Heavy Majorana Neutrino Production at $e^−\gamma$ Colliders

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ABSTRACT

We study signatures of heavy Majorana neutrinos at $e^−\gamma$ colliders. Since these particles violate lepton number, they can give rise to the reactions $e^−\gamma \rightarrow W^−W^−e^+$ and $e^−\gamma \rightarrow W^+\mu^−\mu^−\nu$. The Standard Model background contains extra light neutrinos that escape detection, and can be reduced significantly after imposing appropriate kinematical cuts. We analyze the physics potential of an $e^−\gamma$ collider, with centre of mass energies $\sqrt{s} = 0.5–1$ TeV and a total integrated luminosity of 100 fb$^{-1}$, for detecting heavy Majorana neutrinos with masses $m_N = 100–400$ GeV. Assuming that heavy neutrinos couple predominantly to only one lepton flavour at a time, we find that an observable signal at the 5$\sigma$ level can be established for $|B_{eN}| > 7.4 \times 10^{-3}$. Instead, if no signal is observed, this will imply that $|B_{eN}| < 2.7 \times 10^{-3}$ and $|B_{\mu N}| < 0.05$ at the 90% confidence level.

1 Introduction

Neutrino oscillation experiments [1–3] have established the fact that the observed light neutrinos are not strictly massless, as predicted in the Standard Model (SM), but have non-zero tiny masses. One well-motivated framework for giving small masses to neutrinos is to extend the SM by adding one right-handed neutrino per family. In such a scenario, the right-handed neutrinos are singlets under the SM gauge group, and so are allowed to have large Majorana masses in the Lagrangian. Apart from the three observable neutrinos $\nu_{1,2,3}$, the resulting spectrum will contain three additional heavy neutrinos $N_{1,2,3}$, whose masses could range from the electroweak to the Grand Unified Theory (GUT) scale [4]. Direct searches at LEP put strong limits on the couplings of heavy Majorana neutrinos weighing less than about 100 GeV [5]. For larger heavy Majorana masses, the indirect limits are still significant and usually arise from the non-observation of sizeable quantum
effects on low-energy observables, as well as from the absence of lepton-flavour-violating decays [6–14], e.g. $\mu \to e\gamma$, $\mu \to eee$ and $\mu \to e$ conversion in nuclei etc.

In this paper we study the distinctive signatures of lepton number violation (LNV) at an $e^-\gamma$ collider, which are mediated by heavy Majorana neutrinos. In particular, we consider the LNV reactions $e^-\gamma \to W^-W^-e^+$ and $e^-\gamma \to W^+\mu^-\mu^-\nu$. These processes are strictly forbidden in the SM, and the corresponding background will always involve additional light neutrinos. An $e^-\gamma$ collider provides a clean environment to look for such particles, since the initial state $e^-\gamma$ has a definite non-zero lepton number and hence any LNV signal can easily be detected. There are already realistic proposals [15] for the construction of an $e^-\gamma$ machine, which will run at 0.5–1 TeV centre of mass system (CMS) energies and a total integrated luminosity of 100 fb$^{-1}$. Here, we will analyze the physics potential of such a machine to discover heavy Majorana neutrinos with masses between 100 and 400 GeV.

Our study has been structured as follows. After briefly reviewing the SM with right-handed neutrinos in Section 2, we discuss in Section 3 the LNV signals due to heavy Majorana neutrinos and the associated backgrounds. In Section 4 we present numerical results for the cross sections of the reactions $e^-\gamma \to W^-W^-e^+$ and $e^-\gamma \to W^+\mu^-\mu^-\nu$. In particular, taking into account the SM backgrounds, we place lower limits on the heavy neutrino couplings to the charged leptons and the $W^\pm$ bosons, by requiring that an observable LNV signal is obtained. Finally, our conclusions are summarized in Section 5.

2 The Standard Model with Right-Handed Neutrinos

We briefly review the relevant low-energy structure of the SM modified by the presence of right-handed neutrinos. The Lagrangian describing the neutrino masses and mixings reads:

$$\mathcal{L}_\nu^{\text{mass}} = -\frac{1}{2} \left( \begin{array}{c} \nu^0_L \\ (\nu^0_R)^c \end{array} \right) \left( \begin{array}{cc} 0 & m_D \\ m_D^T & M_R \end{array} \right) \left( \begin{array}{c} (\nu^0_L)^c \\ \nu^0_R \end{array} \right) + \text{h.c.},$$  \hspace{1cm} (2.1)

where $\nu^0_L = (\nu_e L, \nu_\mu L, \nu_\tau L)^T$ and $\nu^0_R = (\nu_e R, \nu_\mu R, \nu_\tau R)^T$ collectively denote the left- and right-handed neutrino fields in the weak basis, and $m_D$ and $M_R$ are 3 $\times$ 3 complex matrices. The latter obeys the Majorana constraint $M_R = M_R^T$ and, without loss of generality, can be assumed to be diagonal and positive. The weak neutrino states are related to the Majorana mass eigenstates through:

$$\begin{pmatrix} \nu_L \\ N_L \end{pmatrix} = U^T \begin{pmatrix} \nu^0_L \\ (\nu^0_R)^c \end{pmatrix},$$  \hspace{1cm} (2.2)

where $U$ is a 6 $\times$ 6 unitary matrix.

