Lepton-specific universal seesaw model with left-right symmetry

Ayon Patra[1,*] and Santosh Kumar Rai[2,†]

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

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We propose a left-right symmetric framework with a universal seesaw mechanism for the generation of masses of the Standard Model quarks and leptons. Heavy vectorlike singlet quarks and leptons are required for generation of Standard-Model-like quark and lepton masses through a 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.

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I. INTRODUCTION

Left-right symmetric (LRS) models [1] are one of the most well-motivated and widely studied extensions of the Standard Model (SM). The popularity of LRS models stems from the fact that in these models it is possible to explain several phenomena which are not very well understood in the framework of the 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 Lagrangian. Since CP violating terms in the color sector are highly constrained from neutron electric dipole measurements, the absence of these terms can solve the strong CP problem [2] naturally without the need to introduce a global Peccei-Quinn symmetry [3]. The gauge structure of these models forces 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 [4].

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 Cabibbo-Kobayashi-Maskawa (CKM) mixings. This multitude of scalar fields makes the scalar sector quite complicated.[1] 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 the 2HDM and the other two their right-handed counterparts). The model we study here is a lepton-specific scenario where one pair of Higgs doublets couples 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 nonzero 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 unstable from a cosmological point of view [6]. 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

*ayon@okstate.edu
†skrai@hri.res.in

[1]A detailed study of various scalar sectors in LRS models is discussed in [5].

lepton rich final states that could be observed or excluded at the LHC.

All the fermion masses in this case are generated through a universal seesaw mechanism [7] 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.

The rest of the paper is organized as follows. In Sec. II, we discuss our model in detail and in Sec. III we discuss the phenomenological implications of our model including the experimental constraints and possible collider signatures. Section IV contains our conclusions and discussion.

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

$$Q = I_{3L} + I_{3R} + \frac{B-L}{2}. \tag{1}$$

The chiral matter fields consist of three families of quarks and leptons:

$$\begin{aligned} Q_L &= \begin{pmatrix} u \\ d \end{pmatrix}_L \sim \left(3, 2, 1, \frac{1}{3}\right), \\ Q_R &= \begin{pmatrix} u \\ d \end{pmatrix}_R \sim \left(3, 1, 2, \frac{1}{3}\right), \\ l_L &= \begin{pmatrix} \nu \\ e \end{pmatrix}_L \sim (1, 2, 1, -1), \\ l_R &= \begin{pmatrix} \nu \\ e \end{pmatrix}_R \sim (1, 1, 2, -1), \end{aligned} \tag{2}$$

where the numbers in the parentheses denote the quantum numbers under $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge groups respectively.

The scalar sector in this model does not contain any bidoublet fields and hence heavy singlet quarks and leptons are necessary for the generation of the quark and lepton masses through the seesawlike mechanism. We introduce heavy up- and down-type quarks given as $U_L(3, 1, 1, \frac{4}{3})$, $U_R(3, 1, 1, \frac{4}{3})$, and $D_L(3, 1, 1, -\frac{2}{3})$, $D_R(3, 1, 1, -\frac{2}{3})$ respectively. The heavy charged leptonic states are $E_L(1, 1, 1, -2)$ and $E_R(1, 1, 1, -2)$ while the heavy neutrino states are given as $N_L(1, 1, 1, 0)$ and $N_R(1, 1, 1, 0)$. It is worth noting that while all the $(B-L)$ charged heavy states can only have Dirac-like terms, the heavy neutrinos can admit both Dirac and Majorana-like terms. Of course one can still argue that the LRS models with triplet Higgs are also lepton specific as they do not couple to the quarks. However one must note that the scalar sector would then define a significantly different phenomenology from that of the standard triplet scenarios, in particular with the absence of a double charged scalar in the spectrum.

The minimal Higgs sector consists of the following fields:

$$\begin{aligned} H_{RQ}(1,1,2,1) &= \begin{pmatrix} H_{RQ}^+ \\ H_{RQ}^0 \end{pmatrix}, \qquad H_{LQ}(1,2,1,1) = \begin{pmatrix} H_{LQ}^+ \\ H_{LQ}^0 \end{pmatrix}, \\ H_{Rl}(1,1,2,1) &= \begin{pmatrix} H_{Rl}^+ \\ H_{Rl}^0 \end{pmatrix}, \qquad H_{Ll}(1,2,1,1) = \begin{pmatrix} H_{Ll}^+ \\ H_{Ll}^0 \end{pmatrix}, \end{aligned} \tag{3}$$

where H_{LQ} and H_{RQ} interact specifically with quarks while H_{Ll} and H_{Rl} only have leptonic interactions. The H_{RQ}^0 and H_{Rl}^0 get nonzero VEVs and are responsible for breaking the right-handed symmetry. The heavy W_R and Z_R gauge boson masses are generated at this scale. The VEVs of the H_{LQ}^0 and H_{Ll}^0 fields on the other hand are the ones responsible for the electroweak symmetry breaking and generation of the W and Z boson masses. Since H_{Rl}^0 and H_{Ll}^0, which are both odd under the Z_2 symmetry, get nonzero VEVs this could lead to formation of domain walls and destabilize the model. This problem is addressed by introducing soft Z_2-breaking terms in the scalar potential which we discuss later.

The VEVs of the Higgs fields are naturally given as (for the universal seesaw mechanism to work)

$$\langle H_{RQ}^0 \rangle = v_{RQ}, \quad \langle H_{Rl}^0 \rangle = v_{Rl}, \quad \langle H_{LQ}^0 \rangle = v_{LQ}, \quad \langle H_{Ll}^0 \rangle = v_{Ll}, \tag{4}$$

with the condition that $v_{LQ}^2 + v_{Ll}^2 = v_{\rm EW}^2$. The hierarchy of the VEVs responsible for symmetry breaking is arranged as

$$v_{RQ}, v_{Rl} \gg v_{LQ} > v_{Ll}. \tag{5}$$

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_{Ll} and H_{Rl} 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

$$\begin{aligned}
D_\mu Q_L &= \left[\partial_\mu - i\frac{g_L}{2}\tau.W_{L\mu} - i\frac{g_V}{6}V_\mu\right]Q_L,\\
D_\mu Q_R &= \left[\partial_\mu - i\frac{g_R}{2}\tau.W_{R\mu} - i\frac{g_V}{6}V_\mu\right]Q_R,\\
D_\mu l_L &= \left[\partial_\mu - i\frac{g_L}{2}\tau.W_{L\mu} + i\frac{g_V}{2}V_\mu\right]l_L,\\
D_\mu l_R &= \left[\partial_\mu - i\frac{g_R}{2}\tau.W_{R\mu} + i\frac{g_V}{2}V_\mu\right]l_R,\\
D_\mu H_R &= \left[\partial_\mu - i\frac{g_R}{2}\tau.W_{R\mu} - i\frac{g_V}{2}V_\mu\right]H_R,\\
D_\mu H_L &= \left[\partial_\mu - i\frac{g_L}{2}\tau.W_{L\mu} - i\frac{g_V}{2}V_\mu\right]H_L,
\end{aligned} \tag{6}$$

