PHYSICAL REVIEW D VOLUME 52, NUMBER 11 1 DECEMBER 1995

Inverse neutrinoless double-β decay in gauge theories with CP violation

J. Gluza* and M. Zrałek†
Department of Field Theory and Particle Physics, Institute of Physics, University of Silesia, Uniwersytecka 4, PL-40-007 Katowice, Poland
(Received 13 February 1995)

We investigate the $e^-e^- \to W_i^- W_j^-$ reactions for the various gauge boson production processes in the frame of the standard model with additional right-handed neutrinos and in the left-right symmetric model. The present bounds on the various model parameters are taken into account. The question of the cross section behavior for large energy and the CP violation problems are discussed.

PACS number(s): 13.10.+q, 11.30.Er, 12.60.Cn, 14.60.St

I. INTRODUCTION

The nature of the known electron, μ and τ neutrinos and the existence of the heavier ($M_\nu > M_Z$) neutrinos are still unresolved problems in particle physics. In the standard model all neutrinos are massless. There is, however, some more or less strongly established evidence which requires massive neutrinos. Problems with (i) the solar (electron) neutrinos, (ii) atmospheric (μ) neutrinos, and (iii) the hot dark matter can be resolved if neutrinos are massive particles. Within the framework of the extended electroweak models neutrinos are usually massive and Majorana-type. At low energies such models can be probed by looking for rare processes, such as the neutrinoless double-β decay.

High-energy e^-e^- collision, a possible option in next linear colliders (NLC), may provide a new test for those $\Delta L = 2$ interactions via the $e^-e^- \to W_i^- W_j^-$ reaction, where W_i may represent standard model gauge bosons or the additional charged gauge bosons. This inverse neutrinoless double-β decay was proposed some time ago [1] and since then has been investigated several times [2]. We have decided to study the process once more for a few reasons. First, we would like to give the numerical values of the total cross sections for the $e^-e^- \to W_i^- W_j^-$ processes, taking into account the up-to-date bounds for the various model parameters. Second, we would like to study the problem of unitarity and to find the conditions for the correct high-energy behavior of the cross sections in various gauge models. For this purpose it is necessary to consider all model's ingredients [for example, in the left-right (L-R) symmetric model both left- and right-handed double-charged gauge bosons $\delta_{L,R}^{\pm\pm}$ have to be considered]. The calculated helicity amplitudes for the process give us the opportunity to have a look at the unitarity cancellations in a very precise way. Finally, we investigate the problem of CP-symmetry breaking in our process and compare its size for various models.

*Electronic address: gluza@usctoux1.cto.us.edu.pl
†Electronic address: zralek@usctoux1.cto.us.edu.pl

In the next section we define the mass Lagrangian and the CP-breaking parameters. In Sec. III we investigate which ingredients of the models are responsible for correct high-energy cross section behavior. The numerical values of total cross sections for model's parameters, which satisfy existing bounds and the size of CP-symmetry breaking, are given in Sec. IV.

II. THE MASS LAGRANGIAN AND CP VIOLATION

In the class of models which we consider the mass Lagrangian for neutrinos and charged leptons is given by

$$L_{\text{mass}} = -\frac{1}{2}(\bar{N}_L^c M_\nu N_R + \bar{N}_R M_\nu^* N_L^c) - (\bar{l}_L M_l l_R + \bar{l}_R M_l^\dagger l_L), \tag{1}$$

where

$$N_R = \begin{pmatrix} i\gamma^2 \nu_L^* \\ \nu_R \end{pmatrix}, \quad N_L^c = \begin{pmatrix} \nu_L \\ i\gamma^2 \nu_R^* \end{pmatrix}$$

are $(n_L + n_R)$ dimensional row vectors of neutrino fields and $l_{L(R)}$ are n_L dimensional charged lepton fields. M_l and M_ν are $n_L \times n_L$ complex and $(n_L + n_R) \times (n_L + n_R)$ symmetric-complex matrices, respectively. The matrix M_ν is usually divided into four parts

$$M_\nu = \begin{array}{cc} \begin{pmatrix} \overbrace{M_L}^{n_L} & \overbrace{M_D}^{n_R} \\ M_D^T & M_R \end{pmatrix} & \begin{array}{l} \}n_L \\ \}n_R \end{array} \end{array}. \tag{2}$$

We know, that without changing the physical meaning (all elements of the Lagrangian, except the mass term [Eq. (1)], will not change) some matrices can be made diagonal. In the models without right-handed charged current interaction two matrices, e.g., M_l and M_R, can be chosen diagonal, in the L-R symmetric models only one, e.g., M_l. In all kinds of models which we consider the basis of lepton fields can be chosen in such a way that the charged lepton mass matrix M_l is real and diagonal. Then all the lepton-violating CP phases are present in the neutrino mass matrix M_ν [Eq. (2)]. The number of CP-breaking phases depends on the model. As an example we consider two kinds of models.

0556-2821/95/52(11)/6238(11)/$06.00

(1) Standard model with the additional three neutral right-handed singlets (RHS's). In this model $n_L = n_R = 3$ and, in our basis,

$$M_\nu = \begin{pmatrix} 0 & M_D \\ M_D^T & M_R \end{pmatrix} , \tag{3}$$

where $M_R = \text{diag}\,(M_1, M_2, M_3)$ are real numbers and

$$M_D = \begin{pmatrix} a_1 e^{i\alpha_1} & a_2 e^{i\alpha_2} & a_3 e^{i\alpha_3} \\ b_1 e^{i\beta_1} & b_2 e^{i\beta_2} & b_3 e^{i\beta_3} \\ c_1 e^{i\gamma_1} & c_2 e^{i\gamma_2} & c_3 e^{i\gamma_3} \end{pmatrix} . \tag{4}$$

The CP symmetry is satisfied if [3]

$$\sin(\alpha_i - \alpha_j) = \sin(\beta_i - \beta_j) = \sin(\gamma_i - \gamma_j) = 0 \quad \text{for } i, j = 1, 2, 3 \tag{5}$$

and six phases $\alpha_1 - \alpha_2 = \chi_1$, $\alpha_1 - \alpha_3 = \chi_2$, $\beta_1 - \beta_2 = \rho_1$, $\beta_1 - \beta_3 = \rho_2$, $\gamma_1 - \gamma_2 = \eta_1$, and $\gamma_1 - \gamma_3 = \eta_2$ break the symmetry.

(2) The left-right symmetric model. We consider two versions of the model. At the beginning we assume that there is the explicit L-R symmetry with all vacuum expectation values real (no spontaneous CP-symmetry breaking) [3]. Then the M_L and M_R matrices are proportional,

$$M_L = \text{const} \times M_R , \tag{6}$$

and M_D is Hermitian:

$$M_D = M_D^\dagger . \tag{7}$$

There are only six CP-violating phases [3] and we can choose them in the following way:

$$M_R = \begin{pmatrix} M_{11} e^{i\alpha_1} & M_{12} & M_{13} \\ M_{12} & M_{22} e^{i\alpha_2} & M_{23} \\ M_{13} & M_{23} & M_{33} e^{i\alpha_3} \end{pmatrix}$$

and (8)

$$M_D = \begin{pmatrix} m_{11} & m_{12} e^{i\beta_1} & m_{13} e^{i\beta_2} \\ m_{12} e^{-i\beta_1} & m_{22} & m_{23} e^{i\beta_3} \\ m_{13} e^{-i\beta_2} & m_{23} e^{-i\beta_3} & m_{33} \end{pmatrix} .$$

This model is known in the literature as manifest (MLRS) or quasimanifest (QMLRS) L-R symmetric one.

We consider also the full version of the L-R model where there is no relation between M_L and M_R matrices and M_D is not Hermitian [nonmanifest L-R symmetric model (NMLRS)]. In this case there are 18 CP-violating phases. The phases of the matrices M_L, M_R, and M_D, which do not satisfy the CP-conserving relations [3],

$$(M_L)_{ij} = |(M_L)_{ij}|\, e^{i\frac{1}{2}(\delta_i + \delta_j)},$$

$$(M_R)_{ij} = |(M_R)_{ij}|\, e^{-i\frac{1}{2}(\delta_i + \delta_j)}, \tag{9}$$

and

$$(M_D)_{ij} = |(M_D)_{ij}|\, e^{-i\frac{1}{2}(\delta_i - \delta_j)}$$

break the CP symmetry.

