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Flavour and polarisation in heavy neutrino production at e^+e^- colliders

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Abstract

We analyse $\ell W \nu$ production at ILC, paying special attention to the role of the final lepton flavour and beam polarisation in the search for a new heavy neutrino N. We show that a sizeable coupling to the electron $V_{eN} \sim 10^{-2}$ is necessary to have an observable signal in any of the channels, despite the fact that the signal may be more visible in muon or tau final states. The non-observation of a heavy neutrino at ILC will improve the present upper bound on its mixing with the electron by more than one order of magnitude, $V_{eN} \leqslant 0.007$ for m_N between 200 and 400 GeV.

1. Introduction

A 500 GeV e^+e^- International Linear Collider (ILC) offers a clean environment for the study of physics beyond the Standard Model (SM) at a scale of few hundreds of GeV. Its potential is not limited to the study of low energy supersymmetry and precision top quark physics [1]. On the contrary, such a machine is a helpful tool for the investigation of less conventional models, and it might even reveal unexpected new physics.

In this last category might be classified the possible existence of heavy neutrinos with masses of few hundreds of GeV. They are absent in the simplest SM extensions, as long as they do not provide a light neutrino mass generation mechanism,[1] nor an explanation

[1] In principle they can give see-saw type contributions to light neutrino masses [2], but these contributions of the order $V_{eN}^2 m_N \sim 10^7$ eV would be too large compared to the typical neutrino mass size $m_\nu \sim 1$ eV. This means that some symmetry or accidental cancellation is required to reproduce the observed neutrino masses [3].

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of the observed baryon to photon ratio of the universe,[2] $\eta_B = 6.5 \times 10^{-10}$ [6]. However, they appear in grand unified theories, in particular those based on SO(10) and larger groups like E_6 [7], and they may acquire masses much smaller than the unification mass scale [8]. Kaluza–Klein towers of neutrinos are also predicted by models with large extra dimensions and bulk fermions [9], being possible to have the lightest heavy modes near the electroweak scale [10]. Their presence is not experimentally excluded and, if they exist and their mixing with the electron is $O(10^{-2})$, they will be produced at ILC. Conversely, if they are not observed, present bounds on their mixing with the electron will be improved by one order of magnitude.

In e^+e^- annihilation the most favourable process for the observation of heavy neutrino singlets N is $e^+e^- \to N\nu \to \ell W\nu$. We will refer to it as $\ell W\nu$ production from now on. This process has been studied by different authors [11–14]. Here we review previous work including the SM background as well as the effects of initial state radiation (ISR) and beamstrahlung, which have a great impact in the observability of the heavy neutrino, e.g., for $m_N = 300$ GeV they reduce the signal to background ratio by more than a factor of two. In addition, we examine the role of flavour and beam polarisation in the search of heavy neutrinos. As we will show in this Letter, an eNW coupling $O(10^{-2})$ is necessary to produce a heavy neutrino at a detectable level in any of the channels. In this situation, the use of beam polarisations $P_{e^-} = -0.8$, $P_{e^+} = 0.6$ improves the statistical significance of the signal, which may be more visible in final states with $\ell = \mu, \tau$. In case that N does not couple to the electron, the opposite polarisations $P_{e^-} = 0.8$, $P_{e^+} = -0.6$ can be used to enhance the signal and reduce the background. However, they do not suffice to make the signal observable unless very large integrated luminosities are collected. The Dirac or Majorana character of N has a negligible effect on its production cross section and hence does not influence the ILC discovery potential for a non-decoupled heavy neutrino. Thus, we restrict ourselves to the case of a Majorana neutrino.

In the following we fix the notation and review present limits on heavy neutrino masses and mixings, specifying also the SM extensions we will consider. Then we discuss in turn the main contributions to the signal, emphasising the phenomenological implications of the final lepton flavour and beam polarisation. After describing the event generation, we obtain the heavy neutrino discovery limits at ILC. In the conclusions we summarise our results and comment on the expected reach of other future experiments.

2. Heavy neutrino mixing with the light leptons

Let us first review some well-known results [15] to make explicit our hypotheses and notation. We assume that besides the three weak isospin $T_3 = 1/2$ fields ν'_{iL} there are three neutrino singlets N'_{Ri}, $i = 1, 2, 3$. The neutrino mass term is

$$\mathcal{L}_M = -\frac{1}{2} (\bar{\nu}'_L \, \bar{N}'_L) \begin{pmatrix} M_L & \frac{v}{\sqrt{2}} Y \\ \frac{v}{\sqrt{2}} Y^T & M_R \end{pmatrix} \begin{pmatrix} \nu'_R \\ N'_R \end{pmatrix}, \tag{1}$$

where $\nu'_{iR} \equiv (\nu'_{iL})^c$, $N'_{iL} \equiv (N'_{iR})^c$ and Y, M_L, M_R are 3×3 matrices. The 6×6 mass matrix $\mathcal{M}$ can be diagonalised by a unitary matrix, $\mathcal{U}^\dagger \mathcal{M} \mathcal{U}^* = \mathcal{M}_{\text{diag}}$, with the mass eigenstates ν, N related to the weak interaction eigenstates ν', N' by

$$\begin{pmatrix} \nu'_L \\ N'_L \end{pmatrix} = \mathcal{U} \begin{pmatrix} \nu_L \\ N_L \end{pmatrix}, \qquad \begin{pmatrix} \nu'_R \\ N'_R \end{pmatrix} = \mathcal{U}^* \begin{pmatrix} \nu_R \\ N_R \end{pmatrix}. \tag{2}$$

