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Right-handed neutrinos and the 2 TeV W' boson

Pilar Coloma[⋆], **Bogdan A. Dobrescu**[⋆], **Jacobo Lopez-Pavon**[⋄]

⋆ *Theoretical Physics Department, Fermilab, Batavia, IL 60510, USA*

⋄ *SISSA and INFN sezione di Trieste, 34136 Trieste, Italy*
& INFN, Sezione di Genova, via Dodecaneso 33, 16146 Genova, Italy

Abstract

The CMS e^+e^-jj events of invariant mass near 2 TeV are consistent with a W' boson decaying into an electron and a right-handed neutrino whose TeV-scale mass is of the Dirac type. We show that the Dirac partner of the right-handed electron-neutrino can be the right-handed tau-neutrino. A prediction of this model is that the sum of the τ^+e^+jj and τ^-e^-jj signal cross sections equals twice that for e^+e^-jj. The Standard Model neutrinos acquire Majorana masses and mixings compatible with neutrino oscillation data.

1 Introduction

Searches for new gauge bosons that carry electric charge ±1, called W' bosons, have been intensely performed at hadron colliders [1]. More than a year ago, the CMS Collaboration [2] has reported an excess of events in the search for a W' boson decaying into a heavy right-handed neutrino and an electron. The observed final state consists of an electron-positron pair and two hadronic jets, with an invariant mass in the 1.8–2.2 TeV range. Although this excess has a statistical significance of only 2.8σ, it is particularly interesting given the theory connection [3] based on the $SU(2)_R$ gauge group between this and other excess events observed at the LHC in the $W' \to jj$ [4], $W' \to WZ$ [5, 6] and $W' \to Wh^0$ [7] channels.

A key feature of the CMS e^+e^-jj excess is that it is not accompanied by an excess in the e^-e^-jj and e^+e^+jj final states. In the traditional $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model, where the right-handed fermions form $SU(2)_R$ doublets [8], the right-handed neutrinos have Majorana masses. As a consequence, the cross section for the final state with same-sign leptons ($e^{\mp}e^{\mp}jj$) is equal to that for opposite-sign leptons (e^+e^-jj) [9, 10]. Given this generic prediction and the smaller backgrounds for same-sign leptons, the ATLAS searches [11] have not included e^+e^-jj resonances.

An explanation for the lack of a same-sign signal has been proposed in [3], where the right-handed neutrino N_R partners with the neutral component of an $SU(2)_R$ doublet fermion and acquires a Dirac mass. This solution works with the minimal Higgs sector [12], namely a bidoublet and an $SU(2)_R$ triplet. A related solution, proposed in [13], is to generate Dirac masses for right-handed neutrinos by introducing three gauge-singlet fermions while extending the Higgs sector with an $SU(2)_R$-doublet scalar. Both these solutions also include subdominant Majorana masses for right-handed neutrinos, so that the physical states are actually pseudo-Dirac. Given that only one of the 14 CMS $eejj$ events has same-sign leptons, which may be due to the background (estimated to be approximately four $eejj$ events), it is safe to neglect the Majorana component.

On the other hand it has been pointed out in [14] that particular values for the Majorana masses of the right-handed neutrinos and their CP-violating phases may allow a suppression of the same-sign $eejj$ signal, without introducing additional fields. This may be counter-intuitive, given that the $W' \to eN$ decays produce very narrow on-shell N particles, whose decay widths should not be sensitive to the CP violating phases. The suppression is provided by interference effects between processes proceeding through different right-handed neutrinos. Other studies of the W' interpretation of the CMS $eejj$ excess can be found in [15].

Here we propose that the Dirac partner of the right-handed electron-neutrino is the right-handed τ-neutrino. We show that the flavor structure required by this mechanism can easily be enforced by a symmetry. The ensuing model is remarkably simple and leads to peculiar signals at the LHC. This mechanism also explains the suppression of same-sign $eejj$ signals observed in the simulations discussed in [14]: our Dirac fermion can be decomposed in two degenerate Majorana fermions whose interactions with the W' boson include imaginary couplings.

