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Quartic Gauge Interactions at NLC *

— Interplay of WWZ/ZZZ-Production and WW-Fusion —

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Abstract. The new physics effects of the strong electroweak symmetry breaking (EWSB) can be parametrized by a complete set of the chiral Lagrangian parameters. Five of them characterize the anomalous quartic gauge interactions which involve pure Goldstone dynamics for the EWSB but are *least* constrained from the current data. After analyzing the different patterns of these parameters in connection with typical underlying resonance/non-resonance models, we perform model-independent systematic study on bounding them via WWZ/ZZZ-production and WW-fusion at the next generation high energy $e^{\pm}e^{-}$ linear colliders (NLC). The main focus is put onto the interplay of these two production mechanisms for achieving a complete probe of the quartic gauge interactions in a multi-parameter analysis. The important roles of both polarized e^{-} and e^{+} beams are revealed.

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1. Introduction

The standard model (SM) has so far been confirmed with great precision at the scale of $O(100)$ GeV [1], although its electroweak symmetry breaking (EWSB) mechanism remains undetermined [2]. The low energy effects of the new physics beyond the SM therefore have to be very small. This basically suggests two main possibilities for the EWSB sector: weakly or strongly interacting. The weakly coupled case ensures the new physics having *decoupling* property [3] at low scales, while the strongly interacting scenario (to be investigated below) parametrizes the new physics as the next-to-leading order (NLO) effects at the TeV scale which is also *generically small* and typically of $O\left(\frac{v^2}{\Lambda^2}\right)\sim O(6\times10^{-3})$ [1]. This situation raises a great challenge to the future high energy colliders for decisively probing the EWSB mechanism in *both scenarios*[2]. The present study is devoted to make the precision test on the strong EWSB scenario at the next generation high energy $e^{\pm}e^{-}$ linear colliders (NLC) which are under current intensive experimental and theoretical investigations [6].

Below the new physics scale Λ, all the new physics effects in the EWSB sector can be parametrized by a complete set of the NLO effective operators of the electroweak chiral Lagrangian (EWCL) [7], in which the $SU(2)_L\otimes U(1)_Y$ gauge symmetry is nonlinearly realized[3]. Without experimental observation on any new light resonance [1], this effective field theory approach [8,4] provides the most economic description of the possible new physics effects. There are a total fifteen NLO operators in the EWCL and their contributions to the various high energy processes at the LHC and the NLC have been systematically classified by means of the global power counting analysis [9]. Among the complete set of the NLO operators, five of them characterize only the quartic gauge interactions, which involve the pure Goldstone boson [10] dynamics (according to the equivalence theorem [11,12]) and are most important for the gauge boson fusion and triple gauge boson production at the TeV scale. They are defined as follows [7]:

$$\left\{\begin{array}{lcl}\mathcal{L}_4 &=& \ell_4\left(\frac{v}{\Lambda}\right)^2[\mathrm{Tr}(\mathcal{V}_\mu\mathcal{V}_\nu)]^2\,,\\ \mathcal{L}_5 &=& \ell_5\left(\frac{v}{\Lambda}\right)^2[\mathrm{Tr}(\mathcal{V}_\mu\mathcal{V}^\mu)]^2\,;\end{array}\right\}\qquad (\ SU(2)_c:\ \surd\ )$$

1) Here $v(\simeq 246$ GeV) is the Fermi scale and $\Lambda(\lesssim 4\pi v\simeq 3.1$ TeV) denotes the new physics scale characterizing the EWSB [4].

2) A possible detection of a light scalar is insufficient to verify the SM Higgs mechanism, and to fully disentangle the imagined complex supersymmetry particle spectrum is not going to be achieved at the LHC alone [5].

3) It is advised that whenever the decoupling theorem [3] becomes ineffective, the nonlinear realization should apply.