The charged lepton mass matrix can be assumed diagonal without loss of generality. The 6×6 matrix $\mathcal{U}$ can be written as

$$\mathcal{U} = \begin{pmatrix} V^{(\nu)} & V^{(N)} \\ V'^{(\nu)} & V'^{(N)} \end{pmatrix}, \tag{3}$$

with $V^{(\nu)}$ ($V'^{(N)}$) describing the mixing between the light (heavy) neutrinos and $V^{(N)}$, $V'^{(\nu)}$ parameterising the light-heavy neutrino mixing. With this ordering the extended Maki–Nakagawa–Sakata (MNS) matrix [16] $V = (V^{(\nu)} \; V^{(N)})$ parameterises the charged and neutral current gauge interactions

$$\mathcal{L}_{\text{CC}}^W = -\frac{g}{\sqrt{2}} \bar{l}_L \gamma^\mu V \begin{pmatrix} \nu_L \\ N_L \end{pmatrix} W_\mu^- + \text{H.c.}, \tag{4}$$

[2] Neutrino singlets with large Majorana masses can produce an excess of lepton number L which can be converted into the observed baryon asymmetry through $B + L$ violating sphaleron interactions [4]. Nevertheless, the heavy neutrinos which may generate a lepton asymmetry large enough have a small mixing with the light fermions, typically of order $\sqrt{\tilde{m}_\nu / m_N} \lesssim 10^{-8}$, with $\tilde{m}_\nu$ the effective light neutrino mass relevant to the process [5], being then their production rates negligible.

$$\mathcal{L}_{\rm NC}^{Z} = -\frac{g}{2\cos\theta_W}(\bar{\nu}_L \bar{N}_L)\gamma^\mu X \begin{pmatrix} \nu_L \\ N_L \end{pmatrix} Z_\mu, \tag{5}$$

where $X = V^\dagger V$. For $V^{(N)} = 0$ the matrix $V^{(\nu)}$ is the usual 3×3 unitary MNS matrix.

The most stringent constraints on neutrino mixing result from tree-level contributions to processes involving neutrinos as external states like $\pi \to \ell\bar{\nu}$ and $Z \to \nu\bar{\nu}$, and from new one-loop contributions to processes with only external charged leptons like $\mu \to e\gamma$ and $Z \to \ell\bar{\ell}'$ [17–23]. These processes constrain the quantities

$$\Omega_{\ell\ell'} \equiv \delta_{\ell\ell'} - \sum_{i=1}^{3} V_{\ell\nu_i} V^*_{\ell'\nu_i} = \sum_{i=1}^{3} V_{\ell N_i} V^*_{\ell' N_i}, \tag{6}$$

because in the former case we must sum over the external light neutrinos (which are not distinguished) and in the latter the sum is over loop contributions. The first type of processes in particular tests universality. A global fit to experimental data gives [21]

$$\Omega_{ee} \leqslant 0.0054, \qquad \Omega_{\mu\mu} \leqslant 0.0096,$$
$$\Omega_{\tau\tau} \leqslant 0.016 \tag{7}$$

with a 90% confidence level (CL). These limits do not depend on the heavy neutrino masses and are model-independent to a large extent. They imply that heavy neutrino mixing with the known charged leptons is very small, $\sum_i |V_{\ell N_i}|^2 \leqslant 0.0054$, 0.0096, 0.016 for $\ell = e, \mu, \tau$, respectively. The bound on Ω_{ee} also guarantees that neutrinoless double beta decay is within experimental limits for the range of heavy neutrino masses we are interested in (larger than 100 GeV) [24].

The second type of processes, involving flavour-changing neutral currents (FCNC), get new contributions only at the one loop level when the SM is extended only with neutrino singlets, as in our case. These contributions, and hence the bounds, depend on the heavy neutrino masses. In the limit $m_{N_i} \gg M_W$, they imply [22]

$$|\Omega_{e\mu}| \leqslant 0.0001, \qquad |\Omega_{e\tau}| \leqslant 0.01,$$
$$|\Omega_{\mu\tau}| \leqslant 0.01. \tag{8}$$

