Universal Seesaw and $0\nu\beta\beta$ in new 3331 Left-Right Symmetric Model

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We consider a class of left-right symmetric model with enlarged gauge group $SU(3) \times SU(3)_L \times SU(3)_R \times U(1)_Y$ without having scalar bitriplet. In the absence of scalar bitriplet, there is no Dirac mass term for fermions including usual quarks and leptons. We introduce new isosinglet vector-like fermions so that all the fermions get their masses through a universal seesaw mechanism. We extend our discussion to neutrino mass and its implications in neutrinoless double beta decay ($0\nu\beta\beta$). We show that for TeV scale $SU(3)_R$ gauge bosons, the heavy-light neutrino mixing contributes dominantly to $0\nu\beta\beta$ that can be observed at ongoing experiments. Towards the end we also comment on different possible symmetry breaking patterns of this enlarged gauge symmetry to that of the standard model.

I. INTRODUCTION

The Standard Model (SM) of particle physics has been the most successful phenomenological theory specially after the discovery of its last missing piece, the Higgs boson at the Large Hadron Collider (LHC) back in 2012 with subsequent null results for Beyond Standard Model (BSM) searches. However, the SM fails to address several observed phenomena as well as theoretical questions. For example, it fails to explain the sub-eV neutrino mass [1–6], the origin of parity violation in weak interactions and the origin of three fermion families. The first two questions can be naturally addressed within the framework of the Left-right symmetric model (LRSM) [7, 8], one of the most widely studied BSM frameworks. These models not only explain tiny neutrino masses naturally through seesaw mechanism but also give rise to an effective parity violating SM at low energy through spontaneous breaking of a parity preserving symmetry at high scale. The conventional LRSM based on the gauge group $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_B - L$ can be enhanced to a more general LRSM based on the gauge group $SU(3)_c \times SU(3)_{L,R} \times SU(3)_R \times U(1)_X$ (or in short 3331). The advantage of such an up gradation of the gauge symmetry is the ability of the latter in providing an explanation to the origin of three fermion families of the SM in addition to having other generic features of the LRSM. In such a model, the number of three fermion generations is no longer a choice, but a necessity in order to cancel chiral anomalies. In such models, where the usual lepton and quark representations are enlarged from a fundamental of $SU(2)_L$ in the SM to a fundamental of $SU(3)_L$, the number of generations must be equal to the number of colors in order to cancel the anomalies [9]. This is in contrast with the SM or the usual LRSM where the gauge anomalies are canceled within each fermion generation separately. One can also build such a sequential 3331 model by including additional chiral fermions. But since such a model does not explain the origin of three families from the anomaly cancellation point of view and contain non-minimal chiral fermion content, we stick to discussing a special type of non-sequential 3331 model here.

There have been a few works [10–14] recently done within the framework of such 3331 models with different motivations. Particularly from the origin of neutrino mass point of view, the work [10] considered a scalar sector comprising of bitriplets plus sextets which gives rise to tiny neutrino masses through canonical type I [15] and type II [16, 17] seesaw. Another recent work [12] studied a specific 3331 model with bitriplet and triplet scalar fields that can explain tiny neutrino masses through inverse [18, 19] and linear seesaw mechanism [19]. The earlier work [13] considered effective higher dimensional operators to explain fermion masses in 3331 models while the recent work [14] studied the model and several of its variants from LHC phenomenology point of view. Here, we simply consider another possible way of generating fermion masses in 3331 models through the universal seesaw mechanism [20, 22] where all fermions acquire their masses through a common seesaw mechanism. Incorporating additional vector like fermion pairs corresponding to each fermion generation, we show that the correct fermion mass spectrum can be generated in such a model with a scalar sector where all of them transform as fundamentals under $SU(3)_{L,R}$ without the need of bi-fundamental and sextet scalars shown in [10, 12] for the implementation of different seesaw mechanism for neutrino masses. We also discuss the possibilities of light neutral fermions apart from sub-eV active neutrinos, their role in neutrinoless double beta decay ($0\nu\beta\beta$) and different possible symmetry breaking chains of the gauge symmetry $SU(3)_c \times SU(3)_L \times SU(3)_R \times U(1)_X$ to that of the SM. We show that for TeV scale $SU(3)_R$ gauge bosons, the right handed neutrinos are constrained to lie around the keV mass regime having interesting conse-

$^{1}$See Refs [23–26] for implementation of universal seesaw mechanism for fermion mass generation within left-right symmetric model.
quences for $0\nu\beta\beta$. We find that although the pure heavy neutrino contribution to $0\nu\beta\beta$ remains suppressed compared to the one from light neutrinos, the heavy-light neutrino mixing which can be quite large in this model without any fine-tuning, gives a large contribution to $0\nu\beta\beta$ keeping it within experimental reach.

