Lepton Number Violating Processes Mediated by Majorana Neutrinos at Hadron Colliders

Sergey Kovalenko, Zhun Lu, and Ivan Schmidt

Departamento de Física,
Universidad Técnica Federico Santa María,
Casilla 110-V, Valparaíso, Chile

and Center of Subatomic Physics,
Valparaíso, Chile

Abstract

We study the Lepton number violating like-sign dilepton processes $h_1h_2 \rightarrow l^\pm l'^\pm jjX$ and $h_1h_2 \rightarrow l^\pm l'^\pm W^\pm X$, mediated by heavy GeV scale Majorana neutrinos. We focus on the resonantly enhanced contributions with a nearly on-mass-shell Majorana neutrino in the s-channel. We study the constraints on like-sign dilepton production at the Tevatron and the LHC on the basis of the existing experimental limits on the masses of heavy neutrinos and their mixings $U_{\alpha N}$ with $\alpha = \nu_e, \nu_\mu, \nu_\tau$. Special attention is paid to the constraints from neutrinoless double beta decay. We note that searches for like-sign $e^\pm e^\pm$ events at Tevatron and LHC may shed light on CP-violation in neutrino sector. We also discuss the conditions under which it is possible to extract individual constraints on the mixing matrix elements in a model independent way.
I. INTRODUCTION

At present there are no doubts that neutrinos are massive particles mixing with each other. Moreover, according to the neutrino oscillation experiments, their masses are extremely small while mixing is nearly maximal. In this respect neutrinos drastically defer from all other known particles, and this difference represents one of the pressing problems for theory. The famous see-saw mechanism showed the way to possible solutions of this problem via introduction of very heavy Majorana particles mixed with ordinary neutrinos, bringing in the observable mass hierarchy between the light neutrinos and other fermions. These heavy particles could be heavy Majorana neutrinos, as in the original formulation of see-saw, or other new particles such as neutralinos, heavy Majorana particles mixing with neutrinos in the framework of SUSY models with R-parity violation. The new heavy particles may lead to observable effects beyond the light neutrino sector, since they could manifest themselves indirectly via their virtual contribution to processes involving ordinary particles. If the masses of some of these new particles are not extremely large, lying within the reach of the running of forthcoming experiments, they can be searched for directly among the products of colliding particles.

Here we consider an extended see-saw scenario, including $n$ species of SM singlet right-handed neutrinos $\nu'_{Rj} = (\nu'_{R1}, \ldots, \nu'_{Rn})$, besides the three left-handed weak doublet neutrinos $\nu'_L = (\nu'_L e, \nu'_L \mu, \nu'_L \tau)$. The general mass term for this set of fields can be written as

$$-rac{1}{2} \overline{\nu'} L \mathcal{M}^{(\nu)} \nu' c + h.c. = -\frac{1}{2} (\overline{\nu'}_L, \nu'_R) \begin{pmatrix} \mathcal{M}_L & \mathcal{M}_D \\ \mathcal{M}_D^T & \mathcal{M}_R \end{pmatrix} \begin{pmatrix} \nu'_L c \\ \nu'_R \end{pmatrix} + h.c. \quad (1)$$

$$= -\frac{1}{2} \left( \sum_{i=1}^{3} m_{\nu_i} \overline{\nu'_i} \nu_i + \sum_{j=1}^{n} m_{\nu'_j} \overline{\nu'_j} \nu'_j \right) + h.c. \quad (2)$$

Here $\mathcal{M}_L, \mathcal{M}_R$ are $3 \times 3$ and $n \times n$ symmetric Majorana mass matrices, and $\mathcal{M}_D$ is a $3 \times n$ Dirac type matrix. Rotating the neutrino mass matrix to the diagonal form by a unitary transformation

$$U^T \mathcal{M}^{(\nu)} U = \text{Diag}\{m_{\nu_1}, \ldots, m_{\nu_{3+n}}\} \quad (3)$$

we end up with $3 + n$ Majorana neutrinos with masses $m_{\nu_1}, \ldots, m_{\nu_{3+n}}$. In this scenario the light neutrinos have the mass scale $\mathcal{M}^2_D/\mathcal{M}_R$, and as a result there should be also heavy Majorana neutrinos ($N$) with mass scale $\mathcal{M}_R$. Due to the fact that Majorana neutrinos are
also their anti-neutrinos, they can produce lepton number \((L)\) violation by two units. One of the decisive processes for probing the Majorana nature of neutrinos is neutrinoless double beta \((0\nu\beta\beta)\) decay which has been extensively studied in the literature \([1, 2]\). Majorana neutrinos may also resonantly contribute to meson \([3]\) and \(\tau\) decays \([4]\).

Another potential process to look for Majorana neutrinos is like-sign dilepton production \(h_1h_2 \rightarrow l^\pm l'^\pm W^\mp X\) at hadron colliders \([5, 6, 7, 8]\). In this paper we give a detailed analysis of like-sign dilepton production processes, evaluating the cross section of \(h_1h_2 \rightarrow l^\pm l'^\pm jjX\) and \(h_1h_2 \rightarrow l^\pm l'^\pm W^\mp X\).

The existing very stringent constraint from \(0\nu\beta\beta\) experiments on the inverse effective mass \(U_{eN}^2/M_N\) of a heavy Majorana neutrino \(N\) is frequently treated as the limit, which makes unrealistic the experimental observation of the above processes with \(e^\pm e^\pm\) pair in the final state. However, we note that CP-violating Majorana phases, if present in the neutrino mixing matrix elements \(U_{e\alpha}\), are able to significantly soften or even completely evade the \(0\nu\beta\beta\)-constraints on the \(e^\pm e^\pm\)-process production. We point out that this can be true even in the presence of only one heavy Majorana neutrino, due to the interference of heavy-light neutrino contribution to \(0\nu\beta\beta\)-decay.