The neutrino mass matrix is diagonalized by the complex orthogonal transformation

$$U^T M_\nu U = \text{diag}[|m_1|, |m_2|, ..., |m_{n_L+n_R}|] \equiv M_{\nu d}. \tag{10}$$

If we denote

$$U = \begin{pmatrix} U_L^* \\ U_R \end{pmatrix}$$

then the mixing matrices K_L and K_R in the physical left- and right-charged current interactions [see Eq. (A2) in the Appendix] are given by

$$K_{L,R} = U_{L,R}^\dagger . \tag{11}$$

There are reversed connections between $M_{L(R)}, M_D$ and $M_{\nu d}$:

$$M_L = K_L^\dagger M_{\nu d} K_L^*, \quad M_D = K_L^\dagger M_{\nu d} K_R,$$

and (12)

$$M_R = K_R^T M_{\nu d} K_R$$

which will be useful in the further considerations.

III. BEHAVIOR OF CROSS SECTIONS FOR HIGH ENERGY

The reduced helicity amplitudes (see the Appendix) have bad high energy behavior. As $\lim_{s\to\infty} \gamma_{1,2} \sim \frac{\sqrt{s}}{2M_{1,2}}$ the amplitudes with $|\Delta\lambda| = 1$ proportional to γ_1 and the amplitudes with $\lambda_1 = \lambda_2 = 0$ (proportional to $\gamma_1\gamma_2$) increase with increasing energy. Of course, these divergences cannot appear in the total cross section so there must exist mechanisms which cause their cancellation. As the cancellations of these divergences have influence on the size of the total cross section it is instructive to show how it happens. There are several reasons why the bad high energy behavior does not appear in the helicity amplitudes for $\sqrt{s} \gg m_a, M_i$, $a = 1, ..., 6$, $i = 1, 2$.

In the L-R symmetric model for $\Delta\sigma = \pm 1$ in the t and the u channels there are (see the Appendix)

$$R_{t(u)} M_{t(u)}\, (\Delta\sigma = \pm 1, |\Delta\lambda| = 1) \longrightarrow \frac{1}{M_i \sqrt{s}\,(1 \mp \cos\Theta)} \sum_a (K_{L(R)}^T)_{ma} m_a^2 \left(K_{R(L)}\right)_{an} \tag{13}$$

and

$$R_{t(u)} M_{t(u)}\, (\Delta\sigma = \pm 1, \lambda_1 = \lambda_2 = 0) \longrightarrow \frac{1}{2 M_i M_j (1 \mp \cos\Theta)} \sum_a (K_{L(R)}^T)_{ma} m_a^2 (K_{R(L)})_{an} . \tag{14}$$

To obtain these formulas we use the orthogonality property of the K_R and K_L matrices:

$$K^T_{L(R)} K_{R(L)} = 0. \tag{15}$$

For $\Delta\sigma = 0$ the problem is only with one amplitude for $\lambda_1 = \lambda_2 = 0$. Separately, each amplitude in the u, t, and s channels is divergent but when we add them together, then

$$\begin{aligned} &B^{ij}(\sigma_1,\sigma_2)\left[R_t M_t + R_s M_s\right] + B^{ij}(\sigma_2,\sigma_1) R_u M_u \\ &\longrightarrow -\frac{\sqrt{2}}{\sqrt{s} M_i M_j} B^{ij}\left(\pm\frac{1}{2},\pm\frac{1}{2}\right)\left[\frac{1}{2\left(1-\cos^2\Theta\right)}\sum_a (K^T_{R(L)})_{ma} m_a^3 \left(K_{R(L)}\right)_{an}\right. \\ &\qquad \left. + \left(M^2_{\delta_{R,L}} - i\Gamma_{\delta_{R,L}} M_{\delta_{R,L}}\right)\sum_a \left(K^T_{R,L}\right)_{ma} m_a \left(K_{R,L}\right)_{an}\right]. \end{aligned} \tag{16}$$

The crucial thing needed to obtain this high energy behavior is existence of the three s, t, and u channels. Without the $\delta^{--}_{R,L}$ bosons exchange in the s channel the unitarity would be violated. And it is important that both left δ^{--}_L and right δ^{--}_R doubly charged Higgs bosons are present. They give contributions to the different helicity amplitudes (see Appendix).

We can see that in the models with L-R symmetry the appropriate high energy behavior is guaranteed because of the following reasons: (i) the left and right mixing matrices in the charged current are orthogonal [Eq. (15)], and (ii) there exist two doubly charged Higgs bosons $\delta^{--}_{L,R}$ with proper relations between various couplings.

We can ask now how it is possible, that in the models without L-R symmetry, where only one charged gauge boson $W^\pm$ exists, the unitarity is also satisfied. In this class of models there is no right-handed charged current. Formally this is equivalent with the assumption that the mixing angle ξ and the mixing matrix K_R vanish:

$$\xi = 0,\; K_R = 0. \tag{17}$$

There are no doubly charged Higgs bosons so there is no contribution to the amplitude from the s channel:

$$M_s = 0. \tag{18}$$

Only one helicity amplitude in the t and u channels $M_{t,u}(-\frac{1}{2},-\frac{1}{2},0,0)$ seems to be divergent. But if we look carefully at the high energy behavior we have

$$\begin{aligned} &R_{t(u)} M_{t(u)}\left(-\frac{1}{2},-\frac{1}{2},0,0\right) \\ &\longrightarrow \frac{1}{\sqrt{2} M_i M_j}\left[\sqrt{s}\sum_a \left(K^T_L\right)_{ma} m_a \left(K_L\right)_{an}\right. \\ &\qquad \left. + \frac{1}{\sqrt{s}}\frac{1}{(1\mp\cos\Theta)}\sum_a \left(K^T_L\right)_{ma} m_a^3 \left(K_L\right)_{an}\right]. \end{aligned} \tag{19}$$

Now from Eq. (12) there is

$$\sum_a \left(K^T_L\right)_{ma} m_a (K_L)_{an} = \left(K^T_L M_{\nu d} K_L\right)_{mn} = \left(M^*_L\right)_{mn}. \tag{20}$$

In the RHS models without doubly-charged bosons the appropriate Yukawa mechanism which generates the mass matrix M_L is not present, so

$$M_L = 0, \tag{21}$$

and the first term in Eq. (19) which is proportional to $\sqrt{s}$ disappears, what guarantees the correct high energy behavior of the cross section. In the models where Higgs triplets are present the unitarity is preserved in similar way as in the L-R symmetric models.

IV. NUMERICAL RESULTS

First we investigate the total cross sections for production of gauge bosons in e^-e^- reaction and their dependence on various model parameters. Then we consider the CP-symmetry breaking.