This letter is organized as follows. In section II we briefly discuss the model with the details of the particle spectrum, fermion masses via universal seesaw and gauge boson masses. In sections III and IV we discuss the contributions to $0\nu\beta\beta$ from purely light (heavy) neutrinos and heavy-light neutrino mixing respectively. Finally we discuss about different possible symmetry breaking chains in section V and then conclude in section VI.

II. THE MODEL FRAMEWORK

A. Particle Spectrum

The usual fermions transform under $SU(3)_C \times SU(3)_L \times SU(3)_R \times U(1)_X$ as

$$\Psi_{aL} = \left( \frac{\nu_{aL}}{Q_{aL}}, \frac{\ell_{aL}}{\epsilon_{aL}} \right), \quad \Psi_{aR} = \left( \frac{\nu_{aR}}{Q_{aR}}, \frac{\ell_{aR}}{\epsilon_{aR}} \right),$$

$$Q_{mL} = \left( \frac{d_{aL}}{u_{aL}}, \frac{a_{aL}}{f_{aL}} \right), \quad Q_{mR} = \left( \frac{d_{aR}}{u_{aR}}, \frac{a_{aR}}{f_{aR}} \right),$$

$$Q_{3L} = \left( \frac{d_{3L}}{u_{3L}}, \frac{a_{3L}}{f_{3L}} \right), \quad Q_{3R} = \left( \frac{d_{3R}}{u_{3R}}, \frac{a_{3R}}{f_{3R}} \right),$$

with $a=1,2,3$ whereas $m=1,2$.

The transformation of the fields under the gauge symmetry $SU(3)_C \times SU(3)_L \times SU(3)_R \times U(1)_X$ are given in table I.

| Particle | $SU(3)_C \times SU(2)_L \times SU(3)_R \times U(1)_X$ |
|----------|--------------------------------------------------|
| $Q_{mL}$ | $(3,3^*,1,-\frac{2}{3})$                        |
| $Q_{mR}$ | $(3,1,3^*,\frac{2}{3})$                         |
| $Q_{3L}$ | $(3,3,1,\frac{2}{3})$                          |
| $Q_{3R}$ | $(3,1,3,\frac{2}{3})$                          |
| $\psi_{aL}$ | $(1,3,1,\frac{2}{3})$                        |
| $\psi_{aR}$ | $(1,1,3,\frac{2}{3})$                        |
| $U_{1L}$ | $(3,1,1,\frac{2}{3})$                          |
| $D_{1L}$ | $(3,1,1,\frac{2}{3})$                          |
| $E_{1L}$ | $(1,1,1,-1)$                                    |
| $N_{1L}$ | $(1,1,1,0)$                                     |
| $\chi_L$ | $(1,3,1,-\frac{2}{3})$                         |
| $\chi_R$ | $(1,1,3,-\frac{2}{3})$                         |
| $\phi_L$ | $(1,3,1,\frac{2}{3})$                          |
| $\phi_R$ | $(1,1,3,\frac{2}{3})$                          |

Assuming $g=0$, if the neutral components of the scalar fields acquire their vacuum expectation value (vev) as

$$\langle \chi_L \rangle \equiv \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_1 \end{pmatrix}, \quad \langle \phi_L \rangle \equiv \frac{1}{\sqrt{2}} \begin{pmatrix} v_2 \omega_L \\ 0 \end{pmatrix},$$

$$\langle \chi_R \rangle \equiv \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_1 \end{pmatrix}, \quad \langle \phi_R \rangle \equiv \frac{1}{\sqrt{2}} \begin{pmatrix} v_2 \omega_R \\ 0 \end{pmatrix}$$

B. Fermion Mass

The Yukawa Lagrangian can be written as

$$\mathcal{L}_Y = - [Y_D]_{ma} (\overline{\nu}_m \phi_L D_{aL} + \overline{\nu}_m \phi_R D_{aR}) - [M_u]_{ab} \overline{U}_{aL} U_{bR}$$

$$- [Y_U]_{ma} (\overline{Q}_{mL} \chi_L U_{aR} + \overline{Q}_{mR} \chi_R U_{aL}) - [M_{\nu}]_{ab} \overline{\nu}_a \nu_b$$

$$- [Y_D']_{ma} (\overline{Q}_{3L} \phi'_L D_{aR} + \overline{Q}_{3R} \phi'_R D_{aL})$$

$$- [Y_U']_{ma} (\overline{Q}_{3L} \phi'_L U_{aR} + \overline{Q}_{3R} \phi'_R U_{aL})$$

$$- [Y_{\nu}]_{ma} (\overline{\nu}_m \phi_L E_{aR} + \overline{\nu}_m \phi_R E_{aL}) - [M_{\nu}]_{ab} \overline{E}_a E_b$$

$$- [Y_{\nu}']_{ma} (\overline{\nu}_m \phi'_L N_{bL} + \overline{\nu}_m \phi'_R N_{bR}) - [M_{\nu}]_{ab} \overline{N}_a N_b$$

$$- [M_{LL}]_{ab} N_{aL} N_{bL} - [M_{RR}]_{ab} N_{aR} N_{bR} + h.c.$$ (3)