Except in the case of the first two families, for which experimental constraints on lepton flavour violation are rather stringent, these limits are of a similar size as for the diagonal elements. An important difference, however, is that (partial) cancellations may operate among heavy neutrino contributions. There can be cancellations with other new physics contributions as well. In this work we are interested in determining the ILC discovery potential and the limits on neutrino masses and mixings which could be eventually established. Then, we must allow for the largest possible neutrino mixing and FCNC, although they require cancellations or fine-tuning. Let us examine in more detail the first bound, which is obtained from present limits on the $e\mu\gamma$ and $e\mu Z$ vertices. The dominant terms involving heavy neutrinos are proportional to [22]

$$\sum_{i=1}^{3} V_{eN_i} V^*_{\mu N_i} \phi\big(m^2_{N_i}/M^2_W\big),$$
$$\sum_{i,j=1}^{3} V_{eN_i} X_{N_i N_j} V^*_{\mu N_j} m_{N_i} m_{N_j} f(m_{N_i}, m_{N_j}), \tag{9}$$

respectively, where

$$\phi(x) = \frac{x(1-6x+3x^2+2x^3-6x^2\log x)}{2(1-x)^4}, \tag{10}$$

$$f(x,y) = \frac{xy\log x^2/y^2}{x^2-y^2}. \tag{11}$$

Then, in principle it is possible to find non-vanishing values of the mixing angles V_{eN_i}, $V_{\mu N_i}$ so that the two sums are cancelled. We can distinguish two cases. In the flavour conserving case each heavy neutrino only mixes with one family and both terms are zero because $V_{eN_i} V^*_{\mu N_i} = 0$ for any i and $X_{N_i N_j}$ is proportional to δ_{ij}. In this case $V_{\ell N_i}$ can saturate Eq. (7). If there is flavour violation and the lightest heavy neutrino mixes with the electron and other charged lepton, a mixing large enough to have an observable signal at ILC requires some fine-tuning to cancel the two terms. For instance, if we assume as in Section 5 that $m_{N_1} = 300$ GeV and $V_{eN_1} = 0.052$, $V_{\mu N_1} = 0.069$, $V_{\tau N_1} = 0.126$, we can saturate Eq. (7) and make both sums in Eqs. (9) negligible taking $V_{eN_2} = -0.052$, $V_{eN_3} = 0.004$, $V_{\mu N_2} = 0.062$, $V_{\mu N_3} = 0.031$, $V_{\tau N_2} = V_{\tau N_3} = 0$, for $m_{N_2} = 500$ GeV, $m_{N_3} = 6$ TeV.

Since in general the mixing between charged leptons and heavy neutrinos is very small, the matrix $V^{(\nu)}$ (which corresponds to the first three columns of V) is approximately unitary, up to order $V^2_{\ell N_i}$. Moreover, for the process under consideration the light neutrino masses can be neglected. Therefore, we can assume in the following $V^{(\nu)} \simeq \mathbb{1}_{3\times 3}$ in Eq. (4), and $X_{\nu_\ell \nu_{\ell'}} \simeq \delta_{\ell\ell'}$, $X_{\nu_\ell N_i} \simeq V_{\ell N_i}$ in Eq. (5).

For the calculation of the $\ell W \nu$ cross section including heavy neutrino contributions the total width Γ_N is needed. The partial widths for N weak decays are

$$\begin{aligned}
&\Gamma\big(N \to W^+ \ell^-\big) \\
&\quad = \Gamma\big(N \to W^- \ell^+\big) \\
&\quad = \frac{g^2}{64\pi} |V_{\ell N}|^2 \frac{m_N^3}{M_W^2} \left(1 - \frac{M_W^2}{m_N^2}\right) \\
&\qquad \times \left(1 + \frac{M_W^2}{m_N^2} - 2\frac{M_W^4}{m_N^4}\right), \\
&\Gamma(N \to Z \nu_\ell) \\
&\quad = \frac{g^2}{64\pi \cos^2\theta_W} |V_{\ell N}|^2 \frac{m_N^3}{M_Z^2} \left(1 - \frac{M_Z^2}{m_N^2}\right) \\
&\qquad \times \left(1 + \frac{M_Z^2}{m_N^2} - 2\frac{M_Z^4}{m_N^4}\right).
\end{aligned} \tag{12}$$

We ignore the decays $N \to H\nu_\ell$, which may take place if $m_H < m_N$ [13,25]. Including these decays in Γ_N slightly decreases the $W\ell$ branching ratios and hence the final signal cross sections, which must be multiplied by a factor ranging from $3/4$ (for $m_H \ll m_N$) to 1 (for $m_H \geqslant m_N$). Independently of the values of the couplings, the branching ratio for charged current decays is $2/3$ ($1/2$ if we include scalar decays and $m_H \ll m_N$).

3. General characteristics of the signal

The discovery of a new heavy neutrino in $e^+e^- \to \ell W \nu$ requires its observation as a peak in the invariant ℓW mass distribution, otherwise the irreducible SM background is overwhelming. This requires to reconstruct the W, what justifies to consider $\ell W \nu$ production (instead of general four fermion production), with W decaying hadronically. In the evaluation of $e^+e^- \to \ell^- W^+ \nu$, with $W^+ \to q\bar{q}'$, we will only consider the contributions from the diagrams in Fig. 1, neglecting diagrams with four fermions $e^- q\bar{q}'\nu$ in the final state but with $q\bar{q}'$ not resulting from a W decay. At any rate, we have checked that the corresponding contributions are negligible in the phase space region of interest.