A. Total cross section in L-R model

To calculate the cross section we need to know the values of the model parameters (see [4]):
the mixing angle ξ

$$\xi \simeq \left(\frac{2k_1k_2}{k_1^2+k_2^2}\right)\frac{M^2_{W_1}}{M^2_{W_2}}; \tag{22}$$

the masses of the gauge bosons,

$$M^2_{W_1} \simeq \frac{g^2}{2}\left(k_1^2+k_2^2\right),\; M^2_{W_2} = \frac{g^2}{2}v_R^2; \tag{23}$$

the masses of doubly-charged Higgs particles,

$$M^2_{\delta_L} \simeq \frac{1}{2}v_R^2,\;\; M^2_{\delta_R} \simeq 2v_R^2. \tag{24}$$

In our numerical analysis we take that $2k_1k_2 \simeq k_1^2+k_2^2$, so practically only one parameter, the mass of heavy gauge boson M_{W_2} is free (M_{W_1} and $g = e/\sin\Theta_W$ are known from the standard model). We do not calculate the decay width for doubly charged Higgs bosons but we put them in the form

$$\Gamma_{\delta^{--}_{L,R}} = \Gamma_{W_1} M_{\delta^{--}_{L,R}}/M_{W_1}. \tag{25}$$

In addition to the above model parameters the cross section depends very strongly on the neutrino masses and mixings. First of all, we can see [Eqs. (15), (A6)–(A8)] that if all neutrinos are massless ($m_a = 0$), then the functions $R_{t,u,s}$ vanish and the cross section is zero. This fact is very well known. For massless neutrinos there is no way to distinguish between the Dirac and the Majorana cases, lepton number conservation is restored, and our process cannot occur.[1]

From existing terrestrial experiments we know that [5]

$$m_{\nu_e} < 5.1 \text{ eV}, \quad m_{\nu_\mu} < 0.27 \text{ MeV}, \quad m_{\nu_\tau} < 31 \text{ MeV} \tag{26}$$

and the heavy neutrinos, if they exist, must have masses $M_N > \frac{M_Z}{2}$ [5] or even $M_N > M_Z$ if additional assumptions about the νNZ coupling are made [6]. There are other laboratory experiments (double-β decay and neutrino oscillation in the vacuum) which tell us something about neutrino masses and mixings. Moreover, there are also solar, astrophysical and cosmological observations which can also give some information about masses and mixings; see, e.g., [7]. It follows from all existing observations that the three known neutrinos (ν_e, ν_μ, ν_τ) should have small, almost degenerate masses in the range between 0 and several eV. The other three neutrinos predicted by the L-R model must have masses above ~ 100 GeV.

One of the possible choices to assure the mentioned pattern of the neutrino masses is to assume that M_D is almost a rank 1 matrix and M_R is almost diagonal with large diagonal elements M_i ($i = 1, 2, 3$) which satisfy the constraints [8]

$$\frac{1}{M_1} + \frac{1}{M_2} + \frac{1}{M_3} \simeq 0. \tag{27}$$

On the other hand, the elements of the M_R matrix are determined by the right-handed vacuum expectation value v_R and should be bounded by the mass of the heavy gauge boson [Eq. (23)]:

$$M_i \leq \frac{2M_{W_2}}{g}. \tag{28}$$

Let us take the matrices M_D and M_R in the form

$$M_D = \begin{pmatrix} 1.0 & 1.0 & 1.0 \\ 1.0 & 1.0 & 1.0 \\ 1.0 & 1.0 & 1.0 - 10^{-6} \end{pmatrix}$$

and

$$M_R = \begin{pmatrix} M_1 & 10^{-6} & 10^{-6} \\ 10^{-6} & M_2 & 10^{-6} \\ 10^{-6} & 10^{-6} & M_3 \end{pmatrix}. \tag{29}$$

Then one neutrino is massless $m_1 = 0$, the other two have very small masses, $m_2 \simeq 0$, $m_3 \simeq 0$, and the masses of the heavy neutrinos are given by the diagonal elements of matrix M_R:

$$m_4 \simeq | M_1 |, \; m_5 \simeq | M_2 |, \; m_6 \simeq | M_3 | \, .$$

For this choice of the neutrino mass matrix the total cross section for various other model parameters is presented on Fig. 1. In the frame of the NMRLS model the present experimental bound on M_{W_2} is not so high and $M_{W_2} \geq 600$ GeV is still a reasonable limit [9]. Then the mixing angle ξ does not suppress so much the helicity amplitudes. Curves depicted by capital A represent the result for $M_{W_2} = 600$ GeV. In the frame of MLRS or QMLRS models the bound on M_{W_2} is larger than $M_{W_2} \geq 1600$ GeV [10] and the cross sections are lower (the line depicted by capital B in Fig. 1). The reasonably high luminosity in the next e^-e^- collision will not give us the possibility to observe the $W_1^- W_1^-$ pair production in this case. The cross section depends also on the heavy neutrinos masses. For $m_4 \sim 100$ GeV (and $m_5 \sim m_6 \sim 200$ GeV) the cross section is small (dashed lines in Fig. 1) and increases with the neutrino mass [the solid lines give the cross section for the biggest acceptable values (Eq. (28)]. On the figure we can also see the influence of the right-handed resonance δ_R^{--}, as for $M_L \sim 0$ the effect of the left-handed resonance δ_L^{--} is not visible (see final remarks in the Appendix).

In the next figures we present the cross sections for production of the light-heavy $W_1^- W_2^-$ gauge bosons (Fig. 2) and two heavy $W_2^- W_2^-$ gauge bosons (Fig. 3). The cross sections are higher but of course the thresholds for

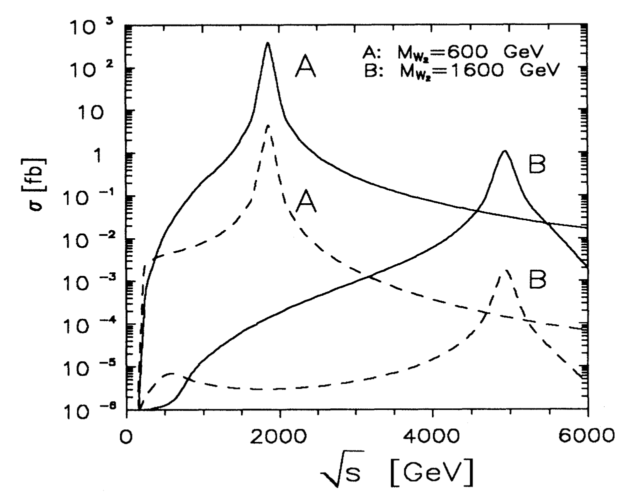

FIG. 1. The total cross section for the process $e^-e^- \to W_1^- W_1^-$ (LR model) as a function of the c.m. energy for various parametrizations of the neutrino mass matrix. Dashed line is for the parametrization given by Eq. (29) with $M_1 = -100$ GeV. Solid line is for the biggest available M_1 [Eq. (28)]. In all cases M_2 and M_3 are chosen in such a way that Eq. (27) is satisfied.

---

[1]Strictly speaking it is impossible to distinguish between the Dirac and Majorana neutrinos for $m_a \to 0$ if there are only left-handed or right-handed current interactions. If both couplings are present (as in the L-R symmetric model) the massless Dirac and Majorana neutrinos are still indistinguishable because of the orthogonality of K_L and K_R matrices ($K_{L,R}^T K_{R,L} = 0$).

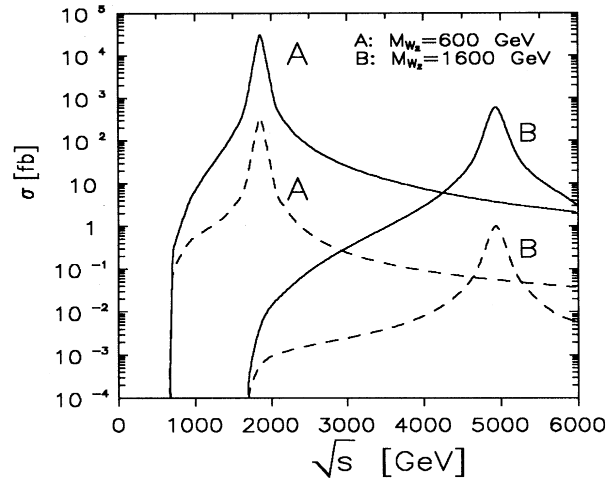

FIG. 2. Process $e^-e^- \to W_1^-W_2^-$ (LR model) as a function of the c.m. energy for various masses of the heavy neutrinos. Notation is the same as for Fig. 1.

