A Lepton-specific Universal Seesaw Model with Left-Right Symmetry

Ayon Patra and Santosh Kumar Rai

1Centre for High Energy Physics, Indian Institute of Science, Bangalore - 560012, India
2Regional Centre for Accelerator-based Particle Physics, Harish-Chandra Research Institute, HBNI, Jhusi, Allahabad - 211019, India

We propose a left-right symmetric framework with universal seesaw mechanism for the generation of masses of the Standard Model quarks and leptons. Heavy vector-like singlet quarks and leptons are required for generation of Standard Model-like quark and lepton masses through seesaw mechanism. A softly broken $Z_2$ symmetry distinguishes the lepton sector and the quark sector of the model. This leads to the presence of some lepton-specific interactions that can produce unique collider signatures which can be explored at the current Large Hadron Collider run and also future colliders.

I. INTRODUCTION

Left-right symmetric (LRS) models are one of the most well motivated and widely studied extensions of the Standard Model (SM). The popularity of LRS models stem from the fact that in these models it is possible to explain several phenomenon which are not very well understood in the framework of SM. Fundamentally parity (P) is a good symmetry in these models and can be spontaneously broken at some high scale leading to a SM-like gauge structure at the electroweak scale. Thus we can understand the origin of parity violation as a spontaneously broken symmetry rather than it being explicitly broken. Parity symmetry also prevents one from writing P and Charge-Parity (CP) violating terms in the Quantum Chromodynamic (QCD) Lagrangian. Since CP violating terms in the color sector are highly constraint from neutron electric dipole measurements, the absence of these terms can solve the strong CP problem naturally without the need to introduce a global Peccei-Quinn symmetry. The gauge structure of these models force us to have a right-handed neutrino in the lepton multiplet. This right-handed neutrino can generate a light neutrino mass through the seesaw mechanism.

Generally, in LRS models, an $SU(2)_R$ triplet Higgs boson is responsible for generation of the right-handed neutrino mass while a bidoublet field is needed to produce the quark and lepton masses and CKM mixings. All these multitude of scalar fields make the scalar sector quite complicated. It would be interesting, on the other hand, to consider a Higgs spectrum consisting purely of doublets. This would be similar to the Two Higgs Doublet model (2HDM) but we would need four doublets instead of two (two similar to 2HDM and other two their right-handed counterparts). The model we study here is a lepton-specific scenario where one pair of Higgs doublets couple only to the leptons. This has the distinct advantage that the quark and charged lepton masses can be generated keeping the Yukawa couplings to be of the same order for each generation. Thus we can easily avoid the large hierarchy observed in the Yukawa sector of the SM.

To arrange the lepton-specific framework, we need to introduce an extra $Z_2$ symmetry under which a couple of Higgs boson doublets as well as the heavy lepton singlet fields are odd, all other fields being even. As these odd-$Z_2$ Higgs bosons get a non-zero vacuum expectation value (VEV), one expects this discrete $Z_2$ symmetry to be spontaneously broken. This could lead to domain walls and can make the model instable from a cosmological point of view. Such instabilities however can be avoided by introducing soft-breaking terms in the scalar potential. The consequence of these terms is that it leads to mixing between the different scalars in the doublets and can lead to interesting phenomenology. With the Large Hadron Collider (LHC) running, it is imperative to consider different scenarios for signals beyond the SM. In that spirit our model within the framework of left-right symmetry proposes new signals arising from a lepton-specific framework which generates all the SM fermion masses and gives lepton rich final states that could be observed or excluded at LHC.

All the fermion masses in this case are generated through universal seesaw mechanism by introducing singlet fermionic states. Most of the charged singlet fermions are quite heavy except the top quark partner to some extent, which is required to be lighter than the others and of the order of a few TeV. The other low lying states are the heavy neutral leptons and some extra scalars in the model.

*Electronic address: ayon@okstate.edu
†Electronic address: skrai@hri.res.in
1 A detailed study of various scalar sectors in LRS models is discussed in [5]
The rest of the paper is organized as follows. In Section II, we discuss our model in detail and in Section III we discuss the phenomenological implications of our model including the experimental constraints and possible collider signatures of our model. Section IV contains our conclusions and discussions.

II. MODEL AND LAGRANGIAN

We consider a left-right symmetric model with the gauge group being $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. An extra $Z_2$ symmetry is introduced which prevents several interactions facilitating a lepton-specific scenario. The charge of a particle in this model is defined as:

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

The hierarchy in the VEVs responsible for symmetry breaking are given as (for the universal seesaw mechanism to work):

$$v_{LQ} = v_{RQ}, \quad \langle H_{LQ}^0 \rangle = \langle H_{RQ}^0 \rangle, \quad \langle H_{LQ}^+ \rangle = \langle H_{RQ}^+ \rangle, \quad \langle H_{LQ}^- \rangle = \langle H_{RQ}^- \rangle,$$

with the condition that $v_{LQ}^2 + v_{Li}^2 = v_{EW}^2$. The hierarchy in the VEVs responsible for symmetry breaking are arranged as

$$v_{RQ}, v_{Li} >> v_{LQ}, v_{Li}, \quad (5)$$

This ensures a naturally heavy mass for the right-handed gauge bosons which have so far eluded any signal at the LHC.

We introduce a lepton-specific $Z_2$ symmetry under which the $E_L, E_R, N_L, N_R, H_L$ and $H_R$ fields are odd.
while all other fields are even. This prevents the $Z_2$-odd Higgs fields from interacting with the quarks. Table I has a list of all the particles along with their respective quantum numbers.

The covariant derivatives appearing in the kinetic terms of the Lagrangian that lead to interaction vertices of the fermions and scalars with the gauge bosons (for all the doublet fields) in this model are defined as

\[
D_\mu Q_L = \left[\partial_\mu - i \frac{g_L}{2} \tau.W_L^\mu - i \frac{g_V}{6} V_\mu\right] Q_L \\
D_\mu Q_R = \left[\partial_\mu - i \frac{g_R}{2} \tau.W_R^\mu - i \frac{g_V}{6} V_\mu\right] Q_R \\
D_\mu l_L = \left[\partial_\mu - i \frac{g_L}{2} \tau.W_L^\mu + i \frac{g_V}{2} V_\mu\right] l_L \\
D_\mu l_R = \left[\partial_\mu - i \frac{g_R}{2} \tau.W_R^\mu + i \frac{g_V}{2} V_\mu\right] l_R \\
D_\mu H_R = \left[\partial_\mu - i \frac{g_R}{2} \tau.W_R^\mu - i \frac{g_V}{2} V_\mu\right] H_R \\
D_\mu H_L = \left[\partial_\mu - i \frac{g_L}{2} \tau.W_L^\mu - i \frac{g_V}{2} V_\mu\right] H_L,
\]

where $V_\mu$ and $g_V$ are the gauge boson and the gauge coupling corresponding to the $U(1)_{B-L}$ gauge group, while $W_L$ and $W_R$ are the gauge bosons corresponding to the $SU(2)_L$ and $SU(2)_R$ gauge groups respectively. The gauge boson masses can be calculated from the kinetic terms for the Higgs boson fields involving the above