In Section 2 we show how a Dirac state arises from two right-handed neutrino flavors, and derive its interactions. The implications for LHC phenomenology are analyzed in Section 3. Mechanisms for generating masses for the Standard Model (SM) neutrinos in this context are discussed in Section 4. We summarize our conclusions in Section 5.

2 Two right-handed neutrinos make a Dirac fermion

We consider an $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge theory with a minimal fermion content: the three generations of SM quarks and leptons plus one right-handed neutrino per generation. The three right-handed neutrinos (N_R^e, N_R^μ, N_R^τ) together with the three right-handed charged leptons (e_R, μ_R, τ_R) form $SU(2)_R$ doublets of $U(1)_{B-L}$ charge -1. We label these by $L_R^e = (N_R^e, e_R)^\top$, $L_R^\mu = (N_R^\mu, \mu_R)^\top$ and $L_R^\tau = (N_R^\tau, \tau_R)^\top$.

The minimal Higgs sector consistent with the W' signals discussed in the Introduction has been analyzed in detail in Ref. [12]. It includes an $SU(2)_R$ triplet scalar T of $U(1)_{B-L}$ charge $+2$, and an $SU(2)_L \times SU(2)_R$ bidoublet scalar Σ that is not charged under $U(1)_{B-L}$. The VEV of T breaks $SU(2)_R \times U(1)_{B-L}$ down to the hypercharge gauge group, $U(1)_Y$, while inducing the required mass splitting between W' and Z' [3].

The VEV of Σ induces a mass mixing between the $SU(2)_L \times SU(2)_R$ gauge bosons that carry electric charge (W_L and W_R), so that the physical spin-1 particles (W and W') couple together to the Z boson. This VEV breaks the electroweak symmetry, and can be parametrized as $\langle \Sigma \rangle = v_H \, \mathrm{diag}\,(\cos\beta, e^{i\alpha_\Sigma} \sin\beta)$, where $v_H \simeq 174$ GeV. The excess events in the $W' \to WZ$ searches [5, 6] indicate a $W'WZ$ coupling consistent with near-maximal $W_L - W_R$ mixing, which corresponds to $\tan\beta \approx 1$.

The right-handed lepton doublets have gauge-invariant Yukawa couplings to the triplet scalar. Our main assumption is that these Yukawa couplings have the following flavor structure:

$$- \frac{y_{\mu\mu}}{2} (\overline{L}_R^\mu)^c \, i\sigma_2 \, T \, L_R^\mu - y_{e\tau} (\overline{L}_R^e)^c \, i\sigma_2 \, T \, L_R^\tau + \text{H.c.} \quad , \tag{2.1}$$

where $y_{\mu\mu}$ and $y_{e\tau}$ are positive dimensionless parameters, and σ_2 is the Pauli matrix acting on $SU(2)_R$ representations. The c label here indicates, as usual, the charge-conjugate spinor [16]. We will show in Section 3 that this non-diagonal structure in flavor space implies the absence of same-sign $eejj$ events at the LHC.

The flavor structure of the Yukawa couplings in Eq. (2.1) can be enforced by a symmetry. An example is a a global $U(1)$ symmetry with the L_R^e, L_R^μ, L_R^τ doublets carrying charges $-1, 0, +1$, respectively; the scalar T must then be neutral under this global symmetry. Even a discrete subgroup (Z_n with $n \geq 3$) of this $U(1)$ symmetry would be sufficient.

Once the scalar triplet T gets a VEV,

$$\langle T\rangle = \begin{pmatrix} 0 & 0 \\ u_T & 0 \end{pmatrix} \ , \tag{2.2}$$

the right-handed neutrinos acquire masses through the following Lagrangian terms:

$$-\, u_T \left(\overline{N}^e_R,\, \overline{N}^\mu_R,\, \overline{N}^\tau_R\right)^c \begin{pmatrix} 0 & 0 & y_{e\tau} \\ 0 & y_{\mu\mu} & 0 \\ y_{e\tau} & 0 & 0 \end{pmatrix} \begin{pmatrix} N^e_R \\ N^\mu_R \\ N^\tau_R \end{pmatrix} \ . \tag{2.3}$$

The parameter u_T is in the 3–4 TeV range [12] in order to accommodate the mass and coupling of the W' boson indicated by the LHC data.

