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Right-handed currents and heavy neutrinos in high-energy ep and e^+e^- scattering

W. Buchmüller
Deutsches Elektronen Synchrotron DESY, Hamburg, Germany

C. Greub *
Institut für Theoretische Physik, Universität Zürich, Switzerland

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Heavy Dirac or Majorana neutrinos can be produced via right-handed charged currents which occur in extensions of the standard model with $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge symmetry. Low-energy processes, Z precision experiments and direct search experiments in $p\bar{p}$ collisions are consistent with W_R bosons heavier than 450 GeV, if the right-handed neutrinos are heavy. We study the production of heavy neutrinos via right-handed currents in e^+e^- annihilation and ep scattering which appears particularly promising. At HERA heavy neutrinos and W_R bosons can be discovered with masses up to 120 GeV and 700 GeV, respectively.

1. Introduction

Within the standard model electroweak interactions are described by the gauge group $SU(2)_L \times U(1)_Y$. However, despite the extraordinary phenomenological success of the standard model it is not excluded that new gauge interactions will become visible already at TeV energies. Such extended gauge theories have indeed been suggested on various theoretical grounds and particular attention has been given to models with right-handed currents based on the symmetry group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [1] and to models with an additional U(1) symmetry contained in the unified group E_6 [2].

Extended gauge theories also predict fermions in addition to the quarks and leptons of the standard model, since all gauged currents have to be anomaly free. In the minimal case, where the extended group is contained in the unified group SO(10), the addition of one "right-handed" neutrino ν_R for each quark-lepton family suffices to satisfy the requirement of anomaly freedom. The spontaneous breaking of the extended gauge group to the standard model group can induce

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Majorana masses for the right-handed neutrinos, which satisfy an upper bound proportional to the mass of the additional neutral vector boson Z′ [3]. If a Z′ vector boson with mass of order 1 TeV is found, also heavy neutrinos with masses in the range from a few tens of GeV to a few hundred GeV are likely to exist.

In this paper we shall study models with right-handed currents, and we shall determine the mass range of heavy neutrinos and "right-handed" charged vector bosons W_R which can be explored at present and projected ep and e^+e^- colliders. Of particular interest are electron–proton colliders where a single heavy neutrino can be produced without suppression by small mixing angles. This extends previous work [4] where models with an additional U(1) factor have been considered.

Stringent bounds on the W_R mass and the W_R–W_L mixing angle can be derived from various low-energy processes. In particular from the K_L–K_S mass difference one obtains the bound $m_{W_R} > (1\text{–}3)$ TeV for the special case of left–right symmetry [5,6]. However, this bound does not hold for all models with $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge symmetry, but requires further assumptions on fermion mass matrices. In models with arbitrary Yukawa couplings the lower bound is reduced to 300 GeV [7]. As we shall see, constraints from LEP data on the mixing of the neutral vector bosons yield the lower bound of 450 GeV for m_{W_R}, which is close to the bound obtained for massless neutrinos [18]. Finally, the lower mass bound $m_{W_R} > 520$ GeV has been obtained from direct production in proton–proton collisions for neutrinos with masses below 15 GeV [8].

The cross section for the production of heavy neutrinos via right-handed currents in ep scattering at HERA energies is known to be small [9], and because of the stringent lower bound on m_{W_R} in left–right symmetric models [5,6] the original interest [10] in this process at HERA waned. However, the current model-independent lower bound on m_{W_R} from low-energy processes and LEP data is only 450 GeV, and therefore it appears appropriate to investigate in some detail the production of right-handed neutrinos at HERA and also at future ep and e^+e^- colliders.

The paper is organized as follows: in sect. 2 we briefly describe the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model, symmetry breaking and neutrino masses. Sect. 3 deals with mass bounds for neutrinos and vector bosons which follow from low-energy processes and from Z physics. In sect. 4 we compute heavy neutrino production cross sections in ep and e^+e^- scattering and discuss discovery limits for HERA, LEP ⊗ LHC, LEP200 and the 500 GeV e^+e^- LINAC. Sect. 5 contains a summary, and in appendix A we have listed the fully differential cross sections needed for Monte Carlo generators.

2. Models with right-handed currents

In models with the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ left- and right-handed leptons and quarks transform as doublets under $SU(2)_L$ and $SU(2)_R$, respectively:

$$\begin{pmatrix} \nu \\ \ell \end{pmatrix}_{\mathrm{L}} \sim \left(\tfrac{1}{2};\, 0;\, -1\right), \qquad \begin{pmatrix} \nu \\ \ell \end{pmatrix}_{\mathrm{R}} \sim \left(0;\, \tfrac{1}{2};\, -1\right), \tag{2.1}$$

$$\begin{pmatrix} u \\ d \end{pmatrix}_{\mathrm{L}} \sim \left(\tfrac{1}{2};\, 0;\, \tfrac{1}{3}\right), \qquad \begin{pmatrix} u \\ d \end{pmatrix}_{\mathrm{R}} \sim \left(0;\, \tfrac{1}{2};\, \tfrac{1}{3}\right). \tag{2.2}$$

Here we have suppressed the index which labels the different generations of weak eigenstates. The symmetry group requires one right-handed neutrino ν_{R} for each generation. In addition we introduce one gauge singlet n_{L} for each generation,

$$n_{\mathrm{L}} \sim (0;\, 0;\, 0), \tag{2.3}$$

as suggested by the embedding into the unified group E_6 where each quark-lepton generation is contained in a 27-plet.

The minimal Higgs sector, which is required in order to generate masses for the fermions (2.1)–(2.3), contains the following doublet and triplet scalar fields (cf. refs. [1,11]):

$$\Phi \sim \left(\tfrac{1}{2};\, \tfrac{1}{2}^{*};\, 0\right), \qquad \chi \sim \left(0;\, \tfrac{1}{2};\, -1\right), \tag{2.4}$$

$$\Delta \sim (0;\, 1;\, 2) \tag{2.5}$$

The vacuum expectation values of χ and Δ break $\mathrm{SU(2)_L \times SU(2)_R \times U(1)}_{B-L}$ to the standard model group which is then further broken to $\mathrm{U(1)_{em}}$ by the vacuum expectation value of Φ. The covariant derivatives of fermion and scalar fields are given by

$$D_\mu \psi_{\mathrm{L,R}} = \left(\partial_\mu + i g_{\mathrm{L,R}} \frac{\tau^a}{2} W^a_{\mathrm{L,R}\mu} + i \frac{\hat{g}}{2} (B-L) C_\mu \right) \psi_{\mathrm{L,R}}, \tag{2.6}$$

$$D_\mu \Phi = \partial_\mu \Phi + \frac{i}{2} g_{\mathrm{L}} \tau^a W^a_{\mathrm{L}\mu} \Phi - \frac{i}{2} g_{\mathrm{R}} \Phi \tau^a W^a_{\mathrm{R}\mu}, \tag{2.7}$$

$$D_\mu \chi = \left(\partial_\mu + \tfrac{1}{2} i g_{\mathrm{R}} \tau^a W^a_{\mathrm{R}\mu} - \tfrac{1}{2} i \hat{g} C_\mu \right) \chi, \tag{2.8}$$

$$D_\mu \Delta = \left(\partial_\mu + i \hat{g} C_\mu\right) \Delta + \tfrac{1}{2} i g_{\mathrm{R}} [\tau^a, \Delta] W^a_{\mathrm{R}\mu}, \tag{2.9}$$

where

$$\Delta = \tau^3 \Delta^+ + \sqrt{2}\left(\tau^+ \Delta^{++} + \tau^- \Delta^0\right). \tag{2.10}$$

