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Lecture AB-III

HEAVY MAJORANA NEUTRINOS IN HIGH-ENERGY LEPTON-PROTON AND PROTON-PROTON COLLISIONS

Based on:

[1] A. Ali, A. V. Borisov, and D. V. Zhuridov, Moscow Univ. Phys. Bull., **59** (1), 19 (2004).
[2] A. Ali, A .V. Borisov, and D. V. Zhuridov, In: Particle Physics in Laboratory, Space and Universe, Proc. 11th Lomonosov Conf. on Elementary Particle Physics (Moscow, 21 – 27 August 2003), Ed. A. I. Studenikin (Singapore, World Scientific, 2005), p. 66.
[3] A. Ali, A. V. Borisov, and D. V. Zhuridov, Phys. At. Nucl., **68** (12), 2061 (2005).
[4] A. Ali, A. V. Borisov, and N. B. Zamorin, hep-ph/0104123.
[5] A. Ali, A .V. Borisov, and N. B. Zamorin, In: Frontiers of Particle Physics, Proc. 10th Lomonosov Conf. on Elementary Particle Physics (Moscow, 23 – 29 August 2001), Ed. A. I. Studenikin (Singapore, World Scientific, 2003), p. 74.

OUTLINE

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Introduction

At the moment there are experimental evidences for nonzero neutrino masses. The nature of neutrino mass, whether it is Dirac or Majorana, is one of the fundamental and still unsolved problems in particles physics. A Dirac neutrino carries a lepton number distinguishing a particle from an antiparticle. In contrast to that, a Majorana neutrino is identical to its own antiparticle. The Majorana mass term in the total Lagrangian does not conserve lepton number, but changes its value by two units. Therefore Majorana neutrinos can lead to various lepton number violating processes. For example, they induce same-sign dilepton production in collisions at high energies, among them

$$e^+p \to \bar{\nu}_e \ell^+ \ell'^+ X, \quad pp \to \ell^+ \ell'^+ X.$$

In theories extending the Standard Model the seesaw mechanism is often used to provide a natural generation of small neutrino masses. Unlike the usual way of Dirac mass generation through weak SU(2)-breaking, this mechanism doesn't need extremely small Yukawa couplings ($y \lesssim 10^{-12}$ for $m_\nu \lesssim 0.1$ eV). For three families of leptons and s right-handed SU(2) singlets the seesaw mechanism leads to 3 light and s heavy massive Majorana neutrino states

$$\nu_\ell = \sum_{i=1}^{3} \tilde{U}_{\ell i}\nu_i + \sum_{j=1}^{s} U_{\ell j} N_j,$$

where ν_ℓ is a neutrino of definite flavor ($\ell = e,\ \mu,\ \tau$), the coefficients $\tilde{U}_{\ell i}$ and $U_{\ell j}$ form the leptonic mixing matrices.

Heavy mass states give a relatively small contribution to neutrino flavor states. Nevertheless effects of light and heavy Majorana neutrinos (HMN) compete in lepton number violating processes, because small values of the mixing parameters $U_{\ell j}$ for heavy neutrinos N_j may be compensated by smallness of the masses of light neutrinos ν_i.

The contributions of light (l) and heavy (h) neutrinos to the cross section of the fundamental subprocess $WW \to \ell\ell'$ can be estimated as follows:

$$\sigma_l \sim \left|\tilde{U}_{\ell i}\tilde{U}_{\ell' i}\right|^2 \left(\frac{m_i^{(l)}}{m_W}\right)^2, \quad m_i^{(l)} \ll m_W,$$

$$\sigma_h \sim \left|U_{\ell j}U_{\ell' j}\right|^2 \left(\frac{s}{m_W m_j^{(h)}}\right)^2, \quad m_j^{(h)} \gg \sqrt{s}.$$

The light neutrino effects are negligible if

$$\left|\frac{\tilde{U}_{\ell i}\tilde{U}_{\ell' i}}{U_{\ell i}U_{\ell' i}}\right|^2 \left(\frac{m_i^{(l)} m_j^{(h)}}{s}\right)^2 \ll 1.$$

The process $e^+p \to \bar{\nu}_e \ell^+ \ell'^+ X$

We will consider the possibilities of observation of the process

$$e^+p \to \bar{\nu}_e \ell^+ \ell'^+ X \tag{1}$$

and its cross symmetric process $\nu_e p \to e\ell^+\ell'^+X$ (X denotes hadron jets) under the conditions of the present ep collider HERA (DESY)[] and of future ep colliders. We assume that these processes at high energies

$$\sqrt{s} \gg m_W$$

are mediated by HMN. The leading-order Feynman diagram for the process (1) is shown in Fig. 1. (There is also a crossed diagram with interchanged lepton lines.)

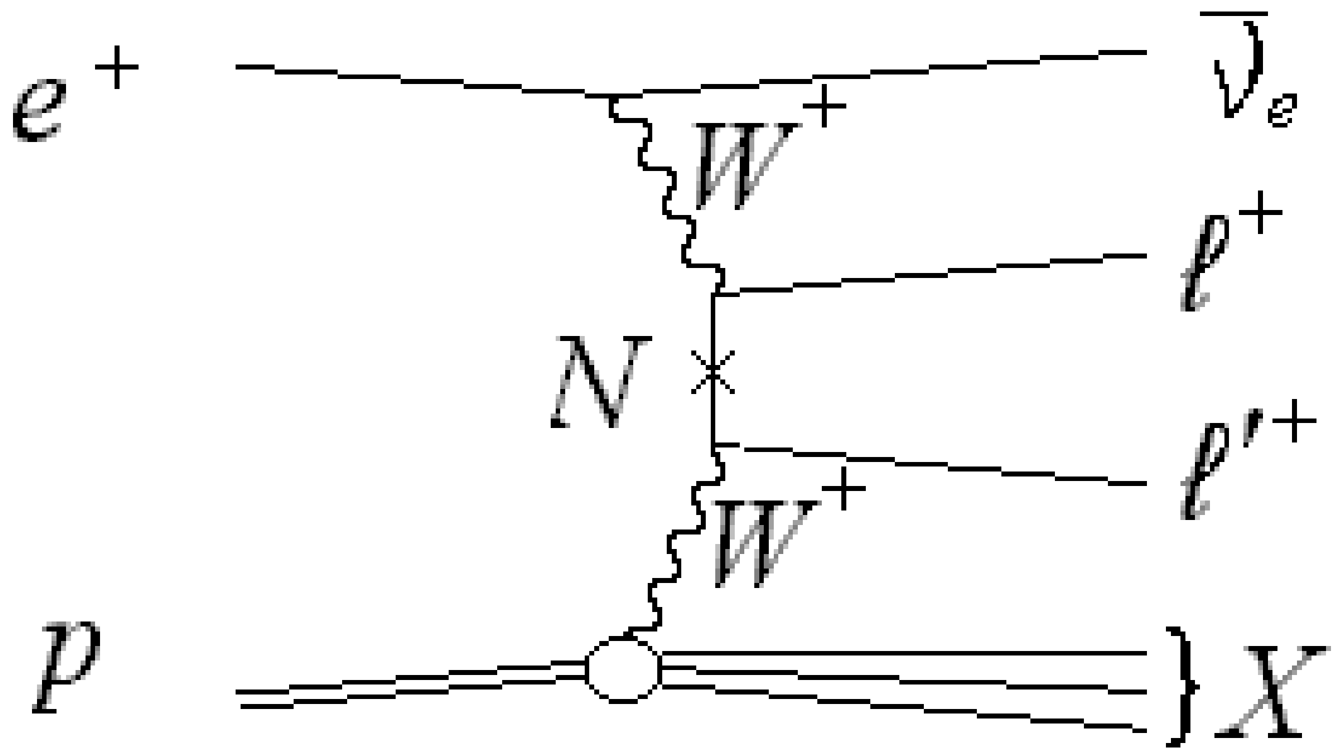

Fig. 1. Feynman diagram for the process $e^+p \to \bar{\nu}_e \ell^+ \ell'^+ X$ via an intermediate heavy Majorana neutrino N.

For calculating the cross sections, we use the leading effective vector-boson (EVB) approximation [M. Chanowitz and M. K. Gaillard (1984); G. L. Kane et al. (1984); S. Dawson (1985); I. Kuss and H. Spiesberger (1996)] neglecting transverse polarizations of W bosons and quark mixing. For this case, cross sections for the process and the crossed channel turn to be equal. As an observation criteria for the process we have chosen the condition

$$\sigma L \geq 1,$$

where σ denotes the cross section and L is the integrated luminosity per year for a collider.