In typical seesaw scenarios [4], the Dirac mass terms $m_D$ are expected to be around the electroweak scale, e.g. $m_D \sim M_l$ or $m_D \sim M_u$, where $M_l$ ($M_u$) is the charged lepton (up-quark) mass matrix, whilst the Majorana mass $M_R$ being singlet under the SM gauge group may be very large, close to the GUT scale. Seesaw models can explain the smallness
of the observed light neutrino masses, generically predicted to be $\sim m_D^2/M_R$, through the huge hierarchy between $m_D$ and $M_R$. However, the couplings of the heavy neutrinos to SM particles will be generically highly suppressed $\sim m_D/M_R$, thus rendering the direct observation of heavy Majorana neutrinos impossible.

Another and perhaps phenomenologically more appealing solution to the smallness of the light neutrino masses may arise due to approximate flavour symmetries that may govern the Dirac and Majorana mass matrices $m_D$ and $M_R$ [16–18]. In such models, the heavy-to-light Majorana couplings are also $\sim m_D/M_R$, but they can be completely unrelated to the light neutrino mass matrix $m_\nu$, which is determined by the relation:

$$m_\nu = -m_D M_R^{-1} m_D^T.$$

In such non-seesaw models, one may have $m_D \sim M_l$ and $M_R \sim 100 \text{ GeV}$, without being in contradiction with neutrino data. However, as we will see below, these models are constrained by electroweak precision data and other low-energy lepton-flavour/number-violating observables. Our study of LNV signatures at an $e^-\gamma$ collider will focus on such non-seesaw realizations.

For our subsequent phenomenological discussion, we now exhibit the interaction Lagrangians of the heavy neutrinos to $W^\pm$, $Z$ and Higgs ($H$) bosons [17]:

$$L_W = -\frac{g}{\sqrt{2}} W^-_{\mu}  \gamma^\mu P_L B_{ij} \left( \nu_N \right)_j + \text{h.c.},$$

$$L_Z = -\frac{g}{4e_w} Z_{\mu} \left( \bar{\nu} \ N \right)_i \gamma^\mu \left( i \text{Im} C_{ij} - \gamma_5 \text{Re} C_{ij} \right) \left( \nu_N \right)_j,$$

$$L_H = -\frac{g}{4M_W} H \left( \bar{\nu} \ N \right)_i \left[ (m_i + m_j) \text{Re} C_{ij} - i\gamma_5 (m_i - m_j) \text{Im} C_{ij} \right] \left( \nu_N \right)_j,$$

where $m_{i,j}$ (with $i, j = 1, 2, \ldots 6$) denote the physical light and heavy neutrino masses, and

$$B_{ij} = \sum_{k=1}^{3} V_{Lki}^* U_{kj}, \quad C_{ij} = \sum_{k=1}^{3} U_{ki} U_{kj}^*.$$

In (2.6), $V_L$ is a $3 \times 3$ unitary matrix relating the weak to mass eigenstates of the left-handed charged leptons.

As was mentioned above, the mixing elements $B_{ij}$ can be constrained from LEP and low-energy electroweak data [6–14]. Following [19], we define

$$\Omega_{ll'} \equiv \delta_{ll'} - \sum_{i=1}^{3} B_{li} B_{l'i}^* = \sum_{i=1}^{3} B_{li} B_{l'i}^*.$$

The above definition is a generalization of the Langacker–London parameters $(s_{L}^{\nu_e,\mu,\tau})^2$ [6], with the identification: $\Omega_{ll} = (s_{L}^{\nu_l})^2$. At the 90% confidence level (CL), the allowed values for the parameters $\Omega_{ll}$ are [12]:

$$\Omega_{ee} \leq 0.0054, \quad \Omega_{\mu\mu} \leq 0.0096, \quad \Omega_{\tau\tau} \leq 0.016.$$
These limits depend only weakly on the heavy neutrino masses.