Here $W_{\mathrm{L}\mu}$, $W_{\mathrm{R}\mu}$ and C_μ denote the $\mathrm{SU(2)_L}$, $\mathrm{SU(2)_R}$ and $\mathrm{U(1)}_{B-L}$ vector fields, respectively, and g_{L}, g_{R} and $\hat{g}$ are the corresponding coupling constants. In eq.

(2.10) the electric charges of the three components of Δ are indicated. Similarly, one has for the Higgs fields Φ and χ:

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{pmatrix}, \qquad \chi = \begin{pmatrix} \chi^0 \\ \chi^- \end{pmatrix}. \tag{2.11}$$

The vacuum expectation values

$$v_1' = \langle \chi^0 \rangle_0, \qquad v_2' = \langle \Delta^0 \rangle_0 \tag{2.12}$$

break $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ to the standard model group $SU(2)_L \times U(1)_Y$, which is further broken by

$$v_1 = \langle \phi_1^0 \rangle_0, \qquad v_2 = \langle \phi_2^0 \rangle_0 \tag{2.13}$$

to the group $U(1)_{e.m.}$ of electromagnetic interactions. The vacuum expectation values (2.12) and (2.13) must satisfy $v' \gg v$, since the allowed contribution of right-handed currents to charged current processes is known to be very small.

Given the Higgs fields Δ, χ and Φ the masses of neutrinos and charged leptons are obtained from the following lagrangian

$$\begin{aligned} -\mathscr{L}_M = & \overline{\ell_L} \Phi g_1 \ell_R + \overline{\ell_L} \tilde{\Phi} g_2 \ell_R + \overline{n_L} \chi^\dagger h_1 \ell_R \\ & + \tfrac{1}{2} i \overline{\ell_R^C} \tau_2 \Delta h_2 \ell_R + \tfrac{1}{2} \overline{n_L^C} m n_L + \text{h.c.}, \end{aligned} \tag{2.14}$$

where

$$\tilde{\Phi} = \tau_2 \Phi^* \tau_2 . \tag{2.15}$$

Here $g_{1,2}$, $h_{1,2}$ and m are 3×3 complex matrices in generation space. The lagrangian (2.14) does not conserve lepton number. Hence, after spontaneous symmetry breaking, one will in general obtain nine Majorana neutrinos as mass eigenstates. In the following we shall restrict ourselves to the case where the additional neutrinos are heavy, since the smallness of the light neutrino masses is then naturally understood without requiring extremely tiny Yukawa couplings. Specifically, we shall assume $m_N > 15$ GeV, so that W_R masses below 520 GeV are still allowed by the present CDF bound [8].

For simplicity, let us now consider two special cases. If $h_2 = 0$ and $m = 0$, lepton number is conserved and the neutrino mass terms read

$$-\mathscr{L}_M = \overline{\nu_L} m_D \nu_R + \overline{n_L} m_1 \nu_R + \text{h.c.}, \tag{2.16}$$

where

$$m_{\mathrm{D}} = g_1 v_1 + g_2 v_2^*, \qquad m_1 = h_1 v_1'^*. \tag{2.17}$$

Since $v' \gg v$, one naturally has $m_1 \gg m_{\mathrm{D}}$. One then obtains as mass eigenstates three Dirac neutrinos,

$$-\mathscr{L}_{\mathrm{M}} = \overline{n_{\mathrm{L}}'} m_{\mathrm{N}} \nu_{\mathrm{R}} + \text{h.c.}, \tag{2.18}$$

$$m_{\mathrm{N}} = m_1 + \mathrm{O}(1/m_1), \tag{2.19}$$

$$\overline{n_{\mathrm{L}}'} = \overline{n_{\mathrm{L}}} + \overline{\nu_{\mathrm{L}}} m_{\mathrm{D}} \frac{1}{m_1} + \mathrm{O}(1/m_1^2), \tag{2.20}$$

and three massless Weyl neutrinos, $\nu_{\mathrm{L}}' = \nu_{\mathrm{L}} + \mathrm{O}(1/m_1)$, which are identified with ν_{e}, ν_μ, ν_τ.

If $h_2 \neq 0$ lepton number is broken. The simplest case of this type corresponds to $h_1 = 0$, where n_{L} represents three Majorana neutrinos which decouple from ν_{L} and ν_{R}. The remaining mass terms are (cf. eq. (2.17))

$$-\mathscr{L}_{\mathrm{M}} = \overline{\nu_{\mathrm{L}}} m_{\mathrm{D}} \nu_{\mathrm{R}} + \tfrac{1}{2} \overline{\nu_{\mathrm{R}}^C} m_2 \nu_{\mathrm{R}} + \text{h.c.}, \tag{2.21}$$

where

$$m_2 = \sqrt{2}\, h_2 v_2'. \tag{2.22}$$

This leads, via the see–saw mechanism [12], to three heavy and three light Majorana neutrinos with mass matrices

$$m_{\mathrm{N}} = m_2 + \mathrm{O}(1/m_2), \tag{2.23}$$

$$m_\nu = -m_{\mathrm{D}} \frac{1}{m_2} m_{\mathrm{D}}^T + \mathrm{O}(1/m_2^2). \tag{2.24}$$

The light Majorana neutrinos correspond to ν_{e}, ν_μ and ν_τ. If v', the scale of $\mathrm{SU(2)_R}$ breaking, is of order 1 TeV, their masses are expected to be close to the present experimental upper bounds.

