TESTING LEFT_RIGHT SYMMETRIC
GAUGE THEORIES AT $e^+e^-$ COLLIDERS

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Abstract

If the Standard Model is embedded in a left–right symmetric gauge theory at the
TeV scale, the pair production of light $W$–bosons in $e^+e^-$ collisions, $e^+e^-\rightarrow W^+W^-$, will be affected by mixings in the gauge and neutrino sectors, and by
the $t$–channel exchange of a heavy right–handed neutrino. The modification of the
cross section by these new effects is studied for high–energy $e^+e^-$ colliders.

1. Introduction. The embedding of the Standard Model (SM) in a left–right (LR) symmetric theory [1] at scales of order 1 TeV is a hypothetical but interesting option. Models
based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ share the success of the SM and, in addition, are able to explain the parity violation in the weak interaction in a dynamical
way. They necessarily incorporate right–handed neutrino states which can be connected
with the non-zero neutrino masses via the see-saw mechanism [2]. Moreover, spontaneous
symmetry breaking may be the origin of CP–violation in LR models [3]. Even the high
left–right breaking scales are generally preferred in theoretical analyses[1] a scale of order one to several TeV is still compatible with all direct experimental observations motivating
collider searches.

Consequences of LR extensions of the SM have been studied in many facets of the theory. In this letter we shall focus on the impact of such extensions on the pair production of the light $W$–bosons at $e^+e^-$ colliders:

$$e^+e^-\rightarrow W^+W^-,$$

However, in supersymmetric LR models [4] there is an upper bound on the right-handed scale set by the scale of supersymmetry.
Fig. 1. Diagrams contributing to the process $ee \rightarrow W^+W^-$ in the LR model.

a precursor to the production of the new heavy charged gauge bosons which may require collider energies in the multi-TeV range. This process is affected in LR symmetric theories by three mechanisms. First, the interpretation of the observed mass eigenstates $W^\pm$ as mixed states of $W^\pm_L$ with a small admixture of $W^\pm_R$, as well as the interpretation of the light electron neutrino state $\nu$ as a mixture of the left–handed neutrino state $\nu_L$ with the heavy right–handed neutrino state $\nu_R$. Second, the $t$–channel exchange of the new heavy, predominantly right–handed neutrino $N$, cf. Fig.1. Third, the mixing in the neutral gauge boson sector and the $s$–channel exchange of the heavy neutral gauge boson $Z'$. These three mechanisms modify the total cross section, the angular distributions as well as the $W$ helicities in the final states. Such effects have been searched for at LEP2 [5], and they will be searched for [6] at future $e^+e^-$ linear colliders [7]. If not discovered, new limits may be derived on the LR parameters. This note expands the earlier analysis of Ref.[8] by taking properly into account the mixings in the charged $W$-boson sector. The effects are in general different from the effects induced by the anomalous gauge boson self–interactions. Other LR analyses in $e^+e^-$ collisions focussed on the exchange of heavy $W'^\pm$ states [9] and on the Higgs phenomena [10].

2. The Model. In the minimal $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model each generation of quarks and leptons carry the quantum numbers $Q_L \sim (1/2, 0, 1/3)$, $Q_R \sim (0, 1/2, 1/3)$, $L_L \sim (1/2, 0, -1)$ and $L_R \sim (0, 1/2, -1)$. The right–handed fields are doublets under $SU(2)_R$ and a right–handed neutrino $N_R$ must exist. The minimal Higgs sector consists of a bidoublet $\phi \sim (1/2, 1/2, 0)$ and two triplets $\Delta_L \sim (1, 0, 2)$ and $\Delta_R \sim (0, 1, 2)$. After the spontaneous symmetry breaking, the phenomenological requirement $|v_R| \gg |k_1|, |k_2| \gg |v_L|$ for the vacuum expectation values $v_{L,R}$ and $k_{1,2}$ of the triplet and doublet Higgs fields, ensures the suppression of the right–handed currents and the smallness of the neutrino mass.