Let us discuss the contributions and sizes of the different diagrams for $e^+e^- \to \ell^- W^+ \nu$ in Fig. 1 and what we can learn from this type of processes at ILC. The first four diagrams are SM contributions. Diagrams 5, 7–9 are present within the SM, mediated by a light neutrino, but they can also involve a heavy one. Diagrams 6 and 10 are exclusive to Majorana neutrino exchange. The SM contribution has a substantial part from resonant W^+W^- production, diagrams 4 and 8, especially for final states with $\ell = \mu, \tau$. The heavy neutrino signal is dominated by diagrams 5 and 6 with N produced on its mass shell, because the Γ_N enhancement of the amplitude partially cancels the mixing angle factor in the decay vertex, yielding the corresponding branching ratio. It must be remarked that the s-channel N production diagram 7 is negligible (few per mille) when compared to the t-, u-channel diagrams 5 and 6. This behaviour is general, because the s-channel propagator is fixed by the large collider energy and suppresses the contribution of this diagram, whereas the t- and u-channel propagators do not have such suppression. Since both diagrams 5 and 6 involve an eNW vertex to produce a heavy neutrino, only in the presence of this interaction the signal is observable. Once the heavy neutrino is produced, it can decay to ℓW with $\ell = e, \mu, \tau$, being the corresponding branching ratios in the ratio $|V_{eN}|^2 : |V_{\mu N}|^2 : |V_{\tau N}|^2$.

In case that the signal is dominated by t and u channel on-shell N production (the only situation in which it is observable), negative electron polarisation and positive positron polarisation increase its statistical significance. For the signal contributions alone we have $\sigma_{e_R^+ e_L^-} : \sigma_{e_L^+ e_R^-} = 1200:1$ (for $m_N = 300$ GeV, $V_{eN} = 0.073$), with $\sigma_{e_R^+ e_R^-} = \sigma_{e_L^+ e_L^-} = 0$. For the SM process, $\sigma_{e_R^+ e_L^-} : \sigma_{e_R^+ e_R^-} : \sigma_{e_L^+ e_R^-} = 150:7:1$, $\sigma_{e_L^+ e_L^-} = 0$. In the limit of perfect beam polarisations $P_{e^-} = -1$, $P_{e^+} = 1$, the signal is enhanced with respect to the background by a factor of 1.05 and, what is more important, the ratio $S/\sqrt{B}$ increases by a factor of two. Using right-handed electrons and left-handed positrons decreases the S/B ratio by a factor of 8, and $S/\sqrt{B}$ by a factor of 50. On the other hand, if the neutrino does not mix with the electron but mixes with the muon or tau, the behaviour is the opposite. Since the only contribution comes from diagram 7 the use of left-handed positrons and right-handed electrons actually increases the signal, while reducing the SM cross

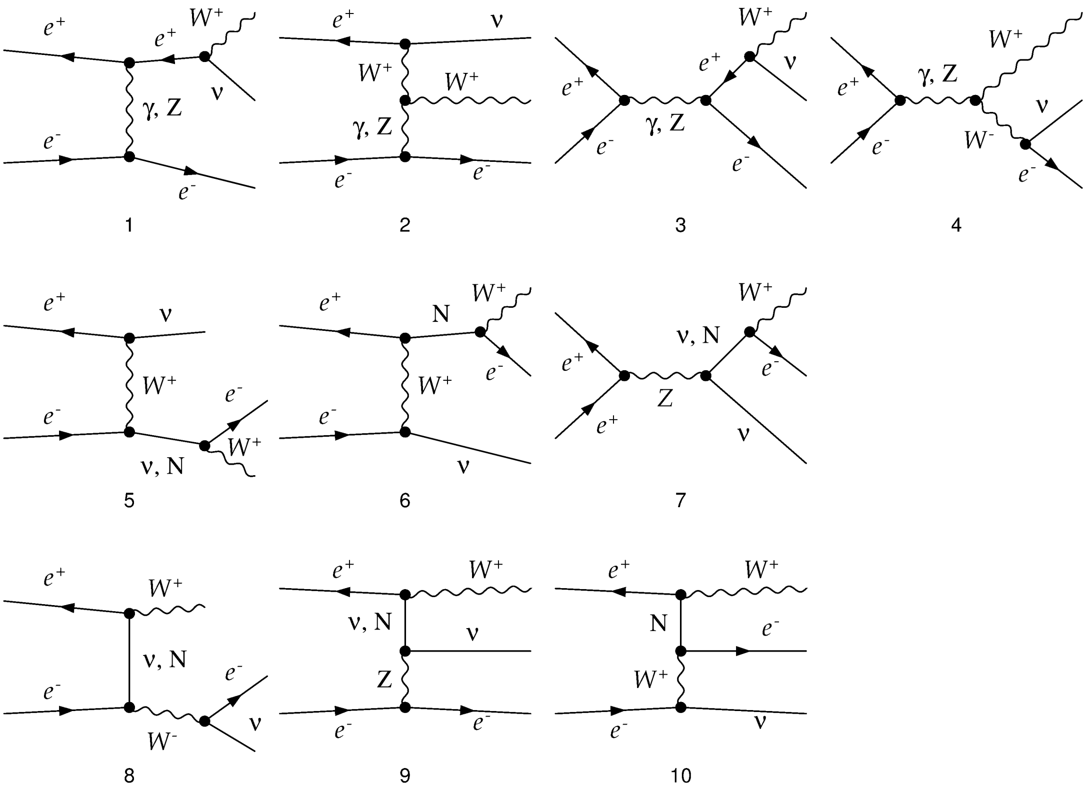

Fig. 1. Diagrams contributing to $e^+e^- \to e^- W^+ \nu$. For $\ell = \mu, \tau$ only diagrams 3–8, 10 contribute.

section for this process [27]. This case is of limited practical interest, since for $V_{eN} = 0$ the signal is barely observable.