Effective Singlet

At first we take the simplest pattern of the HMN mass spectrum

$$m_1 \ll m_2 < m_3 \cdots$$

$(m_{N_i} \equiv m_i)$ assuming the condition to be held

$$\sqrt{s} \ll m_2.$$

The cross section

$$\sigma_1 = C \left(1 - \frac{1}{2}\delta_{\ell\ell'}\right) |U_{\ell 1} U_{\ell' 1}|^2 \left(\frac{m_1}{m_W}\right)^2 \times \int\limits_{y_0}^{1} \frac{dy}{y} \int\limits_{y}^{1} \frac{dx}{x} p(x, xs) h\left(\frac{y}{x}\right) \omega\left(\frac{ys}{m_1^2}\right) \qquad (2)$$

for the process is determined by the convolution of three functions:

$$p\left(x,Q^2\right)=x\sum_i q_i\left(x,Q^2\right)=x\left(u+c+t+\bar{d}+\bar{s}+\bar{b}\right),$$

the quark distribution density having a fraction x of the proton momentum evaluated at the scale $Q^2=xs$,

$$h\left(z\right)=-\left(1+z\right)\ln z-2\left(1-z\right),$$

the normalized luminosity of W^+W^+ pairs in the quark-lepton system, and

$$w\left(t\right)=2+\frac{1}{t+1}-\frac{2\left(2t+3\right)}{t\left(t+2\right)}\ln\left(t+1\right),$$

the normalized cross section for *the fundamental lepton number violating subprocess*

$$W^+W^+\to\ell^+\ell'^+.$$

Here, $y_0=4m_W^2/s$ and the characteristic constant C has the value

$$C=G_F^4m_W^6/(8\pi^5)=0.80\text{ fb}. \tag{3}$$

In the numerical calculation we have used the set of parton distributions CTEQ6 []. We have used the bounds on the mixing parameters $U_{\ell N}$ from precision electroweak data []

$$\sum_N\left|U_{eN}\right|^2<6.6\times10^{-3},\sum_N\left|U_{\mu N}\right|^2<6.0\times10^{-3},$$

$$\sum_N |U_{\tau N}|^2_{eff} < 3.1 \times 10^{-3} \tag{4}$$

and the constraint from the neutrinoless double beta decay []

$$\left| \sum_N U^2_{eN} m^{-1}_N \right| < 5 \times 10^{-5}\ \mathrm{TeV}^{-1}$$

(the sum is over the heavy neutrinos).

The effective value

$$|U_{\tau N}|^2_{eff} = \mathrm{B}_{\tau\mu} \, |U_{\tau N}|^2 \tag{5}$$

with $\mathrm{B}_{\tau\mu} = \mathrm{Br}\left(\tau^- \to \mu^- \overline{\nu}_\mu \nu_\tau\right) = 0.1737$, as this τ-decay mode is most suitable for the like-sign dilepton detection at LHC.

We find that the process is practically unobservable at HERA even with a very optimistic luminosity ($\sqrt{s} = 318$ GeV, $L = 1$ fb^{-1}) and also at the projected supercollider VLHC []

$$\sqrt{s} = 6320\ \mathrm{GeV},\ L = 1.4\ \mathrm{fb}^{-1}.$$

For example, for $m_1 \sim 1$ TeV, we get

$$\sigma L \sim 10^{-10}\ (10^{-3})\ \text{for HERA (VLHC).}$$

For a possible detection of the process, the luminosity and/or the energy of the ep-collider should be substantially increased.

For instance, let the luminosity $L = 100$ fb^{-1} be achieved. In Fig. 2, the discovery limits are shown in the $(m_N, W_{\ell\ell'})$ plane, where

$$m_N \equiv m_1, W_{\ell\ell'} = \sqrt{2 - \delta_{\ell\ell'}}\, |U_{\ell N} U_{\ell' N}| \, .$$

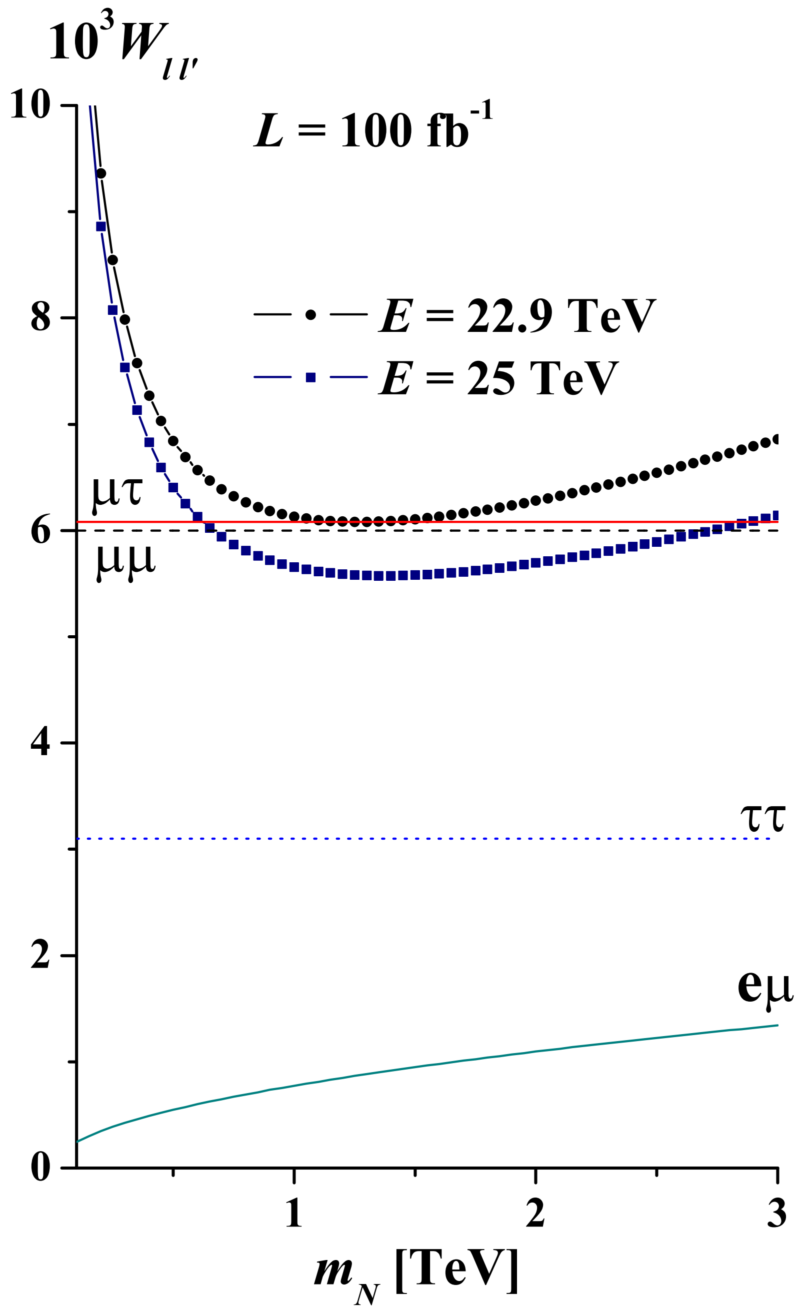

Fig. 2.

The observation of the most probable events ($\mu\tau$ and $\mu\mu$) is possible if $\sqrt{s} > 23$ TeV. For $\sqrt{s} = 25$ TeV such a collider will be sensitive to a range of neutrino masses about 1–3 TeV.

Effective Doublet

We consider also the neutrino mass spectrum of the effective doublet type

$$m_1 < m_2 \ll m_3 \cdots$$

with the bound on energy $\sqrt{s} \ll m_3$. In this scenario, the cross section for the process (1)

$$\sigma_2 = \frac{C}{2} \int_{y_0}^{1} \frac{dy}{y} \int_{y}^{1} \frac{dx}{x} p(x, xs) h\left(\frac{y}{x}\right) W\left(\frac{ys}{m_1^2}, \frac{ys}{m_2^2}\right) \tag{6}$$

includes the normalized cross section for the subprocess

$$\begin{aligned} W(t_1, t_2) = m_W^{-2} \left[\rho_1^2 m_1^2 \omega(t_1) + 2c\rho_1\rho_2 m_1 m_2 \Omega(t_1, t_2) \right. \\ \left. + \rho_2^2 m_2^2 \omega(t_2)\right], \end{aligned} \tag{7}$$

which contains the individual contributions of the neutrinos N_1 and N_2, $\omega(t_i)$, and the interference of the two, $\Omega(t_1, t_2)$, where

$$\begin{aligned} \Omega(t_1, t_2) = 2 - \frac{1}{t_1 + t_2 + t_1 t_2} \left[\frac{t_2(t_1^2 - 2t_1 t_2 - 2t_2)}{t_1(t_1 - t_2)} \right. \\ \left. \times \ln(1 + t_1) + (t_1 \leftrightarrow t_2)\right]; \end{aligned} \tag{8}$$

$$\omega(t) = \lim_{t' \to t} \Omega(t, t') = 2 + \frac{1}{1+t} - \frac{2(3 + 2t)}{t(2 + t)} \ln(1 + t).$$

The mixing parameters for different $\ell\ell'$ channels of the process are

$$\rho_i = \sqrt{2-\delta_{\ell\ell'}}\,|U_{\ell i}U_{\ell' i}|\,, \quad c = \cos\delta_{\ell\ell'}$$

with

$$\delta_{\ell\ell'} = \phi_1 - \phi_2 \in [0, 2\pi), \quad \phi_i = \arg\left(U_{\ell i}U_{\ell' i}\right).$$

The phases $\delta_{\ell\ell'}$ carry information about CP-violation.