After integrating out the heavy fermions, we can write

![Feynman diagram for Dirac mass of fermions within $SU(3)_C \times SU(3)_L \times SU(3)_R \times U(1)_X$ model.](image)

down the effective Yukawa terms for charged fermions of the standard model as follows

$$y_u = Y_u \frac{v_1}{M_D} Y'_u^T,$$

$$y_d = Y_d \frac{v_2}{M_D} Y'_d^T,$$

$$y_e = Y_e \frac{v_1}{M_E} Y'_e^T$$ (4)
Similarly, the heavy neutral singlet fields $N_{L,R}$ can be integrated out to generate the effective mass matrix of neutrinos $\nu_{L},\nu_{R}$ which contains a Dirac mass term and two Majorana mass terms. The effective Dirac mass as well as Majorana mass terms are given by

$$
M_D = Y_N \frac{1}{M_{RR}} M_{LR}^T \frac{1}{M_{RR}} Y_N^T \nu_{L2} \nu_{R2},
$$

$$
M_L = Y_N \frac{1}{M_{RR}} Y_N^T \nu_{L2}^2,
$$

$$
M_R = Y_N \frac{1}{M_{RR}} Y_N^T \nu_{R2}^2.
$$

There are additional neutrino leptons $\xi_{L}, \xi_{R}$ which acquire Dirac and Majorana masses similar to $\nu_{L,R}$ shown above. They are given by

$$
M_{\xi_{L,D}} = Y_N \frac{1}{M_{RR}} M_{LR}^T \frac{1}{M_{RR}} Y_N^T \omega_{L,R},
$$

$$
M_{\xi_{L,L}} = Y_N \frac{1}{M_{RR}} Y_N^T \omega_{L}^2,
$$

$$
M_{\xi_{R,R}} = Y_N \frac{1}{M_{RR}} Y_N^T \omega_{R}^2.
$$

The origin of the Dirac masses can be understood from

$$
\langle \phi_{L} \rangle \rightarrow \langle \phi_{L} \rangle, \quad \langle \phi_{L} \rangle \rightarrow \langle \phi_{L} \rangle, \quad \langle \phi_{L} \rangle \rightarrow \langle \phi_{L} \rangle,
$$

where the factor $W_{\mu}$ is defined as

$$
W_{\mu} = \left( \begin{array}{c}
\frac{g_L}{2} W^{L}_{\mu} + g_X (-2/3) B_{\mu} \\
\frac{g_R}{2} W^{R}_{\mu} + g_X (-2/3) B_{\mu}
\end{array} \right) \chi^L \chi^R.
$$

As a result of this particular structure of the mass matrix, both the mass eigenvalues for left-handed and right-handed neutrinos are proportional to each other. The two mixing matrices are related as

$$
V_{\nu_{L},N} = V_{\nu_{L},R} \equiv U,
$$

where $U$ is the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) leptonic mixing matrix. For a representative set of input model parameters like $v_{L} \approx 174$ GeV, $v_{R} \approx 10$ TeV, and $v_{1} \approx 0.1$ eV, the right-handed neutrino masses lie in the range of keV scale. Such light keV scale right handed neutrinos can also have very interesting implications for cosmology.

### C. Gauge Boson Mass

The relevant kinetic terms leading to gauge boson masses are given by

$$
\mathcal{L}_{G.B.} \supset \frac{g_{L}}{2} W_{\mu}^{L} + g_X (-2/3) B_{\mu} \chi^L \chi^R + \frac{g_{R}}{2} W_{\mu}^{R} + g_X (-2/3) B_{\mu} \chi^L \chi^R + \frac{g_{L}}{2} W_{\mu}^{L} + g_X (1/3) B_{\mu} \phi^L + \frac{g_{R}}{2} W_{\mu}^{R} + g_X (1/3) B_{\mu} \phi^R
$$

where the factor $W_{\mu}^{L,R}$ is defined as

$$
W_{\mu}^{L,R} = \sum_{i=1}^{8} W_{L,R,\mu,i} A_{i} W^{L,R} = \left( \begin{array}{c}
W_{\mu} - W_{\mu}^{+} \frac{1}{\sqrt{3}} W_{\mu}^{+} \frac{1}{\sqrt{3}} W_{\mu}^{+} V_{\mu}^{-q}
\end{array} \right)_{L,R}
$$

Using respective vev's for scalar fields shown in equation (9) we can derive the gauge boson masses for the present model. In the gauge boson spectrum, we have
One massless photon $A$,

- Four neutral gauge bosons $Z_{L,R}, Z'_{L,R}$,
- Four charged gauge bosons $W_{L,R}^\pm$,
- Four gauge bosons with charge $q + 1$, $X_{L,R}^{(1+q)}$,
- Four gauge bosons with charge $q$, $Y_{L,R}^{\pm}$.