Lepton-flavour-violating processes, e.g. $\mu \rightarrow e\gamma$ [10], $\mu \rightarrow eee$, $\tau \rightarrow eee$ [11, 13, 14], can be induced by heavy neutrino quantum effects. Such processes do depend on the heavy neutrino masses and Yukawa couplings [11]. For $m_N \gg M_W$ and $m_D \ll M_W$, the limits are

$$|\Omega_{e\mu}| \lesssim 0.0001, \quad |\Omega_{e\tau}| \lesssim 0.01, \quad |\Omega_{\mu\tau}| \lesssim 0.01.$$ \hspace{1cm} (2.9)

Finally, we present the partial decay widths of a heavy Majorana neutrino $N$ for its dominant decay channels [17]:

$$\Gamma(N \rightarrow l^\pm W^\mp) = \frac{\alpha_w |B_{lN}|^2}{16 M_W^2 m_N^2} (m_N^2 - M_W^2)^2 (m_N^2 + 2M_W^2),$$ \hspace{1cm} (2.10)

$$\Gamma(N \rightarrow \nu_i Z) = \frac{\alpha_w |C_{\nu_i N}|^2}{16 M_W^2 m_N^2} (m_N^2 - M_Z^2)^2 (m_N^2 + 2M_Z^2),$$ \hspace{1cm} (2.11)

$$\Gamma(N \rightarrow \nu_i H) = \frac{\alpha_w |C_{\nu_i N}|^2}{16 M_W^2 m_N^2} (m_N^2 - M_H^2)^2,$$ \hspace{1cm} (2.12)

where $\alpha_w = g^2/(4\pi)$. Notice that up to negligible corrections $O(m_D^4/M_H^4)$, it can be shown that $\sum_i |B_{lN}|^2 = \sum_i |C_{\nu_i N}|^2$. Hence, for heavy neutrinos, with $m_N \gg M_H$, decaying into all charged leptons and light neutrinos, one obtains the useful relation among the branching fractions: $B(N \rightarrow l^+W^-) = B(N \rightarrow l^-W^+) = B(N \rightarrow \nu Z) = B(N \rightarrow \nu H) = 1/4$.

### 3 Lepton Number Violating Signatures

Previous analyses of single heavy neutrino production have mainly been focused on $e^+e^-$ linear colliders [19–21], where the process $e^+e^- \rightarrow N\nu \rightarrow l^\pm W^\mp\nu$ is considered. Such a production channel, however, has two problems not associated with an $e^+\gamma$ collider. Firstly, since the light neutrinos escape detection with their chirality undetermined, the possible Majorana nature of the heavy neutrinos has very little effect on the signal, which makes the reduction of the contributing SM background very hard. Secondly, an observable signal would require a heavy neutrino coupling to the electron and the $W^{\pm}$ bosons of reasonable strength, i.e. $B_{eN} \gtrsim 10^{-2}$.

Another possible option which might have the same advantages as an $e^-\gamma$ collider is a $\gamma\gamma$ collider. Here, the heavy Majorana neutrinos can be produced and observed via $\gamma\gamma \rightarrow W^+Nl^- \rightarrow W^+W^+l^-l^-$ [22]. However, like the case of the $e^+e^-$ collider, one still has to assess the impact of the contributing SM background.

Depending on the relative strength of the heavy neutrino coupling to the electron $B_{eN}$, we consider the two reactions: (i) $e^-\gamma \rightarrow W^-W^+l^+$ and (ii) $e^-\gamma \rightarrow W^+l^-l^-\nu$. Both processes are manifestations of LNV, and as such, they allow to eliminate major part of the SM background. The process (i) dominates for relatively large values of the mixing factor $B_{eN}$, i.e. $B_{eN} \gtrsim 10^{-2}$, whereas the process (ii) becomes only relevant if $B_{eN}$ is unobservably small, e.g. $B_{eN} \lesssim 10^{-3}$. 

Figure 1: Feynman diagrams for the process $e^-\gamma \rightarrow W^-W^-l^+$.  

3.1 $B_{eN} \neq 0$

We first consider the process $e^-\gamma \rightarrow W^-W^- l^+$, which becomes the dominant signal for $B_{eN} \gtrsim 10^{-2}$. The Feynman diagrams pertinent to this process are shown in Fig. 1 \(^*\). We improve upon an earlier study of this reaction \([24]\), by carefully considering the contributing SM background (see our discussion in Section 3.3).

The reaction $e^-\gamma \rightarrow W^-W^- l^+$ is dominated by the graphs (a) and (b) of Fig. 1, where the heavy Majorana neutrinos occur in the $s$-channel. Considering therefore the $2 \rightarrow 2$ subprocess $e^-\gamma \rightarrow W^-N$, Fig. 2 gives the cross section $\sigma(e^-\gamma \rightarrow W^-N)$ as functions of $\sqrt{s}$ and $m_N$. The differential cross sections related to Fig. 2, as well as to Figs. 4 and 7, have been calculated using the FeynCalc package \([25]\). We have also checked that our results are in agreement with \([24]\).