where V_μ is the $U(1)_{B-L}$ gauge boson and g_V its gauge coupling, while W_L, W_R and g_L, g_R are the gauge bosons and gauge couplings 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 covariant derivatives. The charged gauge boson mass-squared matrix in the basis $(W_R^\pm, W_L^\pm)$ is given as

$$\begin{bmatrix} \frac{1}{2}g_R^2(v_{RQ}^2+v_{Rl}^2) & 0 \\ 0 & \frac{1}{2}g_L^2(v_{LQ}^2+v_{Ll}^2) \end{bmatrix}. \tag{7}$$

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

TABLE I. Particle spectrum for the lepton-specific LR universal seesaw model.

Field	$SU(3)_C$	$SU(2)_L$	$SU(2)_R$	$U(1)_{B-L}$	Z_2
$Q_L = \begin{pmatrix} u \\ d \end{pmatrix}_L$	3	2	1	$\frac{1}{3}$	+
$Q_R = \begin{pmatrix} u \\ d \end{pmatrix}_R$	3	1	2	$\frac{1}{3}$	+
$l_L = \begin{pmatrix} \nu \\ e \end{pmatrix}_L$	1	2	1	−1	+
$l_R = \begin{pmatrix} \nu \\ e \end{pmatrix}_R$	1	1	2	−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	−
$H_{RQ} = \begin{pmatrix} H_{RQ}^+ \\ H_{RQ}^0 \end{pmatrix}$	1	1	2	1	+
$H_{LQ} = \begin{pmatrix} H_{LQ}^+ \\ H_{LQ}^0 \end{pmatrix}$	1	2	1	1	+
$H_{Rl} = \begin{pmatrix} H_{Rl}^+ \\ H_{Rl}^0 \end{pmatrix}$	1	1	2	1	−
$H_{Ll} = \begin{pmatrix} H_{Ll}^+ \\ H_{Ll}^0 \end{pmatrix}$	1	2	1	1	−

$$M_{W_R^\pm}^2 = \frac{1}{2}g_R^2(v_{RQ}^2+v_{Rl}^2), \quad M_{W^\pm}^2 = \frac{1}{2}g_L^2(v_{LQ}^2+v_{Ll}^2). \tag{8}$$

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

$$\begin{bmatrix} \frac{1}{4}g_R^2(v_{RQ}^2+v_{Rl}^2) & 0 & -\frac{1}{4}g_R g_V(v_{RQ}^2+v_{Rl}^2) \\ 0 & \frac{1}{4}g_L^2(v_{LQ}^2+v_{Ll}^2) & -\frac{1}{4}g_L g_V(v_{LQ}^2+v_{Ll}^2) \\ -\frac{1}{4}g_L g_V(v_{LQ}^2+v_{Ll}^2) & -\frac{1}{4}g_R g_V(v_{RQ}^2+v_{Rl}^2) & \frac{1}{4}g_V^2(v_{RQ}^2+v_{Rl}^2+v_{LQ}^2+v_{Ll}^2) \end{bmatrix}. \tag{9}$$

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

$$\begin{aligned}
M_{Z_R}^2 &\simeq \frac{1}{2}\left[(g_R^2+g_V^2)(v_{RQ}^2+v_{Rl}^2) + \frac{g_V^4(v_{LQ}^2+v_{Ll}^2)}{g_R^2+g_V^2}\right],\\
M_Z^2 &\simeq \frac{1}{2}(g_L^2+g_Y^2)(v_{LQ}^2+v_{Ll}^2),
\end{aligned} \tag{10}$$

with the effective SM $U(1)_Y$ gauge coupling given as

$$g_Y = \frac{g_L g_V}{\sqrt{g_L^2+g_V^2}}. \tag{11}$$

Quite clearly, in this 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

$$\begin{aligned}\mathcal{L}_Y = (&Y_{uL}\bar{Q}_L\tilde{H}_{LQ}U_R + Y_{uR}\bar{Q}_R\tilde{H}_{RQ}U_L + Y_{dL}\bar{Q}_L H_{LQ}D_R \\ &+ Y_{dR}\bar{Q}_R H_{RQ}D_L + Y_{\nu L}\bar{l}_L\tilde{H}_{Ll}N_R + Y_{\nu R}\bar{l}_R\tilde{H}_{Rl}N_L \\ &+ Y_{eL}\bar{l}_L H_{Ll}E_R + Y_{eR}\bar{l}_R H_{Rl}E_L + M_U\bar{U}_L U_R \\ &+ M_D\bar{D}_L D_R + M_E\bar{E}_L E_R + M_N\bar{N}_L N_R + \text{H.c.}) \\ &+ M_L N_L N_L + M_R N_R N_R \end{aligned} \tag{12}$$

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

$$\tilde{H}_{L/R} = i\tau_2 H^*_{L/R}. \tag{13}$$

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 6×6 mass matrix quite similar to what happens in the seesaw mechanism. The up quark mass terms in this model can be written as

$$\mathcal{L}_u = (\bar{u} \quad \bar{U})(M_u P_L + M_u^T P_R)\begin{pmatrix} u \\ U \end{pmatrix}, \tag{14}$$

where

$$M_u = \begin{pmatrix} 0 & Y_{uR}v_{RQ} \\ Y_{uL}^T v_{LQ} & M_U \end{pmatrix} \tag{15}$$

is the 6×6 up quark mass matrix while Y_{uL}, Y_{uR} and M_U are all 3×3 matrices. The first 3×3 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_{uL} and Y_{uR} terms of Eq. (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 the SM.

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

$$M_d = \begin{pmatrix} 0 & Y_{dR}v_{RQ} \\ Y_{dL}^T v_{LQ} & M_D \end{pmatrix}, \tag{16}$$

where the first 3×3 block is again zero due to the absence of a bidoublet scalar while the off-diagonal blocks arise from the Yukawa couplings. The M_D term is the mass term for the heavy down-type quarks. In the down sector too, we keep the right-handed 3×3 Yukawa matrix Y_{dR} and the M_D matrix to be diagonal while only the left-handed Yukawa matrix Y_{dL} is nondiagonal and sufficient to generate the correct CKM mixings for the SM quarks.