The paper is organized as follows. In the next section we present our approach to calculation of the cross sections of the \(h_1h_2 \rightarrow l^\pm l'^\pm jjX\) and \(h_1h_2 \rightarrow l^\pm l'^\pm W^\mp X\) processes. In Sec. III we calculate the decay width of the heavy Majorana neutrino entering in the above cross sections. Sec. IV is devoted to a discussion of the existing limits on the parameters of heavy Majorana neutrino and their impact on the prospects for searches for like-sign dileptons at Tevatron and LHC. Here we comment on implications of the \(0\nu\beta\beta\)-decay constraints and CP-violation in the neutrino sector and discuss the possibility of extracting the heavy-light neutrino mixing matrix elements \(U_{eN}, U_{\mu N}, U_{\tau N}\) in a model independent way.

**II. THEORETICAL FRAMEWORK**

The processes we study are like-sign dilepton inclusive production in high energy \(pp\) or \(p\bar{p}\) collisions

\[
h_1(P_1) + h_2(P_2) \rightarrow l^\pm(l_1) + l'^\pm(l_2) + jj(W^\mp) + X, \tag{4}
\]
FIG. 1: (a) resonant and (b) non-resonant Majorana neutrino contributions to $h_1(P_1) + h_2(P_2) \rightarrow l^+(l_1) + l'^+(l_2) + X$ process. In the diagrams we indicate the momenta of the corresponding particles.

where $jj$ and $X$ denotes two quark jets and undetected hadronic states. The two leptons $l$ and $l'$ can have the same or different lepton flavors. In these processes the total lepton number $L$ is violated in two units $\Delta L = 2$. If one assumes the existence of a Majorana neutrino $N$, the process can be realized through the diagrams shown in Fig. 1. We focus on the resonantly enhanced diagram Fig. 1(a) with the nearly on mass-shell neutrino in the kinematical region of the studied processes.

In the diagram Fig. 1(a) there is a $W$ boson in the final state, which decays to hadron or lepton states. Here we will consider the situation in which $W^-$ decays to two quark jets, as shown in Fig. 1(b). According to the factorization theorem, we can express the cross section of the process as

$$\sigma_{h_1+h_2-l^++l'^+} = \sum_{k,l} \int dx_1 \int dx_2 \{q_{a/h_1}(x_a) \bar{q}_{b/h_2}(x_b) \times \sigma_{q_a\bar{q}_b-l^++q_k} + (x_a \leftrightarrow x_b) \}, \quad (5)$$

here $q_{a/h_1}(x_a)$ and $\bar{q}_{b/h_2}(x_b)$ denote the densities of partons inside each hadron, and $d\sigma_{q_a\bar{q}_b-l^++q_k}$ is the differential cross section of the partonic subprocess

$$q_a(k_a) + \bar{q}_b(k_b) \rightarrow (l_1) + l' + W(P_W) \rightarrow l(l_1) + l' + q_k(q_1) + \bar{q}_l(q_2) \quad \text{(6)}$$

which has the form

$$d\sigma_{q_a\bar{q}_b-l^++q_k} = \frac{1}{2^8} |M_{q_a\bar{q}_b-l^++q_k}|^2 (2\pi)^4 \delta^4(k_a + k_b - l_1 - l_2 - q_1 - q_2)$$

$$\times \frac{d^3l_1}{(2\pi)^3 2E_1} \frac{d^3l_2}{(2\pi)^3 2E_2} \frac{d^3q_1}{(2\pi)^3 2E_{q_1}} \frac{d^3q_2}{(2\pi)^3 2E_{q_2}}. \quad (7)$$
The next step is to calculate the amplitude squared $|\mathcal{M}|^2$, which can be written as

$$
|\mathcal{M}|^2 = 16G_F^4 M_W^6 \times \frac{(|V_{ud}|^2 u(x_a)\bar{d}(x_b) + |V_{us}|^2 u(x_a)\bar{s}(x_b)) H^{\mu\nu,\alpha\beta} L_{\mu\nu,\alpha\beta}}{((s-M_W^2)^2 + \Gamma_W^2 M_W^2)((P_W^2 - M_W^2)^2 + \Gamma_W^2 M_W^2)} + (x_a \leftrightarrow x_b),
$$

with the tensors having the following forms:

$$
H^{\mu\nu,\alpha\beta} = 16(k^\mu k^\nu + k^\mu k^\nu - g^{\mu\nu}k_a \cdot k_b)(q_1^\alpha q_2^\beta + q_1^\beta q_2^\alpha - g^{\mu\nu}q_1 \cdot q_2) \tag{9}
$$

$$
L^{\mu\nu,\alpha\beta} = \left[ \sum_N \frac{m_N U_{IN} U_{FN}^*}{(k-l_1)^2 - m_N^2} \bar{v}(l_1)\gamma^\mu\gamma^\nu P_R v^c (l_2) \right] \left[ \sum_N \frac{m_N U_{IN} U_{FN}^*}{(k-l_1)^2 - m_N^2} \bar{v}(l_1)\gamma^\alpha\gamma^\beta P_R v^c (l_2) \right]^\dagger \tag{10}
$$

The contraction of the above two tensors yields

$$
H^{\mu\nu,\alpha\beta} L_{\mu\nu,\alpha\beta} = 128 \sum_N \frac{m_N^2 U_{IN} U_{FN}^*}{(k-l_1)^2 - m_N^2} \left[(l_1 \cdot k_b)(l_2 \cdot q_2)(k_a \cdot q_1) + (l_1 \cdot k_b)(l_2 \cdot q_1)(k_a \cdot q_2) + (l_1 \cdot k_a)(l_2 \cdot q_2)(k_b \cdot q_1) + (l_1 \cdot k_a)(l_2 \cdot q_1)(k_b \cdot q_2)\right] + \text{interference term} \tag{11}
$$

