Multilepton collider signatures of heavy Dirac and Majorana neutrinos

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We discuss the possibility of observing multi-lepton signals at the Large Hadron Collider (LHC) from the production and decay of heavy standard model (SM) singlet neutrinos added in extensions of SM to explain the observed light neutrino masses by seesaw mechanism. In particular, we analyze two ‘‘smoking gun’’ signals depending on the Dirac or Majorana nature of the heavy neutrino: (i) for Majorana case, the same-sign di-lepton signal which can be used as a probe of lepton-number violation, and (ii) for Dirac case, the tri-lepton signal which conserves lepton number but may violate lepton flavor. Within a minimal Left-Right symmetric framework in which these additional neutrino states arise naturally, we find that in both cases, the signals can be identified with virtually no background beyond a TeV, and the heavy gauge boson W_R can be discovered in this process. This analysis also provides a direct way to probe the nature of seesaw physics involving the SM singlets at TeV-scale, and in particular, to distinguish type-I seesaw with purely Majorana heavy neutrinos from inverse seesaw with pseudo-Dirac counterparts.

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I. INTRODUCTION

One of the major evidences for the existence of new physics beyond the standard model (SM) is the discovery of nonzero neutrino mass from the observation of neutrino flavor oscillation phenomenon (for a recent update on the global neutrino data analysis, see Ref. [1]). In the SM, the left-handed (LH) neutrinos are massless mainly due to the absence of their right-handed (RH) counterparts (hence no Dirac mass) as well as the conservation of a global $B - L$ symmetry (hence no Majorana mass). Therefore, in order to generate nonzero neutrino masses, one must extend the SM sector by either adding three RH neutrinos (one per family) or by introducing $(B - L)$-breaking fields or both [2]. If we just add RH neutrinos (N) while keeping the $B - L$ symmetry unbroken, then the observed smallness of LH-neutrino masses require that the new Yukawa couplings (h_ν) involving the interaction of the N's with the LH-doublet (L) given by $h_\nu \bar{L} H N$ (where H is the SM Higgs doublet) must be extremely small, i.e. $h_\nu \lesssim 10^{-12}$ for sub-eV LH neutrino mass. In the absence of any obvious compelling arguments for such a tiny Yukawa coupling, the alternative path of generating nonzero neutrino masses by breaking $B - L$ symmetry seems more natural.

The simplest way to parameterize the $B - L$ breaking effects in SM extensions is through Weinberg's dimension-5 operator [3]

$$\mathcal{L}_{\rm eff} = \lambda_{ij} \frac{L_i L_j H H}{M} \qquad (i, j = e, \mu, \tau) \tag{1}$$

added to the SM Lagrangian, where M is the scale of new physics. After electroweak symmetry breaking (EWSB) due to the Higgs vacuum expectation value (vev) $\langle H \rangle \equiv v_{\rm wk}$, this operator leads to a nonzero neutrino mass of the form $m_\nu = \lambda v_{\rm wk}^2 / M$, and hence, $M/\lambda \gtrsim 10^{14}$ GeV for sub-eV neutrino mass. Thus, the new physics scale M depends on the effective Yukawa coupling λ (which is model-dependent), and can be in the TeV range to be directly accessible at colliders provided λ is very small.

There are both tree- and loop-level realizations of the dimension-5 operator given by Eq. (1) to generate nonzero neutrino masses [4]. The tree-level realization is the well-known seesaw mechanism in which the heavy particles associated with the new physics, after being integrated out, lead to the effective operator in Eq. (1). The simplest such model is the type I seesaw [5] in which the heavy particles are SM singlet Majorana fermions, usually known as the RH neutrinos (N), which couple to LH-doublets through Dirac Yukawa:

$$\mathcal{L}_Y = (h_\nu \bar{L} H N + \text{H.c.}) + N M_N N, \tag{2}$$

and M_N is the Majorana mass of N. After EWSB, this leads to the neutrino mass matrix of the form

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & M_D \\ M_D^T & M_N \end{pmatrix}, \tag{3}$$

where $M_D = v_{\rm wk} h_\nu$. The light mass eigenvalues are given by

$$m_\nu = - v_{\rm wk}^2 h_\nu M_N^{-1} h_\nu^T. \tag{4}$$

It is clear from Eq. (4) that for TeV-scale M_N, the Dirac Yukawa $h_\nu \lesssim 10^{-6}$, unless there are cancellations to get small neutrino masses from large Dirac masses using

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symmetries [6]. The heavy RH neutrinos, being SM singlets, can be produced at colliders only via $\nu - N$ mixing after virtual $W(Z)$'s produced in parton collision decay to $\ell(\bar{\nu}) + \nu$. Once produced, the N's decay equally likely to both charged leptons and anti-leptons (due to their Majorana nature), thus giving the distinct collider signature of like-sign di-lepton final states.[1] However, the mixing in type-I seesaw is typically given by $\theta_{\nu N} \sim \sqrt{m_\nu M_N^{-1}} \lesssim 10^{-6}$ (again barring cancellations), and hence, the production of the N's is highly suppressed. A detailed collider simulation shows that the minimal type I seesaw can be tested at colliders only if $\theta_{\nu N}$ is large ($\gtrsim 10^{-2}$) or M_N is small (up to a few hundred GeV) [7–9].

