Like-sign dileptons in the framework of the Manifest Left-Right Symmetric Model

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\begin{abstract}
The Left right symmetric model may present evidence of new physics at the LHC era. We use its framework to investigate the lepton number violating signal of like-sign dileptons and two jets at the 14 TeV LHC, i.e., \(pp \rightarrow e^\pm e^\pm jj + X\). We demonstrate that for an integrated luminosity of 300 fb\(^{-1}\), the right-handed boson \(W_R\) together with the right-handed electron neutrino \(N_e\) could be observed if their masses are smaller than 5.2 TeV and 3.6 TeV, respectively. For a doubly-charged Higgs \(\delta_{R}^{\pm\pm}\) of 500 GeV, for instance, the discovery range can be somewhat pressed to 5.3 TeV for \(W_R\) and 4 TeV for \(N_e\). We point out that the contribution of \(\delta_{R}^{\pm\pm}\) to the \(e^\pm e^\pm jj\) signal for a luminosity of 300 fb\(^{-1}\) can be detected for a doubly charged Higgs mass which is lower than \(M_{\delta_{R}^{\pm\pm}} \simeq 1.3\) TeV.

\textit{Keywords:} left-right model; like sign dileptons; Majorana neutrinos; doubly charged Higgs;
\end{abstract}

1. introduction

The main goal of the LHC is to search for signals of new physics beyond the Standard Model (SM), motivated by the shortcomings of the SM. In particular, The SM is incapable of explaining a number of fundamental issues, such as the hierarchy problem (resulting from the large difference between the weak force and the gravitational force), dark matter and the number of families in the quark and lepton sector. It is, therefore, widely believed that new physics beyond the SM will be discovered in the coming years. Among the possible attractive platforms for new physics are left-right symmetric models (LRSM)\cite{1, 2}.

The LRSM addresses two specific shortcomings of the SM: (i) Parity violation in the weak interactions, and (ii) non-zero neutrino masses implied by the experimental evidence of neutrino oscillation \cite{3}. In particular, the left-right symmetry which underlies LRSM restores Parity symmetry at energies appreciably higher than the electroweak scale, resulting in the addition of three new gauge boson fields, \(W_{R1,2,3}\). Furthermore, in LRSM the neutrinos are massive and their nature (i.e., whether they are of Majorana or

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Dirac type) depends on the details of the LRSM (for further discussion on this subject see [4]).

Early constructions of the LRSM comprise a Higgs sector with a Higgs bidoublet and two Higgs doublets [1]. In such a setup, the neutrinos are of Dirac type and no natural explanation for their small masses is provided. A later version, the manifest LRSM (MLRSM), incorporates a Higgs bidoublet and two Higgs triplets, which leads to Majorana type neutrinos [2]. The LRSM therefore provides a natural setup for the smallness of neutrino masses, relating their mass scale to the large left-right symmetry breaking scale through the (type I) see-saw mechanism [5].

An interesting case of direct detection of both parity breakdown and the see-saw mechanism at the LHC is the lepton number violating (LNV) signal (first pointed out by Keung and Senjanović [7])

\[ pp \rightarrow l^\pm l^\pm + 2j + X. \]  

(1.1)

Namely, two same sign leptons (like-sign dileptons) and two jets. This signal can be produced in a number of ways, in all of which a few-TeV \( W_R \) is produced and subsequently decays either via a right-handed Majorana neutrino or via a doubly charged Higgs and a \( W_R \). This signal allows one to

1. detect the right handed gauge boson \( W_R \),
2. trace the see-saw mechanism and the Majorana nature of the neutrino by detecting a heavy right handed neutrino, which is associated to the spontaneous left-right symmetry breaking scale by its heavy mass.
3. establish the existence of the charged Higgs bosons and further confirm the Higgs triplet nature.

We investigate this signal at the 14 TeV LHC in the framework of the MLRSM. We use for this purpose a computerized model implementation file developed by us in a former work [8]. Former studies of the like sign dilepton signal within the LRSM focused on a Drell-Yan production of a right handed neutrino and a lepton via an s-channel production of \( W_R \), followed by its decay through \( W_R \) to a second same sign lepton and two jets, i.e. \( u\bar{d} \rightarrow W_R^+ \rightarrow Nl^+ \) followed by \( N \rightarrow W_R^{-}l^+ \rightarrow jjl^+ \) (for positively charged leptons) [9, 10, 11]. In the present work, we extend these studies by including all possible diagrams and analysing the dominant contributions among the possible MLRSM amplitudes leading to this signature, including the doubly charged Higgs mediated channel \( u\bar{d} \rightarrow W_R^+ \rightarrow W_R^{-}\delta_R^{++} \rightarrow l^+l^+jj \). Indeed we find that within some range of the available LRSM parameters, the doubly-charged Higgs contribution can significantly contribute to the process \( pp \rightarrow e^\pm e^\pm jj + X \).

\[ \text{For a discussion about the see-saw mechanism (type I and II) in other models see e.g. [6].} \]

\[ \text{For a more general discussion about probing the LRSM Higgs sector at hadron colliders see e.g. [12].} \]
The work is organized as follows. In section 2 we describe the the LRSM Lagrangian structure. In section 3 we briefly summarize the constraints on the $W_R$ mass and the doubly charged Higgs, as well as on known Yukawa matrix elements. In section 4 we check, in the framework of the LRSM, the effects of the leading first order terms which contribute to electroweak precision quantities. We show that the radiative corrections calculated using the benchmark parameter taken in this work are small enough to remain within the current precision electroweak data. In section 5 we investigate (again, using the same benchmark parameter set) the like-sign dilepton plus two jets signal. We analyse in that section the signal versus the possible SM background, by reconstructing the heavy gauge boson $W_R^+$, the right handed electron neutrino $N_e$ and the doubly charged Higgs $\delta_{R}^{++}$, and examining the discovery potential of these particles. Finally, in section 6, we summarize.

2. General model description

The Lagrangian of the MLRSM at tree level can be divided into four terms:

$$\mathcal{L} = \mathcal{L}_{\text{kinetic}} + \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{Yukawa}} + \mathcal{L}_{\text{Higgs}}.$$  
(2.1)

The $\mathcal{L}_{\text{kinetic}}$ part contains the interactions between fermions and gauge bosons which are invariant under $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. In particular, the fermionic kinetic terms take the following form

$$L_f = i \sum \bar{\psi} \gamma^\mu D_\mu \psi$$

$$= \bar{L}_L \gamma^\mu \left( i \partial_\mu + g_L \frac{\sigma}{2} \cdot \vec{W}_L + g_R \frac{\sigma}{2} \cdot \vec{W}_R - \frac{g'_{B-L}}{2} B_\mu \right) L_L$$

$$+ \bar{L}_R \gamma^\mu \left( i \partial_\mu + g_R \frac{\sigma}{2} \cdot \vec{W}_R + g'_2 B_\mu \right) L_R$$

$$+ \bar{Q}_L^{\alpha \beta} \left[ \left( i \partial_\mu + g_L \frac{\sigma}{2} \cdot \vec{W}_L + \frac{g'_B}{6} B_\mu \right) \delta_{\alpha \beta} + \frac{g_s}{2} \lambda_{\alpha \beta} \cdot G_\mu \right] Q_L^\beta$$

$$+ \bar{Q}_R^{\alpha \beta} \left[ \left( i \partial_\mu + g_R \frac{\sigma}{2} \cdot \vec{W}_R + \frac{g'_B}{6} B_\mu \right) \delta_{\alpha \beta} + \frac{g_s}{2} \lambda_{\alpha \beta} \cdot G_\mu \right] Q_R^\beta,$$  
(2.2)

The appropriate coupling constants of the $G_\mu^a$, $\vec{W}_{L,R}\mu$ and the $B_\mu$ fields are $g_s$, $g_{L,R}$ and $g' = g_{B-L}$, respectively. The requirement that the Lagrangian is invariant under the left-right symmetry

$$\psi_L \leftrightarrow \psi_R, \quad \vec{W}_L \leftrightarrow \vec{W}_R,$$  
(2.3)

leads to

$$g_L = g_R.$$  
(2.4)
The gauge bosons kinetic terms and inner interactions are

\[ L_{\text{gauge}} = -\frac{1}{4} G_{\alpha}^{\mu\nu} G_{\alpha}^{\mu\nu} - \frac{1}{4} W_{L_{iL}}^{\mu\nu} W_{L_{iL}}^{\mu\nu} - \frac{1}{4} W_{R_{iR}}^{\mu\nu} W_{R_{iR}}^{\mu\nu} - \frac{1}{4} B_{\mu\nu} B_{\mu\nu}, \]  

where \( G_{\alpha}^{\mu\nu}, \ W_{L_{iL}}^{\mu\nu}, \ W_{R_{iR}}^{\mu\nu}, \) and \( B_{\mu\nu} \) are the field strength tensors of the \( SU(3)_C \), \( SU(2)_{L,R} \) gauge fields and the \( U(1)_{B-L} \) gauge field, respectively. They are defined as follows:

\[ G_{\alpha}^{\mu\nu} = \partial^{\mu} G_{\alpha}^{\nu} - \partial^{\nu} G_{\alpha}^{\mu} - g_s f_{abc} G_{b}^{\mu} G_{c}^{\nu}, \]  
\[ W_{L_{iL}}^{\mu\nu} = \partial^{\mu} W_{L_{iL}}^{\nu} - \partial^{\nu} W_{L_{iL}}^{\mu} + g_{L} \varepsilon^{ijk} W_{L_{jL}}^{\mu} W_{L_{kL}}^{\nu}, \]  
\[ W_{R_{iR}}^{\mu\nu} = \partial^{\mu} W_{R_{iR}}^{\nu} - \partial^{\nu} W_{R_{iR}}^{\mu} + g_{R} \varepsilon^{ijk} W_{R_{jL}}^{\mu} W_{R_{kL}}^{\nu}, \]  
\[ B_{\mu\nu} = \partial^{\mu} B^{\nu} - \partial^{\nu} B^{\mu}. \]  

The Yukawa interactions part, \( L_{\text{Yukawa}} \), consists of the most general possible couplings of the Higgs multiplets to bilinear fermion field products which form singlets under \( SU(2)_{L} \times SU(2)_{R} \times U(1)_{B-L} \):

\[ L_{\text{Yukawa}} = -\sum_{i,j} \left[ \bar{L}_{iL} \left( (h_{L})_{ij} \phi + (\tilde{h}_{L})_{ij} \tilde{\phi} \right) L_{jR} - \bar{Q}_{iL} \left( (h_{Q})_{ij} \phi + (\tilde{h}_{Q})_{ij} \tilde{\phi} \right) Q_{jR} \right. \]

\[ \left. - \left( L_{iL} \right)^{c} \Sigma_{M}(h_{M})_{ij} L_{jL} - \left( L_{iL} \right)^{c} \Sigma_{L}(h_{M})_{ij} L_{jL} \right] + h.c. \]  

where \( \tilde{\phi} = \sigma_{2} \phi^{*} \sigma_{2}, \Sigma_{L,R} = i \sigma_{2} \Delta_{L,R} \) and \( h_{Q}, \ h_{L}, h_{M}, \tilde{h}_{Q}, \tilde{h}_{L} \) are 3 × 3 Yukawa matrices in flavor space. The Higgs Lagrangian term consists of the Higgs kinetic terms and the potential of the Higgs multiplets:

\[ L_{\text{Higgs}} = \sum_{i} Tr\left| D_{\mu} \Theta_{i} \right|^{2} - V_{\text{Higgs}} \]  

where \( \Theta_{i} = \{ \phi, \Delta_{L}, \Delta_{R} \} \). As mentioned above, the covariant derivatives for the Higgs multiplets are given in the adjoint representation. Under the \( SU(2)_{L} \times SU(2)_{R} \times U(1)_{B-L} \) symmetry they are given by

\[ D_{\mu} \phi = \partial_{\mu} \phi - ig_{L} \left( \bar{\sigma} \cdot \tilde{W}_{L_{\mu}} \right) \phi + ig_{R} \left( \bar{\sigma} \cdot \tilde{W}_{R_{\mu}} \right) \]  
\[ D_{\mu} \Delta_{L,R} = \partial_{\mu} \Delta_{L,R} - ig_{L,R} \tilde{W}_{L,R_{\mu}} \cdot \left( \bar{\sigma}, \Delta_{L,R} \right) - ig'B_{\mu} \Delta_{L,R}. \]
neutral gauge bosons acquire masses through the kinetic terms in Eq.(2.8).