$$\left\{\begin{array}{lcl} \mathcal{L}_6 &=& \ell_6\left(\frac{v}{\Lambda}\right)^2 [\mathrm{Tr}(\mathcal{V}_\mu\mathcal{V}_\nu)]\mathrm{Tr}(\mathcal{T}\mathcal{V}^\mu)\mathrm{Tr}(\mathcal{T}\mathcal{V}^\nu)\ , \\ \mathcal{L}_7 &=& \ell_7\left(\frac{v}{\Lambda}\right)^2 [\mathrm{Tr}(\mathcal{V}_\mu\mathcal{V}^\mu)]\mathrm{Tr}(\mathcal{T}\mathcal{V}_\nu)\mathrm{Tr}(\mathcal{T}\mathcal{V}^\nu)\ , \\ \mathcal{L}_{10} &=& \ell_{10}\left(\frac{v}{\Lambda}\right)^2 \frac{1}{2}[\mathrm{Tr}(\mathcal{T}\mathcal{V}^\mu)\mathrm{Tr}(\mathcal{T}\mathcal{V}^\nu)]^2\ ; \end{array}\right\} \quad (\ SU(2)_c:\ \times\ ) \quad (1)$$

where $\mathcal{V}_\mu \equiv (D_\mu U)U^\dagger$, $D_\mu U = \partial_\mu U + \mathbf{W}_\mu U - U\mathbf{B}_\mu$, $\mathbf{W}_\mu \equiv igW_\mu^a\tau^a/2$, $\mathbf{B}_\mu \equiv ig'B_\mu\tau^3/2$, $U = \exp[i\tau^a\pi^a/v]$ (with π^a the would-be Goldstone boson field), and $\mathcal{T} \equiv U\tau_3U^\dagger$ is the custodial $SU(2)_c$-violation operator. In (1), the operators $\mathcal{L}_{4,5}$ conserve $SU(2)_c$ while $\mathcal{L}_{6,7,10}$ violate $SU(2)_c$. Here, the dependence on v and Λ is factorized out so that the dimensionless coefficient ℓ_n of the operator $\mathcal{L}_n$ is naturally of $O(1)$ [4]. Because they contain *only* quartic gauge couplings (QGCs), these five operators cannot be directly tested via their tree-level contributions at low energies and are therefore least constrained from the current data. So far, only some rough estimates have been made by inserting them into the one-loop corrections and keeping the log-terms only. Here is an updated estimate at 90% C.L. by choosing $\Lambda = 2$ TeV and setting only one parameter nonzero at a time [9]:

$$\begin{array}{c} -4 \le \ell_4 \le 20\ , \quad -10 \le \ell_5 \le 50\ ; \\ -0.7 \le \ell_6 \le 4\ , \quad -5 \le \ell_7 \le 26\ , \quad -0.7 \le \ell_{10} \le 3\ ; \end{array} \quad (2)$$

From (2), we see that the bounds on the $SU(2)_c$ symmetric parameters $\ell_{4,5}$ are about an order of magnitude above their natural size; while the allowed range for the $SU(2)_c$-breaking parameters ℓ_{6-10} is about a factor of $O(10-100)$ larger than that for $\ell_0 = \frac{\Lambda^2}{2v^2}\Delta\rho \left(= \frac{\Lambda^2}{2v^2}\alpha T\right)$ derived from the ρ (or T) parameter: $0.052 \le \ell_0 \le 0.12$ [9], for the same Λ and confidence level. To directly test the EWSB dynamics, it is therefore important to probe these QGCs at future high energy scattering processes where their contributions can be greatly enhanced due to the sensitive power-dependence on the scattering energy [9].

2. Quartic Gauge Interactions and Underlying Models

Though the true fundamental theory behind this effective EWCL is yet unknown, it is important to examine how the typical underlying resonance/non-resonance models contribute to these EWSB parameters. Knowing the theoretical sizes and patterns of these parameters tells how to use the phenomenological bounds (cf. sections 3 and 4 below) for discriminating different new physics models. We shall concentrate on the quartic gauge interactions (1) and consider typical models such as a heavy scalar (S), a vector (V_μ^a) and an axial vector (A_μ^a) for the resonance scenario, and the one flavor new heavy doublet fermions for the non-resonance scenario. The effective Lagrangian for these new physics models can be formulated as follows:

$$\mathcal{L}_{\rm eff}^{\rm new} = \mathcal{L}_{\rm eff}^{S} + \mathcal{L}_{\rm eff}^{VA} + \mathcal{L}_{\rm eff}^{F} + \cdots \quad (3)$$