The $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge symmetry group implies that the usual left–handed gauge bosons $W^i_L$ ($i = 1, 2, 3$), their right–handed counterparts $W^i_R$ and the $U(1)$ gauge boson $Y$ combine to form the physical charged and neutral gauge bosons and the photon. In general, the strength of the gauge interactions of these bosons is described by the coupling constants $g_L, g_R$ and $g'$, respectively. However, strict LR symmetry $\Psi_L \leftrightarrow \Psi_R, \Delta_L \leftrightarrow \Delta_R, \phi \leftrightarrow \phi^\dagger$ [with $\Psi$ denoting any fermion] leads to the relation $g_L = g_R \equiv g$, which will be assumed throughout this paper.
The weak eigenstates $W_L^\pm$ and $W_R^\pm$ mix in the mass eigenstates $W^\pm$ and $W'^\pm$; assuming CP invariance, the mixing matrix is defined by the angle $\chi_W$:

\[
W^\pm = \cos \chi_W W_L^\pm + \sin \chi_W W_R^\pm \\
W'^\pm = -\sin \chi_W W_L^\pm + \cos \chi_W W_R^\pm
\] (2)

The weak eigenstate $W_L$ can be identified with the pure SM gauge boson. Similarly the neutrino mass eigenstates are mixtures of the weak eigenstates, parametrized by the angle $\chi_N$:

\[
\nu = \cos \chi_N \nu' + \sin \chi_N N' \\
N = -\sin \chi_N \nu' + \cos \chi_N N'
\] (3)

$\nu$ and $N$ are the light and heavy neutrino mass eigenstates, and $\nu' = \nu_L + \nu'_L$ and $N' = \nu_R + \nu'_R$ are the usual self-conjugate spinors; intergenerational mixings are irrelevant for the present analyses.

The charged–current interaction vertices for the left–chiral and the right–chiral currents are given by

\[
\langle \nu_L | W_L^- | e^- \rangle = \frac{g}{2\sqrt{2}} W_L^\mu \overline{\nu_L} \gamma_\mu (1 - \gamma_5) e \\
\langle N_R | W_R^- | e^- \rangle = \frac{g}{2\sqrt{2}} W_R^\mu \overline{N_R} \gamma_\mu (1 + \gamma_5) e
\] (4)

The charged–current interactions for the mixed mass eigenstate $s$ can easily be obtained from these matrix elements.

The neutral gauge bosons in LR models are mixtures of $W_{L,R}^3$ and $Y$. The mixing between the massive neutral gauge bosons relevant for our analyses can be parametrized as

\[
Z = \cos \chi_Z Z_1 + \sin \chi_Z Z_2 \\
Z' = -\sin \chi_Z Z_1 + \cos \chi_Z Z_2
\] (5)

where $Z$ and $Z'$ denote the mass eigenstates, and $Z_1$ and $Z_2$ denote the weak eigenstates of the massive neutral bosons. The field $Z_1$ can be identified as the corresponding SM boson.

The tree-level neutral current Lagrangian for the physical $Z, Z'$ bosons is of the form

\[
\mathcal{L}_{NC} = \frac{g}{2\cos \theta_W} \left[ \overline{f} \gamma_\mu \left( g_V^f - g_A^f \gamma_5 \right) f Z_\mu + \overline{f} \gamma_\mu \left( g_V^f - g_A^f \gamma_5 \right) f Z'_\mu \right]
\] (6)

where
\begin{align*}
g_V^f &= \cos \chi_Z g_V^{0f} + \sin \chi_Z g_V^{if} \\
g_A^f &= \cos \chi_Z g_A^{0f} + \sin \chi_Z g_A^{if}
\end{align*}

and

\begin{align*}
g_V^{0f} &= I_3^f - 2Q^f \sin^2 \theta_W, & g_V^{if} &= 1/\sqrt{\cos 2\theta_W} \cdot g_V^{0f} \\
g_A^{0f} &= I_3^f, & g_A^{if} &= -\sqrt{\cos 2\theta_W} \cdot g_A^{0f}
\end{align*}

Here we have taken into account that in the LRSM always $I_{3R}^f = I_{3L}^f \equiv I_3^f$ for the third components of the L/R isospin for a given fermion flavor $f$. For simplicity, we have expressed the couplings in terms of $I_3^f$ and the electric charge $Q_f$.