A second way to write the Weinberg operator in Eq. (1) is $(L^T \vec{\tau} L) \cdot (H^T \vec{\tau} H)/M$ where τ^i's are the usual Pauli matrices. This can be implemented by adding an $SU(2)_L$ bosonic triplet $\vec{\Delta} \equiv (\Delta^{++}, \Delta^{+}, \Delta^0)$ coupled to SM leptons through Majorana type couplings. This is known as the type II seesaw mechanism [10]. The Δ's, being SM non-singlets, can couple directly to the SM gauge bosons (W, Z and γ), and can be easily produced at colliders if their masses are in the sub-TeV to TeV range. The presence of doubly and singly charged scalars in the triplet lead to a very rich collider phenomenology of such models [11] which can be easily explored at the LHC.

Yet another way to write the effective Weinberg operator in Eq. (1) is $(L^T \vec{\tau} H)^2/M$ which can be implemented by adding an $SU(2)_L$ fermionic triplet ($\vec{\Sigma}$) coupled to leptons through Dirac Yukawas, just like the singlet ones in type I. This is known as the type III seesaw [12], and has very similar collider signatures as in type I case, except for the fact that the triplet fermions in this case couple directly to the SM W-boson, which makes them easier to search for at colliders up to about a TeV mass [9,13].

A completely different realization of the seesaw mechanism is the so-called inverse seesaw mechanism [14], where instead of one set of SM singlet Majorana fermions, one introduces two sets of them: N (Dirac) and S (Majorana). The resulting Lagrangian is given by

$$\mathcal{L}_Y = (h_\nu \bar{L} H N + N M_N N S + \text{h.c.}) + S \mu S \tag{5}$$

Because of the existence of the second set of singlet fermions (and perhaps additional gauge symmetries), the neutrino mass formula in these models has the form

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & M_D & 0 \\ M_D^T & 0 & M_N \\ 0 & M_N^T & \mu \end{pmatrix} \tag{6}$$

In the limit $\mu \ll v_{\rm wk} \lesssim M_N$, the lightest mass eigenvalues are given by

$$m_\nu \simeq v_{\rm wk}^2 h_\nu M_N^{-1} \mu (M_N^{-1})^T h_\nu^T \equiv F \mu F^T \tag{7}$$

where μ breaks the lepton number. Because of the presence of this new mass scale in the theory which is directly proportional to the light neutrino mass, the seesaw scale M_N can be naturally very low (within the range of colliders) even for "large" Dirac Yukawa couplings. This also allows for a large mixing $\theta_{\nu N} \simeq v_{\rm wk} h_\nu M_N^{-1}$, and makes the collider tests of this possibility much more feasible. However, due to the pseudo-Dirac nature of the RH neutrinos, the "smoking gun" signal for type I seesaw, namely, the lepton number violating same-sign di-lepton signal [8] is absent in this case. Instead, the lepton flavor violating tri-lepton signal [15,16] can be used to test these models. In this paper, we will mainly focus on these SM singlet RH neutrinos and present a detailed collider study of the multi-lepton final states in order to distinguish the heavy Dirac neutrinos from their Majorana counterparts at the LHC.[2]

Since the testability of seesaw is intimately related to the magnitude of the seesaw scale and the couplings of the new heavy particles with the SM particles, a key question of interest is whether there could be any theoretical guidelines for this new physics. A well-known example that explains the seesaw scale as a result of gauge symmetry breaking is the Left-Right (LR) Symmetric Theory based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [18]. Apart from restoring the parity symmetry at high energy, this theory provides a natural explanation of the seesaw scale as connected to the $SU(2)_R \times U(1)_{B-L}$-breaking scale. Also, the smallness of the neutrino mass is connected to the extent to which the RH-current is suppressed at low energy. Thus, the LR-symmetry provides a well-defined theory of neutrino masses [19] and can be used as a guide to study seesaw physics at the LHC [20]. Moreover, it provides a very attractive low-energy realization of a Grand Unified Theory (GUT) such as $SO(10)$, which is arguably the simplest GUT scenario for seesaw mechanism [2] as it automatically predicts the existence of RH neutrinos (along with the SM fermions). The $SO(10)$ embedding of TeV-scale LR models have been discussed in literature for both type I [21] and inverse seesaw [22]. Also, in case of inverse seesaw, as pointed out in Ref. [23], the LR gauge symmetry is essential to stabilize the form of the neutrino mass matrix given by Eq. (6). Therefore, in this paper, we work within the framework of the minimal LR-symmetric theory at TeV-scale.

This paper is organized as follows: in Sec. II we briefly summarize the main features of the minimal LR model, including the mixing between light and heavy neutrinos as well as gauge bosons, relevant for our analysis; in Sec. III, we discuss the production and decay of a heavy SM singlet neutrino at the LHC; in Sec. IV, we perform a detailed collider simulation of the multi-lepton events; in Sec. V, we summarize our results.