| Field | $SU(3)_C$ | $SU(2)_L$ | $SU(2)_R$ | $U(1)_{B-L}$ | $Z_2$ |
|-------|-----------|-----------|-----------|--------------|-------|
| $Q_L$ = $\left(\begin{array}{c} u \\ d \end{array}\right)_L$ | 3 | 2 | 1 | $\frac{1}{3}$ | + |
| $Q_R$ = $\left(\begin{array}{c} u \\ d \end{array}\right)_R$ | 3 | 1 | 2 | $\frac{1}{3}$ | + |
| $l_L$ = $\left(\begin{array}{c} \nu \\ e \end{array}\right)_L$ | 1 | 2 | 1 | -1 | + |
| $l_R$ = $\left(\begin{array}{c} \nu \\ e \end{array}\right)_R$ | 1 | 2 | 1 | -1 | + |
| $U_L, U_R$ | 3 | 1 | 1 | $\frac{4}{3}$ | + |
| $D_L, D_R$ | 3 | 1 | 1 | $-\frac{2}{3}$ | + |
| $E_L, E_R$ | 1 | 1 | 1 | -2 | - |
| $N_L, N_R$ | 1 | 1 | 1 | 0 | - |
| $H_{RQ}$ = $\left(\begin{array}{c} H_{RQ}^+ \\ H_{RQ}^0 \end{array}\right)$ | 1 | 1 | 2 | 1 | + |
| $H_{LQ}$ = $\left(\begin{array}{c} H_{LQ}^+ \\ H_{LQ}^0 \end{array}\right)$ | 1 | 2 | 1 | 1 | + |
| $H_{Rl}$ = $\left(\begin{array}{c} H_{Rl}^+ \\ H_{Rl}^0 \end{array}\right)$ | 1 | 1 | 2 | 1 | - |
| $H_{Ll}$ = $\left(\begin{array}{c} H_{Ll}^+ \\ H_{Ll}^0 \end{array}\right)$ | 1 | 2 | 1 | 1 | - |

TABLE I: Particle spectrum for lepton-specific LR Universal Seesaw Model
covariant derivatives. The charged gauge boson mass-squared matrix in the basis $(W_L^\pm, W_R^\pm)$ is given as:

$$
\begin{bmatrix}
\frac{1}{2}g_R^2(v_{RQ}^2 + v_{Rl}^2) & 0 \\
0 & \frac{1}{2}g_L^2(v_{LQ}^2 + v_{Ll}^2)
\end{bmatrix}
$$

(7)

We can clearly see that unlike the case of LRS with bidoublet scalar fields, there is no mixing between the two $W$ boson states in this case. The mass of the heavy $W_R$ gauge boson and the SM $W_L$ gauge boson states are thus trivially given as:

$$
M_{W_R^\pm}^2 = \frac{1}{2}g_R^2(v_{RQ}^2 + v_{Rl}^2), \quad M_{W_L^\pm}^2 = \frac{1}{2}g_L^2(v_{LQ}^2 + v_{Ll}^2).
$$

(8)

The neutral gauge boson mass-squared matrix in the basis $(W_{3R}, W_{3L}, V)$ is given as:

$$
\begin{bmatrix}
\frac{1}{2}g_R^2(v_{RQ}^2 + v_{Rl}^2) & 0 & \frac{1}{4}g_R^2g_V(v_{RQ}^2 + v_{Rl}^2) \\
0 & \frac{1}{2}g_L^2(v_{LQ}^2 + v_{Ll}^2) & \frac{1}{4}g_L^2g_V(v_{LQ}^2 + v_{Ll}^2) \\
\frac{1}{2}g_Rg_V(v_{RQ}^2 + v_{Rl}^2) & \frac{1}{4}g_Lg_V(v_{LQ}^2 + v_{Ll}^2) & \frac{1}{4}g_V(v_{RQ}^2 + v_{Rl}^2 + v_{LQ}^2 + v_{Ll}^2)
\end{bmatrix}
$$

(9)

This matrix has a zero eigenvalue corresponding to the massless photon state and two other non-zero eigenvalues corresponding to the $Z$ and the $Z_R$ bosons. In the limit $v_{EW} \ll v_{RQ}, v_{Rl}$ and keeping only terms up to $v_{EW}^2/v_{RQ(Rl)}^2$, the masses of the two massive neutral gauge bosons are given by:

$$
M_{Z_R}^2 \simeq \frac{1}{2} \left[ (g_R^2 + g_V^2)(v_{RQ}^2 + v_{Rl}^2) + \frac{g_V^2(v_{LQ}^2 + v_{Ll}^2)}{g_R^2 + g_V^2} \right], \quad M_Z^2 \simeq \frac{1}{2}(g_R^2 + g_V^2)(v_{RQ}^2 + v_{Rl}^2),
$$

(10)

with the gauge couplings satisfying the following relation

$$
g_V = \frac{g_Lg_V}{\sqrt{g_R^2 + g_V^2}}.
$$

(11)

Quite clearly, in the model the $Z_R$ is heavier than the $W_R$ and therefore a strong limit on the $W_R$ mass from experiments would mean an indirect bound exists on the $Z_R$ gauge boson too.

A. Fermion masses and mixings

We now look at the mass of the matter fields in the model. The gauge invariant Yukawa Lagrangian respecting the additionally imposed $Z_2$ symmetry in this model is given as:

$$
\mathcal{L}_Y = \left( Y_{uL}Q_L\bar{H}_{LQ}U_R + Y_{uR}\bar{Q}_R\bar{H}_{RQ}U_L + Y_{dL}Q_L\bar{H}_{LQ}D_R + Y_{dR}\bar{Q}_R\bar{H}_{RQ}D_L \\
+ Y_{eL}\bar{\bar{\nu}}_{LL}N_R + Y_{eR}\bar{\nu}_{RL}\bar{H}_{Rl}N_L + Y_{eL}\bar{\bar{\nu}}_{LL}E_R + Y_{eR}\bar{\nu}_{RL}\bar{H}_{Rl}E_L + M_U\bar{U}_LU_R + M_D\bar{D}_LD_R \\
+ M_E\bar{E}_LE_R + M_N\bar{N}_LN_R + H.C. \right) + M_{LN}N_LN_L + M_{RR}N_RN_R
$$

(12)

where $Y_{iA}$’s are the Yukawa couplings matrices and $M_X$’s are the singlet mass terms allowed by gauge symmetry. The conjugated scalar fields are defined as

$$
\bar{H}_{L/R} = i\tau_2 H^*_L/R.
$$

(13)

It is easy to see that the quark and charged lepton mass matrices will consist of off-diagonal terms proportional to left and right-handed VEVs while diagonal terms exist for only the heavy fields. Thus the quark and charged lepton mass matrices would be very similar in form to the Type-I seesaw neutrino mass matrix and all fermions have the same mechanism of mass generation in this framework.
1. Quarks

The quark masses in this model are obtained by diagonalizing a $6 \times 6$ mass matrix quite similar to what happens in seesaw mechanism. The up quark mass terms in this model can be written as

$$\mathcal{L}_u = (\bar{u} \, U) \left( M_u P_L + M_u^T P_R \right) \begin{pmatrix} u \end{pmatrix},$$

where

$$M_u = \begin{pmatrix} 0 & Y_u^{vR} & Y_u^{vLQ} \\ Y_u^{vLQ}^T & M_U \end{pmatrix}$$

is the $6 \times 6$ up quark mass matrix with $Y_u^{vL}, Y_u^{vR}$ and $M_U$ are all $3 \times 3$ matrices. The first $3 \times 3$ block corresponding to the light up-type quark is zero due to the absence of a bidoublet field in the scalar spectrum. The off-diagonal terms are obtained from the $Y_u^{vL}$ and $Y_u^{vR}$ terms of Eqn. [12] which involve the mixing of the light and heavy states through the Higgs doublet field, while the $M_U$ matrix is the mass term for the heavy up-type quarks. For simplicity we will choose all the Yukawa and heavy mass matrices to be diagonal in the up sector. This would mean that the CKM mixings will be generated entirely from the down sector which is exactly what we do for SM.

