Heavy Majorana Neutrinos at e^+e^- Colliders

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We investigate possibilities for detecting heavy Majorana neutrinos (N's) in e^+e^- at LEP200 and future linear colliders. We concentrate on the processes where the pairs of intermediate heavy N's produce a clear signal of total lepton number violation ($e^+e^- \to NN \to W^+l^-W^+l'^-$). Such a signal is not possible if the heavy neutrinos are of Dirac nature. Our approach is general in the sense that the intermediate N's can be either on shell or off shell. Discussion of the relative numerical importance of the s and the $t + u$ channels of the NN production is also included. [S0031-9007(99)09285-6]

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There has been a significant amount of activity in the high energy physics community towards discerning the nature of the neutrino sector. A basic question is: Are neutrinos Dirac or Majorana particles? If there are no right-handed currents, then it is virtually impossible to discern the nature of the light neutrinos [1]. If there are heavy neutrinos ($M \sim 10^2$ GeV), then present and future experiments offer a realistic prospect of establishing their nature. The production cross section of heavy Majorana neutrinos (N's), mostly via the e^+e^- collisions, has therefore been investigated in the past [2]. Most of these works have been done within specific (classes of) models, and it has been assumed that the center-of-mass (CMS) energy $\sqrt{s}$ in the process is high enough for the production of on-shell (OS) heavy N's. The effects of the *off-shell* (nOS) N's have been ignored. Moreover, to our knowledge, various distributions of N's decay products [$N \to W^\pm \ell^\mp (\to \text{jets} + \ell^\mp)$], which are produced in the full reaction and which can actually be detected, have not been investigated in a quantitative way. The main reason for this was that the expressions for invariant amplitudes with two *on-shell* N's apparently do not allow a straightforward calculation of such distributions. We note that the detection of events for the reactions $e^+e^- \to NN \to W^\pm \ell^\mp W^\pm \ell'^\mp \to \text{jets} + \ell^\mp \ell'^\mp$, which violate the total lepton number, would be a clear signal of the Majorana character of the intermediate neutrinos.

We present some results of calculations for the aforementioned reactions. We do not restrict ourselves to any specific (classes of) models. In contrast to the available literature, our approach allows us to account also for the effects of off-shell intermediate N's on the cross sections σ. This enables us to investigate deviations from the previously known σ's, in the "2OS" kinematic region ($\sqrt{s} > 2M > 2M_W$) where both intermediate N's can, but need not, be on shell—these deviations are termed "finite width effects." Our approach allows us to calculate the σ's, and various distributions, also in the "1OS" region ($2M > \sqrt{s} > M + M_W$), where at most one intermediate N can (but need not) be on shell, and in the "nOS" region ($M + M_W > \sqrt{s}$), where both N's always have to be off shell. Our approach makes possible a straightforward calculation of various distributions of final particles. As an illustration, we include an angular distribution of the final state leptons ℓ, ℓ'.

We start with rather general Lagrangian densities for the couplings of N with Z, and for the coupling of N with W and light leptons ℓ_j ($\ell_1^- = e^-, \ell_2^- = \mu^-, \ell_3 = \tau^-$):

$$\mathcal{L}_{NNZ}(x) = -\frac{g}{4\cos\theta_W} A_{NZ} \overline{N}(x) \gamma^\mu \gamma_5 N(x) Z_\mu \,, \qquad (1)$$

$$\mathcal{L}_{N\ell W}(x) = -\sum_{j=1}^{3} \frac{g B_L^{(j)}}{2\sqrt{2}} \overline{\ell}_j(x) \gamma^\mu \gamma_- N(x) W_\mu^- + \text{H.c.}\,, \qquad (2)$$

