Majorana neutrinos and lepton-number-violating signals in top-quark and W-boson rare decays

Shaouly Bar-Shalom\textsuperscript{a}\textsuperscript{*}, Nilendra G. Deshpande\textsuperscript{b}, Gad Eilam\textsuperscript{c}, Jing Jiang\textsuperscript{b} and Amarjit Soni\textsuperscript{d}

\textsuperscript{a}Physics Department, Technion-Institute of Technology, Haifa 32000, Israel
\textsuperscript{b}Institute for Theoretical Science, University of Oregon, OR 97403, USA
\textsuperscript{c}Theory Group, Brookhaven National Laboratory, Upton, NY 11973, USA

(Dated: March 26, 2022)

We discuss rare lepton-number-violating top-quark and W-boson four-body decays to final states containing a same-charge lepton pair, of the same or of different flavors: $t \to bW^{-} \ell_{i}^{+} \ell_{j}^{-}$ and $W^{+} \to J J' \ell_{i}^{+} \ell_{j}^{-}$, where $i \neq j$ or $i = j$ and $J J'$ stands for two light jets originating from a $u \bar{d}$ or a $c \bar{s}$ pair. These $\Delta L = 2$ decays are forbidden in the Standard Model and may be mediated by exchanges of Majorana neutrinos. We adopt a model independent approach for the Majorana neutrinos mixing pattern and calculate the branching ratios (BR) for these decays. We find, for example, that for $O(1)$ mixings between heavy and light Majorana neutrinos (not likely but not ruled out) and if at least one of the heavy Majorana neutrinos has a mass of $\lesssim 100$ GeV, then the BR’s for these decays are: $BR(t \to b\ell_{i}^{+} \ell_{j}^{+} W^{-}) \sim 10^{-4}$ and $BR(W^{+} \to \ell_{i}^{+} \ell_{j}^{+} J J') \sim 10^{-7}$ if $m_{N} \sim 100$ GeV and $BR(t \to b\ell_{i}^{+} \ell_{j}^{+} J J') \sim BR(W^{+} \to \ell_{i}^{+} \ell_{j}^{+} J J') \sim 0.01$ if $m_{N} \lesssim 50$ GeV. Taking into account the present limits on the neutrino mixing parameters, we obtain more realistic values for these BR’s: $BR(t \to b\ell_{i}^{+} \ell_{j}^{+} W^{-}) \sim 10^{-6}$ and $BR(W^{+} \to \ell_{i}^{+} \ell_{j}^{+} J J') \sim 10^{-10}$ for $m_{N} \sim 100$ GeV and $BR(t \to b\ell_{i}^{+} \ell_{j}^{+} J J') \sim BR(W^{+} \to \ell_{i}^{+} \ell_{j}^{+} J J') \sim 10^{-6}$ for $m_{N} \lesssim 50$ GeV.

PACS numbers:

The recent discovery of neutrino oscillations which indicates mixing between massive neutrinos \[1\], was a major turning point in modern particle physics, since it stands as the first direct evidence for physics beyond the Standard Model (SM). Thus, it is now clear that the SM has to be expanded to include massive neutrinos that mix. Since there is still no understanding of the nature of these massive neutrinos, \textit{i.e.}, Majorana or Dirac-like, the extension of the SM can basically go either way. In particular, a simple way to consistently include sub-eV massive Majorana neutrinos in the SM is to add superheavy right-handed neutrinos with GUT-scale masses and to rely on the seesaw mechanism \[2\], which yields the desired light neutrinos mass scale: $m_{\nu} \sim M_{EW}^{2}/M_{GUT} \sim 10^{-2}$ eV, $M_{EW}$ being the electroweak (EW) scale. The seesaw mechanism, therefore, links neutrino masses with new physics at the GUT-scale, which is well motivated theoretically. On the other hand, a simple way to include massive Dirac neutrinos within the SM is to add Higgs-neutrinos Yukawa couplings which are more than 8 orders of magnitude smaller than the Higgs-electron one. Consequently, the Yukawa couplings of fermions (in the SM) unnaturally span over more than 13 orders of magnitude. Thus, within these simple extensions to the SM, the Majorana neutrinos seem to be favored from the theoretical point of view.