Similarly the down-type quark mass matrix can be written as

$$M_d = \begin{pmatrix} 0 & Y_d^{vR} \\ Y_d^{vR}^T & M_D \end{pmatrix},$$

where the first $3 \times 3$ is again zero due to the absence of a bidoublet scalar, the off-diagonal blocks arise from the Yukawa couplings and the $M_D$ term is the mass term for the heavy down-type quarks. Again in the down sector we keep the right-handed $3 \times 3$ Yukawa matrix $Y_d^{vR}$ and the $M_D$ matrix to be diagonal while only the left-handed Yukawa matrix $Y_d^{vL}$ is non-diagonal and sufficient to generate the correct CKM mixings for the SM quarks.

To diagonalize these non-symmetric matrices we require bi-unitary transformations. For the up-type quark mass matrix we have

$$M_u^{\text{diag}} = U_{uL} M_u U_{uR}^\dagger,$$

where $U_{uL}$ and $U_{uR}$ are the left and right-handed rotation matrices respectively. Similarly for the down sector

$$M_d^{\text{diag}} = U_{dL} M_d U_{dR}^\dagger.$$  

As can be easily seen that we will get two CKM mixing matrices in this case – a left-handed and a right-handed which are given as

$$U_L^{\text{CKM}} = U_{uL}^\dagger U_{dL}$$

and

$$U_R^{\text{CKM}} = U_{uR}^\dagger U_{dR}$$

respectively. These will be $6 \times 6$ matrices whose top-left (bottom-right) $3 \times 3$ block will correspond to the light CKM mixings for ascending (descending) arrangement of eigenvalues by mass. For our choice of parameters with only the left-handed Yukawa being non-diagonal, the right-handed CKM matrix would be almost diagonal with the mixings being quite small while the left-handed CKM mixings must be the same as the experimentally measured values.

The mixing between the heavy singlet quarks and the SM quarks are determined by the magnitude of the Yukawa terms in comparison to the singlet mass terms. For light quarks and even for the b quark, the Yukawa terms are much smaller than the bare mass term and hence the mixing is very small. For the top quarks though, because of its heavy mass compared to the other SM quarks, the mixings can be quite significant.
2. Charged Lepton

The charged lepton mass matrix is given as

\[ M_e = \begin{pmatrix} 0 & Y_{eR} v_{RI} \\ Y_{eL} v_{LI}^T & M_E \end{pmatrix}. \]  

(21)

This is very similar to the quark mass matrix with \( Y_{eL} \) and \( Y_{eR} \) being the 3 \( \times \) 3 Yukawa matrices while \( M_E \) is the heavy lepton mass matrix. Here we will choose all the matrices to be diagonal to prevent charged lepton flavor violation at the tree level. The mixing between the heavy singlet leptons and the SM charged leptons are almost negligible due to the hierarchical structure of the diagonal and the off-diagonal elements required for generation of correct lepton masses.

3. Neutrino

The neutrino matrix, on the other hand, would be quite different due to the Majorana-like \( M_L \) and \( M_R \) terms that could be written for the heavy neutrino states. The neutrino mass matrix in the basis \( (\nu^* L, N_R, \nu_R, N_L^*) \) is given as

\[ \begin{pmatrix} 0 & Y_{\nu L} v_{LI} & 0 & 0 \\ Y_{\nu L} v_{LI}^T & M_R & 0 & M_N^T \\ 0 & M_N & Y_{\nu R} v_{RI} & M_L \\ 0 & 0 & M_N^T & M_L \end{pmatrix}. \]  

(22)

Thus we see that all the fermion masses in this model arise from seesaw-like mass generation mechanism. Hence this model is also popularly known as the universal seesaw model. It is worth noting here that the neutrino mass matrix is actually symmetric (if the Yukawa and mass heavy mass matrices are symmetric) and can be diagonalized by a simple unitary transformation.

The neutrino mass matrix allows for a number of very unique scenarios in the neutrino sector. Firstly, there is the possibility that only the three left-handed doublet neutrinos are light and everything else is heavy. We explore such a scenario where we get three light neutrinos, three others with mass of electroweak scale and the rest at the TeV scale. We will refer to this scenario as Majorana case for obvious reasons. Similar to the previous cases we will again choose most of the Yukawa and mass matrices here to be diagonal except \( Y_{\nu L} \) which is chosen to be a non-diagonal symmetric matrix to explain the experimentally observed Pontecorvo-Maki-Nakagawa-Sakata (PMNS) neutrino mixing matrix elements. The mixing between the light and the heavy states are again quite small here leading to no significant limits from experimental observations. The heavy singlet states though can mix among themselves due to the presence of the \( M_N \) term but due to their masses being at the TeV scale, no observable effects have been discovered so far.

The second case that we can get is when the singlet neutrino Dirac mass term \( M_N \) is zero. In this case the neutrino mass matrix becomes block-diagonal with \( (\nu_L, N_R, \nu_R, N_L) \) bases being the diagonal blocks. In this case we get pseudo-Dirac like states with both the left and right doublet neutrinos being degenerate and light while the heavy singlet states may or may not be degenerate depending upon the choice of parameters. We refer to this scenario as pseudo-Dirac case. If we choose both \( M_L \) and \( M_R \) to be diagonal and equal we are forced to choose both \( Y_{\nu L} \) and \( Y_{\nu R} \) to be non-diagonal. Furthermore, in order to get equal masses for the now light left-handed and right-handed doublet lepton neutral components (neutrinos) we get the condition

\[ Y_{\nu Rij} = \frac{v_{LI}}{v_{RI}} Y_{\nu Lij} \]  

(23)

given we take \( M_L = M_R \). The mixing between the light and heavy states in each block diagonal sub matrix are still very small due to the fact that the Yukawa terms are now extremely small compared to the mass terms required to generate the light neutrino masses. The heavy states in this case do not mix as the \( M_N \) term is also absent.
B. Scalar masses and mixing

The full gauge invariant scalar potential for our model is given as:

\[
V(H) = \sum_{i=1}^{4} \mu_{ii} H_i^4 + \sum_{i,j=1 \atop ij}^{4 \times 4} l_{ij} H_i^3 H_j + \left[ \alpha_1 H_{LQ}^* H_{L} H_{RQ}^* H_{R}^* + \alpha_2 H_{LQ}^* H_{L} H_{R}^* H_{RQ}^* \right] \\
+ \mu_{12}^2 H_{LQ}^* H_{L} + \mu_{34}^2 H_{RQ}^* H_{R} + H.C.
\]

where

\[
H_1 = H_{LQ}, \quad H_2 = H_{L}, \quad H_3 = H_{RQ}, \quad H_4 = H_{R}.
\]