where $\gamma_- = (1 - \gamma_5)$; $B_L^{(j)}$'s are, at first, free parameters; g and θ_W are the standard $\mathrm{SU}(2)_L$ gauge coupling parameter and the Weinberg angle, respectively. The vector part is absent in (1) because N's are Majorana. The right-handed parts were neglected in (2). The other relevant coupling is e^+e^-Z which we consider to be the one of the standard model (SM). We also set $A_{NZ} = -1$, i.e., by replacing, in the SM density for $\nu\nu Z$, the massless Dirac neutrino ν by the (heavy) Majorana neutrino N. These choices would suggest that the considered N is made up primarily of a sequential neutrino with the standard $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ assignments. However, these choices may also represent an approximation to other scenarios (cf. [2–4]). Further, parameters $B_L^{(j)}$ in (2) will affect the final results only via the combinations

$$H1 \equiv |B_L^{(1)}|^2, \qquad H \equiv \sum_{j=1}^{3} |B_L^{(j)}|^2. \qquad (3)$$

We restrict ourselves to the aforementioned reactions $e^+e^- \to NN \to W^\pm \ell_i^\mp W^\pm \ell_j^\mp (\to \text{jets} + \ell_i^\mp \ell_j^\mp)$ with

light leptons ($\ell_1 = e, \ell_2 = \mu, \ell_3 = \tau$). They involve the s and the $t + u$ (shortly: tu) channel—cf. Fig. 1.

For the calculation of the invariant amplitude $\mathcal{M}_{fi}$ (shortly: $\mathcal{M}$ [5]) for various channels, we used the 4-component spinors $u^{(\alpha)}(q) \equiv u(q\alpha)$ and $v^{(\alpha)}(q) \equiv v(q\alpha)$ as defined in [6], but with the normalization convention as given in [5]. For the quark spin components we use the notation $\tilde{\alpha} = 1, 2 \Leftrightarrow \alpha = 2, 1$. In the s channel, the resulting amplitude is

$$\mathrm{i}\mathcal{M}^{(s)} = \frac{4MA^{(s)}}{[s - M_Z^2 + \mathrm{i}\Gamma_Z M_Z]}(-1)^{\overline{\alpha}_\ell}[\overline{v}(\overline{p}\,\overline{\alpha})\gamma_\mu(A_+^{(e)} - A_-^{(e)}\gamma_5)u(p\alpha)] \times \{P_N(p_\ell p_w)P_N(\overline{p}_\ell\overline{p}_w)\overline{u}(p_\ell\alpha_\ell)\tilde{C}^\mu(p_w\overline{p}_w p_\ell\overline{p}_\ell)(1 + \gamma_5)v(\overline{p}_\ell\tilde{\overline{\alpha}}_\ell) + (p_w \leftrightarrow \overline{p}_w)\}. \quad (4)$$

Here, we use notations of Fig. 1; $s = (p + \overline{p})^2$; M_Z and Γ_Z are, respectively, the mass and the total decay width of Z; $A_+^{(e)}$ and $A_-^{(e)}$ are the vector and axial-vector coupling parameters of the e^+e^-Z coupling of the SM, respectively ($A_+^{(e)} = 4\sin^2\theta_W - 1, A_-^{(e)} = -1$). $\tilde{C}^\mu$'s are

$$\tilde{C}^\mu(p_w\overline{p}_w p_\ell\overline{p}_\ell) = \not{\epsilon}[(\not{p}_\ell + \not{p}_w)\gamma^\mu + \gamma^\mu(\not{\overline{p}}_\ell + \not{\overline{p}}_w)]\not{\overline{\epsilon}}, \quad (5)$$

where $\not{q} \equiv q_\nu\gamma^\nu$; $\varepsilon_\nu \equiv \varepsilon_\nu^{(\lambda)}(p_w)$ and $\overline{\varepsilon}_\nu \equiv \varepsilon_\nu^{(\overline{\lambda})}(\overline{p}_w)$ are the real polarization vectors [6] of the final W's, with polarizations $\lambda, \overline{\lambda} = 1, 2, 3$. M is the mass of N's; $A^{(s)}$ is

$$A^{(s)} = g^2 B_L^{(i)} B_L^{(j)} A_{NZ}\mathrm{i}\lambda_M/(128\cos^2\theta_W), \quad (6)$$

where λ_M is the phase factor in the Fourier decomposition of the Majorana field $N(x)$ (cf. [7]; $|\lambda_M|^2 = 1$). P_N in (4) is the (scalar) denominator of the propagator of N

$$P_N(p_\ell p_w) = 1/[(p_\ell + p_w)^2 - M^2 + \mathrm{i}M\Gamma_N], \quad (7)$$

where Γ_N is the total decay width of N.