Approximate values for the cross section $\sigma(e^-\gamma \rightarrow W^-W^- l^+)$ can be obtained by multiplying the $y$-axis values from Fig. 2 with the branching fraction $B(N \rightarrow W^- l^+)$. However, as well as ignoring contributions from off-shell heavy neutrinos, this approximation makes the imposition of kinematical cuts on the three-momentum of the final charged lepton rather difficult. Therefore, in Section 4, we calculate the LNV signatures by considering the complete $2 \rightarrow 3$ process, as it is described by the full set of Feynman diagrams depicted in Fig. 1. To this end, we have extended CompHEP \([26]\) to include heavy Majorana neutrino interactions.

3.2 $B_{eN} = 0$

If the coupling of the heavy neutrino $N$ to the electron is either zero or very small, then the dominant process is $e^-\gamma \rightarrow W^+ l^-l^-\nu$, where $l = \mu, \tau$. The Feynman diagrams for this

\(^*\)All diagrams, including Figs. 3 and 5, are produced with Axodraw \([23]\).
process are shown in Fig. 3. Although the cross section of the new reaction is much smaller in this case, so is the background. Hence, the predicted LNV signal could still be within observable reach. This is a major benefit compared to an $e^+e^-$ collider, which requires sizeable heavy neutrino couplings to the electron for an observable signal. To the best of our knowledge, this is a novel possibility which has not been investigated before.

The cross section for the LNV reaction $e^-\gamma\rightarrow W^-N$ is dominated by the resonant $s$-channel exchange graphs (a)–(c) of the heavy Majorana neutrino $N$, as shown in Fig. 3. Thus, we present in Fig. 4 the heavy neutrino production cross section $\sigma(e^-\gamma\rightarrow N\mu^-\nu)$ as functions of $\sqrt{s}$ and $m_N$. The 4-dimensional phase-space integration involved in the $2 \rightarrow 3$ process was performed numerically using Bases [27]. The cross section values obtained for this process are about 2 orders of magnitude smaller than those presented in Fig. 2, for the reaction $e^-\gamma\rightarrow W^-N$. An approximate estimate of $\sigma(e^-\gamma\rightarrow W^+\mu^-\nu)$ may be obtained by $\sigma(e^-\gamma\rightarrow N\mu^-\nu) \times B(N\rightarrow W^+\mu^-)$. As was mentioned in Section 3.1, this method has certain limitations when differential cross sections are considered. For this reason, the LNV signal in Section 4 is computed using the full set of the diagrams displayed in Fig. 3 and our extended version of CompHEP that includes heavy Majorana neutrino interactions.

### 3.3 The Standard Model Background

The LNV reactions we have been considering are strictly forbidden in the SM. The contributing SM background that may mimic the signal will always involve additional light neutrinos. Specifically, the SM background arises from the resonant production and decay

![Figure 2: Numerical estimates of the cross section $\sigma(e^-\gamma\rightarrow W^-N)$ as functions of $\sqrt{s}$ (left frame) and $m_N$ (right frame), with $B_eN = 0.07$.](image-url)
Figure 3: Feynman diagrams for the process $e^-\gamma \rightarrow W^+l^-\nu$, with $B_{eN} = 0$.

Figure 4: Numerical estimates of the cross section $\sigma(e^-\gamma \rightarrow N\mu^-\nu)$ versus $\sqrt{s}$ (left frame) and $m_N$ (right frame), where $B_{eN} = 0$, $B_{\mu N} = 0.1$ and an infrared angle cut $-0.99 \leq \cos(\theta_{e^-\mu^-}) \leq 0.99$ is used.
The background to the first LNV reaction $e^-\gamma \rightarrow W^-W^-l^+\nu_l$ originates from the decay $W^+ \rightarrow l^+\nu_l$. Correspondingly, the background to the second LNV process $e^-\gamma \rightarrow W^+l^-l^-\nu$ results from the leptonic decays $W^- \rightarrow l^-\bar{\nu}_l$ for both of the $W^-$ bosons in the final state. To compute the background process $e^-\gamma \rightarrow W^-W^-W^+\nu_e$, we first use CompHEP [26], which also generates the appropriate weighted events, and then use PYTHIA [28] interfaced via the CPyth program [29]. The last step is necessary in order to properly describe the Lorentz-boosted $W^\pm$-boson decay products on which appropriate kinematical cuts were placed. These kinematical cuts will be presented in detail in Section 4. Although our method uses the branching fractions for the $W$ decays [30] as an input, it proves by far more practical than generating events for the complete set of $2 \rightarrow 7$ background processes, where a vast number of off-resonant amplitudes give negligible contributions.