The extended gauge symmetry leads to new charged and neutral current interactions which give the dominant contributions to the production of heavy neutrinos in ep and e^+e^- scattering. Ignoring the small mixings with standard model gauge bosons which are proportional to $(v/v')^2$, one obtains

$$\mathscr{L}_{\mathrm{I}} = J_{\mathrm{R}}^\mu W_{\mathrm{R}\mu}^- + J_{\mathrm{R}}^{\dagger\mu} W_{\mathrm{R}\mu}^+ + J'^\mu Z_\mu', \tag{2.25}$$

where

$$J_{\mathrm{R}}^{\mu}=\frac{g_{\mathrm{R}}}{\sqrt{2}}\left(\bar{d}U^{\mathrm{R}}\gamma^{\mu}\frac{1+\gamma_5}{2}u+\bar{e}V^{\mathrm{R}}\gamma^{\mu}\frac{1+\gamma_5}{2}\nu\right), \tag{2.26}$$

$$J_{\mathrm{R}}^{\prime\mu}=g'\sum_{\psi}\bar{\psi}\left(\cot\alpha Y-\frac{1}{\sin 2\alpha}(B-L)\right)\psi, \tag{2.27}$$

$$Y=T_{3\mathrm{R}}+\frac{B-L}{2}. \tag{2.28}$$

Here Y is the standard model hypercharge, g_{R} and g' are the $SU(2)_{\mathrm{R}}$ and $U(1)_Y$ gauge couplings, and $\tan\alpha$ is the ratio of the $U(1)_{B-L}$ and $SU(2)_{\mathrm{R}}$ coupling constants. ψ denotes the leptons and quarks listed in eqs. (2.1) and (2.2), $T_{3\mathrm{R}}$ is the third component of the right-handed isospin, U^{R} and V^{R} are the hadronic and leptonic Kobayashi–Maskawa type matrices of the right-handed charged current. In the special case, where the gauge couplings g_{L} and g_{R} of the gauge groups $SU(2)_{\mathrm{L}}$ and $SU(2)_{\mathrm{R}}$ are equal, the neutral current reads

$$J_{\mathrm{NC}}^{\prime\mu}=g'\sum_{\psi}\left(\frac{\sqrt{\cos 2\theta_{\mathrm{W}}}}{\sin\theta_{\mathrm{W}}}\bar{\psi}\gamma^{\mu}\frac{1+\gamma_5}{2}\frac{\tau^3}{2}\psi-\frac{\sin\theta_{\mathrm{W}}}{\sqrt{\cos 2\theta_{\mathrm{W}}}}\bar{\psi}\gamma^{\mu}\frac{B-L}{2}\psi\right), \tag{2.29}$$

where θ_{W} is the weak angle. In our discussion on heavy neutrino production in e^+e^- scattering we will restrict ourselves to the case $g_{\mathrm{L}}=g_{\mathrm{R}}$.

3. Bounds on the mass of W_{R}

The subject of this paper is the production of heavy neutrinos in ep and e^+e^- collisions via right-handed charged and neutral currents. The production cross sections are strongly dependent on the W_{R} and Z' masses, which are constrained by low-energy processes and by precision measurements of Z boson properties.

Bounds on the W_{R} mass and the W_{L}–W_{R} mixing angle ζ have been studied in great detail by Langacker and Uma Sankar [7] and, in connection with *CP* violation in the B-system, by London and Wyler [13]. In the case of left–right symmetry, where the hadronic KM mixing matrices U^{L} and U^{R} for left- and right-handed currents are identical, the stringent bound $m_{W_{\mathrm{R}}}>1$–3 TeV has been obtained [5,6]. Of importance are also B_d–$\bar{B}_d$ mixing and the semileptonic branching ratio of b-decays. However, if one relaxes the strong assumption $U^{\mathrm{L}}=U^{\mathrm{R}}$, all processes are compatible with $m_{W_{\mathrm{R}}}>300$ GeV [7].

For heavy Majorana neutrinos an important bound follows from the seach for neutrinoless double-beta decay ($\beta\beta_{0\nu}$) of ^{76}Ge [14]. From the analysis of Langacker and Uma Sankar [7] we obtain

$$\left(\sum_i V_{ei}^{R^2}\frac{m_L}{m_{N_i}}\right)\left(\frac{m_L}{m_R}\right)^4 < 8.9\times 10^{-6}\left(\frac{g_L}{g_R}\right)^4 |U_{ud}^R|^{-2}\left(10^{24}\text{yr}/\tau_{1/2}\right)^{1/2}, \quad (3.1)$$

where $m_{L,R}$ are the $W_{L,R}$ masses. We have included the sum of all intermediate heavy Majorana neutrinos N_i. The V_{ei}^R are the corresponding leptonic KM matrix elements and 10^{24} yr is the current experimental bound on the $\beta\beta_{0\nu}$ lifetime. Assuming $V_{ei}^R \sim \delta_{1i}$, $m_{N_1} \sim m_R$, $g_L = g_R$ and $|U_{ud}^R| \sim 1$ one obtains the rather large lower bound $m_R > 800$ GeV. However, there can also be cancellations among the different contributions from heavy neutrinos and, hence, no model independent bound on m_R can be derived from eq. (3.1).

For the simplest Higgs sector, described in the previous chapter, bounds on W_L–W_R and Z–Z′ mixings also imply bounds on the W_R and Z′ masses. The mass matrices for charged and neutral vector bosons read (cf. eqs. (2.11), (2.12)):

$$M_W^2 = \begin{pmatrix} \frac{1}{2}g_L^2 v^2 & -g_L g_R v_1 v_2^* \\ -g_L g_R v_1^* v_2 & \frac{1}{2}g_R^2\left(v^2 + |v_1'|^2 + 2v_2'^2\right) \end{pmatrix}, \quad (3.2)$$

$$M_Z^2 = \frac{1}{2}\begin{pmatrix} g_L^2 v^2 & -g_L g_R v^2 & 0 \\ -g_L g_R v^2 & g_R^2(v^2 + v'^2) & -g_R^2 \tan\alpha v'^2 \\ 0 & -g_R^2 \tan\alpha v'^2 & g_R^2 \tan^2\alpha v'^2 \end{pmatrix}, \quad (3.3)$$

$$v^2 = |v_1|^2 + |v_2|^2, \qquad v'^2 = |v_1'|^2 + 4v_2'^2. \quad (3.4)$$

Here tan α is the ratio of the $U(1)_{B-L}$ and $SU(2)_R$ gauge coupling constants. For $v = 0$, diagonalization of the W_R^3–C submatrix of (3.3) yields one massless vector boson which couples to the standard model hypercharge current and the heavy vector boson Z′ which couples to the neutral current J_{NC}' (cf. eq. (2.27)). The mixing ζ_Z between Z′ and the standard model neutral vector boson Z, and the mixing ζ_W between W_L and W_R are of order (v^2/v'^2), and therefore small. One easily finds

$$|\zeta_Z| \approx \frac{\left(g_L^2 + g_R^2 \sin^2\alpha\right)^{1/2}\cos^3\alpha}{g_R}\frac{v^2}{v'^2}, \quad (3.5)$$

$$|\zeta_W| \approx 2\frac{g_L}{g_R}\frac{|v_1 v_2|}{|v_1'|^2 + 2v_2'^2}. \quad (3.6)$$

The corresponding vector boson masses are

$$m_L^2 \approx \tfrac{1}{2} g_L^2 v^2, \qquad m_R^2 \approx \tfrac{1}{2} g_R^2 (|v_1'|^2 + 2v_2'^2), \tag{3.7}$$

$$m_Z^2 \approx \tfrac{1}{2}(g_L^2 + g_R^2 \sin^2\alpha) v^2, \qquad m_{Z'}^2 \approx \frac{1}{2} \frac{g_R^2}{\cos^2\alpha} (|v_1'|^2 + 4v_2'^2), \tag{3.8}$$

where we have neglected terms of relative order v^2/v'^2.