3. LRSM Constraints and parameter settings

The currently most stringent lower bound on the mass of $W_R$ is obtained from the $K_L - K_S$ mass difference, to be $M_{W_R} \geq 2.5 \text{ TeV}$ [13] (direct collider bounds on $M_{W_R}$, although competitive with this indirect one, are less stringent\(^3\)). As for the masses of the heavy neutrinos: these are constrained by vacuum stability, which leads to the following higher limit on the heavy Majorana neutrino masses [15]:

$$\left[ \sum_{i=e,\mu,\tau} M_{N_i}^4 \right]^{1/4} \lesssim 1.18 M_{W_R}. \quad (3.1)$$

If the right handed neutrinos of all three generations have the same mass, then from Eq.(3.1) $M_N \lesssim M_{W_R}$. Here we will always use $M_N < M_{W_R}$. There is also a lower limit on heavy neutrino masses derived from direct searches for pair production of neutral heavy leptons and a single production of excited neutral leptons. These searches result in a lower limit of 90 GeV for a heavy Majorana neutrino mass [16]\(^4\).

Let us briefly mention at this point the most stringent experimental limits on the right handed and left handed doubly charged Higgs masses\(^5\). These limits where obtained by direct searches for a signal of pair-produced ”left handed” states $\delta_{L}^{\pm\pm}$ and ”right handed” states $\delta_{R}^{\pm\pm}$, $pp \rightarrow \delta_{L,R}^{\pm\pm} \rightarrow l_1^\pm l_2^\pm l_3^\pm l_4^\pm$, in which an excess in like-sign dileptons was examined [19]. In particular, the currently most stringent bound on the mass of $\delta_{L}^{\pm\pm}$ is 551 GeV, assuming a 100% branching ratio to $e^\pm e^\pm$ pairs. For $\delta_{R}^{\pm\pm}$ the limit is 438 GeV and is obtained assuming a 100% branching ratio to $\mu^\pm \mu^\pm$ pairs. These bounds do not necessarily apply for the parameter choice in the analysis performed in this work. For instance, assuming a relatively more realistic scenario, in which $BR(\delta_{R}^{\pm\pm} \rightarrow \mu^\pm \mu^\pm) = \frac{1}{3}$ (as in the case of $M_{N_e} = M_{N_{\mu}} = M_{N_\tau}$ where the $\delta_{R}^{\pm\pm}$ decays evenly to $e^\pm e^\pm$, $\mu^\pm \mu^\pm$ and $\tau^\pm \tau^\pm$, see sec. 5\(^7\)), gives a more flexible bound of $M_{\delta_{R}^{\pm\pm}} > 320 \text{ GeV}$. Finally, we note that the corresponding lower bounds which come from the single production of a doubly charged Higgs is substantially lower: 141 GeV [20]\(^8\).

\(^3\)A search for excess in the SM background of two reconstructed leptons and at least one hadronic jet gives a lower limit of $M_{W_R} > 2.3 \text{ TeV}$ [14].

\(^4\)There is also an indirect lower bound derived from the neutrinoless double beta decay half-life of Ge, namely $M_{N_e} > 63 \text{ GeV} \left( \frac{1.6 \text{ TeV}}{M_{W_R}} \right)$ [17].

\(^5\)For a theoretical analysis of the LRSM Higgs mass spectrum see [18].

\(^6\)The terms ”left handed” and ”right handed” come from the chirality (left or right) of the weak isospin $T_3$ coupled to the doubly charged Higgs: $\delta_{L}^{\pm\pm}$ coupled to either $l_L^\pm l'^-_L$ or $l'^+_L l''_L$ ($(T_3)_L = \pm 1$), and $\delta_{R}^{\pm\pm}$ coupled to either $l''_R l'^-_R$ or $l'^+_R l''_R$ ($(T_3)_R = \pm 1$).

\(^7\)The doubly charged Higgs boson appearing in the signal examined in this work is right-handed. The relevant coupling to its left handed counterpart vanishes.
We now turn back to theoretical bounds on the electroweak scale Higgs VEVs. After SSB, the Higgs bidoublet VEVs $k_1$ and $k_2$ constitute the up- and down-type quark Yukawa mass terms, giving (see Eq. (2.7))

$$
M_U = k_1 h_L + k_2 \tilde{h}_L,
$$

$$
M_D = k_2 e^{i\alpha} h_L + k_1 \tilde{h}_L. \tag{3.2}
$$

Assuming there is no fine tuning involved, and considering the fact that the top quark is much heavier than the bottom quark, one can conclude that $k_1$ and $k_2$ are not of the same order (and neither are $h_L$ and $\tilde{h}_L$). We, therefore, use $k_1 \gg k_2$ (and $h_L \gg \tilde{h}_L$) (see also [22]). This setting also satisfies a theoretical bound on the charged gauge boson mixing angle which can be derived from the Schwarz inequality [23]:

$$
|\xi| \simeq \frac{|k_1 k_2|}{v_R^2} < \frac{M_{W_L}^2}{M_{W_R}^2}, \tag{3.3}
$$

which, for $m_{W_R} > 2.5 \text{ TeV}$ (see [13] and discussion above), implies $|\xi| < 0.00103$.$^9$

Let us turning now to the heavy non-SM neutral Higgs particles $H_0^1$ and $A_0^1$, for which $M_{H_0^1, A_0^1} \simeq \frac{1}{2} \alpha_3 v_R^2 \frac{k_1^2 + k_2^2}{k_1^2 - k_2^2}$, where $v_R$ is the VEV of the right Higgs triplet ($v_R \propto M_{W_R}$) and $\alpha_3$ is one of the Higgs potential parameters. The bound on these masses is derived from LRSM flavor changing neutral current contributions to Kaon mixing shown in Fig. 1. The corresponding Kaon mass splitting is given by

$$
\Delta M_K = 2\Re \left[ \left \langle K^0 \right \rangle \mathcal{H}_{\text{eff}} \left \langle K^0 \right \rangle \left \langle K^0 \right \rangle \mathcal{H}_{\text{eff}} \left \langle \overline{K}^0 \right \rangle \right ]^{1/2}. \tag{3.4}
$$

The observed mass splitting is, however, already generated by the SM box diagrams, so that the LRSM contributions (part of which contains the Higgs exchanges in Fig. 1)

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$^8$A concise review on single production of doubly charged Higgs in the LRSM framework can be found in [21].

$^9$A direct limit which is far less stringent comes from measurement of the $\xi$ parameter in $\mu$ decay (in order to search for deviations in the V-A theory), resulting in $\xi \leq 0.035$ [24].
should be controlled. The corresponding bounds were calculated in \[5\]:

\[ M_{H^0_{1},A^0_{1}} \gtrsim 10 \text{ TeV}. \]  

(3.5)

Applying this lower limit in the expression for the $H^0_1$ and $A^0_1$ mass terms (see above), leads to an interplay between the Higgs potential parameter $\alpha_3$ and the Higgs right triplet VEV $v_R$, namely

\[ \alpha_3 \geq \frac{2(k_1^2 - k_2^2)}{(k_1^2 + k_2^2)v_R^2} \cdot 100 \text{ TeV}^2. \]  

(3.6)

Setting for instance the Higgs bidoublet VEV to $k_1 \lesssim 246.22$, ($k_2 = \sqrt{246.22^2 - k_1^2}$) and the Higgs potential parameter $\alpha_3 = 3.6$, yields $M_{W_R} \geq 3.5 \text{ TeV}$. This setting is in fact used in section 5 to explore the like-sign dilepton signal in the LHC. It should be noted that a larger value of the (lower) bound on the heavy neutral Higgs masses would lead to a larger value of $\alpha_3$, if one requires $W_R$ to have a mass of order of a few TeV (i.e., within the reach of the LHC), and thus $\alpha_3$ may become non-perturbative. Based on our result in the next section, we take the maximum mass of an LHC-probed $W_R$ to be $\sim 5.3 \text{ TeV}$, which leads to a lower bound on $\alpha_3$ from Eq. (3.6): $\alpha_3 \gtrsim 1.57$, well within the perturbative regime.

Let us point out that taking into account the above mentioned experimental lower and higher bounds for the doubly charged Higgs $\delta^{\pm \pm}_R$ and for the right handed $W_R$, respectively, one can also obtain a lower bound on the potential parameter $\rho_2$ from the mass expression of the right handed doubly charged Higgs, given by (see \[25\]):

\[ M^2_{\delta^{\pm \pm}_R} = 2\rho_2 v_R^2 + \frac{1}{2} \alpha_3 (k_1^2 - k_2^2). \]  

(3.7)

A lower bound on $\rho_2$ is then given by

\[ (\rho_2)_{\text{min}} = \frac{1}{2 \left((v_R)_{\text{max}}\right)^2} \left[ \left(M_{\delta^{\pm \pm}_R}\right)_{\text{min}}^2 - \frac{1}{2} (k_1^2 - k_2^2)(\alpha_3)_{\text{max}}\right], \]  

(3.8)

where $(M_{W_R})_{\text{max}}$ and $(M_{\delta^{\pm \pm}_R})_{\text{min}}$ are the aforementioned experimental bounds\[10\]. In addition, one has to consider the upper bound $(\alpha_3)_{\text{max}}$ beyond which $\alpha_3$ is no longer within the perturbative regime. This is illustrated in Fig. 2, where we give a contour plot of the mass of the right handed doubly charged Higgs $\delta^{\pm \pm}_R$ in the $\alpha_3 - \rho_2$ plane (a similar plot for the left handed $\delta^{\pm \pm}_L$ is shown in \[27, 28\]). In the plot, we set $(\alpha_3)_{\text{max}} = 4.5$, which can be justified by observing that $\frac{(\alpha_3)_{\text{max}}}{4\pi} \sim O(1)$\[11\]. We have also shaded the

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\[9\]For a more general analysis of Higgs potential parameters see, e.g. \[26\], where RG equations are employed.

\[10\]As also suggested by Prof. R. Mohapatra in a private communication.
regions of the parameter space which are excluded by either FCNC constraints (i.e. \(M_{H^0_1, A^0_1} \gtrsim 10\) TeV), by direct LHC searches and by the combined requirement derived from Eq. (3.8)\(^{12}\).

**Figure 2:** The mass of \(\delta_R^{\pm \pm}\) (in TeV) as a function of the potential parameters \(\alpha_3\) and \(\rho_2\). The contour plot is divided into mass regions shown in the legend. The point marked by ★ corresponds to the benchmark parameter set chosen in this work (see Appendix A), taking \(M_{\delta_R^{\pm \pm}} = 500\) GeV and \(M_{W_R} = 3.5\) TeV. The shaded regions are excluded from FCNC constraints (the region under the horizontal dashed red line), from direct LHC searches (the region bound left to the dashed red curve) and due to the combined requirement from Eq. (3.8) (the region left to the vertical dashed red line).