Elaborate analyses of high-precision data have been presented in the past years on the LR gauge sector [11,12], constraining the $Z'$ mass to values above $\mathcal{O}(1)$ TeV scale and the mixing among the neutral gauge bosons to values below $\mathcal{O}(10^{-4})$. In the present context these effects in the neutral currents are expected to be subleading compared with the possible effects in the charged current interactions.

The lower bound on the $W'$ mass derived from the $K_L-K_S$ mass difference is quite stringent [13], $M_{W'} \gtrsim 1.6$ TeV (being, however, subject to uncertainties from low energy QCD in the kaon system); the bound on the mixing angle $\chi_W$ is as low as $\chi_W \lesssim 0.013$ [14]. The direct searches for $W'$ at the Tevatron yield bounds $M_{W'} \gtrsim 720$ GeV assuming a light keV–range $N$, and $M_{W'} \gtrsim 650$ GeV assuming $M_N < M_{W'}/2$ [15]. These bounds are weakened considerably for more general LR models [16].

The least tested components of the LR model are the masses and mixings of neutrinos. Analyses of the precision data that constrain fermion mixings [17] have given a 90% CL bound $|\chi_N| \lesssim 0.081$ for the electron neutrinos. As no new particles have been discovered at LEP2 it is plausible to assume that the mass of the heavy R–type neutrino may exceed about 100 GeV.

The mixing angle $\chi_W$ is, even in the simplest models of the Higgs representations, independent of the neutral current parameters. Also the prediction on the mixing angle $\chi_N$ depends strongly on the scenario in which the neutrino mixings and masses are generated [18], allowing mixings still as large as $\chi_N \sim 0.1$.

3. Pair Production of $W$-Bosons. In the LR symmetric theory the process $e^+e^- \rightarrow W^+W^-$ is built up by $s$–channel exchanges of $\gamma$, $Z$ and $Z'$ bosons and $t$–channel neutrino $\nu$ and $N$ exchanges, cf. Fig.1. As a result, the cross section and the distributions in the process $e^+e^- \rightarrow W^+W^-$ are modified by $W$ and $\nu$ mixings and $N$ exchange, and by the mixings
in the neutral boson sector and the $Z'$ exchange.

The helicity amplitudes of the process (1) in the LR model can be written as

$$
\mathcal{M}(\sigma, \bar{\sigma}; \lambda, \bar{\lambda}) = \mathcal{M}_\gamma + \mathcal{M}_Z + \mathcal{M}_{Z'} + \mathcal{M}_\nu + \mathcal{M}_N
$$

where $\sigma/2$ and $\bar{\sigma}/2$ denote the helicities of the incoming electron and positron, respectively; $\lambda$ and $\bar{\lambda}$ are the helicities of the $W^-$ and $W^+$ bosons, respectively. In the following analysis we closely follow the notation of Ref.[19] by writing the helicity amplitudes in terms of reduced amplitudes $\tilde{M}$ which are defined by the relation

$$
\mathcal{M}_X(\sigma, \bar{\sigma}; \lambda, \bar{\lambda}|\theta) = \sqrt{2} \sigma e^{i2} \tilde{\mathcal{M}}_X(\sigma, \bar{\sigma}; \lambda, \bar{\lambda}|\theta) d^{J_0}_{\Delta\sigma, \Delta\lambda}(\theta)
$$

where $X = \gamma, Z, Z', \nu, N$ and

$$
\Delta\sigma = \frac{1}{2}(\sigma - \bar{\sigma}), \quad \Delta\lambda = (\lambda - \bar{\lambda}), \quad J_0 = \max(|\Delta\sigma|, |\Delta\lambda|)
$$