[1] This is a collider analogue of neutrinoless double beta decay to probe the lepton number violation.

[2] For a discussion on collider signals in other seesaw models, see Refs. [15,17].

II. THE MINIMAL LEFT-RIGHT MODEL

In this section, we review the minimal LR model based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [18] and discuss the mixing between the light and heavy neutrinos as well as gauge bosons. We also set the notations for the following sections.

In the LR model, the quarks and leptons are assigned to the following irreducible representations of the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$:

$$Q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}:(2,1,1/3), \qquad Q_R = \begin{pmatrix} u_R \\ d_R \end{pmatrix}:(1,2,1/3),$$
$$L_L = \begin{pmatrix} \nu_e \\ e_L \end{pmatrix}:(2,1,-1), \qquad L_R = \begin{pmatrix} N_e \\ e_R \end{pmatrix}:(1,2,-1)$$

and similarly for second and third generations. The minimal Higgs sector consists of a bi-doublet Φ:(2, 2, 0) and two triplets Δ_L:(3, 1, 2) and Δ_R:(1, 3, 2). After the spontaneous breaking of the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ symmetry to $U(1)_Q$ by the vev $v_{L,R}$ and κ, κ' of the Higgs fields $\Delta^0_{L,R}$ and Φ respectively, the phenomenological requirement $v_L \ll \kappa$, $\kappa' \ll v_R$ ensures the suppression of the RH-currents and the smallness of the neutrino mass. Also, the LR symmetry $\psi_L \leftrightarrow \psi_R$ for fermions and $\Delta_L \leftrightarrow \Delta_R$, $\Phi \leftrightarrow \Phi^\dagger$ for the Higgs fields leads to the relations $g_L = g_R = e/\sin\theta_W$ and $g' = e/\sqrt{\cos 2\theta_W}$ for the coupling strengths of the gauge bosons $W_{L,R}$ and Z' corresponding to the $SU(2)_{L,R}$ and $U(1)_{B-L}$ gauge symmetries, respectively, (where θ_W is the Weinberg angle and e is the electric charge of proton).

A. Mixing in the gauge sector

The charged gauge bosons $W^\pm_{L,R}$ in the weak eigenstate mix in the mass eigenstates W, W':

$$\begin{aligned} W &= \cos\zeta_W W_L + \sin\zeta_W W_R, \\ W' &= -\sin\zeta_W W_L + \cos\zeta_W W_R, \end{aligned} \tag{8}$$

where $\tan 2\zeta_W = 2\kappa\kappa'/(v_R^2 - v_L^2)$. The current bound on the mixing angle is as low as $\zeta_W < 0.013$ [24,25]; hence for our purposes, we can safely assume the mass eigenstates as the weak eigenstates, and recognize W_L as the pure SM W-boson. The lower bound on the W' mass comes from a variety of low-energy constraints, e.g. $K_L - K_S$ mass difference, $B_{d,s} - \bar{B}_{d,s}$ mixing, weak CP violation etc (For a recent update on the old results, see Ref. [26,27]). The most stringent limit on W_R mass in LR models is for the case of same CKM mixing angles in the left and right sectors: $M_{W_R} > 2.5$ TeV [26]; however, this limit can be significantly lowered if there is no correlation between the mixing angles in the two sectors [24,28]. The current collider bound on W' mass is around 1 TeV [29].

The neutral gauge bosons in LR model are mixtures of $W^3_{L,R}$ and B and the mixing between the weak eigenstates of these massive neutral bosons is parameterized as

$$Z = \cos\zeta_Z Z_1 + \sin\zeta_Z Z_2, \qquad Z' = -\sin\zeta_Z Z_1 + \cos\zeta_Z Z_2 \tag{9}$$

where Z, Z' are the mass eigenstates, and in the limit $v_L \ll \kappa$, $\kappa' \ll v_R$, the mixing angle is given by $\tan 2\zeta_Z \simeq 2\sqrt{\cos 2\theta_W}(M_Z/M_{Z'})^2$. Current experimental data constrain the mixing parameter to $<\mathcal{O}(10^{-4})$ and the Z' mass to values $>\mathcal{O}(\text{TeV})$ [25,30]. The current collider limit on the LR Z' mass is >998 GeV [29].

B. Mixing in the neutrino sector

In the neutrino sector of LR models, due to the presence of the RH neutrinos, the neutrino mass eigenstates (ν_i, N_i) are mixtures of the flavor eigenstates (ν_α, N_α) where $i = 1, 2, 3$ and $\alpha = e, \mu, \tau$ for three generations. For type I seesaw with only one additional set of SM singlets, this mixing can be parameterized as

$$\begin{pmatrix} \nu_\alpha \\ N_\beta \end{pmatrix} = \mathcal{V}_1 \begin{pmatrix} \nu_i \\ N_j \end{pmatrix} \tag{10}$$

where $\mathcal{V}_1$ is a 6×6 unitary matrix diagonalizing the full neutrino mass matrix in Eq. (3). Similarly, for inverse seesaw case in which we have two sets of SM singlet heavy neutrinos, the mixing is given by