We finally point out that the signal cross section exhibits little dependence on the heavy neutrino mass, except close to the kinematical limit [13], and the final results are almost independent of m_N within the range 200–400 GeV [26]. For our calculations we take $m_N = 300$ GeV. In contrast with what has been claimed in the literature [14], we find equal production cross sections for Dirac and Majorana neutrinos to a very good approximation. The reason is easy to understand: while in the present case the signal is strongly dominated by diagrams 5 and 6 (which give equal contributions to the cross section and do not interfere because light neutrino masses can be safely neglected), for a Dirac neutrino only diagram 5 is present. On the other hand, the width of a Dirac neutrino is one half of the width of a Majorana neutrino with the same mixing angles [26].

4. Generation of signals

The matrix elements for $e^+e^- \to \ell^- W^+ \nu \to \ell^- q \bar{q}' \nu$ are calculated using HELAS [28], including all spin correlations and finite width effects. We sum SM and heavy neutrino-mediated diagrams at the amplitude level. The charge conjugate process is included in all our results unless otherwise noted. We assume a CM energy of 500 GeV, with electron polarisation $P_{e^-} = -0.8$ and positron polarisation $P_{e^+} = 0.6$. The luminosity is taken as 345 fb^{-1} per year [29]. In our calculations we take into account the effects of ISR [30] and beamstrahlung [31,32]. For the design luminosity at 500 GeV we use the parameters $\Upsilon = 0.05$, $N = 1.56$ [29]. The actual expressions for ISR and beamstrahlung used in our calculation are collected in Ref. [33]. We also include a beam energy spread of 1%.

In final states with τ leptons, we select τ decays to π, ρ and a_1 mesons (with a combined branching fraction of 55% [34]), in which a single ν_τ is produced, discarding other hadronic and leptonic decays. We simulate the τ decay assuming that the meson and τ momenta are collinear (what is a good approximation for high τ energies) and assigning a random fraction x of the τ momentum to the meson, according to the probability distributions [35]

$$P(x) = 2(1-x) \tag{13}$$

for pions, and

$$P(x) = \frac{2}{2\zeta^3 - 4\zeta^2 + 1}\left[\left(1-2\zeta^2\right) - (1-2\zeta)x\right] \tag{14}$$

for ρ and a_1 mesons, where $\zeta = m^2_{\rho,a_1}/m^2_\tau$. We assume a τ jet tagging efficiency of 50%.

We simulate the calorimeter and tracking resolution of the detector by performing a Gaussian smearing of the energies of electrons, muons and jets, using the specifications in Ref. [36],

$$\frac{\Delta E^e}{E^e} = \frac{10\%}{\sqrt{E^e}} \oplus 1\%, \qquad \frac{\Delta E^\mu}{E^\mu} = 0.02\% E^\mu,$$
$$\frac{\Delta E^j}{E^j} = \frac{50\%}{\sqrt{E^j}} \oplus 4\%, \tag{15}$$

respectively, where the two terms are added in quadrature and the energies are in GeV. We apply kinematical cuts on transverse momenta, $p_T \geqslant 10$ GeV, and pseudorapidities $|\eta| \leqslant 2.5$, the latter corresponding to polar angles $10° \leqslant \theta \leqslant 170°$. To ensure high τ momenta (so that the meson resulting from its decay is effectively collinear) we require $p_T \geqslant 30$ GeV for τ jets. We reject events in which the leptons or jets are not isolated, requiring a "lego-plot" separation $\Delta R = \sqrt{\Delta\eta^2 + \Delta\phi^2} \geqslant 0.4$. For the Monte Carlo integration in 6-body phase space we use `RAMBO` [37].

In final states with electrons and muons the light neutrino momentum p_ν is determined from the missing transverse and longitudinal momentum of the event and the requirement that $p_\nu^2 = 0$ (despite ISR and beamstrahlung, the missing longitudinal momentum approximates with a reasonable accuracy the original neutrino momentum). In final states with τ leptons, the reconstruction is more involved, due to the additional neutrino from the τ decay. We determine the "first" neutrino momentum and the fraction x of the τ momentum retained by the τ jet using the kinematical constraints

$$E_W + E_\nu + \frac{1}{x}E_j = \sqrt{s},$$
$$\vec{p}_W + \vec{p}_\nu + \frac{1}{x}\vec{p}_j = 0,$$
$$p_\nu^2 = 0, \tag{16}$$

in obvious notation. These constraints only hold if ISR and beamstrahlung are ignored, and in the limit of perfect detector resolution. When solving them for the generated Monte Carlo events we sometimes obtain $x > 1$ or $x < 0$. In the first case we arbitrarily set $x = 1$, and in the second case we set $x = 0.55$, which is the average momentum fraction of the τ jets. With the procedure outlined here, the reconstructed τ momentum reproduces with a fair accuracy the original one, while the obtained p_ν is often quite different from its original value.

