Tests of Anomalous Quartic Couplings at the NLC

O. J. P. Éboli¹*, M. C. Gonzalez-Garcia¹,²,†, and J. K. Mizukoshi³,‡

¹ Instituto de Física Teórica, Universidade Estadual Paulista
Rua Pamplona, 145, 01405–900 – São Paulo, Brazil

² Instituto de Física Corpuscular – IFIC/CSIC, Departament de Física Teòrica
Universitat de València, 46100 Burjassot, València, Spain

³ Instituto de Física, Universidade de São Paulo
C. P. 66.318, 05315–970 – São Paulo, SP, Brazil.

Abstract

We analyze the potential of the Next Linear $e^+e^-$ Collider to study anomalous quartic vector–boson interactions through the processes $e^+e^- \rightarrow W^+W^-Z$ and $ZZZ$. In the framework of $SU(2)_L \otimes U(1)_Y$ chiral Lagrangians, we examine all effective operators of order $p^4$ that lead to four-gauge-boson interactions but do not induce anomalous trilinear vertices. In our analysis, we take into account the decay of the vector bosons to fermions and evaluate the efficiency in their reconstruction. We obtain the bounds that can be placed on the anomalous quartic interactions and we study the strategies to distinguish the possible couplings.

*Email: eboli@ift.unesp.br; † Email: concha@axp.ift.unesp.br; ‡ Email: mizuka@fma.if.usp.br
I. INTRODUCTION

The impressive agreement of the Standard Model (SM) predictions for the fermion–vector boson couplings with the experimental results has been a striking confirmation of the $SU_L(2) \times U_Y(1)$ gauge structure of the model in that sector \cite{1}. However, we still lack the same accuracy tests for the structure of the bosonic sector. If the gauge and symmetry breaking sectors are invariant under the $SU_L(2) \times U_Y(1)$ gauge group, the structure of the triple and quartic vector-boson is completely determined. Thus a detailed study of these interactions can either confirm the local gauge invariance of the theory or indicate the existence of new physics beyond the SM.

Presently, we have only started to probe directly the triple gauge–boson couplings at the Tevatron \cite{2,3} and LEP \cite{4} through the production of pairs of vector bosons. Notwithstanding, the constraints on these couplings are still very loose. Future hadron \cite{5} and $e^+e^-$, $e\gamma$, and $\gamma\gamma$ \cite{6} colliders will provide further information on these couplings and improve significantly our knowledge of possible anomalous gauge-boson interactions.

If the $SU(2)_L \otimes U(1)_Y$ symmetry of the model is to be linearly realized, these studies of the triple gauge–boson couplings will be able to furnish information on the gauge–boson four–point functions provided that dimension 8 and higher anomalous operators are suppressed. This is the case when the breaking of the $SU(2)_L \otimes U(1)_Y$ symmetry takes place via the Higgs mechanism with a relatively light elementary Higgs boson. If, on the other hand, no fundamental light Higgs particle is present in the theory, one is led to consider the most general effective Lagrangian which employs a nonlinear representation of the broken $SU(2)_L \otimes U(1)_Y$ gauge symmetry. In this case the relation between the structure of the three– and four–point functions of the gauge bosons does not hold already at $p^4$ order, leaving open the question of the structure of the quartic vector–boson interactions.

At present the only information on quartic gauge–boson interactions is obtained indirectly as they modify the gauge–boson two–point functions at one loop \cite{7}. The precise electroweak measurements both at low energy and at the $Z$ pole, constrains the quartic anomalous
couplings to be smaller than $10^{-3} - 10^{-1}$ depending on the coupling.

Direct studies of quartic vector–boson interactions cannot be performed at the present colliders since the available centre-of-mass energy is not high enough for multiple vector–boson production. This crucial test of the gauge structure of the SM will only be possible at the CERN Large Hadron Collider (LHC) through the reaction $pp \rightarrow V_L V_L X$ or at the next linear collider (NLC) through the processes $e^+ e^- \rightarrow VVV$ [10], $e^+ e^- \rightarrow FFVV$ [14], $e^+ e^- \rightarrow FFVV$ [13], $e \gamma \rightarrow VVF$ [16], $\gamma \gamma \rightarrow VV$ [17], and $\gamma \gamma \rightarrow VVV$ [18], where $V = Z, W^\pm$ or $\gamma$ and $F = e$ or $\nu_e$.