$$\begin{pmatrix} \nu_\alpha \\ N_\beta \\ S_\gamma \end{pmatrix} = \mathcal{V}_2 \begin{pmatrix} \nu_i \\ N_j \\ N_k \end{pmatrix} \tag{11}$$

where $\mathcal{V}_2$ is a 9×9 unitary matrix diagonalizing the neutrino mass matrix in Eq. (6). Thus the weak interaction currents of light and heavy neutrinos are modified as follows:

$$\begin{aligned} \mathcal{L}_{\rm CC} &= \frac{g}{\sqrt{2}}[W^\mu_L \bar{\ell}_\alpha \gamma^\mu P_L \nu_\alpha + W^\mu_R \bar{\ell}_\beta \gamma^\mu P_R N_\beta] + \text{h.c.} \\ &= \frac{g}{\sqrt{2}}[W^\mu_L \bar{\ell}_\alpha \gamma^\mu P_L(\mathcal{V}_{\alpha i}\nu_i + \mathcal{V}_{\alpha j} N_j) \\ &\quad + W^\mu_R \bar{\ell}_\beta \gamma^\mu P_R(\mathcal{V}_{\beta i}\nu_i + \mathcal{V}_{\beta j} N_j)] + \text{h.c.}, \end{aligned} \tag{12}$$

$$\begin{aligned} \mathcal{L}_{\rm NC} &\simeq \frac{g}{2\cos\theta_W}[Z_\mu \bar{\nu}_\alpha \gamma^\mu P_L \nu_\alpha + \sqrt{\cos 2\theta_W} Z'_\mu \bar{N}_\beta \gamma^\mu P_R N_\beta] \\ &= \frac{g}{2\cos\theta_W}[Z_\mu\{\mathcal{V}^*_{\alpha i_1}\mathcal{V}_{\alpha i_2}\bar{\nu}_{i_1}\gamma^\mu P_L \nu_{i_2} + (\mathcal{V}^*_{\alpha i}\mathcal{V}_{\alpha j}\bar{\nu}_i \gamma^\mu P_L N_j + \text{H.c.}) + \mathcal{V}^*_{\alpha j_1}\mathcal{V}_{\alpha j_2}\bar{N}_{j_1}\gamma^\mu P_L N_{j_2}\} \\ &\quad + \sqrt{\cos 2\theta_W} Z'_\mu\{\mathcal{V}^*_{\beta j_1}\mathcal{V}_{\beta j_2}\bar{N}_{j_1}\gamma^\mu P_R N_{j_2} + (\mathcal{V}^*_{\beta i}\mathcal{V}_{\beta j}\bar{\nu}_i\gamma^\mu P_R N_j + \text{H.c.}) + \mathcal{V}^*_{\beta i_1}\mathcal{V}_{\beta i_2}\bar{\nu}_{i_1}\gamma^\mu P_R \nu_{i_2}\}], \end{aligned} \tag{13}$$

where we have dropped the subscript for $\mathcal{V}$ which now generically represents both $\mathcal{V}_1$ and $\mathcal{V}_2$ in Eqs. (10) and (11) respectively. Thus, in general, $\mathcal{V}$ is a $(3+n)\times(3+n)$ unitary matrix, where n stands for the number of SM singlets (3 for type-I and 6 for inverse seesaw). This can be decomposed into the following blocks:

$$\mathcal{V} = \begin{pmatrix} U_{3\times 3} & V_{3\times n} \\ X_{n\times 3} & Y_{n\times n} \end{pmatrix} \tag{14}$$

where U is the usual PMNS mixing matrix for the light neutrinos. The unitarity of $\mathcal{V}$ implies that

$$\begin{aligned} UU^\dagger + VV^\dagger &= U^\dagger U + X^\dagger X = I_{3\times 3}, \\ XX^\dagger + YY^\dagger &= V^\dagger V + Y^\dagger Y = I_{n\times n}. \end{aligned} \tag{15}$$

with $UU^\dagger$, $Y^\dagger Y \sim \mathcal{O}(1)$ and $VV^\dagger$, $X^\dagger X \sim \mathcal{O}(m_\nu/M_N)$. Thus in Eqs. (12) and (13), the mixing between the light states, $\mathcal{V}_{\alpha i} \equiv U_{\alpha i}$, and between the heavy states, $\mathcal{V}_{\beta j} \equiv Y_{\beta j}$ both are of order $\mathcal{O}(1)$, whereas the mixing between the light and heavy states, $\mathcal{V}_{\alpha j} \equiv V_{\alpha j}$, $\mathcal{V}_{\beta i} \equiv X_{\beta i} \sim \mathcal{O}(M_D M_N^{-1})$ for both type I and inverse seesaw cases, which, in principle, could be large for TeV mass RH neutrinos and large Dirac Yukawa case. Henceforth, we will generically denote this mixing between light and heavy neutrinos by $V_{\ell N}$, and assume the other mixing elements in Eqs. (12) and (13) to be $\mathcal{O}(1)$.