The fact that a Majorana mass term violates lepton number by two units, \textit{i.e.} $\Delta L = \pm 2$, has dramatic phenomenological signatures that can be used to distinguish Majorana neutrinos from Dirac neutrinos within many extensions of the SM. The most extensively studied process is neutrinoless double beta decay $(A, Z) \to (A, Z + 2) + e^{-} + e^{-}$ \[3\]. Also interesting are the $\Delta L = 2$ lepton-number-violating (LNV) processes in various high-energy collisions such as: $e^{-} \gamma \to 2 \ell$, $pp \to 2 \ell$, $e^{\pm} \gamma \to 2 \ell$, $e^{\pm} e^{-} \to 2 \ell$, $e^{-} e^{-} \to 2 \ell$, $e^{+} e^{-} \to 2 \ell$, and rare charged meson decays \[7\].

In this letter we explore two additional LNV decay channels of the real top-quark and of the real W-boson, to like-sign lepton pairs:

\begin{equation}
\text{(1)} \quad t \to b\ell_{i}^{+} \ell_{j}^{+} W^{-},
\end{equation}

\begin{equation}
\text{(2)} \quad W^{+} \to \ell_{i}^{+} \ell_{j}^{+} f \bar{f}'.
\end{equation}
These decays are induced by heavy Majorana neutrino exchanges and may, therefore, serve as important tests of the neutrino sector and as a possible evidence for the existence of Majorana-type heavy neutrinos with masses at the EW scale. The Feynman diagrams for these decays are depicted in Fig. 1. Both decays emanate from the same “kernel” process: \( W^\pm W^\mp \to \ell^+_i \ell^+_j \), with a t (or u)-channel exchange of a Majorana neutrino. This is the same kernel that induces double beta decay. However, in contrast to the double beta decay case, the decays in (1) and (2) are dominated by the exchanges of heavy (EW scale) neutrinos instead of the solar and atmospheric sub-eV neutrinos. Note that in (1) the top-quark decays to a real (i.e., on-shell) W-boson with a “wrong” charge. The W-boson in (2) can decay both purely leptonically if \( f' f' = \ell' \bar{\nu} \ell \) or semileptonically if \( f' f' = J J' \), where \( J \) and \( J' \) are light-quark jets coming from either a \( \bar{u}d \) or a \( \bar{c}s \) pair. In what follows we will concentrate only on the semileptonic decay of the W, since it is not possible (experimentally) to determine whether the leptonic channel is LNV or not. That is, a Dirac neutrino exchange will lead to the same observable final state with a \( W \) escapes detection, the leptonic channel signature is not unique to Majorana neutrino exchanges. Besides, the BR for the semileptonic channel is two times larger than the purely leptonic channel.

The kernel amplitude \( (W^\pm W^\mp \to \ell^+_i \ell^+_j) \) for the decays in (1) and (2) arises from the following Lagrangian term:

\[
\mathcal{L} = -\frac{g}{2\sqrt{2}} B_{\alpha iN} W^-_\mu \ell_i \gamma^\mu (1 - \gamma_5) n_\alpha + \text{H.c.}
\]  

where the index \( \alpha = 1 - 6 \) denotes the six Majorana neutrino states, i.e., three light ones and three heavy ones. Also, \( B_{\alpha iN} \) is a \( 3 \times 6 \) matrix, defined as \( B_{\alpha iN} \equiv \sum_{k=1}^{3} V_{\alpha i}^{L} V_{kN}^{L\ast} \), where \( V^L \) is the \( 3 \times 6 \) unitary mixing matrix of the left-handed charged leptons and \( U \) is the \( 6 \times 6 \) unitary mixing matrix in the neutrino sector, see e.g., (4) (in what follows \( n \) means \( n_\alpha \)). In the simplest scenario which relies on the classic seesaw mechanism (2), the couplings of heavy neutrinos \( (N) \) to SM particles, e.g., \( B_{\alpha iN} \) in (3), are highly suppressed by \( \sqrt{m_\nu/m_N} \), where \( m_\nu \) is the mass of the light neutrinos typically of the order of the solar or atmospheric neutrino masses. However, the possibility of non-seesaw realizations or internal symmetries in the neutrino sector, that may decouple the heavy-to-light neutrino mixing from the neutrino masses, cannot be excluded (3). This motivates us to adopt a purely phenomenological approach by assuming no a-priori relation between the mixing angles in \( B_{\alpha iN} \) and the neutrino masses. Within such a model independent approach, the elements in \( B_{\alpha iN} \) need only be bounded by existing model independent experimental constraints. For example, the 95% CL mass limits from LEP are \( m_N \gtrsim 80 - 90 \text{ GeV} \), depending on whether it couples to an electron, muon or a tau (14). For such heavy Majorana neutrinos and assuming the dominance of only one heavy neutrino \( N \) (see discussion below), the limits on its couplings to the charged leptons can be expressed in terms of the products \( \Omega_{\ell\nu} \equiv B_{\ell iN} B_{\ell'N} \) (see (4) and references therein). In particular, the limits on its flavor-diagonal couplings come from precision electroweak data, and at 90% CL are:

\[
\Omega_{ee} \leq 0.012 \, , \, \Omega_{\mu\mu} \leq 0.0096 \, , \, \Omega_{\tau\tau} \leq 0.016 \, ,
\]

while the limits on its flavor-changing couplings come from limits on rare flavor-violating lepton decays such as \( \mu \to e\gamma, \mu, \tau \to eee \) (3):

\[
|\Omega_{e\mu}| \leq 0.0001 \, , \, |\Omega_{e\tau}| \leq 0.02 \, , \, |\Omega_{\mu\tau}| \leq 0.02 \, .
\]
Using (3), the kernel amplitude is given by (including both t and u-channel diagrams):

\[ iM(W^-W^+ \to \ell^-_i \ell^-_j) = W^-W^+E_i E_j B_{in} B_{jn} m_n \bar{u}_j (1 + \gamma_5) \left[ \frac{\gamma^\mu \gamma^\nu}{p_{n(t)} - m_n^2 + im_n \Gamma_n} + \frac{\gamma^\nu \gamma^\mu}{p_{n(u)} - m_n^2 + im_n \Gamma_n} \right] v_i , \]  

where \( p_{n(t)} = p_W - p_{\ell_i} \), \( p_{n(u)} = p_W - p_{\ell_j} \), \( m_n \) and \( \Gamma_n \) are the Majorana neutrino t and u-channel 4-momenta, mass and total width, respectively. From (1) it is evident that for \( m_n^2 >> m_W^2 \) (recall that for the top-quark and W-boson decays the momentum transfer scale is of \( \mathcal{O}(m_W) \)), the amplitude is proportional to \( \prod_{k} \Gamma(k) \). On the other hand, for the sub-eV light neutrinos, i.e., \( m_n \) of order of the solar and atmospheric mass scales, the kernel amplitude is proportional to \( \prod_{k} \Gamma(k) \). Thus, these decays are by far dominated by the exchanges of the heavy Majorana neutrinos, if their masses are of \( \mathcal{O}(m_W) \). In what follows, we will, therefore, focus on the effect of the heavy Majorana neutrinos, with a further simplifying assumption that the kernel amplitude is dominated by an exchange of only one heavy Majorana neutrino, \( N \), which maximizes the quantity \( \prod_{k} \Gamma(k) \). In this way, the width for the partial decays \( N \to \ell \nu \) in (1) or the W-boson in (2):

\[ \Gamma \approx \frac{g^2}{64\pi m_W^2 m_N^2} \sum_k |B_{kN}|^2 . \]  

For the dominant decays of \( N \) we take \( N \to \ell_k W^\mp, \nu_k Z, \nu_k H \), where \( \nu_k, k = 1 - 3 \), are the three light sub-eV neutrinos. The partial widths for these decay channels are given by (see e.g., [4]):

\[ \sum_k \Gamma(N \to \ell_k W^\mp) \approx C(m_N^2 + 2m_W^2)(m_N^2 - m_W^2)^2 , \]

\[ \sum_k \Gamma(N \to \nu_k Z) \approx C(m_N^2 + 2m_Z^2)(m_N^2 - m_Z^2)^2 , \]

\[ \sum_k \Gamma(N \to \nu_k H) \approx C m_N^2 (m_N^2 - m_H^2)^2 , \]  

where

\[ C \equiv \frac{g^2}{64\pi m_W^2 m_N^2} \sum_k |B_{kN}|^2 . \]  

We note that our results depend very weakly on \( m_H \). Nonetheless, for definiteness, we will set \( m_H = 120 \text{ GeV} \) throughout our analysis. Also, the widths for the partial decays \( N \to \nu_k Z \) and \( N \to \nu_k H \) depend on the neutral couplings \( C_{\nu N} \) which appear in the interaction terms of \( Z \) and \( H \) with a pair of Majorana neutrinos. In Eq. (7) we have used the approximate relation between the charged and neutral couplings of \( N \) to the gauge-bosons: \( \sum_k |B_{kN}|^2 \approx \sum |C_{\nu N}|^2 \), see e.g., [4].