The 2-component fermion N^μ_R is already in the mass eigenstate basis, and has a Majorana mass $m_{N^\mu} = y_{\mu\mu} u_T$. More importantly, we find that N^e_R and N^τ_R form a 4-component fermion N_1 of Dirac mass

$$m_{N_1} = y_{e\tau} u_T \ . \tag{2.4}$$

The interactions of the W' boson with the right-handed neutrinos in the gauge eigenstate basis are given by

$$\frac{g_{\rm R}}{\sqrt{2}} W'_\nu \left(\overline{N}^e_R \gamma^\nu e_R + \overline{N}^\mu_R \gamma^\nu \mu_R + \overline{N}^\tau_R \gamma^\nu \tau_R\right) + \text{H.c.} \ , \tag{2.5}$$

where the $g_{\rm R}$ coupling is equal to the $SU(2)_R$ gauge coupling up to negligible corrections. We identify the N^τ_R and N^e_R fields with the left- and right-handed components of the Dirac fermion N_1:

$$N^\tau_R \equiv N^c_{1_L} \quad , \quad N^e_R \equiv N_{1_R} \ . \tag{2.6}$$

The spinors satisfy the usual [16] relation $N^c_{1_L} \equiv (N_{1_L})^c = (1/2)(1+\gamma_5) i\gamma^2 N^*_1$. The interactions of N_1 with W' take the following form:

$$\frac{g_{\rm R}}{\sqrt{2}} W'_\nu \left(\overline{N}_{1_R} \gamma^\nu e_R + \overline{N}^c_{1_L} \gamma^\nu \tau_R\right) + \text{H.c.} \tag{2.7}$$

The interaction of N^μ_R remains as in Eq. (2.5).

We will assume in what follows that the gauge and mass eigenstates of the charged leptons are identical. In Section 4 we will show that this can be a consequence of the same global $U(1)$ symmetry responsible for the flavor structure in Eq. (2.5).

3 Predictions for the LHC

The couplings of the 4-component N_1 fermion (formed of the 2-component fermions N_R^e and N_R^τ) to the W' boson displayed in Eq. (2.7) are peculiar: the $\overline{N}_1$ antifermion interacts with the electron, while the N_1 fermion interacts with the τ. This has profound implications for the LHC phenomenology.

For the mass ordering of interest here, $m_{N_1} < M_{W'} \lesssim m_{N^\mu}$, the W'^+ boson may undergo leptonic decays only into $e^+ N_1$ or $\tau^+ \overline{N}_1$, with equal branching fractions. Similarly, the only leptonic decay channels of the W'^- boson are $e^- \overline{N}_1$ and $\tau^- N_1$.

The N_1 fermion predominantly decays via an off-shell W' into a quark-antiquark pair and an e^- or a τ^+, again with equal branching fractions. Note that N_1 decays into a quark-antiquark pair and an e^+ or a τ^- are forbidden by the $U(1)$ flavor symmetry introduced in Section 2, while the $\overline{N}_1$ antifermion decays into these final states are allowed. Consequently, the following cascade decays have equal branching fractions:

$$\begin{aligned} B(W'^+ \!\to e^+ N_1 \to e^+ e^- jj) &= B(W'^+ \!\to \tau^+ \overline{N}_1 \to \tau^+ \tau^- jj) \\ &= B(W'^+ \!\to e^+ N_1 \to e^+ \tau^+ jj) \\ &= B(W'^+ \!\to \tau^+ \overline{N}_1 \to \tau^+ e^+ jj) \ . \end{aligned} \tag{3.1}$$

CPT invariance implies that these branching fractions are also equal to those for the cascade decays of W'^- into e^+e^-jj, $\tau^+\tau^-jj$, $e^-\tau^-jj$, and τ^-e^-jj.

The W' widths into same-sign $eejj$ or $\tau\tau jj$, as well as into opposite-sign $e^\pm\tau^\mp jj$ vanish. These are immediate consequences of the W' couplings shown in Eq. (2.7) in the Dirac basis.[1] Both the same-sign and opposite-sign $\mu\mu jj$ signals are negligible as long as m_{N^μ} is not smaller than $M_{W'}$ by more than 5% or so.