• **A Non-SM Singlet Scalar** Up to dimension-4 and including both $SU(2)_c$ conserving and breaking effects, we can write down the most general Lagrangian for a singlet scalar which is invariant under the SM gauge group $\mathcal{G}^{\rm SM} = SU(2)_L \otimes U(1)_Y$:

$$\begin{array}{rcl} \mathcal{L}_{\rm eff}^{S} &=& \frac{1}{2}\left[\partial^\mu S\partial_\mu S - M_S^2 S^2\right] - V(S) \\ && -\left[\frac{\kappa_s}{2}vS + \frac{\kappa_s'}{4}S^2\right]\mathrm{Tr}\left[\mathcal{V}_\mu\mathcal{V}^\mu\right] - \left[\frac{\tilde\kappa_s}{2}vS + \frac{\tilde\kappa_s'}{4}S^2\right]\left[\mathrm{Tr}\mathcal{T}\mathcal{V}_\mu\right]^2 \end{array} \quad (4)$$

where $V(S)$ contains only Higgs self-interactions. The SM Higgs scalar only corresponds to a special parameter choice: $\kappa_s = \kappa_s' = 1$, $\tilde\kappa_s = \tilde\kappa_s' = 0$ and $V(S) = V(S)_{\rm SM}$. A heavy scalar can be integrated out from low energy spectrum with its effects formulated in the heavy mass expansion:

$$\widehat{\mathcal{L}}_{\rm eff}^{S} = \frac{v^2}{8M_S^2}\left[\kappa_s\mathrm{Tr}\left(\mathcal{V}_\mu\mathcal{V}^\mu\right) + \tilde\kappa_s\mathrm{Tr}\left(\mathcal{T}\mathcal{V}_\mu\right)^2\right]^2 + O\left(\frac{1}{M_S^4}\right) \quad (5)$$

With the identification $\Lambda = M_S$, the corresponding contributions to (1) are derived as follows

$$\ell_4^s = 0\ , \quad \ell_5^s = \frac{\kappa_s^2}{8} \ge 0\ ; \qquad \ell_6^s = 0\ , \quad \ell_7^s = \frac{\kappa_s\tilde\kappa_s}{4}\ , \quad \ell_{10}^s = \frac{\tilde\kappa_s^2}{8} \ge 0\ . \quad (6)$$

Besides contributing to the quartic gauge interactions, the parameters $\{\kappa_s, \tilde\kappa_s\}$ are also related to the physical partial width of S decaying into two longitudinal gauge bosons ($\Gamma_S[S \to W_L W_L, Z_L Z_L]$) for a given mass M_S:

$$2\kappa_s^2 + (\kappa_s - 2\tilde\kappa_s)^2 = \frac{32\pi v^2\Gamma_S}{M_S^3}\ . \quad (7)$$

In (6)-(7), the deviation from $\kappa_s = 1$ and $\tilde\kappa_s = 0$ signals a *non-SM Higgs boson*.

• **Vector and Axial-Vector Bosons** The LEP measurement on the S-parameter disfavors the naive QCD-like dynamics for the EWSB [13], where the vector $\rho_{\rm TC}$ is the lowest new resonance in the TeV regime. This suggests the necessity of including the axial-vector boson in a general formalism for modeling the non-QCD-like dynamics. We consider the vector V_μ^a and axial-vector A_μ^a fields as the weak isospin triplets of (custodial) $SU(2)_c$. $\{V, A\}$ transform under the SM global $SU(2)_c$ as

$$\widehat{V}_\mu \Rightarrow \widehat{V}_\mu' = \Sigma_v\widehat{V}_\mu\Sigma_v^\dagger\ , \qquad \widehat{A}_\mu \Rightarrow \widehat{A}_\mu' = \Sigma_v\widehat{A}_\mu\Sigma_v^\dagger\ . \quad (8)$$

where $\widehat{V}_\mu \equiv V_\mu^a\tau^a/2$, $\widehat{A}_\mu \equiv A_\mu^a\tau^a/2$, and $\Sigma_v \in SU(2)_c$. If $\{V, A\}$ are further regarded as gauge fields of a new local hidden symmetry group $\mathcal{H} = SU(2)_L' \otimes$

$SU(2)'_R$ (with a discrete left-right parity) [14], we can write down the following general Lagrangian (up to two derivatives), in the *unitary gauge* of the group $\mathcal{H}$ [4] and with both $SU(2)_c$-conserving and -breaking effects included,