All the combinations of attaching the four final particles to the two N's are accounted for in amplitude (4). Further, (4) had originally four instead of two terms in the curly brackets; however, two of them were reduced to the other two, by using the general identities

$$-\mathrm{i}\gamma^2 u(q\alpha)^* = (-1)^\alpha v(q\tilde{\alpha}), \quad -\mathrm{i}\gamma^2 v(q\alpha)^* = (-1)^{\tilde{\alpha}} u(q\tilde{\alpha}), \quad (8)$$

and $(\not{a}\not{b})^T = -\gamma^0\gamma^2(\not{b}\not{a})\gamma^2\gamma^0$, where the Dirac basis and the conventions of [6] are used for γ^μ's. Using (8), we can further rewrite (4) into a form involving $v(p_\ell\tilde{\alpha}_\ell)$ and $\overline{u}(\overline{p}_\ell\overline{\alpha}_\ell)$ instead of $\overline{u}(p_\ell\alpha_\ell)$ and $v(\overline{p}_\ell\tilde{\overline{\alpha}}_\ell)$. The complex conjugate (c.c.) of this alternative form *and* of the form (4) are needed to calculate later the s-tu interference term of the full $\langle|\mathcal{M}|^2\rangle$, where $\langle\cdots\rangle$ stands for summation over the final and average over the initial polarizations. For the s-s part of $\langle|\mathcal{M}|^2\rangle$, the form (4) is used.

The tu-channel amplitude $\mathcal{M}^{(tu)}$ turns out to be

$$\mathrm{i}\mathcal{M}^{(tu)} = \frac{4MA^{(tu)}P_N(p_\ell p_w)P_N(\overline{p}_\ell\overline{p}_w)}{[(p - p_\ell - p_w)^2 - M_W^2 + \mathrm{i}\Gamma_W M_W]}(-1)^{\overline{\alpha}_\ell}\{[\overline{u}(p_\ell\alpha_\ell)\tilde{\mathcal{A}}(p_w p_\ell)\gamma^\nu\gamma_- u(p\alpha)][\overline{v}(\overline{p}\,\overline{\alpha})\gamma_\nu\not{\overline{\epsilon}}\gamma_+ v(\overline{p}_\ell\tilde{\overline{\alpha}}_\ell)] + \frac{(p_\ell + p_w)^2}{M_W^2}[\overline{u}(p_\ell\alpha_\ell)\not{\epsilon}\gamma_- u(p\alpha)][\overline{v}(\overline{p}\,\overline{\alpha})\mathcal{A}(\overline{p}_w\overline{p}_\ell)\gamma_+ v(\overline{p}_\ell\tilde{\overline{\alpha}}_\ell)]\} + \ldots, \quad (9)$$

where $\tilde{\mathcal{A}}(p_w p_\ell) = (\not{\epsilon}\not{p}_w + 2p_\ell\cdot\varepsilon)$, $\mathcal{A}(\overline{p}_w\overline{p}_\ell) = (\not{\overline{p}}_w\not{\overline{\epsilon}} + 2\overline{p}_\ell\cdot\overline{\varepsilon})$, $\gamma_\pm = (1 \pm \gamma_5)$. The dots at the end of (9) stand for three analogous terms, obtained from the above explicit expression by replacements: (I) $(p_w, \varepsilon) \leftrightarrow (\overline{p}_w, \overline{\varepsilon})$; (II) $(p_\ell, \alpha_\ell) \leftrightarrow (\overline{p}_\ell, \overline{\alpha}_\ell)$ and overall factor (-1); (III) combined replacements (I) and (II). $A^{(tu)}$ in (9) is

$$A^{(tu)} = g^4|B_L^{(1)}|^2 B_L^{(i)} B_L^{(j)}\mathrm{i}\lambda_M/64. \quad (10)$$