It is important to clarify at this point that observation of $e^-\gamma \rightarrow W^+l^-l^-\nu$ does not constitute by itself a signature for LNV, even if $l \neq e$. If the heavy neutrino couples to the electron as well as to the muon or tau lepton, then the occurrence of such a reaction will only signify lepton flavour violation. If one now assumes that $e^-\gamma \rightarrow W^-W^-l^+$ is not observed, there are two possibilities that could result in an observation of $e^-\gamma \rightarrow W^+l^-l^-\nu$. Either the heavy neutrino is a Majorana particle that does not couple predominantly to the electron, or it is a Dirac particle that does. To distinguish between these two possibilities, one has to look for $e^-\gamma \rightarrow W^-W^+l^-$, where both $W^\pm$ bosons decay hadronically. This process cannot occur, unless the heavy neutrino has a relevant coupling to the electron. For a heavy Dirac neutrino, however, the latter will always be larger than $e^-\gamma \rightarrow W^+l^-l^-\nu$. Therefore, if the
reactions $e^-\gamma \rightarrow W^-W^+l^\pm$ were not detected, observation of $e^-\gamma \rightarrow W^+l^-l^-\nu$ would be a clear manifestation of a LNV signature, which is mediated by a heavy Majorana neutrino.

4 Numerical Results

There are many theoretical parameters in the SM with right-handed neutrinos that could vary independently, such as heavy Majorana neutrino masses and couplings. In addition, the machine parameters of any future $e^-\gamma$ collider are still under discussion. Since it would be counter-productive to explore all possible situations that may occur, we will, instead, analyze two representative scenarios for each of the processes: (i) $e^-\gamma \rightarrow W^-W^-e^+$ and (ii) $e^-\gamma \rightarrow W^+\mu^-\mu^-\nu$.

In our analysis, we take $|B_{eN}|^2 \approx |C_{\nu_iN}|^2$, i.e. we assume that only one flavour will couple predominantly to the heavy Majorana neutrino $N$. The Higgs-boson mass value $M_H = 120$ GeV is used throughout our estimates, but our results depend only very weakly on $M_H$. Finally, we use the global kinematical cut,

$$-0.99 \leq \cos \theta_{e^-l^\pm} \leq 0.99,$$  \hspace{1cm} (4.1)

for the angle $\theta_{e^-l^\pm}$ between the $e^-$ in the initial state and the final-state charged leptons $l^\pm$. The global cut (4.1) was used to ensure that the produced charged leptons $l^\pm$ are detected.

4.1 $e^-\gamma \rightarrow W^-W^-e^+$

We will consider the simplest case where $B_{eN}$ is sizeable, but $B_{\mu N} = B_{\tau N} = 0$. We analyze two scenarios:

Scenario (1): $\sqrt{s} = 500$ GeV, $m_N = 200$ GeV, $B_{eN} = 0.07$, predicting $\Gamma_N = 0.04$ GeV and $B(N \rightarrow W^-e^+) \approx 0.3$,

and

Scenario (2): $\sqrt{s} = 1$ TeV, $m_N = 400$ GeV, $B_{eN} = 0.07$, predicting $\Gamma_N = 0.4$ GeV and $B(N \rightarrow W^-e^+) \approx 0.25$.

If the two $W^-$ bosons decay hadronically, the LNV signal will have no missing momentum, whilst the SM background $e^-\gamma \rightarrow W^-W^-e^+\nu_e\nu_e$ has two light neutrinos that will escape detection. Making use of this fact, one may apply the kinematical cut on the missing transverse momentum $p_T$:

$$p_T \leq p_T^{\text{max}} = 10 \text{ GeV}.$$  \hspace{1cm} (4.2)

This has no effect on the signal, but reduces the background significantly. The value 10 GeV has been chosen simply as a reasonable limit of the detector resolution [15]. Figure 6 shows the dependence of the background on $p_T^{\text{max}}$. The inaccuracies for low values of $p_T^{\text{max}}$ that
are apparent in Fig. 6 should be regarded as numerical artifacts due to the small fraction of events that have passed the $p_T^{\text{max}}$ selection criterion.

In Table 1 we present comparative numerical values of cross sections for the signal and the background in Scenarios (1) and (2), with and without the $p_T$ cut. Our numerical estimates also include the branching fraction for the two $W^-$ bosons to decay into hadrons. Table 1 shows that the background can be drastically reduced to an almost unobservable level. For a given scenario, this enables one to place sensitivity limits to the mixing factor $B_{eN}$. Since the LNV cross sections exhibited in Table 1 are roughly equal in Scenarios (1) and (2), the limits on $|B_{eN}|$ will be very similar. For a 5σ discovery, we require $S/\sqrt{S+B} \geq 5$, where $S$ is the number of signal events and $B$ the number of background events. Assuming an integrated luminosity of 100 fb$^{-1}$, this requires $|B_{eN}| \geq 7.4 \times 10^{-3}$. For a 3σ signal, one would need $|B_{eN}| \geq 4.6 \times 10^{-3}$. Finally, if no signal is observed, this will place the upper limit $|B_{eN}| \leq 2.7 \times 10^{-3}$ at 90% CL. Observe that the sensitivity of an $e^-\gamma$ collider to $B_{eN}$ is better than that of an $e^+e^-$ linear collider [19] due to the far smaller SM background.