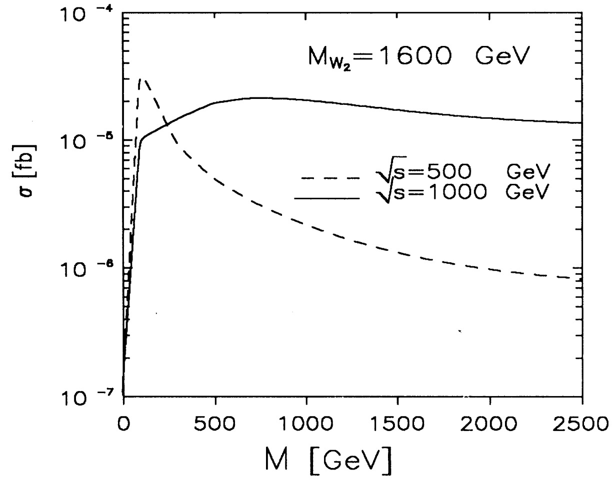

FIG. 4. Process $e^-e^- \to W_1^-W_1^-$ (LR model) as a function of the heavy neutrino mass for NLC, $\sqrt{s} = 500$ GeV (dashed line), and $\sqrt{s} = 1$ TeV (solid line) with parametrization of the heavy neutrino masses by Eq. (29), with $M_1 = -M$, $M_2 = M_3 = 2M$.

these production processes are too high to hope that the appropriate colliders will be built in the near future.

Figure 4 gives the behavior of the total cross section for two light gauge boson production $e^-e^- \to W_1^-W_1^-$ as a function of heavy neutrino masses. We take the following parametrization of the neutrino mass matrix in Eq. (29):

$$M_1 = -M, \quad M_2 = M_3 = 2M. \tag{30}$$

The behavior of $\sigma(M)$ agrees with our previous discussion: $\sigma \to 0$ for $M \to 0$. The cross section increases with the increasing neutrino masses (not shown on the figure) if we go with the mass M above the limit given by Eq. (28).

The dependence of the total cross section on the gauge boson mass M_{W_2}, $\sigma(M_{W_2})$, for various CM energies and different mass matrix parametrizations is depicted in Fig. 5. As was discussed before, the cross section depends strongly on M_{W_2} and this behavior is not influenced too much by the neutrino matrix parametrization.

B. Total cross section in the models with RHS's

We consider the model with three additional right-handed neutrinos which are singlets of the gauge group.

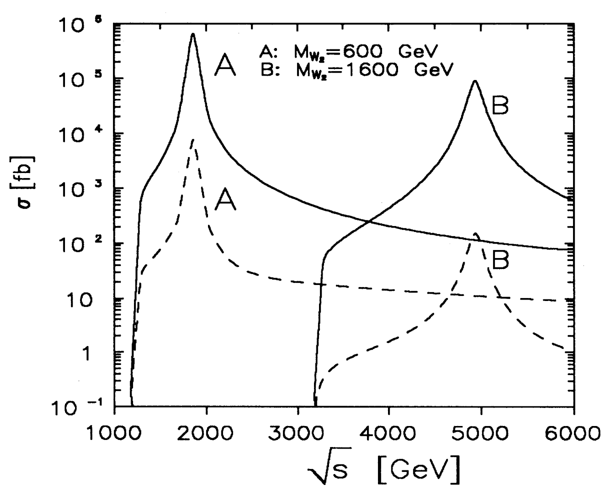

FIG. 3. Process $e^-e^- \to W_2^-W_2^-$ (LR model) as a function of the c.m. energy for various masses of the heavy neutrinos. Notation is the same as for Fig. 1.

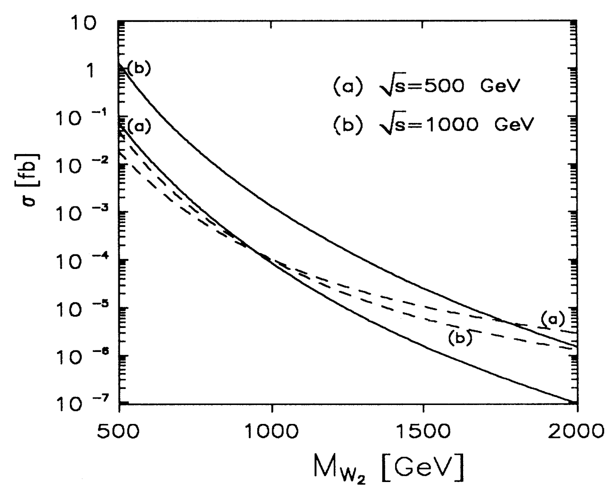

FIG. 5. Process $e^-e^- \to W_1^-W_1^-$ (LR model) as a function of the heavy gauge boson mass for different energies, (a) for $\sqrt{s} = 500$ GeV and (b) for $\sqrt{s} = 1000$ GeV. Again dashed and solid lines are for the neutrino mass paramatrization as in Fig. 1.

At the beginning we assume that the neutrino mass matrix is pure real (CP is conserved) and is the same as before [Eq. (29)]. There are no doubly-charged Higgs bosons so only the t and u channels contribute to the light gauge boson production amplitude $e^-e^- \to W_1^-W_1^-$ ($\xi = 0$ and $K_R = 0$). Only one helicity amplitude with $\sigma_1 = \sigma_2 = -1/2$ gives contribution to the total cross section. The appropriate K_L mixing matrix elements, $(K_L)_{ei}$ ($i = 4, 5, 6$), are decreasing functions of the heavy neutrino mass and the same can be observed in the total cross section (Fig. 6). The $\sigma\left(e^-e^- \to W_1^-W_1^-\right)$ dependence on the c.m. energy is presented in Fig. 7. The cross section is smaller than in the L-R symmetric model and decreases with energy. This is caused by the destructive interferences between heavy neutrinos. It is also very sensitive to the neutrino mass matrix parametrization. If we take, for example,

$$M_D = \begin{pmatrix} 1.0 & 1.0 & 0.9 \\ 1.0 & 1.0 & 0.9 \\ 0.9 & 0.9 & 0.95 \end{pmatrix}, \; M_R = \begin{pmatrix} M_1 & 10 & 20 \\ 10 & M_2 & 10 \\ 20 & 10 & M_3 \end{pmatrix}, \tag{31}$$

the cross section is much larger (dashed line in Fig. 7) and its decreasing with the energy is smaller (destructive interferences between heavy neutrinos are absent although such interferences can appear—see the next section), visible only for higher energies (not shown in Fig. 7). For $\sqrt{s} = 500$ GeV $\sigma \simeq 0.25 \times 10^{-4}$ fb with parametrization from Eq. (29) and $\sigma \simeq 0.5 \times 10^{-3}$ fb with parametrization given by Eq. (31). The cross sections which we get are much smaller than that obtained in other papers (see, e.g., Heusch and Minkowski paper in [2]). The reason follows from the way in which the mixing matrix is treated. In our approach the small masses for the known neutrinos are obtained by the seesaw mechanism. We calculate the K_L matrix by diagonalization of the neutrino mass matrix M_ν which has strong model foundation and gives realistic spectrum of neutrino masses (three light and three heavy neutrinos). Then the mixing matrix elements have the "seesaw" type

$$(K_L)_{ea} \sim \frac{\langle M_D \rangle}{m_a}, \tag{32}$$

and are small for large neutrino mass m_a. The small mixing matrix causes the cross section also to be small. However, such classical "seesaw" mechanism is not the only scenario to explain the small masses of known neutrinos. There are models [11,12] where relation (32) does not work and the mixing matrix elements can be large even for large masses of heavy neutrinos. Then the smallness of the known neutrino mass is guaranteed by the special symmetry arguments which are applicable if number of right-handed neutrinos is larger than 1 (see, e.g., [12]). In such models the mixing matrix elements $(K_L)_{ea}$ are independent parameters and can be bounded by existing experimental data. The present experimental bounds on K_L are such that the cross section $\sigma(e^-e^- \to W^-W^-)$ can be much larger, e.g., $\sigma = 4(64)$ fb for $\sqrt{s} = 0.5(1)$ TeV (see Heusch and Minkowski paper in [2] and for detailed discussion see also [13]).