In the special case $g_L = g_R$ one has $\cos^2\alpha = 1 - \tan^2\theta_W$ which, together with eqs. (3.7) and (3.8) yields the inequalities

$$\tfrac{1}{2}(1 - \tan^2\theta_W) \leqslant \frac{m_R^2}{m_{Z'}^2} \leqslant 1 - \tan^2\theta_W. \tag{3.9}$$

The Z–Z′ mixing angle reads in this case

$$|\zeta_Z| = \sqrt{\cos 2\theta_W} \left(\frac{m_Z}{m_{Z'}} \right)^2. \tag{3.10}$$

From an analysis of recent LEP data the lower bound $m_{Z'} > 800$ GeV has been derived [15,16] which, together with the inequalities (3.9) implies

$$m_R > 450 \text{ GeV}. \tag{3.11}$$

Hence, LEP data yield a more stringent lower bound on the W_R mass than low-energy processes.

4. Production of heavy neutrinos

Production and decay of heavy Dirac or Majorana neutrinos in ep scattering proceeds through the charged current processes shown in fig. 1. The production cross section is essentially identical to the one for heavy Majorana neutrino

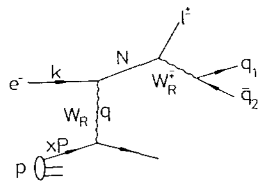

Fig. 1. Production and decay of heavy neutrinos in ep scattering.

production via W_L exchange which has been studied in detail in ref. [4]. For $g_L = g_R$, one easily obtains

$$\frac{d\sigma}{dx\,dy} = \frac{G_F^2}{2\pi}\frac{m_L^4}{\left(y\hat{s}+m_R^2\right)^2}\Big[\left(\hat{s}-m_N^2\right)\left(u(x,\mu^2)+c(x,\mu^2)\right)$$
$$+(1-y)\left(\hat{s}(1-y)-m_N^2\right)\left(\bar{d}(x,\mu^2)+\bar{s}(x,\mu^2)\right)\Big], \qquad (4.1)$$

where x, y and $\hat{s}=xs$ are the usual kinematical variables (cf. fig. 1)

$$s=(P+k)^2, \qquad Q^2=-q^2, \qquad x=\frac{Q^2}{2P\cdot q}, \qquad y=\frac{P\cdot q}{P\cdot k}. \qquad (4.2)$$

which are restricted to the intervals

$$\frac{m_N^2}{s}\leqslant x\leqslant 1, \qquad 0\leqslant y\leqslant 1-\frac{m_N^2}{xs}. \qquad (4.3)$$

u, c, d and s are the densities of up, charm, down and strange quarks in the proton, which depend on the renormalization scale μ. Note, that the cross section is not suppressed by small mixing angles but rather by the W_R mass which enters the propagator and which has to be larger than 450 GeV (cf. eq. (3.11)).

The total cross section is easily obtained after performing numerical integration over x and y. The result is shown in figs. 2–4 as function of m_N for different W_R masses and for three different center-of-mass energies which correspond to HERA,

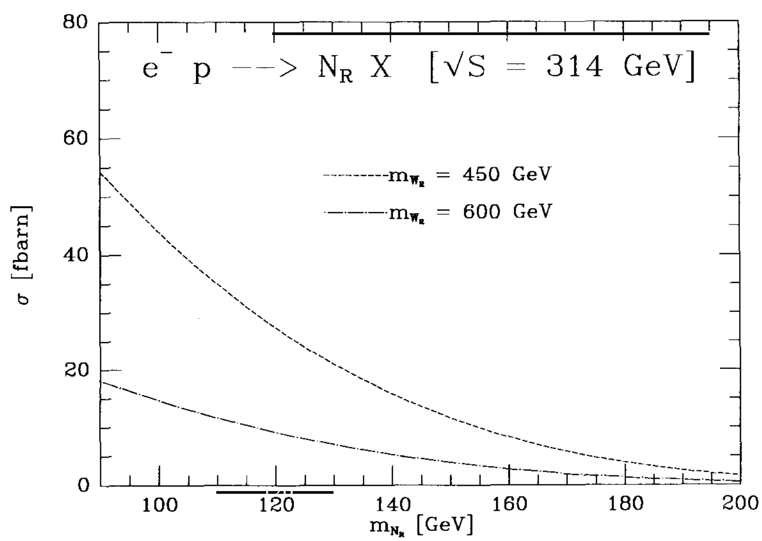

Fig. 2. Total cross section for ep → NX at HERA for different values of the W_R mass.

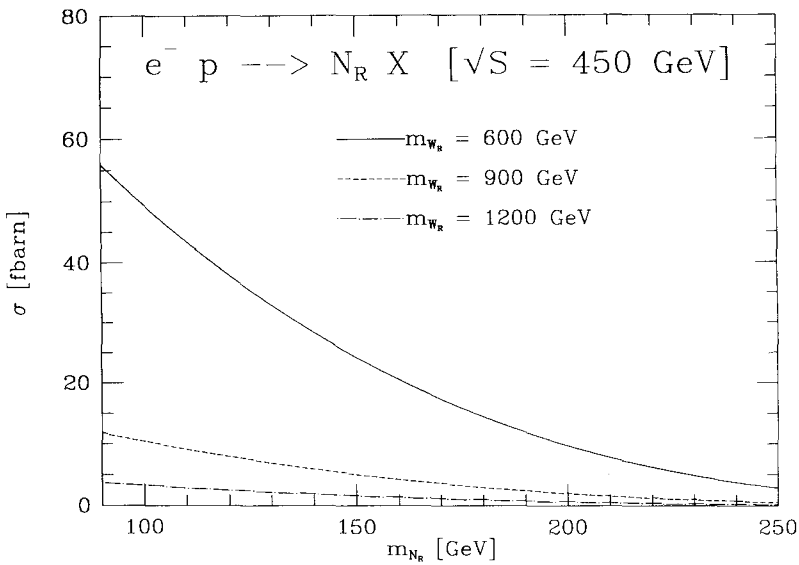

Fig. 3. Total cross section for ep → NX at a HERA upgrade (see text) for different values of the W_R mass.