The electroweak precision data constrain the mixing $V_{\ell N}$ involving a single charged lepton [31] and the current 90% C.L. limits are summarized below:

$$\sum_{i=1}^{3} |V_{eN_i}|^2 \leq 3.0\times 10^{-3}, \qquad \sum_{i=1}^{3} |V_{\mu N_i}|^2 \leq 3.2\times 10^{-3}, \qquad \sum_{i=1}^{3} |V_{\tau N_i}|^2 \leq 6.2\times 10^{-3} \tag{16}$$

These limits are crucial for our analysis since they determine the decay rate of the heavy neutrinos to multi-lepton final states, as discussed in next section. One can also get constraints on the mixing involving two charged leptons from lepton-flavor violating (LFV) processes [32][3]:

$$\left|\sum_{i=1}^{3} V_{eN_i}V^*_{\mu N_i}\right| \leq 1.0\times 10^{-4},$$
$$\left|\sum_{i=1}^{3} V_{eN_i}V^*_{\tau N_i}\right| \leq 1.0\times 10^{-2},$$
$$\left|\sum_{i=1}^{3} V_{\mu N_i}V^*_{\tau N_i}\right| \leq 1.0\times 10^{-2}$$

For the heavy neutrino mass below 100 GeV, the updated limits are summarized in Ref. [33].

[3]However, these constraints can be easily evaded if, for example, each heavy neutrino mixes with a different charged lepton.

Another constraint for the manifest LR model comes from neutrinoless double beta decay as there is a new contribution involving the heavy gauge boson W_R and RH Majorana neutrino [34]. For a TeV mass RH neutrino, this puts a lower bound on $M_{W_R} \geq 1.1$ TeV which increases as $M_N^{-1/4}$ for smaller RH neutrino mass. In this paper, we therefore mainly focus on a TeV mass RH neutrino.

III. PRODUCTION AND DECAY OF HEAVY NEUTRINOS

At a proton-proton collider, a single heavy neutrino can be produced at the parton-level, if kinematically allowed, in

$$q\bar{q}' \to W_L^*/W_R \to \ell^+ N(\ell^- \bar{N}), \tag{17}$$

which has lepton-number conserving (LNC) or violating (LNV) decay modes depending on whether N is Dirac or Majorana.[4] Since τ-lepton identification may be rather complicated in hadron colliders [35], we restrict our analysis to only the light charged-leptons ($\ell = e, \mu$). The parton-level production cross sections, generated using CalcHEP [36] and with the CTEQ6L parton distribution function [37], are shown in Fig. 1 as a function of the mass of N for 1.5, 2 and 2.5 TeV W_R mass (solid lines) at $\sqrt{s} = 14$ TeV LHC. We also show the normalized production cross section $\sigma/|V_{\ell N}|^2$ (normalized to $|V_{\ell N}|^2 = 1$) for SM W_L-boson mediation (dashed line), which is generated only through the mixing $V_{\ell N}$ between the LH and RH neutrinos. We can clearly see that the W_L-mediated production is highly suppressed by the mixing; even for large mixing, the cross section for a heavy RH neutrino with $M_{W_R} > M_N \gg M_{W_L}$ is mostly dominated by the W_R-channel because W_R can always decay on-shell whereas the W has to be highly off-shell to produce N.

The heavy RH neutrino decays to SM leptons plus a gauge or Higgs boson through its mixing with the left sector: $N \to \ell W$, νZ, νH. So all these decay rates are suppressed by the mixing parameter $|V_{\ell N}|^2$. In LR models, N can also have a three-body decay mode: $N \to \ell W_R^* \to \ell jj$ (and similarly for Z') which is not suppressed by mixing, but by mass of W_R. Note that the decay mode $N \to \ell W_R^* \to \ell\ell\nu$ will be suppressed by mixing as well as W_R-mass and hence the di-jet mode is always the dominant final state for the three-body decay of N. The various partial decay widths of N are shown in Fig. 2 for a mixing parameter $|V_{\ell N}|^2 = 0.001$ and Higgs mass of 125 GeV. It is clear that for mixing larger than $\mathcal{O}(10^{-4})$, N mainly decays into the SM gauge or Higgs boson which could subsequently lead to multi-lepton final states.

[4]In Eq. (17) and following, $\bar{N}$ should be replaced by N for a Majorana RH neutrino.

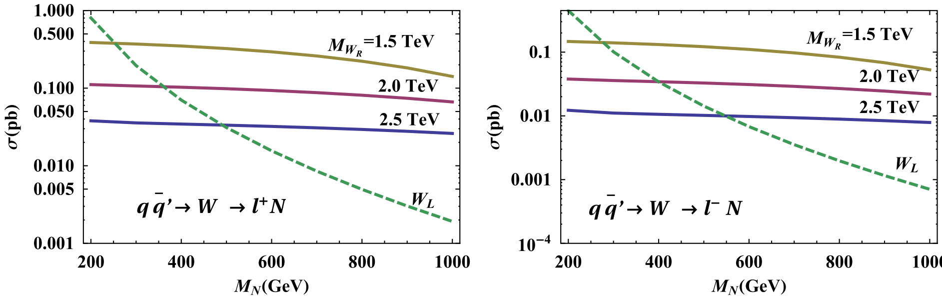

FIG. 1 (color online). The cross section for $q\bar{q}' \to W_L^*/W_R \to N\ell^\pm$ for various values of W_R mass (solid lines). Also shown is the normalized cross section $\sigma/|V_{\ell N}|^2$ for W_L-mediated s-channel (dashed line).