We assume the saturation of the upper bound

$$B = 6.0 \times 10^{-3}$$

in the second sum in (4) only by the first two terms, i.e.,

$$|U_{\mu 1}|^2 = rB,\ |U_{\mu 2}|^2 = (1-r)B$$

with $r \in [0, 1]$. Then for the most probable $\mu\mu$ channel we obtain

$$\sigma_2 = A(r^2 f_1 + 2cr\bar{r}F_{12} + \bar{r}^2 f_2), \quad \bar{r} = 1 - r, \quad (9)$$

where $f_i = f(s, m_i)$ and $F_{ij} = F(s, m_i, m_j)$ are expressed through obvious convolutions of the functions ω and Ω with h and p, respectively (see Eq. (6)), the constant A has the value

$$A = 1.4 \times 10^{-5}\ \text{fb}.$$

For $r = 1$ ($r = 0$), only a single neutrino N_1 (N_2) contributes to the cross section which is reduced to the form given in Eq. (2). Generalization to the case of n neutrinos is straightforward.

In our calculations, we have chosen the following values for the parameters:

$$\sqrt{s} = 25 \text{ TeV}, \; m_1 = 1.3 \text{ TeV}, \; c = 1.$$

The cross section σ_2 (9) as a function of m_2 for various fixed values of r is shown in Fig. 3.

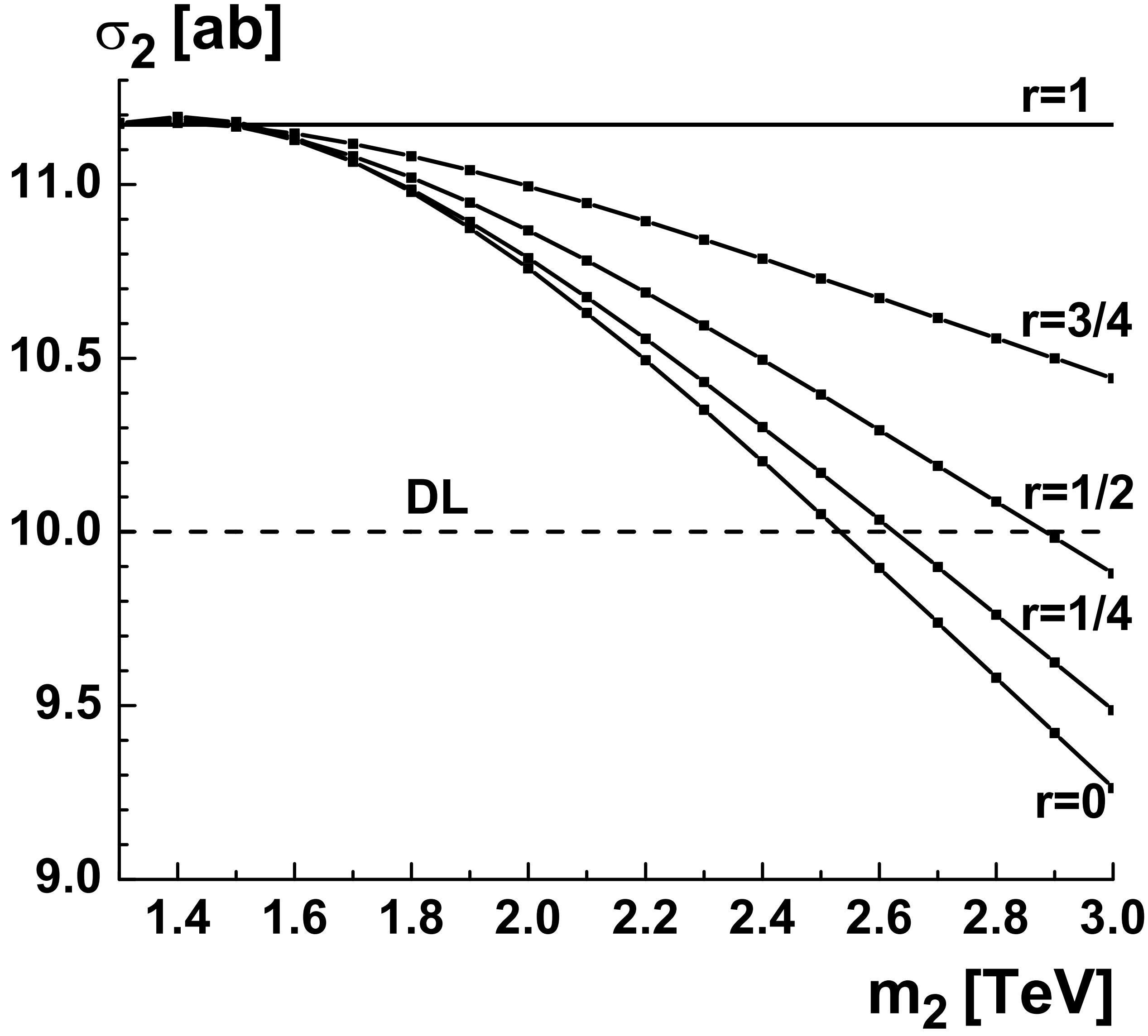

Fig. 3. σ_2 of m_2 dependence plots for $r = 0;\ 1/4;\ 1/2;\ 3/4;\ 1$ with $\sqrt{s}$=25 TeV, m_1=1.3 TeV, $c = 1$. Line DL is the discovery limit.

For the almost degenerate doublet ($m_1 \simeq m_2$) case and/or for the case of small mixing with N_2 ($r \simeq 1$), the cross section σ_2 is close to σ_1, the cross-section for the effective singlet case. But we should note that for the case of *destructive interference* of the two almost degenerate massive states (e.g., for $m_2 \simeq m_1$, $r = 1/2$, $c = -1$), the cross section is vanishingly small.

In Fig. 4 we show the discovery domain in the (m_1, m_2) plane for the $\mu\mu$ process at a collider with $L = 100\,\mathrm{fb}^{-1}$, $\sqrt{s} = 25$ TeV. The favorable values of the mixing parameters have been chosen:

$$\left|U_{\mu 1}\right|^2 = \left|U_{\mu 1}\right|^2 = 3 \times 10^{-3},\ c_{12} = 1.$$

The cross section as a function of the neutrino masses m_1 and m_2 is shown in Fig. 5.