We have checked that the interference term gives a very small contribution compared to the first term in (11), and therefore we ignore it in the following calculations. The 4-body phase space in Eq. (7) can be rewritten as

$$
\Phi_4 = \delta^4(k_a + k_b - l_1 - l_2 - P_W) \frac{d^3 l_1}{(2\pi)^3 2E_{l_1}} \frac{d^3 l_2}{(2\pi)^3 2E_{l_2}} \frac{d^3 P_W}{(2\pi)^3 2E_{P_W}} \times \delta^4(P_W - q_1 - q_2) \frac{d^3 q_1}{(2\pi)^3 2E_{q_1}} \frac{d^3 q_2}{(2\pi)^3 2E_{q_2}} \frac{\lambda^{1/2}(P_W^2, m_{q_1}^2, m_{q_2}^2)}{8(2\pi)^6 P_W^2} d\Omega, \tag{12}
$$

where $P_W = q_1 + q_2$ denotes the momentum of the final virtual $W$ boson. In the c.m. frame of $q_1$ and $q_2$, the two body phase space can be expressed as

$$
\delta^4(P_W - q_1 - q_2) \frac{d^3 q_1}{(2\pi)^3 2E_{q_1}} \frac{d^3 q_2}{(2\pi)^3 2E_{q_2}} \lambda^{1/2}(P_W^2, m_{q_1}^2, m_{q_2}^2) \sim P_W^2. \tag{13}
$$

where $d\Omega = d\cos \theta_1 d\phi_1$ is the solid angle of $q_1$, and the form of $\lambda(P_W^2, m_{q_1}^2, m_{q_2}^2)$ is given in Eq. (A5). In the case $P_W^2 \gg m_{q_1}^2$ and $m_{q_2}^2$, $\lambda^{1/2}(P_W^2, m_{q_1}^2, m_{q_2}^2) \sim P_W^2$. After integrating over $d\Omega$, Eq. (11) turns to

$$
\int d\Omega H^{\mu\nu,\alpha\beta} L_{\mu\nu,\alpha\beta} = \frac{128\pi}{3} \sum_N \frac{m_N^2 U_{IN} U_{FN}^*}{(k-l_1)^2 - m_N^2} P_W^2 \left[(l_1 \cdot k_b)(2E_{l_2} E_{k_a} + l_2 \cdot k_a) + (l_1 \cdot k_a)(2E_{l_2} E_{k_b} + l_2 \cdot k_b)\right], \tag{14}
$$

and the four-body phase space in Eq. (12) is reduced to a three-body phase space.
As shown in Fig. 2, it is convenient to use the following Lorentz invariants in order to express the kinematical variables:

\[
\hat{s} = (k_a + k_b)^2 = x_a x_b s, \quad s_a = (l_1 + P_W)^2, \\
s_b = (l_2 + P_W)^2, \quad t_a = (k_a - l_2)^2, \quad t_b = (k_b - l_1)^2. \tag{15}
\]

With the above invariant variables, we can apply the phase space transformation given in Ref. [9] to rewrite the cross section given in Eq. (5) as

\[
\sigma = a \int dx_a \int dx_b \int ds_a \int ds_b \int dt_a \int dt_b \int dP_W^2 \\
\times \frac{\Theta(-\Delta_4)}{s^2 \sqrt{-\Delta_4}} \mathcal{M}(x_a, x_b, s_a, s_b, t_a, t_b, P_W^2), \tag{16}
\]

where \(a = 2G_F^4 M_W^8 / (3(2\pi)^6)\), and

\[
\mathcal{M}(x_a, x_b, s_a, s_b, t_a, t_b) = \frac{|V_{ud}|^2 u(x_a) \bar{d}(x_b) + |V_{us}|^2 u(x_a) \bar{s}(x_b) + (x_a \leftrightarrow x_b)}{((\hat{s} - M_W^2)^2 + \Gamma_W^2 M_W^2)((P_W^2 - M_W^2)^2 + \Gamma_W^2 M_W^2)} \\
\times \sum_N \frac{m_N U_{In} U_{lN}^*}{s_b - m_N^2} \mathcal{M}_1(\hat{s}, s_a, s_b, t_a, t_b, P_W^2) + \text{interference term}, \tag{17}
\]

where

\[
\mathcal{M}_1(\hat{s}, s_a, s_b, t_a, t_b, P_W^2) = (\hat{s} - s_b + t_a - m_a^2)((s_b - P_W^2 - m_{l_2}^2)(s_a + t_a - t_b - m_{l_1}^2) \\
+ P_W^2(s_a + t_b - m_{l_2}^2) - (t_a - m_a^2 - m_{l_1}^2)((s_b - P_W^2 - m_{l_1}^2) \\
\times (s_b - t_a + t_b - m_{l_2}^2) - P_W^2(t_b - m_{l_2}^2 - m_{l_1}^2)), \tag{18}
\]

The phase space can be constrained to the physical region by the following inequalities for the Gram determinants

\[
\Delta_3(p_1, p_2, p_3) \geq 0, \quad \Delta_4(p_1, p_2, p_3, p_4) \leq 0, \tag{19}
\]

FIG. 2: Definition of the invariant variables.
where \{p_i\} denotes any subset of momenta \{k_a, k_b, l_1, l_2, P_W\}. The resulting integration limits \int_{s_a}^{s_a} ds_a, \int_{t_a}^{t_a} dt_a, \int_{t_b}^{t_b} dt_b are given in Appendix.