It should be emphasized here that the LR symmetry provides a unique channel for the production of RH neutrino through the W_R gauge boson, without any mixing suppression, and multi-lepton final states through the decay of N to SM gauge bosons, which even though suppressed by the mixing, still offer a promising channel to study the Dirac or Majorana nature of N. Without the LR symmetry (and hence W_R), the production of N (through SM W/Z) will also be suppressed by mixing, which limits its observability to only a few hundred GeV masses, mainly due to the large SM background [8]. On the other hand, LR-symmetric models provide much higher mass reach at the LHC in the multi-lepton channel, as we discuss in the next section.

We further note that a single N can also be produced in $q\bar{q} \to Z^*/Z' \to \bar{\nu}N$ but the resulting final state has either one charged lepton or opposite-sign di-leptons, and is buried under the huge LHC background.[5] One could also produce the RH neutrinos in pairs through a Z'-exchange: $q\bar{q} \to Z' \to N\bar{N}$, if kinematically allowed; however, the decay of two N's will be suppressed by $|V_{\ell N}|^4$, and hence, negligible.

Thus we conclude from this study that for a hadron collider analysis, the most suitable production channel for a Dirac RH neutrino in LR models is through W_R-exchange and the N decay mode through SM W. We note that this particular channel was not considered in the previous studies of RH neutrino signals in LR models [27,40], because they only considered a heavy Majorana neutrino (in type I seesaw) for which the golden channel is the same-sign di-lepton mode in Fig. 3(a): $q\bar{q}' \to W_R^\pm \to N\ell^\pm \to W_R^*\ell^\pm l^\pm \to jj\ell^\pm\ell^\pm$ [41]. In this case, the 3-body decay mode of $N \to \ell W_R^* \to \ell jj$ is dominant over the 2-body decay $N \to \ell W$ because the latter is suppressed by mixing which is usually very small in type I seesaw. However, for a heavy Dirac neutrino, this same-sign di-lepton mode is absent and the corresponding opposite-sign di-lepton mode $q\bar{q}' \to W_R^\pm \to N\ell^\pm \to W_R^*\ell^\mp l^\pm \to jj\ell^\mp\ell^\pm$ has large SM background. So the golden channel for a heavy Dirac neutrino is the tri-lepton mode in Fig. 3(b) where the W/W_R^* decays to leptonic final states: $pp \to W_R^\pm \to N\ell^\pm \to W/W_R^*\ell^\mp\ell^\pm \to \nu\ell^\pm\ell^\mp\ell^\pm$ [15,16]. As discussed earlier in this section, the N decay to SM W is dominant over the 3-body decay through W_R for mixing $|V_{\ell N}| \lesssim 10^{-4}$, which is easily satisfied in inverse seesaw models, for instance. This is also true for type I seesaw with large mixing [6,9], in which case the 2-body decay of N to SM gauge bosons (W, Z, H) will be dominant over the three-body decay through a virtual W_R.

IV. MULTI-LEPTON SIGNALS AND SM BACKGROUND

We perform a full LHC analysis of the multi-lepton final states given in Fig. 3 and the SM background associated with it. The signal and background events are calculated at parton-level using CalcHEP [36] which are then fed into PYTHIA [42] to add initial and final state radiation and pile up, and perform hadronization of each event. Finally, a fast detector simulation is performed using PGS [43] to simulate a generic LHC detector. We use the more stringent L2 trigger [44] in order to reduce the SM background. We note

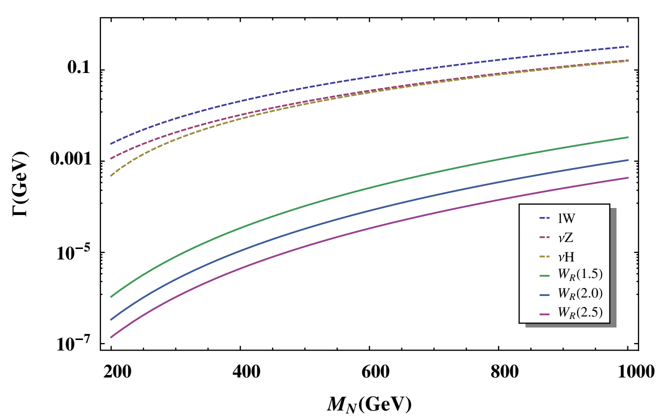

FIG. 2 (color online). The partial decay widths of the RH neutrino into ℓW, νZ, νH (dashed lines) as a function of its mass for a mixing parameter $|V_{\ell N}|^2 = 0.001$. Also shown are the three-body decay widths for $N \to \ell W_R \to \ell jj$ for $M_{W_R} = 1.5$, 2.0 and 2.5 TeV.

[5] This could, however, be important in cleaner environments, e.g. e^+e^- [38] and $e\gamma$ [39] colliders.