Finally, there are also constraints on the Yukawa couplings of the doubly charged Higgs to leptons. The relevant Yukawa interactions are given by the following Lagrangian terms:

\[
L_Y = - \sum_{i,j} \left[ (L^c_i R) R_i (h_M)^{ij} L_{jR} + (L^c_i L) L_i (h_M)^{ij} L_{jL} \right] + h.c. \tag{3.9}
\]

The Yukawa matrix \(h_M\) can be expressed in terms of the neutrino mass matrices and

\(^{12}\)A possible alternative to the above scenario is to take \((M_{H^0_1, A^0_1})_{\text{min}} = 14\) TeV. This will result in elevating the \(\alpha_3\) exclusion bound to 3. In that case, setting e.g. \(\alpha_3 = 5\) (i.e. still within the perturbative regime) will allow for \(M_{W_R} \geq 4.1\) TeV, thus meeting the Higgs bound in [29] for the relevant range of \(M_{W_R}\).
CKM type mixing matrices (assuming, for simplicity, a manifest LR symmetric model):

\[ h_M = \frac{1}{\sqrt{2}v_R} K_R^T M_{\text{diag}}^\nu K_R, \]  

where \( K_R \) is a CKM-type mixing matrix in the lepton-sector and \( M_{\text{diag}}^\nu \) is the 6 \( \times \) 6 neutrino diagonal mass matrix (see [25]). The constraints on the matrix elements of \( h_M \) originate from different experimental processes, as shown in Table 1 [30].

| Constraint source | Constraint |
|-------------------|------------|
| \( \mu \to eee \) | \( h_{Me\mu} \cdot h_{Mee} \lesssim 3.2 \times 10^{-11} \text{GeV}^{-2} \cdot M_{H_R^{++}}^2 \) |
| Bhabha scattering | \( h_{Mee}^2 \lesssim 9.7 \times 10^{-6} \text{GeV}^{-2} \cdot M_{H_R^{++}}^2 \) |
| Extra coupling to \((g - 2)\mu\) | \( h_{M\mu\mu}^2 \lesssim 2.5 \times 10^{-5} \text{GeV}^{-2} \cdot M_{H_R^{++}}^2 \) |
| Muonium \((\mu^+e^-)\) transformation to anti-muonium | \( h_{M\mu\mu}^2 \lesssim 5.8 \times 10^{-5} \text{GeV}^{-2} \cdot M_{H_R^{++}}^2 \) |
| Non-observation of \( \mu \to e\gamma \) decay | \( h_{Me\mu} \cdot h_{M\mu\mu} \lesssim 2 \times 10^{-10} \text{GeV}^{-2} \cdot M_{H_R^{++}}^2 \) |
| Vacuum stability | \( h_{Mee}, h_{M\mu\mu} \lesssim 1.2 \) |

Table 1: Experimental constraints on the Yukawa couplings of the doubly charged Higgs to leptons.

All of the constraints discussed in this section are implemented in the benchmark parameter set of the present work (see Appendix A), and are included in the cross section plots.

4. Oblique corrections and precision measurements in the LRSM

Apart from the fermion masses, the CKM-mixing angles and the Higgs mass, the EW sector of the SM has three fundamental parameters: the gauge coupling constants \( g \) and \( g' \) and the Higgs VEV \( v \). These parameters are conventionally replaced by parameters which can be directly measured in physical processes, where a specific choice of experimentally measured input parameters defines a normalization scheme.

\[ ^{13} \text{For a more recent discussion on neutrino mixings, see [31].} \]
One natural choice, replacing $g$, $g'$ and $v$, would be the fine structure constant and the gauge-boson masses $\alpha$, $M_W$ and $M_Z$ (i.e., disregarding for now the fermion and SM Higgs masses $m_f$ and $m_H$). The Fermi constant can then be calculated in terms of $\alpha$, $M_W$ and $M_Z$, yielding at tree-level:

$$G_F = \frac{\pi \alpha}{\sqrt{2} \left( 1 - \frac{M_W^2}{M_Z^2} \right) M_W^2},$$

(4.1)

where all the above quantities had been measured with extreme accurately $^{[32]}$:

$$G_F = 1.1663787(6) \times 10^{-5} \text{ GeV}^{-2},$$
$$\alpha \equiv \frac{e^2}{4\pi} = 1/137.035999074(44),$$
$$M_W \equiv 80.315 \pm 0.015 \text{ GeV},$$
$$M_Z \equiv 91.1876 \pm 0.0021 \text{ GeV}.$$

(4.2)

When radiative corrections are included, Eq. 4.1 is modified as follows (see e.g., $^{[33]}$):

$$\sqrt{2} G_F M_W^2 \left( 1 - \frac{M_W^2}{\rho M_Z^2} \right) = \pi \frac{\alpha}{1 - \Delta r},$$

(4.3)

where $\rho$ is fixed to its tree-level value, $\rho_{\text{tree}} = 1$, such that $\sin^2 \Theta_W = 1 - \frac{M_W^2}{M_Z^2}$, which is the on-shell definition of the EW mixing angle. Thus, Eq. 4.3 provides a useful relation between the EW gauge-boson masses, the fine structure constant and the Fermi constant, which by itself is linked to the muon lifetime.

The EW precision parameter $\Delta r$ (first calculated by Sirlin $^{[34]}$), collects the quantum corrections to the muon decay process and it plays an important role in placing bounds or searching for new physics that couples to the SM fields. In particular, it is given by:

$$\Delta r \equiv \frac{\hat{\Sigma}_W(0)}{M_W^2} + \Delta r_{\text{vert,box}},$$

(4.4)

where $\hat{\Sigma}_W(k^2)$ is the on-shell renormalized W-boson self-energy, which accounts for the universal ("oblique") part of the EW radiative corrections to the muon decay. The corrections which are non-universal (i.e., process-dependent) constitute the non-leading term $\Delta r_{\text{vert,box}}$. In particular, the experimental values for $M_Z$, $M_W$, $G_F$ and $\alpha$ can be used as an input for the evaluation of the experimental value $\Delta r_{\text{exp}}$. Plugging the
corresponding central values (shown in Eq. (4.2)) into Eq. (4.3), we obtain\(^{14}\):

\[
\Delta r_{\text{exp}} = \frac{\sqrt{2} G_F}{\pi \alpha} M_W^2 \left( 1 - \frac{M_W^2}{M_Z^2} \right) - 1 = 36.322 \times 10^{-3}. \tag{4.5}
\]

Another useful precision parameter, known as the \(\delta \rho\)-parameter, which is also very accurately measured \((\rho = \frac{G_{\text{NC}}}{G_F} = 1 + \delta \rho)\), is defined through the ratio between the neutral and charged weak currents as \(^{[33]}\):

\[
\delta \rho = \frac{\Sigma Z(0)}{M_Z^2} - \frac{\Sigma W(0)}{M_W^2}, \tag{4.6}
\]

so that \(\delta \rho \neq 0\) is caused by mass splitting between partners of a given weak isospin doublet, therefore, tracing the degree of departure from the global custodial SU(2) invariance of the SM Lagrangian. In particular, in the SM it is dominated by the top-quark loops, giving \(^{[35]}\):

\[
\delta \rho_{\text{SM}}^{[t]} = \frac{3 G_F m_t^2}{8 \sqrt{2} \pi^2}. \tag{4.7}
\]

Unlike the above quadratic top quark mass contribution to \(\delta \rho\), the Higgs contribution to \(\delta \rho\) in the SM is "screened" leading to a milder logarithmic behaviour\(^{15}\) \(^{[36]}\):

\[
\delta \rho_{\text{SM}}^{[H]} \approx -\frac{3 \sqrt{2} G_F M_W^2}{16 \pi^2} \frac{s_W^2}{c_W^2} \left\{ \ln \frac{M_H^2}{M_W^2} - \frac{5}{6} \right\} + ... \tag{4.8}
\]

where \(s_W(c_W) \equiv \sin \Theta_W(\cos \Theta_W)\). This logarithmic behaviour of the SM Higgs contribution to \(\delta \rho\) does not hold when extended Higgs sectors are involved, as will be shortly shown for our case of the LRSM \(^{16}\).

\(^{14}\)Since the quantities \(G_F\) and \(\alpha\) are known to extreme accuracy, measuring \(\Delta r\) and its uncertainty are equivalent to measuring the masses and uncertainties of the \(W\) and \(Z\) bosons.

\(^{15}\)It should be noted that the Higgs boson contribution to \(\delta \rho\) in Eq. 4.8 is not gauge invariant on its own and should be combined with the remaining bosonic contributions.

\(^{16}\)The \(\delta \rho_{\text{H}}\) term, although small in comparison to \(\delta \rho_{\text{exp}}\), has importance in measuring the deviation from custodial symmetry. The custodial symmetry in the Higgs sector corresponds to the limit \(g' \to 0\), in which \(W^+ W^- \) and \(Z\) form a triplet of an unbroken global symmetry, resulting from the following symmetry breaking pattern: \(SU(2)_L \times SU(2)_R \to SU(2)_{L+R}\). Beyond tree level, this remaining symmetry means that radiative corrections to the \(\rho\) parameter \((\rho = \frac{G_{\text{NC}}}{G_F} = \frac{M_W^2}{M_Z^2})\) as a result of the gauge and Higgs bosons must be proportional to \(g'^2 (M_Z^2 - M_W^2 \propto g'^2, s_W^2 \to 0 \text{ as } g' \to 0)\). Thus, the correction terms of \(\delta \rho_{\text{H}}\) in Eq. 4.8 have to be either small or disappear in the limit \(g' \to 0\), and cannot contain a quadratic mass term of the Higgs. Protecting the \(\rho\) parameter from large radiative corrections is known as screening, which is a feature of the SM Higgs boson \(^{[35]}\). In other models with Higgs sectors comprising also other entities except doublets, however, this property does not necessarily hold \(^{[33]}\).
It is, therefore, useful to recast the precision parameter $\Delta r$, so that its dependence on $\delta \rho$ is explicitly manifest [37]:

$$\Delta r = \Delta \alpha - \frac{c_w^2}{s_w^2} \delta \rho + \Delta r_{\text{rem}},$$

(4.9)

where $\Delta \alpha$ accounts for the (leading) light-fermion logarithmic corrections to the photon vacuum polarization (therefore unchanged by new heavy physics), given by [33]:

$$\Delta \alpha = \frac{\alpha}{3\pi} \sum_f Q_f^2 N_{cf} \left( \ln \frac{M_Z^2}{m_f^2} - \frac{5}{3} \right) \simeq 0.06637.$$  

(4.10)

$\Delta r_{\text{rem}}$ is the so-called "remainder" term which contains the remaining (typically smaller) contributions. In particular, in the SM $\Delta r_{\text{rem}} \simeq 0.01$, whereas the top quark contribution to $\Delta r$ from $\delta \rho$ is $\Delta r(\text{top}, \text{SM}) = -\frac{c_w^2}{s_w^2} \delta \rho_{\text{SM}} \approx -0.04$. The SM Higgs contribution to $\Delta r$ comes from $\delta \rho_{\text{SM}}^{[H]}$ in Eq. 4.8 and from a similarly structured term from $\Delta r_{\text{rem}}$, giving [36]:

$$\Delta r^H \approx \frac{\sqrt{2} G_F M_W^2}{16 \pi^2} \left\{ \ln \frac{M_Z^2}{M_W^2} - \frac{5}{6} \right\} + \ldots.$$  

(4.11)