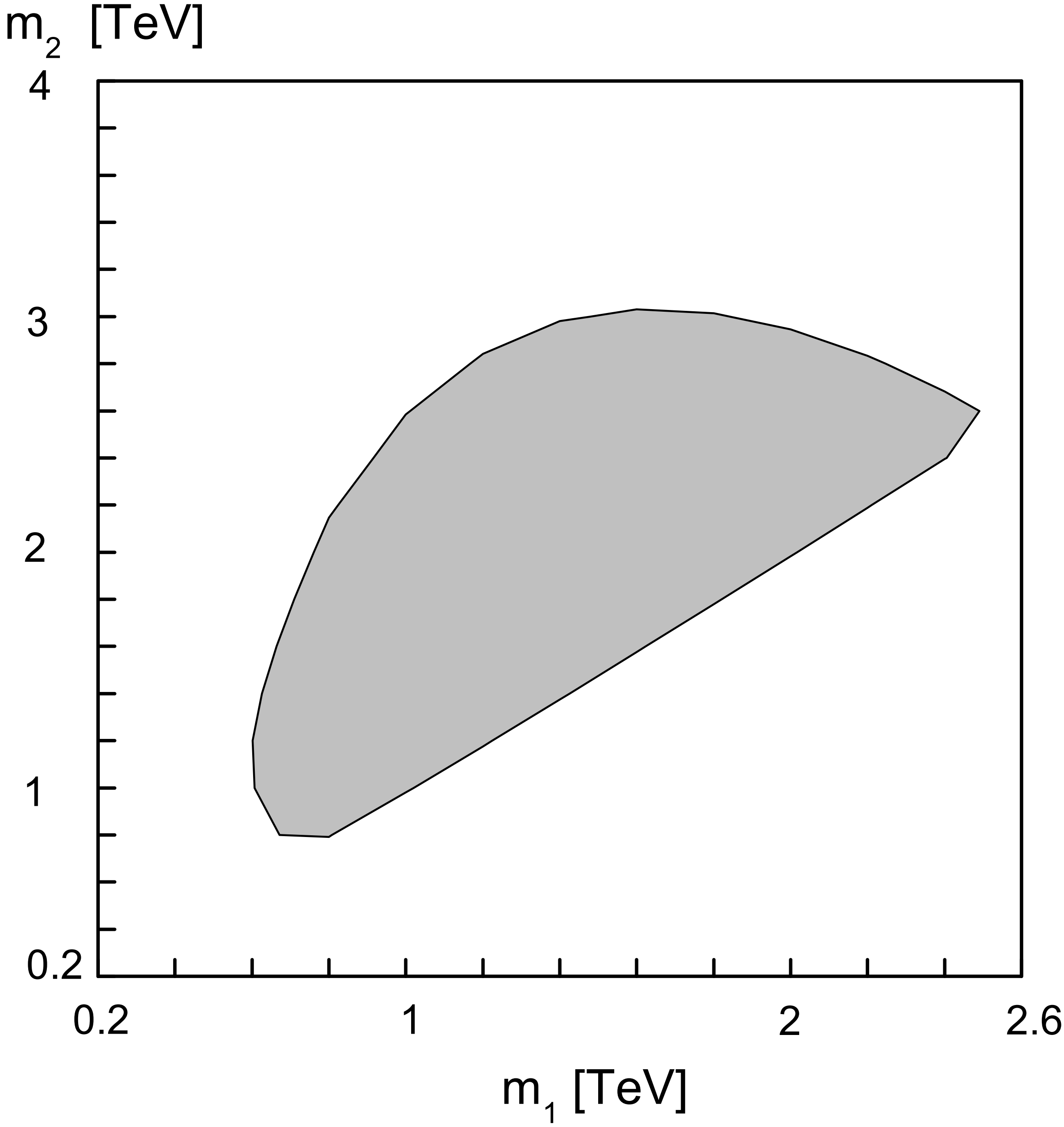

Fig. 4.

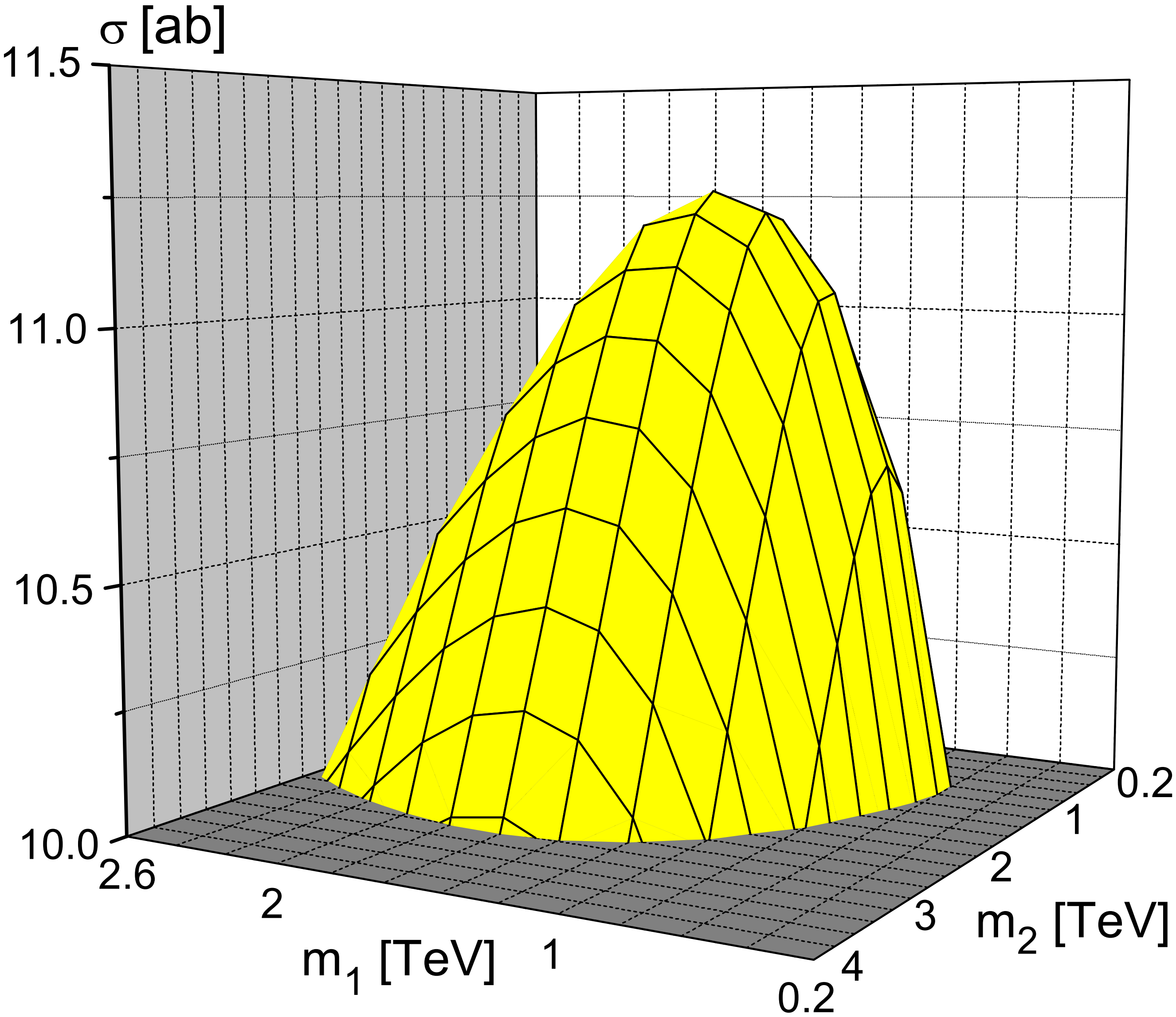

Fig. 5.

Effective triplet

In this case, the mass spectrum has the form

$$m_1 \lesssim m_2 \lesssim m_3 \ll m_4 \cdots,$$

and three HMN are effective. The cross section is

$$\sigma_3 = \frac{1}{2}\, C \left[\rho_1^2 f_1 + \rho_2^2 f_2 + \rho_3^2 f_3 + 2c_{12}\rho_1\rho_2 F_{12} + \right.$$
$$\left. + 2c_{13}\rho_1\rho_3 F_{13} + 2c_{23}\rho_2\rho_3 F_{23}\right], \qquad (10)$$

where $c_{ij} = \cos(\phi_i - \phi_j)$.

For

$$m_1 < m_2 \simeq m_3,$$

$F_{23} \simeq f_2 \simeq f_3$, $F_{12} \simeq F_{13}$, and the cross section σ_3 is similar to σ_2:

$$\sigma_3 \simeq \tfrac{1}{2}\, C \left[\rho_1^2 f_1 + 2\rho_1(c_{12}\rho_2 + c_{13}\rho_3) F_{12} \right.$$
$$\left. + (\rho_2^2 + 2c_{23}\rho_2\rho_3 + \rho_3^2) f_2\right].$$

In the simplest case of approximate degeneracy,

$$m_1 \simeq m_2 \simeq m_3; \rho_1 = \rho_2 = \rho_3,$$

the cross section

$$\sigma_3 \simeq \frac{2}{9}\sigma_1 \left(\frac{3}{2} + c_{12} + c_{13} + c_{23}\right),$$

where

$$\sigma_1 = C\rho_1^2 f_1/2.$$

For $\phi_1 = \phi_2 = \phi_3$ (all $c_{ij} = 1$), the cross section reaches the maximal value $\sigma_3 = \sigma_1$ (constructive interference). For destructive interference ($c_{12} + c_{13} + c_{23} =$

$-3/2$, e.g., at $\phi_1 - \phi_2 = \phi_2 - \phi_3 = 2\pi/3$), the cross section σ_3 is vanishingly small.

For arbitrary number n of HMN, the relation $\sigma_n \leq \sigma_1$ is valid.