Let us now define a generic “reduced” amplitude squared:

\[ \bar{\mu}^2_{nn'} \equiv \frac{1}{\text{pol}} \sum_{\text{pol}} \mu_n \mu_{n'}^* , \]  

where \( \text{pol} \) is the number of polarization states of the decaying particle (\( \text{pol} = 2 \) and \( \text{pol} = 3 \) for the top-quark and the W-boson decays, respectively), \( \mu_n \) is the top-quark or W-boson decay amplitude for an exchange of a Majorana neutrino \( n \) (see Fig. 1), and \( n, n' = 1 - 6 \) are indices of the six Majorana neutrino states.

Then, using (9) and summing over all intermediate Majorana neutrino states, we obtain the total amplitude squared:

\[ |\bar{\mathcal{M}}|^2 = \sum_{n=1}^6 \bar{\mu}^2_{nn} + \sum_{n<n'} 2 \text{Re}(\mu^*_{nn'} ) , \]

and the decay width for either the top-quark in (1) or the W-boson in (2):

\[ \Gamma = \frac{(1 - \delta_{kk})}{2M(2\pi)^8} \int \prod_{k=1}^4 \frac{d^4 p_k}{2E_k} \delta^4(P - \sum_{k=1}^4 p_k) |\bar{\mathcal{M}}|^2 , \]  

(11)
where $M$ and $P$ are the mass and 4-momentum of the decaying particle, $i, j$ are flavor indices of the lepton pair in the final state of both decays and $p_k$ are the momenta of the final state particles. The reduced amplitude squared for the top-quark ($\bar{\mu}(t)_{nn'}^2$) and for the W-boson decays ($\bar{\mu}(W)_{nn'}^2$) are given by (neglecting the masses of the final state fermions)

$$\bar{\mu}(t)_{nn'}^2 = 8 A_{nn'}^{ij} p_t \cdot p_{i} \left\{ \frac{m_W^2}{m_W^2} p_t \cdot p_b \left( \frac{(p_W \cdot p_W')^2 - m_W^2}{m_W^2} \right) + \frac{2}{m_W^2} p_t \cdot p_W p_b \cdot p_W' \right\} + 2 \frac{m_W^2}{m_W^2} \left[ p_b \cdot p_W' - \frac{1}{m_W^2} p_b \cdot p_W p_W' \right],$$

$$\bar{\mu}(W)_{nn'}^2 = \frac{16}{3} A_{nn'}^{ij} p_t \cdot p_{i} \left\{ p_f \cdot p_{f'} + \frac{2}{m_W^2} p_W \cdot p_f p_W' \cdot p_{f'} \right\},$$

where

$$A_{nn'}^{ij} \equiv \left( \frac{g}{\sqrt{2}} \right)^6 |\Pi_W|^2 |B_{in}|^2 |B_{jn}|^2 |m_n|^2 \Pi_n \Pi_{n'}^*,$$

and $\Pi_x \equiv (p_x^2 - m_x^2 + i m_x \Gamma_x)^{-1}$, where $p_x$, $m_x$ and $\Gamma_x$ are the 4-momentum, mass and width of the particle $x$, respectively. Also, $p_W'$ which appears in (12) and (13) is the 4-momentum of the virtually exchanged W-boson in both the top-quark and W-boson decays (see Fig. 1).

Note that, within our assumption of a single-N dominated amplitude, we get:

$$|\mathcal{M}|^2 = \bar{\mu}_N^2,$$

$$A_{NN'}^{ij} = \left( \frac{g}{\sqrt{2}} \right)^6 |\Pi_W|^2 |\Pi_N|^2 m_N^2 |B_{in}|^2 |B_{jn}|^2.$$

We will first consider the case $m_N > m_W$ and then discuss the implications of a “light” $m_N$, $m_N < m_W$, on the top and W decays under investigation. In Fig. 2 we plot the BR’s for both the top-quark and the W-boson decays, scaled by the neutrino mixing parameters, i.e., setting $B_{ijN} = B_{ijN} = 1$, as a function of the Majorana neutrino mass, $m_N$, in the mass range $m_N > m_W$. We see that, for both decays, a sizable and experimentally accessible BR can arise only for $m_N$ values around 100 GeV, for which we obtain:

$$\frac{BR(t \rightarrow bW^- \ell^+_i \ell^+_j)}{|B_{in}|^2 |B_{jn}|^2} \sim 10^{-4},$$

$$\frac{BR(W^+ \rightarrow J^+ \ell^+_i \ell^+_j)}{|B_{in}|^2 |B_{jn}|^2} \sim 10^{-7}.$$