To diagonalize these nonsymmetric matrices we require biunitary transformations. For the up-type quark mass matrix we have

$$M_u^{\text{diag}} = U_{uL}M_u U_{uR}^\dagger, \tag{17}$$

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

$$M_d^{\text{diag}} = U_{dL}M_d U_{dR}^\dagger. \tag{18}$$

We will get two CKM mixing matrices in this case—one for the left-handed quarks and another for the right-handed quarks—given by

$$U_L^{\text{CKM}} = U_{uL}U_{dL}^\dagger \tag{19}$$

and

$$U_R^{\text{CKM}} = U_{uR}U_{dR}^\dagger \tag{20}$$

respectively. These will be 6×6 matrices whose top-left (bottom-right) 3×3 block will correspond to the light CKM mixings for ascending (descending) arrangement of eigenvalues by mass. For our choice of parameters and with only the left-handed Yukawa being nondiagonal, 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 is 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

$$M_e = \begin{pmatrix} 0 & Y_{eR}v_{Rl} \\ Y_{eL}^T v_{Ll} & M_E \end{pmatrix}. \tag{21}$$

This is very similar to the quark mass matrix with Y_{eL} and Y_{eR} being the 3×3 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 is 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

$$\begin{pmatrix} 0 & Y_{\nu L} v_{Ll} & 0 & 0 \\ Y_{\nu L}^T v_{Ll} & M_R & 0 & M_N^T \\ 0 & 0 & 0 & Y_{\nu R}^T v_{Rl} \\ 0 & M_N & Y_{\nu R} v_{Rl} & M_L \end{pmatrix}. \tag{22}$$

Thus we see that all the fermion masses in this model arise from the seesawlike 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 couplings and 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 with mass at around the electroweak scale and the rest at the TeV scale. We will refer to this scenario as the 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 nondiagonal symmetric matrix to explain the experimentally observed Pontecorvo-Maka-Nakagawa-Sakata (PMNS) neutrino mixing matrix elements. The mixing between the light and the heavy states is 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 (ν_L, N_R) and (ν_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 the 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 nondiagonal. 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

$$Y_{\nu R_{ij}} = \frac{v_{Ll}}{v_{Rl}} Y_{\nu L_{ij}}, \tag{23}$$

given we take $M_L = M_R$. The mixing between the light and heavy states in each block diagonal submatrix 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

$$\begin{aligned} V(H) = {} & \sum_{i=1}^{4} \mu_{ii} H_i^\dagger H_i + \sum_{\substack{i,j=1\\ i\leq j}}^{4} \lambda_{ij} H_i^\dagger H_i H_j^\dagger H_j \\ & + \Big[\alpha_1 H_{LQ}^\dagger H_{Ll} H_{RQ}^\dagger H_{Rl} + \alpha_2 H_{LQ}^\dagger H_{Ll} H_{Rl}^\dagger H_{RQ} \\ & + \mu_{12}^2 H_{LQ}^\dagger H_{Ll} + \mu_{34}^2 H_{RQ}^\dagger H_{Rl} + \text{H.c.}\Big] \end{aligned} \tag{24}$$

where

$$H_1 = H_{LQ}, \quad H_2 = H_{Ll}, \quad H_3 = H_{RQ}, \quad H_4 = H_{Rl}. \tag{25}$$

The last two terms involving μ_{12} and μ_{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 minimization conditions

$$\begin{aligned} \mu_{11} &= \frac{1}{v_{LQ}} (\alpha_{12}^+ v_{Ll} v_{Rl} v_{RQ} + 2\lambda_{11} v_{LQ}^3 + \lambda_{12} v_{Ll}^2 v_{LQ} \\ &\quad + \lambda_{13} v_{LQ} v_{RQ}^2 + \lambda_{14} v_{LQ} v_{Rl}^2 - \mu_{12}^2 v_{Ll}), \\ \mu_{22} &= \frac{1}{v_{Ll}} (\alpha_{12}^+ v_{LQ} v_{Rl} v_{RQ} + \lambda_{12} v_{Ll} v_{LQ}^2 + 2\lambda_{22} v_{Ll}^3 \\ &\quad + \lambda_{23} v_{Ll} v_{RQ}^2 + \lambda_{24} v_{Ll} v_{Rl}^2 - \mu_{12}^2 v_{LQ}), \\ \mu_{33} &= \frac{1}{v_{RQ}} (\alpha_{12}^+ v_{Ll} v_{LQ} v_{Rl} + \lambda_{13} v_{LQ}^2 v_{RQ} + \lambda_{23} v_{Ll}^2 v_{RQ} \\ &\quad + 2\lambda_{33} v_{RQ}^3 + \lambda_{34} v_{Rl}^2 v_{RQ} - \mu_{34}^2 v_{Rl}), \\ \mu_{44} &= \frac{1}{v_{Rl}} (\alpha_{12}^+ v_{Ll} v_{LQ} v_{RQ} + \lambda_{14} v_{LQ}^2 v_{Rl} + \lambda_{24} v_{Ll}^2 v_{Rl} \\ &\quad + \lambda_{34} v_{Rl} v_{RQ}^2 + 2\lambda_{44} v_{Rl}^3 - \mu_{34}^2 v_{RQ}), \end{aligned} \tag{26}$$

where $\alpha_{12}^+ = (\alpha_1 + \alpha_2)$.

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 bosons 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 and Z_R. The charged Higgs mass-squared matrix in this case is a 4×4 block diagonal matrix with two blocks of 2×2. In the basis $(H^+_{LQ}, H^+_{Ll}, H^+_{RQ}, H^+_{Rl})$ the charged Higgs mass-squared matrix is given as

$$\begin{pmatrix} \frac{v_{Ll}}{v_{LQ}}\{\mu^2_{12} - \alpha^+_{12} v_{Rl} v_{RQ}\} & \alpha^+_{12} v_{Rl} v_{RQ} - \mu^2_{12} & 0 & 0 \\ \alpha^+_{12} v_{Rl} v_{RQ} - \mu^2_{12} & \frac{v_{LQ}}{v_{Ll}}\{\mu^2_{12} - \alpha^+_{12} v_{Rl} v_{RQ}\} & 0 & 0 \\ 0 & 0 & \frac{v_{Rl}}{v_{RQ}}\{\mu^2_{34} - \alpha^+_{12} v_{Ll} v_{LQ}\} & \alpha^+_{12} v_{Ll} v_{LQ} - \mu^2_{34} \\ 0 & 0 & \alpha^+_{12} v_{Ll} v_{LQ} - \mu^2_{34} & \frac{v_{RQ}}{v_{Rl}}\{\mu^2_{34} - \alpha^+_{12} v_{Ll} v_{LQ}\} \end{pmatrix}. \tag{27}$$