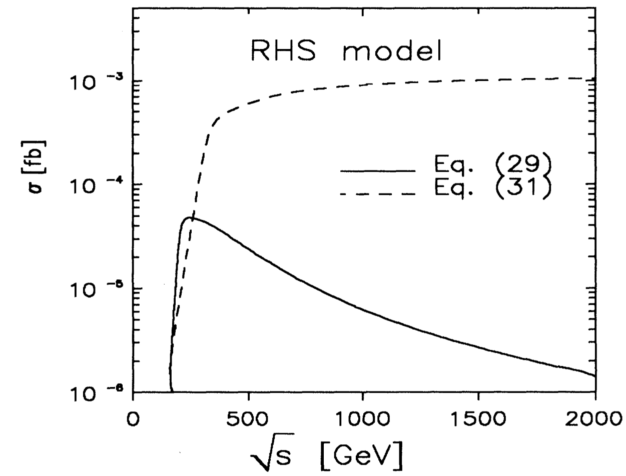

FIG. 7. Process $e^-e^- \to W_1^-W_1^-$ (RHS model) as a function of the c.m. energy with parametrization of the heavy neutrino masses given by Eqs. (29), (30) with M=100 GeV—solid line and by Eq. (31) with $M_1 = 100$ GeV, $M_2, M_3 = 200$ GeV—dashed line.

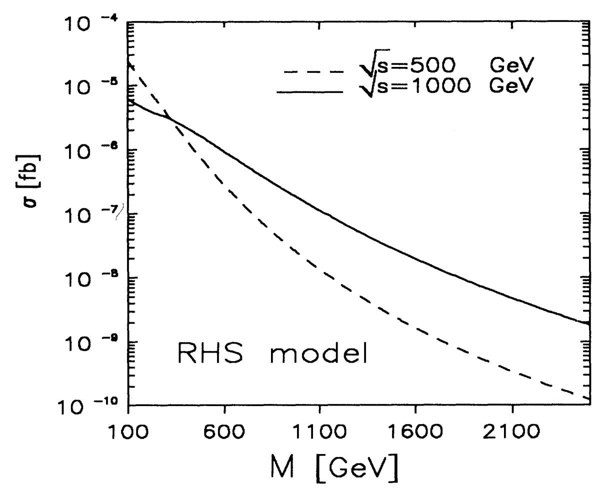

FIG. 6. Process $e^-e^- \to W_1^-W_1^-$ (RHS model) as a function of heavy neutrino mass for NLC (dashed line) and 1 TeV (solid line) colliders with parametrization of the heavy neutrino masses as in Fig. 4.

C. CP-symmetry breaking

There are many phases which can break the CP symmetry [see Eqs. (4),(8)]. We don't study the cross section as a function of all of them. We choose only three phases and we assume that the neutrino mass matrix is given by Eq. (31), where the following changes are made:

$$M_1 \to e^{i\alpha}M, \; M_2 \to 2e^{i\beta}M, \; M_3 \to 3e^{i\gamma}M. \tag{33}$$

If all other matrix elements in Eq. (31) are real then the phases α, β, γ which are equal to 0 or π, correspond to the CP-conserving case.

If all diagonal elements M_i ($i = 1,2,3$) are real and positive ($\alpha = \beta = \gamma = 0$) then the eigenvalues of the neutrino mass matrix are also positive and the CP symmetry is conserved if the CP parities of heavy neutrinos are the same and equal:

$$\eta_{CP}(N_4) = \eta_{CP}(N_5) = \eta_{CP}(N_6) = +i. \tag{34}$$

The same happens when all masses M_1, M_2, M_3 are real negative ($\alpha = \beta = \gamma = \pi$). Then CP is also conserved if neutrinos CP parities are negative imaginary:

$$\eta_{CP}(N_4) = \eta_{CP}(N_5) = \eta_{CP}(N_6) = -i. \tag{35}$$

In both cases above the mixing matrix elements $(K_{L,R})_{ae}$, $a = 4,5,6$ are either pure real for $\alpha = \beta = \gamma = 0$ or pure imaginary if $\alpha = \beta = \gamma = \pi$. We have checked that for the $e^-e^- \to W_i^- W_j^-$ processes the helicity amplitudes with $\Delta\sigma = 0$, proportional to the heavy neutrino masses, are dominant. These helicity amplitudes include either square of the K_L mixing matrix elements (for $\sigma_1 = \sigma_2 = -1/2$) or square of the K_R ones (for $\sigma_1 = \sigma_2 = +1/2$)—see Eqs. (A6)–(A8). That is why, summing the helicity amplitude over neutrino masses [Eq. (A5)], we get in both cases mentioned constructive contributions from all of them (solid line in Fig. 8).

In the case of mixing CP parities the CP symmetry is also conserved but the destructive interference between contributions from various neutrinos causes that the cross section decreases. In this case some of $(K_{L,R})_{ae}$, $a = 4,5,6$ matrix elements are real and some are imaginary.

To obtain these CP-violating effects several K_L or K_R matrix elements must interfere in the same helicity amplitude. The structure of the chosen neutrino mass matrix causes the $(K_L)_{ae}$, $a = 4,5,6$ matrix elements to be of a similar order,

$$\mid (K_L)_{4e} \mid \simeq \frac{1}{m_4} > \mid (K_L)_{5e} \mid \simeq \frac{1}{m_5} > \mid (K_L)_{6e} \mid \simeq \frac{1}{m_6}\ , \tag{36}$$

but only one suitable element of the K_R matrix is large:

$$\mid K_R \mid_{4e} \sim 1 \gg \mid (K_R)_{ae} \mid\ , \quad a = 5,6. \tag{37}$$

This property causes that only if the K_L matrix elements [Eq. (36)] contribute to the cross section in the visible way, can the CP breaking be seen. It is just the case for two light gauge bosons $W_1^- W_1^-$ production where $\cos^2\xi$ multiplies the K_L matrix contributions in $\sigma_1 = \sigma_2 = -1/2$ helicity amplitude. If the contribution from $\sigma_1 = \sigma_2 = +1/2$ helicity amplitude (where $\sin^2\xi$ is multiplied by $K_R^T K_R$) becomes important, then the CP-symmetry-breaking effect decreases. That is why the effect is visible only outside the δ_R^{--} resonance region where only one right-handed mixing matrix K_R gives an essential contribution [Eq. (37)] and the interference has no importance. It also means that the CP-symmetry breaking is more visible for the larger M_{W_2}, when $\sin\xi \to 0$. For the same reason the CP-breaking effect will not be seen in the heavy gauge boson production ($\cos^2\xi$ is multiplied by the K_R matrix contributions) and weakly visible in the light-heavy gauge bosons production ($\sin\xi\cos\xi$ multiplies $K_L^T K_L$).

If the CP symmetry is violated (phases α, β, $\gamma \neq 0, \pi$), the cross sections lie between two limiting lines in Fig. 8.

The effects of the CP violation in the standard model with the additional right-handed neutrinos are depicted in Fig. 9. For the same energy with the same mass matrix parametrization the effects are almost as large as those in the L-R symmetric models.

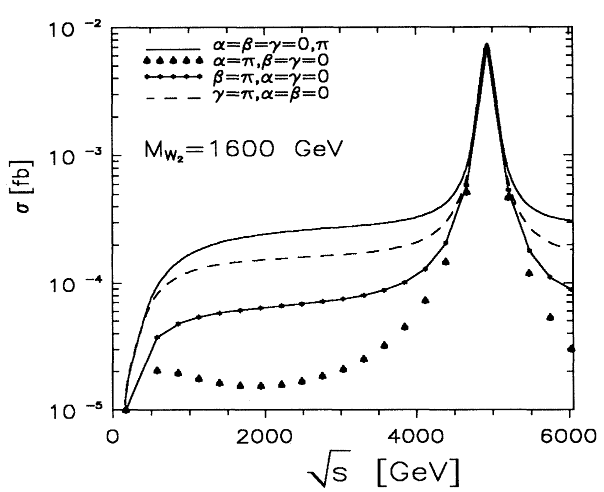

FIG. 8. Process $e^-e^- \to W_1^- W_1^-$ (LR model) for $M_1 = 200, M_2 = 400, M_3 = 600$ GeV [Eq. (32)] and $M_{W_2} = 1600$ GeV when CP parity is conserved in the lepton sector.