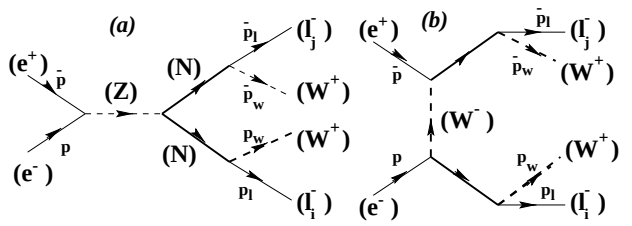

FIG. 1. An s-channel (a) and a tu-channel (b) diagram for $e^-(p\alpha)e^+(\overline{p}\,\overline{\alpha}) \to NN \to W^+(p_w\lambda)W^+(\overline{p}_w\overline{\lambda})\ell_i^-(p_\ell\alpha_\ell)\ell_j^-(\overline{p}_\ell\overline{\alpha}_\ell)$.

As in the s-channel case, we can reexpress any of the terms in $\mathcal{M}^{(tu)}$ in alternative forms, by applying transformations (8)—e.g., if we want to use, in scalar expressions in square brackets of (9), $u(\overline{p}\,\tilde{\overline{\alpha}})$ and $\overline{u}(\overline{p}_\ell\overline{\alpha}_\ell)$ instead of $\overline{v}(\overline{p}\,\overline{\alpha})$ and $v(\overline{p}_\ell\tilde{\overline{\alpha}})$. Such transformations are convenient when we calculate $\langle|\mathcal{M}|^2\rangle \equiv \langle|\mathcal{M}^{(s)} + \mathcal{M}^{(tu)}|^2\rangle$. Then we can always end up with traces involving $u(q,\beta)\overline{u}(q,\beta) = \not{q}$ and/or $v(q,\beta)\overline{v}(q,\beta) = \not{q}$ ($q = p, \overline{p}, \ldots; \beta = \alpha, \overline{\alpha}, \ldots = 1, 2$).

The integrand $\langle|\mathcal{M}|^2\rangle$ is long—the s-tu and (even more so) tu-tu parts extend over tens of pages when printed out. Numerical integration of $\langle|\mathcal{M}|^2\rangle$ over (parts of) the final phase space leads to the cross sections. This general (nOS) program, as mentioned, accounts for the effects of off-shell and on-shell N's.

The input were values of $\sqrt{s}$, M, $H1$, and H [cf. Eqs. (2) and (3)]. $H1$ measures the eWN coupling and affects the tu amplitude ($\propto H1$). H affects the total $\langle|\mathcal{M}|^2\rangle$ which is then *formally* $\propto H^2$ (if $H1$ is fixed). In $\langle|\mathcal{M}|^2\rangle$, we average over the initial $(\alpha, \overline{\alpha})$, and sum

over the final polarizations $(\lambda, \overline{\lambda}; \alpha_\ell, \overline{\alpha}_\ell)$ *and* over the flavors $(i, j = 1, 2, 3)$ of the two final light leptons. In the general (nOS) expression, an additional factor $1/4$ is included in $\langle|\mathcal{M}|^2\rangle$ to avoid double-counting of the two W^+'s and of the final leptons when integrating over the phase space.