The last two terms involving \( \mu_{12} \) and \( \mu_{34} \) are responsible for breaking the discrete \( Z_2 \) symmetry softly without introducing any domain walls which could otherwise destabilize the model. We minimize this potential and get the following conditions

\[
\mu_{11} = \frac{1}{v_{L}} \left( \alpha_1 v_{Li} v_{Ri} v_{RQ} + \alpha_2 v_{Li} v_{Ri} v_{RQ} + 2l_{11} v_{LQ}^2 + l_{12} v_{Li}^2 v_{LQ} + l_{13} v_{LQ} v_{RQ} + l_{14} v_{LQ} v_{RQ}^2 - \mu_{12}^2 v_{Li} \right) \\
\mu_{22} = \frac{1}{v_{Li}} \left( \alpha_1 v_{Li} v_{Ri} v_{RQ} + \alpha_2 v_{Li} v_{Ri} v_{RQ} + 2l_{22} v_{Li}^2 v_{LQ} + l_{23} v_{Li} v_{RQ} + l_{24} v_{Li} v_{RQ}^2 - \mu_{12}^2 v_{Li} \right) \\
\mu_{33} = \frac{1}{v_{RQ}} \left( \alpha_1 v_{Li} v_{Ri} v_{RQ} + \alpha_2 v_{Li} v_{Ri} v_{RQ} + 2l_{33} v_{Li} v_{RQ} + l_{34} v_{RQ}^2 - \mu_{34}^2 v_{Li} \right) \\
\mu_{44} = \frac{1}{v_{RQ}} \left( \alpha_1 v_{Li} v_{Ri} v_{RQ} + \alpha_2 v_{Li} v_{Ri} v_{RQ} + 2l_{44} v_{RQ}^2 - \mu_{34}^2 v_{Li} \right).
\]

The Higgs boson spectrum in this case is significantly large and consists of four CP-even states, two CP-odd states and two charged Higgs bosons. Two charged goldstone boson are eaten up by the \( W_L \) and \( W_R \) gauge boson to give them mass while two neutral goldstone states give mass to the \( Z_L \) and \( Z_R \). The charged Higgs mass-squared matrix in this case is a \( 4 \times 4 \) block diagonal matrix with two blocks of \( 2 \times 2 \). In the basis \((H_{LQ}^+, H_{L}, H_{RQ}^+, H_{R}^+)\) the charged Higgs mass-squared matrix is given as:

\[
\begin{pmatrix}
\frac{v_{L}^2}{v_{LQ}} \left( \mu_{12}^2 - \alpha_{12}^2 v_{Ri} v_{RQ} \right) & \frac{v_{L}^2}{v_{LQ}} \left( \mu_{12}^2 - \alpha_{12}^2 v_{Ri} v_{RQ} \right) & 0 & 0 \\
\frac{v_{L}^2}{v_{LQ}} \left( \mu_{12}^2 - \alpha_{12}^2 v_{Ri} v_{RQ} \right) & \frac{v_{RQ}}{v_{L}} \left( \mu_{34}^2 - \alpha_{12}^2 v_{Ri} v_{LQ} \right) & 0 & 0 \\
0 & 0 & \frac{v_{RQ}}{v_{Li}} \left( \mu_{34}^2 - \alpha_{12}^2 v_{Ri} v_{LQ} \right) & 0 \\
0 & 0 & 0 & \frac{v_{RQ}}{v_{Li}} \left( \mu_{34}^2 - \alpha_{12}^2 v_{Ri} v_{LQ} \right)
\end{pmatrix}
\]

where \( \alpha_{12}^2 = (\alpha_1 + \alpha_2) \). Diagonalizing this matrix we get two goldstone states which are given as

\[
G_1^+ = \frac{1}{v_{L}^2 + v_{LQ}^2} (v_{LQ}, v_{Li}, 0, 0)^T, \quad G_2^+ = \frac{1}{v_{RQ}^2 + v_{Ri}^2} (0, 0, v_{RQ}, v_{Ri})^T.
\]

The two physical charged Higgs boson masses are

\[
m_{H_1^+}^2 = \frac{v_{Ri}^2 + v_{RQ}^2}{v_{Ri} v_{RQ}} \left( \mu_{34}^2 - \alpha_{12}^2 v_{Ri} v_{LQ} \right), \quad m_{H_2^+}^2 = \frac{v_{Li}^2 + v_{LQ}^2}{v_{Li} v_{LQ}} \left( \mu_{12}^2 - \alpha_{12}^2 v_{Ri} v_{RQ} \right),
\]

with the eigenstates being

\[
H_1^+ = \frac{1}{v_{Ri}^2 + v_{RQ}^2} (0, 0, -v_{Ri}, v_{RQ})^T, \quad H_2^+ = \frac{1}{v_{Li}^2 + v_{LQ}^2} (v_{LQ}, v_{Li}, 0, 0)^T.
\]

It is easy to see here that if we choose \( \mu_{12}^2 = \mu_{34}^2 \approx v_{EW}^2 \) then the right-handed charged Higgs boson is indeed the lightest state. This is due to the fact that the left-handed charged state has an additional enhancement of
$v_{LQ}/v_{LI}$ except for a very fine-tuned region around

$$\alpha_1 + \alpha_2 \approx \frac{\mu_{12}^2}{v_{RQ}v_{RQ}}.$$  

(31)

The CP-odd Higgs boson mass-squared matrix in the basis $(\text{Im}H_{LQ}^0, \text{Im}H_{LI}^0, \text{Im}H_{RQ}^0, \text{Im}H_{Rl}^0)$ is

$$
\begin{pmatrix}
\frac{\mu_{12}}{v_{LQ}}\left(\mu_{12}^2 - \alpha_{12}v_{RL}v_{RQ}\right) & \alpha_{12}v_{RL}v_{RQ} - \mu_{12}^2 & v_{LI}v_{RL}(\alpha_2 - \alpha_1) & v_{LI}v_{RL}(\alpha_2 - \alpha_2) \\
\alpha_{12}v_{RL}v_{RQ} - \mu_{12}^2 & \frac{\mu_{12}^2}{v_{LQ}} & \alpha_{12}v_{RL}v_{RQ} - \mu_{12}^2 & v_{LQ}v_{RL}(\alpha_1 - \alpha_2) \\
v_{LQ}v_{RL}(\alpha_2 - \alpha_1) & \alpha_{12}v_{RL}v_{RQ} - \mu_{12}^2 & \frac{\mu_{12}^2}{v_{LQ}} & \alpha_{12}v_{LQ}v_{LQ} - \mu_{12}^2 \\
v_{LQ}v_{RL}(\alpha_2 - \alpha_2) & \alpha_{12}v_{LQ}v_{LQ} - \mu_{12}^2 & \alpha_{12}v_{LQ}v_{LQ} - \mu_{12}^2 & \frac{\mu_{12}^2}{v_{LQ}}
\end{pmatrix}.
$$

(32)

This again will have two zero eigenstates corresponding to the two goldstone bosons required for $Z$ and $Z_R$ mass generation. It is also to see here that in the case where $\alpha_1 = \alpha_2$ this CP-odd mass-squared matrix would reduce to the block diagonal charged Higgs boson mass-squared matrix.