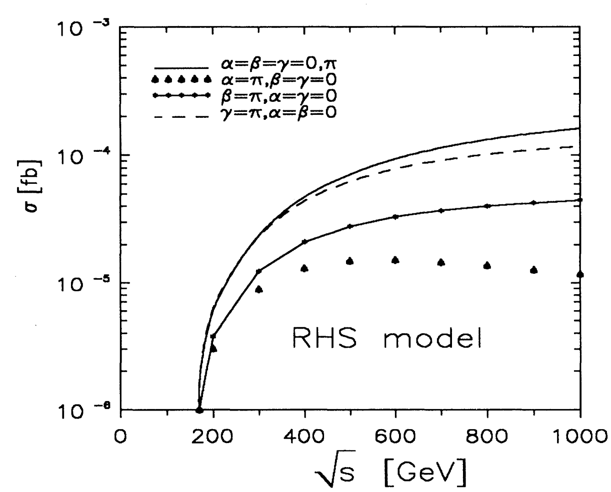

FIG. 9. Process $e^-e^- \to W_1^- W_1^-$ (RHS model) for $M_1 = 200, M_2 = 400, M_3 = 600$ GeV [Eq. (32)] when CP parity is conserved in the lepton sector.

In contrast with the case of the CP violation in the quark sector where the effect is small the violation of the CP symmetry in the lepton sector with the Majorana neutrino can be very large. The CP violation phases can change the cross section by even more than one order of the magnitude.

V. CONCLUSIONS

We have calculated the total cross section for various gauge bosons production processes in e^-e^- scattering. Exact calculations have been done in the frames of two versions of the left-right symmetric model and the standard model with additional right-handed neutrinos. We have checked the high energy behavior for the cross sections, $\lim_{s\to\infty} \sigma(s)$. Having calculated the helicity amplitudes we have been able to show which ingredients of the models are responsible for the correct high energy limit $\sigma(s \to \infty) = 0$. In the left-right symmetric model the unitarity is satisfied because there are s channels with doubly charged Higgs bosons $\delta^{--}_{L,R}$ and the left- and right-handed mixing matrices are orthogonal. In the standard model with additional RHS's but without Higgs triplets the correct high energy is guaranteed because the fragment of the neutrino mass matrix identically equals zero, $M_L = 0$. We have calculated this cross section taking into account existing bounds on the parameters of the considered models. Optimistically, with not very restrictive bounds on the models parameters, we have found that for $\sqrt{s} = 500$ GeV in the L-R symmetric model $\sigma\left(e^-e^- \to W_1^- W_1^-\right) \sim 0.01$ fb and for the standard model with a RHS $\sigma\left(e^-e^- \to W_1^- W_1^-\right) \sim 0.001$ fb. The small values of given cross sections are caused by the small mixing matrix elements K_{ae} between electron and heavy neutrinos calculated in our seesaw-type models. For nonseesaw models where mixing matrices $K_{L,R}$ are not calculated and are free parameters bounded only by existing data, the cross sections can be much higher.

In the L-R symmetric model the cross sections for

TABLE I. The reduced helicity amplitudes for all polarizations of the $l_m^- l_n^- \to W_i^- W_j^-$ process $[c = \cos\Theta\tilde{\eta} = \lambda_1\Delta\sigma,\ \eta = \Delta\sigma\Delta\lambda,\ \breve{\eta} = 2\sigma_1\lambda_1,\ \hat{\eta} = (\sigma_1+\sigma_2)(\lambda_1+\lambda_2)]$.

λ_1	λ_2	$\Delta\sigma = \pm 1$	$\Delta\sigma = 0$
		$M_t(\Delta\sigma;\lambda_1,\lambda_2)$	
1	1	$[(x-\tilde{\eta}y_1)+\frac{\tilde{\eta}-c}{2}]$	$\frac{1}{\sqrt{2}}(1+\breve{\eta}c)$
-1	-1		
1	0	$\gamma_2[(1-\eta\beta_2)(x-\eta y_1)+\eta-c]$	$\frac{1}{\sqrt{2}}\gamma_2\hat{\eta}(1+\hat{\eta}\beta_2)$
-1	0		
0	1	$\gamma_1[(1-\eta\beta_1)(x+\eta y_1)+\beta_1-c]$	$\frac{1}{\sqrt{2}}\gamma_1\hat{\eta}(1+\hat{\eta}\beta_1)$
0	-1		
1	-1	$-\frac{1}{\sqrt{2}}$	0
-1	1		
0	0	$\gamma_1\gamma_2 *$ $[y_1(\beta_2-\beta_1)+x(1-\beta_1\beta_2)+\beta_1-c]$	$-\frac{1}{\sqrt{2}}\gamma_1\gamma_2 *$ $[1+\beta_1\beta_2-c(\beta_1+\beta_2)]$
		$M_u(\Delta\sigma;\lambda_1,\lambda_2)$	
1	1	$-[(x-\tilde{\eta}y_2)+\frac{\tilde{\eta}+c}{2}]$	$\frac{1}{\sqrt{2}}(1-\breve{\eta}c)$
-1	-1		
1	0	$-\gamma_2[(1+\eta\beta_2)(x-\eta y_2)+\beta_2+c]$	$-\frac{1}{\sqrt{2}}\gamma_2\hat{\eta}(1+\hat{\eta}\beta_2)$
-1	0		
0	1	$-\gamma_1[(1+\eta\beta_1)(x+\eta y_2)-\eta+c]$	$-\frac{1}{\sqrt{2}}\gamma_1\hat{\eta}(1+\hat{\eta}\beta_1)$
0	-1		
1	-1	$-\frac{1}{\sqrt{2}}$	0
-1	1		
0	0	$-\gamma_1\gamma_2 *$ $[y_2(\beta_1-\beta_2)+x(1-\beta_1\beta_2)+\beta_2+c]$	$-\frac{1}{\sqrt{2}}\gamma_1\gamma_2 *$ $[1+\beta_1\beta_2+c(\beta_1+\beta_2)]$
		$M_s(\Delta\sigma;\lambda_1,\lambda_2)$	
1	1	0	$2\sqrt{2}$
-1	-1		
1	0	0	0
-1	0		
0	1	0	0
0	-1		
1	-1	0	0
-1	1		
0	0	0	$-2\sqrt{2}(1+\beta_1\beta_2)\gamma_1\gamma_2$

production of the non-standard gauge bosons $e^-e^- \to W_1^-W_2^-, W_2^-W_2^-$ are larger but then also the thresholds for these processes are higher. For $\sqrt{s}=1$ TeV the cross section $\sigma(e^-e^- \to W^-(80)W^-(600)) \sim 10$ fb and for $\sqrt{s}=1.5$ TeV $\sigma(e^-e^- \to W^-(600)W^-(600)) \sim 5$ pb.

We have also checked how the CP-violating parameters influence the total cross section. We have found that $\sigma_{\rm tot}$ is the biggest if CP is conserved and if all heavy neutrinos CP parities η_{CP} are the same, equal $+i$ or $-i$. If heavy neutrinos have mixing CP parities (for some $\eta_{CP}=+i$ and for others $\eta_{CP}=-i$) or if CP is violated then the cross section is smaller. The effect of the CP-symmetry breaking can be very large. This supports the statement that the size of the CP violation can be large if in the lepton sector the Majorana neutrinos are present.

ACKNOWLEDGMENTS

This work was supported by Polish Committee for Scientific Researches under Grant Nos. 2P30225206/93, PB 136/IF/95, and University of Silesia internal Grant.