In this work we analyze in detail the processes $e^+ e^- \rightarrow W^+ W^- Z$ and $ZZZ$ in order to assess the potential of the NLC, with and without polarized beams, to study anomalous quartic couplings of vector bosons. These reactions will be the most important processes to study the quartic gauge couplings at the NLC up to energies of the order of 1 TeV, where the processes $e^+ e^- \rightarrow VVF$ start to become important [19]. We work in the framework of chiral Lagrangians, and we study all $p^4$ operators that lead genuine quartic gauge interactions, i.e. these operators do not give rise to triple gauge–boson vertices, and consequently are not bounded by the study of the production of gauge–boson pairs. We extend the analysis of Ref. [12] for the custodial $SU(2)_C$ conserving operators taking into account realistic cuts and detection efficiencies. Moreover, we also study the non-conserving $SU(2)_C$ interactions.

II. THEORETICAL FRAMEWORK

If the electroweak symmetry breaking is due to a heavy (strongly interacting) Higgs boson, which can be effectively removed from the physical low–energy spectrum, or to no fundamental Higgs scalar at all, one is led to consider the most general effective Lagrangian which employs a nonlinear representation of the broken $SU(2)_L \otimes U(1)_Y$ gauge symmetry [20]. The resulting chiral Lagrangian is a non–renormalizable non–linear $\sigma$ model coupled in a gauge–invariant way to the Yang–Mills theory. This model independent approach in-
corporates by construction the low-energy theorems [21], that predict the general behavior of Goldstone boson amplitudes irrespective of the details of the symmetry breaking mechanism. Notwithstanding, unitarity implies that this low–energy effective theory should be valid up to some energy scale smaller than $4\pi v \simeq 3$ TeV [22], where new physics would come into play.

To specify the effective Lagrangian one must first fix the symmetry breaking pattern. We consider that the system presents a global $SU(2)_L \otimes SU(2)_R$ symmetry that is broken to $SU(2)_C$. With this choice, the building block of the chiral Lagrangian, in the notation of Ref. [20], is the dimensionless unimodular matrix field $\Sigma(x)$, which transforms under $SU(2)_L \otimes SU(2)_R$ as $(2, 2)$:

$$\Sigma(x) = \exp \left( i \frac{\varphi^a(x) \tau^a}{v} \right). \quad (1)$$

The $\varphi^a$ fields are the would-be Goldstone fields and $\tau^a$ ($a = 1, 2, 3$) are the Pauli matrices. The $SU(2)_L \otimes U(1)_Y$ covariant derivative of $\Sigma$ is defined as

$$D_\mu \Sigma \equiv \partial_\mu \Sigma + ig \frac{\tau^a}{2} W^a_\mu \Sigma - ig' \Sigma \frac{\tau^3}{2} B_\mu. \quad (2)$$

The lowest-order terms in the derivative expansion of the effective Lagrangian are

$$\mathcal{L}^{(2)} = \frac{v^2}{4} \text{Tr} \left[ (D_\mu \Sigma)^\dagger (D^\mu \Sigma) \right] + \beta_1 g'^2 \frac{v^2}{4} \left( \text{Tr} [TV_{\mu}] \right)^2. \quad (3)$$

where we have introduced the auxiliary quantities $T \equiv \Sigma \tau^3 \Sigma^\dagger$ and $V_\mu \equiv (D_\mu \Sigma) \Sigma^\dagger$ which are $SU(2)_L$–covariant and $U(1)_Y$–invariant. Notice that $T$ is not invariant under $SU(2)_C$ custodial due to the presence of $\tau^3$.