The one-loop heavy fermion contribution to $\Delta r$, is dominated by the top-quark contribution to $\delta \rho$, i.e., $\delta \rho_{\text{SM}}^{[t]}$ in Eq. 4.7 (due to the large mass splitting within the top and bottom quark SU(2) doublet), $\Delta r$ also contains non-negligible logarithmic terms from the remainder term, $\Delta r_{\text{rem}}$, giving in total [36]:

$$\Delta r^{\text{top}} = -\frac{\sqrt{2} G_F M_W^2}{16 \pi^2} \left[ 3 \frac{c_w^2}{s_w^2} \frac{m_t^2}{M_W^2} + 2 \left( \frac{c_w^2}{s_w^2} - \frac{1}{3} \right) \ln \frac{m_t^2}{M_W^2} + 4 \ln \frac{c_w^2}{s_w^2} + \frac{7}{9} \right]$$

$$\approx -10.431 G_F m_t^2/8 \sqrt{2} \pi^2 = -0.0327 (m_t/173.24 \text{GeV})^2.$$  

(4.12)

Turning now to the LRSM, one can use a renormalization procedure analogous to the SM, in which (apart from the fermion masses, the CKM-like mixing parameters and the Higgs parameters), the following set of physical parameters is used as an input [38, 39, 40, 41, 42]:

$$e, M_W, M_{W_R}, M_Z, M_{Z_L},$$

(4.13)

so that, similar to the SM case, $\Delta r_{\text{LR}}$ is also extracted from Eq. 4.3 with $\Delta r \to \Delta r_{\text{LR}}$. It should be noted that, within the LRSM, the tree level contribution to the muon decay which involves $W_R$ can be neglected [43].

\[17\] It should be noted that, within the LRSM, the tree level contribution to the muon decay which involves $W_R$ can be neglected [43].
In particular, in the LRSM we have:

$$\Delta r_{LR} = \Delta \alpha + \Delta r_{LR, heavy} + \Delta r_{LR, rem},$$

(4.14)

where $\Delta \alpha$ remains unchanged and is given in Eq. 4.10, $\Delta r_{LR, rem}$ is the ”remainder” term in the LRSM and $\Delta r_{LR, heavy}$ contains the dominant (oblique) contribution to $\Delta r_{LR}$ from the heavy particles in the theory. In particular, the 1-loop oblique corrections from the top-quark, heavy neutrinos and Higgs particles were calculated in [38, 39, 40, 41]:

$$\left(\Delta r\right)_{LR}^{\text{top}} = -\frac{\sqrt{2} G_F c_W^2}{8\pi^2} \left(\frac{c_W^2}{s_W^2} - 1\right) \frac{M_W^2}{M_{WR}^2} - M_W^2 - 3m_t^2,$$

(4.15)

$$\left(\Delta r\right)_{LR}^{N} = \sum_{N=\text{heavy}} \frac{\sqrt{2} G_F}{16\pi^2} (1 - 2s_W^2) \frac{M_W^2}{M_{WR}^2} - M_W^2 - m_N^2,$$

(4.16)

$$\left(\Delta r\right)_{LR}^{\text{Higgs}} = -\frac{\sqrt{2} G_F}{48\pi^2} \left(\frac{M_W^2}{M_{WR}^2} + \frac{1}{s_W^2} (4c_W^2 - 1)\right) M_{\text{Higgs}}^2,$$

(4.17)

A few comments are in order regarding the above $\Delta r_{LR}$ terms:

- The top-quark contribution in LRSM, $\left(\Delta r\right)_{LR}^{\text{top}}$, no longer exhibits the dominant SM’s quadratic behaviour, since it is suppressed by $M_W^2/M_{WR}^2$. In particular, even for moderate values of $M_{WR} \sim 500$ GeV, the top contribution, $\left(\Delta r\right)_{LR}^{\text{top}}$, is smaller than the SM logarithmic terms.

- The contributions from the Higgs particles in the theory is no longer screened by the custodial symmetry of the SM, so that a ”dangerous” quadratic dependence on the Higgs masses emerges in LRSM, in particular, since low-energy flavor physics impose stringent bounds on the heavy Higgs sector, e.g., $M_{H_1^0} \gtrsim 10$ TeV (see previous section). Nonetheless, here also, the Higgs contributions are suppressed by factors of $O(\frac{v_{EW}}{v_R})^2$, keeping them consistent with PEWD (see below).

- The heavy neutrino contribution, $\left(\Delta r\right)_{LR}^{N}$, is quadratic in the heavy neutrino masses, but also suffers from a $M_W^2/M_{WR}^2$ suppression. For $M_N \sim M_{WR} \sim O(\text{TeV})$, it is still much smaller than the SM top-quark term.

In Fig. 3 we plot $\Delta r_{LR}^{(1)}$, which consists of the sum of the dominant one-loop contributions in Eq. 4.14. The uncertainty here lies in $\Delta r_{LR, rem}$, which contains corrections from vertex and box diagrams that are expected to be sub-leading compared to the oblique corrections. Since the calculation of these type of 1-loop corrections is beyond the scope of this work and since, to the best of our knowledge, these terms were not calculated before, we will assume that $\Delta r_{LR, rem}$ is an order of magnitude smaller than the oblique corrections, setting $\Delta r_{LR, rem} = 0.01$, i.e., similar in size to the SM remainder terms,
Figure 3: Evaluation of the leading contributions to $\Delta r$ in the LRSM (solid line, in black), as a function of the mass ratio between the LRSM heavy-neutrino and the heavy charged gauge boson ($W_R$). The calculation is performed for two cases: i) the masses of the three heavy Majorana neutrinos are equal, $M_{N_e} = M_{N_\mu} = M_{N_\tau}$, and ii) the heavy electron neutrino is heavier than the other heavy neutrinos: $M_{N_\mu} = M_{N_\tau} = 0.2 M_{N_e}$, and using the LRSM parameter set given in Appendix A. We also give the SM prediction $\Delta r_{SM}$ (dashed line, in blue) and the experimental value $\Delta r_{exp}$ (dashed line, in red). The $1\sigma$ and $2\sigma$ C.L. regions are explicitly indicated.

$\Delta r_{rem} \approx 0.01$. In this respect, we note that an $O(0.01)$ shift to $\Delta r_{LR}$ does not change the main results of this paper. Results for $\Delta r_{LR}$ in Fig. 3 are displayed alongside the (one loop) SM value $\Delta r_{SM}$ and the experimental value $\Delta r_{exp}$ as obtained from Eq.(4.5). Confidence levels for $\Delta r_{exp}$ are calculated from experimental uncertainties of the input parameters (i.e., $\delta M_W$, $\delta M_Z$, etc...) and are also included in the plot; corresponding to $1\sigma$ (yellow region) and $2\sigma$ (green region).

We see that $\Delta r_{LR}$ vary in a smooth and monotonous manner, over an $O(10^{-3})$ range in the heavy (electron) neutrino mass range $0 \leq M_{N_e} \leq 1$. As mentioned above, the top quark contribution (Eq.(4.15)) is small (of $O(10^{-5})$) due to the dependence on LR symmetry breaking scale ($W_R$) in the denominator. The heavy Higgs boson contribution (Eq.(4.17)) contains the LR scale both in the numerator and in the denominator (through mass dependency) and is, therefore, roughly flat with respect to $M_{W_R}$. For our benchmark parameter set, it gives to the dominant term in $\Delta r_{LR}$ with a sign opposite to $\Delta \alpha$. 

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The evaluation of $\Delta r$ can also be translated into a theoretical prediction for the W-boson mass, and can, therefore, be used to further confront the experimental result. This can be done by solving Eq. 4.3 for $M_W^2$:

$$M_W^2 = \frac{1}{2} M_Z^2 \left[ 1 + \sqrt{1 - \frac{4\pi\alpha}{\sqrt{2} G_F M_Z^2}} (1 + \Delta r(M_{Z, Z_2}, M_{W, W_2}, s_W, m_t, M_{\text{Higgs}})) \right].$$ (4.18)

Figure 4: $M_W$ as a function of $s^2_W$ including the one-loop corrections in the LRSM, using Eq. 4.18. The yellow (green) shaded region represents the 1σ (2σ) error for the experimental data of $M_{\text{exp}}^W$ and $s_{\text{exp}}^2$. We show the results for our main benchmark parameter set (BPS), used throughout this work (see Appendix A) as an example of representative Higgs spectra. We also included in Fig. 4 two additional BPS’s, BPS1 and BPS2, with slightly larger Higgs masses in order to demonstrate the sensitivity of this analysis to the masses the Higgs particles in the model. In particular, set BPS1
corresponds to the following values for potential parameters: $\rho_3 = 2.8$ and $\rho_1 = 1.3$, thus yielding larger masses for the neutral Higgs particles $H_2^0$, $H_3^0$ and $A_2^0$ as well as for the charged Higgs particles $H_1^\pm$ and $\delta_L^\pm$ (i.e., $\sim 14.8$ TeV instead of $\sim 12.8$ TeV in our main BPS for $H_2^0$, and $\sim 1.5$ TeV instead of $\sim 0.7$ TeV in our main BPS for the other particles), while keeping the rest of the Higgs masses unchanged. The second subsidiary set, BPS2, corresponds to: $\alpha_3 = 4.2$, which yields larger masses for the neutral Higgs particles $H_1^0$ and $A_1^0$ and for the charged Higgs $H_2^\pm$ ($\sim 14$ TeV instead of 13 TeV in the main BPS)\(^\text{18}\). The increased Higgs masses are associated with larger weak mixing angle, which may imply the necessity to fine-tune Higgs potential parameters. A possible way out of this undesired situation is to look for an underlying higher symmetry (e.g., GUT or SUSY) beyond the context of the LRSM (see also discussion in [5]), that would more naturally restrict the Higgs parameters.

Finally, let us use the formalism of the oblique parameters $S, T$ and $U$ to examine whether the LRSM with our benchmark parameter set remains within the limit of the PEWD constraints. In this formalism (introduced by Peskin and Takeuchi [44]) we use the above one-loop precision quantities as input. Interestingly, this can be done in a non-direct manner by using yet another alternative set of parameters, the $\varepsilon_{1,2,3}$ parameters, which are given by [45]:

\[
\begin{align*}
\varepsilon_1 &= \delta \rho , \\
\varepsilon_2 &= c_W^2 \Delta \rho + \frac{s_W^2 \Delta r}{c_W^2 - s_W^2} - 2 s_W^2 \Delta k , \\
\varepsilon_3 &= c_W^2 \delta \rho + (c_W^2 - s_W^2) \Delta k ,
\end{align*}
\] (4.19)

where $\Delta k$ relates the precision observable $\sin^2 \theta_{\text{eff}}$ to the $s_0^2$ parameter:

\[
\sin^2 \theta_{\text{eff}} = \frac{1}{1 - \Delta k} s_0^2 ,
\] (4.20)

with

\[
s_0^2 c_0^2 = \frac{\pi \alpha (m_Z)}{\sqrt{2} G_F m_Z^2} .
\] (4.21)

The $S, T$ and $U$ parameters are then defined such that the SM contributions are subtracted from the $\varepsilon$ parameters, and hence are denoted as $S_{\text{new}}, T_{\text{new}}$ and $U_{\text{new}}$, re-

\(^{18}\)The increase of $\alpha_3$ to 4.2 also leads to a larger mass of the right handed doubly charged Higgs $\delta_R^\pm$, which increases by $\sim 3\%$
respectively. They are given by [46]:

\[
S_{\text{new}} = \varepsilon_3 \frac{4 \sin^2 \Theta_G}{\alpha(M_Z^2)} - d_S,
\]

\[
T_{\text{new}} = \varepsilon_1 \frac{1}{\alpha(M_Z^2)} - d_T,
\]

\[
U_{\text{new}} = -\varepsilon_2 \frac{4 \sin^2 \Theta_G}{\alpha(M_Z^2)} - d_U,
\]

(4.22)

where

\[
2 \sin^2 \Theta_G = 1 - \sqrt{1 - \sqrt{8\pi \alpha(M_Z^2)/(G_F M_Z^2)}},
\]

(4.23)

and \(d_i\) are the SM prediction for a chosen \(M_H\) and \(m_t\) values. The \(S_{\text{new}}, T_{\text{new}}, U_{\text{new}}\) parameters vanish for \(\text{SM}_{\text{ref}}\), which is defined for \(m_t = 173\) GeV and \(M_H = 125\) GeV\(^{19}\), and thus measure deviations from the chosen \(\text{SM}_{\text{ref}}\).