The CP-even scalar Higgs boson mass-squared matrix elements in the basis $(\text{Re}H_{LQ}^0, \text{Re}H_{LI}^0, \text{Re}H_{RQ}^0, \text{Re}H_{Rl}^0)$ can be expressed in terms of the CP-odd Higgs mass-squared matrix elements as

$$M_{ij,\text{CP-Even}}^2 = M_{ij,\text{CP-Odd}}^2 + 2S_{ij}v_{LQ}v_{LQ},$$

(33)

where $i, j = 1, 2, 3, 4$, $v_1 = v_{LQ}$, $v_2 = v_{LI}$, $v_3 = v_{RQ}$, $v_4 = v_{Rl}$ and

$$S_{ij} = \begin{cases} 2, & \text{if } i = j \\ 1, & \text{otherwise}. \end{cases}$$

(34)

We choose our parameters such that the lightest eigenvalue of this CP-even Higgs mass-squared matrix is the one corresponding to the SM-like Higgs with mass of 125 GeV, while all the other states are chosen to be much heavier. The lightest charged and pseudo-scalar Higgs boson masses come out to be around a few 100 GeV while the heavier ones are around a few TeV. In Table II we give a list of the physical Higgs boson masses and the respective eigenstates for a sample benchmark point. Note that unlike the case of 2HDM models, here only the pseudoscalar and charged Higgs bosons are light while all CP even scalars turn out to be very heavy. In addition both the light pseudoscalar and charged scalar are admixtures of the right sector scalar doublets.

| Particle | Mass (GeV) | Eigenstate |
|----------|------------|------------|
| $H_1$    | 125.5      | 0.996 Re($H_{LQ}^0$) - 0.010 Re($H_{RQ}^0$) + 0.080 Re($H_{Rl}^0$) + 0.027 Re($H_{RL}^0$) |
| $H_2$    | 2543.9     | -0.0289 Re($H_{LQ}^0$) - 0.381 Re($H_{RQ}^0$) + 0.001 Re($H_{RL}^0$) + 0.924 Re($H_{RI}^0$) |
| $H_3$    | 4229.0     | 0.001 Re($H_{LQ}^0$) - 0.924 Re($H_{RQ}^0$) + 0.005 Re($H_{RL}^0$) - 0.381 Re($H_{RI}^0$) |
| $H_4$    | 7127.5     | 0.080 Re($H_{LQ}^0$) - 0.005 Re($H_{RQ}^0$) - 0.997 Re($H_{RL}^0$) + 0.003 Re($H_{RI}^0$) |
| $A_1$    | 217.0      | 0.001 Im($H_{LQ}^0$) - 0.707 Im($H_{RQ}^0$) - 0.008 Im($H_{RL}^0$) + 0.707 Im($H_{RI}^0$) |
| $A_2$    | 7127.7     | -0.080 Im($H_{LQ}^0$) - 0.006 Im($H_{RQ}^0$) + 0.997 Im($H_{RL}^0$) + 0.006 Im($H_{RI}^0$) |
| $H_1^+$  | 225.8      | -0.707$H_{RQ}^0$ + 0.707$H_{RL}^0$ |
| $H_2^+$  | 7127.5     | 0.080$H_{LQ}^0$ - 0.997$H_{RI}^0$ |

TABLE II: Scalar Eigenstates for $\alpha_1 = -0.3$, $\alpha_2 = 0.1$, $l_{11} = 0.172$, $l_{12} = 0.8$, $l_{13} = 0.05$, $l_{14} = -0.1$, $l_{22} = 0.5$, $l_{23} = 0.1$, $l_{24} = 0.1$, $l_{33} = 0.2$, $l_{34} = 0.1$, $l_{44} = 0.1$, $\mu_{12}^2 = 2.5 \times 10^4$, $\mu_{34}^2 = 2.5 \times 10^4$

III. PHENOMENOLOGICAL IMPLICATIONS

It is worth pointing out here that to generate the model and study the spectrum we have implemented the model into SARAH and use the generated SPHENO file to analyse the spectrum and decay probabilities of
the various particles of the model.

We now look at the phenomenological implications of this model: the allowed parameters, experimental constraints and unique signals of this model which may be studied at the colliders. As has been discussed before almost all of the matrices are taken to be diagonal except $Y_{uL}$ and the neutrino left-handed Yukawa matrices. These are necessarily off-diagonal in order to generate the CKM and PMNS mixing. Unlike SM where the Yukawa couplings can range from $10^{-6}$ to 1, this model requires a much smaller range of Yukawa couplings ranging from $10^{-3}$ to 1 for all the charged particles. In general we have chosen

$$Y_{uL}^{11} \sim Y_{dL}^{11} \sim Y_{eL}^{11} \approx 10^{-2}, \quad Y_{uL}^{22} \sim Y_{dL}^{22} \sim Y_{eL}^{22} \approx 10^{-1}, \quad Y_{uL}^{33} \sim Y_{dL}^{33} \sim Y_{eL}^{33} \approx 1,$$

while the other elements in the $Y_{dL}$ matrix are of the order of $10^{-3}$. We further choose $Y_{dR}^{33} = 0.023$ and $Y_{eL}^{33} = 0.26$ so that the third generation heavy fermions are all of the order of a few TeV. With this kind of a Yukawa structure we can easily get the correct masses of all the fermions by choosing appropriate values of the heavy masses. The left-handed CKM matrix mixing elements are obtained entirely in the down sector similar to the SM, while the right-handed down quark mixings are very small due to the diagonal structure of the $Y_{dR}$ matrix. The vacuum values of the fields are taken to be

$$v_{RQ} = v_{RI} = 4.5 \text{ TeV}, \quad v_{LQ} = 173.4 \text{ GeV}, \quad v_{Ll} = 14 \text{ GeV}.$$  

The bare masses of the singlet heavy vector-like fermions are chosen accordingly so as to get the correct masses for the SM-like fermions. Here we note that since the third generation fermions are the heaviest followed by the second generation and then the first generation fermion masses, the reverse order is generally followed by the vector-like singlet fermion masses. For each type of fermions (up quark, down quark and charged leptons), our choices are such that third generation vector-like fermions are indeed lightest. This can be understood easily as in the seesaw formula the mass of the light state is inversely proportional to the heavy mass in the seesaw matrix for the same value of off-diagonal terms. Though it is not strictly valid for this case as the off-diagonal Dirac masses are also higher for the third generation, we choose our Yukawa couplings so that the third generation vector fermions are indeed lightest. Table III lists the masses of all the new fermions in our model and it can be seen that indeed the third generation-like vector fermions are lightest in each case. With the strong sector exotic quarks and charged leptons having mass above $3.5 \text{ TeV}$ it would be quite impossible to observe any signals for these fermions at the current LHC energies. However they could be more copiously produced at future $100 \text{ TeV}$ machines such as the FCC-hh collider [10].

