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# Single $W_R$ Production in $e^-e^-$ Collisions at the NLC\*

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## Abstract

Single  $W_R$  production in  $e^-e^-$  collisions at the NLC can be used to probe the Majorana nature of the heavy neutrinos present in the Left-Right Symmetric Model below the kinematic threshold for their direct production. For colliders in the  $\sqrt{s} = 1 - 1.5$  TeV range, typical cross sections of order  $1 - 10fb$  are obtained, depending on the specific choice of model parameters. Backgrounds arising from Standard Model processes are shown to be small. This analysis greatly extends the kinematic range of previous studies wherein the production of an on-shell, like-sign pair of  $W_R$ 's at the NLC was considered.

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One of the most attractive explanations for the apparently small magnitude of neutrino masses is the see-saw mechanism[1] which can be easily implemented within the framework of extended electroweak gauge theories(EEGT). Such a scheme naturally leads to the prediction that the ordinary Standard Model(SM) neutrinos are Majorana particles and that there must also exist heavy, neutral,  $SU(2)_L$ -isosinglet fields ( $N$ ) which are also Majorana particles. Due to the Majorana nature of both sets of neutrinos, one expects that the Lagrangian containing the corresponding mass terms can result in the existence of new lepton-number violating,  $\Delta L = 2$  interactions at a level that might be experimentally observable. At low energies, the best example of such an interaction is neutrinoless double- $\beta$  decay [2] which has been sought for some time. In fact, existing limits from searches for such decays must be used as input by model builders to construct consistent scenarios for neutrino masses. At higher energies, these  $\Delta L = 2$  interactions can manifest themselves in many ways, *e.g.*, once produced, the  $N$ 's decay to charged leptons of both signs with equal rates. One interesting possibility[3], that has recently been revitalized[4], is the production of a like-charged pair of gauge bosons at an  $e^-e^-$  collider. As has been much emphasized, high energy  $e^-e^-$  collisions in the 0.5 – 1.5 TeV range are a possible option at the NLC[5]. As in ordinary neutrinoless double- $\beta$  decay( $\beta\beta_{0\nu}$ ), this reaction can *only* occur if the massive neutrinos are Majorana particles, as lepton number is violated by two units. In some sense, this reaction is really just the *inverse* process to  $\beta\beta_{0\nu}$ .

The Left-Right Symmetric Model(LRM)[6], which is based on the extended weak gauge group  $SU(2)_L \times SU(2)_R \times U(1)$ , provides a very natural setting for the above scenario wherein the  $N$ 's are identified with heavy right-handed neutrinos and  $SU(2)_R$ -breaking occurs via isotriplet scalars. In this model, the cross section for  $e^-e^- \rightarrow W_L^-W_L^-$  is quite small (due to the fact that left-handed neutrino masses are tiny and the various mixing angles are small) and so we are left considering the process  $e^-e^- \rightarrow W_R^-W_R^-$  as was done in Refs.[3, 4].

One difficulty, however, is that it may not be possible to produce two, on-shell  $W_R$ 's at the NLC with a center of mass energy in the range  $0.5 \leq \sqrt{s} \leq 1.5$  TeV as  $W_R$ 's are generally expected to be rather heavy. It may, however, be possible to examine the corresponding single-production process  $e^-e^- \rightarrow W_R^-(W_R^-)^* \rightarrow W_R^- + jj$  at the NLC, since it only requires the  $W_R$  to have a mass somewhat below  $\sqrt{s}$  for this final state to be kinematically accessible. If the resulting cross section was found to be sufficiently large, this would greatly extend the capability of the NLC to probe the Majorana nature of the heavy neutrinos in this model as well as to examine the interplay between the  $SU(2)_R$  symmetry breaking and Majorana mass generating mechanisms. Of course, to produce only one on-shell  $W_R$  we must pay the price of the square of an additional electroweak gauge coupling, as well as three-body phase space, both of which result in a corresponding rate reduction in comparison to the two-body  $W_R^-W_R^-$  process. However, since typical  $e^-e^- \rightarrow W_R^-W_R^-$  cross sections were found to be on the order of  $1pb$  or more, this additional cost may not be too high if integrated luminosities in the  $100fb^{-1}$  range are obtainable at the NLC for center of mass energies at or above 1 TeV. In addition, this single-production process is expected to be rather background-free since the on-shell  $W_R$  can be easily reconstructed from its decay products. It is the purpose of this paper to calculate the cross section for this single-production process at NLC energies and determine the corresponding event rates given the range of anticipated NLC luminosities. Note that this process will allow us to probe the Majorana nature of  $N$  even though we may be below threshold for its direct production.