The observed excess events in the e^+e^-jj final state [2] can be explained in this theory, together with the other hints for a W' near 2 TeV discussed in [3] for $g_{\rm R} \approx 0.5$. Note first that compared to the case where the Dirac partner of N_R^e is a new fermion, the e^+e^-jj rate predicted in the theory discussed here has one less parameter because the N_1 coupling to an electron and a W' is fixed in Eq. (2.7). The ensuing branching fraction

[1] The same result can be obtained using the Majorana basis. After diagonalizing the mass matrix shown in Eq. (2.3), the 2-component mass eigenstates are given by $i(N_R^e - N_R^\tau)/\sqrt{2}$ and $(N_R^e + N_R^\tau)/\sqrt{2}$. The destructive interference between decay amplitudes proceeding through these two Majorana states gives vanishing widths for $W'^\mp \to e^\mp e^\mp jj$. Feynman rules for Majorana fermions useful for this computation are given in [17, 10].

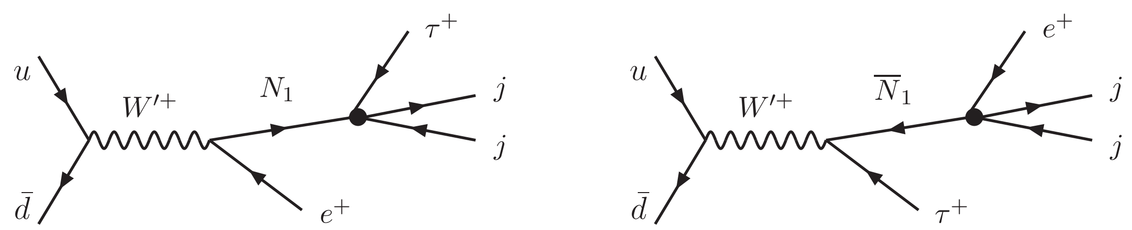

Figure 1: Processes for resonant $e^+\tau^+jj$ production at the LHC. The Dirac fermion N_1, formed of the N_R^e right-handed neutrino and the charge-conjugate of N_R^τ, has interactions with τ that violate fermion number (note the arrows) according to Eq. (2.7). The $\bullet$ represents the 4-fermion interaction mediated by an off-shell W'^{-}.

of N_1 into an electron and a quark-antiquark pair is larger by a factor of 2 than in the baseline model analyzed in [3], while the branching fraction of $W' \to eN_1$ is larger by a factor of 4. As these changes increase the e^+e^-jj rate (by a factor of 8), they can easily be compensated by the phase-space suppression of the $W' \to eN_1$ width. We find that the predicted e^+e^-jj rate is consistent with the CMS excess when the N_1 mass satisfies 1.4 TeV $\lesssim m_{N_1} \lesssim$ 1.7 TeV for $M_{W'} \approx 1.9$ TeV.

As a result of the equalities between branching fractions, there are a few striking predictions for the LHC. First the cross section for $pp \to W' \to e^+e^-jj$ is equal to that for $pp \to W' \to \tau^+\tau^-jj$. A comparison of the same-sign $e\tau jj$ cross sections with the opposite sign $eejj$ cross section is less straightforward because the production cross section in pp collisions for W'^{+} is larger than that for W'^{-}. There is, however, a simple relation: the $pp \to W' \to e^+e^-jj$ cross section is exactly one half of the sum of the $pp \to W'^{+} \to e^+\tau^+jj$ and $pp \to W'^{-} \to e^-\tau^-jj$ cross sections. Thus,

$$\begin{aligned} \sigma(pp \to W' \to e^+e^-jj) &= \sigma(pp \to W' \to \tau^+\tau^-jj) \\ &= \frac{1}{2}\Big[\sigma(pp \to W' \to e^+\tau^+jj) + \sigma(pp \to W' \to e^-\tau^-jj)\Big] \quad . \qquad (3.2) \end{aligned}$$