Diagonalizing this matrix we get two Goldstone states which are given as

$$G^+_1 = \frac{1}{v^2_{Ll} + v^2_{LQ}} (v_{LQ}, v_{Ll}, 0, 0)^T,$$
$$G^+_2 = \frac{1}{v^2_{Rl} + v^2_{RQ}} (0, 0, v_{RQ}, v_{Rl})^T. \tag{28}$$

The two physical charged Higgs boson masses are

$$m^2_{H^+_1} = \frac{v^2_{Rl} + v^2_{RQ}}{v_{Rl} v_{RQ}} (\mu^2_{34} - \alpha^+_{12} v_{Ll} v_{LQ}),$$
$$m^2_{H^+_2} = \frac{v^2_{Ll} + v^2_{LQ}}{v_{Ll} v_{LQ}} (\mu^2_{12} - \alpha^+_{12} v_{Rl} v_{RQ}), \tag{29}$$

with the eigenstates being

$$H^+_1 = \frac{1}{v^2_{Rl} + v^2_{RQ}} (0, 0, -v_{Rl}, v_{RQ})^T,$$
$$H^+_2 = \frac{1}{v^2_{Ll} + v^2_{LQ}} (-v_{Ll}, v_{LQ}, 0, 0)^T. \tag{30}$$

It is easy to see here that if we choose $\mu^2_{12} = \mu^2_{34} \sim v^2_{\mathrm{EW}}$ 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_{Ll} except for a very fine-tuned region around

$$\alpha_1 + \alpha_2 \approx \frac{\mu^2_{12}}{v_{Rl} v_{RQ}}. \tag{31}$$

The CP-odd Higgs boson mass-squared matrix in the basis $(\mathrm{Im}H^0_{LQ}, \mathrm{Im}H^0_{Ll}, \mathrm{Im}H^0_{RQ}, \mathrm{Im}H^0_{Rl})$ is

$$\begin{pmatrix} \frac{v_{Ll}}{v_{LQ}}(\mu^2_{12} - \alpha^+_{12} v_{Rl} v_{RQ}) & \alpha^+_{12} v_{Rl} v_{RQ} - \mu^2_{12} & v_{Ll} v_{Rl}(\alpha_2 - \alpha_1) & v_{Ll} v_{RQ}(\alpha_1 - \alpha_2) \\ \alpha^+_{12} v_{Rl} v_{RQ} - \mu^2_{12} & \frac{v_{LQ}}{v_{Ll}}(\mu^2_{12} - \alpha^+_{12} v_{Rl} v_{RQ}) & v_{LQ} v_{Rl}(\alpha_1 - \alpha_2) & v_{LQ} v_{RQ}(\alpha_2 - \alpha_1) \\ v_{Ll} v_{Rl}(\alpha_2 - \alpha_1) & v_{LQ} v_{Rl}(\alpha_1 - \alpha_2) & \frac{v_{Rl}}{v_{RQ}}(\mu^2_{34} - \alpha^+_{12} v_{Ll} v_{LQ}) & \alpha^+_{12} v_{Ll} v_{LQ} - \mu^2_{34} \\ v_{Ll} v_{RQ}(\alpha_1 - \alpha_2) & v_{LQ} v_{RQ}(\alpha_2 - \alpha_1) & \alpha^+_{12} v_{Ll} v_{LQ} - \mu^2_{34} & \frac{v_{RQ}}{v_{Rl}}(\mu^2_{34} - \alpha^+_{12} v_{Ll} v_{LQ}) \end{pmatrix}. \tag{32}$$

This again will have two zero eigenstates corresponding to the two Goldstone bosons required for Z and Z_R mass generation. It is also easy 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 $(\mathrm{Re}H^0_{LQ}, \mathrm{Re}H^0_{Ll}, \mathrm{Re}H^0_{RQ}, \mathrm{Re}H^0_{Rl})$ can be expressed in terms of the CP-odd Higgs mass-squared matrix elements as

$$M^2_{ij,\text{CP-Even}} = M^2_{ij,\text{CP-Odd}} + 2S_{ij}\lambda_{ij} v_i v_j, \tag{33}$$

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

$$S_{ij} = \begin{cases} 2, & \text{if } i = j \\ 1, & \text{otherwise.} \end{cases} \tag{34}$$

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. This state is consistent with the SM Higgs properties in all its decay channels and branching ratios, while all the other states are chosen to be much heavier. Note that we have implemented the model in SARAH [8] and

TABLE II. Scalar eigenstates for $\alpha_1 = -0.2, \alpha_2 = 0.1, \lambda_{11} = 0.168, \lambda_{12} = 0.8, \lambda_{13} = 0.05, \lambda_{14} = -0.1, \lambda_{22} = 0.5, \lambda_{23} = 0.1, \lambda_{24} = 0.1, \lambda_{33} = 0.2, \lambda_{34} = 0.1, \lambda_{44} = 0.1, \mu_{12}^2 = 2.5\times 10^4, \mu_{34}^2 = 2.5\times 10^4$, $v_{RQ} = v_{Rl} = 6.0$ TeV, $v_{LQ} = 173.4$ GeV, $v_{Ll} = 14$ GeV.

Particle	Mass (GeV)	Eigenstate
H_1	125.2	$0.996\mathrm{Re}(H_{LQ}^0) - 0.008\mathrm{Re}(H_{RQ}^0) + 0.080\mathrm{Re}(H_{Ll}^0) + 0.019\mathrm{Re}(H_{Rl}^0)$
H_2	3386.1	$-0.0209\mathrm{Re}(H_{LQ}^0) - 0.381\mathrm{Re}(H_{RQ}^0) + 0.001\mathrm{Re}(H_{Ll}^0) + 0.924\mathrm{Re}(H_{Rl}^0)$
H_3	5638.1	$0.001\mathrm{Re}(H_{LQ}^0) - 0.924\mathrm{Re}(H_{RQ}^0) + 0.008\mathrm{Re}(H_{Ll}^0) - 0.381\mathrm{Re}(H_{Rl}^0)$
H_4	6772.5	$0.080\mathrm{Re}(H_{LQ}^0) - 0.007\mathrm{Re}(H_{RQ}^0) - 0.997\mathrm{Re}(H_{Ll}^0) + 0.004\mathrm{Re}(H_{Rl}^0)$
A_1	214.8	$0.001\mathrm{Im}(H_{LQ}^0) - 0.707\mathrm{Im}(H_{RQ}^0) - 0.010\mathrm{Im}(H_{Ll}^0) + 0.707\mathrm{Im}(H_{Rl}^0)$
A_2	6772.7	$-0.080\mathrm{Im}(H_{LQ}^0) - 0.007\mathrm{Im}(H_{RQ}^0) + 0.997\mathrm{Im}(H_{Ll}^0) + 0.006\mathrm{Im}(H_{Rl}^0)$
H_1^+	224.7	$-0.707H_{RQ}^+ + 0.707H_{Rl}^+$
H_2^+	6772.4	$0.080H_{LQ}^+ - 0.997H_{Ll}^+$