5. Heavy neutrino discovery at ILC

Following the discussion in Sections 2 and 3 we can distinguish two interesting scenarios for our analysis: (i) the heavy neutrino only mixes with the electron; (ii) it mixes with e and either μ, τ, or both. A third less interesting possibility is that the heavy neutrino does not mix with the electron. We discuss these three cases in turn.

5.1. Mixing only with the electron

A heavy neutrino coupling to the electron yields a peak in the distribution of the ejj invariant mass m_{ejj}, plotted in Fig. 2(a) for $V_{eN} = 0.073$. The solid line corresponds to the SM plus a 300 GeV Majorana neutrino, being the dotted line the SM prediction. The width of the peak is due to energy smearing applied in our Monte Carlo and not to the intrinsic neutrino width $\Gamma_N = 0.14$.

This already striking signal can be enhanced applying a veto cut on the $e\nu$ invariant mass $m_{e\nu}$, shown in Fig. 2(b) for the SM (dashed line) and the new heavy neutrino signal alone (solid line). The two contributions have been separated for clarity. The SM and SM plus heavy neutrino cross sections are collected in Ta-

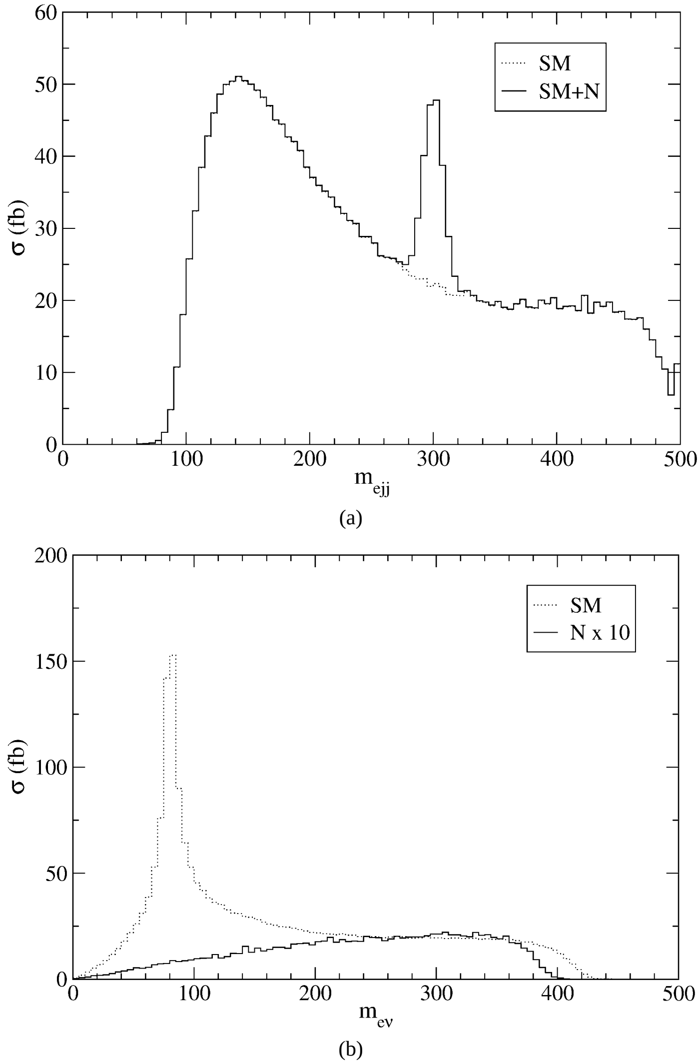

Fig. 2. Kinematical distributions of the ejj invariant mass (a) and the $e\nu$ invariant mass (b).

Table 1
Cross sections (in fb) for $e^+e^- \to e^\mp W^\pm \nu$ before and after the kinematical cuts in Eqs. (17)

	No cuts	m_{ejj}	$m_{e\nu}$	$m_{ejj}, m_{e\nu}$
SM	2253	89.1	1387	53.6
SM + N	2339	173.7	1489	130.8

ble 1, before and after the kinematical cuts

$$290 \leqslant m_{ejj} \leqslant 310 \text{ GeV},$$
$$m_{e\nu} \leqslant 40 \text{ GeV} \quad \text{or} \quad m_{e\nu} \geqslant 110 \text{ GeV}. \tag{17}$$

The new neutrino is said to be discovered when the excess of events (the signal S) in the peak region amounts to more than 5 standard deviations of the number of expected events (the background B), that is, $S/\sqrt{B} \geqslant 5$.[3] This ratio is larger than 5 for $V_{eN} \geqslant 1.2 \times 10^{-2}$, which is the minimum mixing angle for which a 300 GeV neutrino can be discovered. If no signal is found, the limit $V_{eN} \leqslant 6.7 \times 10^{-3}$ can

[3] It must be noted that the SM cross section at the peak can be calculated and normalised using the measurements far from this region.