Constraints on the mass of the  $W_R$  ( $M_R$ ) arise from many sources but are subject to various different assumptions about the magnitudes of the undetermined parameters of the LRM[7]:  $\kappa = g_R/g_L$ , the ratio of the  $SU(2)_{L,R}$  gauge couplings, the elements of the right-handed Cabibbo-Kobayashi-Maskawa mixing matrix,  $V_R$ , and the masses for the right-handed neutrinos,  $M_N$ . The most important of the existing constraints on  $M_R$  are all severely

weakened if we are free to allow these unknown parameters to vary significantly from such simplifying assumptions as  $\kappa = 1$ ,  $V_R = V_L$ , and that the  $N$ 's are light. For example, the polarized  $\mu$ -decay bounds[8] (originating from both  $K$  and  $\pi$  decays) are trivially avoided if  $M_N > 50$  MeV. Likewise, the Tevatron bounds[9] can be easily satisfied if *any* of the following are true: (i) the  $N$ 's decay in the detector, (ii) if the  $W_R$ 's were to decay to many non-SM final states, or (iii) if  $(V_R)_{ud}$  were much smaller than that suggested by the  $V_R = V_L$  relationship. The well-known limit on  $M_R$  from the  $K_L - K_S$  mass difference[10] is also greatly weakened if  $V_R = V_L$  is not assumed and  $\kappa < 1$ . (Grand unified models generally require[11] that  $\kappa \leq 1$  while consistency of the couplings in the Left-Right Model requires  $\kappa^2 \geq \frac{x_w}{1-x_w} \simeq 0.303$ , where  $x_w = \sin^2\theta_w$ . Based on naturalness assumptions alone, we might expect that  $\kappa$  does not differ from unity by more than a factor of two.) Putting all this together one finds that  $M_R > 300$  GeV[12], in agreement with Ref. [7], but we should anticipate that  $W_R$ 's may be significantly heavier than this weak lower bound would indicate, *i.e.*, it is more than likely that if  $W_R$ 's do exist their masses should be at least several times larger than this bound suggests. Clearly,  $W_R$ 's more massive than about 700 GeV would be too heavy to pair produce (on shell) at a  $\sqrt{s} = 1.5$  TeV collider with a significant cross section; this forces us to consider the single production scenario. We remind the reader that the mass of the  $Z'$  in this model is highly correlated with that of the  $W_R$  and is a decreasing function of  $\kappa$ ; we find  $M_{Z'}/M_R = 3.55(1.69, 1.47)$  for  $\kappa = 0.6(1, 2)$ .

An interesting case to consider is the possibility that a  $W_R$  will not be discovered until the NLC turns on, *i.e.*, the LHC does not see  $W_R$ 's. This can happen in the following way. Several analyses have shown[13] that  $W_R$ 's should be observable at the LHC in the mass range of our interest for all reasonable values of  $\kappa$  and the elements of  $V_R$  *provided* the decay  $W_R \rightarrow eN$  is kinematically allowed. Essentially, the reason for this is that both the machine and parton luminosities are sufficient large (for  $W_R$ 's with masses of order 1 TeV)

that reduced effective couplings or leptonic branching fractions will not prevent the  $W_R$ 's observation even if  $N$ 's are allowed to decay within the detector volume. However, it may be that  $M_N > M_R$  so that the  $W_R$  possesses *no* leptonic decay modes. In this case searches in the dijet final state at the LHC would be necessary and, due to the tremendous QCD backgrounds, it is unlikely that a  $W_R$  could be discovered by such searches[14]. (Of course, a  $W_R$  might be observable in the dijet mode *provided* its mass were already known from measurements in other channels and sufficiently good mass resolution was available.) If such a scenario were realized, the NLC would play the role of discovery machine for  $W_R$ 's, most likely in the  $e^+e^-$  collider mode.

Unfortunately, to calculate the  $e^-e^- \rightarrow W_R^-(W_R^-)^* \rightarrow W_R^- + jj$  cross section at the NLC, an additional free parameter is introduced in the form of the doubly-charged Higgs scalar( $\Delta$ ) mass,  $M_\Delta$ . The reason for this is that the  $W_R^-(W_R^-)^*$  state is produced by  $N$ -exchange in the  $t$ - and  $u$ -channels together with a  $\Delta$ -exchange in the  $s$ -channel. All three contributions are required to maintain unitarity via gauge cancellations as discussed in Ref.[3, 4]. While the couplings of the  $\Delta$  to both  $e^-e^-$  and  $W_R^-W_R^-$  are fixed by the  $SU(2)_R$  gauge symmetry breaking, the value of  $M_\Delta$  remains a free parameter in analogy to the Higgs boson mass in the SM. Thus the set of parameters we must consider when calculating the cross section are  $\kappa$ , and  $M_{N,R,\Delta}$ . Fortunately, the cross section itself scales with an overall factor of  $\kappa^6$ , which helps simplify our results.