The first term in Eq. (3) is responsible for giving mass to the $W^\pm$ and $Z$ gauge bosons for $v = (\sqrt{2} G_F)^{-1}$. The second term violates the custodial $SU(2)_C$ symmetry and contributes to $\Delta \rho$ at tree level, being strongly constrained by the low–energy data. This term can be understood as the low-energy remnant of a high–energy custodial symmetry breaking physics, which has been integrated out above a certain scale $\Lambda$. Moreover, at the one–loop level, this term is also required in order to cancel the divergences in $\Delta \rho$, arising from
diagrams containing a hypercharge boson in the loop. This subtraction renders $\Delta \rho$ finite, although dependent on the renormalization scale [20].

At the next order in the derivative expansion, $D = 4$, several operators can be written down [20]. We shall restrict ourselves to those containing genuine quartic vector-boson interactions, which are

$$L_4^{(4)} = \alpha_4 [\text{Tr} (V_\mu V_\nu)]^2,$$

$$L_5^{(4)} = \alpha_5 [\text{Tr} (V_\mu V^\mu)]^2,$$

$$L_6^{(4)} = \alpha_6 \text{Tr} (V_\mu V_\nu) \text{Tr} (TV^\mu) \text{Tr} (TV^\nu),$$

$$L_7^{(4)} = \alpha_7 \text{Tr} (V_\mu V^\mu) [\text{Tr} (TV^\nu)]^2,$$

$$L_{10}^{(4)} = \frac{1}{2} \alpha_{10} [\text{Tr} (TV_\mu) \text{Tr} (TV_\nu)]^2.$$ (8)

In an arbitrary gauge, these Lagrangian densities lead to quartic vertices involving gauge bosons and/or Goldstone bosons. In the unitary gauge, these effective operators give rise to anomalous $ZZZZ$ (all operators), $W^+W^-ZZ$ (all operators except $L_4^{(4)}$), and $W^+W^-W^+W^-$ ($L_4^{(4)}$ and $L_5^{(4)}$) interactions. Moreover, the interaction Lagrangians $L_6^{(4)}$, $L_7^{(4)}$, and $L_{10}^{(4)}$ violate the $SU(2)_C$ custodial symmetry due to the presence of $T$ in their definitions. Notice that quartic couplings involving photons remain untouched by the genuinely quartic anomalous interactions at the order $D = 4$. The Feynman rules for the quartic couplings generated by these operators can be found in the last article of Ref. [20].

In chiral perturbation theory, the $p^4$ contribution to the processes $e^+e^- \rightarrow W^+W^-Z$ and $ZZZ$ arises from the tree level insertion of $p^4$ operators, as well as from one-loop corrections due to the $p^2$ interactions, which renormalize the $p^4$ operators [20]. However, the loop corrections to the scattering amplitudes are negligible in comparison to the $p^4$ contributions for the range of values of the couplings and center–of–mass energies considered in this paper. Therefore, numerically, our analysis is consistent even though we neglected the loop corrections and kept only the tree–level $p^4$ contributions.
III. LIMITS ON QUARTIC COUPLINGS

In order to study the quartic couplings of vector bosons we analyzed the processes

\[ e^+e^- \rightarrow W^+W^-Z , \]
\[ e^+e^- \rightarrow ZZZ , \]

which may receive contributions from anomalous \( WWZZ \) and \( ZZZZ \) interactions. We included in our calculations all SM and anomalous contributions that lead to these final states. Therefore, we consistently considered the effect of all interferences between the anomalous and SM amplitudes. The scattering amplitudes were generated using Madgraph [23] in the framework of Helas [24], with the anomalous couplings arising from the Lagrangians (4-8) being implemented as Fortran routines. Moreover, we include in our calculation the \( W \)’s and \( Z \)’s decays taking into account the gauge boson widths, spin structures, and correlations of the scattering amplitude.