By the same described methods we also calculated $\mathcal{M}$ and $\langle|\mathcal{M}|^2\rangle$ when one N or both N's are explicitly put on shell (1OS, 2OS expressions, respectively). The 1OS $\langle|\mathcal{M}|^2\rangle$, for the sum of reactions $e^+e^- \to NN_{\rm OS} \to W^+\ell_i^- N_{\rm OS}$ $(i = 1, 2, 3)$, was then multiplied by the branching ratio BR for the sum of the decay modes $N_{\rm OS} \to W^+\ell_j^-$ $(j = 1, 2, 3)$; the 2OS $\langle|\mathcal{M}|^2\rangle$ for $e^+e^- \to N_{\rm OS}N_{\rm OS}$ was multiplied by $({\rm BR})^2$.

Γ_N, appearing in (7), was determined at the tree level, assuming that the only (dominant) decay modes are $N \to W^\pm\ell_j^\mp$ $(j = 1, 2, 3)$ $[\Rightarrow \Gamma_N \propto H]$. Then BR $= 1/2$.

Numerical calculations were performed in various kinematic regions (nOS, 1OS, 2OS regions) with the general (nOS) expression [cf. (7)]. In the 1OS and 2OS regions, the 1OS expression was also used. In the 2OS region, the 2OS expression was also used. Results are depicted in Figs. 2 and 3. The Γ_N parameter H was set $H = 1$ in all these figures.

Figure 2 shows the M dependence of the cross section σ, at fixed $\sqrt{s}$. The difference between the results of the general (nOS) and the 2OS program, for $\sqrt{s} =$ 300 GeV, is less than 10% over most of the 2OS kinematic region $(M_W < M < \sqrt{s}/2)$, except near the threshold $M \approx \sqrt{s}/2$, where the results of the nOS program are significantly higher. The difference between the results of the 1OS and 2OS programs is less than 3% in most of the 2OS kinematic region. However, in the 1OS region $(\sqrt{s}/2 < M < \sqrt{s} - M_W)$, the results of the 1OS program are usually by several factors lower than those of the general nOS program, except near the threshold $M \approx \sqrt{s}/2$, where they differ only little. All these differences are in general less pronounced when the *tu* channel is excluded $(H1 = 0.0)$. For the chosen *tu*-strength value $H1 = 0.25$, the contributions of the *tu*-*tu* channel are at least by 1 order of magnitude larger than those of the *s*-*s* channel. Each curve has two slope increases: at $M \approx \sqrt{s}/2$ and at $M \approx \sqrt{s} - M_W$ (onset of the 2OS, 1OS kinematic region, respectively). If we take the integrated luminosity at LEP200 ($\sqrt{s} = 200$ GeV) to be 500 pb^{-1}, Fig. 2 implies that the maximal number of events cannot exceed 17 and 112 if $H1 = 0.0, 0.25$, respectively, and we assume $M > 85$ GeV.

In Fig. 3 we show the $\sqrt{s}$ dependence of σ, at fixed M. Most of the remarks about Fig. 2 apply also to Fig. 3. The differences between the results of the nOS and 2OS programs (finite width effects) are very significant when $M = 200$ GeV (and $H1 = 0.25$, $H = 1.0$), because $\Gamma_N = 4.8$ GeV is large then $(\Gamma_N/M \approx 2.5\%$; for $M = 150$ GeV: $\Gamma_N/M \approx 1.1\%)$.

Our general nOS program can be applied also to calculation of various distributions in the process. As an illustration, we present in Fig. 4 an angular distribution of the final leptons ℓ^-, ℓ'^-. The corresponding total cross sections are $\sigma = 0.280$ pb, 0.007 pb, for $M = 200$, 255 GeV, and the kinematic regions are 2OS, 1OS, respectively. If linear colliders at $\sqrt{s} = 500$ GeV achieve the integrated luminosity of 10^4 pb^{-1}, and most of the final states $W^+W^+\ell_i^-\ell_j^-$ can be identified, then these σ's will correspond to 2800 and 70 events, respectively.

In the 2OS kinematic region, the σ's and distributions as given numerically by the general (nOS) program depend on H weakly. Parameter H (note: $\Gamma_N \propto H$) is responsible in the 2OS region for the deviation of the full σ from the pure 2OS σ. In the 1OS region, H dependence of the full σ becomes quite strong (approximately $\sigma \propto H$), and in the nOS region even more so $(\propto H^2)$. In Figs. 2–4, we chose $H = 1$.