Denoting the incoming  $e^-$  momenta by  $p_{1,2}$ , the out-going, on-shell  $W_R$  momenta by  $P$ , and the final state fermion momenta by  $q_{1,2}$ , the spin-summed, squared matrix element for  $e^-e^- \rightarrow W_R^-(W_R^-)^* \rightarrow W_R^- + f\bar{f}'$  is given by

$$|\bar{\mathcal{M}}|^2 = 4N_c \kappa^6 (g^2 M_N/4)^2 \left(\frac{g}{2}\right)^2 \text{Tr}(\not{p}_2 \Gamma_{\mu\nu} \not{p}_1 \Gamma_{\nu'\mu'}) \left(-g^{\mu\mu'} + \frac{P^\mu P^{\mu'}}{M_R^2}\right) \left(g^{\nu\lambda} - \frac{k^\nu k^\lambda}{M_R^2}\right) \quad (1)$$

$$\left(g^{\nu'\lambda'} - \frac{k^{\nu'} k^{\lambda'}}{M_R^2}\right) [(k^2 - M_R^2)^2 + (\Gamma_R M_R)^2]^{-1} (q_{1\lambda} q_{2\lambda'} + q_{2\lambda} q_{1\lambda'} - g_{\lambda\lambda'} q_1 \cdot q_2),$$

where  $k = q_1 + q_2$ ,  $g = g_L$ ,  $N_c$  is a color factor, and  $\Gamma_R$  is the width of  $W_R$ . Here,  $\Gamma_{\mu\nu}$  is given by

$$\Gamma_{\mu\nu} = \frac{\gamma_\nu \gamma_\mu}{t - M_N^2} + \frac{\gamma_\mu \gamma_\nu}{u - M_N^2} + \frac{4g_{\mu\nu}}{s - M_\Delta^2}, \quad (2)$$

with  $s = (p_1 + p_2)^2$ ,  $t = (p_1 - P)^2$ , and  $u = (p_2 - P)^2$ . Assuming massless final state fermions, we integrate over their momenta and sum over all possible flavour and color combinations leading to the differential cross section for the  $jj$  final state:

$$\frac{d\sigma}{dz} = 9 \frac{(G_F M_W^2/\pi)^3}{24} \frac{\kappa^6 M_N^2}{s} \int_\delta^{1+\delta^2/4} dx \frac{1-x+\delta^2/4}{x^2-\delta^2} \frac{1-x+\delta^2/4}{(1-x)^2 + (\Gamma_R M_R/s)^2} R, \quad (3)$$

with  $M_W$  being the SM  $W$  mass,  $z = \cos\theta$ , and  $(x, \delta) = 2(E, M_R)/\bar{s}$ , where  $E$  is the energy of the on-shell  $W_R$ . We define the angle  $\theta$  to be that between the three-vector components of  $P$  and  $p_1$ . It is important to note that the cross section is directly proportional to  $M_N^2$ , thus it vanishes as the Majorana mass of the heavy neutrino tends to zero. This is as expected since the reaction we're considering is a  $|\Delta L| = 2$  process. The expression for  $R$  is rather complicated; let us first define the following combinations of kinematic variables in order to simplify the various contributions that appear below:

$$k^2 = s + M_R^2 - 2E \bar{s},$$

$$\Sigma = s - M_R^2 - k^2,$$

$$\begin{aligned}
s_\Delta &= s - M_\Delta^2, \\
t_{R,N} &= t - M_{R,N}^2, \\
u_{R,N} &= u - M_{R,N}^2, \\
G_\Delta &= M_\Delta \Gamma_\Delta, \\
f &= (k^2 M_R^2)^{-1},
\end{aligned} \tag{4}$$

with  $\Gamma_\Delta$  being the width of the  $\Delta$  which we obtain by summing over the  $e^-e^-$  and  $W_R^-W_R^-$  decay modes. In terms of the  $W_R$  energy and scattering angle, the kinematic quantities  $t$  and  $u$  are given by

$$\begin{aligned}
t &= -\bar{s}E(1 - \beta z) + M_R^2, \\
u &= -\bar{s}E(1 + \beta z) + M_R^2, \\
\beta &= \frac{\sqrt{x^2 - \delta^2}}{x}.
\end{aligned} \tag{5}$$