We required the visible final state fermions to be in the rapidity region \(|\eta| < 3\) and separated by \( \Delta R = \sqrt{\Delta \eta^2 + \Delta \phi^2} > 0.7\). Furthermore, we also folded in the experimental resolution factors associated to the electromagnetic and hadronic calorimetry:

\[ \delta E_E \bigg|_{em} = \frac{0.12}{\sqrt{E}} \oplus 0.01 \text{ electromagnetic ,} \]
\[ \delta E_E \bigg|_{had} = \frac{0.25}{\sqrt{E}} \oplus 0.02 \text{ hadronic .} \]

The momentum carried out by neutrinos was obtained using energy–momentum conservation after smearing the momenta of final state quarks and charged leptons. As illustration, we show in Fig. 1 the effect of the smearing on the reconstructed difermion invariant masses for \( W \)’s and \( Z \)’s.

Difermion final states \( (jj, \ell^+\ell^-, \ell^+\nu, \text{ and } \nu\nu; \text{ with } \ell = e, \mu) \) were identified as being a \( W \) or a \( Z \) provided their invariant masses after the smearing were in the range \([23]\), respectively,

\[ \left[ 0.85M_W , \frac{1}{2}(M_W + M_Z) \right] , \left[ \frac{1}{2}(M_W + M_Z) , 1.15M_Z \right] . \]
In what follows we present our results for two different center–of–mass energies, 500 GeV and 1 TeV, assuming an integrated luminosity of 100 fb$^{-1}$ for both energies. We also study the impact of a 80% left–handed and 80% right–handed polarized electron beam while keeping the positron beam unpolarized.

\textbf{A. }$e^+e^- \to W^+W^-Z$

We identified $W^+W^-Z$ events through the topologies $6j$, $4j+2\ell$, $4j+2\nu$, and $4j+\ell+\nu$, requiring two difermion systems with invariant masses compatible with the $W$ mass — see Eq. (13) — and one difermion system with an invariant mass consistent with it being a $Z$. We show in Table I the fraction of $WWZ$ events that are reconstructed as $WWZ$ and $ZZZ$ for center–of–mass energies of 0.5 and 1 TeV. It is interesting to notice that the reconstruction probabilities are basically independent of the $e^-$ polarization. Furthermore, the fraction of $WWZ$ events reconstructed as $ZZZ$ generates a background for the study of anomalous couplings in $ZZZ$ production.

The cross section for $W^+W^-Z$ ($ZZZ$) is a quadratic function of the anomalous couplings $\alpha_i$, \textit{i.e.}

$$\sigma_{\text{tot}} = \sigma_{\text{sm}} + \sum_i \alpha_i \sigma_{\text{int}}^{\alpha_i} + \sum_{ij} \alpha_i \alpha_j \sigma_{\text{ano}}^{\alpha_i\alpha_j},$$

(14)

where $\sigma_{\text{sm}}$ stands for the SM cross section and $\sigma_{\text{int}}^{\alpha_i}$ ($\sigma_{\text{ano}}^{\alpha_i\alpha_j}$) is the interference (pure anomalous) contribution. In Table II, we present our results after the cuts on $\eta$ and $\Delta R$, but before $W$ and $Z$ identification. Therefore, these results should be multiplied by the efficiencies given in Table II. Notice that there are only two independent Lorentz invariant structures for the $WWZZ$ vertices at $p^4$ order, which implies that the couplings $\alpha_5$ and $\alpha_7$ ($\alpha_4$ and $\alpha_6$) give rise to identical contributions to $\sigma_{\text{int}}^{\alpha_i}$ and $\sigma_{\text{ano}}^{\alpha_i\alpha_j}$ in $W^+W^-Z$ production. From this table we can witness that the SM contributions are a slowly varying functions of the center–of–mass energy, while the anomalous contributions grow rapidly, as one could naively expect. Moreover, the SM background can be efficiently reduced using right–handed polarized beams.
electrons as this polarization eliminates almost completely the contribution where the $W^−$ couples directly to the $e^−$ fermion line.

