Heavy Majorana Neutrino Production at Future $ep$ Colliders

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Abstract

The heavy singlet Majorana neutrinos are introduced to generate the neutrino mass in the framework of canonical type-I seesaw mechanism. The phenomena induced by the heavy Majorana neutrinos are important to search for new physics. In this paper, we explore the heavy Majorana neutrino production and decay at future $e^-p$ colliders. The corresponding cross sections via $W$ and photon fusion are predicted for different collider energies. Combined with the results of the heavy Majorana neutrino production via single $W$ exchange, this work can provide helpful information to search for heavy Majorana neutrinos at future $e^-p$ colliders.

Keywords: type-I seesaw, Majorana neutrino, $e^-p$ collision

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1 Introduction

The neutrino oscillation experiments show that the neutrinos have minor non-zero masses, which are compelling for new physics beyond the standard model. The phenomena induced by the neutrino mass generating mechanism are important to search for new physics and attract more and more attention. To generate the neutrino mass, the natural way is to introduce the seesaw mechanism. Among them, one simple model is to introduce the heavy right-handed Majorana neutrinos $N_R$, which is known as the famous type-I seesaw mechanism [1–5]. The key point to test the type-I seesaw model is to search for the existence of $N_R$. In this kind of model, besides the Dirac neutrino mass term $\bar{\nu}_L M_D N_R$, which comes from the Yukawa interactions, the right-handed neutrino $N_R$ and its charged-conjugate counterpart $(N_R)^c$ can also form a Majorana neutrino mass term $(N_R)^c M_R N_R$, where the effective neutrino mass matrix can be given by the seesaw formula $M_\nu \approx -M_D M_R^{-1} M_D^T$, and the smallness of $M_\nu$ can be attributed to the large mass scale of $M_R$. Explicitly, the Majorana neutrino mass term $(N_R)^c M_R N_R$ violates the lepton-number by two units ($\Delta L = 2$), so we can probe the Majorana neutrino production signal through the lepton-number violating processes.

The search for Majorana neutrinos via the lepton-number violating process has been studied. At the low Majorana neutrino mass region, the experiments of the interesting processes, such as the well known neutrinoless double beta decays ($N(A,Z) \rightarrow N(A,Z+2) + 2e^-$) [6], the rare decays of the meson ($M_1^\pm \rightarrow M_2^\mp \ell^+_1 \ell^-_2$) [7] and tau decays (e.g. $\tau^\pm \rightarrow \ell^\mp M_1^\pm M_2^\pm$) [8] are explored to set the strong constraints on both the heavy Majorana neutrino mass and the related mixing parameters with the charged-leptons [9]. For the Majorana neutrino mass above the electroweak scale, its production has been investigated at various collider experiments (for a review, we refer to Ref. [10]). At hadron colliders, the most widely studied mode is the Drell-Yan process via a single virtual $W$ boson [11–13]. Due to the collinear logarithmic enhancement in $t$-channel exchange of massless gauge boson, the vector boson fusion process is important for the higher Majorana neutrino mass [14–16]. At higher collider energies, the gluon luminosity grows faster, so that the heavy Majorana neutrino production via the gluon fusion process becomes interesting [17,18]. At $e^+e^-$ colliders, the Majorana neutrino production can be studied via the processes, e.g. $t$-channel $W^*$ exchange, $s$-channel $Z^*$ exchange [19,20]. Complementary to search for the lepton-number violating processes at hadron colliders and $e^+e^-$ colliders, the $t$-channel $W^*$ exchange process at $e^-p$ colliders has been studied [21–23]. In this work, we focus on the heavy Majorana neutrino production and decay in the context of $W^*\gamma$ interaction at $e^-p$ colliders. As shown in Ref. [2], the general amplitude of $\Delta L = 2$ process is proportional to the Majorana neutrino mass, and the light Majorana neutrino contribution is strongly suppressed due to the small neutrino mass, we only consider the heavy Majorana neutrino. The mixing parameter between electron and the heavy Majorana neutrino is strictly constrained [24,25] and the tau lepton is hard to be reconstructed, therefore we purely concentrate on the di-muon production channel at future $e^-p$ colliders, e.g. LHeC [26], FCC-ep [27], ILC$\otimes$FCC [28].

