac{M_i}{p^2 - M_i^2} \approx -\frac{1}{M_i}, \quad (27)$$

and the effective parameter for right-handed neutrino exchange yields

$$\eta_N = m_p \left(\frac{g_R}{g_L}\right)^4 \left(\frac{M_W}{M_{W_R}}\right)^4 \sum_{i=1}^3 \frac{U_{ei}^2}{M_i} \propto \eta_\nu(m_i^{-1}), \quad (28)$$

where in the expression for $\eta_\nu(m_i^{-1})$ the individual neutrino masses are replaced by their inverse values. Such a contribution clearly becomes suppressed the smaller the right-handed neutrino masses are.

On the other hand, if the mass of the neutrino is much less than its typical momentum, $M_i \ll |p|$, the propagator simplifies in the same way as for the light neutrino exchange,

$$P_R \frac{\not{p} + M_i}{p^2 - M_i^2} P_R \approx \frac{M_i}{p^2}, \quad (29)$$

because both currents are right-handed. As a result, the $0\nu\beta\beta$ decay contribution leads to the dimensionless parameter

$$\eta_N = \frac{m_p}{|p|^2} \left(\frac{g_R}{g_L}\right)^4 \left(\frac{M_W}{M_{W_R}}\right)^4 \sum_{i=1}^3 U_{ei}^2 M_i \propto \eta_\nu. \quad (30)$$

This is proportional to the standard parameter η_ν but in the case of very light right-handed neutrinos, e.g. $M_i \approx m_i$, the contribution becomes negligible because of the strong suppression with the heavy right-handed W boson mass.

In general, we consider right-handed neutrinos both lighter and heavier than 100 MeV and use (26) to calculate the contribution. In addition, the relevant nuclear matrix element changes; for $M_i \gg 100$ MeV it approaches $\mathcal{M}'_N \rightarrow \mathcal{M}_N$ whereas for $M_i \ll 100$ MeV it approaches $\mathcal{M}'_N \rightarrow \mathcal{M}_\nu$. For intermediate values, we use a simple smooth interpolation scheme within the regime 10 MeV—1 GeV, which yields a sufficient accuracy for our purposes.

C. Right-handed triplet scalar

Finally, the exchange of a doubly charged right-handed triplet scalar gives

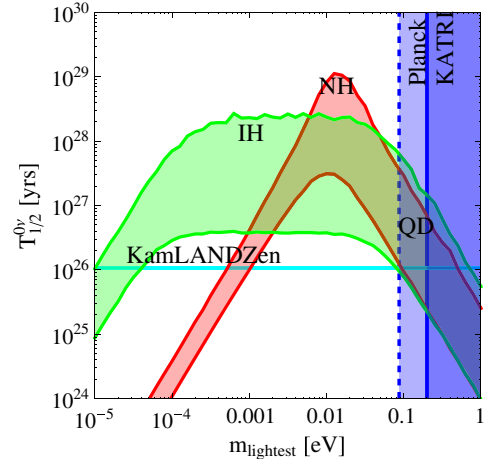


FIG. 2. $0\nu\beta\beta$ decay half-life as a function of the lightest neutrino mass in the case of normal hierarchical (NH) and inverse hierarchical (IH) light neutrinos in red and green bands respectively. The other parameters are fixed as $M_{W_R} = 5$ TeV, $M_{\delta_R^-} \approx 5$ TeV and the heaviest right-handed neutrino mass is 1 TeV. The gauge couplings are assumed universal, $g_L = g_R$, and the intermediate values for the nuclear matrix elements are used, $\mathcal{M}_\nu = 4.5$, $\mathcal{M}_N = 270$. The bound on the sum of light neutrino masses from the KATRIN and Planck experiments are represented as vertical lines. The bound from KamLAND-Zen experiment is presented in horizontal line for Xenon isotope. The bands arise due to 3σ range of neutrino oscillation parameters and variation in the Majorana phases from $0 - 2\pi$.

TABLE II. Phase space factor G_{01} and ranges of nuclear matrix elements for light and heavy neutrino exchange for the isotopes ^{76}Ge and ^{136}Xe [47].

Isotope	G_{01} (yr^{-1})	\mathcal{M}_ν	\mathcal{M}_N
^{76}Ge	5.77×10^{-15}	2.58–6.64	233–412
^{136}Xe	3.56×10^{-14}	1.57–3.85	164–172

$$\eta_\Delta = \frac{m_p}{M_{\delta_R}^2} \left(\frac{g_R}{g_L}\right)^4 \left(\frac{M_W}{M_{W_R}}\right)^4 \sum_{i=1}^3 U_{ei}^2 M_i \propto \eta_\nu. \quad (31)$$

This expression is also proportional to the standard η_ν because the relevant coupling of the triplet scalar is proportional to the right-handed neutrino mass.