| Up-type Quark | Down-type Quark | Charged Lepton | Neutrino |
|---------------|----------------|---------------|----------|
| $M_T = 4.51 \text{ TeV}$ | $M_B = 3.97 \text{ TeV}$ | $M_{E_4} = 6.13 \text{ TeV}$ | $M_{\nu_4} = 136 \text{ GeV}$, $M_{\nu_5} = 258 \text{ GeV}$, $M_{\nu_6} = 317 \text{ GeV}$, $M_{\nu_7} = 9.07 \text{ TeV}$, $M_{\nu_8} = 9.13 \text{ TeV}$, $M_{\nu_9} = 9.16 \text{ TeV}$, $M_{\nu_{10}} = 11.06 \text{ TeV}$, $M_{\nu_{11}} = 11.1 \text{ TeV}$, $M_{\nu_{12}} = 11.2 \text{ TeV}$ |
| $M_C = 6.17 \text{ TeV}$ | $M_S = 10.4 \text{ TeV}$ | $M_{E_2} = 9.92 \text{ TeV}$ | $M_{\nu_4} = 200.0 \text{ GeV}$, $M_{\nu_5} = 300.0 \text{ GeV}$, $M_{\nu_6} = 400.0 \text{ GeV}$ |
| $M_U = 30.0 \text{ TeV}$ | $M_D = 17.2 \text{ TeV}$ | $M_{E_1} = 12.3 \text{ TeV}$ | $M_{\nu_4} = 136 \text{ GeV}$, $M_{\nu_5} = 258 \text{ GeV}$, $M_{\nu_6} = 317 \text{ GeV}$, $M_{\nu_7} = 9.07 \text{ TeV}$, $M_{\nu_8} = 9.13 \text{ TeV}$, $M_{\nu_9} = 9.16 \text{ TeV}$, $M_{\nu_{10}} = 11.06 \text{ TeV}$, $M_{\nu_{11}} = 11.1 \text{ TeV}$, $M_{\nu_{12}} = 11.2 \text{ TeV}$ |

**TABLE III:** Fermion masses

The mixing between the heavy singlet-like states and the light SM-like states are very low ($\lesssim 1\%$) except for the top sector which behaves quite differently. To get the correct top quark mass we need to take $M_u$ to be quite small to be around $350 \text{ GeV}$. The heavy top partner mass almost entirely comes from the right handed top quark contribution and hence the right-handed CKM mixing of the top quark is almost entirely coming from the heavy singlet top partner. Thus the decay of the heavy gauge bosons which belong to the $SU(2)_R$ do not couple with the same strength to the third generation SM quarks as they do to the first two generations. This effect will quite clearly be seen in the $W_R$ decay modes where its branching ratio $W_R^+ \rightarrow t\bar{b}$ is significantly suppressed ($\sim 0.2\%$) while for the other generation light quarks its around $33\%$ each. This in turn would make the bounds on the $W_R$ gauge boson much stronger from existing dijet data than that of conventional LRS models which have slightly lower branchings into light jets. We show the $W_R$ decay channels in Fig. [1].
FIG. 1: Branching ratio for $W_R$ boson as a function of its mass.

| $\Delta m_{21}^2$ | $\Delta m_{31}^2$ | $\sin^2 \theta_{12}$ | $\sin^2 \theta_{23}$ | $\sin^2 \theta_{13}$ |
|-------------------|-------------------|----------------------|----------------------|----------------------|
| $7.03 \times 10^{-5}$ eV$^2$ | $8.09 \times 10^{-5}$ eV$^2$ | $0.271$ | $0.385$ | $0.01934$ |
| $2.407 \times 10^{-3}$ eV$^2$ | $2.643 \times 10^{-3}$ eV$^2$ | $0.345$ | $0.635$ | $0.02392$ |

$U_{PMNS}$

| $Y_{\nu L}$ | $Y_{\nu R}$ | $Y_{\nu \tau}$ |
|-------------|-------------|-------------|
| $0.800 \to 0.844$ | $0.515 \to 0.581$ | $0.139 \to 0.155$ |
| $0.229 \to 0.516$ | $0.438 \to 0.699$ | $0.614 \to 0.790$ |
| $0.249 \to 0.528$ | $0.462 \to 0.715$ | $0.595 \to 0.776$ |

TABLE IV: Experimental $3\sigma$ ranges for light neutrino parameters

In the neutrino sector there are two specific scenarios (Majorana and pseudo-Dirac) as has been discussed earlier. Here we will only discuss the case of normal hierarchy for the neutrino masses. For the Majorana case, we fit the experimental data for neutrino oscillation parameters \[11\] with the variations being within the $3\sigma$ ranges of the respective central values obtained in the global fits. We have listed the values used for the fit in Table IV while scanning the parameter space of our model. We choose all the matrices to be diagonal except $Y_{\nu L}$ which is a symmetric matrix. We choose

$$M_R = M_L = \text{Diag} \left( 10^4, 10^4, 10^4 \right), \quad M_N = \text{Diag} \left( 10^3, 10^3, 10^3 \right),$$

while the elements of the Yukawa coupling matrix $Y_{\nu R}$, $Y_{\nu L}$, and $Y_{\nu \tau}$ are $10^{-5}$. This choice gives us our desired masses with the three light neutrino masses around 0.001 to 0.01 eV, the next three states being around a few 100 GeV and the rest being around 10 TeV. The three light neutrino physical states are almost entirely from the three generations of $\nu_L$, the next three lighter ones (masses of a few 100 GeV) are mostly from $\nu_R$. This is very similar to Type-I seesaw in conventional LRS models and would give very similar phenomenology.

---

\[2\] It is worth noting that an arrangement for inverted hierarchy of the neutrino masses is equally possible in our model, which we have not considered here.
with same-sign lepton signals, neutrinoless double-beta decays. However, the modified scalar sector interaction with the heavy neutrinos lead to much different collider signals which we shall discuss later. The much heavier eigenstates which would be beyond the reach of current accelerator energies are a mixture of \(N_L\) and \(N_R\). In Fig. 2(a) we show the allowed parameters which gives us the correct neutrino mass-squared differences and the correct PMNS mixing angles for normal hierarchy. We can see that \(Y_{\nu_{13}}\) is indeed the largest owing to \(\nu_3\) being the largest in normal hierarchy case, while \(Y_{\nu_{13}}\) is the smallest in magnitude as required to explain the small value of the mixing \(\theta_{13}\).

Similarly for this case the \(12 \times 12\) neutrino mass matrix is symmetric and hence it can be diagonalized with a simple unitary transformation. The first \(3 \times 3\) block corresponds to the three light neutrinos and satisfies the experimental \(3\sigma\) bounds of the PMNS matrix. The other mixing of the light neutrinos with the heavier ones are extremely small with \(\sin \theta_{ij} \lesssim 10^{-8}\) where \(\theta_{ij}\) is the mixing angle between the heavy and light states. Hence there are no bound coming from lepton flavor changing processes. The three mass eigenstates with masses of a few 100 GeV are almost the same as the flavor eigenstates of \(\nu_R\) with small mixing (\(\sim 1\%\)) with \(N_L\). The six heavier states of masses around 10 TeV almost equally constitute of \(N_{L_i}\) and \(N_{R_i}\), where \(i = 1, 2, 3\). Table III gives a list of all the neutrino masses in this scenario for a particular benchmark point.