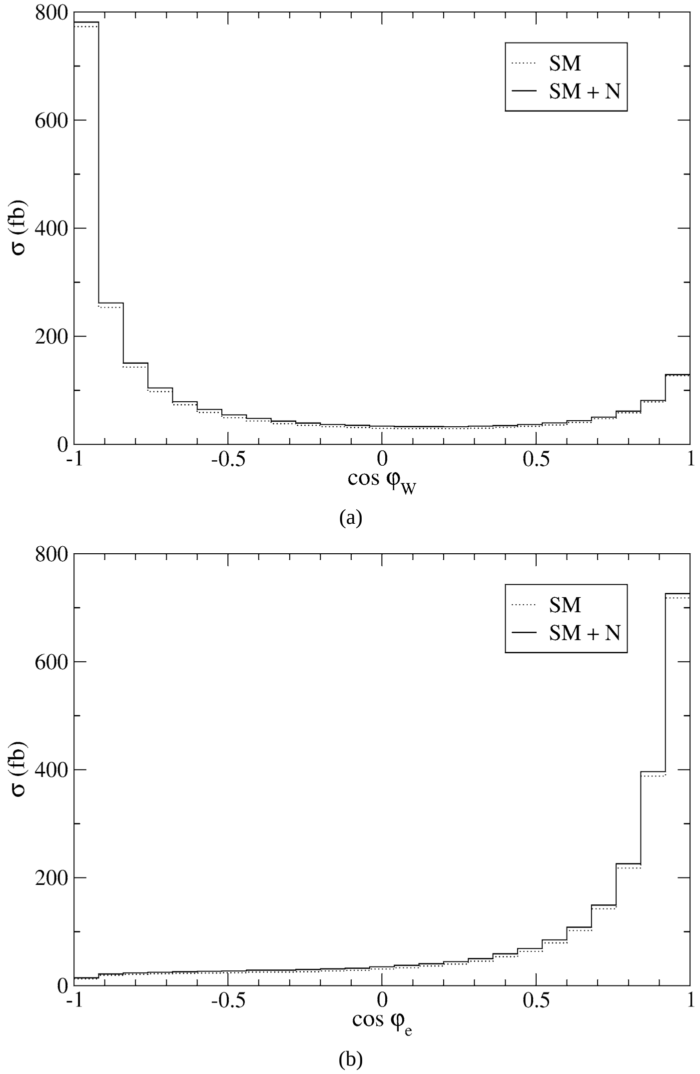

Fig. 3. Dependence of the cross section on the angles φ_W (a) and φ_e (b), for the SM and the SM plus a heavy Majorana neutrino.

be set at 90% confidence level (CL), improving the present limit $V_{eN} \leqslant 0.073$ by a factor of ten.

We have also examined the potential of the angular distributions of the produced $W^\pm$, $e^\mp$ to signal the presence of a heavy neutrino. We define the angle φ_W as the polar angle between the W and the electron (in $e^- W^+ \nu$ final states) or positron (for the charge conjugate process). The angle φ_e is defined analogously. Their kinematical distributions are shown in Fig. 3. From the comparison of these plots with Fig. 2(a) it is apparent that the best kinematical variable to signal the presence of a heavy neutrino is the ejj invariant mass. For these two angular distributions the deviation of the SM prediction amounts to $\chi^2/\text{d.o.f.} \simeq 10\,000/25$, $10\,500/25$, respectively. Dividing the m_{ejj} distribution in 25 bins for better comparison, the corresponding deviation is $\chi^2/\text{d.o.f.} \simeq 64\,000/25$.

The quantitative results obtained here hold for heavy neutrino masses in the range 200–400 GeV to a very good approximation [26]. Although for heavier N the signal cross sections are smaller, the SM background decreases for larger m_{ejj} as well, as can be seen in Fig. 2, and the two effects compensate. For $m_N > 400$ GeV, the cross sections decrease quickly

Table 2
Cross sections (in fb) for $e^+e^- \to \ell^\mp W^\pm \nu$, for $\ell = e, \mu, \tau$, including the kinematical cuts in Eqs. (17)

	e	μ	τ
SM	53.0	31.8	9.7
SM + N	57.4	39.1	13.8

Table 3
$e^+e^- \to \ell^\mp W^\pm \nu$ cross sections (in fb) for a heavy neutrino coupling only to the muon (first column, $\ell = \mu$) or coupling only to the tau (second column, $\ell = \tau$), and the kinematical cuts in Eqs. (17)

	$N - \mu$	$N - \tau$
SM	1.10	0.389
SM + N	1.20	0.414

and thus the limits obtained for the mixing angles are worse.