Extracting the CP violating phases

Generically, the leptonic mixing matrix can be written in the form

$$U = D_R V D_M,$$

where

$$D_R = \mathrm{diag}(e^{i\varphi_e}, e^{i\varphi_\mu}, \ldots),$$
$$D_M = \mathrm{diag}(e^{i\alpha_1}, e^{i\alpha_2}, \ldots)$$

are diagonal phase matrices, and V is a (non-generic) unitary mixing matrix. Within the Standard model, D_R can be eliminated be redefining the phases of the charged lepton fields. The phases α_i are Majorana phases which are potentially observable. In particular, if there are only three Majorana neutrinos, U is a 3×3 matrix parameterized by six independent parameters: three mixing angles, one Dirac and two Majorana phases (an overall phase, common to all neutrinos, is not physical, and one is only sensitive to phase differences). In the general case of n neutrino flavors, one has $n(n-1)/2$ mixing angles, $(n-1)(n-2)/2$ Dirac phases and $n-1$ Majorana phases. In this case, the charged current leptonic mixing matrix U is a $3 \times n$ submatrix of the unitary $n \times n$ matrix, and there are $n-3$ sterile neutrinos.

For a multiplet HMN mass spectrum, the cross section of the process in question carries information about the CP violating phases due to interference of mass states.

For the effective doublet spectrum, five parameters

$$m_1, m_2, \rho_1, \rho_2, c_{12}$$

can be extracted from the measurements of the differential cross section

$$\frac{d^3\sigma}{dy dy' dp_\perp} = \frac{C}{8} \frac{S}{m_W^2 p_\perp \cosh^2 y_*} \int_{x_0}^{1} \frac{dx}{x} p(x, xs) \times R(x, s, y, y', p_\perp). \quad (11)$$

Here $y,\ y',\ p_\perp$ — rapidities and the transverse momentum of final leptons;

$$C = G_F^4 m_W^6/(8\pi^5), x_0 = 4p_\perp^2 \cosh^2 y_*/s,$$

and

$$S \equiv S_2 = \left| \sum_i M_i \rho_i e^{i\phi_i} g_i \right|^2$$
$$= \frac{1}{m_W^2} (\rho_1^2 M_1^2 g_1^2 + 2c_{12}\rho_1\rho_2 M_1 M_2 g_1 g_2 + \rho_2^2 M_2^2 g_2^2).$$

is the reduced differential cross section of the subprocess $W^+W^+ \to \ell^+\ell'^+$, where

$$g_i = \frac{1-z}{1-z+k_i} + \frac{1+z}{1+z+k_i}, \quad k_i = \frac{M_i}{2p_\perp^2 \cosh^2 y_*},$$
$$z \equiv \cos\theta_* = \tanh y_*,$$

θ_* is the scattering angle of the leptons in the center-of-mass frame of the W-pair;

$$R = 1 - \sqrt{\frac{x_0}{x}} \cosh Y + \frac{x_0}{x}$$

is the normalized spectrum of W-pair emission;

$$y_* = \frac{1}{2}(y-y'), Y = \frac{1}{2}(y+y').$$

The ranges of variables are

$$m_W / \cosh y_* \leq p_\perp \leq \sqrt{s}/(2 \cosh y_*),$$
$$-\infty < y, y' < \infty.$$

Fixing $\sqrt{s}$, the collider energy, and two of three kinematic variables, we obtain three different functions of one variable and five parameters $m_1, m_2, \rho_1, \rho_2, c_{12}$. The parameters in question can be extracted in principle from suitable measurements of each of these functions.

The cross sections of six possible $\ell\ell'$ processes are determined by twenty parameters:

$$2(m_1, m_2) + 18(\rho_1^{\ell\ell'}, \rho_1^{\ell\ell'}, c_{12}^{\ell\ell'}).$$

There are only nine independent mixing parameters (with $\ell = \ell'$) that determine nine parameters of the mixing matrix:

$$U_{\ell j} = \sqrt{\rho_j^{\ell\ell}} e^{i\varphi_j^{\ell\ell}/2}, \ \ell = e, \mu, \tau; \ j = 1, 2.$$

They are $\rho_j^{\ell\ell}$ (six) and $\Delta\varphi_\ell = \varphi_1^{\ell\ell} - \varphi_2^{\ell\ell}$ (three).

The process $pp \to \ell^+\ell'^+X$

We have calculated the cross section for the process at high energies,

$$\sqrt{s} \gg m_W, \tag{12}$$

via an intermediate heavy Majorana neutrino N in the leading effective vector-boson approximation neglecting transverse polarizations of W bosons and quark mixing. We use the effective singlet scenario for neutrino mass spectrum

$$m_{N_1} \equiv m_N \ll m_{N_2} < m_{N_3}, ...,$$

assuming

$$\sqrt{s} \ll m_{N_2}.$$

The cross section for the process

$$\sigma\left(pp \to \ell^+\ell'^+X\right) = C\left(1 - \tfrac{1}{2}\delta_{\ell\ell'}\right)\left|U_{\ell N}U_{\ell' N}\right|^2 \times F\left(E, m_N\right)\ , \tag{13}$$

with

$$C = \frac{G_F^4 m_W^6}{8\pi^5} = 0.80 \text{ fb}\ ,$$

and

$$F\left(E, m_N\right) = \left(\frac{m_N}{m_W}\right)^2 \int_{z_0}^1 \frac{dz}{z} \int_z^1 \frac{dy}{y} \int_y^1 \frac{dx}{x} \times p\left(x, xs\right) p\left(\tfrac{y}{x}, \tfrac{y}{x}s\right) h\left(\tfrac{z}{y}\right) w\left(\tfrac{s}{m_N^2}z\right)\ . \tag{14}$$

Here, $z_0 = 4m_W^2/s$, $E = \sqrt{s}$, and

$$w\left(t\right) = 2 + \frac{1}{t+1} - \frac{2\left(2t+3\right)}{t\left(t+2\right)} \ln\left(t+1\right)$$

is the normalized cross section for the subprocess $W^+W^+ \to \ell^+\ell'^+$ (in the limit $\sqrt{s} \gg m_W$, it is obtained from the well-known cross section for $e^-e^- \to W^-W^-$ [] using crossing symmetry). The function $h(r)$ defined as

$$h\left(r\right) = -\left(1+r\right)\ln r - 2\left(1-r\right)$$

is the normalized luminosity (multiplied by r) of W^+W^+ pairs in the two-quark system [], and

$$p\left(x, Q^2\right) = x \sum_i q_i\left(x, Q^2\right) = x\left(u + c + t + \bar{d} + \bar{s} + \bar{b}\right)$$

is the corresponding quark distribution in the proton.

The reduced cross section $F(E, m_N)$ as a function of the neutrino mass m_N is shown in Figs. 6 and 7 for the LHC energy $E = \sqrt{s} = 14$ TeV.