4.2 \( e^-\gamma \rightarrow W^+\mu^-\mu^-\nu \)

As was mentioned in Section 3, this process becomes relevant, only when the mixing factor $|B_{eN}| \lesssim 10^{-3}$. We therefore consider cases where $B_{\mu N}$ is sizeable, but $B_{eN} = B_{\tau N} = 0$. As
Table 1: Cross sections (in fb) for the process $e^-\gamma \rightarrow W^-W^-e^+$ and its SM background. The branching fraction $B(W^-W^- \rightarrow \text{hadrons})$ is included.

| Process                        | Scenario (1) | Scenario (2) |
|--------------------------------|--------------|--------------|
| Signal (without cut)           | 22.9         | 23.3         |
| Background (without cut)       | 0.11         | 1.54         |
| Signal (with cut)              | 22.9         | 23.3         |
| Background (with cut)          | 0.002        | 0.01         |

before, we analyze the following two scenarios:

Scenario (3): $\sqrt{s} = 500$ GeV, $m_N = 200$ GeV, $B_{\mu N} = 0.1$, yielding $\Gamma_N = 0.08$ GeV and $B(N \rightarrow W^+\mu^-) \approx 0.3$,

and

Scenario (4): $\sqrt{s} = 1$ TeV, $m_N = 400$ GeV, $B_{\mu N} = 0.1$, yielding $\Gamma_N = 0.8$ GeV and $B(N \rightarrow W^+\mu^-) \approx 0.26$.

Since the heavy Majorana neutrino will have a small decay width, one expects the invariant mass $m_{\mu^-W^+}$ of one of the muons with the $W^+$ boson to be very close to the mass of the heavy neutrino. The other muon, which does not come from the decaying heavy neutrino, will generally have a preference to a small scattering angle from the direction of the incoming photon, due to the infrared property of the exchanged muon in the $t$-channel. This last feature is also shown in Fig. 7. In view of all the above reasons, the following cuts prove very effective:

$$m_{\mu^-W^+} = m_N \pm \Delta m,$$

for one of the muons,

$$\cos\theta_{e^-\mu^-} \leq \cos\theta_{\text{min}} = -0.5,$$

for the other muon. (4.3)

Here $\Delta m = 10$ GeV is again set by the detector resolution, which is much greater than the actual heavy neutrino decay width $\Gamma_N$. Figure 8 shows how the background would be affected by different choices of these cuts.

Table 2: Cross sections (in fb) for $e^-\gamma \rightarrow W^+\mu^-\mu^-\nu$ and its SM background, where $B(W^+ \rightarrow \text{hadrons})$ is included.
Figure 7: Differential cross section for the process $e^- \gamma \rightarrow N\mu^-\nu$ ($B_{eN} = 0$, $B_{\mu N} = 0.1$). Our cuts veto the region to the right of the dotted line.

Figure 8: The dependence of the background cross section $\sigma(e^-\gamma \rightarrow W^+\mu^-\nu_e\bar{\nu}_\mu) \times B(W^+ \rightarrow \text{hadrons})$ on the kinematical cuts $\Delta m$ (left panel) and $\cos\theta_{\text{min}}$ (right panel).
In Table 2 we summarize our results for the signal and the background, before and after the kinematical cuts (4.3) have been implemented. In particular, it can be seen that the selected kinematical cuts were very effective to drastically reduce the background by 2 orders of magnitude, without harming much the signal. The limits that can be derived for Scenarios (3) and (4) are very similar. To make a $5\sigma$ discovery with $|B_{\mu N}| = 0.1$, one would need a total integrated luminosity of about 200 fb$^{-1}$, because of the low signal cross sections. Instead, if the expected total integrated luminosity of 100 fb$^{-1}$ is assumed, observation of a $3\sigma$ signal will require a coupling of $|B_{\mu N}| \geq 0.08 \sim 0.09$. Finally, the absence of any LNV signal would put the upper limit $|B_{\mu N}| \leq 0.05$ at the 90% CL. It is important to remark that the sensitivity limits that can be placed on $|B_{\mu N}|$ at an $e^{-}\gamma$ collider will be much higher than the previously considered studies at $e^{+}e^{-}$ and $\gamma\gamma$ colliders [19, 22].