The angle $\theta$ is the $W^-$ production angle with respect to the electron momentum. The function $d^{J_0}_{\Delta\sigma, \Delta\lambda}(\theta)$ denotes the angular–momentum wave–function associated with the minimum angular momentum $J_0$ in the production process. The relevant $d$ functions are collected in Table 1 which extends the corresponding table of Ref.[19].

a) For the $W^\pm$ helicity combinations $|\lambda - \bar{\lambda}| = 2$, i.e. $(\lambda, \bar{\lambda}) = (+, -)$ and $(-, +)$, only $t$-channel neutrino exchanges contribute. The reduced helicity amplitudes are given by

$$
\begin{align*}
\tilde{\mathcal{M}}_\nu(\sigma = -\bar{\sigma} = -; \lambda = -\bar{\lambda} = \pm) &= -\frac{\sqrt{2}}{\sin^2\theta_W (1 + \beta^2 - 2\beta \cos\theta)} \frac{\cos^2\chi_W \cos^2\chi_N}{(1 + \beta^2 - 2\beta \cos\theta + \mu_N^2)} \\
\tilde{\mathcal{M}}_N(\sigma = -\bar{\sigma} = -; \lambda = -\bar{\lambda} = \pm) &= -\frac{\sqrt{2}}{\sin^2\theta_W (1 + \beta^2 - 2\beta \cos\theta + \mu_N^2)} \\
\tilde{\mathcal{M}}_N(\sigma = -\bar{\sigma} = +; \lambda = -\bar{\lambda} = \pm) &= \frac{\sqrt{2}}{\sin^2\theta_W (1 + \beta^2 - 2\beta \cos\theta + \mu_N^2)} \\
\tilde{\mathcal{M}}_N(\sigma = \bar{\sigma} = \pm; \lambda = -\bar{\lambda} = \pm) &= 0
\end{align*}
$$

The $W$ velocity is given by $\beta = \sqrt{1 - 4M_W^2/s}$ and $\mu_N = 2M_N/s$ is the scaled mass parameter of the heavy neutrino. The amplitude $\tilde{\mathcal{M}}_\nu$ describes the light neutrino exchange which is reduced by the coefficient $\cos^2\chi_W \cos^2\chi_N$ compared with the SM amplitude [i.e. $\sim (1 - \chi_W^2 - \chi_N^2)$ in the leading order of the mixing angles]. The amplitudes $\tilde{\mathcal{M}}_N$ are generated by the heavy neutrino exchange, characteristic of the LR model. The amplitude $\tilde{\mathcal{M}}_N(\sigma = -, \bar{\sigma} = +)$ is induced by the left–handed interaction in both $W\nu e$ vertices, suppressed by the neutrino mixing angle squared $\sin^2\chi_N$. The amplitude $\tilde{\mathcal{M}}_N(\sigma = +, \bar{\sigma} = -)$, by contrast, is induced by the right–handed interaction in both charged–current vertices,
suppressed by the mixing angle squared \( \sin^2 \chi_W \). The amplitude \( \tilde{M}_N(\sigma = \sigma) \) for equal electron/positron helicities is built up by a mixture of left– and right–handed interactions, which vanishes for \( |\Delta \lambda| = 2 \).