5.2. Mixing with the three charged leptons

In the most general case that a heavy neutrino mixes simultaneously with the three charged leptons, there may be in principle signals in the e, μ, τ channels. The three of them must be experimentally analysed in the search for a heavy neutrino. As we will show in the following, it is possible that the clearest signals come from the μ or τ channels, even despite the fact that an eNW coupling is necessary to observe them. To prove it we choose the values $V_{eN} = 0.073/\sqrt{2}$, $V_{\mu N} = 0.098/\sqrt{2}$, $V_{\tau N} = 0.13$. These figures are conservative in the sense that the heavy neutrino mainly decays to τ leptons, which are harder to see experimentally, and the cleanest electron and muon channels are relatively suppressed. The cross sections after the kinematical cuts in Eqs. (17) can be found in Table 2 for the SM and the SM plus a heavy neutrino, and for the three modes (in the τ channel we do not apply the veto cut on $m_{\tau\nu}$). After one year of running, the heavy neutrino signal could be seen with 11σ, 24σ, 24σ in the e, μ, τ final states, respectively.

For equal couplings $V_{eN} = V_{\mu N}$ the statistical significance $S/\sqrt{B}$ of the electron and muon signals is similar. Therefore, mixing with the muon does not reduce the sensitivity to V_{eN} [26]. This follows from the fact that for a fixed V_{eN} the N production cross section is independent of $V_{\mu N}$, while the branching ratios in these two channels are in the relation $|V_{eN}|^2 : |V_{\mu N}|^2$. For $V_{\mu N} = 0$ the charged current decays reduce to $N \to eW$, while for $V_{\mu N}$ larger than V_{eN} $N \to \mu W$ dominates, being the combined statistical significance of both channels similar to the one of the electron channel alone for $V_{\mu N} = 0$. The same argument shows that mixing with the tau lowers the sensitivity to V_{eN}, because the observation of the heavy neutrino in the $\tau W \nu$ channel is more difficult.

5.3. Heavy neutrinos not coupling to the electron

We have previously argued that a heavy neutrino signal in the μ or τ channels is observable only if the neutrino also mixes with the electron. We now quantify this statement. We consider a heavy neutrino coupling only to the muon, with $V_{\mu N} = 0.098$, or only to the tau, with $V_{\tau N} = 0.13$. The beam polarisations $P_{e^-} = 0.8$, $P_{e^+} = -0.6$, opposite to the previous ones, are used to enhance the signal and reduce the SM background. The cross sections for the SM and SM plus a heavy neutrino are shown in Table 3 for these two cases. For N mixing only with the muon, the statistical significance of the signal is $S/\sqrt{B} = 1.85$ for one year of running, and 7.3 years would be necessary to observe a 5σ deviation. If the heavy neutrino only mixes with the τ, the statistical significance is only $S/\sqrt{B} = 0.76$, in which case a luminosity 43 times larger is required to achieve a 5σ evidence.

6. Conclusions

Heavy neutrinos with masses near the electroweak scale and large mixing angles ~ 0.1–0.01 with the SM leptons are observable at ILC if they exist. Here we have studied the ILC potential for their detection in the process $e^+e^- \to \ell W \nu$, taking into account the SM background and the effects of ISR and beamstrahlung, paying special attention to the relevance of the final state lepton flavour and initial beam polarisation. Using a parton simulation it has been shown that it is possible to observe a heavy neutrino signal in this final state if it has a mixing with the electron $V_{eN} \gtrsim 10^{-2}$. Although a mixing with the electron of this size is necessary to observe a heavy neutrino at ILC, the signal may be more visible in the muon or tau channel if it also has a relatively large coupling to them. The production cross sections and then the discovery limits do not depend on the Dirac or Majorana nature of the heavy neutrino.

These non-decoupled heavy neutrinos do not explain the observed light neutrino masses nor the baryon asymmetry of the universe. In this sense this search for heavy neutrinos at large lepton colliders is complementary to the joint experimental effort for determining the light neutrino properties, and in particular the neutrino mixing matrix [38]. We could also know about non-decoupled heavy neutrinos if the MNS matrix was found to be non-unitary or CP violation beyond the allowed limits for the minimal SM extension with light Dirac or Majorana masses was measured [39]. (CP violation is unobservable in $e^+e^- \to N\nu \to \ell W\nu$ at ILC because the possible effects and statistics are small.) Other signals of heavy neutrinos like pair production [40] are suppressed by extra powers of the small mixing angles and the center of mass energy threshold. In models with extra matter or interactions further new physics signatures are possible [41]. For example, in left–right symmetric models the new gauge bosons can mediate heavy neutrino single and pair production [42], but we assume that they are too heavy to produce a signal at ILC. At any rate, we are interested in the production of heavy neutrinos having relatively large mixings with the SM fermions.

Finally, other future experiments might exhibit indirect signals of these non-decoupled heavy neutrinos, or limit the possibility of their observation at large colliders. If no deviation from the SM predictions for the invisible Z width is observed in the giga$Z$ option of an e^+e^- collider [1], the bound on their mixing with the light leptons will be reduced by more than one order of magnitude. Analogously, future improvements on the limits on flavour violating processes like $\mu \to e\gamma$ [43] or μ–e conversion [44] by several orders of magnitude would also reduce the corresponding bounds on neutrino mixing by more than one order of magnitude, implying a reduction of the possible signals at ILC or requiring a more delicate fine-tuning. In all cases the eventual improvement on mixing angle constraints is comparable to the ILC potential.

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