This paper is organized as follows. A simple model is briefly introduced in Section 2 and the numerical results and discussions are obtained in Section 3. Finally, a short summary is given.
2 The model

Within the standard model, neutrinos are massless in the absence of right-handed neutrinos. However, recent neutrino oscillation experiments have clearly shown that neutrinos are massive. In order to explain the smallness of neutrino masses, many new physics models have been proposed. A simple extension of the standard model is to introduce three heavy right-handed neutrino singlets $N_R$ and the gauge-invariant Lagrangian relevant for the neutrino masses can be written as

\begin{equation}
-L_{\text{neutrino}} = \ell_L Y_\nu \tilde{H} N_R + \frac{1}{2} (N_R)^c M_R N_R + \text{h.c. ,} \tag{1}
\end{equation}

where $\ell_L$ and $\tilde{H} \equiv i \sigma_2 H^*$ respectively denote the left-handed lepton doublet and Higgs doublet, and $N_R$ the right-handed neutrino singlet. $Y_\nu$ is the $3 \times 3$ neutrino Yukawa coupling matrix and $M_R$ the symmetric right-handed Majorana neutrino mass matrix. After the spontaneous gauge symmetry breaking, the neutrino mass terms appear as

\begin{equation}
-L_{\text{mass}} = \overline{\nu}_L M_D N_R + \frac{1}{2} (N_R)^c M_R N_R + \text{h.c.}
= \frac{1}{2} (\nu_L)^c \begin{pmatrix} 0 & M_D & (\nu_L)^c \\ M_D^T & M_R & \end{pmatrix} \begin{pmatrix} N_R \end{pmatrix} + \text{h.c. .} \tag{2}
\end{equation}

Here $(\nu_L)^c$ and $(N_R)^c$ are respectively defined as $(\nu_L)^c = CV_L^T$ and $(N_R)^c = CN_R^T$ with $C$ being the charge-conjugation operator. $M_D = Y_\nu \langle H \rangle$ is the Dirac neutrino mass matrix with $\langle H \rangle \approx 174$ GeV being the Higgs vacuum expectation value. Since the right-handed neutrinos $N_R$ are $SU(2)_L$ gauge singlets and thus the Majorana neutrino mass term $(N_R)^c M_R N_R$ is not subject to the gauge symmetry breaking scale, the absolute scale of the right-handed Majorana neutrino mass matrix $M_R$ can naturally be much higher, $M_R \gg \langle H \rangle$. In the second line of Eq. (2), we have used the relationship $(N_R)^c M_D^T (\nu_L)^c = \overline{\nu}_L M_D N_R$.

The overall mass matrix in Eq. (2) is symmetric and can be diagonalized by one unitary transformation

\begin{equation}
\begin{pmatrix} V & R \\ S & U \end{pmatrix}^\dagger \begin{pmatrix} 0 & M_D \\ M_D^T & M_R \end{pmatrix} \begin{pmatrix} V & R \\ S & U \end{pmatrix}^* = \begin{pmatrix} \tilde{M}_\nu & 0 \\ 0 & \tilde{M}_N \end{pmatrix}, \tag{3}
\end{equation}

where $\tilde{M}_\nu = \text{Diag}\{m_1, m_2, m_3\}$ and $\tilde{M}_N = \text{Diag}\{M_1, M_2, M_3\}$ denote the mass eigenvalues of light and heavy Majorana neutrinos, respectively. In the limit of $M_D \ll M_R$, the effective neutrino mass matrix can be of the order of $M_\nu \approx -M_D M_R^{-1} M_D^T$, and the smallness of $M_\nu$ can be attributed to the largeness of $M_R$. According to the unitary condition of the transformation matrix in Eq. (3), the matrices $V, R, S, U$ can satisfy

\begin{equation}
VV^\dagger + RR^\dagger = SS^\dagger + UU^\dagger = 1 , \\
V^\dagger V + S^\dagger S = R^\dagger R + U^\dagger U = 1 , \tag{4}
\end{equation}

with $VV^\dagger \sim UU^\dagger \sim O(1)$ and $RR^\dagger \sim S^\dagger S \sim O(M_D^2/M_R^2)$. 