b) For \( |\lambda - \bar{\lambda}| = 0,1 \) also \( s \)-channel exchanges of vector bosons are possible. The corresponding reduced amplitudes can be written as

\[
\tilde{M}_\gamma(\sigma = -\sigma = \pm; \lambda, \bar{\lambda}) = -\beta A_{\lambda \bar{\lambda}} \\
\tilde{M}_Z(\sigma = -\sigma = \pm; \lambda, \bar{\lambda}) = \beta A_{\lambda \bar{\lambda}} \left[ \left( 1 - \frac{\sin^2 \chi_W}{\cos^2 \theta_W} \right) \cos \chi_Z + \frac{\sqrt{\cos 2\theta_W}}{\cos \theta_W} \sin^2 \chi_W \sin \chi_Z \right] \\
\tilde{M}_{Z'}(\sigma = -\sigma = \pm; \lambda, \bar{\lambda}) = \beta A_{\lambda \bar{\lambda}} \left[ \left( 1 - \frac{\sin^2 \chi_W}{\cos^2 \theta_W} \right) \sin \chi_Z + \frac{\sqrt{\cos 2\theta_W}}{\cos \theta_W} \sin^2 \chi_W \cos \chi_Z \right] \\
\tilde{M}_\nu(\sigma = -\sigma = \pm; \lambda, \bar{\lambda}) = \frac{\cos^2 \chi_W \cos^2 \chi_N}{2\beta \sin^2 \theta_W} \left[ B_{\lambda \bar{\lambda}} - \frac{1}{1 + \beta^2 - 2\beta \cos \theta + \mu_N^2 C_{\lambda \bar{\lambda}}^N} \right] \\
\tilde{M}_N(\sigma = -\sigma = \pm; \lambda, \bar{\lambda}) = -\frac{\sin^2 \chi_N}{2\beta \sin^2 \theta_W} \left[ B_{\lambda \bar{\lambda}} - \frac{1}{1 + \beta^2 - 2\beta \cos \theta + \mu_N^2 C_{\lambda \bar{\lambda}}^N} \right] \\
\tilde{M}_N(\sigma = \bar{\sigma} = \pm; \lambda, \bar{\lambda}) = -\frac{\sqrt{2} \sin \chi_N \sin \chi_W}{2\beta \sin^2 \theta_W} \mu_N \left[ D_{\lambda \bar{\lambda}} - \frac{1}{1 + \beta^2 - 2\beta \cos \theta + \mu_N^2 E_{\lambda \bar{\lambda}}} \right]
\]

The subamplitudes \( A_{\lambda \bar{\lambda}} \) to \( E_{\lambda \bar{\lambda}} \) are collected in Table 2.

Evidently, the process \( e^+e^- \to W^+W^- \) in the LR symmetric model receives non-zero contributions from all combinations of the electron/positron helicities while the electron/positron helicities are always opposite in the SM. The photonic \( s \)-channel contribution is not modified in the LR model. The \( Z' \) exchange effect is small because it is suppressed both by heavy propagator and small mixing effects. While the \( s \)-channel \( Z \) boson contribution is changed by an overall mixing factor, the \( t \)-channel neutrino contribution receives non-trivial new contributions from the exchange of the heavy neutrino. The angular distributions of the final state \( W \)-bosons are therefore modified for all \( W^+W^- \) helicity combinations. Thus, not only the total cross section but also the angular distributions of the \( W \)-bosons and their decay products are modified in the LR models.
Table 1
Angular momentum $d$–functions of the process $e^+e^- \rightarrow W^+W^-$ as defined in Eq.(9); $\eta_\sigma = \text{sign}(\Delta \sigma)$ and $\eta_\lambda = \text{sign}(\Delta \lambda)$.

| $\Delta \sigma$ | $\Delta \lambda$ | $J_0$ | $d_{J_0}^{\Delta \sigma,\Delta \lambda}$ |
|----------------|-----------------|-------|---------------------------------|
| 1              | 2               | 2     | $\eta_\lambda(1 + \eta_\sigma \eta_\lambda \beta) \sin \theta/2$ |
| 1              | 1               | 1     | $(1 + \eta_\sigma \eta_\lambda \beta)/2$ |
| 1              | 0               | 1     | $-\eta_\sigma \sin \theta/\sqrt{2}$ |
| 0              | 2               | 2     | 0 |
| 0              | 1               | 1     | $\sin \theta/\sqrt{2}$ |
| 0              | 0               | 0     | 1 |