To obtain more realistic BR’s we can use the bounds on the neutrino mixing couplings in (14) and (15). For the W-boson decay (this decay is essentially insensitive to the heavy neutrino width), the largest BR subject to the constraints in (14) and (15) is of order of $10^{-10}$. This is too small to be observed at the Large Hadron Collider (LHC), where about $10^9 - 10^{10}$ inclusive on-shell W’s are expected to be produced through pp $\rightarrow W + X$, at an integrated luminosity of $O(100)$ fb$^{-1}$.

However, as was shown in (16), for on-shell production of N via $ud \rightarrow W^* \rightarrow \ell N$, the sensitivity to the heavy Majorana neutrino can be significantly enhanced. Indeed, in this case, the s-channel $W^*$ “decays”, as in (12), to $W^* \rightarrow J^+ \ell^+_i \ell^+_j$, by first decaying to an on-shell Majorana neutrino $W^* \rightarrow \ell N$, followed by $N \rightarrow \ell W \rightarrow \ell J^+ J^+$. The production of an on-shell $N$ substantially enhances the cross-section and makes this process, i.e., $ud \rightarrow W^* \rightarrow J^+ \ell^+_i \ell^+_j$, easily accessible at the LHC.

As will be shown below, a similar enhancement occurs in the cascade decay of an on-shell W, $W^+ \rightarrow \ell^+_i N \rightarrow \ell^+ J^+ J^+$ if $m_N < m_W$.

In the case of the top-quark decay $t \rightarrow bW^- \ell^+_i \ell^+_j$, taking $m_N \sim 100$ GeV and using the limits from (14) and (15), the BR’s for the various $\ell^+_i \ell^+_j$ channels are given in Table 1.

Note that, for $m_N < m_t$, the $BR(t \rightarrow bW^- \ell^+_i \ell^+_j)$ can be approximated by:
FIG. 2: The scaled BR (i.e., the BR with $B_{iN} = B_{jN} = 1$) for the flavor and lepton number violating decays $t \to bW^- \ell_i^+ \ell_j^+$ (solid line) and $W^+ \to J\bar{J}' \ell_i^+ \ell_j^+$ (dashed line), $i \neq j$ and $J$, $J'$=light jets, as a function the heavy neutrino mass, $m_N$. The right figure focuses on the range $m_N < m_t$.

TABLE I: $BR(t \to bW^- \ell_i^+ \ell_j^+)$ in the various $\ell_i \ell_j$ channels, subject to the bounds in [4] and [5] on the neutrino mixing parameters, for $m_N = 90$ and 100 GeV. For the flavor changing channels we set $\Omega_{\ell\ell'} = \sqrt{\Omega_{\ell\ell}} \times \sqrt{\Omega_{\ell'\ell'}}$ in case the limits on the couplings in [5] are weaker than the individual limits given in [4].

| $\ell_i \ell_j$ | $ee$ | $\mu\mu$ | $\tau\tau$ | $e\mu$ | $e\tau$ | $\mu\tau$ |
|----------------|-----|--------|--------|------|------|--------|
| $m_N = 90$ GeV | 1.4 | 1.1 | 1.9 | $1.1 \times 10^{-4}$ | 1.6 | 1.4 |
| $m_N = 100$ GeV | 0.6 | 0.5 | 0.8 | $0.4 \times 10^{-4}$ | 0.7 | 0.6 |

$$BR(t \to bW^- \ell_i^+ \ell_j^+) \approx BR(t \to b\ell_i^+ N) \times BR(N \to W^- \ell_j^+) + (i \leftrightarrow j \text{ for } i \neq j),$$

(17)

where, for $m_N \sim 100$ GeV we obtain:

$$BR(t \to b\ell_i^+ N) \frac{1}{|B_{iN}|^2} \sim 10^{-4},$$

(18)

and

$$BR(N \to W^- \ell_j^+) \sim 0.5 \times \frac{|B_{jN}|^2}{|B_{iN}|^2 + |B_{jN}|^2}.$$