In order to quantify the effect of the new couplings, we defined the statistical significance $S$ of the anomalous signal

$$S = \frac{\sigma_{\text{tot}} - \sigma_{\text{sm}}}{\sqrt{\sigma_{\text{sm}}}} \sqrt{L\epsilon},$$

which can be easily evaluated using the parametrization (14) with the coefficients given in Table II. In the above expression $\epsilon$ stands for the detection efficiencies presented in Table I.

Table III contains the values of the quartic anomalous couplings that lead to an increase in the total number of events smaller than 3$\sigma$, assuming an integrated luminosity of 100 fb$^{−1}$ and that only one anomalous coupling is non-vanishing. These limits were obtained combining events reconstructed as $W^+W^−Z$ from all the topologies. It is interesting to notice that having right–handed polarized electrons improves the bounds in 20–30% with respected to the results for unpolarized beams, while the use of left–handed electrons weakens the limits. This result is in agreement with Ref. [12]. Moreover, the bounds improve as the center–of–mass energy increases since the anomalous contributions grow with energy.

In general, more than one anomalous coupling might be non-vanishing. In this case the correlation among the anomalous couplings can be easily taken into account using the full expression of Eq. (14) and Table II.

In order to discriminate between the different couplings we studied the kinematical distributions of the final gauge bosons. Figure 2 displays the $W^+W^−$ invariant mass spectrum and the $p_T$ distribution of the $Z$ in the $W^+W^−Z$ production with unpolarized beams at $\sqrt{s} = 500$ GeV. We plotted in this figure the standard model prediction (dotted line) as well as the predictions for $\alpha_4 = 0.61$ (dashed line) and $\alpha_5 = 0.38$ (solid line), which are the values that lead to a 3$\sigma$ signal in the total number of events for unpolarized beams. As we can see, the $W^+W^−$ invariant mass distribution for $\alpha_4$ presents a larger contribution at small values of the $W^+W^−$ invariant mass, while $\alpha_5$ gives rise to more events with larger invariant masses. In principle we can use this distribution not only to distinguish the anomalous couplings,
but also to increase the sensitivity to the signal. However, this can only be accomplished with higher integrated luminosity. On the other hand the $p_T$ distributions of the $Z$ are very similar in the SM and in presence of the anomalous couplings, being the only difference the larger number of events in the latter case.

**B. $e^+e^- \to ZZZ$**

The production cross section for $ZZZ$ final states is much smaller than the one for $W^+W^-Z$, and consequently just a few fermionic topologies can be used to identify these events. We considered only the final states $6j$, $4j+2\ell$, and $4j+2\nu$, and required three difermion systems with invariant masses compatible with the $Z$ one according to the prescription given in Eq. (13). We present in Table [IV] the efficiency for the reconstruction of the $ZZZ$ final state for the above topologies and center–of–mass energies of 0.5 and 1.0 TeV. Analogously to the $W^+W^-Z$ case, these efficiencies are independent of the polarization of the $e^-$. Table [V] contains the values of $\sigma_{sm}$, $\sigma_{\alpha_i}^{\alpha_i}$, and $\sigma_{\alpha_i\alpha_j}^{\alpha_i\alpha_j}$ for $ZZZ$ production, taking into account the $\eta$ and $\Delta R$ cuts, but not the reconstruction efficiencies. At $p^4$ order in chiral perturbation theory, all the anomalous interactions are proportional to each other since there is only one possible Lorentz structure for the vertex which is multiplied by $\alpha_4+\alpha_5+2(\alpha_6+\alpha_7+\alpha_{10})$. Therefore, we only present the results for $\alpha_4$, being straightforward the generalization to the other cases. From this table we can see that most of the reconstructed $ZZZ$ events will be observed in the $6j$ and $4j+2\nu$ topologies. Furthermore, the largest anomalous contribution comes from $\sigma_{\alpha_i\alpha_j}^{\alpha_i\alpha_j}$, being the interference with the SM of the same order of the SM contribution but with the opposite sign. Analogously, in $W^+W^-Z$ production, the anomalous contributions grow substantially with the increase of the center–of–mass energy, while the SM cross section decreases slightly.