5 Conclusions

We have analyzed the phenomenological consequences of electroweak-scale heavy Majorana neutrinos at an $e^{-}\gamma$ collider. We have found that heavy Majorana neutrinos, with masses $m_N = 100$–400 GeV and couplings $B_{eN} \sim 10^{-2}$ and $B_{\mu N} \sim 10^{-1}$, will become easily observable at an $e^{-}\gamma$ collider, with CMS energies $\sqrt{s} = 0.5$–1 TeV and a total integrated luminosity of 100 fb$^{-1}$. Specifically, we have computed the cross sections for the LNV reactions $e^{-}\gamma \rightarrow W^{-}W^{-}e^{+}$ and $e^{-}\gamma \rightarrow W^{+}\mu^{-}\mu^{-}\nu$. In our analysis, we have also considered the contributing SM background, which involves additional light neutrinos in the final state. After imposing realistic missing $p_T$, invariant mass and angle cuts, we have been able to suppress the SM background by 2 orders of magnitude to an unobservable level, without harming much the signal.

Observation of the reaction $e^{-}\gamma \rightarrow W^{-}W^{-}e^{+}$ would be a clear signal for LNV and hence for unraveling the possible Majorana nature of heavy neutrinos. The process $e^{-}\gamma \rightarrow W^{+}\mu^{-}\mu^{-}\nu$ is also a clear signal for LNV in the absence of observations of $e^{-}\gamma \rightarrow W^{+}W^{-}\mu^{-}$. The search for heavy Majorana neutrinos at an $e^{-}\gamma$ collider adds particular value to analogous searches at an $e^{+}e^{-}$ linear collider. Although the latter may be built first, it would be more difficult to discern the Dirac or Majorana nature of possible heavy neutrinos. Another obstacle facing $e^{+}e^{-}$ colliders is the large irreducible SM background which has similar kinematical characteristics as the signal. Instead, an $e^{-}\gamma$ collider provides a cleaner environment and has a unique potential to identify the possible Majorana nature of the heavy neutrinos, as well as probe lower coupling strengths thanks to its highly reducible background.

The analysis presented in this paper may contain some degree of model dependence. For example, models that include new heavy gauge bosons may add new relevant LNV interactions to the Lagrangian. As long as these additional contributions are negligible, our results will be robust. In conclusion, an $e^{-}\gamma$ collider would provide valuable information about the properties of possible heavy neutrinos that could be relevant for low-scale
resonant leptogenesis [31], and also shed light on the structure of the light neutrino masses and mixings.

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References

[1] S. Fukuda et al. [Super-Kamiokande Collaboration], Phys. Lett. B 539, 179 (2002) [arXiv:hep-ex/0205075]; Phys. Rev. Lett. 86, 5656 (2001) [arXiv:hep-ex/0103033]; Phys. Rev. Lett. 86, 5651 (2001) [arXiv:hep-ex/0103032]; Y. Fukuda et al. [Super-Kamiokande Collaboration], Phys. Rev. Lett. 82, 2644 (1999) [arXiv:hep-ex/9812014]; Phys. Rev. Lett. 82, 1810 (1999) [arXiv:hep-ex/9812009].

[2] S. N. Ahmed et al. [SNO Collaboration], Phys. Rev. Lett. 92, 181301 (2004) [arXiv:nucl-ex/0309004]; Q. R. Ahmad et al. [SNO Collaboration], constraints Phys. Rev. Lett. 89, 011302 (2002) [arXiv:nucl-ex/0204009]; Phys. Rev. Lett. 89, 011301 (2002) [arXiv:nucl-ex/0204008]; Phys. Rev. Lett. 87, 071301 (2001) [arXiv:nucl-ex/0106015].

[3] K. Eguchi et al. [KamLAND Collaboration], Phys. Rev. Lett. 90, 021802 (2003) [arXiv:hep-ex/0212021].

[4] P. Minkowski, Phys. Lett. B 67, 421 (1977); M. Gell-Mann, P. Ramond and R. Slansky in Supergravity, p. 315, edited by F. Nieuwenhuizen and D. Friedman, North Holland, Amsterdam, 1979; T. Yanagida, Proceedings of the Workshop on the Unified Theories and Baryon Number of the Universe, edited by O. Sawada and A. Sugamoto, KEK, Japan 1979; R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980).

[5] P. Achard et al. [L3 Collaboration], Phys. Lett. B 517 (2001) 67 [arXiv:hep-ex/0107014].

[6] P. Langacker and D. London, Phys. Rev. D 38, 886 (1988).