We now can write  $R$  as

$$\begin{aligned}
R &= a + b + c, \\
a &= 16s(s_\Delta^2 + G_\Delta^2)^{-1}(2 + f\Sigma^2/4), \\
b &= 8s_\Delta(s_\Delta^2 + G_\Delta^2)^{-1}[s(2 + f\Sigma^2/4)(t_N^{-1} + u_N^{-1}) + f\Sigma(t_N^{-1} - u_N^{-1})(t_R^2 - u_R^2)/4], \\
c_1 &= s[(t_N^{-1} + u_N^{-1})^2 + 4t_N^{-1}u_N^{-1}], \\
c_2 &= f(k^2 + M_R^2)u_R t_R (t_N^{-2} + u_N^{-2}), \\
c_3 &= f\Sigma[(t_R t_N^{-1})^2 + (u_R u_N^{-1})^2], \\
c_4 &= f\Sigma(s\Sigma - t_R^2 - u_R^2)u_N^{-1}t_N^{-1}, \\
c &= c_1 + c_2 + c_3 + c_4.
\end{aligned} \tag{6}$$

Here, ‘ $a$ ’ arises from the pure  $s$ -channel  $\Delta$  exchange, ‘ $c$ ’ is the summed contribution of both

the  $u$ - and  $t$ -channel  $N$  exchanges, while ‘ $b$ ’ is the interference between the  $s$ - and  $u, t$ -channels. As expected, the differential cross section is symmetric under the interchange of  $u$  and  $t$ . The angular distribution itself is generally found to be quite flat owing to the rather large masses involved in the propagators and the unitarity cancellation among the three exchanges. (This lack of sensitivity to  $z$  is found to be essentially independent of the choice of particle masses so long as we restrict ourselves to parameter space regions that yield large cross sections.) This implies that mild acceptance cuts will not lead to any significant alterations in the event rates we obtain below. This will be shown explicitly after a brief discussion of the total cross section for  $W_R + jj$  production.

Integrating over the  $W_R$  production angle yields the total event rates found in Figs. 1 and 2, in which we have set  $\kappa = 1$  and scaled by an integrated luminosity of  $100 fb^{-1}$ . Fig. 1a shows the number of expected  $W_R + jj$  events, as a function of  $M_R$ , at a  $\sqrt{s} = 1$  TeV  $e^-e^-$  collider for different choices of  $M_N$  and  $M_\Delta$ . The results are seen to be quite sensitive to the values of these mass parameters even when  $M_R$  is fixed. In Fig. 1b(c), we fix  $M_R = 700$  GeV and plot the event rate as a function of  $M_N(M_\Delta)$  for various values of  $M_\Delta(M_N)$ . Typically, we see event rates of order several hundred/yr except near the  $\Delta$  resonance (where very large rates are obtained) or when  $M_N$  is small (as the cross section vanishes for massless  $N$  since it probes the  $N$ ’s Majorana nature). Increasing  $\sqrt{s}$  to 1.5 TeV, as shown in Fig. 2a, we see substantial cross sections are obtainable even assuming  $W_R$ ’s in the 1-1.2 TeV mass range for some parameter choices. Fixing  $M_R = 1$  TeV in Figs. 2b and c, we again see reasonable event rates for most choices of  $M_N$  and  $M_\Delta$  assuming  $\sqrt{s} = 1.5$  TeV. The exact rate is, however, a sensitive probe of both the  $N$  and  $\Delta$  masses. It is important to remember that the  $\Delta$  can appear as a resonance in the  $e^-e^-$  channel.

To verify our claim that the  $W_R$  angular distribution is quite flat for choices of  $M_N$  and  $M_\Delta$  which yield large cross sections, we show in Fig. 3 the angular distribution of a



$W_R$  with a mass of 700 GeV at a  $\sqrt{s} = 1$  TeV  $e^-e^-$  collider. For most choices of the input masses we obtain extremely flat distributions, however, when  $N$  is light a significant angular dependence is observed. This is simply a result of the  $t$ - and  $u$ - channel poles which develop as  $M_N$  tends to zero. Of course, small  $M_N$  also leads to a small cross section, as shown in Figs. 1 and 2, as might be expected since the matrix element vanishes in this massless limit. This substantiates our claim above that when the cross section is large the corresponding angular distribution is flat.