Moreover, the relation between the neutrino flavor eigenstates $\nu_\alpha$ (for $\alpha = e, \mu, \tau$) and the mass eigenstates $\nu_i$ (for $i = 1, 2, 3$) can be given by

$$
\begin{pmatrix}
\nu_e \\
\nu_\mu \\
\nu_\tau
\end{pmatrix}_L = V \begin{pmatrix}
\nu_1 \\
\nu_2 \\
\nu_3
\end{pmatrix}_L + R \begin{pmatrix}
N_1 \\
N_2 \\
N_3
\end{pmatrix}_L.
$$

Therefore the standard weak charged-current interaction Lagrangian of leptons in terms of the mass eigenstates can be written as

$$
-L_{cc} = \frac{g}{\sqrt{2}} \left[ (e, \mu, \tau)_L V^\gamma \mu \begin{pmatrix}
\nu_1 \\
\nu_2 \\
\nu_3
\end{pmatrix}_L W^\mp + (e, \mu, \tau)_L R^\gamma \mu \begin{pmatrix}
N_1 \\
N_2 \\
N_3
\end{pmatrix}_L W^\mp \right] + \text{h.c}.
$$

It is worth mentioning that we have already chosen the basis where the flavor eigenstates of three charged-leptons are identified with their mass eigenstates.

The matrix $V$ in Eq. (6) is the so-called Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix [29,30], denotes the mixing between charged-leptons and light Majorana neutrinos and can be measured from the oscillation experiments. The standard parametrization of the matrix $V$ can be expressed as [31]

$$
V = \begin{pmatrix} c_{13}c_{12} & c_{13}s_{12} & s_{13}e^{-i\delta} \\
-s_{12}c_{23} - c_{12}s_{13}s_{23}e^{i\delta} & c_{12}c_{23} - s_{12}s_{13}s_{23}e^{i\delta} & c_{13}s_{23} \\
+s_{12}s_{23} - c_{12}s_{13}c_{23}e^{i\delta} & -c_{12}s_{23} - s_{12}s_{13}c_{23}e^{i\delta} & c_{13}c_{23} \end{pmatrix} \times \text{diag}\{e^{i\rho}, e^{i\sigma}, 1\},
$$

where $c_{ij} \equiv \cos\theta_{ij}$, $s_{ij} \equiv \sin\theta_{ij}$, $\delta$ is the Dirac CP violation phase and $\rho$, $\sigma$ are two additional Majorana CP violation phases. The matrix $R$ indicates the mixing between charged-leptons and heavy Majorana neutrinos, which can be determined from the $0\nu\beta\beta$-decay experiments or possible collider experiments.

In our previous work [32], we proposed a simple scenario to explicitly break the $S_{3L} \times S_{3R}$ flavor symmetry in the framework of the canonical seesaw model, where we can naturally explain the realistic lepton mass spectra and flavor mixing pattern as well as the cosmological matter-antimatter asymmetry via resonant leptogenesis. In that case, however, the mass scale of $M_R$ is too high to be detected experimentally. In order to probe the heavy Majorana neutrino production signal, the scale of $M_R$ must be at the order of $\mathcal{O}(1)$ TeV or smaller and the strength of matrix $R$, which scales as $\mathcal{O}(M_D/M_R)$, must be large enough. In Ref. [33], a significant “structural cancellation” has been imposed on the textures of $M_D$ and $M_R$ to get $M_\ell \sim \mathcal{O}(10^{-2})$ eV with $M_D \sim \mathcal{O}(10^3)$ GeV and $M_R \sim \mathcal{O}(10^3)$ GeV. Both of the two necessary conditions $M_R \sim \mathcal{O}(10^3)$ GeV and $R \sim \mathcal{O}(M_D/M_R) \sim \mathcal{O}(0.1)$ can be satisfied spontaneously.