We present in Table [VI] the $3\sigma$ allowed for genuinely quartic couplings that can be obtained from the non observation of deviations from the SM in $ZZZ$ production. In our
analysis of the $6j$ topology, we included as background the $6j$ events coming from $W^+W^-Z$ that are identified as $ZZZ$. Despite the reduced number of events in the $ZZZ$ channel, the bounds on the quartic couplings are at least a factor of 2 better than the ones drawn from the $W^+W^-Z$ channel due to the smaller size of the background. On the other hand, contrary to the $W^+W^-Z$ channel beam polarization does not lead to a substantial improvement on the attainable limits.

We display in Fig. 3 the pseudorapidity and transverse momentum distribution of the $Z$’s in unpolarized $ZZZ$ production at $\sqrt{s} = 500$ GeV. As we can see, the anomalous quartic interactions leads to more centrally produced $Z$’s (smaller $|\eta_z|$) which have a slightly harder $p_T$ spectrum. However, the number of reconstructed events is not large enough to allow the use of cuts to enhance the anomalous contributions.

IV. DISCUSSION AND CONCLUSIONS

$W^+W^-Z$ and $ZZZ$ are the best channels for direct study of quartic gauge–boson couplings in $e^+e^-$ colliders with center–of–mass energies up to 1 TeV. At higher energies the most $e^+e^- \rightarrow W^+W^-f\bar{f}$ becomes important process [14]. We showed in this work that the NLC will be able to uncover the existence of anomalous quartic couplings of the order $\mathcal{O}(10^{-1})$ for center–of–mass energies up to 1 TeV and an integrated luminosity of 100 fb$^{-1}$; see Tables III and VI. Despite these limits being weaker or of the order of the present indirect bounds [7], the above processes will provide a direct test of the quartic interactions among the electroweak gauge bosons. We have also shown that the use of a right–handed polarized electron beam leads to better limits on the anomalous interactions from the $W^+W^-Z$ production due to the substantial reduction of the SM backgrounds.

It is also important to devise a strategy to disentangle the anomalous couplings in case a departure from the SM prediction is observed. In $W^+W^-Z$ production, the analysis of the $W^+W^-$ invariant mass distribution, see Fig. 2, can be used to distinguish between the two possible structures for the $WWZZ$ vertex, one associated to $\alpha_{4,6}$ and the other
related to $\alpha_{5,7}$. However, we are still left some two possibilities in both cases. At this point it is important to use the information from the $ZZZ$ reaction, because the $SU(2)_C$ violating interactions leads to a much larger excess of events for the same value of the anomalous coupling, due to the coupling structure $\alpha_4 + \alpha_5 + 2(\alpha_6 + \alpha_7 + \alpha_{10})$. Therefore, the combination of the $W^+W^-$ distribution and the total number of events in both reactions are a powerful tool to separate the effects of the different anomalous couplings provided there is enough statistics. Moreover, the comparison between the $W^+W^-Z$ event rates for different polarizations can also be used to further distinguish between the couplings $\alpha_{4,6}$ and $\alpha_{5,7}$, since the latter are less sensitive to the electron polarization. Finally the anomalous coupling $\alpha_{10}$ has the distinguished characteristics of modifying only the $ZZZ$ production.

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TABLES

| Topology | $F_{WWZ\%}$ | $F_{ZZZ\%}$ |
|----------|--------------|--------------|
| $6j$     | 64./52.      | 0.8/1.2      |
| $4j + 2\ell$ | 66./55.   | 0.8/1.2      |
| $4j + 2\nu$ | 28./8.     | 0.2/0.2      |
| $4j + \ell \nu_\ell$ | 20./5.      | 0./0.        |