In the pseudo-Dirac neutrino case the singlet neutrino mixing term \(M_N\) is taken to be zero. Then we choose our parameters as \(M_R = M_L = \text{Diag}(200, 300, 400)\). To get the correct light neutrino masses and mixings for this choice of \(M_L\) and \(M_R\) we are forced to choose both \(Y_{\nu_L}\) and \(Y_{\nu_R}\) to have non-zero off-diagonal elements but they can still be chosen to be symmetric matrices. Now we get two block diagonal \(6 \times 6\) symmetric matrices each of whose three light eigenvalues should be equal and satisfy the experimentally observed mass-squared differences for them to form pseudo-Dirac-like states. Using Eqn. 23 along with the observed mass-squared differences and mixing constraints lead to the following choice of the matrix elements \(Y_{\nu_{L,ij}} \sim 10^{-6}\) and \(Y_{\nu_{R,ij}} \sim 10^{-9}\). The neutrino mixing in this case again have to satisfy the experimental PMNS mixing limits and have to be the same for both the sectors which is easily satisfied by using the condition given in eq. 23. Fig. 2(b) gives some allowed parameters for this case which satisfy the neutrino mass-squared differences and the mixing angles for normal hierarchy. Like the previous case again we can see that \(Y_{\nu_{13}}\) is usually the largest while the elements for \(Y_{\nu_{13}}\) are the smallest for most of the points.

The mixing between the states of \(\nu_L\) and \(\nu_R\) which are now of equal masses in this scenario are not quite as small as the previous case here. The mixing angle \(\theta\) between two light states of equal masses are typically such that \(\sin \theta \sim 10^{-2}\). The heavy states again have negligibly small mixing with the light states like in the previous case. Here since we have taken \(M_N = 0\) there is no mixing between the heavy states and they are purely consisting of \(N_{L_i}\) or \(N_{R_i}\). As a result their decay widths in this case becomes very small, which can lead to observable displaced vertex signals. We have listed the sub-TeV heavy Majorana neutrino decay widths and the most important decay channels in Tab. IV which clearly shows the very small decay widths of many of the heavy neutrinos which mostly decay into three body final states which are charged and may be observed in the experiments.
12

| Particle | Width (GeV) | Important Decay Channels |
|----------|-------------|--------------------------|
| $\nu_4$  | $7.12 \times 10^{-9}$ | $e^+\pi d (25.1\%), \ e u \bar{d} (25.1\%), \ e^+ \tau s (24.5\%), \ e c \bar{\pi} (24.5\%)$ |
| $\nu_5$  | $4.57 \times 10^{-4}$ | $\mu H^+_1 (50\%), \ \mu^+ H^-_1 (50\%)$ |
| $\nu_6$  | $2.97 \times 10^{-3}$ | $\tau H^+_1 (50\%), \ \tau^+ H^-_1 (50\%)$ |

TABLE V: Heavy Neutrino Decays

A. Experimental Constraints

The scalar sector of the model discussed in this paper may be considered as a left-right extension of the lepton specific two Higgs doublet model (2HDM). As such there are a number of flavor constraints which restrict the parameter space of the 2HDM. Most stringent of these constraints come from the $b \to s \gamma$ process and constrains the charged Higgs mass such that $m_{H^\pm} > 460 \text{ GeV}$ [12] in Type III 2HDM. The main process responsible for $b \to s \gamma$ in 2HDM is given in fig 3. In our case though this process is present, the lighter charged Higgs boson of mass around 200 GeV actually corresponds to the right-handed charged Higgs boson as can be seen from eq. 30. As has been discussed before, the right-handed down type Yukawa coupling matrix is diagonal in this model and hence the CKM mixings are really small. This results in a much weaker bound on the lightest charged Higgs boson mass in this case. There are no significant bounds from the flavor observable on the pseudoscalar masses in the lepton-specific 2HDM and hence there are no bounds in this model as well.

![FIG. 3: $b \to s \gamma$ through charged Higgs](image)

There is also bound on the 2HDM pseudoscalar Higgs boson mass from the single production and associated production of the pseudoscalar decaying into two $\tau$ final state [13]. This gives a lower limit on the pseudoscalar mass as a function of the $\sigma \times BR(A \to \tau\tau)$ for both the single and the associated production mode. For both these production channels the important couplings would be $A_1 q\bar{q}$ where $q$ is a quark. To get significant production, the third generation quarks are the most important but here again the couplings of the pseudoscalar with the third generation quarks will be much weaker than in the case of 2HDM. This is because the coupling here would be

$$f^L_{A_1}(q\bar{q}) = \frac{\gamma^{33}_{uL(dL)} \times Z^L_{A_1} \times Z^Q q_{uR(dR)}}{m_{A_1}^2}$$

(38)

where $Z^L_{A_1}$ is the amount of $H^{0}_{LQ}$ contained in the eigenstate of $A$ and $Z^Q q_{uR(dR)}$ is the mixing of the right-handed heavy and the light up-type (down-type) quarks. Similar formula can be written for the right-handed pseudoscalar coupling with two quarks with $L \leftrightarrow R$ in Eq. [38]. The light pseudoscalar in this model is coming from the right-handed doublets and its couplings with the third-generation quarks come out to be much weaker than the 2HDM case. So for our model this limit will not be applicable because the production cross-section of the pseudoscalar will be much smaller than in 2HDM.

B. New Collider Signals

This model can lead to a number of interesting new signals at accelerator experiments. We primarily focus in mentioning the ones from the scalar sector in the form of charged Higgs as well as the heavy Majorana neutrinos which can be accessible to the current run of the Large Hadron Collider (LHC). Note that the other exotics
such as the heavy quarks and leptons are beyond the reach of LHC because of their extremely heavy masses.

Unlike other models for heavy neutrinos including left-right symmetric models we find that the decay modes of the heavy Majorana neutrinos, as listed in Table [V] are quite different. Note that the heavy neutrino decays are again driven by their composition and therefore could be either singlet dominated or even SU(2)R doublet dominated. As the Yukawa couplings $Y_{\nu_R i} \sim 0.1$ we find that the dominantly right-sector charged Higgs which is light would couple to the heavy Majorana neutrinos which have dominant right-handed components as well as singlet components over the left-handed one’s. This plays a crucial role in deciding the decay of the heavy neutrinos in the model. The heavy neutrinos prefer to decay via the off-shell charged Higgs while the subleading contributions come from the decay via off-shell $W_R$. Although both the mediating particles would contribute, a quick look at the decay probabilities in Table [V] shows the absence of the leptonic modes which are highly suppressed. This indicates that the decay is driven by the off-shell charged Higgs over the much heavier $W_R$ gauge boson. The challenge for observing signals for these heavy neutrinos would be dictated by the production mechanism. At LHC, it would mean that they could be produced via exchange of $W_R$ and $Z_R$ in the $s$-channel. This would give a resonant production of the heavy neutrinos and therefore the dominant channel. So we can produce the heavy neutrinos as

$$\bullet \quad pp \rightarrow W_R^+ \rightarrow \ell_i^+ \nu_j$$

$$\bullet \quad pp \rightarrow Z_R \rightarrow \nu_j \nu_j$$

where $i = 1, 2, 3$ and $j = i + 3$.