III. $0^{\nu}\beta\beta$ WITH PURELY LIGHT (HEAVY)
NEUTRINO CONTRIBUTIONS

We find that the light sub-eV scale left-handed neutrinos ($\nu_L$) and the heavy right-handed neutrinos ($\nu_R$) with keV scale masses can give sizable contributions to neutrinoless double beta decay. Since the bitriplet scalar is absent in the present left-right symmetric 3331 model, there are no Dirac mass term for light neutrinos at tree level. However one can eventually generate Majorana masses for $\nu_L$ and $\nu_R \equiv N_R$ through universal seesaw, see Fig. 2. Such Majorana nature of neutrinos violate lepton number by two units and thus, contributes to $0^{\nu}\beta\beta$ decay. The Feynman diagram for $0^{\nu}\beta\beta$ decay is depicted in Fig. 4 due to exchange of left-handed as well as right-handed neutrinos. The corresponding Feynman amplitudes due to exchange of left-handed and right-handed neutrinos are given by

$$A_{\nu_L} \propto G_F^2 \frac{U_{ei}^2 m_i}{p^2}$$
$$A_{\nu_R} \propto G_F^2 \left( \frac{M_{W_L}^2}{M_{W_R}^2} \right) U_{ei}^2 \frac{M_i}{p^2}$$  \hspace{1cm} (11)

The inverse half-life for a given isotope for $0^{\nu}\beta\beta$ decay due to exchange of left-handed light neutrinos via left-handed currents, and right-handed neutrinos via right-handed currents is given by

$$[T_{1/2}^{0^{\nu}}]^{-1} = G_{01} \left( |M_{\nu_L} \eta_L|^2 + |M_{\nu_R} \eta_R|^2 \right),$$  \hspace{1cm} (12)

where $G_{01}$ is $0^{\nu}\beta\beta$ phase space factor, $M_i$ correspond to the nuclear matrix elements (NME) and $\eta_i$ is the corresponding dimensionless particle physics parameter. Since we have $M_i \approx 1 - 10$ MeV masses of right-handed neutrinos in the present model and satisfying $|M_{\nu_R}^2| \ll p^2$ where $p$ being the neutrino virtually momentum around 100 MeV, the NMEs for right-handed neutrinos and left-handed neutrinos are same i.e., $M_{\nu_L} = M_{\nu_R}$.

| $\eta_i$ | Effective Mass Parameter |
|---|---|
| $\eta_{\nu_L} \approx \frac{1}{m_e} \sum_{i=1}^3 U_{ei}^2 m_i$ | $m_{\nu_L}^{ee} \approx \frac{1}{m_e} \sum_{i=1}^3 U_{ei}^2 m_i$ |
| $\eta_{\nu_R} \approx \frac{1}{m_e} \left( \frac{M_{W_L}^2}{M_{W_R}^2} \right)^4 \sum_{i=1}^3 U_{ei}^2 M_i$ | $m_{\nu_R}^{ee} \approx \frac{1}{m_e} \left( \frac{M_{W_L}^2}{M_{W_R}^2} \right)^4 \sum_{i=1}^3 U_{ei}^2 M_i$ |

TABLE II: Dimensionless particle physics parameters due to exchange of left-handed and right-handed neutrinos and the corresponding effective mass parameters.

A. Standard mechanism via left-handed neutrinos $\nu_L$

The standard mechanism for neutrinoless double beta decay due to exchange of left-handed neutrinos via left-handed currents gives dimensionless particle physics parameter as,

$$\eta_{\nu_L} = \frac{1}{m_e} \sum_{i=1}^3 U_{ei}^2 m_i = \frac{m_{\nu_L}^{ee}}{m_e}.$$  \hspace{1cm} (13)

Here, $m_e$ is the electron mass. The effective mass parameter for standard mechanism is explicitly given by

$$m_{\nu_L}^{ee} = |c_{12}^2 s_{13}^2 m_1 + s_{12}^2 c_{13}^2 m_2 e^{i\alpha} + s_{13}^2 m_3 e^{i\beta}|.$$  \hspace{1cm} (14)