TABLE I. Fraction of $WWZ$ events that are reconstructed as $WWZ$ and $ZZZ$ for several topologies and center–of–mass energies of 0.5/1 TeV.
| Topology | $\sigma_{\text{sm}}$ (fb) | $-\sigma_{\text{int}}^{\alpha_4}$ (fb) | $\sigma_{\text{ano}}^{\alpha_4\alpha_4}$ (fb) | $\sigma_{\text{int}}^{\alpha_4}$ (fb) | $\sigma_{\text{ano}}^{\alpha_5}$ (fb) | $\sigma_{\text{ano}}^{\alpha_5\alpha_5}$ (fb) |
|----------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| $\sqrt{s}$/TeV | 0.5 (1) | 0.5 (1) | 0.5 (1) | 0.5 (1) | 0.5 (1) | 0.5 (1) |
| 6j | unpol | 7.41 (8.09) | 0.12 (0.12) | 2.42 (6.30) | 0.47 (0.21) | 4.53 (14.89) | 1.18 (-4.71) |
| | pol − | 13.4 (14.59) | 0.0 (0.0) | 2.81 (7.38) | 0.57 (0.) | 5.26 (17.36) | 1.37 (-5.52) |
| | pol + | 1.61 (1.60) | 0.22 (0.24) | 2.0 (5.2) | 0.34 (0.38) | 3.77 (12.4) | 0.99 (-3.93) |
| 4j + 2ℓ | unpol | 0.74 (0.80) | 0.01 (0.01) | 0.24 (0.63) | 0.01 (0.02) | 0.45 (1.48) | 0.12 (-0.47) |
| | pol − | 1.32 (1.43) | 0.0 (0.0) | 0.28 (0.72) | 0.05 (0.0) | 0.52 (1.73) | 0.14 (-0.53) |
| | pol + | 0.159 (0.163) | 0.023 (0.023) | 0.20 (0.52) | 0.03 (0.04) | 0.375 (1.23) | 0.01 (-0.39) |
| 4j + 2ν | unpol | 3.06 (2.88) | 0.01 (0.09) | 0.94 (3.68) | 0.22 (0.16) | 2.09 (6.87) | 0.66 (-2.73) |
| | pol − | 5.44 (5.14) | 0.0 (0.05) | 1.09 (4.30) | 0.29 (0.12) | 2.44 (7.95) | 0.76 (-3.21) |
| | pol + | 0.65 (0.58) | 0.06 (0.15) | 0.78 (3.08) | 0.15 (0.18) | 1.74 (5.7) | 0.55 (-2.27) |
| 4j + ℓ νℓ | unpol | 5.95 (6.38) | 0.12 (0.11) | 1.97 (7.31) | 0.39 (0.48) | 3.45 (19.74) | 0.98 (-3.63) |
| | pol − | 10.63 (11.6) | 0.0 (0.0) | 2.29 (8.60) | 0.49 (0.43) | 4.03 (23.12) | 1.14 (-4.27) |
| | pol + | 1.28 (1.28) | 0.2 (0.25) | 1.64 (6.14) | 0.26 (0.56) | 2.89 (16.4) | 0.81 (-3.04) |

**TABLE II.** Values for the standard model, pure anomalous and interference cross sections (see Eq. [14]) for the $W^+W^-Z$ production and several center–of–mass energies and $e^-$ polarizations. The pol − (pol +) lines correspond to 80% left–handed (right–handed) electron beam polarization.
| $\sqrt{s}$ GeV | $e^-$ polarization (%) | $\alpha_{4,6}$   | $\alpha_{5,7}$   |
|--------------|------------------------|------------------|------------------|
| 500          | 0                      | $(-0.56, 0.61)$  | $(-0.48, 0.38)$  |
| 500          | $-80$                  | $(-0.63, 0.63)$  | $(-0.51, 0.40)$  |
| 500          | 80                     | $(-0.39, 0.49)$  | $(-0.37, 0.28)$  |
| 1000         | 0                      | $(-0.38, 0.40)$  | $(-0.26, 0.25)$  |
| 1000         | $-80$                  | $(-0.43, 0.43)$  | $(-0.28, 0.28)$  |
| 1000         | 80                     | $(-0.27, 0.32)$  | $(-0.20, 0.17)$  |

TABLE III. 3\(\sigma\) allowed values of the quartic anomalous couplings obtained from the reaction 
\(e^+e^- \rightarrow W^+W^-Z\).