At present, the constraint on the mixing between heavy Majorana neutrinos and electrons can be derived from the $0\nu\beta\beta$-decay experiments [24,25]

$$
\sum_N \frac{|R_{\ell N}|^2}{m_N} < 5 \times 10^{-5} \text{ TeV}^{-1}, \quad \text{for } m_N \gg 1 \text{ GeV}.
$$

For the mixing between heavy Majorana neutrinos and muons, the most stringent bound comes from the LHC experiments [34,35]

$$
|R_{\mu N}|^2 < 2 \times 10^{-3} - 0.5 \text{ (at 95\% C.L.)}, \quad \text{for } m_N = 100 - 500 \text{ GeV}.
$$
From the lepton universality tests [36], the mixing between heavy Majorana neutrinos and charged-leptons can be restricted to

\[ |R_{eN}|^2 < 5.9 \times 10^{-3}, |R_{\mu N}|^2 < 2.5 \times 10^{-3}, |R_{\tau N}|^2 < 5.9 \times 10^{-3} \text{(at 90% C.L.)}, \]

for \( m_N > m_W \). (10)

As shown in Ref. [9,37], the total decay width of the heavy Majorana neutrino can be expressed approximately as

\[
\Gamma_N \simeq \begin{cases} 
\left( \frac{3G_F^2 m_N^5}{32\pi^3} \right) \sum_{\alpha=e,\mu,\tau} |R_{\alpha N}|^2, & m_N < m_W, \\
\left( \frac{3G_F m_N^3}{8\pi\sqrt{2}} \right) \sum_{\alpha=e,\mu,\tau} |R_{\alpha N}|^2, & m_N > m_W.
\end{cases} \tag{11}
\]

3 Heavy Majorana neutrino phenomena at ep colliders

We start by considering the process

\[ e^- + p \rightarrow e^- + \ell^\pm_\alpha + N + X \rightarrow e^- + \ell^\pm_\alpha + \ell^\pm_\beta (\ell^\mp_\beta) + X. \tag{12}\]

The relevant process at the parton level (Fig. 1) is

\[ \gamma(p_1) + q(p_2) \rightarrow q'(p_3) + \ell^\pm_\alpha (p_4) + \ell^\pm_\beta (\ell^\mp_\beta)(p_5) + q_1(p_6) + q_2(p_7), \tag{13}\]

where \( p_i \) (for \( i = 1, \cdots, 7 \)) is the four-momentum of the corresponding particle. The photon is emitted from the electron and can be described by the photon density function [38]

\[
f_{\gamma/e^-}(x) = \frac{\alpha}{2\pi} \left[ 1 + \frac{(1-x)^2}{x} \ln \frac{Q_{\text{max}}^2}{Q_{\text{min}}^2} + 2m^2_e x \left( \frac{1}{Q_{\text{max}}^2} - \frac{1}{Q_{\text{min}}^2} \right) \right]. \tag{14}\]

Here \( x = E_\gamma/E_e \) with \( E_\gamma \) and \( E_e \) the energies of the photon and electron, respectively. \( \alpha \) is the fine structure constant and \( m_e \) the mass of electron. \( Q_{\text{min}}^2 = m_e^2 x^2/(1-x) \) and \( Q_{\text{max}}^2 = (\theta_c E_e)^2(1-x) + Q_{\text{min}}^2 \) with \( \theta_c \) the cut of the electron scattering angle.

Figure 1: Feynman diagrams at the parton level for the process \( \gamma q \rightarrow \ell^\pm_\alpha \ell^\pm_\beta + X \).