(19)

We recall that the cross-section for $t\bar{t}$ production at the LHC is $\sim 850$ pb [16], yielding about $10^8$ $t\bar{t}$ pairs at an integrated luminosity of $O(100)$ fb$^{-1}$. Thus, a $BR(t \to bW^- \ell_i^+ \ell_j^+) \sim 10^{-6}$ that can arise in most $\ell_i^+ \ell_j^+$ channels (see Table I), should be accessible at the LHC. In particular, the flavor conserving channels $t \to bW^- e^+ e^+$ and $t \to bW^- \mu^+ \mu^+$ are expected to be more effective, since the channels involving the $\tau$-lepton will suffer from a low $\tau$ detection efficiency.
Let us now consider the case of a lighter $N$ with a mass $m_N < m_W$. Such a “light” Majorana neutrino is not excluded by LEP data if its couplings/mixings with the SU(2) leptonic doublets are small enough \[17\]. For example, $N$ can have a mass in the range $5 \text{ GeV} \lesssim m_N \lesssim 50 \text{ GeV}$ if $|B_{iN}|^2 \sim 10^{-4}$. In this mass range the 5-body cascade top decay $t \to bW^+ \to b\ell_i^+ N \to b\ell_i^+ \ell_j^+ J \bar{J}'$ (recall that $J \bar{J}'$ stands for a pair of light jets originating from a $u\bar{d}$ or $e\bar{e}$ pair) has a much larger width than the 4-body decay $t \to b\ell_i^+ \ell_j^+ W^-$, since the intermediate $N$ can not decay to an on-shell $W$. Thus, for $m_N < m_W$ both $t \to b\ell_i^+ \ell_j^+ J \bar{J}'$ and $W^+ \to \ell_i^+ \ell_j^+ J \bar{J}'$ originate from the cascade decay $W^+ \to \ell_i^+ N$ followed by $N \to \ell_j^+ J \bar{J}'$ and, for $BR(t \to bW^+) \sim 1$, they have equal branching ratios since:

$$BR(t \to b\ell_i^+ \ell_j^+ J \bar{J}') \sim BR(t \to bW^+) \times BR(W^+ \to \ell_i^+ N) \times BR(N \to \ell_j^+ J \bar{J}') + (i \leftrightarrow j \text{ for } i \neq j),$$

$$BR(W^+ \to \ell_i^+ \ell_j^+ J \bar{J}') \sim BR(W^+ \to \ell_i^+ N) \times BR(N \to \ell_j^+ J \bar{J}') + (i \leftrightarrow j \text{ for } i \neq j),$$

(20)

where the partial width for $W^+ \to \ell_i^+ N$ is:

$$\Gamma(W^+ \to \ell_i^+ N) = \frac{g^2}{96\pi} |B_{iN}|^2 m_W \cdot (2 - 3 \frac{m_N^2}{m_W^2} + \frac{m_N^6}{m_W^6}),$$

(21)

Also, in the mass range $10 \text{ GeV} \lesssim m_N \lesssim m_W$, the BR for $N \to \ell_j^+ J \bar{J}'$ is:

$$BR(N \to \ell_j^+ J \bar{J}') \approx \frac{\Gamma(N \to \ell_j^+ d\bar{u}) + \Gamma(N \to \ell_j^+ s\bar{c})}{\Gamma_N(Z) + \Gamma_N(H) + \Gamma_N(W)} \approx \frac{1}{4},$$

(22)

where

$$\Gamma_N(Z, H) = \sum_f \Gamma(N \to \nu_j Z^* (H^*) \to \nu_j f \bar{f}) ; f = u, d, c, s, b, e, \mu, \tau, \nu_e, \nu_{\mu}, \nu_{\tau};$$

$$\Gamma_N(W) = \sum_{(f\bar{f})} \Gamma(N \to \ell_j^+ W^+ \bar{\tau} \to \ell_j^+ (f \bar{f}^*)^\tau) ; (f \bar{f}^*)^\tau = (d\bar{u}), (s\bar{c}), (e\nu_e), (\mu\nu_{\mu}), (\tau\nu_{\tau}).$$

(23)