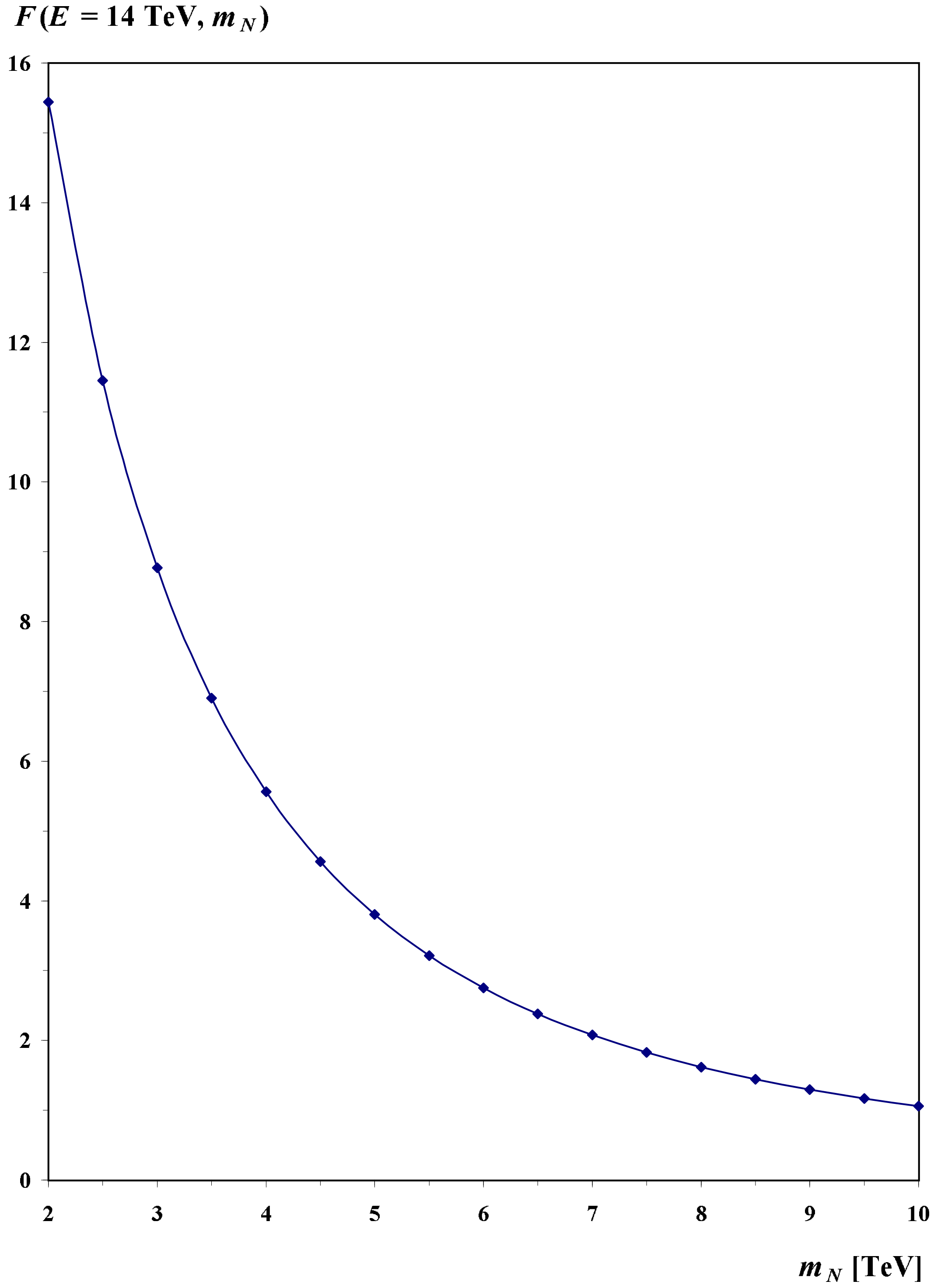

Fig. 6. The reduced cross section $F(E, m_N)$ for dilepton production as a function of the heavy Majorana mass m_N at LHC with $E = 14$ TeV.

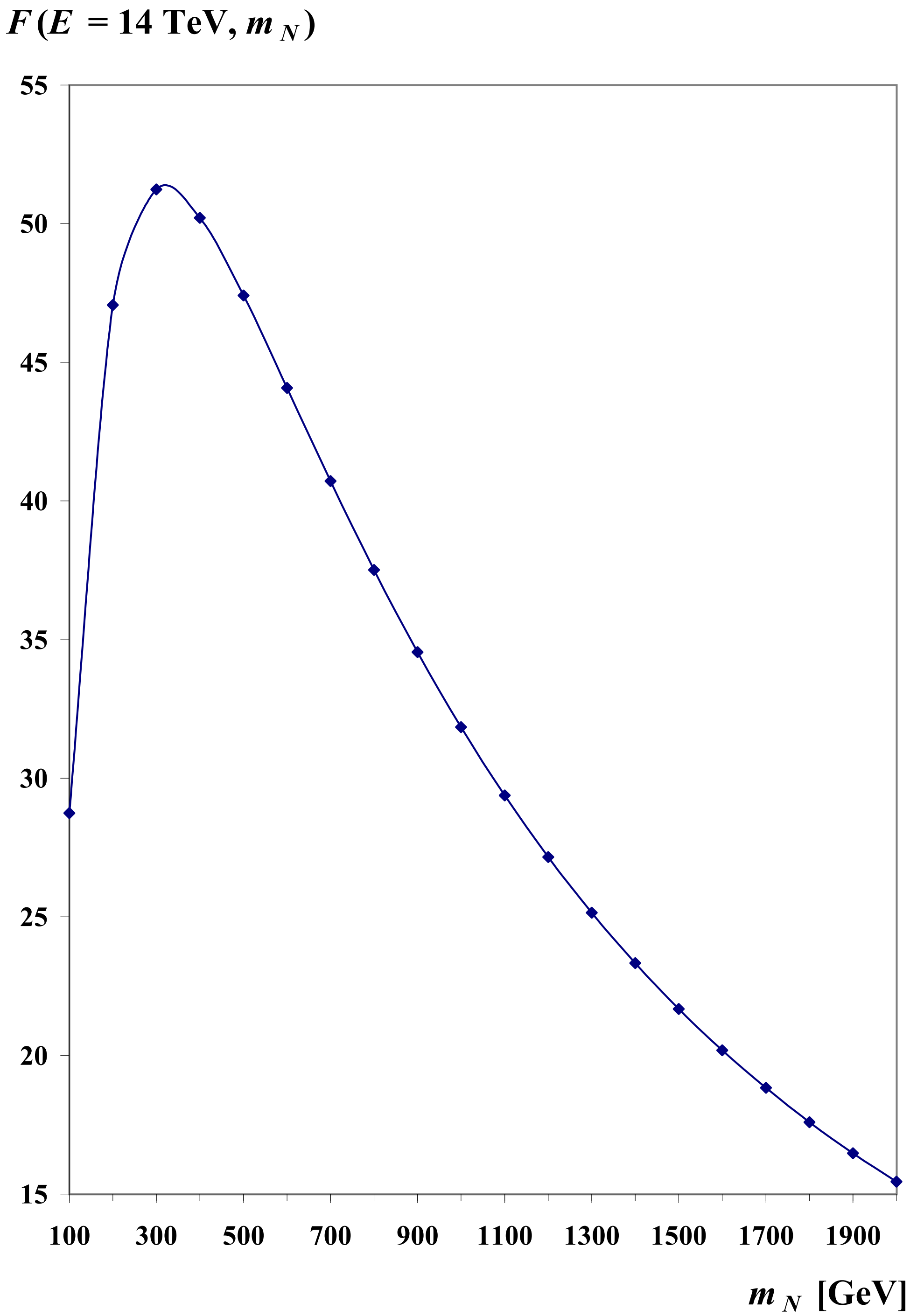

Fig. 7. The same as Fig. 6 but for lighter Majorana neutrinos.

We assume the mixing constraints obtained from the precision electroweak data

$$\Sigma\,|U_{eN}|^2 < 6.6 \times 10^{-3}, \quad \Sigma\left|U_{\mu N}\right|^2 < 6.0 \times 10^{-3},$$
$$\Sigma\,|U_{\tau N}|^2_{eff} < 3.1 \times 10^{-3}, \tag{15}$$

and from $0\nu\beta\beta$:

$$\left|\Sigma_{N(heavy)} U_{eN}^2 \frac{\eta_N}{m_N}\right| < 5 \times 10^{-5}\ \mathrm{TeV}^{-1}. \tag{16}$$

Taking these constraints a nominal luminosity $L = 100\ \mathrm{fb}^{-1}$, we obtain

$$\sigma L < 1 \text{ for all } \ell\ell' \text{ channels.}$$

Therefore at the LHC there is no room for observable signals for same-sign dilepton processes $pp \to \ell\ell' X$ $(\ell, \ell' = e, \mu, \tau)$ *due to the existing constraints for the mixing elements* $|U_{\ell N}|^2$ *from the precision electroweak data and neutrinoless double beta decay.*

Let us take into account a recent proposal to increase the instantaneous LHC luminosity $\mathcal{L}$ to a value of 10^{35} cm^{-2}s^{-1}, i.e., a total luminosity $L = \mathcal{L} \times$ 1 year $\simeq 3200$ fb^{-1}.

Using the upgraded LHC luminosity and demanding $n = 1,\ 3$ events for discovery (i.e., $\sigma L > n$), we present the two-dimensional plot for the discovery limits in Figs. 8 and 9 for the case of identical same-sign leptons ($\ell = \ell'$) and in Figs. 10 and 11 for the case of distinct same-sign leptons, ($\ell\ell' = e\mu,\ e\tau,\ \mu\tau$).