use the generated SPHENO [9] code to obtain the model spectrum and calculate the decay of various particles. We have checked that the light Higgs boson of 125 GeV is consistent with the expected branching ratios as well as the total decay width of the Standard Model Higgs boson. If we consider the $H_1 \to \gamma\gamma$ decay channel for instance, it gives us a branching ratio of 2.27×10^{-3}. We can write the partial decay width as [10]

$$\Gamma_{H_1\to\gamma\gamma} = \frac{G_F\alpha^2 m_{H_1}^3}{128\sqrt{2}\pi^3}\left|\sum_f N_c Q_f^2 g_{H_1 ff} A_{1/2}^{H_1}(\tau_f) + g_{H_1 VV} A_1^{H_1}(\tau_W) + \frac{M_W^2 \lambda_{H_1 H_1^+ H_1^-}}{2c_W^2 M_{H_1^\pm}^2} A_0^{H_1}(\tau_{H_1^\pm}) + \frac{M_W^2 \lambda_{H_1 H_2^+ H_2^-}}{2c_W^2 M_{H_2^\pm}^2} A_0^{H_1}(\tau_{H_2^\pm})\right|^2 \quad (35)$$

which can be reexpressed as

$$\Gamma_{H_1\to\gamma\gamma} = \frac{G_F\alpha^2 m_{H_1}^3}{128\sqrt{2}\pi^3}\left|\sum\nolimits_f F_{1/2}^{H_1}(\tau_f, g_{H_1 ff}) + F_1^{H_1}(\tau_W, g_{H_1 VV}) + \sum\nolimits_i F_0^{H_1}(\tau_{H_i^\pm}, \lambda_{H_1 H_i^+ H_i^-})\right|^2 \quad (36)$$

giving us a better idea of the relative contribution from each sector. The benchmark point that we have chosen (Table II) gives us $\lambda_{H_1 H_1^+ H_1^-} = 0.034$, $\lambda_{H_1 H_2^+ H_2^-} = 1.20$, $M_{H_1^\pm} = 224.7$ GeV, $M_{H_2^\pm} = 6772.4$ GeV. The terms in Eq. (36) are thus

$$F_{1/2}^{H_1}(\tau_t, g_{H_1 t\bar{t}}) = 1.835,$$
$$F_1^{H_1}(\tau_W, g_{H_1 VV}) = -8.324,$$
$$F_0^{H_1}(\tau_{H_1^\pm}, \lambda_{H_1 H_1^+ H_1^-}) = 0.0013,$$
$$F_0^{H_1}(\tau_{H_2^\pm}, \lambda_{H_1 H_2^+ H_2^-}) = 3.6\times 10^{-5},$$

where the four contributions are from the top quark, W boson and the two charged Higgs states respectively. Note that the contributions from the charged Higgs bosons are orders of magnitude lower compared to the top quark and gauge boson contributions and do not affect the $H_1 \to \gamma\gamma$ branching ratio.

The lightest charged and pseudoscalar 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 2HDMs, here only the pseudoscalar and charged Higgs bosons are light while all other 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. To check whether the Higgs potential is stable for our choice of benchmark points, we examine the copositivity conditions for the stability of the potential when the couplings are negative [11]. The condition for the stability of the potential for the negative coupling λ_{14} is given by

$$\lambda_{11} \geq 0, \qquad \lambda_{44} \geq 0, \qquad \lambda_{14} \geq -\sqrt{\lambda_{11}\lambda_{44}}, \quad (37)$$

which are easily satisfied. We checked the condition for α_1 numerically by constructing the principal submatrices and found that it satisfies the criteria for stability as well.

III. PHENOMENOLOGICAL IMPLICATIONS

We now look at the phenomenological implications of this model, viz. 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 the Y_{dL} and Y_{ν_L} matrices. These are necessarily off diagonal in order to generate the CKM and PMNS mixing. Unlike the 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

$$Y^{11}_{u(L,R)} \sim Y^{11}_{d(L,R)} \sim Y^{11}_{e(L,R)} \approx 10^{-2},$$
$$Y^{22}_{u(L,R)} \sim Y^{22}_{dR} \sim Y^{22}_{e(L,R)} \approx 10^{-1},$$
$$Y^{33}_{u(L,R)} \sim Y^{33}_{dL} \sim Y^{33}_{eR} \approx 1, \tag{38}$$

while the other elements in the Y_{dL} matrix are of the order of 10^{-3}. We further choose $Y^{33}_{dR} = 0.023$ and $Y^{33}_{eL} = 0.26$ so that the third-generation heavy fermion masses are all of the order of a few TeV. With this type of Yukawa structure we can easily get the correct masses of all the fermions by choosing the appropriate values of the heavy masses. The left-handed CKM matrix 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 VEVs are taken to be

$$v_{RQ} = v_{Rl} = 6.0\,\text{TeV}, \quad v_{LQ} = 173.4\,\text{GeV}, \quad v_{Ll} = 14\,\text{GeV}. \tag{39}$$

The bare masses of the singlet heavy vectorlike 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 vectorlike singlet fermion masses. For each type of fermion (up quark, down quark and charged leptons), our choices are such that the third-generation vectorlike fermion is the lightest while the first generation is the heaviest. 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. In Table III we list the mass of all the new fermions in our model. With the strong sector exotic quarks and charged leptons having masses above 3.5 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 TeV machines such as the FCC-hh Collider [12].