($\sqrt{s}$ = 314 GeV), an upgraded version of HERA, ($\sqrt{s}$ = 450 GeV) and LEP ⊗ LHC ($\sqrt{s}$ = 1300 GeV). We have used set 1 of Duke–Owens densities [17] with scales corresponding to the average transverse momentum of the produced heavy neu-

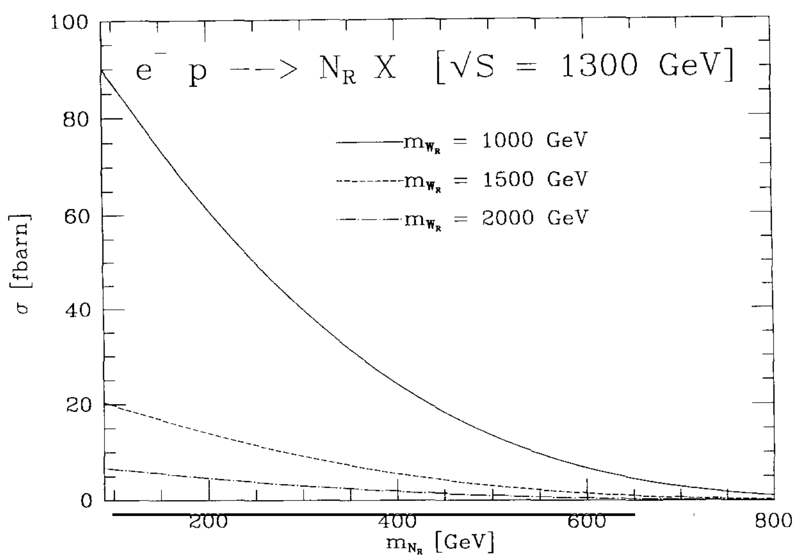

Fig. 4. Total cross section for ep → NX at LEP ⊗ LHC for different values of the W_R mass.

TABLE 1
Discovery limits for W_R masses (m_R) and heavy neutrino masses (m_N) for three ep colliders with different center-of-mass energies and integrated luminosities per year. 5 events are required

	HERA	HERA upgrade	LEP⊗LHC
$\sqrt{s}$ [GeV]	314	450	1300
L [pb^{-1}/y]	200	4000	2000
m_R [GeV]		m_N [GeV]	
450	120	–	–
600	80	270	–
900	–	220	750
1200	–	170	640
1500	–	–	540
2000	–	–	350

trino, i.e. $\mu^2 = 10^3$ GeV^2 ($\sqrt{s} = 314$ GeV), $\mu^2 = 2 \times 10^3$ GeV^2 ($\sqrt{s} = 450$ GeV), $\mu^2 = 5 \times 10^3$ GeV^2 ($\sqrt{s} = 1300$ GeV). At the three machines one year of running is expected to yield the integrated luminosities 200 pb^{-1}, 4000 pb^{-1} and 2000 pb^{-1}, respectively. A rough estimate of the discovery limits for neutrino and W_R masses is obtained by requiring five events. The corresponding values of m_N and m_R are listed in table 1 for the three different machines.

A thorough study of the discovery limits has to take the decays of the heavy neutrinos into account. For $m_N < m_R$ these will be three-body decays into two jets and one charged lepton which, for Dirac neutrinos, always carries negative charge, whereas for Majorana neutrinos positively and negatively charged leptons occur with equal probability. The average transverse momentum of the charged lepton will be smaller than in the case of a two-body decay into charged lepton and W_L boson, which was considered in ref. [4]. However, the signature should still be sufficiently spectacular to allow for a clear separation from background. A detailed study of the final states is most efficiently carried out by means of a Monte Carlo event generator. In the appendix we have listed the necessary formulae, i.e. the

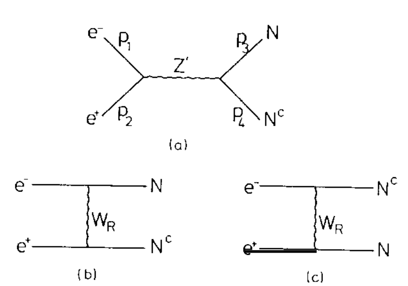

Fig. 5. Pair production of heavy Dirac (a, b) or Majorana neutrinos (a, b, c) in e^+e^- annihilation.

fully differential cross section for the production of a heavy neutrino followed by its three-body decay.

In e^+e^- scattering heavy Majorana neutrinos can be pair produced via the charged and neutral current processes depicted in fig. 5. We consider the case $g_L = g_R$ for which the relevant couplings are given by eqs. (2.25)–(2.28). From the different s-, t- and u-channel contributions one obtains the differential cross section:

$$\frac{d\sigma}{dt} = \frac{G_F^2 m_L^4}{4\pi s^2}\left\{\left[\left(\frac{1}{t-m_R^2} - \frac{\eta}{s-m_{Z'}^2}\right)^2 + \frac{\lambda^2}{\left(s-m_{Z'}^2\right)^2}\right]\left(\left(u-m_N^2\right)^2 - s m_N^2\right)\right.$$
$$+\left[\left(\frac{1}{u-m_R^2} - \frac{\eta}{s-m_{Z'}^2}\right)^2 + \frac{\lambda^2}{\left(s-m_{Z'}^2\right)^2}\right]\left(\left(t-m_N^2\right)^2 - s m_N^2\right)$$
$$\left.+\left(\frac{1}{u-m_R^2} - \frac{1}{t-m_R^2}\right)^2 m_N^2 s\right\}, \tag{4.4}$$

where s, t and u are the kinematical variables (cf. fig. 5)

$$s = (p_1 + p_2)^2, \qquad t = (p_1 - p_4)^2, \qquad u = (p_2 - p_4)^2 \tag{4.5}$$

$$\lambda = \frac{\sin^2\theta_W}{2\cos 2\theta_W}, \qquad \eta = \tfrac{1}{2} - \lambda. \tag{4.6}$$

For heavy Dirac neutrinos, and assuming $L_N = L_{e^-}$, only the processes shown in fig. 5a, 5b contribute to the production cross section, and one obtains (cf. eq. (4.6))

$$\frac{d\sigma}{dt} = \frac{G_F^2 m_L^4}{2\pi s^2}\left[\left(\frac{1}{u-m_R^2} - \frac{\eta}{s-m_{Z'}^2}\right)^2 \left(t-m_N^2\right)^2\right.$$
$$\left.+\frac{\lambda^2}{\left(s-m_{Z'}^2\right)^2}\left(u-m_N^2\right)^2\right]. \tag{4.7}$$

Compared to the Majorana case, eq. (4.4), the t-channel contribution and the corresponding interference terms are missing.