B. New contribution from right-handed neutrinos $\nu_R$

In the present left-right symmetric 3331 model, we found that the right-handed neutrino mass lies around a few keV (for TeV scale $W_R$ which is much less than its momentum, $M_i \ll |p|$). Under this condition, the propagator simplifies in a similar way as for the light neutrino exchange,

$$P_R \frac{p^4 + M_i}{p^2 - M_R^2} \approx M_i \frac{p^4}{p^2}.$$  \hspace{1cm} (15)

This results dimensionless particle physics parameter $\eta_{\nu_R}$ due to exchange of right-handed neutrinos via right-handed currents as,

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

where the proportionality relation between $\eta_{\nu_R}$ and $\eta_{\nu_L}$ appears at the last step due to the proportionality between heavy and light neutrino mass matrix discussed.
above. After a little simplification, the effective mass parameter due to exchange of right-handed neutrinos can be expressed as,

$$m_{ee} \propto \left( \frac{M_{el}}{M_{el}} \right)^4 \sum_{i=1}^{3} U_{ei}^2 M_i \propto m_e. \quad (17)$$

It is clear from eq. (6) that both light and heavy neutrino mass eigenvalues are proportional to each other as

$$M_i \propto v^2 m_i. \quad (18)$$

Thus, one can express $M_R$ mass as

$$M_R^2 \approx \frac{1}{4} v^2 g_R = \frac{1}{4} g_R v^2 m_i. \quad (19)$$

As a result, the effective Majorana mass parameter– with $g_L \approx g_R$ and $M_W \approx 2 g_L v_L$ – is modified to

$$m_{ee} \propto \left( \frac{m_1}{M_1} \right)^2 \sum_{i=1}^{3} U_{ei}^2 M_i. \quad (20)$$

Comparing this with the light neutrino contribution $m_{ee} = \sum_{i=1}^{3} U_{ei}^2 m_i$, it is straightforward to estimate that the heavy neutrino contribution is suppressed by a factor of $m_i/M_1 = (v_L/v_R)^2$ compared to the light neutrino contribution.

C. Numerical Results

The total contribution to inverse half-life for neutrino-less double beta decay for a given isotope in the present left-right symmetric 3331 model due to exchange of left-handed as well as right-handed neutrinos is given by

$$[T_{1/2}^{0\nu}]^{-1} = G_0 \frac{M_{el}}{m_e} |m_{ee}^\text{eff}|^2, \quad (21)$$

where,

$$|m_{ee}^\text{eff}|^2 = |U_{ei}^2 m_i |^2 + \left| \frac{M_{el}^4}{M_{el}^4} U_{ei}^2 M_1 \right|^2, \quad (22)$$

1. For NH Pattern

For normal hierarchical (NH) pattern of light neutrinos we consider the following mass structures for left-handed and right-handed neutrinos,

$$m_1 = m_{\text{lightest}}, \quad m_2 = \sqrt{m_1^2 + \Delta m^2_{\text{sol}}},$$

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

$$M_2 = M_3, \quad M_1 = \frac{m_1}{m_3} M_3, \quad M_2 = \frac{m_2}{m_3} M_3. \quad (23)$$

where $M_3$ is fixed around few keV range, as a result of choosing the $W_R$ mass scale at a few TeV. The analytic form for effective mass parameters due to exchange of right-handed neutrinos is given by

$$m_{ee}^\text{NH} = \left| c_{12}^2 c_{13}^2 m_1 + s_{12}^2 c_{13}^2 m_2 e^{i\alpha} + s_{13}^2 m_3 e^{i\beta} \right|,$$

$$|m_{ee}^\text{NH}| = \left| \frac{m_1}{M_3} \right|^2 M_3 \left| \frac{m_1}{m_3} c_{12} c_{13}^2 m_2 e^{i\alpha} + \frac{m_2}{m_3} s_{12}^2 c_{13}^2 e^{i\alpha} + s_{13}^2 e^{i\beta} \right|. \quad (24)$$

2. For IH Pattern

Similarly for inverse hierarchical (IH) pattern of light neutrinos, the masses for light left-handed and right-
The new physics contributions to $0\nu\beta\beta$ decay arising from purely left-handed currents due to exchange of $\nu_3$ and $N_R$ scale right-handed neutrinos results in the following effective mass parameter, 

$$M_{\nu} = \begin{pmatrix} Y_{N} \frac{1}{M_{RR}} Y_{N}^{T} v_{2L}^{2} & Y_{N} \frac{1}{M_{RR}} Y_{N}^{T} v_{2L} v_{2R} \\ Y_{N} \frac{1}{M_{RR}} Y_{N}^{T} v_{2R} & Y_{N} \frac{1}{M_{RR}} Y_{N}^{T} v_{2L}^{2}R \end{pmatrix}.$$  

(28)

In the limit $M_{L} \ll M_{D} \ll M_{R}$, the type-I seesaw contribution to light neutrino mass is given by 

$$M_{\nu}^{I} = -M_{D} \frac{1}{M_{R}} M_{R},$$  

(29)

and the light-heavy neutrino mixing is proportional to $M_{D}/M_{R} \approx v_{2L}/v_{2R}$. With $v_{2L} \approx 174$ GeV and $v_{2R}$ around few TeV, we find that light-heavy neutrino mixing is large of the order of $\lesssim 0.1$. This large value of light-heavy neutrino mixing where heavy neutrinos are fixed at few keV scale, can contribute to $0\nu\beta\beta$ significantly.