The cross section for the process in Eq. (12) can be written as

\[
\sigma(e^- p \rightarrow \ell^\pm_\alpha \ell^\pm_\beta + X) = \sum_q \int dx_1 dx_2 f_{\gamma/e^-}(x_1)f_{q/p}(x_2, \mu^2) \cdot \tilde{\sigma}(\gamma q \rightarrow \ell^\pm_\alpha \ell^\pm_\beta + X), \tag{15}\]

5
where \( f_{q/p}(x_2, \mu^2) \) is the parton distribution function with \( x_2 \) the energy fraction of \( q \), and \( \mu \) the factorization scale. Here, we employ the CT14QED \cite{39} for the photon distribution function and parton distribution functions in proton. \( \sigma \) is the partonic cross section

\[
\hat{\sigma}(\gamma q \rightarrow \ell_\alpha \ell_\beta + X) = \frac{1}{2\hat{s}} \int |\mathcal{M}|^2 d\mathcal{L}_{ips}^5 .
\]  

Here \( \hat{s} = x_1 x_2 s \) is the flux factor with \( \sqrt{s} \) the electron-proton center-of-mass energy. \( d\mathcal{L}_{ips}^5 \) represents the five-body Lorentz invariant phase space of the final particles, and \( |\mathcal{M}|^2 \) is the squared scattering amplitude averaged (summed) over the initial (final) particles.

For the convenience of the discussions on the numerical results, all the input parameters used in our numerical analysis are listed as follows,

\[
\alpha = 1/137, \ m_e = 0.51 \text{ MeV}, \ \theta_c = 32 \text{ mrad}, \ \mu = m_N ,
\]

\[
m_W = 80.385 \text{ GeV}, \ \Gamma_W = 2.085 \text{ GeV}, \ G_F = 1.166 \times 10^{-5} \text{ GeV}^{-2}, \ \sin^2 \theta_W = 0.231 ,
\]

\[
|R_{eN}|^2 = 5.0 \times 10^{-6}, \ |R_{\mu N}|^2 = 2.0 \times 10^{-3}, \ |R_{\tau N}|^2 = 5.9 \times 10^{-3} .
\]  

\[
(17)
\]

In this work, we only consider the contribution of a single heavy Majorana neutrino and concentrate on the di-muon production channel. We obtain the cross section for the process in Eq. (12) at LHeC with \( \sqrt{s} = 1.3 \text{ TeV} \), FCC-ep with \( \sqrt{s} = 3.5 \text{ TeV} \) and ILC \( \otimes \) FCC with \( \sqrt{s} = 10 \text{ TeV} \). The cross sections in final states with like-sign dileptons are shown in Fig. 2 as a function of the heavy Majorana neutrino mass. The difference between \( \sigma(e^- p \rightarrow e^- \mu^- \mu^- + X) \) and \( \sigma(e^- p \rightarrow e^- \mu^- \mu^- + X) \) can be attributed to the role of parton distribution function \( f_{q/p}(x, \mu^2) \) and induce the charge asymmetry. To investigate this charge asymmetry, we define

\[
\delta = \frac{\sigma(e^- p \rightarrow e^- \mu^- \mu^- + X) - \sigma(e^- p \rightarrow e^- \mu^- \mu^- + X)}{\sigma(e^- p \rightarrow e^- \mu^- \mu^- + X) + \sigma(e^- p \rightarrow e^- \mu^- \mu^- + X)} .
\]  

\[
(18)
\]

The numerical results of \( \delta \) as a function of \( m_N \) are listed in Table 1.

\[\text{Figure 2: The cross section for (a) } e^- p \rightarrow e^- \mu^- \mu^- + X, \ (b) e^- p \rightarrow e^- \mu^- \mu^- + X \text{ as a function of } m_N.\]
We also investigate the process \( e^- p \to e^- \mu^- \mu^+ + X \) of the unlike-sign dileptons, the corresponding cross sections as a function of \( m_N \) are displayed in Fig. 3. For this process, the standard model process for the \( \mu^- \mu^+ \) production via a \( Z^0 \) or a virtual photon is the dominant background and can be greatly reduced by the constraint for the invariant mass of the \( \mu^- \mu^+ \) pair.

Similar as Eq. (12), we study the process \( e^- p \to \nu_e + \ell^- + N + X \to \nu_e + \ell^- + \ell^\pm + X \), its cross section can be written as

\[
\sigma(e^- p \to \ell^- + N + X) = \int dx_1 f_{\gamma/p}(x_1, \mu^2) \cdot \hat{\sigma}(\gamma e \to \ell^- + N + X), \tag{19}
\]

where the photon is emitted from the proton and can be described by the photon distribution function \( f_{\gamma/p}(x, \mu^2) \). The results of the corresponding cross sections related to \( \mu^- \mu^\pm \) channels are shown in Fig. 3.