The mixing between the heavy singletlike states and the light SM-like states is 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^{33}_u to be quite small to be around 350 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 is clearly visible in the W_R decay modes where its branching ratio $W^+_R \to t\bar{b}$ is significantly suppressed (~0.2%) while for the first two-generation light quarks it is 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. The current bound on a heavy SM-like W' from the LHC is 2.6 TeV [13]. At a mass of 2.6 TeV, a SM-like W' goes into light quarks with a branching ratio of 47.6%. The W_R in our case has a 66% branching ratio into light quarks and therefore the limits would be stronger. Using the experimental bound on the cross section×branching ratio ($\sigma \times BR$) we get a lower bound on W_R mass of 2.75 TeV in our model. For our choice of benchmark points, the mass of the W_R boson comes out to be ~4 TeV and is safe from dijet bounds. It is also allowed from heavy neutrino searches [14], provided the heavy neutrino mass is not very heavy. We show the W_R decay modes and the branching probabilities as a function of its mass in Fig. 1.

In the neutrino sector there are two specific scenarios (Majorana and pseudo-Dirac) as has been discussed earlier.

TABLE III. Fermion masses.

Up-type quark	Down-type quark	Charged lepton	Neutrino: Majorana	Neutrino: Pseudo-Dirac
$M_T = 4.51$ TeV, $M_C = 6.17$ TeV, $M_U = 30.0$ TeV	$M_B = 3.97$ TeV, $M_S = 10.4$ TeV, $M_D = 17.2$ TeV	$M_{E_3} = 6.13$ TeV, $M_{E_2} = 9.92$ TeV, $M_{E_1} = 12.3$ TeV	$M_{\nu_4} = 136$ GeV, $M_{\nu_5} = 258$ GeV, $M_{\nu_6} = 317$ GeV, $M_{\nu_7} = 9.07$ TeV, $M_{\nu_8} = 9.13$ TeV, $M_{\nu_9} = 9.16$ TeV, $M_{\nu_{10}} = 11.06$ TeV, $M_{\nu_{11}} = 11.1$ TeV, $M_{\nu_{12}} = 11.2$ TeV	$M_{\nu_4} = 200.0$ GeV, $M_{\nu_5} = 300.0$ GeV, $M_{\nu_6} = 400.0$ GeV

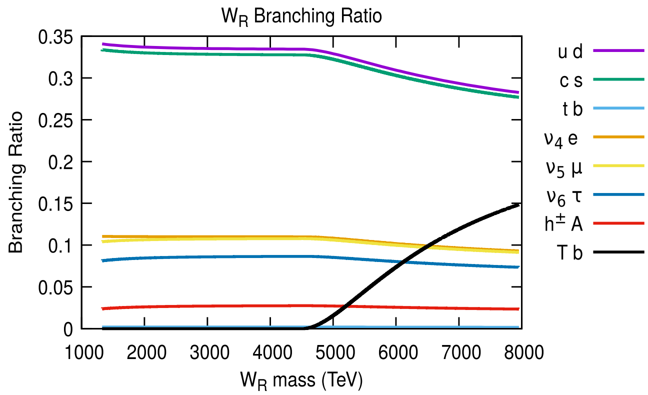

FIG. 1. Branching ratio for the W_R boson as a function of its mass.

Here we will only discuss the case of normal hierarchy for the neutrino masses.[2] For the Majorana case, we fit the experimental data for neutrino oscillation parameters [15] with the variations being within the 3σ range of their 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

$$M_R = M_L = \mathrm{Diag}(10^4, 10^4, 10^4),$$
$$M_N = \mathrm{Diag}(10^3, 10^3, 10^3), \tag{40}$$

while the elements of the Yukawa coupling matrix $Y_{\nu R_{ii}} \sim 0.1$ and $Y_{\nu_{L_{ij}}} \sim 10^{-5}$. This choice gives us our desired neutrino masses with the three light neutrino states lying between 0.001 to 0.05 eV while the next three heavy states have masses around a few 100 GeV. The rest of the physical states have mass around 10 TeV. The three light neutrino physical states are almost entirely from the three generations of ν_L and the next three (masses of a few 100 GeV) come mostly from ν_R. This is very similar to the type-I seesaw in conventional LRS models and would give very similar phenomenology with same-sign lepton signals and neutrinoless double-beta decay. However the modified scalar sector interaction with the heavy neutrinos leads 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 give us the correct neutrino mass-squared differences and the correct PMNS mixing angles for normal hierarchy. We can see that $Y_{\nu_{33}}$ is indeed the largest owing to ν_3 being the heaviest in the normal hierarchy case, while $Y_{\nu_{13}}$ is the smallest in magnitude as required to explain the small value of the mixing angle θ_{13}.

[2]It is worth noting that an arrangement for the inverted hierarchy of the neutrino masses is equally possible in our model, which we have not considered here.

TABLE IV. Experimental 3σ ranges for light neutrino parameters.

$$7.03 \times 10^{-5}\ \mathrm{eV}^2 < \Delta m^2_{21} < 8.09 \times 10^{-5}\ \mathrm{eV}^2$$
$$2.407 \times 10^{-3}\ \mathrm{eV}^2 < \Delta m^2_{31} < 2.643 \times 10^{-3}\ \mathrm{eV}^2$$
$$0.271 < \sin^2\theta_{12} < 0.345$$
$$0.385 < \sin^2\theta_{23} < 0.635$$
$$0.01934 < \sin^2\theta_{13} < 0.02392$$
$$U_{PMNS}$$
$$\begin{pmatrix} 0.800 \to 0.844 & 0.515 \to 0.581 & 0.139 \to 0.155 \\ 0.229 \to 0.516 & 0.438 \to 0.699 & 0.614 \to 0.790 \\ 0.249 \to 0.528 & 0.462 \to 0.715 & 0.595 \to 0.776 \end{pmatrix}$$

The 12×12 neutrino mass matrix is symmetric and hence it can be diagonalized with a simple unitary transformation. The first 3×3 block corresponds to the three light neutrinos and satisfies the experimental 3σ bounds of the PMNS matrix. The other mixing of the light neutrinos with the heavier ones is extremely small with $\sin\theta_{ij} \lesssim 10^{-8}$ where θ_{ij} is the mixing angle between the heavy and light states. Hence there are no bounds 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 ν_{R_i} with small mixing ($\sim 1\%$) with N_{L_i}. The six heavier states of masses around 10 TeV are almost equally constituted 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 $M_R = M_L = \mathrm{Diag}(200, 300, 400)$. 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 nonzero off-diagonal elements (still being symmetric matrices). Now we get two block diagonal 6×6 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 Eq. (23) along with the observed mass-squared differences and mixing constraints gives us the following choice of the matrix elements: $Y_{\nu L_{ij}} \sim 10^{-6}$ and $Y_{\nu R_{ij}} \sim 10^{-9}$. The neutrino mixings in this case again have to satisfy the experimental PMNS mixing limits and have to be the same for both sectors. This is easily satisfied by using the condition given in Eq. (23). Figure 2(b) gives some allowed parameters for this case which satisfy the neutrino mass-squared differences and the mixing angles for normal hierarchy. As observed in the previous case, we find that $Y_{\nu_{33}}$ 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 ν_L and ν_R which are now of equal masses in this scenario are not quite as small as the previous case. Typically, the mixing angle θ between