These predictions can be tested in various ways in Run 2 of the LHC. The production cross section for a W' of mass near 2 TeV grows by a factor of 5 at $\sqrt{s} = 13$ TeV compared to $\sqrt{s} = 8$ TeV. Thus, the approximately ten $eejj$ signal events observed by CMS with 20 fb^{-1} in Run 1 imply about 500 $e^\pm\tau^\pm jj$ signal events (see Figure 1) with 100 fb^{-1} in Run 2, if the efficiency of the event selection is not modified. Even though the backgrounds increase in Run 2, that number of events would allow the observation of same-sign $e\tau jj$ signals independently with leptonic or hadronic τ decays.

So far we focused on N_1 decays into a lepton and two light quarks, which have the dominant branching fractions. Other N_1 decay modes potentially observable in Run 2 include a lepton plus $t\bar{b}$ or WZ, which occur via an off-shell W', as well as a lepton plus a W boson. All these decays lead to the same cross-section relations as those involving light quarks. For example, $\sigma(pp \to e^+\tau^+\bar{t}\,b) + \sigma(pp \to e^-\tau^-\bar{t}\,b) = 2\sigma(pp \to e^+e^-\bar{t}\,b)$.

The $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge group implies that besides the W' boson there is a Z' boson. The breaking of $SU(2)_R \times U(1)_{B-L}$ by an $SU(2)_R$-triplet VEV, as assumed here, implies that the Z' is substantially heavier than the W' boson. For $M_{W'} \approx 1.9$ TeV and $g_{\rm R}$ in the 0.45–0.6 range, the Z' mass satisfies 3.4 TeV $< M_{Z'} <$ 4.5 TeV [3], a range that will be probed in Run 2 of the LHC. Our right-handed neutrino sector leads to testable predictions. The $Z' \to \bar{N}_1 N_1$ decay followed by N_1 decays to an electron or τ and two jets leads to final states with two leptons and four jets. The flavor structure of the model implies

$$B(Z' \to e^-\tau^- + 4j) = B(Z' \to e^+\tau^+ + 4j) = B(Z' \to e^+e^- + 4j) = B(Z' \to \tau^+\tau^- + 4j) \ . \tag{3.3}$$

An additional prediction of our model is that the same-sign same-flavor channels $Z' \to \tau^\pm\tau^\pm + 4j$ and $Z' \to e^\pm e^\pm + 4j$ are forbidden.

4 Masses for Standard Model neutrinos

Both the charged lepton masses and Dirac neutrino masses may be generated, in principle, by the following $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$-invariant Yukawa couplings to the bidoublet Σ and its charge-conjugate state $\widetilde{\Sigma}$:

$$-\,\overline{L}_L^\alpha \left(y_{\alpha\beta}\,\Sigma + \tilde{y}_{\alpha\beta}\,\widetilde{\Sigma} \right) L_R^\beta + \text{H.c.} \tag{4.1}$$

where α, β are flavor indices, $L_L^\alpha = (\nu_L^\alpha, \ell_L^\alpha)^\top$ are the $SU(2)_L$ lepton doublets, and y and $\tilde{y}$ are dimensionless coefficients. However, depending on the charges of Σ and L_L^α under the $U(1)$ global symmetry discussed in Section 2, the terms shown in Eq. (4.1) may be forbidden.

The Dirac masses that link the SM neutrinos to the right-handed neutrinos cannot be larger than the MeV scale if the right-handed neutrino masses are at the TeV scale (otherwise the SM neutrino masses, obtained from the seesaw mechanism, are too large). Since the nonzero components of the bidoublet VEV are of the same order (*i.e.*, $\tan\beta \sim 1$

as discussed in Section 2), a fine-tuned cancellation between the y and $\tilde{y}$ coefficients would be needed to keep these Dirac masses much smaller than the τ mass. We therefore choose global $U(1)$ charges that forbid the interactions in Eq. (4.1); for example, L_L^e, L_L^μ and L_L^τ have same charges as the corresponding right-handed leptons $(-1,\ 0,\ +1)$, while Σ has charge $+3$.