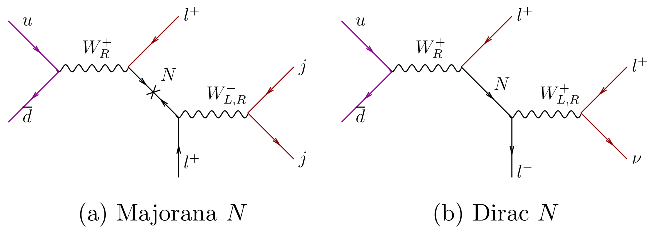

FIG. 3 (color online). The golden channels for heavy Majorana and Dirac neutrino signals at the LHC.

that the signal strength remains the same, if we use the low threshold L1 trigger, which is very close to the actual values used by the CMS detector. The L2 trigger has high enough thresholds to reduce all the SM background below the signal and therefore we do not need to impose any additional cuts on the events.

The major SM background for the di-lepton signal comes from the semi-leptonic decay of a $t\bar{t}$ pair,

$$q\bar{q}, gg \to t\bar{t} \to W^+ bW^- \bar{b} \to jjb\ell^- \bar{\nu} \bar{b}, \qquad (18)$$

and the b-quark giving the second lepton: $b \to c\ell\nu$. Similarly, tri-lepton background is produced in the fully leptonic decay of $t\bar{t}$ and the third lepton coming from b-quark. Though the charged leptons from b-quark decay typically have small transverse momentum, the large $t\bar{t}$ production cross section (compared to the production of N) is responsible for the dominant background, and must be taken into account in the detector simulation. The other dominant SM backgrounds for multi-lepton channels at the LHC arise from the production of WZ, WW, ZZ, WWW, $Wt\bar{t}$, $Zb\bar{b}$, $Wb\bar{b}$ etc.. A detailed discussion of the background analysis for multi-lepton final states can be found in Ref. [15,45]. We find that by implementing the L2 trigger, most of this SM background can be eliminated, and the remaining background is dominantly due to $t\bar{t}$, WW, WZ and ZZ (which we collectively denote as "SM background" in the following).

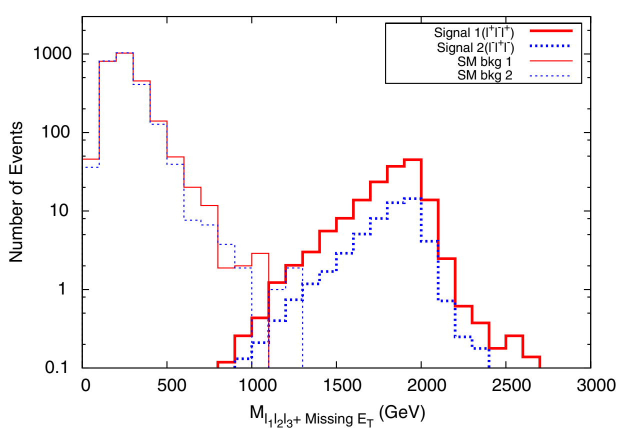

FIG. 4 (color online). Selected events for the tri-lepton final state as a function of the invariant mass of $\ell^\pm\ell^\mp\ell^\pm + \not{E}_T$ (100 GeV bins) for $\sqrt{s} = 14$ TeV and $\mathcal{L} = 8$ fb^{-1}. We have chosen $M_{W_R} = 2$, $M_N = 1$ TeV and $|V_{\ell N}|^2 = 0.0025$ for this plot. The dominant SM background ($t\bar{t} + WW + WZ + ZZ$) is also shown here.

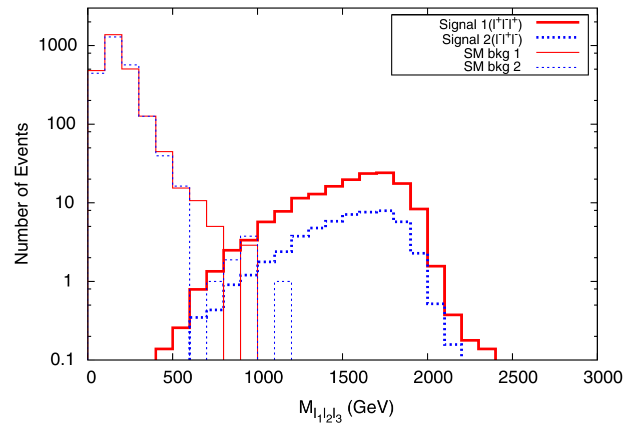

FIG. 5 (color online). Selected events for tri-lepton final state as a function of the invariant mass of $\ell^\pm\ell^\mp\ell^\pm$ for the same parameters as in the Fig. 4 caption.