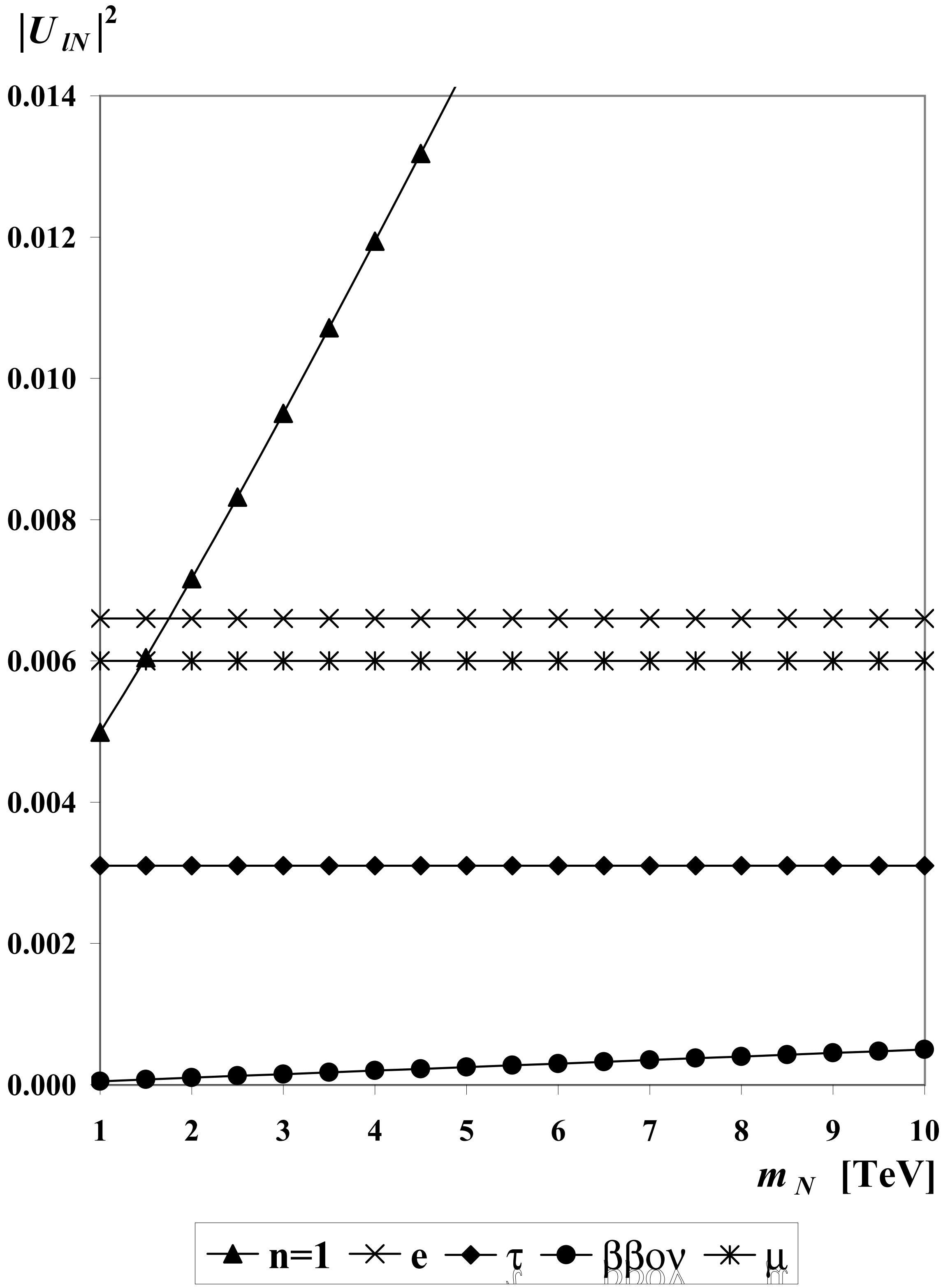

Fig. 8. Discovery limits for $pp \to \ell^+\ell^+X$ as functions of m_N and $|U_{\ell N}|^2$ for $E = 14$ TeV, $L = 3200\ \mathrm{fb}^{-1}$. We also superimpose the experimental limit from $\beta\beta_{0\nu}$, as well as the experimental limits on $|U_{\ell N}|^2$.

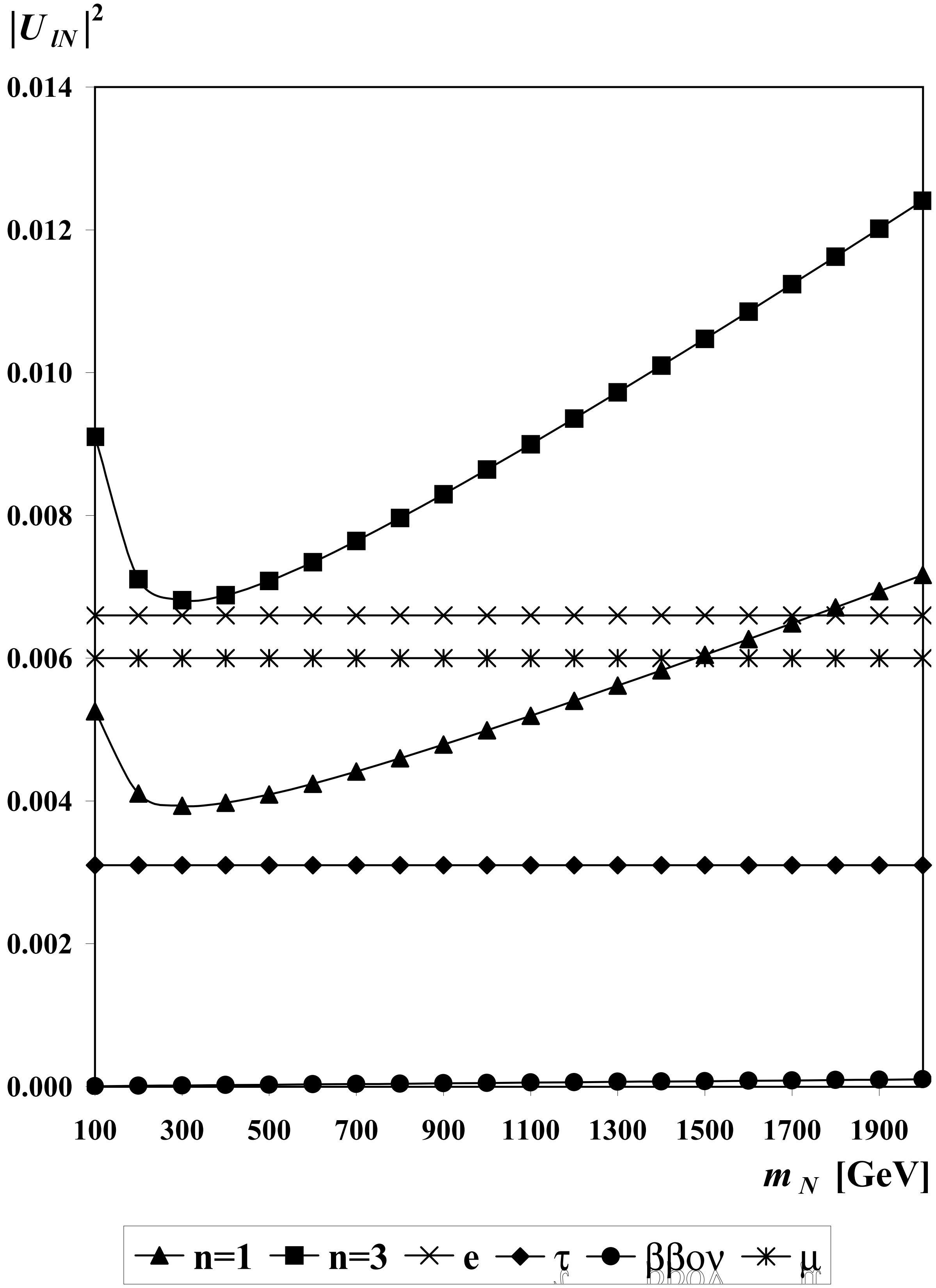

Fig. 9. The same as Fig. 8 but for lighter Majorana neutrinos.