An electroweak observable is thus given by the sum of its \(\text{SM}_{\text{ref}}\) projection and new physics contributions parametrized by \(S, T, U\), as follows:

\[
O = O_{\text{SM,ref}}(M_H, m_t) + c_S S_{\text{new}} + c_T T_{\text{new}} + c_U U_{\text{new}}.
\]

(4.24)

In Fig. 5 we plot \((S_{\text{new}}, T_{\text{new}})\) in the LRSM for various values of the mass ratio \(M_{N_e}/M_{W_R}\) in the aforementioned two configurations, namely \(M_{N_e} = M_{N_\mu} = M_{N_\tau}\) and \(M_{N_\mu} = M_{N_\tau} = 0.2 M_{N_e}\). In general, the effects of the LRSM enter through the one-loop corrections in \(\Delta r_{\text{LR}}\) (see Eq. (4.14) where, as explained above, \(\Delta r_{\text{LR,rem}}\) is taken as 0.01) and through \(\delta \rho_{\text{LR}}\), which can be calculated using Eq. (4.9). Following the recipe of Eq. (4.22), we subtract the predictions of the SM from their LRSM counterparts\(^{20}\), and compare the results with the experimental fit. We obtain that, for the current parameter setting (see Appendix A), the LRSM prediction for the \(0 \leq M_{N_e}/M_{W_R} \leq 1\) range is in agreement with data in both the above cases.

5. The signal

In this section we will examine the production of the like-sign dilepton signal \(\ell^\pm \ell^\pm jj\) at the LHC, where for definiteness we will set the charged leptons \(\ell^\pm\) to be an electron and positron. Although the expected cross section is similar for the three lepton generation pairs (as opposed to the background which is unnecessarily the same), we choose to deal with electrons and positrons in order to correspond with former works related to this signal, see \([9, 10]\). For simplicity we also disregard mixing in the lepton sector, in

\(^{19}\)The current experimental measurements give \(S = 0.06 \pm 0.11, T = 0.09 \pm 13, \ U = 0.01 \pm 0.11\)\(^{[47]}\).

\(^{20}\)The relation between the one-loop input quantities \(\Delta r\) and \(\delta \rho\) is given in Eq. (4.9).
agreement with a negligible intergenerational heavy neutrino mixing which was derived from the ratio between the lepton flavor changing processes $\mu \rightarrow e\gamma$ and $\mu \rightarrow e\nu\bar{\nu}_e$ [7]. The like-sign lepton production proceeds through several channels, where in general they can be divided into s-channel processes and t-channel processes. In each channel there are two topological groups, which are characterized by processes mediated by a Majorana neutrino exchange and processes mediated by a doubly charged Higgs, respectively. The corresponding diagrams are depicted in Figs. 6 and 7.\footnote{For the sake of simplicity and without loss of generality, we will address the two positron channel, although our discussion applies for the sign-reversal channels leading to two electrons as well.}

Within the range of masses used\footnote{The dominant contributions to the processes considered arise from the diagrams with the right handed gauge boson $W_R$, the three (heavy) right handed neutrinos $N_{e,\mu,\tau}$ and the "right handed" doubly charged Higgs $\delta^{\pm\pm}_R$. The diagrams involving other relevant particles, such as the "left handed" doubly charged Higgs $\delta^{\pm\pm}_L$ or the singly charged Higgs $H^{\pm}_R$ are negligible since these particles have either zero coupling, as will be explained, or are much heavier. The relevant parameters of the LRSM (see Appendix A) are set subject to the constraints detailed in sec.3. In particular, the mass ranges considered are: $W_R$ of a few TeV, right handed neutrinos $N_{e,\mu,\tau} < W_R$ and a doubly charged Higgs $\delta^{++}_R$ lighter than 1.5 TeV. We choose to scenarios for the heavy neutrinos mass spectrum: one in which all three heavy neutrinos are degenerate and another one, in which $N_\mu$ and $N_\tau$ are either lighter or heavier than $N_e$.}, we find that the contribution from t-channel
Figure 6: The s-channel processes leading to like-sign dilepton signature. In the diagram, two lines with common endpoints represent two alternative possible particles, each with a different diagrammatic representation (if the two alternative particles have the same diagrammatic representation their names are separated by a comma). Also, the general $h^\pm$ field represents possible singly charged Higgs gauge eigenstates (i.e. $\phi^+_{1,2}$ or $\delta^+_R$, see Appendix B).

Figure 7: The t-channel processes leading to like-sign dilepton signature. See also caption to figure 6.
diagrams is negligible, since the s-channel diagrams (as opposed to the t-channel one) can be divided into the production and the decay of the on-shell neutrino and doubly charged Higgs, respectively.

5.1. Doubly charged Higgs mediated diagrams

Focusing on the s-channel diagrams which, as mentioned above, give the dominating contributions to \( \sigma(pp \rightarrow e^+e^+jj) \), we start by studying the possible diagrams in the doubly-charged Higgs mediated processes, examining first the channel which contains an s-channel \( W_R^+ \) "decaying" into \( W_R^- \delta_R^{++} \):

\[
(a) \quad u \bar{d} \rightarrow W_R^+ \rightarrow W_R^- \delta_R^{++} \rightarrow e^+e^+jj
\]

The \( W_{L,R} \) coupling to fermions is governed by the covariant derivative in the fermion-gauge interaction term in the Lagrangian

\[
L_f = \sum_{\Psi=(Q),(L)} \bar{\Psi} L \gamma^\mu \left( \partial_\mu - ig_L \frac{\vec{\sigma}}{2} W_{L\mu} - ig'Y B_\mu \right) \Psi_L + (L \rightarrow R). \tag{5.1}
\]

The left and right electroweak gauge couplings of the MLRSM considered here are equal, \( g_L = g_R \equiv g \), leading to a significant production rate of the SM \( W_L \) and the new heavy \( W_R \) at the 14 TeV LHC. For instance, at an integrated luminosity of \( \mathcal{O}(100) \) fb\(^{-1}\) a rate of \( 10^9 - 10^{10} \) on-shell \( W_L \) and \( \mathcal{O}(10^4) \) on-shell \( W_R \) (assuming \( W_R = 3 \) TeV) are expected to be produced [48]. The coupling \( W_R^- W_R^- \delta_R^{++} \) originates from the Higgs kinetic term \( (D_\mu \Delta_R)^\dagger D^\mu \Delta_R \), and is given by

\[
\frac{1}{\sqrt{2}} g^2 v_R W_R^- W_R^- \delta_R^{++}. \tag{5.2}
\]

Since we take the symmetry breaking scale, \( v_R \), to be a few TeV (so that \( W_R \) is light enough to be observable at the LHC), the above coupling is significant. The corresponding cross section is, however also suppressed by inverse powers of the \( W_R \) mass since it is exchanged via the s-channel.

Moving on to examine the relative partial widths and the branching ratios of the decay of the produced \( \delta_R^{++} \) into two positrons, the coupling \( \delta_R^{++} l^-l^- \) is given by

\[
\delta_R^{++} l^- l^- \ L_{MlR}^l. \tag{5.3}
\]

23The contributions stemming from the diagonal terms\(^{24}\) for each t-channel diagram are at best of order of \( 10^{-8} \) fb, which is too small to be detected at the LHC. In the same manner, the non-diagonal terms (i.e. the interference terms) of the t-channel diagrams were found to give negligible contributions.

24For example, in a cross section composed of two possible amplitudes/diagrams \( \sigma \propto |M_1 + M_2|^2 \), the contributions from the diagonal terms are \( \sigma_{\text{diag}} \propto |M_1|^2 + |M_2|^2 \).
where the Yukawa matrix $h_M$ is (recalling the $6 \times 3$ CKM-type mixing matrix of the lepton-sector $K_R$ and the $6 \times 6$ neutrino diagonal mass matrix $M_\nu$, see for example [25]):

$$h_M = \frac{1}{\sqrt{2}v_R} K_R^T M_\nu^{\text{diag}} K_R. \quad (5.4)$$

Eq.(5.4) contains a ratio between heavy neutrino masses and the mass scale of the heavy gauge bosons. For instance, taking for simplicity a 2-generation basis, $(\nu_e \nu_\mu, N_e N_\mu)$ (where $\nu$ and $N$ are light and heavy neutrino, respectively) and neglecting neutrino intergeneration mixing, one obtains for the Yukawa electron-electron (or positron-positron) coupling

$$h_{M_{ee}} = \frac{1}{\sqrt{2}v_R} \left(K_{R\nu e}^2 M_{\nu e} + K_{R\nu \mu}^2 M_{\nu \mu} + K_{R N_e e}^2 M_{N e} + K_{R N_\mu e}^2 M_{N \mu}\right) \simeq \frac{K_{R N_e e}^2 M_{N e}}{\sqrt{2}v_R}, \quad (5.5)$$

and similarly,

$$h_{M_{\mu\mu}} \simeq \frac{K_{R N_\mu \mu}^2 M_{N \mu}}{\sqrt{2}v_R},$$

$$h_{M_{e\mu}} \approx 0. \quad (5.6)$$

Assuming negligible mixing ($K_{R N_e e} \simeq K_{R N_\mu \mu}$) and for instance $M_{N e} = M_{N_\mu} = M_{N_\tau}$, this gives an equal branching ratio to flavor diagonal lepton pairs. This is to be compared with $\delta^{++} R \rightarrow W^+_L W^+_L$ which can be implied from Eq.(5.2). In particular, replacing $W_R$ with $W_L$ forces one to suppress the coupling of the vertex by multiplying it with the squared $W_L - W_R$ small mixing factor $\xi$:

$$\tan 2\xi = -\frac{2k_1 k_2}{v_R^2}. \quad (5.7)$$

Now, the fact that $v_R \propto M_{W_R}, M_{Z'}$ implies that

$$v_R \gg k_1, k_2, \quad (5.8)$$

and therefore $\Gamma(\delta^{++}_R \rightarrow l^+ l^+) > > \Gamma(\delta^{++}_R \rightarrow W^+_L W^+_L)$. In the case of equal heavy neutrino masses: $M_{N e} = M_{N_\mu} = M_{N_\tau}$, we have $BR(\delta^{++}_R \rightarrow e^+ e^+) = BR(\delta^{++}_R \rightarrow \mu^+ \mu^+) = BR(\delta^{++}_R \rightarrow \tau^+ \tau^+) = \frac{1}{3}$. In the case of $M_{N_\mu} = M_{N_\tau} = 0.2 M_{N e}$ (and no mixing between the heavy neutrinos is assumed), the branching ratios are $BR(\delta^{++}_R \rightarrow e^+ e^+) = 0.926, BR(\delta^{++}_R \rightarrow \mu^+ \mu^+) = BR(\delta^{++}_R \rightarrow \tau^+ \tau^+) = 0.037$. In each case the decay to the $W_L^-$ bosons is highly suppressed.