The invariant mass of the final state particles is used to reconstruct the mass of W_R. The selected events for the tri-lepton $(\ell^\pm\ell^\mp\ell^\pm) + \not{E}_T$ final state is shown in Fig. 4 (thick lines) as a function of the invariant mass (100 GeV bins) for $\sqrt{s} = 14$ TeV LHC and integrated luminosity, $\mathcal{L} = 8$ fb^{-1}. The expected SM background events ($t\bar{t} + VV$) are also shown (thin lines). Here we have chosen $M_{W_R} = 2$ TeV and $M_N = 1$ TeV. We have also taken the mixing parameter $V_{\ell N}$ just below the experimental upper bound: $|V_{\ell N}|^2 = 0.0025$ (For a lower value of mixing, the cross section and hence the total number of events, will decrease as $|V_{\ell N}|^2$). We find that the invariant mass of W_R is reconstructed nicely and the tri-lepton channel is virtually background free above 1 TeV or so. We also plot the invariant mass of $(\ell^\pm\ell^\mp\ell^\pm)$ in Fig. 5 which has the sharp end point at W_R mass. We note here that the tri-lepton final

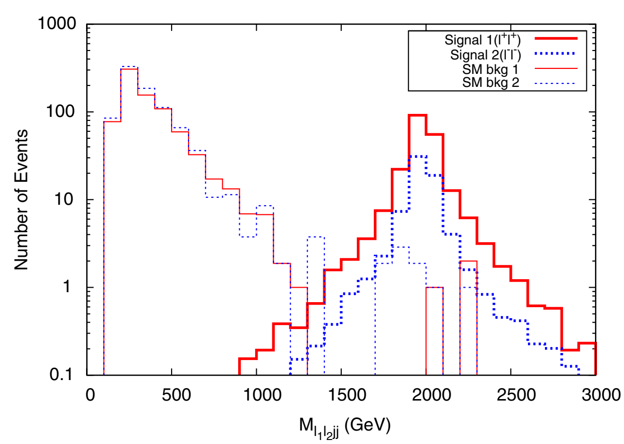

FIG. 6 (color online). Selected events for the same-sign di-lepton final state as a function of the invariant mass of $\ell^\pm\ell^\pm jj$ for the same parameters as in the Fig. 4 caption.

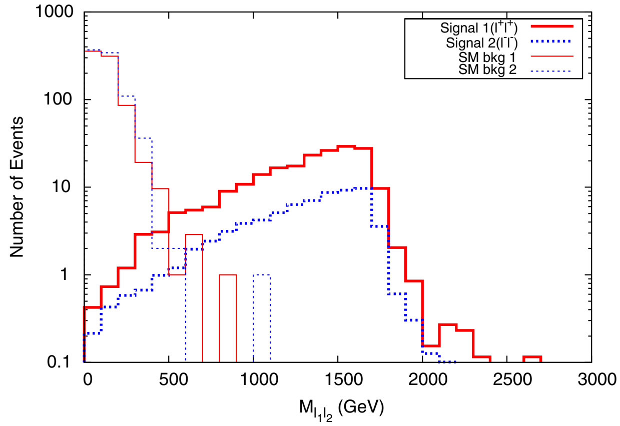

FIG. 7 (color online). Selected events for same-sign di-lepton final state as a function of the invariant mass of $\ell^\pm\ell^\pm$ for the same parameters as in the Fig. 4 caption.

states with two positively charged (anti)leptons has more likelihood to be produced than those with one positively charged, which is naively expected for a proton-proton collision.

For comparison, we have also performed similar analysis for a heavy Majorana neutrino, similar to those in Ref. [27,40], but with a large mixing $|V_{\ell N}|^2 = 0.0025$. Hence, as we discussed in Sec. III, N mostly decays to SM gauge bosons and charged leptons, and not through the 3-body decay involving W_R. The resulting events are shown in Figs. 6 and 7 for the invariant mass of $\ell^\pm\ell^\pm jj$ and $\ell^\pm\ell^\pm$ respectively. The parameters chosen are the same as for Figs. 4 and 5. We note that the number of same-sign di-lepton events passing the L2 trigger are roughly 1 order of magnitude larger than the tri-lepton events. This is because of the overall enhancement of the cross section for the di-lepton final state because the branching fraction for hadronic decay modes of $W \to jj$ is roughly thrice that of the light leptonic decay modes $W \to \ell\nu$.

V. CONCLUSION

We have discussed the collider signatures of a heavy SM singlet neutrino in a minimal LR framework, which can be of either Majorana or Dirac nature depending on the mechanism for neutrino mass generation. In particular, we have analyzed the multi-lepton signals to distinguish a TeV-scale Dirac neutrino from a Majorana one at the LHC. We perform a detailed collider simulation to show that, in LR models, a TeV-scale heavy neutrino can be produced at the LHC dominantly through a W_R exchange, which subsequently decays dominantly via SM gauge boson exchange. The invariant mass of the final state particles can be used to nicely reconstruct the mass of W_R in multi-lepton channels which are virtually background free above a TeV. We observe that if the heavy neutrino is of Majorana-type, there will be distinct lepton-number violating signals, including the same-sign di-lepton signal discussed here. However, in the absence of the same-sign di-lepton signal, the tri-lepton signal can be used to establish the Dirac nature of the heavy neutrino. This provides a direct way of probing the seesaw mechanism and the associated new physics at TeV-scale, and can be used to distinguish type-I seesaw (with purely Majorana heavy neutrinos) from inverse seesaw (with pseudo-Dirac ones) at the LHC.

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