D. Numerical estimate

In the following, we numerically estimate the half-life for $0\nu\beta\beta$ decay of the isotope ^{136}Xe as shown in Fig. 2. We use the current values of masses and mixing parameters from neutrino oscillation data reported in the global fits taken from Ref. [46]. For the $0\nu\beta\beta$ phase space factors and nuclear matrix elements we use the values given in Table II. In Fig. 2, we show the dependence of the $0\nu\beta\beta$ decay half-life on the lightest neutrino mass, i.e. m_1 for normal and m_3 for inverse hierarchical neutrinos. The other model parameters are fixed as

$$g_R = g_L, \quad M_{W_R} = M_{\delta_R} \approx 5 \text{ TeV}, \quad M_N^{\text{heaviest}} = 1 \text{ TeV}. \quad (32)$$

The lower limit on lightest neutrino mass is derived to be $m_{<} \approx 0.9 \text{ meV}$, 0.01 meV for NH and IH pattern of light neutrino masses respectively by saturating the KamLAND-Zen experimental bound.

As for the experimental constraints, we use the current best limits at 90% C.L., $T_{1/2}^{0\nu}(^{136}\text{Xe}) > 1.07 \times 10^{26} \text{ yr}$ and $T_{1/2}^{0\nu}(^{76}\text{Ge}) > 2.1 \times 10^{25} \text{ yr}$ from KamLAND-Zen [3] and the GERDA Phase I [48], respectively. Representative for the sensitivity of future $0\nu\beta\beta$ experiments, we use the expected reach of the planned nEXO experiment, $T_{1/2}^{0\nu}(^{136}\text{Xe}) \approx 6.6 \times 10^{27} \text{ yr}$ [49]. As for the other experimental probes on the neutrino mass scale, we use the future sensitivity of the KATRIN experiment on the effective single β decay mass $m_\beta \approx 0.2 \text{ eV}$ [50] and the current limit on the sum of neutrino masses from cosmological observations, $\Sigma_i m_i \lesssim 0.7 \text{ eV}$ [51].

For a better understand of the interplay between the left- and right-handed neutrino mass scales, we show in Fig. 3 the $0\nu\beta\beta$ decay half-life as a function of the lightest neutrino mass and the heaviest neutrino mass for a normal (left) and inverse (right) neutrino mass hierarchy. The other model parameters are fixed, with right-handed gauge boson and doubly-charged scalar masses of 5 TeV. The oscillation parameters are at their best fit values and the Majorana phases are always chosen to yield the smallest rate at a given point, i.e. the longest half life. The nuclear matrix employed are at the lower end in Table II. This altogether yields the longest, i.e. most pessimistic, prediction for the $0\nu\beta\beta$ decay half-life. The red-shaded area is already excluded with a predicted half life of 10^{26} yr or faster. As expected, this sets an upper limit on the lightest neutrino mass $m_{\text{lightest}} \lesssim 1 \text{ eV}$, but it also puts stringent constraints on the mass scale of the right-handed neutrinos. For an inverse hierarchy, the range $50 \text{ MeV} \lesssim M_2 \lesssim 5 \text{ GeV}$ is excluded whereas in the normal hierarchy case, large M_3 can be excluded if there is a strong hierarchy, $m_1 \rightarrow 0$. This is due the large contribution of the lightest heavy neutrino N_1 in such a case.

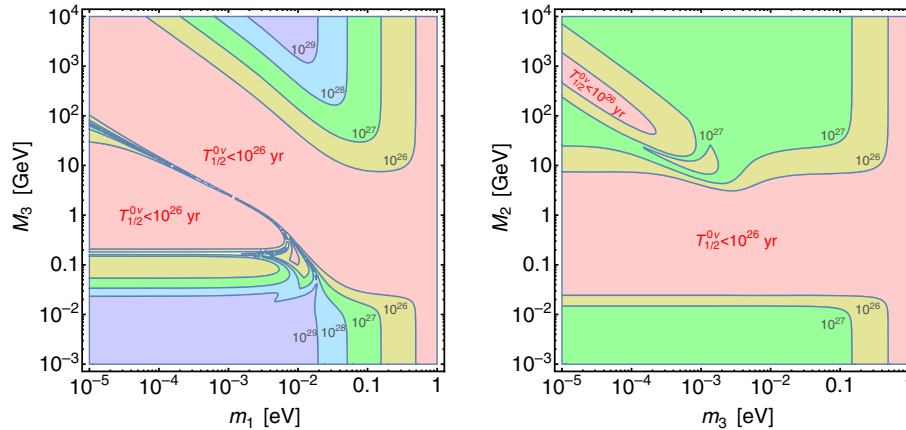


FIG. 3. Half-life of $0\nu\beta\beta$ decay in Xe as a function of the lightest and the heaviest neutrino mass for a normal (left) and inverse (right) neutrino mass hierarchy. The contours denote the half-life in years. Best-fit oscillation data are used and the Majorana phases are chosen to yield the longest half-life. Likewise, the smallest values of the nuclear matrix elements in Table II are employed. The other model parameters are chosen as $g_R = g_L$ and $M_{W_R} = M_\Delta = 5 \text{ TeV}$.

VII. CONCLUSION

We have presented a left-right symmetric model with additional vectorlike fermions in order to simultaneously explain the charged fermion and Majorana neutrino masses. The quark and charged lepton masses and mixing is realized via a universal seesaw mechanism. Although spontaneous symmetry breaking is achieved with two doublet Higgs fields with nonzero $B - L$ charge, we have introduced scalar triplets with small induced VEVs such that they give Majorana masses to light as well as heavy neutrinos. The Majorana nature of these neutrinos leads to $0\nu\beta\beta$ decay and it is found that the right-handed currents

play an important role in discriminating between the mass hierarchy as well as the absolute scale of light neutrinos.

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