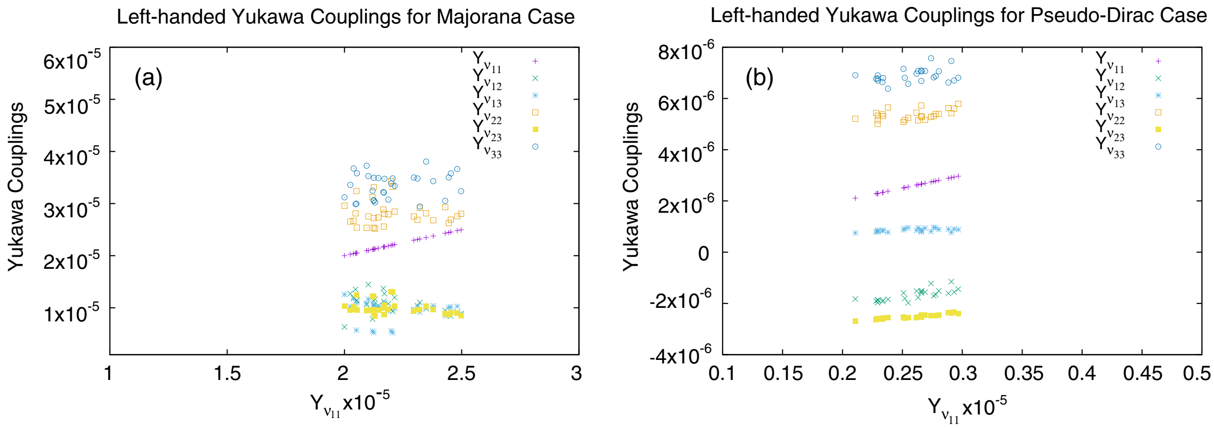

FIG. 2. (a) Benchmark points satisfying the neutrino masses and PMNS mixings for the Majorana case. (b) Benchmark points satisfying the neutrino masses and PMNS mixings for the pseudo-Dirac case.

two light states of equal masses is such that $\sin\theta \sim 10^{-2}$. The heavy states again have negligibly small mixing with light states like in the previous case. Note that as we have taken $M_N = 0$ there is no mixing between the heavy states and they are purely compositions of N_{L_i} or N_{R_i}. As a result their decay width in this case becomes very small and with suitable choice of parameters it may lead to observable displaced vertex signals. It is worth noting that when the charged Higgs is lighter than the heavy neutrino states, it becomes the primary channel of decay and therefore can provide for interesting signal channels for the model at collider experiments, which we discuss later in more detail.

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 2HDM. As such there are a number of flavor constraints which restrict the parameter space of the 2HDM. The most stringent of these constraints comes from the $b \to s\gamma$ process and constrains the charged Higgs mass to $m_{H^\pm} > 460$ GeV [16] in the type-III 2HDM. The main process responsible for $b \to s\gamma$ in the 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 the right-handed down-type Yukawa coupling matrix is diagonal in this model, the CKM mixings are really small. This results in a much weaker bound on the lightest charged Higgs boson

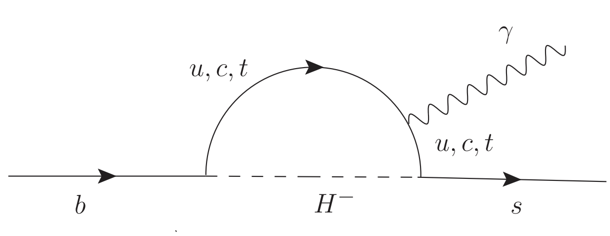

FIG. 3. $b \to s\gamma$ through the charged Higgs.

mass in this case. There are no significant bounds from the flavor observable on the pseudoscalar mass in the lepton-specific 2HDM and hence there are no bounds in this model as well.

There is a bound on the 2HDM pseudoscalar Higgs boson mass from the single production and associated production of the pseudoscalar decaying into two τ final states [17]. This gives a lower limit on the pseudoscalar mass as a function of the $\sigma \times \mathrm{BR}(A \to \tau\tau)$ 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 the 2HDM. This is because the coupling here would be

$$f^L_{Au\bar{u}(Ad\bar{d})} = Y^{33}_{uL(dL)} \times Z^L_{Ai} \times Z^{Qq}_{uR(dR)} \tag{41}$$

where Z^L_{Ai} is the amount of H^0_{LQ} contained in the eigenstate of A and $Z^{qQ}_{uR(dR)}$ is the mixing of the right-handed heavy and the light up-type (down-type) quarks. A similar formula can be written for the right-handed pseudoscalar coupling with two quarks with $L \leftrightarrow R$ in Eq. (41). 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 the 2HDM.

B. New collider signals

This model can lead to a number of interesting new signals at accelerator experiments. We primarily focus on 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 LHC. Note

TABLE V. Representative benchmark points of the particle spectrum where the massive neutrino states are Majorana type. We also illustrate the dominant decay modes of the lightest charged Higgs and heavy neutrino states. The heavy gauge bosons for both BP1 and BP2 are the same with $M_{W_R} = 4$ TeV and $M_{Z_R} = 4.7$ TeV.

Particle		Width (GeV)	Important decay channels
BP1	$M_{\nu_4} = 136.4$ GeV	7.12×10^{-9}	$\nu_4 \to e^\pm jj \sim 99.3\%$
	$M_{\nu_5} = 258.3$ GeV	4.57×10^{-4}	$\nu_5 \to \mu^\pm H_1^\mp \sim 100\%$
	$M_{\nu_6} = 317.0$ GeV	2.97×10^{-3}	$\nu_6 \to \tau^\pm H_1^\mp \sim 100\%$
	$M_{H_1^\pm} = 224.7$ GeV	4.6×10^{-4}	$H_1^\pm \to \nu_4 e^\pm \sim 99\%$
BP2	$M_{\nu_4} = 317.0$ GeV	2.11×10^{-3}	$\nu_4 \to e^\pm H_1^\mp \sim 100\%$
	$M_{\nu_5} = 550.9$ GeV	3.03×10^{-2}	$\nu_5 \to \mu^\pm H_1^\mp \sim 100\%$
	$M_{\nu_6} = 837.6$ GeV	1.04×10^{-1}	$\nu_6 \to \tau^\pm H_1^\mp \sim 100\%$
	$M_{H_1^\pm} = 224.7$ GeV	5.78×10^{-6}	$H_1^+ \to t\bar{b} \sim 92.9\%$

that the other exotics such as the heavy quarks and leptons are beyond the reach of the LHC because of their extremely heavy masses.