(b) $u\bar{d} \rightarrow W^+_R \rightarrow \delta^{*+}_R \delta^{++}_R \rightarrow e^+ e^+ jj$

The case where an s-channel $W_R$ ”decays” via $W^-_R \rightarrow \delta^{*+}_R \delta^{++}_R$ involve the $\delta^{*+}_R \delta^{++}_R W_R^-$
coupling which also originates from the Higgs kinetic term \((D^\mu \Delta_R)^\dagger D_\mu \Delta_R\) and is given by

\[ -ig \left[ (\partial^\mu \delta^-_2) \delta^{++}_R - (\partial^\mu \delta^{++}_R) \delta^-_2 \right] W^+_R. \]  

(5.9)

This diagram, however, also gives a negligible contribution to the \(\delta^{++}_R\) production channel for several reasons: (i) the gauge eigenstate \(\delta^{++}_R\) consists of a suppressed fraction of the physical (massive) Higgs eigenstate \(H^\pm_2\):

\[ \delta^{\pm}_R = \frac{1}{\sqrt{1 + \left(\frac{\sqrt{2k_L v_R}}{k^-}\right)^2}} H^\pm_2. \]  

(5.10)

Since, as mentioned above, \(v_R \gg k_1, k_2\), the coupling to the physical charged Higgs is therefore suppressed compared to the \(W^- R W^- R \delta^{++}_R\) coupling (which is proportional to \(v_R\)), (ii) the mass of the physical particle \(H^\pm_2\) is

\[ M^2_{H^\pm_2} = \frac{1}{2} a_3 [v_R^2 \frac{k^+_2}{k^-_2} + \frac{1}{2} k^2_2]. \]  

(5.11)

For the parameters we use in this work (see Appendix A) we have \(M^2_{H^\pm_2} \sim 1.3 M_{W_R}\) which further reduces the kinematical viability of the processes involving \(H^\pm_2\), (iii) the Yukawa coupling of the physical \(H^\pm_2\) to the quarks is proportional to the quark masses, which for the light quarks (our case) is small compared to the coupling \(g\) of the charged gauge bosons to light quarks (again, suppressing this channel).

(c) \(u \bar{d} \rightarrow \phi^+_{1,2} \rightarrow \phi^-_2 \delta^{++}_R \rightarrow e^+ e^+ jj\)

In the case of an s-channel singly charged Higgs "decaying" via \(\phi^+_{1,2} \rightarrow \phi^-_2 \delta^{++}_R\) we need to consider the \(\phi^-_2 \delta^{++}_R\) coupling, which arises from the \(\beta_i\) terms in the scalar potential:

\[ \beta_1 Tr \left[ \phi \Delta_R \phi^\dagger L \right] + \beta_2 Tr \left[ \bar{\phi} \Delta_R \phi^\dagger L \right] + \beta_3 Tr \left[ \phi \Delta_R \phi^\dagger L \right], \]  

(5.12)

which gives rise to the following \(\phi^-_1 \phi^-_2 \delta^{++}_R\) coupling:

\[ \frac{1}{\sqrt{2}} v_L \delta^{++}_R \left( \beta_1 \phi^-_1 \phi^-_2 - \beta_2 \phi^-_1 \phi^-_2 - \beta_3 \phi^-_1 \phi^-_2 \right). \]  

(5.13)

The VEV \(v_L\) originates from the left-handed Higgs triplet and, as described in [17], vanishes due to the combined constraints which originate from the scalar potential minimization conditions and its explicit CP conservation. Moreover, the \(\beta_i\) parameters are set to zero in the MLRSM in order to reduce the mass scale of the non-SM gauge bosons and thus to theoretically allow possible observation at the LHC. This diagram is therefore also negligible.
(d) $u\bar{d} \rightarrow W_L^+ \rightarrow W_L^{-*}\delta_L^{++} \rightarrow e^+e^+jj$

Switching from the right handed doubly charged Higgs $\delta_R^{++}$ production to the production of its left-handed counterpart $\delta_L^{++}$ through an s-channel $W_L$ which then "decays" (i.e., off-shell) into a $\delta_L^{++}$ via $W_L^- \rightarrow W_L^-\delta_L^{++}$. For this process we need to consider the interaction vertex of two gauge bosons and the doubly charged Higgs $\delta_L^{++}$

$$\frac{1}{\sqrt{2}}g^2v LW_L^-W_L^{-}\delta_L^{++},$$

(5.14)

which is again proportional to $v_L$, and thus its contribution vanishes in the framework of the MLRSM. As opposed to the $\delta_R^{++}$ amplitudes above, one cannot replace one of the $W_L$ with a singly charged Higgs since, although the Lagrangian contains the vertex $W_L^-\delta_L^\mp\delta_L^{++}$, the singly charged Higgs field $\delta_L^\mp$ does not couple to quarks (as the quark Yukawa coupling originates from the Higgs bidoublet). Replacing the two $W_L$ with two singly charged Higgs fields is of course ruled out for the same reason.

We therefore conclude that the dominant diagram involving the decay of a doubly charged Higgs to leptons is diagram (a): $u\bar{d} \rightarrow W_R^+ \rightarrow W_R^{-*}\delta_R^{+++} \rightarrow e^+e^+jj$.

5.2. Majorana neutrino mediated diagrams

We consider below the (Majorana) neutrino mediated processes, which are characterized by a Drell-Yan production of $W_{L,R}$ boson which decays through a right handed heavy electron neutrino $N_e$.

(a) $u\bar{d} \rightarrow W_R^+ \rightarrow e^+N_e \rightarrow e^+e^+W_{L,R}^{-*} \rightarrow e^+e^+jj$

Let us first consider the case of an s-channel $W_R$ production, followed by a decay to a positron and a right handed (on-shell, massive and lighter) electron-neutrino $N_e$, which in turn decays through $W_R$ to a second positron and two jets. The coupling of the produced $W_R$ to the relevant fermions is not suppressed (as opposed to the alternatives below), and $W_R$ is dominantly on-shell in the mass range considered here. With nothing to suppress its occurrence $^{25}$, this diagram will turn out to give the dominant contribution among the neutrino mediated diagrams. We also point out that the produced $N_e$, although mostly right-handed, can decay substantially through both left handed and right handed W bosons, as explained below (see 5.2.1).

(b) $u\bar{d} \rightarrow W_L^+ \rightarrow e^+\nu_e \rightarrow e^+e^+W_L^{-*} \rightarrow e^+e^+jj$

Another possible Drell-Yan production which contributes to the neutrino mediated signal is the s-channel production of the SM $W_L$ which is part of the left-chiral

$^{25}$For a discussion about the Majoran neutrino-charged lepton coupling,
charged current. The $W_L$ can "decay" to either light or heavy neutrino, i.e. via $W_L^\pm \nu e^\mp$ or $W_L^\pm N e^\mp$ interactions, respectively. That said, we note that the coupling of $W_L^\pm N e^\mp$ is highly suppressed by a factor of $\sqrt{M_\nu(\text{light neut.})/M_N(\text{heavy neut.})}$ compared to $W_L^\pm \nu e^\mp$ and to $W_R^\pm N e^\mp$. In addition, even the diagram consisting of the coupling $W_L^\pm \nu e^\mp$ is suppressed, because an on-shell light neutrino cannot decay to a positron plus jets. The contribution of this diagram is therefore negligible.

\[(c)\ u\bar{d} \to \phi_{1,2}^+ \to e^+(N,\nu)e \to e^+e^+W_{L,R}^- \to e^+e^+jj\]

While the couplings involved in the Drell-Yan production of a gauge boson originate through the covariant derivative in the kinetic terms of the fermions, in the case of the singly charged Higgs Drell-Yan production the Yukawa coupling involved is, again, proportional to the masses of the light quarks, and thus negligible compared to the gauge coupling $g$. Moreover, since the singly charged Higgs fields (i.e. the fields $\phi_{1,2}^\pm$, see Appendix B) are composed of a very massive physical particle ($H_2^\pm$) of $O(10)$ TeV (in the mass range considered in this work), its Drell-Yan production has a negligible contribution to our signal.

We therefore conclude that the dominant contribution to our signal from the neutrino mediated diagrams comes from diagram (a): $u\bar{d} \to W_R^+ \to e^+N_e \to e^+e^+W_R^+ \to e^+e^+jj$.

5.2.1. The heavy neutrino decay

As mentioned above, the produced heavy neutrino $N_e$ (which is mostly right handed) can decay to a positron and two jets via $W_R$, $W_L$ or a singly charged Higgs, i.e. $N_e \to e^\pm W_{L,R}^\pm$ or $N_e \to e^\pm \phi_{1,2}^\pm$ followed by $(W_{L,R}^\pm$ or $\phi_{1,2}^\pm) \to jj$. The dominant decay is carried out via $N_e \to e^\pm W_R^\pm \to e^\pm jj$ for the same reason as before: the coupling of the right handed charged current $N_e e^\pm W_R^\pm$ is by far dominant compared to the couplings of the alternative decay channels (i.e. $N_e e^\pm W_L^\pm$ and $N_e e^\pm \phi_{1,2}^\pm$). Furthermore, the heavy neutrino decay via the singly charged Higgs is extremely small, as the mass of the singly charged Higgs is of $O(10)$ TeV (see also Eq.(5.11)). For instance, setting the parameter of the Higgs potential to $\alpha_3 = 4$ (in light of the above described bound on the neutral Higgs mass) leads to a lower bound of $M_{W_R} \geq 2.4$ TeV and, thus, the branching ratio of a heavy neutrino decay $N_e \to H_2^\pm e^\pm \to jj e^\pm$ upon setting e.g. $M_N = 2$ TeV is of $O(10^{-19})$. The relative weight of the neutrino decay through $W_L$ is determined by the interplay between the small coupling of the left handed charged current $N_e e^\pm W_L^\pm$ on one side and the enhancement due to the on-shell formation of $W_L$ (which is lighter than $N_e$) on the other side. Figure 8 presents the branching ratios of the right handed electron neutrino to an on-shell $W_L$ plus electron/positron and to a three body final
state through an off-shell $W_R$ (set as 4.5 TeV).\textsuperscript{26} In the figure, the three body decay of $N_e$ through $W_R$ is directly related to the mass of $N_e$. The increasing $N_e$ mass also reduces the $\sqrt{M_L(\text{light neut.})/M_N(\text{heavy neut.})}$ coupling of the charged current $N e^\pm W^\mp_L$, thereby decreasing the $W_L$ decay channel.

5.3. \textit{The case of equal neutrino masses} $M_{N_e} = M_{N_\mu} = M_{N_\tau}$

The cross section of the two like-sign leptons plus two jets signal, $e^\pm e^\pm jj$, for the case of degenerate heavy neutrinos is shown in figure 9. The Majorana neutrino mediated processes are clearly dominant in the lighter neutrino mass range ($M_{N_e} \lesssim 0.7 M_{W_R}$). However, as the neutrino mass increases, the branching ratio of $W_R \rightarrow N_e e$ decreases due to phase-space, causing a monotonic reduction in the cross section. The influence of the varying neutrino mass on the doubly charged Higgs mediated process is more complex. As shown before, the expression for the Yukawa coupling of the doubly charged Higgs to a pair of charged leptons depends on the heavy neutrino masses and the right handed lepton CKM-type mixing:

$$h_M = \frac{1}{\sqrt{2} v_R} K_R^T M'_{\text{diag}} K_R.$$ \hfill (5.4)

\textsuperscript{26}Other decay channels of $N_e$ are highly suppressed. For instance the decay $N_e \rightarrow Z \nu_e$ depends on the heavy-light mixing terms in the $K_R$ and $K_L$ matrices, which are, as mentioned above, extremely small.
Figure 9: Production cross sections for $pp \rightarrow e^\pm e^\pm jj + X$ process at the 14 TeV LHC. The masses of the three right-handed Majorana neutrinos are identical.
As a result, the Higgs-lepton coupling is directly related to the neutrino mass (see also Eq. (5.5)). This also leads, however, to a larger decay width of $\delta R^\pm$ as $N_e$ becomes heavier (reducing the Breit-Wigner distribution arising from the $\delta R^\pm$ propagator). Overall, this turns out to give a very moderate increase in the cross section of the $\delta R^\pm$ channel as $N_e$ becomes heavier, resulting eventually in dominating the signal production (in the highest $M_{N_e}/M_{W_R}$ range), as shown in figure 9. In the figure, the three panels represent three different values of the doubly charged Higgs mass, $M_{\delta R^\pm} = 350, 500, 700$ GeV. For a given $M_{W_R}$, the contribution of the $\delta R^\pm$ mediated diagram to the $e^\pm e^\pm jj$ signal becomes evident as $M_{N_e}$ approaches $M_{W_R}$. The rapid decline of the $N_e$ mediated channel is then compensated by the contribution of the $\delta R^\pm$ mediated one.