If there exists one Majorana neutrino with the mass \(m_N\) in the region \(s_b < m_N^2 < s_b^+\), the integrand shown in Eq. (16) has a pole at \(s_b = m_N^2\). Introducing a decay width \(\Gamma_N\) for \(N\) through the substitution \(m_N \rightarrow m_N - (i/2)\Gamma_N\), the Majorana neutrino propagator can be written as \(1/(s - m_N^2 + i\Gamma_N\). As we will show in the next section, \(\Gamma_N \ll m_N\) for \(m_N\) in range we are studying in the present paper. Therefore, the cross section can be written as

\[
\sigma^{h_1h_2 \rightarrow l^+l^+X} (m_N) \approx a \pi \int dx_a \int dx_b \int_{s_a}^{s_a} ds_a \int_{t_a}^{t_a} dt_a \int_{t_b}^{t_b} dt_b \frac{\Theta(-\Delta_4)}{s^2 \sqrt{-\Delta_4}}
\]

\[
\times \frac{1}{\Gamma_N} \frac{f(s_b)}{(s_b - m_N^2)^2 + m_N^2 \Gamma_N^2} = \pi \frac{f(m_N)}{m_N \Gamma_N}, \quad s_b < m_N^2 < s_b^+ \quad \text{and} \quad \Gamma_N << m_N. \quad (21)
\]

Therefore, with \(\Gamma_N\) in the denominator, the cross section can be strongly enhanced through the resonant production of a Majorana neutrino. If there are several Majorana neutrinos in the range \(s_b < m_N^2 < s_b^+\), there should be a sum over \(N_i\) in Eq. (20).

In the case \(m_N > M_W\), one can detect the \(W\) boson together with the lepton pair \(l^+l^+\) in the final state. Therefore we also give the formula to calculate the cross section of the process \(h_1 + h_2 \rightarrow l^+ + l^+ + W^- + X\),

\[
\sigma^{h_1h_2 \rightarrow l^+l^+W^-X} (m_N) \approx a' \pi \int dx_a \int dx_b \int_{s_a}^{s_a} ds_a \int_{t_a}^{t_a} dt_a \int_{t_b}^{t_b} dt_b \frac{\Theta(-\Delta_4)}{s^2 \sqrt{-\Delta_4}}
\]

\[
\times \frac{1}{\Gamma_N} \frac{f(s_b)}{(s_b - m_N^2)^2 + m_N^2 \Gamma_N^2} M_1(\hat{s}, s_a, m_N^2, t_a, t_b, P_W^2), \quad (22)
\]

where

\[
M_1(\hat{s}, s_a, s_b, t_a, t_b) = (s_a + t_a - t_b - m_{l_1}^2)(\hat{s} - s_b + t_a - m_{a_0}^2)(s_b - M_W^2 - m_{l_2}^2)/M_W^2
\]

\[
- (t_a - m_{a_0}^2 - m_{l_1}^2)(s_b - t_a + t_b - m_{a_0}^2)(s_b - M_W^2 - m_{l_2}^2)/M_W^2
\]

\[
+ (m_{a_0}^2 + m_{l_2}^2 - t_b)(m_{a_0}^2 + m_{l_2}^2 - t_a)
\]

\[
+ (\hat{s} - s_b + t_b - m_{b_0}^2)(\hat{s} - s_b + t_a - m_{a_0}^2), \quad (23)
\]
and \( a' = 2\sqrt{2} G_F^2 M_W^6 / (6(2\pi)^4) \).

### III. NEUTRINO DECAY WIDTH

As shown in Eq. (20) and (22), the cross sections of the process \( h_1 h_2 \rightarrow l^+ l^+ X \), and \( h_1 h_2 \rightarrow l^\pm l^\pm W^\mp X \) mediated by the resonant production of Majorana neutrino depend on the neutrino decay width \( \Gamma_N \). For \( m_N \ll M_W \) the decay width \( \Gamma_N \) receives contributions from leptonic and semi-leptonic channels

\[
N \rightarrow l_1 l_2 \nu, \; \nu_i (\nu^c_i) l_1, \; l^\pm q_1 \bar{q}_2, \; \nu_i (\nu^c_i) q \bar{q}.
\]

(24)

We approximate the semi-leptonic decays by inclusive quark-antiquark pair production. In such an approach [4] based on the Bloom-Gilman duality [10] the total decay width of the heavy Neutrino \( N \) to hadrons is reproduced in average with an accuracy sufficient for our estimates. One of the advantages of this approach is that it does not require knowledge of meson masses \( M_M \) and decay constants \( F_M \) necessary for the calculation of \( N \rightarrow l\mathcal{M} \) partial widths which then are summed up in order to derive the total decay width of heavy neutrino \( N \) in the channel-by-channel approach [6]. For heavy mesons these parameters are purely known and may introduce a significant uncertainty.

The leading order decay rates for the channels listed in (24) can be found for the leptonic and semileptonic decays in Refs. [4, 6, 8] and [4] respectively. We summarize the corresponding formulas neglecting lepton and quark masses. In this approximation the partial decay widths of Majorana neutrino in the region \( m_N \leq M_W \) are

\[
\Gamma(N \rightarrow l_1 l_2 \nu) = |U_{l_1N}|^2 \frac{G_F^2}{192\pi^3} m_N^5 F_W(m_N) \equiv |U_{l_1N}|^2 \Gamma_l^{\ell\nu},
\]

(25)

\[
\Gamma(N \rightarrow \nu_i l_1 \bar{l}_2) = |U_{l_1N}|^2 \frac{G_F^2}{96\pi^3} m_N^5 F_Z(m_N) \times
\]

\[
\times [(g_R^l)^2 + (g_R^l)^2 + g_L^l g_R^l + \delta_{l_1 l_2} (g_R^l + 2g_L^l)] \equiv |U_{l_1N}|^2 \Gamma_l^{\nu l},
\]

(26)

\[
\sum_{l_2 = e, \mu, \tau} \Gamma(N \rightarrow \nu_i l_1 \bar{l}_2) = |U_{l_1N}|^2 \frac{G_F^2}{96\pi^3} m_N^5 F_Z(m_N) \equiv |U_{l_1N}|^2 \Gamma_l^{3\nu},
\]

(27)

\[
\Gamma(N \rightarrow l_1 \bar{u} d) = |U_{l_1N}|^2 (|V_{ud}|^2 + |V_{us}|^2 + |V_{cs}|^2) \frac{G_F^2}{64\pi^3} m_N^5 F_W(m_N) \equiv |U_{l_1N}|^2 \Gamma_l^{lud},
\]

(28)

\[
\Gamma(N \rightarrow \nu_i q \bar{q}) = |U_{l_1N}|^2 \frac{G_F^2}{32\pi^3} m_N^5 F_Z(m_N) \times
\]

\[
\times [(g_L^q)^2 + (g_R^q)^2 + g_L^q g_R^q] = |U_{l_1N}|^2 \Gamma_l^{\nu q},
\]

(29)
In Eq. (28) we neglected the small charged current $cb$ contribution. Here $g_L^l = -1/2 + \sin^2\theta_W, g_R^l = \sin^2\theta_W, g_L^u = 1/2 - (2/3)\sin^2\theta_W, g_R^u = -(2/3)\sin^2\theta_W, g_L^d = -1/2 + (1/3)\sin^2\theta_W, g_R^d = (1/3)\sin^2\theta_W$ are the SM neutral current lepton and quark couplings where $l = e, \mu, \tau$ and $u = u, c, t; d = d, s, b$.