To highlight the signals for the model at the LHC we choose two representative points in the model parameter space as BP1 and BP2 and list them in Table V. 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 $Y_{\nu R_{ii}} \sim 0.1$ 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 components. 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 for BP1 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 of observing the signals of these heavy neutrinos would be dictated by the production mechanism. At the 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

$$pp \to W_R^\pm \to \ell_i^\pm \nu_j,$$

$$pp \to Z_R \to \nu_j \nu_j$$

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

We find that a W_R of mass 4 TeV consistent with current experimental limits 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 two generations ~11%. Similarly a Z_R of mass around 4.8 TeV has a combined ~20%–22% branching probability to decay in the pair of ν_4, ν_5, ν_6 for BP1 and BP2, respectively. Resonant production of heavy neutrinos can be useful to have appreciable rates of production [18] 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 in association with jets with a 100% branching probability for BP1, leading to same-sign dilepton signals with jets in the final state provided the charged Higgs in the model is heavier. In fact one can also have the interesting signal where

$$pp \to \nu_5\nu_5 \to \mu^\pm \mu^\pm H_1^\mp H_1^\mp \to 2\mu^\pm + 2e^\mp + 4j.$$

The pair production cross sections at the LHC for the heavy neutrinos with $\sqrt{s} = 13$ TeV are

$$\begin{aligned} \text{BP1:} \quad & \sigma(pp \to \nu_4\nu_4) = 0.121 \text{ fb}, \\ & \sigma(pp \to \nu_5\nu_5) = 0.101 \text{ fb}, \\ & \sigma(pp \to H_1^+ H_1^-) = 5.27 \text{ fb}. \end{aligned}$$

The ν_4 production gives the familiar same-sign lepton signal

$$pp \to \nu_4\nu_4 \to 2e^\pm + 4j.$$

By slightly changing our parameters we get a spectrum where all the heavy neutrinos are actually heavier than the lightest charged Higgs represented by BP2 as shown in Table V. We have only modified the neutrino sector making sure that the new set of parameters is still consistent with the neutrino oscillation data, while keeping the other sectors almost the same as before. In this case where all the Majorana neutrinos are heavier than the lightest charged Higgs which in turn is heavier than the top quark, 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 $H_1^+ \to t\bar{b}$ (see Table V). 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 quarks in the final state. This is completely free from any SM background and would be a unique signal for discovery. Thus we have for example

$$pp \to \nu_4\nu_4 \to e^\pm e^\pm H_1^\mp H_1^\mp \to 2e^\pm + 2\bar{t}/t + 2b/\bar{b}$$

when the ν_4 is pair produced. Again one gets $2\mu^\mp + 2t/\bar{t} + 2\bar{b}/b$ when ν_5 is pair produced. The total cross sections for the pair production of the heavy neutrino pairs at the LHC with $\sqrt{s} = 13$ TeV are

$$\begin{aligned} \text{BP2:}\ &\sigma(pp \to \nu_4\nu_4) = 0.096\ \text{fb}, \\ &\sigma(pp \to \nu_5\nu_5) = 0.076\ \text{fb}. \end{aligned}$$

Subsequent decay probabilities for heavy neutrino decay as well as charged Higgs decay are almost 100%. To compare it with the SM background, there are no subprocesses that can contribute directly to similar final states. Due to charge mismeasurements we can consider the process $pp \to e^+e^- t\bar{t}b\bar{b}$ as a possible background. The cross section for this process at the LHC with $\sqrt{s} = 13$ TeV is around 0.16 fb. The charge mismeasurement probabilities are much below 10^{-3} and therefore this would hardly give any event even with an integrated luminosity of 3000 fb^{-1}. The signal would still yield a handsome 516 events for BP2 where we add the contributions for both the ν_4 and ν_5 channels in the case of BP2. Although the signal would be difficult to observe in the near future at both ATLAS and CMS, but with the very high luminosity option at the LHC, it would be a very unique channel to observe.

For the charged Higgs ($H_1^\pm$) 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 the 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 a same-sign lepton along with a top and bottom quark, where $\sigma(pp \to W_R) = 3.2$ fb for $M_{W_R} = 4$ TeV. 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 the LHC here.

In addition, if the charged Higgs's are heavier than the heavy neutrinos, then they would dominantly decay into them and the corresponding charged lepton. This can lead to a 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, the presence of a light Majorana neutrino completely overwhelms the tb decay option. The charged Higgs production would mostly be via the photon exchange,

$$pp \to H_1^+ H_1^- \to \nu_j \nu_j \ell_i^+ \ell_i^-$$

where again $i = 1, 2, 3$ while $j = i + 3$. The heavy neutrino would decay via the off-shell charged Higgs in the three-body decay channel $\nu_j \to \ell^\pm jj'$. This quite clearly gives a multilepton 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. For example in the case of BP1 there is a four-lepton channel contribution coming from the pair production of charged Higgs

$$pp \to H_1^+ H_1^- \to \nu_4 e^+ \nu_4 e^- \to 2e^\pm + 4j + e^+e^-.$$

Here again this is a very interesting signal channel in the form of three same-sign electrons which has negligible SM background. This would prove to be an interesting signal [19] to look for the charged Higgs search in this model.

Another unique signal of this model which differentiates it from other LR models involves the W_R decay channels. Figure 1 gives a plot of the various W_R decay branching ratios in this model. In general LR models an important decay channel for the W_R boson is a $(t\bar{b})$ 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 $(T\bar{b})$. Thus the model opens the possibility of some very interesting signal topologies which are quite nonstandard and can give surprisingly different and unique signals from the production of the heavy Majorana neutrino as well as charged the Higgs boson at the LHC.

IV. CONCLUSION

In this work we have proposed a model for SM fermion mass generation through a 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 vectorlike 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 the 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 τ final state decay modes. These bounds 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 be easily 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. We have considered two benchmark points in our model to highlight their signal strengths. In addition to an interesting signal from heavy neutrino production where one gets same-sign leptons and a same-sign top quark pair in the same event, the model also gives a very unique and different signal for the charged Higgs in the model. As no search has been performed at either ATLAS or CMS for such event topologies, the observation of such nonstandard signal events at the LHC could provide hints on new physics with an underlying model quite different from the popular left-right models.

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