### A. Purely left-handed current effects

The new physics contributions to $0\nu\beta\beta$ decay arising from purely left-handed currents due to exchange of $\nu_{\ell}$ and $N_{R}$ scale right-handed neutrinos results in the following effective mass parameter,

$$m_{\nu_{\ell},LL}^{N} = \sum_{i=1}^{3} \frac{V_{\nu_{N}}^{T} M_{i}}{v_{\ell}}.$$  

(30)

here $M_{i}$ is in keV range and $V_{\nu_{N}}^{T}$ is the light-heavy neutrino mixing. Since $V_{\nu_{N}} \times v_{2L}/v_{2R} \approx 0.01$, the effective mass parameter for $0\nu\beta\beta$ is estimated to be $m_{\nu_{\ell},LL}^{N} = (0.01)^{2} \cdot 10^{3}$ eV which is of the order of 0.1 eV, saturating the KamLAND-Zen [30] bound.

### B. From $\lambda -$ diagram

The effective mass parameters due to the $W_{L} - W_{R}$ mediated diagrams (known as $\lambda$ diagrams) shown in figure 6.

FIG. 6: $0\nu\beta\beta$ decay diagrams due to purely left-handed charge current interaction and with the exchange of $\nu_{\ell}$ and $N_{R}$.
In the scalar potential written above, the discrete left-right symmetry is assumed which ensures the equality of left and right sector couplings. However, as shown in earlier works [21] in the context of usual LRSM with universal seesaw that the scalar potential of such a model with exact discrete left-right symmetry is too restrictive and gives to either parity preserving \((v_L = v_R)\) solution or a solution with \((v_R \neq 0, v_L = 0)\) at tree level. While the first one is not phenomenologically acceptable the latter solution can be acceptable if a non-zero vev \(v_L \neq 0\) can be generated through radiative corrections [32]. While it may naturally explain the smallness of \(v_L\) compared to \(v_R\), it will constrain the parameter space significantly [32].}

Another way of achieving a parity breaking vacuum is to consider softly broken discrete left-right symmetry by considering different mass terms for the left and right sector scalars [8] [21]. As it was pointed out by the authors of [8], such a model which respects the discrete left-right symmetry everywhere except in the scalar mass terms, preserve the naturalness of the left-right symmetry in spite of radiative corrections. Another interesting way is to achieve parity breaking vacuum is to decouple the scale of parity breaking and gauge symmetry breaking by introducing a parity odd singlet scalar [33]. While we do not perform a detailed analysis of different possible symmetry breaking chains and their constraints on the parameter space of the model, we outline them pictorially in the cartoon shown in figure [8]. As can be seen from figure [8] there are seven different symmetry breaking chains through which the gauge symmetry of the model \(SU(3)_c \times SU(3)_L \times SU(3)_R \times U(1)_X\) can be broken down to that of the SM as summarized below.

- **One step breaking:** The vev’s satisfy \(v_{1L,2L} \ll v_{1R,2R} \approx \omega_L \approx \omega_R\) in this case.

- **Two step breaking:** The vev’s satisfy either \(v_{1L,2L} \ll \omega_L \ll v_{1R,2R} \approx \omega_R\) or \(v_{1L,2L} \ll v_{1R,2R} \ll \omega_R\) or \(v_{1L,2L} \ll \omega_R \ll v_{1R,2R} \approx \omega_L\). The usual 331 model presumes an intermediate stage in the first case while the usual LRSM or 3221 symmetry arises an intermediate symmetry in the second case. In the third case, the 3231 symmetry assumes an intermediate stage. The phenomenology of such asymmetric LRSM was discussed recently by [14].

- **Three step breaking:** This is possible in three different ways when the vev’s satisfy \(v_{1L,2L} \ll v_{1R,2R} \ll \omega_R \ll \omega_L\) or \(v_{1L,2L} \ll v_{1R,2R} \ll \omega_R \ll \omega_L\) or \(v_{1L,2L} \ll \omega_L \ll v_{1R,2R} \ll \omega_R\). One can have both the usual LRSM or 3221 or asymmetric LRSM (3321 or 3231) or the usual 331 model as an intermediate stage.