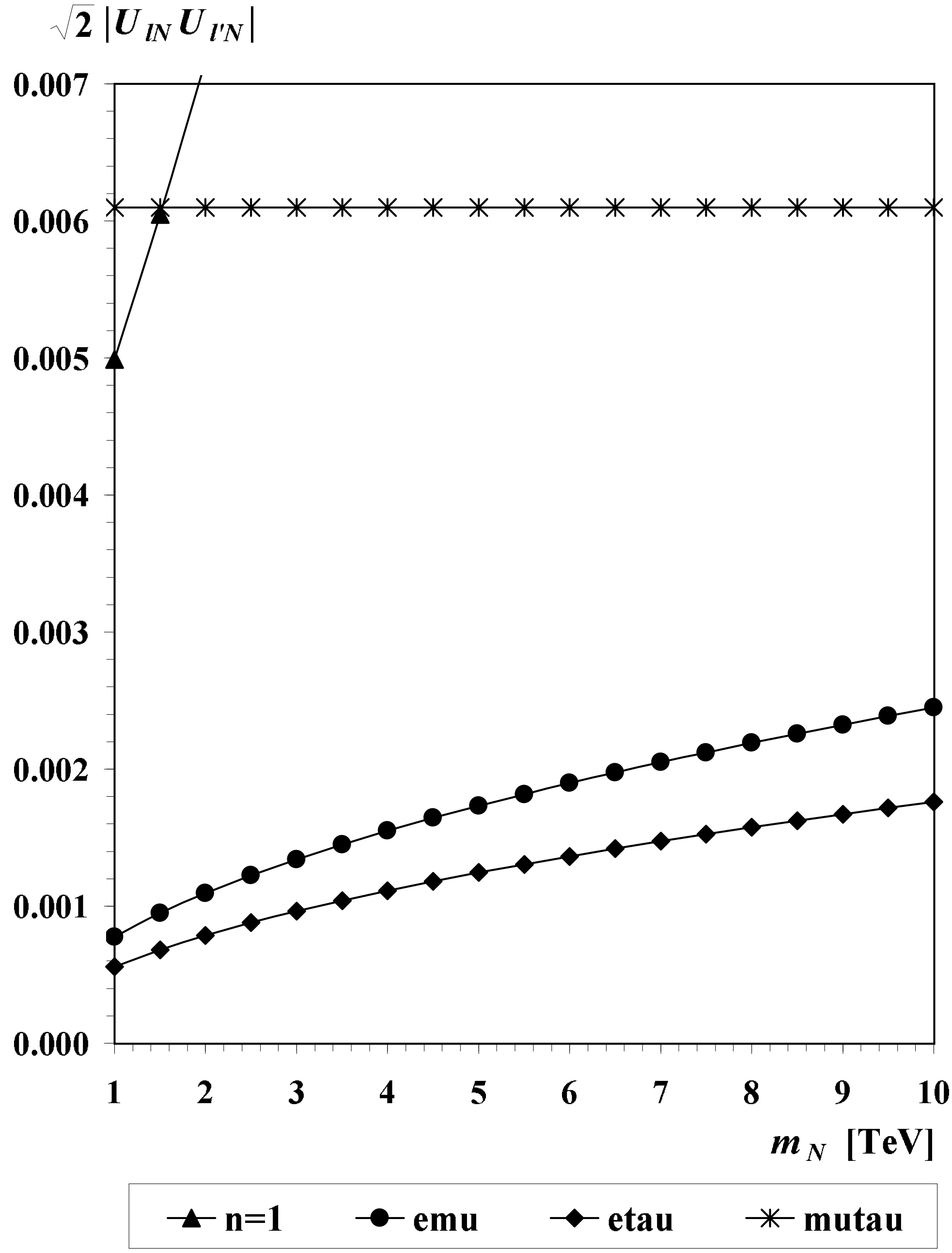

Fig. 10. Discovery limits for $pp \to \ell^+\ell'^+X$, $\ell\ell' = e\mu,\ e\tau,\ \mu\tau$. We also superimpose the limits on $\sqrt{2}\,|U_{\ell N}U_{\ell' N}|$ obtained from the experimental bounds.

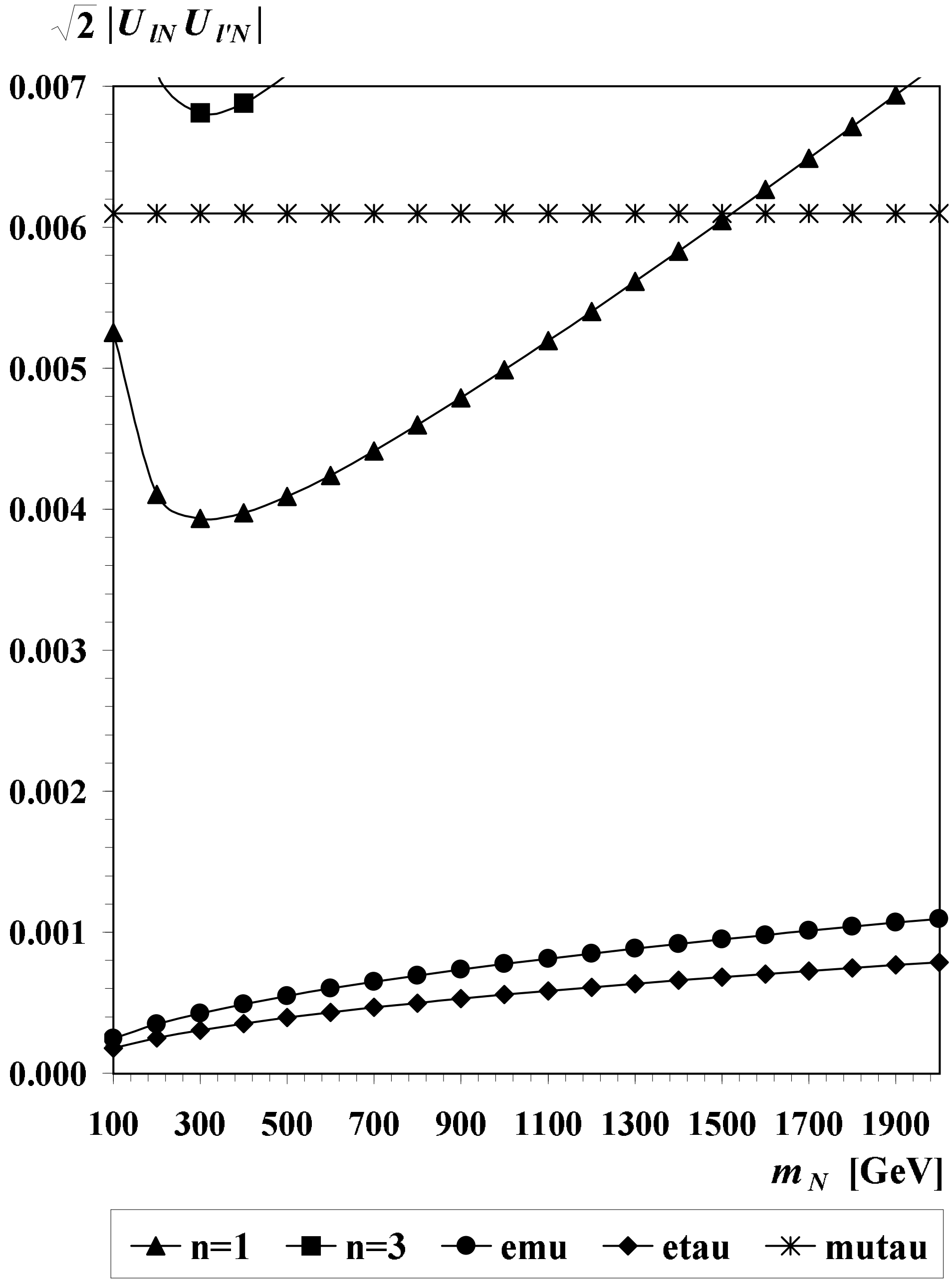

Fig. 11. The same as Fig. 10 but for lighter Majorana neutrinos.

From Figs. 8 – 11 we see that the existing strong constraints for the mixing elements $|U_{\ell N}|^2$ allow a possibility to observe only the same-sign $\mu\mu$ and $\mu\tau$ processes after increasing the nominal LHC luminosity $L = 100$ fb^{-1} by a factor of about 30. For this case, LHC experiments will have a sensitivity to heavy Majorana neutrinos of mass $m_N \lesssim 1.5$ TeV.

Conclusion

We have discussed the prospects of detecting the processes $e^+p \to \bar{\nu}_e \ell^+ \ell'^+ X$ and $\nu_e p \to e \ell^+ \ell'^+ X$ ($\ell, \ell' = e, \mu, \tau$) under the conditions of the present ep collider HERA and of future colliders. These high-energy processes are assumed to be mediated by the exchange of heavy Majorana neutrinos (HMN). We find that these processes are practically unobservable at HERA. For their possible detection, the luminosity and/or the energy of the ep-collider should be substantially increased. We have considered three simple scenarios for the HMN mass spectrum: the effective singlet ($m_1 \ll m_2 \lesssim m_3$ ...), doublet ($m_1 \lesssim m_2 \ll m_3$ ...), and triplet ($m_1 \lesssim m_2 \lesssim m_3 \ll m_4$ ...). For the two latter cases, the cross sections include information about CP-violating phases due to the effect of interference of neutrino mass states.

We have discussed how in principle neutrino masses and mixing parameters (including the phases) can be extracted from suitable data.

We have examined the process $pp \to \ell^+ \ell'^+ X$ mediated by Majorana neutrinos at the LHC energy. We find that at LHC for the nominal luminosity $L = 100$ fb^{-1} there is no room for observable same-sign dilepton signals due to the existing constraints from the precision electroweak data and neutrinoless double beta decay. But increasing the nominal LHC luminosity by a factor of about 30 will allow a possibility to observe the same-sign $\mu\mu$ and $\mu\tau$ processes mediated by a heavy Majorana neutrino of mass $m_N \lesssim 1.5$ TeV.