In Eqs. (29)-(32) we use in our analysis of the like-sign dilepton production. Numerically $\Gamma \ll m_N$ holds for 1GeV $\leq m_N \leq 1$TeV and, therefore, Eqs. (20) and (22) are good approximations for the cross sections of the process $h_1h_2 \rightarrow l^+l^+jj(W)X$. 

The above formulas we use in our analysis of the like-sign dilepton production. Numerically $\Gamma \ll m_N$ holds for 1GeV $\leq m_N \leq 1$TeV and, therefore, Eqs. (20) and (22) are good approximations for the cross sections of the process $h_1h_2 \rightarrow l^+l^+jj(W)X$. 

The corresponding partial widths of Majorana neutrino are given by \[6, 11\]:

$$\Gamma_N(m_N \leq m_W) \approx \left( \sum_{l_i=e,\mu,\tau} |U_{lN}|^2 \right) \left[ 2\Gamma_1^l + 2\Gamma^{ud}_1 + \Gamma^{\nu}_2 + \Gamma^{3\nu} + \sum_{q=u,d,s,c,b} \Gamma^{\nu q} \right].$$

The factor 2 in the first two terms is due to Majorana nature of heavy neutrino $N$ which can decay into two non-equivalent charge conjugate final states. For $m_N > M_W$ the total decay of heavy neutrino is determined by the decay channels:

$$N \rightarrow l^\pm W^\mp, \quad N \rightarrow \nu Z, \quad N \rightarrow \nu H^0.$$ 

The corresponding partial widths of Majorana neutrino are given by \[6, 11\]:

$$\Gamma(N \rightarrow l^\pm W^\mp) = |U_{lN}|^2 \frac{G_F}{8\sqrt{2}\pi} m_N^3 (1 + \frac{2M_W^2}{m_N^2}) (1 - \frac{M_W^2}{m_N^2})^2 \theta(m_N - M_W) = |U_{lN}|^2 \Gamma^{(lW)},$$

$$\Gamma(N \rightarrow \nu_1(\nu_1^c)Z^0) = |U_{\nu_1N}|^2 \frac{G_F}{8\sqrt{2}\pi} m_N^3 (1 + \frac{2M_Z^2}{m_N^2}) (1 - \frac{M_Z^2}{m_N^2})^2 \theta(m_N - M_Z) = |U_{\nu_1N}|^2 \Gamma^{(\nu Z)},$$

$$\Gamma(N \rightarrow \nu_1(\nu_1^c)H^0) = |U_{\nu_1N}|^2 \frac{G_F}{8\sqrt{2}\pi} m_N^3 (1 - \frac{M_H^2}{m_N^2})^2 \theta(m_N - M_H) = |U_{\nu_1N}|^2 \Gamma^{(\nu H)}.$$ 

In the same way we can write the total decay width of a Majorana neutrino at $m_N > M_W$ as

$$\Gamma_N(m_N > M_W) \approx \left( \sum_{l=e,\mu,\tau} |U_{lN}|^2 \right) (2\Gamma^W + \Gamma^Z + \Gamma^H)$$

The above formulas we use in our analysis of the like-sign dilepton production. Numerically $\Gamma \ll m_N$ holds for 1GeV $\leq m_N \leq 1$TeV and, therefore, Eqs. (20) and (22) are good approximations for the cross sections of the process $h_1h_2 \rightarrow l^+l^+jj(W)X$. 

9
IV. LIMITS ON HEAVY MAJORANA NEUTRINO SECTOR

The remaining ingredient, which determines in the most crucial way the event rate of the like-sign dilepton production, is the heavy Majorana neutrino masses and their mixing with $\nu_e$, $\nu_\mu$, $\nu_\tau$. In the literature there are various limits on $M_N$ and $U_{\alpha N}$ for $\alpha = e, \mu, \tau$, extracted from direct experimental searches \[12\] for these particles in a wide region of masses, from the non-observation of lepton number violating decays, as well as from precision measurements of certain observable quantities \[13\]. In the present context we are interested in these limits for the mass range $M_N \geq 1\text{GeV}$. From the global fit of the electroweak precision measurements, including LEP data, it was found \[13\]:

\[|U_{eN}|^2 \leq 0.0052, \quad |U_{\mu N}|^2 \leq 0.0001, \quad |U_{\tau N}|^2 \leq 0.01\]  (37)

The LEP data on direct searches of heavy leptons \[14\] shows that

\[|U_{lN}|^2 \leq 10^{-4} - 10^{-5} \quad \text{for} \quad 3\text{GeV} \leq M_N < 80\text{GeV}\]  (38)

where $l = e, \mu, \tau$.