In figs. 6 and 7 the total cross sections are shown as function of the neutrino mass m_N for $\sqrt{s} = 200$ GeV and different W_R masses with $m_{Z'} = 2\ m_R$, for Dirac neutrinos and Majorana neutrinos, respectively. The cross section for Dirac neutrinos is slightly larger than the one for Majorana neutrinos. The difference is

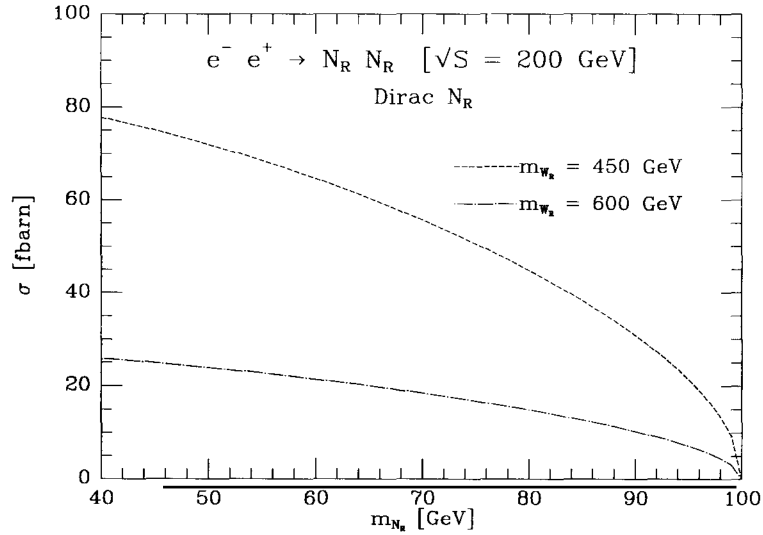

Fig. 6. Total pair production cross section for Dirac neutrinos at LEP200.

particularly significant for neutrino masses close to the kinematic limit $m_N = \frac{1}{2}\sqrt{s}$ due to the well known β^3 factor in the Majorana case. The corresponding cross sections for $\sqrt{s} = 500$ GeV are shown in figs. 8 and 9 for different values of m_N.

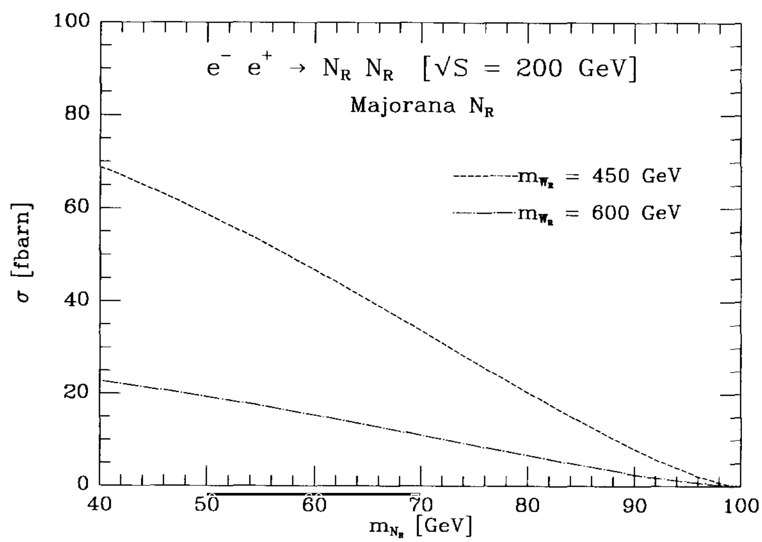

Fig. 7. Total pair production cross section for Majorana neutrinos at LEP200.

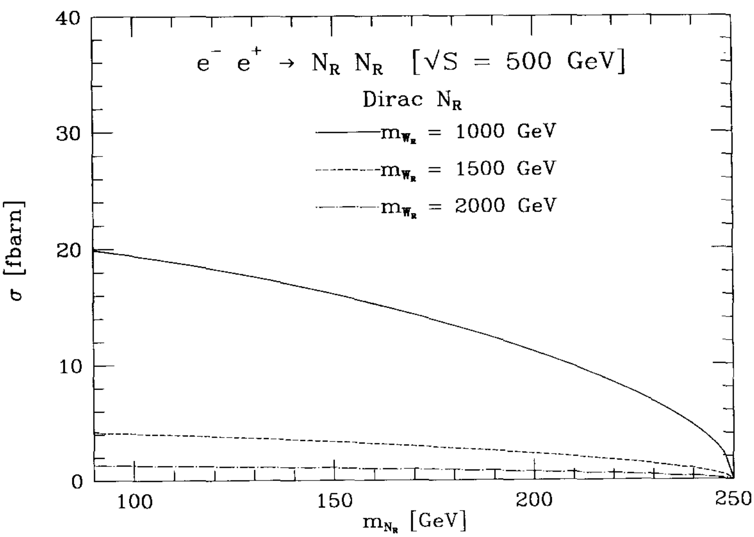

Fig. 8. Total pair production cross section for Dirac neutrinos at NLC.

Estimates of discovery limits for neutrino masses and W_R masses for LEP200 ($\sqrt{s} = 200$ GeV, $L = 500\ \mathrm{pb}^{-1}/\mathrm{y}$) and NLC ($\sqrt{s} = 500$ GeV, $L = 10\ \mathrm{fb}^{-1}/\mathrm{y}$) can be obtained from figs. 6–9 by requiring 10 events. Some representative values are listed in table 2.

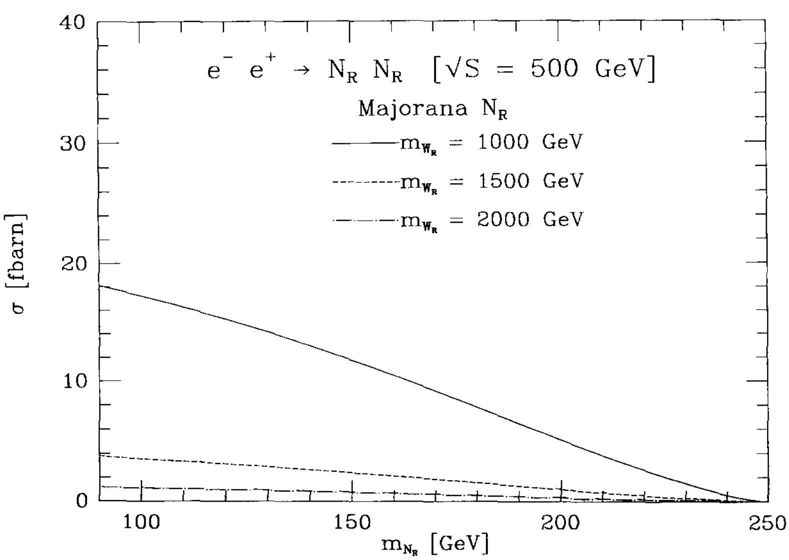

Fig. 9. Total pair production cross section for Majorana neutrinos at NLC.

TABLE 2
Discovery limits for W_R masses (m_R) and heavy Dirac and Majorana neutrino masses (m_N) for two e^+e^- colliders with different center-of-mass energies and integrated luminosities per year. 10 events are required

	LEP200		NLC	
$\sqrt{s}$ [GeV]	200		500	
L [pb^{-1}/y]	500		10000	
m_R [GeV]	m_N [GeV]			
	Dirac	Majorana	Dirac	Majorana
450	90	80	–	–
600	70	50	–	–
1000	–	–	250	240
1500	–	–	240	200
2000	–	–	170	130

5. Conclusions

New gauge interactions beyong the strong and electroweak forces with gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$ are predicted by all unified theories, and it is conceivable that an extension of the standard model gauge group becomes visible already at TeV energies. In the past particular attention has been given to models with right-handed currents based on the electroweak symmetry group $SU(2)_R \times SU(2)_R \times U(1)_{B-L}$ which also predict right-handed neutrinos.