Thus, combining the scaled $BR(W^+ \to \ell_j^+ J \bar{J}') / |B_{iN}|^2$ calculated from (21) with the BR of the 3-body $N$ decay $BR(N \to \ell_j^+ J \bar{J}')$ given in (22), we plot in Fig. 3 the scaled BR’s for the top and $W$ decays in (20). We see that for e.g., $5 \text{ GeV} \lesssim m_N \lesssim 50 \text{ GeV}$ with $|B_{iN}|^2 \sim 10^{-4}$, not excluded by LEP \[17\], we obtain $BR(t \to b\ell_i^+ \ell_j^+ J \bar{J}') \sim BR(W^+ \to \ell_i^+ \ell_j^+ J \bar{J}') \gtrsim 10^{-6}$. For the $W$-decay, this rather large BR will be well within the reach of the LHC, which as mentioned above, is expected to produce $10^9 - 10^{10}$ inclusive on-shell $W$’s through $pp \to W + X$.

Before summarizing let us add a few comments:

- The heavy Majorana neutrino induced $\Delta L = 2$ branching ratios considered here i.e., $t \to b\ell_i^+ \ell_j^+ W^-$ (or $t \to b\ell_i^+ \ell_j^+ J \bar{J}'$ if $m_N < m_W$) and $W^+ \to \ell_i^+ \ell_j^+ J \bar{J}'$ are of course forbidden in the SM, thus a sighting of each constitutes a spectacular signal of lepton flavor violation (as well as LNV). Take for example the above top decay which produces 2 same-charge leptons, possibly of a different flavor, in addition to a wrong charge $W^-$ (recall that in the SM its dominant decay mode is $t \to bW^+$). Of course, these decays are not “stand-alone” since the $t$-quarks and $W$-bosons are created and decay in a specific accelerator and measured by a specific detector. Therefore, the background is highly accelerator and detector dependent - a detailed discussion of which is beyond the scope of this letter. As mentioned above, this top decay is unique since it has both a pair of same-charge leptons and a “wrong” charge $W$ i.e., $W^-$, unlike the positively charged $W$ produced in the dominant $t \to bW$ decay. The observation of this $\Delta L = 2$ top quark decay would therefore be a clear signal for LNV.

- A Majorana exchange is not necessarily the only mechanism leading to $\Delta L = 2$ processes. One can envisage, for instance, a situation in which another type of new physics contributes together with the heavy Majorana exchange. Viable examples are R-parity violating supersymmetry \[18\], or leptoquark exchanges \[19\]. In cases like these it is in principle possible to obtain destructive interference between the different mechanisms, thus evading the limits in \[21\] and \[22\], leaving the Majorana exchange significant for at least the top-quark decay considered here. Therefore, the rather sizable branching ratios in \[15\] obtained for $O(1)$ mixing angles cannot be excluded.
There are some discussions about a Super LHC (SLHC) in which the luminosity of the LHC would increase by about factor of 10. There is also some mention of an energy upgrade from $\sqrt{s} = 14$ TeV to 25-28 TeV, which may require a new machine. Such an upgrade in both luminosity and energy would yield more than an order of magnitude increase in the number of $t\bar{t}$ pairs and $W$’s produced, making the $O(10^{-6})$ BR of the top-quark decay in question easily accessible to this machine.

To summarize: we have discussed the $\Delta L = 2$ decays of the top-quark and of the $W$-boson, where both are mediated by a heavy Majorana neutrino $N$. Our main results appear in Figs. 2 and 3 and in the Table and are significant for both the top-quark case if $m_N < \sim 100$ GeV and the $W$-boson case if $m_N < m_W$.

Acknowledgments

S.B.S thanks the hospitality of the theory group in Brookhaven National Laboratory where part of this study was performed. The work of S.B.S. and of A.S. was supported in part by US DOE under Grants Nos. DE-FAG02-94ER40817 (USA) and DE-AC02-98CH10886 (BNL). The work of N.G.D and of J.J. was supported in part by US DOE under Grant No. DE-FAG02-96ER40969. G.E. would like to thank Vernon Barger, Tao Han and Tom Rizzo for helpful discussions. The research of G.E. was supported in part by the Israel Science Foundation and by the Fund for Promotion of Research at the Technion.