APPENDIX: HELICITY AMPLITUDES FOR THE $l_m^-l_n^- \to W_iW_j$ PROCESS

Helicity amplitudes for two gauge bosons $W_i^-W_j^-$ $(i,j=1,2)$ production in two charged leptons $l_m^-l_n^-$ scattering processes $(m,n=e,\mu,\tau)$ are described generally by the Feynman diagrams in the s, t, and u channels (Fig. 10). The vertices for the t and u channels are determined by the charged current Lagrangian (α and β are spinor indices)

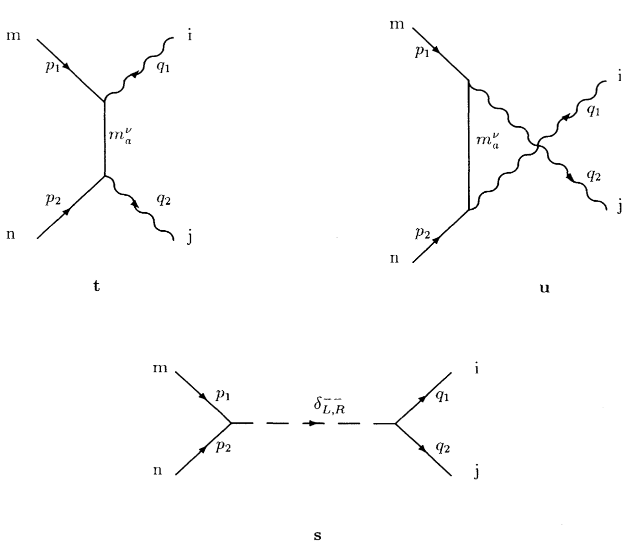

FIG. 10. The tree level Feynman diagrams which contribute to the $l_m^-l_n^- \to W_i^-W_j^-$ process $(m,n=e,\mu,\tau;$ $i,j=1,2)$. In the LR model all three channels are present, while in the RHS model there is no channel s.

$$L=\sum_{a=1,\dots 6,b=m,n}\left[\bar{N}_{a\alpha}\left(\Gamma_{ab}^{(1)\mu}\right)_{\alpha\beta}l_{b\beta}W_\mu^{(1)}\right.$$
$$\left.+\bar{N}_{a\alpha}\left(\Gamma_{ab}^{(2)\mu}\right)_{\alpha\beta}l_{b\beta}W_\mu^{(2)}\right], \qquad \text{(A1)}$$

where

$$\left(\Gamma_{ab}^{(1)\mu}\right)_{\alpha\beta}=\frac{g}{\sqrt{2}}\{\cos\xi(\gamma^\mu P_L)_{\alpha\beta}(K_L)_{ab}$$
$$-\sin\xi(\gamma^\mu P_R)_{\alpha\beta}(K_R)_{ab}\},$$
$$\left(\Gamma_{ab}^{(2)\mu}\right)_{\alpha\beta}=\frac{g}{\sqrt{2}}\{\sin\xi(\gamma^\mu P_L)_{\alpha\beta}(K_L)_{ab}$$
$$+\cos\xi(\gamma^\mu P_R)_{\alpha\beta}(K_R)_{ab}\}. \qquad \text{(A2)}$$

For the LR model the mixing angle ξ and 6×3 lepton

TABLE II. High energy behavior of the reduced helicity amplitudes for all polarizations of the $l_m^-l_n^- \to W_i^-W_j^-$ process [$\tilde\eta=\lambda_1\Delta\sigma$, $\breve\eta=2\sigma_1\lambda_1$, $\eta=\Delta\sigma\Delta\lambda, \hat\eta=(\sigma_1+\sigma_2)(\lambda_1+\lambda_2)$].

λ_1	λ_2	$\Delta\sigma=\pm1$	$\Delta\sigma=0$
		$M_t\,(\Delta\sigma;\lambda_1,\lambda_2)$	
1	1	$\frac{1}{2}(1-c)$	$\frac{1}{\sqrt{2}}(1+\breve\eta c)$
-1	-1		
1	0	$\frac{\sqrt{s}}{2M_2}(1-c)$	$\frac{\sqrt{s}}{2\sqrt{2}M_2}\hat\eta(1+\hat\eta\beta_2)$
-1	0		
0	1	$\frac{\sqrt{s}}{2M_1}(1-c)$	$\frac{\sqrt{s}}{2\sqrt{2}M_1}\hat\eta(1+\hat\eta\beta_1)$
0	-1		
1	-1	$-\frac{1}{\sqrt{2}}$	0
-1	1		
0	0	$\frac{s}{4M_1M_2}(1-c)$	$-\frac{s}{2\sqrt{2}M_1M_2}(1-c)$
		$M_u\,(\Delta\sigma;\lambda_1,\lambda_2)$	
1	1	$-\frac{1}{2}(1+c)$	$\frac{1}{\sqrt{2}}(1-\breve\eta c)$
-1	-1		
1	0	$-\frac{\sqrt{s}}{2M_2}(1+c)$	$-\frac{\sqrt{s}}{2\sqrt{2}M_2}\hat\eta(1+\hat\eta\beta_2)$
-1	0		
0	1	$-\frac{\sqrt{s}}{2M_1}(1+c)$	$-\frac{\sqrt{s}}{2\sqrt{2}M_1}\hat\eta(1+\hat\eta\beta_1)$
0	-1		
1	-1	$-\frac{1}{\sqrt{2}}$	0
-1	1		
0	0	$-\frac{s}{4M_1M_2}(1+c)$	$-\frac{s}{2\sqrt{2}M_1M_2}(1+c)$
		$M_s\,(\Delta\sigma;\lambda_1,\lambda_2)$	
1	1	0	$2\sqrt{2}$
-1	-1		
1	0	0	0
-1	0		
0	1	0	0
0	-1		
1	-1	0	0
-1	1		
0	0	0	$-\frac{\sqrt{2}s}{M_1M_2}$

mixing matrices $K_{L,R}$ are defined in [4], for example. In the RHS model with n_R right neutrino singlets $\xi = 0$, $K_R = 0$ and the K_L matrix has dimensions $(3+n_R, 3)$.

In the model where the doubly charged Higgs bosons δ^{++} exist (e.g., models with the Higgs triplets) the considered process has the s-channel diagram too, determined by the $\delta^{++}W^-W^-$ and the Yukawa Lagrangian couplings. In the LR model,

$$\begin{aligned}L_{HWW} = &-\frac{g^2 v_L}{\sqrt{2}}\delta_L^{++}\left(\cos^2\xi W_1^- W_1^- + \sin^2\xi W_2^- W_2^- + \sin 2\xi W_1^- W_2^-\right)\\&-\frac{g^2 v_R}{\sqrt{2}}\delta_R^{++}\left(\sin^2\xi W_1^- W_1^- + \cos^2\xi W_2^- W_2^- - \sin 2\xi W_1^- W_2^-\right) + \text{H.c.},\end{aligned} \tag{A3}$$

$$L_{\text{Yukawa}} = \sum_{a,m,n}\left(\frac{1}{\sqrt{2}v_L}\left[\delta_L^{++} l_m^T C P_L (K_L^T)_{ma} m_a (K_L)_{an} l_n\right] + \frac{1}{\sqrt{2}v_R}\left[\delta_R^{++} l_m^T C P_R (K_R^T)_{ma} m_a (K_R)_{an} l_n\right]\right) + \text{H.c.}$$