Low-energy processes require the charged W_R boson to be heavier than 300 GeV. From Z precision experiments one finds the lower bound of 450 GeV for equal gauge couplings g_L and g_R of $SU(2)_L$ and $SU(2)_R$, respectively; in this case the Z–Z′ mixing angle is fixed in terms of the neutral vector boson masses. The direct search for W_R bosons in pp collisions yields the current lower bound of 520 GeV for right-handed neutrinos lighter than 15 GeV.

Our present knowledge about W_R bosons and heavy neutrinos will be significantly improved by current and projected ep and e^+e^- colliders. HERA can search for W_R bosons with masses up to 700 GeV and for heavy neutrinos with masses up to 120 GeV. At LEP ⊗ LHC W_R and ν_R masses of order 1 TeV appear accessible. LEP200 can test for W_R masses up to 600 GeV and the projected 500 GeV e^+e^- linear collider (NLC) should be able to reach 1 TeV for the W_R boson.

In the search for right-handed currents and heavy neutrinos ep colliders have two main advantages compared to e^+e^- colliders: First, in ep scattering only a single heavy neutrino is produced and, consequently, much larger neutrino masses are accessible. Second, in ep collisions a violation of lepton number in the decays of Majorana neutrinos can be discovered by just measuring a positive charge of the final state lepton, whereas in e^+e^- annihilation the angular distribution of the final state leptons has to be studied.

We would like to thank G. Ingelman for helpful discussions on Monte Carlo generators. We are also indebted to J. Polak and M. Zrałek for pointing out a numerical error in the original version of the paper.

Appendix A

THE FULLY DIFFERENTIAL CROSS SECTIONS

In this appendix we give formulae for the fully differential cross sections for the reactions

$$e^- p \to X_h N \to X_h \ell^- W_R^+ \to X_h \ell^- q_1 \bar{q}_2, \tag{A.1}$$

$$e^- p \to X_h N \to X_h \ell^+ W_R^- \to X_h \ell^+ q_1 \bar{q}_2. \tag{A.2}$$

In eqs. (A.1) and (A.2) X_h denotes hadronic matter stemming from the proton remnants and the scattered quark (or antiquark). It turns out that the differential cross sections, when expressed in terms of variables defined in the centre-of-mass frame of the heavy neutrino, and also the kinematical ranges of these variables have a relatively simple form. For both processes we use the following variables: x, y – deep-inelastic variables defined in eq. (4.2); E_ℓ – energy of ℓ in rest-frame of heavy neutrino; $E_{\bar{q}}$ – energy of $\bar{q}_2$ in rest-frame of heavy neutrino; θ_ℓ, $\theta_{\bar{q}}$ – polar angles of ℓ and $\bar{q}_2$ in rest-frame of heavy neutrino with respect to proton direction; $\phi_{\bar{q}}$ – azimuthal angle of $\bar{q}_2$ in rest-frame of heavy neutrino.

For given values of E_ℓ, $E_{\bar{q}}$, θ_ℓ, $\theta_{\bar{q}}$, $\phi_{\bar{q}}$ the azimutal angle ϕ_ℓ of ℓ in the rest-frame of the heavy neutrino N is then fixed to be either

$$\phi_\ell = \phi_{\bar{q}} + \arccos C \qquad \text{or} \qquad \phi_\ell = \phi_{\bar{q}} + 2\pi - \arccos C, \tag{A.3}$$

with

$$C = \frac{m_N^2 - 2m_N(E_\ell + E_{\bar{q}}) + 2E_\ell E_{\bar{q}}(1 - \cos\theta_\ell \cos\theta_{\bar{q}})}{2E_\ell E_{\bar{q}} \sin\theta_\ell \sin\theta_{\bar{q}}}. \tag{A.4}$$

Hence, by specifying E_ℓ, $E_{\bar{q}}$, θ_ℓ, $\theta_{\bar{q}}$, $\phi_{\bar{q}}$ and by choosing ϕ_ℓ according to (A.3), the four-momenta of ℓ and $\bar{q}_2$ in the rest-frame of N are uniquely fixed, and by energy-momentum conservation also the four-momentum of q_1 is determined.

In order to write down the cross section in a compact form, we also use (as auxiliary variables) the energy (E_N), the momentum (p_N) and the polar angle (θ) of the heavy neutrino in the laboratory frame. x, y and the laboratory beam energies E_e of the electron and E_p of the proton fix longitudinal (p_N^L) and transverse ($p_N^\perp$) parts of the momentum of the heavy neutrino N:

$$(p_N^\perp)^2 = y[(1-y)\hat{s} - m_N^2],$$

$$p_N^L = \frac{\hat{s}y + m_N^2 - 4(1-y)E_e^2}{4E_e}, \qquad \hat{s} = xs = 4xE_pE_e. \tag{A.5}$$

Hence, also θ, p_N and E_N are fixed:

$$p_N = \left[(p_N^L)^2 + (p_N^\perp)^2\right]^{1/2}, \qquad \cos\theta = \frac{p_N^L}{p_N}, \qquad E_N = \sqrt{p_N^2 + m_N^2}. \tag{A.6}$$

We now give the fully differential cross section $d\sigma_-$ for reaction (A.1) with a negatively charged lepton in the final state. Note, that one obtains the same result for Dirac and Majorana neutrinos:

$$\frac{d\sigma_-}{dx\, dy\, dE_\ell\, dE_{\bar{q}}\, d\theta_\ell\, d\theta_{\bar{q}}\, d\phi_{\bar{q}}} = V_- \times \left\{ A_-\left[u(x,\mu^2) + c(x,\mu^2)\right] + \frac{\hat{s}(1-y) - m_N^2}{\hat{s}} B_-\left[\bar{d}(x,\mu^2) + \bar{s}(x,\mu^2)\right] \right\}, \tag{A.7}$$

with

$$V_- = \frac{G_F^4 m_L^8 m_N E_{\bar{q}}(m_N - 2E_{\bar{q}}) \sin\theta_\ell \sin\theta_{\bar{q}}}{16\pi^6 \Gamma_N J(xys + m_R^2)^2 (m_N^2 - 2m_N E_\ell - m_R^2)^2}, \tag{A.8}$$

$$A_- = (\hat{s} - m_N^2) - 2\cos\theta_{\bar{q}}\left[E_e(p_N + E_N\cos\theta) + xE_p(p_N - E_N\cos\theta)\right] - 2m_N(xE_p - E_e)\sin\theta\sin\theta_{\bar{q}}\cos\phi_{\bar{q}}, \tag{A.9}$$

$$B_- = \hat{s}(1-y) - 2xE_p\cos\theta_{\bar{q}}(p_N - E_N\cos\theta) - 2xm_N E_p\sin\theta\sin\theta_{\bar{q}}\cos\phi_{\bar{q}}, \tag{A.10}$$

$$J = \left[\cos(\theta_\ell - \theta_{\bar{q}}) - Z\right]^{1/2}\left[Z - \cos(\theta_\ell + \theta_{\bar{q}})\right]^{1/2}, \tag{A.11}$$

$$Z = \frac{2E_\ell E_{\bar{q}} - 2m_N(E_\ell + E_{\bar{q}}) + m_N^2}{2E_\ell E_{\bar{q}}}. \tag{A.12}$$