In the following, taking the process Eq. (12) as an example, we investigate the reconstructed invariant mass distributions and transverse momentum distributions of the final state like-sign dileptons and jets for the process \( e^- p \to e^- \mu^- \mu^+ + X \). The invariant mass of the heavy Majorana neutrino can be reconstructed from the four-momenta of the final state charge-leptons and jets. Since the final state charged-leptons \( (\ell^\pm_\alpha \text{ and } \ell^\pm_\beta) \) are indistinguishable, we define the normalized differential distribution \( 1/\sigma d\sigma/dM_{\ell\bar{\ell}jj} = 1/\sigma(d\sigma/dM_{\ell\bar{\ell}jj} + d\sigma/dM_{\ell\ell jj})/2 \) for the reconstructed invariant mass. In Fig. 4, the results of the normalized invariant mass distributions \( 1/\sigma d\sigma/dM_{\ell\ell jj} \) for various heavy Majorana neutrino mass are displayed. The peak positions imply the mass of heavy Majorana neutrino and can be reconstructed effectively. We also calculate the normalized distributions for the invariant mass of the lepton pair in Fig. 6. When \( m_N \) is much lower than \( m_W \), e.g. \( m_N = 20 \text{ GeV} \), a peak of \( m_{\ell\ell} \) appears due to the resonant production of \( W \) boson. Analogously,
Figure 4: The cross section for (a) $e^- p \rightarrow \nu_e \mu^- \mu^- + X$, (b) $e^- p \rightarrow \nu_e \mu^- \mu^+ + X$ as a function of $m_N$.

we define the normalized differential distribution $1/\sigma \frac{d\sigma}{dp_{T,\ell,j}^{\alpha,j1}} = 1/\sigma (d\sigma/dp_{T,\ell,j}^{\alpha,j1} + d\sigma/dp_{T,\ell,j}^{\beta,j2})/2$ for the transverse momentum of the final state charged-leptons and jets, and display the normalized transverse momentum distributions $1/\sigma \frac{d\sigma}{dp_{T,\ell,j}^{\alpha,j1}}$ for $m_N = 60$ GeV (Fig. 7).

Figure 5: The normalized invariant mass distributions $1/\sigma \frac{d\sigma}{dM_{\ell\ell}}$ for $m_N = 20, 40, 60, 80$ GeV.

Figure 6: The normalized invariant mass distributions $1/\sigma \frac{d\sigma}{dM_{\ell\ell}}$ of (a) like-sign dileptons, (b) unlike-sign dileptons for $m_N = 20, 80$ GeV.
Figure 7: The normalized transverse momentum distributions (a) $1/\sigma d\sigma/dp_T^\ell$ and (b) $1/\sigma d\sigma/dp_T^j$ for $m_N = 60$ GeV at 1.3 TeV LHeC (solid), 3.5 TeV FCC-ep (dash), 10 TeV ILC$\otimes$FCC (dot).

4 Summary

The heavy Majorana neutrinos $N_R$ are introduced to explain the minor neutrino mass in the framework of type-I seesaw mechanism. Due to the existence of the Majorana neutrino mass term, we can search for the Majorana neutrinos via the lepton-number violating processes, which have been studied in various experiments. In this work, we explore the heavy Majorana neutrino production and decay in the context of $W^*\gamma$ interaction and investigate the related dilepton production process at future $e^-p$ colliders. The cross sections for the processes $e^-p \rightarrow e^-\mu^+\mu^\pm + X$ and $e^-p \rightarrow \nu_e\mu^-\mu^\pm + X$ at future LHeC, FCC-ep and ILC$\otimes$FCC are predicted. We further investigate the process $e^-p \rightarrow e^-\mu^+\mu^\pm + X$ in detail, and obtain several differential distributions. Combined with the results of the heavy Majorana neutrino production via single $W$ exchange, this work can provide helpful information to search for the heavy Majorana neutrinos at future $e^-p$ colliders.

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