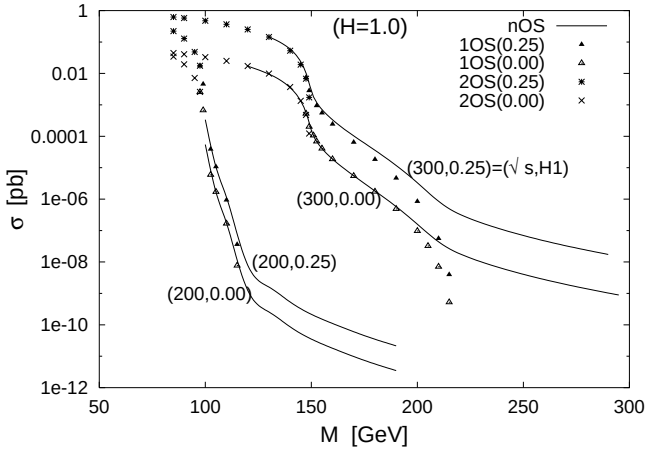

FIG. 2. Sum of cross sections for $e^+e^- \to NN \to W^+W^+\ell_i^-\ell_j^-$ $(\ell_1 = e, \ell_2 = \mu, \ell_3 = \tau)$, as function of neutrino mass M, for $\sqrt{s} = 200$, 300 GeV and the *tu* strengths $H1 = 0.0$, 0.25. Full lines are results of the general (nOS) program. Triangles and crosses are results of the 1OS and 2OS programs, respectively.

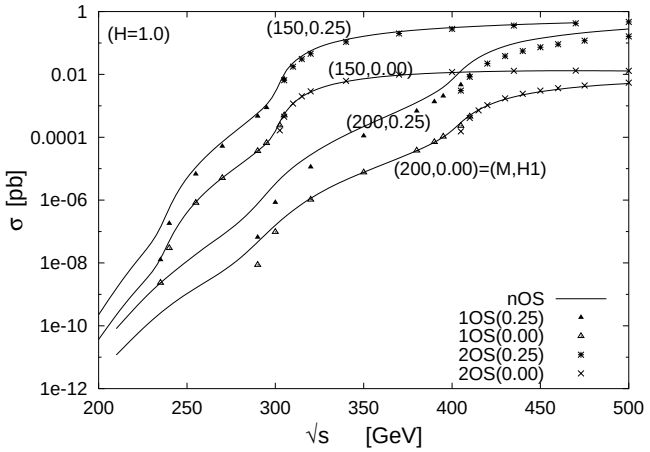

FIG. 3. Sum of the cross sections for the mentioned reactions, as function of the CMS energy $\sqrt{s}$, at fixed $M = 150$, 200 GeV. Again $H1 = 0.0$, 0.25. Results of various programs (nOS, 1OS, 2OS) are displayed in the kinematic regions where they are applicable.

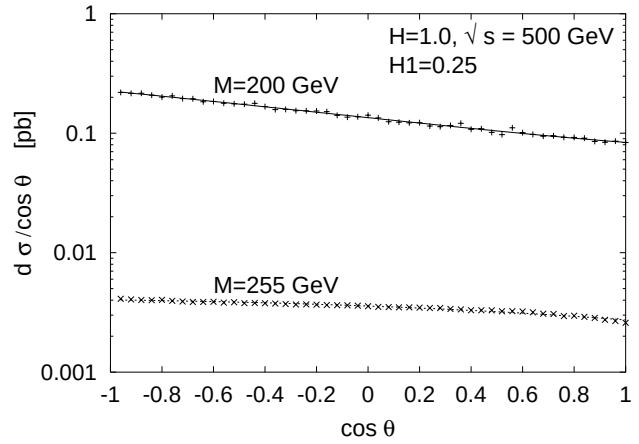

FIG. 4. $d\sigma/d\cos\theta$, where θ is the CMS angle between ℓ_i, ℓ_j. Input values $(H, H1, \sqrt{s}, M)$ are denoted. The points are results of calculations, and the curves are parabolas (in $\cos\theta$) fitted to the points with equal weights. Fluctuations are due to the limited statistics of the Monte Carlo integration.