In general, for a heavier $W_R$, the $\delta R^\pm$ mediated channel contribution is more significant due to the dependency of the $W_R^\pm W_R^\mp \delta R^\pm$ coupling in $v_R$, as mentioned above. On the other hand, this contribution becomes numerically smaller as the masses of $W_R$ and $\delta R^\pm$ increase. For instance, for $M_{W_R} = 4$ TeV and $M_{N_e} = 0.7 M_{W_R}$ (and a cross section of $\sigma(N_e \text{ mediated}) = 0.364$ fb), the ratio $\frac{\sigma(\delta R^\pm \text{ mediated})}{\sigma(N_e \text{ mediated})}$ is 0.18, 0.12 and 0.07 for $M_{\delta R^\pm} = 350, 500, 700$ GeV, respectively. The corresponding ratio for $M_{W_R} = 5$ TeV and $M_{N_e} = 0.7 M_{W_R}$ (and $\sigma(N_e \text{ mediated}) = 0.045$ fb) is 0.3, 0.19 and 0.13, respectively. That is, while for a given $M_{\delta R^\pm}$ the relative contribution of the $\delta R^\pm$ channel increases with $M_{W_R}$, for any given $M_{W_R}$ the numerical cross section of the $\delta R^\pm$ channel decreases for a larger $M_{\delta R^\pm}$ and, therefore, harder to detect.

5.4. Non-equal neutrino masses: two benchmark cases

In general, the masses of the right handed neutrinos should not necessarily be identical. In order to investigate additional degrees of freedom, we examine two more benchmark cases with $M_{\delta R^\pm} = 500$ GeV: (i) $M_{N_\mu} = M_{N_e} = 0.2 M_{N_e}$ and (ii) $M_{N_\mu} = M_{N_e} = 2 M_{N_e}$. The results of case (i) are shown in figure 10. The dominance of the $N_e$ mediated channel is again evident in the lower neutrino mass range. However its cross section turns out to be slightly lower than in the equal masses case, since the two other right handed neutrinos (i.e. $N_\mu$ and $N_e$) are now lighter, which implies a larger $W_R$ decay width into these neutrinos and, therefore, a smaller BR to decay to an on-shell $N_e$. Moreover, the contribution of the $\delta R^\pm$ mediated channel turns out to be $\sim 2.5$ times stronger than in the equal masses case. This results from the change in the decay channels of $\delta R^\pm$, which in this case decays mainly to $e^\pm e^\pm$ since its coupling to the other charged leptons become substantially weaker (see discussion leading to Eqs. (5.4) and (5.5)). Using again the same setting as in the above example for the degenerate neutrinos case (i.e. $M_{W_R} = 5$ TeV, $M_{N_e} = 0.7 M_{W_R}$ and taking $M_{\delta R^\pm} = 500$ GeV), one obtains that $\frac{\sigma(\delta R^\pm \text{ mediated})}{\sigma(N \text{ mediated})}$ is increased from 0.19 in the equal masses case to 0.48 in the current case.

The results of case (ii) are shown in figure 11, where the mass range of the neutrino stems from the restriction of Eq. (3.1), which implies an upper bound in the setting of case (ii): $M_{N_e} \lesssim 0.5 M_{W_R}$. The main difference from the former cases is that the lower allowed mass range prevents the $N_e$ mediated channel from rapidly decreasing when
approaching the $W_R$ mass, thus diminishing the relative contribution of the $\delta_{R}^{±±}$ mediated channel. Moreover, the increment in the masses of the two additional right handed neutrinos enhances their coupling to $\delta_{R}^{±±}$, thereby increasing the $\delta_{R}^{±±}$ decay width and decreasing $BR(\delta_{R}^{±±} \rightarrow e^+e^-)$, therefore reducing the $\delta_{R}^{±±}$ mediated contribution.

5.5. Background analysis and realistic sensitivity estimates

The SM processes which have to be considered as a potential background for the two positron plus two jets signal are those which contain at least two positrons and two jets in the final state. The leading background processes turn out to be (see also [9])

$$pp \rightarrow W^+W^-b\bar{b} + X$$
$$pp \rightarrow ZW^± + X.$$  \hspace{1cm} (5.15)

These processes do not violate lepton number and therefore contain an additional electron/positron and/or neutrinos in the final state. For instance, the $WZ$ production shown in figure 12 contains two positrons and two jets, but also two neutrinos, and thus its final state contains missing energy, as opposed to the signal.

In order to correspond with former works [9, 10] we chose the case of degenerate neutrino masses $M_{N_\mu} = M_{N_\tau} = 0.2M_{N_e}$ (for $M_{\delta_{R}^{±±}} = 500$ GeV) and investigated the observability of the signal at the 14 TeV LHC. For the signal and background generation

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**Figure 10**: Production cross sections for $pp \rightarrow e^±e^-jj + X$ process at the 14 TeV LHC. The mass setting of the three right-handed Majorana neutrinos is $M_{N_\mu} = M_{N_\tau} = 0.2M_{N_e}$. 

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Figure 11: Production cross sections for $pp \rightarrow e^\pm e^\mp jj + X$ process at the 14 TeV LHC. The mass setting of the three right-handed Majorana neutrinos is $M_{N_\mu} = M_{N_\tau} = 2 M_{N_e}$. We used the implementation of our model into the CALCHEP [49] and MADGRAPH [50] softwares. We used Pythia [51] for the shower, fragmentation and hadronizations.

The K-factor for the signal was calculated using FEWZ 2.1 [52] to be around 1.3 in the $M_{W_R} = 2 - 7$ TeV range. The K-factor for the background processes was taken to be 1.4 [53].

We selected events with two isolated electrons and at least two jets in the final state, using some basic detector cuts on the pseudorapidity ($|\eta| < 3$) and on the transverse energy ($p_T > 5$ TeV). We used PGS [54] with the LHC card for the detector simulation. The efficiency of the signal event selection for $M_{N_e}/M_{W_R} = 0.1, 0.5, 0.9$ (and $M_{W_R} = 3$ TeV) is given in table 2 for both the Majorana neutrino mediated and the doubly charged Higgs mediated channels. The relatively low efficiency value in the case of $M_{N_e}/M_{W_R} \ll 1$ follows from the fact that for such low values of $M_{N_e}$ compared to $M_{W_R}$, $N_e$ is highly boosted, leading to difficulties in separating its decay products in the detector. This effect is absent from the Higgs channel, where the masses of the participating particles are the same for different values of $M_{N_e}$.

In order to reduce the background without considerably affecting the signal, the following cuts were applied:
**Figure 12:** An example of the SM background process $pp \rightarrow WZ \rightarrow e^+e^-jj + \text{missing energy}$ to the like-sign dileptons plus two jets signature; the process contains missing energy, unlike the signal.

$$M_{N_e}/M_{WR} \ (M_{WR} = 3 \text{TeV}) \quad 0.1 \quad 0.5 \quad 0.9$$

| Channel                  | 0.1  | 0.5  | 0.9  |
|--------------------------|------|------|------|
| $N_e$ channel            | 40.5%| 68.6%| 67.1%|
| $\delta_R^{\pm\pm}$ (500 GeV) channel | 72.3%| 68.7%| 67.3%|
| $\delta_R^{\pm\pm}$ (700 GeV) channel | 71.6%| 70.3%| 69.4%|

**Table 2:** The efficiency of the event selection criteria for the signal ($N_e$ and the $\delta_R^{\pm\pm}$ mediated channels) for different values of $M_{N_e}/M_{WR}$ (with $M_{WR} = 3 \text{TeV}$ and $M_{\delta_R^{\pm\pm}} = 500 \text{GeV, } 700 \text{GeV}$).

- Each of the two jets is required to have $E_T > 100 \text{GeV}$,
- The invariant mass of the $ee$ system is required to be larger than 200 GeV,
- The missing transverse energy does not exceed 100 GeV.

In figure 13 we present the like-sign dilepton signal and compare it to the SM background. In the three panels of the figure we plot the number of events as a function of the invariant masses of $e^\pm e^\pm jj$, $e^\pm jj$ and $e^\pm e^\pm$, corresponding to the reconstruction of the three masses of $W_R$, $N_e$ and $\delta_R^{\pm\pm}$, respectively (see also figure 6). It is evident from these figures that this process has practically zero background in the signal region of each of the three mass searches, and the signals of $W_R$, $N_e$ and of $\delta_R^{\pm\pm}$ are very distinct and within the LHC discovery limits.

We can now determine the sensitivity of the LHC to our signal by comparing the signal events which passed the above selection criteria and cuts to the corresponding
Figure 13: Mass reconstructions of $M_{W_R}$, $M_{N_e}$ and $M_{\delta_R^{\pm\pm}}$ from the selected events for $e^+e^-jj$ (upper panel), $jj e^\pm$ (middle panel) and $e^\pm e^\pm$ (lower panel) as a function of their invariant mass at the 14 TeV LHC, for $L = 300 \text{ fb}^{-1}$. The three heavy right-handed Majorana neutrinos are assumed to be degenerate.
background events. The two criteria which we define for discovery (in selected mass windows) are

- at least 10 signal events,
- the significance of the signal should be $S/\sqrt{B} \geq 5$.

The results for the benchmark case of $\delta_{R}^{\pm\pm} = 500$ GeV and degenerate heavy neutrino masses ($M_{N_e} = M_{N_u} = M_{N_\tau}$) are shown in figure 14. From the figure it is seen that after data-taking at low luminosity of 30 fb$^{-1}$, the discovery reach approaches 4.1 TeV for the $W_R$ and 2.9 TeV for the $N_e$ (with an insignificant excess, below 4%, in the reach of the $N_e$ mediated channel due to the contribution from the $\delta_{R}^{\pm\pm}$ mediated channel). After data-taking at high luminosity of 300 fb$^{-1}$, the discovery limits grow to 5.2 TeV and 3.6 TeV for $W_R$ and $N_e$, respectively, when considering only the $N_e$ mediated channel. Accounting also for the $\delta_{R}^{\pm\pm}$ mediated channel contribution, those limits slightly grow further, approaching 5.3 TeV and 4 TeV for $W_R$ and $N_e$, respectively.

**Figure 14:** Discovery potential for $pp \to W_R \to e^\pm e^\pm jj + X$ at the 14 TeV LHC, for integrated luminosities of 30 fb$^{-1}$ and 300 fb$^{-1}$ ($M_{\delta_{R}^{\pm\pm}} = 500$ GeV). The excluded area (see [55]) is a result of searching an excess of like-sign dilepton events with respect to the LRSM in the case of degenerate neutrino masses at the 8 TeV LHC.
Due to the the (above discussed) difficulty to separate between highly boosted neutrinos at lower $M_{N_e}/M_{W_R}$ ranges, the lower $M_{N_e}$ discovery limit is elevated by up to $\sim 25\%$ in the higher $W_R$ mass range, where the cross section is lower in general, and changes more moderately as a function of $M_{N_e}$ (the $\delta_R^{\pm\pm}$ mediated channel enhances the signal in that region by a tiny measure). In the higher range of $M_{N_e}/M_{W_R}$ the effect of the $\delta_R^{\pm\pm}$ compensates the decline in the $N_e$ channel and thereby enhances the signal by up to 10%.