All these different symmetry breaking chains can not only provide a different phase transition history in cosmology but also give rise to different particle spectra including gauge bosons as well as neutral fermions which could be tested in different experiments.
We have demonstrated a class of left-right symmetric model with extended gauge group $SU(3)_c \times SU(3)_L \times SU(3)_R \times U(1)_X$ with a universal seesaw mechanism for fermion masses and mixing and the implications for neutrinoless double beta ($0\nu\beta\beta$) decay. The novel feature of the model is that masses and mixing for left-handed and right-handed neutrinos are exactly determined by oscillation parameters and lightest neutrino mass. This forces the heavy neutrino masses to lie in keV regime if the $W_R$ mass is fixed at a few TeV. We show that for such a case, the heavy neutrino contribution to $0\nu\beta\beta$ remains suppressed compared to the usual light neutrino contribution. We also show that for such a TeV scale model, the heavy-light neutrino mixing can be quite large and can contribute substantially to $0\nu\beta\beta$ diagrams, keeping it within experimental reach. In the end we have discussed the scalar potential and possible symmetry breaking patterns that can be allowed for spontaneous breaking of the 3331 gauge symmetry to that of the standard model.

**VI. CONCLUSION**

We have demonstrated a class of left-right symmetric model with extended gauge group $SU(3)_c \times SU(3)_L \times SU(3)_R \times U(1)_X$ with a universal seesaw mechanism for fermion masses and mixing and the implications for neutrinoless double beta ($0\nu\beta\beta$) decay. The novel feature of the model is that masses and mixing for left-handed and right-handed neutrinos are exactly determined by oscillation parameters and lightest neutrino mass. This forces the heavy neutrino masses to lie in keV regime if the $W_R$ mass is fixed at a few TeV. We show that for such a case, the heavy neutrino contribution to $0\nu\beta\beta$ remains suppressed compared to the usual light neutrino contribution. We also show that for such a TeV scale model, the heavy-light neutrino mixing can be quite large and can contribute substantially to $0\nu\beta\beta$ diagrams, keeping it within experimental reach. In the end we have discussed the scalar potential and possible symmetry breaking patterns that can be allowed for spontaneous breaking of the 3331 gauge symmetry to that of the standard model.

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[1] S. Fukuda et al. (Super-Kamiokande), Phys. Rev. Lett. 86, 5656 (2001), hep-ex/0103033; Q. R. Ahmad et al. (SNO), Phys. Rev. Lett. 89, 171801 (2002), hep-ex/0204008; Phys. Rev. Lett. 90, 011302 (2003), nucl-ex/0204009; J. N. Bahcall and C. Pena-Garay, New J. Phys. 6, 63 (2004), hep-ph/0404061; K. Nakamura et al., J. Phys. G37, 075021 (2010).

[2] K. Abe et al. [T2K Collaboration], Phys. Rev. Lett. 107, 041801 (2011), arXiv:1106.2822 [hep-ex].

[3] Y. Abe et al., Phys. Rev. Lett. 108, 131801 (2012), arXiv:1112.6553 [hep-ex].

[4] F. P. An et al. [DAYA-BAY Collaboration], Phys. Rev. Lett. 108, 171803 (2012), arXiv:1203.1669 [hep-ex].

[5] J. K. Ahn et al. [RENO Collaboration], Phys. Rev. Lett. 108, 191802 (2012), arXiv:1204.0626 [hep-ex].

[6] P. Adamson et al. (MINOS), Phys.Rev.Lett. 110, 171801 (2013).

[7] J. C. Pati and A. Salam, Phys. Rev. D10, 275 (1974); G. Senjanovic and R. N. Mohapatra, Phys. Rev. D12, 1502 (1975); R. N. Mohapatra and R. E. Marshak, Phys. Rev. Lett. 44, 1316 (1980); J. F. Gunion, J. Grifols, A. Mendez, B. Kayser and F. I. Olness, Phys. Rev. D40, 1546 (1989); N. G. Deshpande, J. F. Gunion, B. Kayser and F. I. Olness, Phys. Rev. D44, 837 (1991).

[8] R. N. Mohapatra and J. C. Pati, Phys. Rev. D11, 2558 (1975).

[9] M. Singer, J. W. F. Valle and J. Schechter, Phys. Rev. D22, 758 (1980); F. Pisano and V. Pleitez, Phys. Rev. D46, 410 (1992); P. H. Frampton, Phys. Rev. Lett. 69, 2889 (1992); R. Foot, O. F. Hernandez, F. Pisano and V. Pleitez, Phys. Rev. D47, 4158 (1993).

[10] M. Reig, J. W. F. Valle and C. A. Vaquera-Araujo, arXiv:1611.02066 [hep-ph].

[11] E. T. Franco and V. Pleitez, arXiv:1611.05658 [hep-ph].