Four Feynman diagrams give contributions to the process $l_m^- l_n^- \to W_i^- W_j^-$ at the tree level. The differential cross section is given by

$$\frac{d\sigma_{mn}^{ij}(\sigma_1,\sigma_2;\lambda_1,\lambda_2)}{d\cos\Theta} = \frac{x}{16\pi s}\left| M(\sigma_1,\sigma_2;\lambda_1,\lambda_2)_{mn}^{ij}\right|^2 , \tag{A4}$$

where $\sigma_1(\sigma_2)$, $\lambda_1(\lambda_2)$ are the helicities of $m(n)$ fermions and $W_i(W_j)$ gauge bosons, respectively. The helicity amplitudes can be written in the form (Θ, ϕ are polar angles of W_i^- in the c.m. frame)

$$\begin{aligned}M\Big(\sigma_1,\sigma_2;\lambda_1,\lambda_2\Big)_{mn}^{ij} = &-\frac{e^2}{\sqrt{2}\sin^2\Theta_W} D^J_{\Delta\sigma,\,\Delta\lambda}(0,\Theta,\phi)\\&\times\left\{B^{ij}(\sigma_1,\sigma_2)M_t(\Delta\sigma;\lambda_1,\lambda_2)\,(R_t)_{mn} + B^{ij}(\sigma_2,\sigma_1)M_u(\Delta\sigma;\lambda_1,\lambda_2)\,(R_u)_{mn}\right.\\&\left.+B^{ij}(\sigma_1,\sigma_2)M_s(\Delta\sigma;\lambda_1,\lambda_2)\,(R_s)_{mn}\right\},\end{aligned} \tag{A5}$$

where

$$(R_t)_{mn} = \sum_a (K^T_{2\sigma_1})_{ma}\frac{\left(\frac{m_a}{\sqrt{s}}\right)^{1-|\Delta\sigma|}}{A + x\cos\Theta - \frac{m_a^2}{2s}}(K_{2\sigma_2})_{an}, \tag{A6}$$

$$(R_u)_{mn} = \sum_a (K^T_{2\sigma_2})_{ma}\frac{\left(\frac{m_a}{\sqrt{s}}\right)^{1-|\Delta\sigma|}}{A - x\cos\Theta - \frac{m_a^2}{2s}}(K_{2\sigma_1})_{an}, \tag{A7}$$

$$\begin{aligned}(R_s)_{mn} = &\,\delta_{\sigma_1,1/2}\frac{1}{1-\frac{M^2_{\delta_R}}{s}+i\frac{\Gamma_{\delta_R}M_{\delta_R}}{s}}\sum_a (K_R^T)_{ma}\left(\frac{m_a}{\sqrt{s}}\right)(K_R)_{an}\\&+\delta_{\sigma_1,-1/2}\frac{1}{1-\frac{M^2_{\delta_L}}{s}+i\frac{\Gamma_{\delta_L}M_{\delta_L}}{s}}\sum_a (K_L^T)_{ma}\left(\frac{m_a}{\sqrt{s}}\right)(K_L)_{an},\end{aligned} \tag{A8}$$

and

$$K_\kappa = \begin{cases} K_R & \text{for } \kappa = +1 \\ K_L & \text{for } \kappa = -1\end{cases},$$

$$\Delta\sigma = \sigma_1 - \sigma_2 \ , \quad \Delta\lambda = \lambda_1 - \lambda_2 \ , \quad J = \max(|\,\Delta\sigma\,|, |\,\Delta\lambda\,|),$$

$$A = \frac{M^2_{W_i} - M^2_{W_j}}{4s^2} - \frac{1+4x^2}{4},$$

$$x = \frac{k}{\sqrt{s}} \ , \quad k = \frac{1}{2\sqrt{s}}\sqrt{s^2 - 2s(M_1^2 + M_2^2) + (M_1^2 - M_2^2)^2}.$$

The factors B are different for various gauge boson productions

$$B^{ij}\,(\sigma_1,\sigma_2) = \begin{cases} -2\sigma_1(\sin\xi)^{||\Delta\sigma|-2\sigma_2|}(\cos\xi)^{||\Delta\sigma|-2\sigma_1|}\ , & i=1\ ,\ j=2, \\ 4\sigma_1\sigma_2(\sin\xi)^{(1+\sigma_1+\sigma_2)}(\cos\xi)^{(1-\sigma_1-\sigma_2)}\ , & i=1\ ,\ j=1, \\ (\sin\xi)^{(1-\sigma_1-\sigma_2)}(\cos\xi)^{(1+\sigma_1+\sigma_2)}\ , & i=2\ ,\ j=2. \end{cases} \quad \text{(A9)}$$

The reduced helicity amplitudes $M_{t(u,s)}\,(\Delta\sigma;\lambda_1,\lambda_2)$ are gathered in Table I, where we use the notation

$$y_{1(2)} = \frac{E_{1(2)}}{\sqrt{s}}\ ,\ \beta_{1(2)} = \frac{k}{E_{1(2)}}\ ,\ \gamma_{1(2)} = \frac{E_{1(2)}}{M_{1(2)}},$$

and

$$E_{1(2)} = \frac{s + M_{2(1)}^2 - M_{1(2)}^2}{2\sqrt{s}},$$

where $E_1(E_2)$ and $M_1(M_2)$ are energies and masses of 1(2) gauge bosons, respectively. The high energy behavior of the helicity amplitudes are given in Table II. It is worthwhile to notice several interesting properties of the helicity amplitudes [Eq. (A5)]: if $M_L = 0$ then δ_L^{--} does not contribute to the process [Eqs. (12),(A8)]; the δ_L^{--} (δ_R^{--}) in the s channel contributes only to the amplitude with the electron helicities $-1/2(+1/2)$; although there are no Majorana neutrinos in the s channel the amplitude in this channel is also proportional to the neutrino masses; the helicity amplitudes with $\Delta\sigma = 0$ are proportional to the neutrino masses m_a.

[1] T. Rizzo, Phys. Lett. **116B**, 23 (1982).

[2] D. London, G. Belanger, and J.N. Ng, Phys. Lett. B **188**, 155 (1987); J. Maalampi, A. Pietilä, and J. Vuori, Nucl. Phys. **B381**, 544 (1992); Phys. Lett. B **297**, 327 (1992); T. Rizzo, in *Proceedings of the Workshop on Physics and Experiments with Linear* e^+e^- *Colliders*, Waikoloa, Hawaii, 1993, edited by F. Harris *et al.* (World Scientific, Singapore, 1994), p. 520; C.A. Heusch and P. Minkowski, Nucl. Phys. **B416**, 3 (1994); P. Helde, K. Huitu, J. Maalampi, and M. Raidal, *ibid.* **B437**, 305 (1995).

[3] F. del Aquilla and M. Zrałek, "CP violation in the lepton sector with Majorana neutrinos," Report No. UG-FT-51/95,SU-TP-2/95 (unpublished).

[4] N. G. Deshpande, J.F. Gunion, B. Kayser, and F. Olness, Phys. Rev. D **44**, 837 (1991); J. Gluza and M. Zrałek, *ibid.* **48**, 5093 (1993); **51**, 4695 (1995).

[5] Particle Data Group, L. Montanet *et al.*, Phys. Rev. D **50**, 1173 (1994).

[6] L3 Collaboration, O. Adriani *et al.*, Phys. Lett. B **295**, 371 (1992).

[7] G. Gelmini and E. Roulet, "Neutrino masses," Report No. UCLA/94/TEP/36, 1994 (unpublished).

[8] W. Buchmuller and C. Greub, Nucl. Phys. **B363**, 345 (1991); G. Ingelman and J. Rathsman, Z. Phys. C **60**, 243 (1993).

[9] F. Abe *et. al.*, Phys. Rev. Lett. **67**, 2609 (1991).

[10] G. Beall, M. Bander, and A. Soni, Phys. Rev. Lett. **48**, 848 (1982).

[11] D. Wyler and L. Wolfenstein, Nucl. Phys. **B218**, 205 (1983); R.N. Mohapatra and J.W.F. Valle, Phys. Rev. D **34**, 1642 (1986); E. Witten, Nucl. Phys. **B268**, 79 (1986); J. Bernabeu *et al.*, Phys. Lett. B **187**, 303 (1987); J. L. Hewett and T.G. Rizzo, Phys. Rep. **183**, 193 (1989); P. Langacker and D. London, Phys. Rev. D **38**, 907 (1988); E. Nardi, *ibid.* **48**, 3277 (1993); D. Tommasini *et al.*, Nucl. Phys. **B444**, 451 (1995).

[12] L.N. Chang, D. Ng, and J.N. Ng, Phys. Rev. D **50**, 4589 (1994).

[13] J. Gluza and M. Zrałek, "On possibility of detecting the $e^-e^- \to W^-W^-$ process in the standard model with additional right-handed neutrinos," Report No. TP-USL 195104, hep-ph19507269 (unpublished).