Let us present a different mechanism for generating the SM lepton masses. We introduce a gauge-singlet scalar ϕ that carries $U(1)$ charge $+3$, and acquires a VEV $\langle\phi\rangle$. Consider the following gauge-invariant operators:

$$-\frac{\tilde{c}_\alpha}{m_f^3}\,\phi\,\overline{L}_L^\alpha\,\widetilde{\Sigma}\,\widetilde{T}^\dagger\widetilde{T}\,L_R^\alpha\ +\text{H.c.}\quad , \tag{4.2}$$

where $\tilde{c}_\alpha$ $(\alpha = e, \mu, \tau)$ are dimensionless couplings and m_f is the mass of some heavy fields, which have been integrated out. These operators generate the charged lepton masses (similarly to the down-type quark masses [12]) once the scalars acquire VEVs:

$$(m_e, m_\mu, m_\tau) = v_H \cos\!\beta\,\frac{\langle\phi\rangle u_T^2}{m_f^3}\,(\tilde{c}_e, \tilde{c}_\mu, \tilde{c}_\tau)\quad . \tag{4.3}$$

This ensures that the gauge and mass eigenstates of the charged leptons coincide, as mentioned in Section 2.

The SM neutrinos get Majorana masses from the following gauge-invariant dimension-6 operators:

$$\frac{\eta_{\alpha\beta}}{M^2}\,(L_L^\alpha)^c\,\Sigma\,T\,\Sigma^\dagger\,L_L^\beta\quad . \tag{4.4}$$

The global $U(1)$ allows the $\eta_{\mu\mu}$, $\eta_{e\tau}$ and $\eta_{\tau e}$ coefficients to be nonzero. Dirac masses between the SM and right-handed neutrinos may also be generated by

$$-\frac{C_{\alpha\beta}}{m_f^3}\,\phi\,\overline{L}_L^\alpha\,\widetilde{\Sigma}\,T^\dagger T\,L_R^\beta + \text{H.c.}\quad , \tag{4.5}$$

with the result

$$m_D = v_H \sin\!\beta\,\frac{\langle\phi\rangle u_T^2}{m_f^3}\,C\quad . \tag{4.6}$$

Here we ignored the complex phase from the Σ VEV, and we collected the $C_{\alpha\beta}$ coefficients in a $3{\times}3$ matrix C. The mass matrices for the charged and neutral leptons are independent of one another and no cancellation is necessary. The full 6×6 mass matrix in the neutrino sector has the following block structure:

$$\mathcal{M}_\nu = \begin{pmatrix} m_L & m_D \\ m_D^\top & M_R \end{pmatrix}\quad , \tag{4.7}$$

where the Majorana mass matrix for the right-handed neutrinos, M_R, is shown in Eq. (2.3), and the Majorana mass matrix for the SM neutrinos, m_L, arises from operators (4.4).

All Lagrangian terms discussed above are invariant under the global $U(1)$. However, the spontaneous breaking of this symmetry implies the existence of a Nambu-Goldstone boson θ_ϕ that couples to leptons. In order to satisfy phenomenological constraints, the global $U(1)$ must also be explicitly broken so that θ_ϕ becomes heavy (alternatively, the global symmetry is a Z_n group, avoiding the presence of Nambu-Goldstone bosons). Note that $\langle\phi\rangle$ may be much larger than the $SU(2)_R$ breaking scale, so that even small explicit $U(1)$ breaking terms may push the mass of θ_ϕ above the reach of the LHC.

Lagrangian terms that explicitly break the global $U(1)$ may also have important contributions to the elements of the m_D and m_L matrices. The smallness of these mass terms implies negligible effects in M_R. Thus, there is enough freedom to accommodate the masses and mixings of the active neutrinos. After integrating out the right-handed neutrinos, whose masses are at the TeV scale, the SM neutrinos acquire Majorana masses.