[12] M. Reig, J. W. F. Valle and C. A. Vaquera-Araujo, arXiv:1611.04571 [hep-ph].

[13] A. G. Dias, C. A. de S.Pires and P. S. Rodrigues da Silva, Phys. Rev. D 82, 035013 (2010); arXiv:1003.3260 [hep-ph].

[14] D. T. Huang and P. V. Dong, Phys. Rev. D93, 095019 (2016).

[15] P. Minkowski, Phys. Lett. B67, 421 (1977); M. Gell-Mann, P. Ramond, and R. Slansky (1980), print-80-1731; T. Yanagida (1979), in Proceedings of the Workshop on the Baryon Number of the Universe and Unified Theories, Tsukuba, Japan, 13-14 Feb 1979; R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980); J. Schechter and J. W. F. Valle, Phys. Rev. D22, 2227 (1980).

[16] R. N. Mohapatra and G. Senjanovic, Phys. Rev. D23, 165 (1981).

[17] G. Lazarides, Q. Shafi and C Wetterich, Nucl. Phys. B181, 287 (1981); C. Wetterich, Nucl. Phys. B187, 343 (1981); J. Schechter and J. W. F. Valle, Phys. Rev. D25, 774 (1982); B. Brahmacari and R. N. Mohapatra, Phys. Rev. D58, 015001 (1998); R. N. Mohapatra, Nucl. Phys. Proc. suppl. 138, 257 (2005); S. Antusch and S. F. King, Phys. Lett. B597, 199 (2004).

[18] R. N. Mohapatra and J. W. F. Valle, Phys. Rev. D34, 1642 (1986); M. Gonzalez-Garcia and J. W. F. Valle, Phys. Lett. B216, 360 (1989); M. E. Catanato, R. Martinez and F. Ochoa, Phys. Rev. D86, 073015 (2012).
[19] E. K. Akhmedov, M. Lindner, E. Schnapka and J. W. F. Valle, Phys. Lett. B368, 270 (1996); E. K. Akhmedov, M. Lindner, E. Schnapka and J. W. F. Valle, Phys. Rev. D53, 2752 (1996); M. Malinsky, J. Romao and J. W. F. Valle, Phys. Rev. Lett. 95, 161801 (2005).
[20] B. Brahmachari, E. Ma and U. Sarkar, Phys. Rev. Lett. 91, 011801 (2003).
[21] A. Davidson and K. C. Wali, Phys. Rev. Lett. 59, 393 (1987); R. N. Mohapatra and Y. Zhang, JHEP 1406, 072 (2014).
[22] K. S. Babu and R. N. Mohapatra, Phys. Rev. Lett. 62, 1079 (1989); P. -H. Gu and M. Lindner, Phys. Lett. B698, 40 (2011).
[23] S. Patra, Phys. Rev. D 87, no. 1, 015002 (2013) doi:10.1103/PhysRevD.87.015002 [arXiv:1212.0612 [hep-ph]].
[24] P. S. B. Dev, R. N. Mohapatra and Y. Zhang, JHEP 1602, 186 (2016) doi:10.1007/JHEP02(2016)186 [arXiv:1512.08507 [hep-ph]].
[25] F. F. Deppisch, C. Hati, S. Patra, P. Pritimita and U. Sarkar, Phys. Lett. B 757, 223 (2016) doi:10.1016/j.physletb.2016.03.081 [arXiv:1601.00952 [hep-ph]].
[26] F. F. Deppisch, C. Hati, S. Patra, P. Pritimita and U. Sarkar, arXiv:1701.02107 [hep-ph].
[27] M. C. Gonzalez-Garcia, M. Maltoni, J. Salvado and T. Schwetz, JHEP 1212, 123 (2012) [arXiv:1209.3023 [hep-ph]].
[28] P. A. R. Ade et al. [Planck Collaboration], Astron. Astrophys. 594, A13 (2016) [arXiv:1502.01589 [astro-ph.CO]].
[29] S. Mertens [KATRIN Collaboration], Phys. Procedia 61, 267 (2015).
[30] A. Gando et al. [KamLAND-Zen Collaboration], Phys. Rev. Lett. 117, no. 8, 082503 (2016) Addendum: [Phys. Rev. Lett. 117, no. 10, 109903 (2016) doi:10.1103/PhysRevLett.117.109903, 10.1103/PhysRevLett.117.082503 [arXiv:1605.02889 [hep-ex]].
[31] D. Borah, Phys. Rev. D94, 075024 (2016).
[32] A. Kobakhidze and A. Spencer-Smith, JHEP 1308, 036 (2013).
[33] D. Chang, R. N. Mohapatra and M. K. Parida, Phys. Rev. Lett. 52, 1072 (1984).