[7] J. G. Korner, A. Pilaftsis and K. Schilcher, Phys. Lett. B 300, 381 (1993) [arXiv:hep-ph/9301290]; J. Bernabeu, J. G. Korner, A. Pilaftsis and K. Schilcher, Phys. Rev. Lett. 71, 2695 (1993) [arXiv:hep-ph/9307295].

[8] C. P. Burgess, S. Godfrey, H. Konig, D. London and I. Maksymyk, Phys. Rev. D 49, 6115 (1994) [arXiv:hep-ph/9312291].

[9] E. Nardi, E. Roulet and D. Tommasini, Phys. Lett. B 327, 319 (1994) [arXiv:hep-ph/9402224].

[10] T. P. Cheng and L. F. Li, Phys. Rev. Lett. 45, 1908 (1980).

[11] A. Ilakovac and A. Pilaftsis, Nucl. Phys. B 437, 491 (1995) [arXiv:hep-ph/9403398].

[12] S. Bergmann and A. Kagan, Nucl. Phys. B 538, 368 (1999) [arXiv:hep-ph/9803305].

[13] J. I. Illana and T. Riemann, Phys. Rev. D 63, 053004 (2001) [arXiv:hep-ph/0010193].

[14] G. Cvetic, C. Dib, C. S. Kim and J. D. Kim, Phys. Rev. D 66, 034008 (2002) [Erratum-ibid. D 68, 059901 (2003)] [arXiv:hep-ph/0202212].
[15] B. Badelek et al. [ECFA/DESY Photon Collider Working Group], Int. J. Mod. Phys. A 19 (2004) 5097 [arXiv:hep-ex/0108012].

[16] E. Witten, Nucl. Phys. B 268 (1986) 1642; R.N. Mohapatra and J.W.F. Valle, Phys. Rev. D 34 (1986) 1642.

[17] A. Pilaftsis, Z. Phys. C 55 (1992) 275 [arXiv:hep-ph/9901206].

[18] J. Gluza, Acta Phys. Polon. B 33 (2002) 1735 [arXiv:hep-ph/0201002]. G. Altarelli and F. Feruglio, New J. Phys. 6 (2004) 106 [arXiv:hep-ph/0405048].

[19] F. del Aguila, J. A. Aguilar-Saavedra, A. Martinez de la Ossa and D. Meloni, Phys. Lett. B 613, 170 (2005) [arXiv:hep-ph/0502189]; F. del Aguila and J. A. Aguilar-Saavedra, JHEP 0505, 026 (2005) [arXiv:hep-ph/0503026].

[20] W. Buchmuller and C. Greub, Nucl. Phys. B 363, 345 (1991); G. Cvetic, C. S. Kim and C. W. Kim, Phys. Rev. Lett. 82, 4761 (1999) [arXiv:hep-ph/9812525].

[21] F. M. L. Almeida, Y. A. Coutinho, J. A. Martins Simoes and M. A. B. do Vale, Phys. Rev. D 63, 075005 (2001) [arXiv:hep-ph/0008201]; J. Gluza and M. Zralek, Phys. Rev. D 55, 7030 (1997) [arXiv:hep-ph/9612227].

[22] J. Peressutti and O. A. Sampayo, Phys. Rev. D 67 (2003) 017302 [arXiv:hep-ph/0211355].

[23] J. A. M. Vermaseren, Comput. Phys. Commun. 83, 45 (1994).

[24] J. Peressutti, O. A. Sampayo and J. I. Aranda, Phys. Rev. D 64, 073007 (2001) [arXiv:hep-ph/0105162].

[25] R. Mertig, M. Bohm and A. Denner, Comput. Phys. Commun. 64, 345 (1991).

[26] E. Boos et al. [CompHEP Collaboration], Nucl. Instrum. Meth. A 534, 250 (2004) [arXiv:hep-ph/0403113]. A. Pukhov et al., arXiv:hep-ph/9908288.

[27] S. Kawabata, Comput. Phys. Commun. 88, 309 (1995).

[28] T. Sjostrand, P. Eden, C. Friberg, L. Lonnblad, G. Miu, S. Mrenna and E. Norrbin, Comput. Phys. Commun. 135, 238 (2001) [arXiv:hep-ph/0010017].

[29] A. S. Belyaev et al., arXiv:hep-ph/0101232.

[30] S. Eidelman et al. [Particle Data Group], Phys. Lett. B 592 (2004) 1.

[31] For most recent studies, see A. Pilaftsis and T. E. J. Underwood, [arXiv:hep-ph/0506107]; G. C. Branco, R. Gonzalez Felipe, F. R. Joaquim and B. M. Nobre, arXiv:hep-ph/0507092; S. Dar, Q. Shafi and A. Sil, arXiv:hep-ph/0508037; E. J. Chun, arXiv:hep-ph/0508050.