Let us comment on the particular case where m_L is negligibly small. Given the particular structure of M_R, the minimal possibility for m_D that reproduces the light neutrino squared-mass differences and mixing angles observed at neutrino oscillation experiments [18] is the following:

$$m_D = \begin{pmatrix} m_{11} & 0 & m_{13} \\ m_{21} & 0 & m_{23} \\ m_{31} & 0 & m_{33} \end{pmatrix} \equiv (m_{D1} \ , \ 0 \ , \ m_{D3}) \quad . \tag{4.8}$$

The Majorana mass matrix for the light active neutrinos is obtained by a TeV-scale seesaw:

$$m_\nu = m_{D1} \frac{1}{m_{N1}} m_{D3}^\top . \tag{4.9}$$

Notice that in this minimal model only the Dirac fermion, formed by N_R^e and $(N_R^\tau)^c$, participates in the light neutrino mass generation, because the Majorana state N_R^μ decouples.

The neutrino mass matrix given in Eq. (4.9) has a zero mode and thus one of the light neutrino masses vanishes, while the other two generate the solar and atmospheric neutrino mass differences. This scenario resembles the minimal linear seesaw [19] studied in detail in [20], where it was shown that the flavor structure of the neutrino Dirac mass terms, m_{D1} and m_{D3}, is completely fixed by neutrino oscillation data up to a global factor.

In order to generate a third light neutrino mass (and have more freedom in the parameter space), the second column of m_D should be switched on. If this is the case, N_R^μ would also participate in the generation of light neutrino masses.

Our TeV-scale right-handed neutrino sector is not currently constrained by low-energy observables such as neutrinoless double beta decay or lepton-flavor violating decays. Generically, the rate for neutrinoless double beta decay mediated by the W' and the right-handed neutrinos is tightly correlated to the $pp \to e^{\mp}e^{\mp}jj$ cross section. Our mechanism, which forbids the same-sign $eejj$ signal at the LHC, also forbids neutrinoless double beta decay mediated by two W' bosons.

There is, however, a new physics contribution to the neutrinoless double beta decay mediated by two W bosons. This arises from one insertion of the mixing between a SM neutrino and a right-handed neutrino [see m_{13} in Eq. (4.8)], and one insertion of the mixing between the $SU(2)_L \times SU(2)_R$ gauge bosons. The latter has an upper limit [12] of $\sin\theta_+ \leq (g_{\rm R}/g)(M_W/M'_W)^2 \approx 10^{-3}$, where g is the SM weak coupling. The contribution of this type [21] to the effective neutrino mass relevant for neutrinoless double beta decay is given by

$$\begin{aligned} m_{\beta\beta} &\simeq \frac{g_{\rm R}}{g}\sin\theta_+ \frac{m_{13}}{m_{N_1}}\langle p\rangle \\ &\approx 0.1\,\text{eV}\left(\frac{\sin\theta_+}{10^{-3}}\right)\left(\frac{m_{13}}{1\,\text{MeV}}\right)\left(\frac{1\,\text{TeV}}{m_{N_1}}\right)\left(\frac{\langle p\rangle}{100\,\text{MeV}}\right) \;, \end{aligned} \tag{4.10}$$

where $\langle p\rangle$ is the momentum transfer in the nuclear transition. For $m_{13} \sim 0.3$ MeV and $m_{N_1} \approx 1.5$ TeV, this would be within the reach of future neutrinoless double beta decay experiments. In addition, there is the usual light neutrino contribution with two W propagators. This is suppressed if the light neutrino spectrum has normal mass ordering, so a neutrinoless double beta decay signal would provide information about N_1. For an inverted ordering (or normal ordering with a quasi-degenerate spectrum), the contributions from N_1 and from the light neutrinos are comparable. Thus, the total rate could be slightly enhanced or even suppressed with respect to the case of only light neutrinos.

Due to the large $e - \tau$ flavor mixing in the right-handed neutrino sector in our model, it is relevant to ask whether there are constraints from charged lepton-flavor violating processes. The $\tau \to e\gamma$ contribution from loops involving the W' and right-handed neutrinos is suppressed because the flavor structure of our model imposes that lepton-flavor violation and lepton-number violation can only occur simultaneously. The main contribution to charged lepton-flavor violating decays in our model is in the $\tau^{\mp} \to \mu^{\mp}\mu^{\mp}e^{\pm}$ channel. This decay proceeds via a box diagram with two W' bosons, where one of the fermion lines involves only muon-flavor states, while the $\tau - e$ flavor transition takes place along the second fermion line via an N_1 exchange. The branching fraction for this process is of the order of $g_{\rm R}^8 M_W^4/(4\pi g M_{W'})^4 \approx 10^{-12}$, where we have taken into account that the $m_{N^\mu}/M_{W'}$ and $m_{N_1}/M_{W'}$ ratios are of order one. This prediction is below the current experimental limit $B(\tau^- \to e^+\mu^-\mu^-) < 1.7 \times 10^{-8}$ [1].