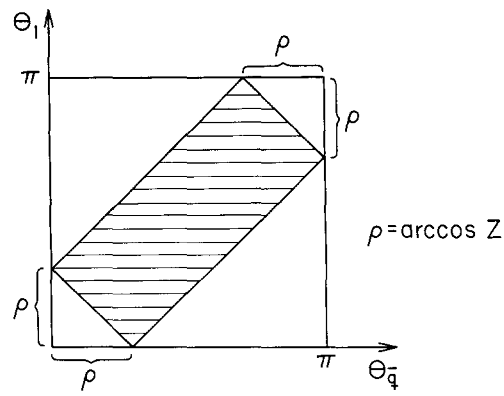

Fig. A.1. The shaded area denotes the allowed range for the variables $\theta_{\bar{q}}$ and θ_ℓ.

For simplicity, in eq. (A.7) we have chosen the Kobayashi–Maskawa type matrices U^R and V^R equal to the unit matrix. The case of nonvanishing mixing angles can be easily implemented. The kinematically allowed ranges of x, y, E_ℓ, $E_{\bar{q}}$ and $\phi_{\bar{q}}$ read:

$$\frac{m_N^2}{s} \leqslant x \leqslant 1, \qquad 0 \leqslant y \leqslant 1 - \frac{m_N^2}{xs}, \tag{A.13}$$

$$E_\ell \in \left[0, \frac{m_N}{2}\right], \qquad E_{\bar{q}} \in \left[\frac{m_N - 2E_\ell}{2}, \frac{m_N}{2}\right], \tag{A.14}$$

$$\phi_{\bar{q}} \in [0, 2\pi]. \tag{A.15}$$

The boundaries of the polar angles θ_ℓ and $\theta_{\bar{q}}$ are given by the inequalities

$$\cos(\theta_\ell + \theta_{\bar{q}}) \leqslant Z \qquad \cos(\theta_\ell - \theta_{\bar{q}}) \geqslant Z, \tag{A.16}$$

with Z given in eq. (A.12). The allowed range in the $(\theta_\ell, \theta_{\bar{q}})$ plane corresponds to the shaded area in fig. A.1.

We now discuss the second process given in eq. (A.2), where the final-state lepton is positively charged. This case only occurs for heavy Majorana neutrinos. The corresponding fully differential cross section $d\sigma_+$ is obtained from eq. (A.7) by replacing V_-, A_- and B_- by V_+, A_+ and B_+, respectively. For these quantities we get:

$$V_+ = \frac{G_F^4 m_L^8 m_N (2E_\ell + 2E_{\bar{q}} - m_N) \sin\theta_\ell \sin\theta_{\bar{q}}}{16\pi^6 \Gamma_N J (xys + m_R^2)^2 (m_N^2 - 2m_N E_\ell - m_R^2)^2}, \tag{A.17}$$

$$A_+ = (\hat{s} - m_N^2)(m_N - E_\ell - E_{\bar{q}})$$
$$- 2(E_\ell \cos\theta_\ell + E_{\bar{q}} \cos\theta_{\bar{q}})[E_e(p_N + E_N \cos\theta) + xE_p(p_N - E_N \cos\theta)]$$
$$- 2m_N \sin\theta(xE_p - E_e)(E_\ell \sin\theta_\ell \cos\phi_\ell + E_{\bar{q}} \sin\theta_{\bar{q}} \cos\phi_{\bar{q}}), \qquad (A.18)$$

$$B_+ = \hat{s}(1-y)(m_N - E_\ell - E_{\bar{q}})$$
$$- 2xE_p(p_N - E_N \cos\theta)(E_\ell \cos\theta_\ell + E_{\bar{q}} \cos\theta_{\bar{q}})$$
$$- 2xm_N E_p \sin\theta(E_\ell \sin\theta_\ell \cos\phi_\ell + E_{\bar{q}} \sin\theta_{\bar{q}} \cos\phi_{\bar{q}}). \qquad (A.19)$$

We now discuss how one has to generate the four-momenta of the final-state particles of an event in the laboratory frame. For a given "point"

$$(x, y, E_\ell, E_{\bar{q}}, \theta_\ell, \theta_{\bar{q}}, \phi_{\bar{q}}) \qquad (A.20)$$

we first choose ϕ_ℓ according to eq. (A.3) with equal probability. Then the four-momenta of the final state particles originating from the heavy neutrino are uniquely fixed in the neutrino rest-frame. In order to get the corresponding four-momenta in the laboratory frame of the ep reaction one has to perform the following Lorentz transformation:

$$v_{LAB} = R_z(\Phi) R_y(\theta) L_z v_{C.M.}, \qquad (A.21)$$

where v generically stands for the four-momentum of ℓ, $\bar{q}_2$ and q_1. The rotations $R_z(\Phi)$, $R_y(\theta)$ and the boost L_z read

$$R_z(\Phi) = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & \cos\Phi & -\sin\Phi & 0 \\ 0 & \sin\Phi & \cos\Phi & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}, \quad R_y(\theta) = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & \cos\theta & 0 & \sin\theta \\ 0 & 0 & 1 & 0 \\ 0 & -\sin\theta & 0 & \cos\theta \end{pmatrix}, \qquad (A.22)$$

$$L_z = \frac{1}{m_N} \begin{pmatrix} E_N & 0 & 0 & p_N \\ 0 & m_N & 0 & 0 \\ 0 & 0 & m_N & 0 \\ p_N & 0 & 0 & E_N \end{pmatrix}. \qquad (A.23)$$

Here θ and Φ are polar and azimuthal angle of the heavy neutrino in the lab-frame, and E_N and p_N are its energy and momentum:

$$p_N^\mu = (E_N, p_N \sin\theta \cos\Phi, p_N \sin\theta \sin\Phi, p_N \cos\theta), \qquad p_N = |\boldsymbol{p}_N|. \qquad (A.24)$$

θ, E_N and p_N are fixed by eqs. (A.5) and (A.6), whereas Φ has to be chosen in the interval $[0, 2\pi]$ with equal probability. The absolute weight of the event is then given by

$$\frac{d\sigma_\pm}{dx\, dy\, dE_\ell\, dE_{\bar{q}}\, d\theta_\ell\, d\theta_{\bar{q}}\, d\phi_{\bar{q}}} \frac{1}{2} \frac{1}{2\pi}, \tag{A.25}$$

where the factors $\frac{1}{2}$ and $1/(2\pi)$ are due to the choices of ϕ_ℓ and Φ, respectively.

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