Table 2
Subamplitudes of the process $e^+e^- \rightarrow W^+W^-$ as defined in Eq.(12) for the $W$ helicities $|\lambda - \bar{\lambda}| = 0, 1$. The parameter $\gamma$ is the boost factor $\gamma = \sqrt{s}/2M_W$.

| $\lambda \bar{\lambda}$ | $A_{\lambda \bar{\lambda}}$ | $B_{\lambda \bar{\lambda}}$ | $C^\nu_{\lambda \bar{\lambda}}$ | $C^N_{\lambda \bar{\lambda}}$ | $D_{\lambda \bar{\lambda}}$ | $E_{\lambda \bar{\lambda}}$ |
|------------------------|------------------------|------------------------|------------------------|------------------------|------------------------|------------------------|
| $\pm \pm$              | 1                      | 1                      | $1/\gamma^2$           | $1/\gamma^2 + \mu_N^2$ | $\lambda/2$            | $\lambda[(1 + \lambda \sigma \beta)^2 + \mu_N^2]/2$ |
| $\pm 0$                | $2\gamma$              | $2\gamma$              | $2(1 + \lambda \beta)/\gamma$ | $2[(1 - \lambda \sigma \beta)/\gamma + \mu_N^2 \gamma]$ | 0                      | $\gamma \beta(1 + \lambda \sigma \beta)$ |
| 0                      | $2\gamma$              | $2\gamma$              | $2(1 - \bar{\lambda} \beta)/\gamma$ | $2[(1 + \bar{\lambda} \sigma \beta)/\gamma + \mu_N^2 \gamma]$ | 0                      | $-\gamma \beta(1 + \bar{\lambda} \sigma \beta)$ |
| 00                     | $1 + 2\gamma^2$        | $2\gamma^2$            | $2/\gamma^2$           | $2(1/\gamma^2 + \mu_N^2 \gamma^2)$ | $\sigma \beta \gamma^2$ | $\sigma \beta \gamma^2 \mu_N^2$ |

4. Phenomenological Analyses. To exemplify the results, we analyze the LR contribution to the total cross section of the process $e^+e^- \rightarrow W^+W^-$. It is important to realize that standard analyses of anomalous static $W^\pm$ parameters cannot be applied due to the heavy neutrino exchange in the $t$–channel which affects angular momentum states $|J_z| \neq 1$. The total cross section is given by the differential cross section integrated over the production angle $\theta$ in the following form:

$$\sigma[e^+e^- \rightarrow W^+W^-] = \frac{\pi \alpha^2}{4s} \beta \int d\cos \theta \sum_{\sigma,\bar{\sigma},\lambda,\bar{\lambda}} |\tilde{M}(\sigma, \bar{\sigma}; \lambda, \bar{\lambda})|^{2} d_{J_0}^{\Delta \sigma,\Delta \lambda}|^2$$

(13)

where $\alpha \approx 1/128$ is the electromagnetic coupling evaluated at the high–energy scale $s$.

By adopting the estimate $\chi_Z \lesssim 10^{-4}$ for $M_{Z'} = 1$ TeV, as required by precision data [11],
Fig. 2. The deviation of the unpolarized cross section of the process $e^+e^- \rightarrow W^+W^-$ at LEP2 from the SM prediction as a functions of the mixing angle $\chi_N$ [figure (A)] and the neutrino mass $M_N$ [figure (B)]. The chosen numerical values are shown in the figure.

the effect of $Z - Z'$ mixing on the cross section is very small, i.e. at the per-mille level, and it can in general be neglected.