Neutrinoless double beta decay ($0\nu\beta\beta$) is known to be a sensitive probe of Majorana neutrino masses and mixing. Presently the best the experimental lower bound \[15\] on the $0\nu\beta\beta$-decay half life was obtained for $^{76}\text{Ge}$:

\[T_{1/2}^{0\nu}(^{76}\text{Ge}) \geq 1.9 \times 10^{25}\text{yrs}.\]  (39)

In Ref. \[2\] this bound was used to constrain the contribution of Majorana neutrinos of arbitrary mass. In a good approximation this constrain reads:

\[\sum_k \frac{|U_{ek}|^2 e^{i\alpha_k} m_{\nu k}}{m_{\nu k}^2 + q_0^2} \leq 5 \times 10^{-8} \text{GeV}^{-1}.\]  (40)

with $q_0 = 105$ MeV. Here $\alpha_k$ denotes the Majorana CP-phase of the Majorana neutrino state $\nu_k$ of mass $m_{\nu k}$. For the light-heavy neutrino scenario one has:

\[\sum_{N=\text{heavy}} \frac{|U_{eN}|^2}{M_N} e^{i\alpha_N} + q_0^{-2} \sum_{i=\text{light}} |U_{ei}|^2 e^{i\alpha_i} m_{\nu i} \leq 5 \times 10^{-8} \text{GeV}^{-1}.\]  (41)

where $M_N \gg q_0$ and $m_{\nu i} \ll q_0$. This stringent constraint, applied to each term separately, seems to make unrealistic observation of like-sign dielectrons (positrons). However, the possibility of CP violation in the neutrino sector should be taken into account. If there are
FIG. 3: The cross section of \( p + p(\bar{p}) \rightarrow e^\pm e^\mp jj + X \) mediated by a heavy Majorana neutrino, at the LHC (solid line) and Tevatron (dashed line), as a function of \( m_N \). The thick lines correspond to the constraint from the 0\( \nu \beta \beta \)-decay experimental data and the thin lines are calculated from the fixed mixing elements \( |U_{\tau N}|^2 \sim |U_{\mu N}|^2 \sim |U_{eN}|^2 = 10^{-3} \).

more than one heavy neutrino state \( N_k \), with different CP phases \( \alpha_N \), then different terms in the first sum of (41) may compensate each other, reducing the individual 0\( \nu \beta \beta \)-constrain on each of them. Note that even with the presence of only one heavy neutrino there may happen the above mentioned compensation between the heavy \( N \) and the three light neutrinos \( \nu_i \), since they all coherently contribute to 0\( \nu \beta \beta \)-decay. Thus searching for like-sign dielectron (positron) events may provide evidence for CP violation in the neutrino sector. In case of observation of these events at a rate larger than that stemming from the 0\( \nu \beta \beta \)-limit applied to each term in (41), this would point to CP-violation in the neutrino sector.

With this in mind we have calculated the maximum of the \( p + p(\bar{p}) \rightarrow e^\pm e^\pm jjX \) and \( p + p(\bar{p}) \rightarrow e^\pm e^\pm W^\pm X \) cross sections at the Tevatron and LHC, consistent with the 0\( \nu \beta \beta \)-limit (41), assuming no CP violation in the neutrino sector. The results are shown by the thick lines in Figs 3 and 4. In the calculation we used the limits

\[
\frac{|U_{eN}|^2}{M_N} \leq 5 \times 10^{-8}\text{GeV}^{-1}, \quad |U_{\mu N}|^2 \leq 10^{-4},
\]

\[
|U_{\tau N}|^2 \leq 10^{-4} \text{ for } 3\text{GeV} \leq m_N \leq 80\text{GeV},
\]

\[
|U_{\tau N}|^2 \leq 10^{-2} \text{ for } 80\text{GeV} > m_N.
\]

(42)

For comparison, we also give the cross section calculated with the assumption \( |U_{\tau N}|^2 \sim \ldots \)
$|U_{\mu N}|^2 \sim |U_{e N}|^2 = 10^{-3}$, shown in Figs. 3 and 4 by thin lines.

With the same constraints (42) we calculated bounds on $p + p(\bar{p}) \to l^\pm l'^\pm j j X$ and $p + p(\bar{p}) \to l^\pm l'^\pm W^\mp X$ cross sections for other lepton flavors, at the LHC and Tevatron, and the results are plotted in Figs. 5 and 6.

In Tables I and II we show the upper limits for $p + p(\bar{p}) \to l^\pm l'^\pm j j X$ event number for all possible lepton flavors, calculated with the constraints in Eq. (42), at the LHC with integrated luminosity 10 fb$^{-1}$, and at the Tevatron with integrated luminosity 2 fb$^{-1}$.

The numbers given in these tables, except for the like-sign $\mu\mu$-production, do not take into account event selection cuts necessary to suppress the background and which may also dramatically affect theoretical predictions for the signal. In the present paper we do not address this issue. In order to evaluate an impact of the experimental cuts on our results we use the results of the recent analysis [5, 6] of the like-sign $\mu\mu$ selection criteria for the case of the LHC and TEVATRON experiments. In Ref. [5] there were found sets of the cuts which permit efficient suppression of the background and confident selection of $\mu^\pm\mu^\mp$-signal. In Tables I and II we indicate in brackets the number of $\mu^\pm\mu^\mp$-events after applying these cuts for the LHC and Tevatron experiments. As seen from these tables, the regions where there is still a chance to observe $\mu^\pm\mu^\pm$-production with statistical significance are 20 GeV $\leq m_N \leq$ 70 GeV for LHC and 20 GeV $\leq m_N \leq$ 50 GeV for Tevatron. As to the other lepton flavors there exist in the literature studies (see, e.g. Refs. [7, 8]) of the corresponding selection cuts, although not so detailed as in Ref. [6]. Unfortunately the effect of these cuts
FIG. 5: The bounds on the cross section of $p + p(\bar{p}) \rightarrow l^\pm l'^\pm jjX$ at the LHC (solid line) and Tevatron (dashed line) as a function of $m_N$ for $ll' = \mu\mu, \tau\tau, \mu\tau, e\mu$ and $e\tau$.

on our results cannot be directly recognized and requires a special analysis to be done elsewhere. Roughly, according to Refs. [7, 8], the cuts reduce the like-sign dilepton event by an order of magnitude for $m_N \leq 100\text{GeV}$, leaving practically unchanged the region of higher $m_N$ values.