Using some generic parameters as discussed in the text earlier, we find that a $W_R$ of mass 3 TeV has a combined branching of nearly 30% to decay to a SM charged lepton and heavy neutrino, with the dominant mode of the three being the decay to the first 2 generations $\sim 11\%$. Similarly a $Z_R$ of mass around 3.5 TeV has a combined $\sim 20\%$ branching probability to decay in the pair of $\nu_4$, $\nu_5$, $\nu_6$ which have mass of 136, 258 and 317 GeV, respectively. Resonant production of heavy neutrinos can be useful to have appreciable rates of production [14] without depending on the active-sterile mixing parameter in the neutrino sector. As the decay in Table [V] suggests, the heavy neutrino decays to a single flavor charged lepton once produced in association with jets with a 100% branching probability, leading to same-sign dilepton signals with jets in the final state provided the charged Higgs in the model is heavier. Note that if the Majorana neutrinos are heavier than the charged Higgs which in turn is heavier than the top quark then a very unique and different signal is produced. With no heavy neutrino decay available to the charged Higgs, it decays to the quark final states with the dominant channel being $H^+ \rightarrow t\bar{b}$. Now as the charged Higgs comes from the decay of a heavy Majorana neutrino then one gets an interesting signal where one has same-sign leptons as well as same-sign top quark in the final state. This is completely free from any SM background and would be an unique signal for discovery. Thus we have for example

$$pp \rightarrow \nu_4\nu_4 \rightarrow e^- e^- H^+ H^+ \rightarrow 2e^- + 2t + 2\bar{b}$$

when the $\nu_4$ is pair produced. Again one gets $2\mu^- + 2t + 2\bar{b}$ when $\nu_5$ is pair produced. For the $H^+$ lighter than the top quark it decays to the light quarks thus giving a more conventional signal of same-sign leptons with multiple jets. However a marked difference is the absence of SM $W$ and $Z$ boson in the decay cascades of the heavy neutrino decay. These modes become available for the scenario with pseudo-Dirac heavy neutrinos. Similarly, the single production of the heavy neutrino through $W_R$ resonance would lead to a signal with same-sign lepton along with a top and bottom quark. We leave a much more detailed signal analysis of the collider signals of the model for future work and focus on pointing out the interesting signals that one can expect to observe at LHC here.

In addition, if the charged Higgs are heavier than the heavy neutrinos, then they would dominantly decay into them and the corresponding charged lepton. This can lead to significantly different search signal for the charged Higgs when compared to conventional ones. Thus even when the charged Higgs is heavier than the top quark, a presence of a light Majorana neutrino completely overwhelms the $t\bar{b}$ decay option. The charged Higgs production would be via the photon exchange mostly:

$$pp \rightarrow H^+ H^- \rightarrow \nu_j \nu_j \ell_i^+ \ell_i^-$$

where again $i = 1, 2, 3$ while $j = i + 3$. The heavy neutrino would decay via the off-shell charged Higgs in the 3-body decay channel $\nu_j \rightarrow \ell^\pm j j'$. This quite clearly gives a multi-lepton signature for the charged Higgs mediated by lepton-number violating interactions which again has very little or no SM background and can be a very unique signal of the model.
Another unique signal of this model which differentiates it from other LR models involves the $W_R$ decay channels. Fig. [1] gives a plot of the various $W_R$ decay branching ratios in this model. In general LR models an important decay channel for $W_R$ boson is a $t, b$ final state but that channel is almost absent in this case owing to the extremely small branching ratio as can be seen in Fig. [1] In fact once the heavy $T$ fermion channel opens up, a significant branching is into this $T, b$. Thus the model opens a possibility of some interesting signal topologies which are quite non-standard and can give surprisingly different and unique signals from production of heavy Majorana neutrino as well as charged Higgs boson at the LHC.

IV. CONCLUSION

In this work we have proposed a model for SM fermion mass generation through universal seesaw mechanism. The model is based on a left-right symmetric framework where all the gauge symmetries are spontaneously broken via $SU(2)$ scalar doublets only. The gauge symmetry is augmented with an additional $Z_2$ discrete symmetry which differentiates the quarks from the lepton sector. Additional heavy vector-like singlet fermions are needed for the generation of the SM quark and lepton masses through a universal seesaw mechanism. The neutrino matrix, on the other hand, can lead to two very interesting physical scenarios – one with Majorana-like neutrinos and the other where the neutrinos are pseudo-Dirac in nature.

The scalar sector here may be considered as a LRS extension of lepton-specific 2HDM. The SM-like neutral Higgs boson (with mass of 125 GeV) and the lightest charged and pseudoscalar Higgs states remain light of the order of a few hundred GeV. The most stringent bounds on the charged and the pseudoscalar Higgs masses in a general 2HDM scenario come from the flavor-changing processes and two $\tau$ final state decay modes. These bound are quite relaxed in this model due to the right-handed nature of both of these light scalars and their much reduced effective couplings in this model. Hence the light charged or pseudoscalar states can easily be accommodated here which can lead to interesting collider signatures.

The model also presents us with some unique collider signatures that could be observed at the LHC. In addition to interesting signal from heavy neutrino production where one gets same-sign leptons and same-sign top quark pair in the same event, the model also gives very unique and different signal for the charged Higgs in the model. The observation of such non-standard signal events at LHC could provide hints on new physics with an underlying model quite different from the popular left-right models.

Acknowledgments

This work was partially supported by funding available from the Department of Atomic Energy, Government of India, for the Regional Centre for Accelerator-based Particle Physics (RECAPP), Harish-Chandra Research Institute. AP was partially funded by the SERB National Postdoctoral fellowship file no. PDF/2016/000202. AP would like to thank RECAPP, Harish-Chandra Research Institute for their hospitality during the visit when a part of this work was done. AP would also like to thank Biplob Bhattacharjee for useful discussions. SKR would like to thank the CERN Theory Group for a short-term visit while this work was being completed and written up.

[1] R. N. Mohapatra and J. C. Pati, Phys. Rev. D 11, 566 (1975); G. Senjanovic and R. N. Mohapatra, Phys. Rev. D 12, 1502 (1975).
[2] M. A. B. Beg and H. -S. Tsao, Phys. Rev. Lett. 41, 278 (1978); R. N. Mohapatra and G. Senjanovic, Phys. Lett. B 79, 283 (1978); K. S. Babu and R. N. Mohapatra, Phys. Rev. D 41, 1286 (1990); S. M. Barr, D. Chang and G. Senjanovic, Phys. Rev. Lett. 67, 2765 (1991); R. N. Mohapatra and A. Rasin, Phys. Rev. Lett. 76, 3490 (1996); R. Kuchimanchi, Phys. Rev. Lett. 76, 3486 (1996); R. N. Mohapatra, A. Rasin and G. Senjanovic, Phys. Rev. Lett. 79, 4744 (1997); K. S. Babu, B. Dutta and R. N. Mohapatra, Phys. Rev. D 65, 016005 (2001); R. Kuchimanchi, Phys. Rev. D 82, 116008 (2010).
[3] R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. 38, 1440 (1977).
[4] P. Minkowski, Phys. Lett. B 67 (1977) 421; T. Yanagida, proceedings of the Workshop on Unified Theories and Baryon Number in the Universe, Tsukuba, 1979, eds; A. Sawada, A. Sugamoto, KEK Report No. 79-18, Tsukuba; S. Glashow, in Quarks and Leptons, Cargèse 1979, eds; M. Lévy. et al., (Plenum, 1980, New York); M. Gell-Mann, P. Ramond, R. Slansky, proceedings of the Supergravity Stony Brook Workshop, New York, 1979, eds. P. Van Nieuwenhuizen, D. Freeman (North-Holland, Amsterdam). R. Mohapatra, G. Senjanović, Phys.Rev.Lett. 44 (1980) 912.