| Topology | $F_{ZZZ}\%$ |
|----------|-------------|
| 6\(j\)  | 59./54.    |
| 4\(j\) + 2\(\ell\) | 62./58. |
| 4\(j\) + 2\(\nu\) | 25./7. |

TABLE IV. Fraction of $ZZZ$ events that are reconstructed as $ZZZ$ for several topologies and 
center–of–mass energies of 0.5/1 TeV.
| Topology | $\sqrt{s}$/TeV | $\sigma_{\text{sm}}$ (fb) | $-\sigma_{\text{int}}^{\alpha_4}$ (fb) | $\sigma_{\text{ano}}^{\alpha_4}$ (fb) |
|----------|----------------|----------------------|---------------------------------|----------------------|
|          |                | 0.5 (1)              | 0.5 (1)                          | 0.5 (1)              |
| 6j       | unpol          | 0.163 (0.145)        | 0.169 (0.119)                    | 2.83 (8.35)          |
|          | pol $-$        | 0.236 (0.21)         | 0.224 (0.158)                    | 3.32 (9.77)          |
|          | pol $+$        | 0.090 (0.081)        | 0.11 (0.082)                     | 2.38 (6.98)          |
| 4j $+$ 2\ell | unpol        | 0.049 (0.043)        | 0.044 (0.034)                    | 0.846 (2.49)         |
|          | pol $-$        | 0.07 (0.063)         | 0.057 (0.047)                    | 0.985 (2.90)         |
|          | pol $+$        | 0.027 (0.024)        | 0.03 (0.023)                     | 0.71 (2.08)          |
| 4j $+$ 2\nu | unpol        | 0.204 (0.144)        | 0.195 (0.118)                    | 3.67 (14.2)          |
|          | pol $-$        | 0.294 (0.207)        | 0.27 (0.156)                     | 4.29 (16.5)          |
|          | pol $+$        | 0.113 (0.080)        | 0.126 (0.077)                    | 3.08 (11.8)          |

TABLE V. Values for the standard model, pure anomalous and interference cross sections (see Eq. [4]) for the $ZZZ$ production and several center–of–mass energies and $e^-$ polarizations.

| $\sqrt{s}$/GeV | $e^-$ polarization (%) | $\alpha_{4,5}$     | $\alpha_{6,7,10}$ |
|----------------|------------------------|--------------------|--------------------|
| 500            | 0                      | $(-0.19, 0.25)$    | $(-0.096, 0.12)$   |
| 500            | 80                     | $(-0.20, 0.26)$    | $(-0.098, 0.13)$   |
| 500            | $-80$                  | $(-0.18, 0.22)$    | $(-0.088, 0.11)$   |
| 1000           | 0                      | $(-0.14, 0.15)$    | $(-0.068, 0.075)$  |
| 1000           | 80                     | $(-0.14, 0.16)$    | $(-0.071, 0.079)$  |
| 1000           | $-80$                  | $(-0.12, 0.13)$    | $(-0.058, 0.063)$  |

TABLE VI. $3\sigma$ allowed values of the quartic anomalous couplings obtained from the reaction $e^+e^- \to ZZZ$. 
FIG. 1. Reconstructed invariant mass distribution for a jet pair from $W$ and $Z$ decays. The full line only includes the effect of the finite width while the dashed line contains also the effect of the smearing due to the experimental resolution.
FIG. 2. $W^+W^-$ invariant mass and $p_{TZ}$ distributions for unpolarized $W^+W^-Z$ production at $\sqrt{s} = 500$ GeV. The dotted line stands for the SM result, while the solid (dashed) line represent the case $\alpha_5 = 0.38$ ($\alpha_4 = 0.61$).
FIG. 3. Pseudorapidity and transverse momentum Z distributions in unpolarized $ZZZ$ production at $\sqrt{s} = 500$ GeV. The dotted line stands for the SM result, while the solid line represents the case $\alpha_4 = 0.24$. 