The following comment is in order here. Notice that because of the involvement of all the mixing matrix elements $U_{eN}, U_{\mu N}, U_{\tau N}$ in the like-sign dilepton production cross section (20)-(22), one needs to rely on various scenarios relating these matrix elements, in order extract limits on them from experimental data. These additional not well motivated assumptions decrease reliability of such results. Naturally it would be desirable to get the possibility of extracting individual limits on each of the three mixing matrix elements $U_{eN}, U_{\mu N}, U_{\tau N}$ without any ad hoc assumptions on their relative sizes. Such an extraction is possible in the studied case of production of like-sign dileptons by simultaneous search for events with different lepton flavors. As follows from (20)-(22)

$$\sigma^{h_1 h_2 \rightarrow l e} + \sigma^{h_1 h_2 \rightarrow l \mu} + \sigma^{h_1 h_2 \rightarrow l \tau} = |U_{lN}|^2$$

(43)

Thus, searching for one flavor diagonal and two flavor non-diagonal processes and establishing upper bounds on each cross section within the same kinematical domain, one can set an upper limit on the corresponding mixing matrix element $|U_{lN}|$. Therefore searches for like-sign dileptons with all the lepton flavors are required for the model independent extraction of each mixing matrix element $|U_{lN}|$. An experimentally challenging point is searching for the
FIG. 6: The bounds on the cross section of $p + p(\bar{p}) \rightarrow l^+l^-W^\pm X$ at the LHC (left figure) and Tevatron (right figure) as a function of $m_N$ for $ll' = \mu\mu, \tau\tau, \mu\tau, e\mu$ and $e\tau$.

| $m_N$ (GeV) | $ee$ | $e\mu$ | $e\tau$ | $\mu\mu$ | $\mu\tau$ | $\tau\tau$ |
|-------------|------|--------|--------|-----------|-----------|-----------|
| 10          | 0    | 44     | 41     | $4.4 \cdot 10^3$ (0) | $8.3 \cdot 10^3$ | $3.8 \cdot 10^3$ |
| 20          | 0    | 80     | 78     | $4.0 \cdot 10^3$ (4) | $7.8 \cdot 10^3$ | $3.8 \cdot 10^3$ |
| 30          | 1    | 105    | 104    | $3.5 \cdot 10^3$ (10) | $6.9 \cdot 10^3$ | $3.4 \cdot 10^3$ |
| 50          | 1    | 101    | 101    | $2.0 \cdot 10^3$ (10$^2$) | $4.0 \cdot 10^3$ | $2.0 \cdot 10^3$ |
| 70          | 0    | 28     | 29     | 401(20) | 801       | $1.0 \cdot 10^3$ |
| 90          | 0    | 0      | 2      | 0        | 34        | 877       |
| 110         | 0    | 0      | 0      | 0        | 7         | 690       |

TABLE I: Maximum event number of $pp \rightarrow l^\pm l'^\pm jjX$ at the LHC $(10 \text{ fb}^{-1}$ data) consistent with the limits in Eq. 42. Indicated in brackets are typical number of events after the event selection cuts of Refs. [5, 6].

events with $\tau$’s since they involve missing momentum. In this respect it is worth mentioning a method for reconstructing $\tau$-momenta on the basis of the analysis of the isolated charged tracks from $\tau$-decays proposed in Ref. 6. The authors have argued their method to be a promising basis for future searches for like-sign dileptons with one or two $\tau$’s.
V. SUMMARY AND CONCLUSIONS

We have studied like-sign $\Delta L = 2$ dilepton inclusive production mediated by heavy Majorana neutrinos. We focused on the dominant mechanism via resonant Majorana neutrino production. We applied the existing limits on heavy Majorana neutrino mass $m_N$ and its mixing with active flavors $\nu_e, \nu_\mu, \nu_\tau$, in order to predict maximal like-sign dilepton production rates for $p + p(p) \to l^\pm l^\mp j j X$ and $p + p(p) \to l^\pm l'^\pm W^\mp X$, at the Tevatron and the LHC. Special emphasis has been made on the stringent limit on the heavy Majorana neutrino derived from $0\nu\beta\beta$-decay experiments. Our results displayed in Figs. 3-6 and Tables I, II have been obtained assuming no CP-violation in the neutrino sector.

As seen from Tables I and II, we expect no good prospects for observation of like-sign $ee$ events from the studied reactions neither at the Tevatron, with 2 fb$^{-1}$ integrated luminosity, nor at the LHC, with an integrated luminosity of 10 fb$^{-1}$. This is a direct consequence of very stringent $0\nu\beta\beta$-decay limits. With an increasing integrated luminosity 100 fb$^{-1}$ at the LHC, the like-sign $ee$ production can reach a few events in the region $20 \text{ GeV} \leq m_N \leq 60 \text{ GeV}$. For the $p + p(p) \to e^\pm e^\pm W^\mp X$ process, in which a real $W$ boson is detected 100 fb$^{-1}$ data at the LHC are sufficient for discovering Majorana neutrino through like-sign $ee$ production. One of the messages of the present paper is that despite of these discouraging predictions the like-sign $ee$ events are worth searching for both at Tevatron and LHC. The point is that any observation of such events in these experiments would point to the CP-violation in the neutrino sector, as explained in sec. IV.

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|c|}
\hline
$m_N$ (GeV) & $ee$ & $\mu\nu$ & $e\tau$ & $\mu\mu$ & $\mu\tau$ & $\tau\tau$ \\
\hline
10 & 0 & 1 & 1 & 104 (2) & 187 & 82 \\
\hline
20 & 0 & 2 & 2 & 95 (4) & 183 & 89 \\
\hline
30 & 0 & 2 & 2 & 82 (5) & 163 & 80 \\
\hline
50 & 0 & 2 & 2 & 47 (6) & 95 & 47 \\
\hline
70 & 0 & 1 & 1 & 9 (0) & 18 & 9 \\
\hline
90 & 0 & 0 & 2 & 0 & 1 & 8 \\
\hline
110 & 0 & 0 & 0 & 0 & 0 & 6 \\
\hline
\end{tabular}
\caption{Same as Table I, but for $p\bar{p} \to l^\pm l^\pm j j X$ at the Tevatron (2 fb$^{-1}$ data).}
\end{table}
Like-sign dilepton production for other flavors at the LHC and Tevatron have better prospects. In the case of $\mu^+\mu^\pm$ production, where we were able to estimate the effect of the event selection cuts, the possibility of their observation is open in the regions $20$ GeV $\leq m_N \leq 70$ GeV for LHC and $20$ GeV $\leq m_N \leq 50$ GeV for the Tevatron.