The deviation from the SM cross section scaled with respect to the improved Born approximation, is exemplified in Figs.2 for LEP2; in (A) as a function of the mixing angle $\chi_N$ for fixed $M_N$, and in (B) as a function of the heavy neutrino mass $M_N$ for fixed $\chi_N$. For the sake of simplicity $\sin{\chi_W}$ is taken $10^{-2}$ in these numerical examples. The collision energy is set to $\sqrt{s} = 200$ GeV and the values of the parameters are indicated in the figure. Evidently, the $WW$ production cross section is reduced for this set of the parameters compared with the SM. The 1% level of the deviation from the SM is reached for the mixing angles $\sin{\chi_N} \sim 0.07$. While the quadratic dependence of $(\sigma - \sigma_{SM})/\sigma_{SM}$ on $\sin{\chi_N}$ is strong, the dependence on $M_N$ is rather weak. The effect of the non-zero mixing angle and the heavy neutrino exchange cancel each other partly. While non-zero mixing angles decrease the cross section, the neutrino $N$ exchange shifts the cross section upward. This correlation is also evident from Fig.3 (A). The area on the $[M_N, \sin{\chi_N}]$ plane in which the $W^+W^-$ production cross section of the LR model deviates by more than 1% from the cross section of the Standard Model lies above the solid curve. As a result, for relatively light right–handed neutrino masses larger mixing angles are allowed than for heavy masses.

We have repeated the analyses for a linear collider energy $\sqrt{s} = 500$ GeV in Fig.4. While for a relatively light right–handed neutrino $N$ the increase in the experimental sensitivity is not remarkable, for a heavy neutrino $N$, by contrast, a high-energy linear collider will test the neutrino mixing angles with significantly higher precision than LEP2. The same
Fig. 3. The sensitivity curves for (A) LEP2 and (B) a linear collider of energy $\sqrt{s} = 500$ GeV. In the region above the curve the total cross section of the process $e^+e^- \rightarrow W^+W^-$ deviates from the SM prediction by more than 1%. Note the non–trivial behavior of the contours which is a consequence of the counteracting mixing and neutrino–exchange effects.

conclusion follows also from Fig.3 (B). The reason for this behavior is the cancellation of the neutrino mixing effect by the heavy neutrino exchange. If the latter is minimized at high neutrino mass, a linear collider will probe the mixing angle $\chi_N$ with high accuracy.

Linear colliders may also have $e^-\gamma$ collision mode in addition to the usual $e^+e^-$ one. Because in $e^-\gamma$ collisions the initial photon can test directly the trilinear gauge boson coupling [20], this collision mode will be ideal to discriminate the new physics discussed in this work from the anomalous gauge boson couplings.

The analysis can be improved by exploiting the angular distributions of the $W$-bosons and their decay products; significant improvements can be expected in the region of small neutrino mass $M_N$. Such an analysis which depends more strongly on the experimental conditions, is beyond of the scope of the present letter.

5. Conclusions. We have shown that given the present experimental bounds on the $W-W'$, $\nu-N$ mixings and on the $N$ mass, observable deviations from the SM cross section and distributions of the process $e^+e^- \rightarrow W^+W^-$ at $e^+e^-$ colliders are possible in LR symmetric models. The dominant new effects are associated with the $t$-channel exchange of neutrinos which differ from the anomalous self–interaction effects of the electroweak gauge bosons. Independent analyses of the angular distributions of the $W$ boson helicity components are needed to explore these new effects. Since no major deviations from the SM have been found at LEP2, new bounds on mixing angles and on the heavy neutrino mass can be derived from the data. High precision experiments at future $e^+e^-$ linear colliders in
Fig. 4. The same analysis as in Fig. 2 for the collision energy $\sqrt{s} = 500$ GeV. 

the TeV range will extend the sensitivity to heavy neutrinos into the multi–TeV region according to the scaling law $M_N \sim s^{1/2} \cdot L^{1/4}$ for the energy squared $s$ and the integrated luminosity $L$; for $\sqrt{s} \sim 1$ TeV and $L \sim 1$ ab$^{-1}$ this corresponds to an increase by more than an order of magnitude compared to LEP2.

Acknowledgements

I would like to thank P. Zerwas with whom most of this work has been done. I am grateful to DESY Theory Group for hospitality. His work is partly supported by the U.S. Department of Energy under Grant No. DE-FG03-94ER40837.

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