In general, for a given $M_{\delta_R^{\pm\pm}}$, the effect of the $\delta_R^{\pm\pm}$ channel is stronger for a larger $BR(\delta_R^{\pm\pm} \rightarrow e^\pm e^\pm)$ due to the relative strength of the Higgs-di-electron (/di-positron) Yukawa coupling, as explained above. This can be significant if $M_{\delta_R^{\pm\pm}}$ is not too high. For instance, in the above mentioned case of $M_{N_{\mu}} = M_{N_e} = 0.2M_{N_e}$ and $M_{\delta_R^{\pm\pm}} = 500$ GeV (see also figure 10), for the mass value $M_{W_R} = 4.8$ TeV, the contribution of the $\delta_R^{\pm\pm}$ channel improves the LHC reach from $M_{N_e} = 0.7M_{W_R}$ by 20% to $M_{N_e} = 0.85M_{W_R}$. This increment, however, is nearly maximal for this specific $M_{\delta_R^{\pm\pm}}$, as in this setting $BR(\delta_R^{\pm\pm} \rightarrow e^\pm e^\pm) = 0.92$, i.e. close to one. The maximal contribution of the $\delta_R^{\pm\pm}$ channel to the signal is assessed in figure 15, by setting $BR(\delta_R^{\pm\pm} \rightarrow e^\pm e^\pm) = 1$ and $M_{W_R} = 3$ TeV.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure15.png}
\caption{The maximal contribution of the $\delta_R^{\pm\pm}$ channel to the $e^\pm e^\\mp jj$ signal, assuming $BR(\delta_R^{\pm\pm} \rightarrow e^\pm e^\pm) = 1$ and $M_{N_e}/M_{W_R} = 1$ (for $M_{W_R} = 3$ TeV).}
\end{figure}

setting $M_{W_R} = 3$ TeV. Since the cross section is determined also by the $\delta_R^{\pm\pm} e^\mp e^\mp$ coupling (which is proportional to the heavy neutrino mass, see Eqs. (5.4) and (5.5)), we also set $M_{N_e}/M_{W_R} = 1$. The cross section is plotted as a function of $M_{\delta_R^{\pm\pm}}$. It can be used, together with the criterion for an effective significance level$[10, 56]$

\begin{equation}
\text{Eff.Sig} = 2(\sqrt{S + B} - \sqrt{B}) \geq 5,
\end{equation}

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in order to estimate a higher mass limit for $\delta_{\pm}^{\pm}$, beyond which its contribution cannot be observed in the $\sigma(pp \rightarrow e^{\pm}e^{\pm}jj)$ signal. Assuming the absence of background events in the signal region (see the discussion above) one can infer from the cross section in figure 15 that, for a luminosity of 300 fb$^{-1}$ and for an assumed signal efficiency of 30%$^{27}$ (see also [55]), the contribution from the doubly charged Higgs process at the limit of $M_{\delta_{\pm}^{\pm}} \gtrsim 1.3$ TeV decreases below the limit of Eq.(5.16), reducing the effectiveness of this contribution.

6. Summary

The importance of the manifest/quasi manifest left-right symmetric model lies mainly in the fact that it restores parity symmetry at higher energy scales and provides a natural setup for the observed neutrino oscillations phenomena, based on the see-saw mechanism. Within the framework of this model we investigated the signal of like-sign dileptons and two jets, $pp \rightarrow e^{\pm}e^{\pm}jj + X$, at the LHC. We examined the dominant diagrams of this process and their behaviour as a function of the relevant physical parameters (i.e. mass, coupling, mixing, decay width). We calculated the cross sections of the signal for a number of benchmark cases, demonstrating the relative contribution of the two dominant diagrams, i.e. the two s-channel diagrams consisting of the Majorana neutrino ($N_e$) and the doubly charged Higgs ($\delta_{\pm}^{\pm}$) mediated diagrams. We showed that, while the $\delta_{\pm}^{\pm}$ contribution becomes more dominant as the mass of the right handed gauge boson $W_R$ grows, still the overall process becomes weaker.

We used the benchmark case of three degenerate heavy neutrino masses $M_{N_e} = M_{N_\nu} = M_{N_\tau}$ (and a doubly charged Higgs mass of $M_{\delta_{\pm}^{\pm}} = 500$ GeV) in order to demonstrate the practically zero background in the signal region of the $W_R$, $\delta_{\pm}^{\pm}$ and $N_e$ mass searches, and the expected positive prospects for the reconstruction and discovery of these particles at the LHC. In particular, it was shown that, for the integrated luminosity of 30 fb$^{-1}$ the $W_R$ boson and the right handed neutrino $N_e$ can be observed if $M_{W_R}$ and $M_{N_e}$ are lighter than 4.1 TeV and 2.9 TeV, respectively. The contribution of the $\delta_{\pm}^{\pm}$ in this case is below 4%. For the high luminosity case of 300 fb$^{-1}$, the discovery range can be further pushed to 5.2 TeV for $M_{W_R}$ and 3.6 TeV for $M_{N_e}$ upon considering only the $N_e$ mediated diagram, and to 5.3 TeV for $M_{W_R}$ and 4 TeV for $M_{N_e}$ when adding the $\delta_{\pm}^{\pm}$ contribution.

We pointed out that, with a luminosity of 300 fb$^{-1}$, the maximal sensitivity of the $e^{\pm}e^{\pm}jj$ signal to $\delta_{\pm}^{\pm}$ is obtained for $M_{\delta_{\pm}^{\pm}} \simeq 1.3$ TeV (this bound, taken for an assumed value of $M_{W_R} = 3$ TeV, can be lower for higher values of $M_{W_R}$). Beyond this limit, the contribution of the $\delta_{\pm}^{\pm}$ channel becomes too suppressed to be observed.

Complementarily to directly estimating the observability of an LRSM signal, we also examined the leading one-loop effects of the LRSM on the electroweak precision quantities $\Delta \nu_{LR}$ and $\delta \rho_{LR}$. We showed that, for the benchmark parameter space used

$^{27}$This is the overall efficiency consisting of applying the aforementioned selection efficiency and cuts.
in the signal analysis, these leading effects are indeed small enough to remain within
the limited precision of the electroweak data.

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Appendix A. Parameter settings used

The parameter settings of the MLRSM model file used in this work are as follows:

• Higgs VEVs (in GeV)

\[ k_1 = 245, \ k_2 = \sqrt{246.22 - k_1^2} = 24.48, \ v_R = 2543.2. \]  \hfill (A.1)

• Parameters in the Higgs potential

\[ \lambda_1 = 2, \ \lambda_2 = 0.05, \ \lambda_3 = -1.6, \ \lambda_4 = 0.05, \]
\[ \rho_1 = 0.9, \ 0.00056 \leq \rho_2 \leq 0.00128, \ \rho_3 = 1.81, \ \rho_4 = 1, \]
\[ \alpha_1 = 0.5, \ \alpha_2 = 0.5, \ \alpha_3 = 3.6. \]  \hfill (A.2)

• Couplings

\[ G_f = 1.166 \times 10^{-5} \text{ GeV}^{-2}, \ \alpha_s(M_Z) = 0.1184, \ \alpha(0) = \frac{1}{137.036}. \]  \hfill (A.3)

• Leptonic mixing matrices

\[ K_{L_i=j} = 1, \ K_{R_i=j+3} = 1, \ K_{L,R_{i\neq j,j+3}} = 0, \]
\[ K_{L1,4} = V_e, \ K_{L2,5} = V_\mu, \ K_{L3,6} = V_\tau, \]
\[ K_{R1,1} = -V_e, \ K_{R2,2} = -V_\mu, \ K_{R3,3} = -V_\tau, \]
\[ V_e = \sqrt{M_{N1}/M_{N4}}, \ V_\mu = \sqrt{M_{N2}/M_{N5}}, \ V_\tau = \sqrt{M_{N3}/M_{N6}}. \]  \hfill (A.4)

Appendix B. Higgs physical eigenstates

The Higgs multiplets consist of 20 degrees of freedom, i.e. 20 real fields. Obtaining the fields eigensystem is done by diagonalizing the squared-mass matrix:

\[ \left. \frac{\partial^2}{\partial \phi_i \partial \phi_j} V \right|_{\phi_i = \phi_j = 0} = m_{i,j}^2. \]  \hfill (B.1)

The eigenstates consist of

1. Four neutral scalar eigenstates \( H, H^0_1, H^0_2, H^0_3, \)

2. Four neutral pseudoscalar eigenstates \( A^0_1, A^0_2, G^0_1, G^0_2 \) (\( G^0_1 \) and \( G^0_2 \) are Goldstone bosons),

\[ \text{The leptonic mixing parameters } K_{L,R} \text{ are } 6 \times 3 \text{ CKM-like parameters in the lepton sector which connect the charged leptons to the six Majorana neutrinos. For more information see [8, 25].} \]
3. Four singly charged scalar eigenstates $H_1^\pm$, $H_2^\pm$, $G_L^\pm$, $G_R^\pm$ ($G_L^\pm$ and $G_R^\pm$ are Goldstone bosons).

4. Two doubly charged scalar eigenstates $H_{L,R}^{\pm\pm}$.

The corresponding eigenvalues/masses are given, e.g., in [25]. The non-physical Higgs fields may be written in terms of the above eigenstates as follows (the $\phi_1^0$, $\phi_2^0$ and $\delta_R^0$ states are given in the approximation $v_R \gg k_+$):

$$
\phi_1^0 \approx \frac{1}{k_+ \sqrt{2}} \left( k_1 k_+ + k_1 H^0 - k_2 H_1^0 - i k_1 G_1^0 + i k_2 A_1^0 \right),
$$

$$
\phi_2^0 \approx \frac{1}{k_+ \sqrt{2}} \left( k_2 k_+ + k_2 H + k_1 H_1^0 + i k_2 G_1^0 + i k_1 A_1^0 \right),
$$

$$
\delta_L^0 = \frac{1}{\sqrt{2}} \left( v_L + H_3^0 + i A_2^0 \right),
$$

$$
\delta_R^0 \approx \frac{1}{\sqrt{2}} \left( v_R + H_2^0 + i G_2^0 \right),
$$

$$
\phi_1^+ = \frac{k_1}{k_+ \sqrt{1 + \left( \frac{k^2}{k^2 + v_R^2} \right)}} H_1^+ - \frac{k_1}{k_+ \sqrt{1 + \left( \frac{\sqrt{2} k_+ v_R}{k^2 + v_R^2} \right)^2}} G_1^+ - \frac{k_2}{k_+} G_L^+,
$$

$$
\phi_2^+ = \frac{k_2}{k_+ \sqrt{1 + \left( \frac{k^2}{k^2 + v_R^2} \right)^2}} H_2^+ - \frac{k_2}{k_+ \sqrt{1 + \left( \frac{\sqrt{2} k_+ v_R}{k^2 + v_R^2} \right)^2}} G_2^+ + \frac{k_1}{k_+} G_L^+,
$$

$$
\delta_L^\pm = H_1^\pm,
$$

$$
\delta_R^\pm = H_2^\pm + \frac{1}{\sqrt{1 + \left( \frac{\sqrt{2} k_+ v_R}{k^2 + v_R^2} \right)^2}} G_2^+,
$$

$$
\delta_{L,R}^{\pm\pm} = H_{L,R}^{\pm\pm}.
$$

(B.2)