In large classes of models, in which heavy neutrinos are sequential or have exotic SU(2) × U(1) assignments, $H1$ and H of (3) are severely restricted by available experimental data (LEP and low-energy data) [3,4]: $H1 < 0.016$, $H < 0.122$. In principle, in certain models these restrictions can be avoided. Taking into account the restriction $H1 < 0.016$, the influence of the tu channel is much weaker than for $H1 = 0.25$, but the s-tu interference term still increases the cross section significantly above the pure s-s contributions (cf. Table I). The restriction $H < 0.122$ would imply that the displayed off-shell (finite width) effects in the 2OS region would be weaker than in the $H = 1$ case, and that the displayed σ's in the 1OS and nOS regions would be reduced, approximately by factors H and H^2, respectively. The numbers in Table I were obtained from the general (nOS) program, except in the case of $M = 85$–95 GeV when the 2OS program was used.

To summarize, we calculated cross sections for $e^+e^- \to NN \to W^+\ell^- W^+\ell'^-$, where N's are Majorana neutrinos (mass $M \sim 10^2$ GeV) and ℓ, ℓ' are light leptons e, μ, τ. In contrast to the calculations available so far, we included the effects of the off-shell intermediate N's. These effects were significant even when $\sqrt{s} > 2M$ $(>2M_W)$. They are more pronounced when the $t + u$ channel contributions are significant. The number of reaction events at LEP200 ($\sqrt{s} = 200$ GeV; integrated luminosity ≈ 500 pb^{-1}) would be low (<20, for $M > 85$ GeV) if the strength of the tu channel is restricted by available experiments *and* by confining ourselves to certain classes of models where the heavy neutrinos are either sequential or have exotic SU(2) × U(1) assignments (i.e., when $H1 < 0.016$). Numbers of such events in general significantly increase at linear colliders ($\sqrt{s} = 500$ GeV; integrated luminosity 10^4 pb^{-1}), and may be significant even when $2M > \sqrt{s}$. Further, our approach allows us to calculate numerically various distributions of the final particles.

TABLE I. Values of cross sections σ, for various values of $\sqrt{s}$ and M, and for the tu-strength parameter $H1 = 0$, and in $(\cdots)$ for $H1 = 0.016$. Given are also the relative increases of σ when $H1 = 0 \mapsto 0.016$. The N-decay parameter H is taken $H = 1$; numbers in $[\cdots]$ are for $H = 0.122$.

$\sqrt{s}$ [GeV]	M [GeV]	σ [pb]	Increase
500	200	0.52×10^{-2} (0.77×10^{-2})	49%
500	200	$[0.53 \times 10^{-2}$ $(0.75 \times 10^{-2})]$	[43%]
500	255	0.86×10^{-4} (1.41×10^{-4})	65%
500	255	$[1.14 \times 10^{-5}$ $(1.88 \times 10^{-5})]$	[65%]
300	145	1.33×10^{-3} (1.69×10^{-3})	27%
300	155	0.41×10^{-4} (0.52×10^{-4})	27%
200	85	0.34×10^{-1} (0.40×10^{-1})	17%
200	95	0.72×10^{-2} (0.84×10^{-3})	18%
200	105	1.70×10^{-5} (1.96×10^{-5})	15%

We ignored the questions connected with the experimental difficulties of detecting the discussed process unambiguously. In particular, there are problems connected with identification of the (on-shell) W's and τ's. Further, we ignored the possibility $M < M_W$ [8]—however, additional problems arise in the identification of the process since the two W^+'s are then intermediate off shell.

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