As the muon-lepton number is violated in Eq. (2.1) by two units, the $U(1)$-invariant contributions to $\mu \to eee$, $\mu \to e$ conversion in nuclei, or $\mu \to e\gamma$ vanish. The explicit breaking of the $U(1)$ flavor symmetry leads to contributions through the mixing of light and heavy neutrinos. Given that in our model flavor-violating processes occur at one loop, the most sensitive probe is $\mu \to e\gamma$. The predicted branching fraction, though, is below the current limit [22] because it is suppressed by $(\alpha/\pi)\sin^2\theta_+ \lesssim 10^{-9}$ as well as $(m_{21}/m_\mu)^2 \lesssim 10^{-5}$.

5 Conclusions

In a SM extension with three right-handed neutrinos of masses at the TeV scale, we have proposed a flavor structure that pairs the right-handed electron-neutrino and the right-handed tau-neutrino, forming a Dirac fermion (N_1) whose mass violates lepton number. The right-handed muon-neutrino (N_R^μ) remains a purely Majorana state. This flavor structure is enforced by a global $U(1)$ symmetry, and leads to peculiar interactions [see Eq. (2.7)] of the W' boson associated with the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge group.

Our flavor symmetry predicts specific relations between the decay widths of the TeV-scale neutrinos. N_1 decays into an e^- or a τ^+ and a W' (which is off-shell for $m_{N_1} < M_{W'}$) or a W (through $W_L - W_R$ mixing). The branching fractions for decays involving e^- are equal to those involving τ^+. The N_R^μ Majorana fermion decays, also with 50% branching fractions, into a μ^- or a μ^+ and a W' (which is on-shell assuming $m_{N^\mu} > M_{W'}$).

Furthermore, the decays of the W' into a SM lepton and a heavy neutrino have tightly correlated branching fractions, as shown in Eq. (3.1). As a consequence, the decay of the W' boson into an electron and a right-handed neutrino produces an e^+e^-jj signal, while the rate for same-sign $e^\mp e^\mp jj$ events automatically vanishes. This provides a compelling explanation for the excess of opposite-sign e^+e^-jj events with an invariant mass near 2 TeV reported by the CMS Collaboration [2].

The flavor symmetry leads to additional predictions at the LHC. First, the opposite-flavor same-sign processes $pp \to W'^+ \to \tau^+e^+jj$ and $pp \to W'^- \to \tau^-e^-jj$ are allowed; the sum or their rates is exactly twice the rate for the same-flavor opposite-sign process

$pp \to W' \to e^+e^-jj$. Second, the rates for the two same-flavor opposite-sign processes, $W' \to e^+e^-jj$ and $W' \to \tau^+\tau^-jj$, are equal. These and other predictions discussed in Section 3 can be tested in Run 2 of the LHC.

The masses for active neutrinos get contributions from both a TeV-scale seesaw mechanism and lepton-number violating dimension-6 operators involving only left-handed fields. Ignoring the latter, there are Dirac mass terms at the MeV scale that link the left- and right-handed neutrinos; after integrating out the right-handed neutrinos at the TeV scale, the left-handed neutrinos acquire Majorana masses at the sub-eV scale. The minimal choice for the Dirac mass matrix that reproduces the neutrino oscillation data leaves one of the light neutrinos massless. In this scenario, the light neutrinos acquire Majorana masses only from their Yukawa interactions with N_1. A more general choice for the Dirac mass matrix would allow the third left-handed neutrino to acquire mass. In that case, the right-handed muon-neutrino would also be involved in the generation of active neutrino masses.

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