We proposed a method of model independent extraction of the heavy Majorana neutrino mixing $U_{lN}$ with the neutrino active flavors, on the basis of searches of certain sets of like-sign dileptons with different lepton flavors. This method will be helpful in reconstructing the structure of heavy neutrino sector.

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APPENDIX A: INTEGRATION LIMITS FOR $s_{a/b}$ AND $t_{a/b}$

The limits of integration over $s_{a/b}, s_{a/b}^+, t_{a/b}^-$ and $t_{a/b}^+$, given in Eqs. (20) and (22), can be obtained in the following way. The fourth order Gram determinant has the form

$$\Delta_4(P_W, l_2, k_a, k_b) = \begin{vmatrix} P_W \cdot P_W & P_W \cdot l_2 & P_W \cdot k_a & P_W \cdot k_b \\ l_2 \cdot P_W & l_2 \cdot l_2 & l_2 \cdot k_a & l_2 \cdot k_b \\ k_a \cdot P_W & k_a \cdot l_2 & k_a \cdot k_a & k_a \cdot k_b \\ k_b \cdot P_W & k_b \cdot l_2 & k_b \cdot k_a & k_b \cdot k_b \end{vmatrix} \quad (A1)$$

From the definitions in Eq. (15), we see that $\Delta_4$ is a function of $\hat{s}, s_a, s_b, t_a$ and $t_b$. Picking $s_b$ as the innermost integration variable, explicit evaluation of $\Delta_4$ yields

$$16\Delta_4 = as_b^2 + bs_b + c = a(s_b - s_b^-)(s_b - s_b^+). \quad (A2)$$

Therefore the $s_b$ limits are

$$s_b^+ = \frac{-b + \sqrt{\Delta}}{2a},$$
$$s_b^- = \frac{c}{a s_b^-} \quad (A3)$$

The $s_a$-integration limits can be obtained from the requirement $\Delta > 0$. Solving this equation yields

$$s_a^+ = (t_b(\hat{s} + m_b^2 - m_a^2) + \hat{s}(m_b^2 - m_{l_2}^2) + m_{l_2}^2(m_a^2 + m_b^2) - m_b^2(m_a^2 - m_{l_2}^2))$$
$$+ \lambda^{1/2}(\hat{s}, m_a^2, m_b^2)\lambda^{1/2}(t_b, m_b^2, m_{l_2}^2)/(2m_b^2),$$

$$s_a^- = (t_b(\hat{s} + m_b^2 - m_a^2) + \hat{s}(m_b^2 - m_{l_2}^2) + m_{l_2}^2(m_a^2 + m_b^2) - m_b^2(m_a^2 - m_{l_2}^2))$$
$$+ \lambda^{1/2}(\hat{s}, m_a^2, m_b^2)\lambda^{1/2}(t_b, m_b^2, m_{l_2}^2)/(2m_b^2).$$

16
\[
s_a^- = t_b + m_a^2 + ((t_1 + m_a^2 - m_b^2)(m_{i_2}^2 - t_a - t_b)
- \lambda^{1/2}(t_a, t_b, P_W^2)\lambda^{1/2}(t_a, m_{a_i}^2, m_{b_i}^2))/2t_a,
\]

where the third order Gram determinant has the form
\[
\lambda(x, y, z) = x^2 + y^2 + z^2 - 2xy - 2yz - 2xz.
\]

Defining
\[
Q = \hat{s}(t_b + m_b^2 - m_{i_2}^2) - m_{a_2}^2 t_b + m_{i_2}^2 (m_a^2 + m_b^2) - m_b^2 (m_b^2 + m_{i_2}^2 + P_W^2)
+ \lambda^{1/2}(\hat{s}, m_a^2, m_b^2)\lambda^{1/2}(t_b, m_{b_2}^2, m_{i_2}^2),
\]
\[
a' = 2m_b^2 (Q + m_a^2 (t_b + m_{a_2}^2 + m_{i_2}^2 + P_W^2))
\]
\[
b' = Q^2 - m_b^4 (t_b - m_{a_2}^2 - P_W^2 + m_{i_2}^2)^2 - 4m_b^4 (m_a^2 + m_{i_2}^2)(t_b + P_W^2)
\]
\[
c' = 2m_b^2 (Q(m_a^2 - m_{i_2}^2)(t_b - P_W^2)
+ m_b^2 (m_a^2 - m_{i_2}^2)^2 (t_b + P_W^2) + m_b^2 (t_b - P_W^2)^2 (m_a^2 + m_{i_2}^2)),
\]

we can arrive at the integration limit for \(t_a\):
\[
t_a^- = \frac{-b' - \sqrt{b'^2 - 4a'c'}}{2a'},
\]
\[
t_a^+ = \frac{c'}{a't_a^-}.
\]

Finally we give the \(t_b\)-integration limits as
\[
t_b^\pm = (m_{i_1} + \sqrt{P_W^2})^2 + m_{a_2}^2 - ((\hat{s} + m_a^2 - m_b^2)(\hat{s} + (m_{i_1} + \sqrt{P_W^2})^2 - m_{i_2}^2)
\pm (\lambda^{1/2}(\hat{s}, m_a^2, m_b^2)\lambda^{1/2}(\hat{s}, (m_{i_1} + \sqrt{P_W^2})^2, m_{i_2}^2)))/(2\hat{s}).
\]

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