[1] For recent reviews see e.g.: K. Long, M. D. Messier and O. Yasuda, Nucl. Phys. Proc. Suppl. 155, 102 (2006); J. W. F. Valle, arXiv:hep-ph/0603223

[2] M. Gell-Mann, P. Ramond and R. Slansky, in Supergravity, ed. by P. Van Nieuwenhuizen and D.Z. Freedman (North Holland, Amsterdam, 1979), p.315; S.L. Glashow, in Quarks and Leptons, Cargese, 1979, ed. by M. Levy et al. (Plenum, New-York, 1980); T. Yanagida, in Proc. of the Workshop on the Unified Theory and Baryon Number in the Universe, ed. by O. Sawada and A. Sugamoto (KEK report 79-18, Tsukuba, Japan, 1979), p.95; R.N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980).

[3] For recent reviews see e.g.: A.S. Barabash, “Double beta decay experiments”, JINST 1, P07002 (2006) [arXiv:hep-ex/0602037]; P. Vogel, Prog. Part. Nucl. Phys. 57, 177 (2006).

[4] S. Bray, J.S. Lee and A. Pilaftsis, Phys. Lett. B628, 250 (2005).

[5] F. del Aguila, J.A. Aguilar Saavedra and R. Pittau, hep-ph/0606198

[6] H. Tso-Hsiu, C. Cheng-Rui and T. Zhi-Jian, Phys. Rev. D42, 2265 (1990); D.A. Dicus and D.D. Karatas, Phys. Rev. D44, 2033 (1991). A. Datta, M. Guchait and D.P. Roy, Phys. Rev. D47, 961 (1993); A. Ferrari et al., Phys. Rev. D62, 013001 (2000).

[7] A. Ali, A.V. Borisov and N.B. Zamorin, Eur. Phys. J. C21, 123 (2001).
[8] F.M.L. Almeida, Y.A. Countinho, J.A. Martins Simões and M.A.B. do Vale, Phys. Rev. D62, 075004 (2000).
[9] T. Han and B. Zhang, hep-ph/0604064
[10] M. Flanz, W. Rodejohann and K. Zuber, Phys. Lett. B473, 324 (2000), Erratum-ibid. B480, 418 (2000); A. Ali, A.V. Borisov and D.V. Zhuridov, hep-ph/0512005.
[11] W. Buchmüller and C. Greub, Nucl. Phys. B363, 345 (1991); F.M.L. Almeida, Y.A. Countinho, J.A. Martins Simões and M.A.B. do Vale, Eur. Phys. J. C30, 327 (2003).
[12] T.G. Rizzo, Phys Let. B116, 23 (1982) and hep-ph/9501261.
[13] See e.g., W. Buchmüller, C. Greub and P. Minkowski, Phys. Lett. B267, 395 (1991); A. Pilaftsis, Z. Phys. C55, 275 (1992), hep-ph/9001206; D. Tommasini, G. Barenboim, J. Bernabeu and C. Jarlskog, Nucl. Phys. B444, 451 (1994); J. Gluza, Acta. Phys. Polon. B33, 1735 (2002); G. Altarelli and F. Feruglio, New J. Phys. 6, 106 (2004).
[14] W.-M. Yao et al., Particle Data Group, J. of Phys. G33, 1 (2006) (URL: http://pdg.lbl.gov).
[15] S. Bergmann and A. Kagan, Nucl. Phys. B538, 368 (1999); for recent limits on sub-eV neutrinos from neutrinoless double-beta see e.g., A. Atre, V. Barger and T. Han, Phys. Rev. D71, 113014 (2005).
[16] See e.g., “Physics of electroweak gauge bosons” and “Heavy quarks and leptons”, in http://atlas.web.cern.ch/Atlas/GROUPS/PHYSICS/TDR/access.html.
[17] O. Adriani et al. (for the L3 Collaboration), Phys. Lett. B295, 371 (1992).
[18] J.D. Vergados, Phys. Lett. B184, 55 (1987) and and Phys. Rep. 361, 1 (2002). See however M. Hirsch, H.V. Klapdor-Kleingrothaus and S.G. Kovalenko, Phys. Rev. D53, 1329 (1996).
[19] M. Hirsch, H.V. Klapdor-Kleingrothaus and S.G. Kovalenko, Phys. Rev. D54, 4207 (1996); For other “contrived” or “fine-tuned” models see: G. Bélanger, F. Boudjemaa, D. London and H. Nadeau, Phys. Rev. D53, 6292 (1996), and O. Panella, M. Cannoni, C. Carimalo and Y.N. Srivastava, Phys. Rev. D65, 035005 and references therein.
[20] See e.g., F. Gianotti, Nucl. Phys. Proc. Suppl. 147, 23 (2005).