Potential backgrounds to the process  $e^-e^- \rightarrow W_R^-(W_R^-)^* \rightarrow W_R^- + jj$  at the NLC are easily controlled and/or removed. One might imagine, for example, some contamination from the SM process  $e^-e^- \rightarrow W_L^-W_L^- \nu\nu$ , but this can be easily eliminated by using missing energy cuts and demanding that the  $W_R$  final state be reconstructed from either the  $jj$  or  $eN \rightarrow eejj$  decay modes. (Since the on-shell  $W_R$  decays to either  $jj$  or  $eN \rightarrow eejj$  there is no missing energy in the signal process.) In addition, with polarized beams, we can take advantage of the fact that  $W_R$  couples via right-handed currents while any SM background must arise only via left-handed currents. Within the LRM itself a possible background could arise from a similar lepton-number conserving processes such as  $e^-e^- \rightarrow W_R^-W_R^- NN$ . The rate for such a process would be highly suppressed since there are several additional powers of the weak coupling and there is now a five-body final state for one virtual  $W_R$ . In addition, the  $N$ 's are quite massive implying that such final states are most likely to be kinematically inaccessible. Even if such a final state could be produced, in comparison to the process we are considering, the subsequent  $N$  decays would lead to a final state with too many charged leptons and/or jets.

In this paper we have addressed the following points:

- (i) While  $e^-e^- \rightarrow W_R^-W_R^-$  is an excellent probe of both the Majorana nature of  $N$  and the symmetry breaking sector of the Left-Right Symmetric Model, it is more than likely

that  $W_R$ 's are too massive to be pair produced at the NLC if  $\sqrt{s} = 1 - 1.5$  TeV. This forces us to consider the production of a single on-shell  $W_R$  via the process  $e^-e^- \rightarrow W_R^-(W_R^-)^* \rightarrow W_R^- + jj$ .

(ii) Since the pair of on-shell  $W_R$ 's cross section was generally very large, we would expect that the single  $W_R$  rate would be significant if integrated luminosities in the  $100 fb^{-1}$  range were available. From the explicit calculations we found that these expectations were indeed valid for most of the model parameter space yielding cross sections of order  $1 - 10 fb^{-1}$ .

(iii) For values of the input parameters that lead to significant rates the  $W_R$  angular distribution was found to be rather flat implying that angular cuts will not significantly reduce the cross sections. The rates themselves were found to be quite sensitive to the particular values of the masses of  $N$  and  $\Delta$ . Masses for both these particles beyond the kinematic reach of the NLC were found to be probed by the single  $W_R$  production process.

(iv) The process  $e^-e^- \rightarrow W_R^-(W_R^-)^* \rightarrow W_R^- + jj$  can be used to probe the Majorana nature of the heavy neutrinos present in the Left-Right Symmetric Model even when they are too massive to be directly produced.

The NLC can provide an excellent probe into the detailed structure of the Left-Right Symmetric Model.

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### Figure Captions (Contact author for figures.)

Figure 1. Event rates per  $100fb^{-1}$  for  $W_R + jj$  production at a 1 TeV  $e^-e^-$  collider assuming  $\kappa = 1$  (a) as a function of  $M_R$  for  $M_N = M_\Delta = 1$  TeV (dots),  $M_\Delta = 1.2$  TeV and  $M_N = 0.4$  TeV (dashes),  $M_\Delta = 0.3$  and  $M_N = 0.1$  TeV (dash-dots),  $M_\Delta = 2$ ,  $M_N = 0.6$  TeV (solid), or  $M_\Delta = 1.8$  and  $M_N = 0.6$  TeV (square dots); (b) with  $M_R = 700$  GeV fixed as a function of  $M_N$  for  $M_\Delta = 0.3(0.6, 1.2, 1.5, 2)$  TeV corresponding to the dotted(dashed, dash-dotted, solid, square-dotted) curve; (c) as a function of  $M_\Delta$  for  $M_N = 0.2(0.5, 0.8, 1.2, 1.5)$  TeV corresponding to the dotted(dashed, dash-dotted, solid, square-dotted) curve.

Figure 2. Same as Fig. 1, but for a 1.5 TeV  $e^-e^-$  collider. In (b) and (c), a  $W_R$  mass of 1 TeV is assumed.

Figure 3. Angular distribution for a  $W_R$  of mass 700 GeV produced at a  $\sqrt{s} = 1$  TeV  $e^-e^-$  collider for the same parameter choices as in Fig. 1a.