$$\begin{aligned}\mathcal{L}_{\rm eff}^{VA} = \mathcal{L}_{\rm kinetic}^{VA} - v^2 \Big[& \kappa_0 {\rm Tr}\overline{\mathcal{V}}_\mu^2 + \kappa_1 {\rm Tr}\left(J_\mu^V - 2V_\mu\right)^2 + \kappa_2 {\rm Tr}\left(J_\mu^A + 2A_\mu\right)^2 + \kappa_3 {\rm Tr} A_\mu^2 \\ & + \tilde{\kappa}_0 \left[{\rm Tr}\widetilde{\mathcal{T}}\overline{\mathcal{V}}_\mu\right]^2 + \tilde{\kappa}_1 \left[{\rm Tr}\widetilde{\mathcal{T}}(J_\mu^V - 2V_\mu)\right]^2 + \tilde{\kappa}_2 \left[{\rm Tr}\widetilde{\mathcal{T}}(J_\mu^A + 2A_\mu)\right]^2 + \tilde{\kappa}_3 \left[{\rm Tr}\widetilde{\mathcal{T}} A\right]^2 \Big]\end{aligned} \tag{9}$$

with

$$\left\{\begin{array}{l} J_\mu^V = J_\mu^L + J_\mu^R \\ J_\mu^A = J_\mu^L - J_\mu^R \end{array}\right. \qquad \left\{\begin{array}{l} J_\mu^L = \xi^\dagger D_\mu^L \xi = \xi^\dagger \left(\partial_\mu \xi + W_\mu \xi\right) \\ J_\mu^R = \xi D_\mu^R \xi^\dagger = \xi\left(\partial_\mu \xi^\dagger + B_\mu \xi^\dagger\right) \end{array}\right. \tag{10}$$

and, by definition, $V_\mu \equiv i\tilde{g}\widehat{V}_\mu = i\tilde{g}V_\mu^a\tau^a/2$, $A_\mu \equiv i\tilde{g}\widehat{A}_\mu = i\tilde{g}A_\mu^a\tau^a/2$, $\overline{\mathcal{V}}_\mu \equiv U^\dagger D_\mu U = U^\dagger \mathcal{V}_\mu U$, $\widetilde{\mathcal{T}} = \tau^3 = U^\dagger \mathcal{T} U$,[5] and $U \equiv \xi^2$. ($\tilde{g}$ is the gauge coupling of the group $\mathcal{H}$.) Under the group $\mathcal{G}_{\rm global}^{\rm SM} \otimes \mathcal{H}_v$ (where $\mathcal{G}_{\rm global}^{\rm SM} = SU(2)_L \otimes SU(2)_R$ and $\mathcal{H}_v = SU(2)'_v \subset \mathcal{H}$), we have the following transformation laws:

$$\begin{aligned} & \xi' = \Sigma_L \xi h_v^\dagger = h_v \xi \Sigma_R^\dagger\,, \quad V'_\mu = h_v V_\mu h_v^\dagger + h_v \partial_\mu h_v^\dagger\,, \quad A'_\mu = h_v A_\mu h_v^\dagger\,; \\ & J_\mu^{V\prime} = h_v J_\mu^V h_v^\dagger + 2h_v\partial_\mu h_v^\dagger\,, \quad J_\mu^{A\prime} = h_v J_\mu^A h_v^\dagger\,; \end{aligned} \tag{11}$$

where $\Sigma_{L(R)} \in SU(2)_{L(R)}$ of $\mathcal{G}_{\rm global}^{\rm SM}$ and $h_v \in \mathcal{H}_v$. Under the choice of $\Sigma_L = \Sigma_R = \Sigma_v \in SU(2)_c$, ξ transforms linearly because of $\Sigma_v = h_v$. This plus eqs. (8) and (11) shows why the last four terms in (9) beak the SM custodial $SU(2)_c$ symmetry while the others conserve it. Among the above four new $SU(2)_c$-conserving parameters κ_n's, κ_0 is determined by normalizing the Goldstone kinematic term: $\kappa_0 = -4\kappa_2\kappa_3/(4\kappa_2 + \kappa_3)$.

For deriving ℓ_n's in (1) from (9), it is technically more convenient to use the matrix U instead of ξ as the variable. This can be done by going to another unitary gauge of $\mathcal{H}$ under a proper hidden local gauge transformation. Then, we derive

$$J_\mu^V = \overline{\mathcal{V}}_\mu + 2{\rm B}_\mu\,, \quad J_\mu^A = \overline{\mathcal{V}}_\mu\,. \tag{12}$$

---

4) By "unitary gauge" we mean a gauge containing no new Goldstone boson other than the three ones for generating the longitudinal components of the *known* W, Z. In fact, it is not essentially necessary to introduce such a new local symmetry $\mathcal{H}$ for $\{V, A\}$ [4] since $\mathcal{H}$ has to be broken anyway and $\{V, A\}$ can be traditionally treated as matter fields [15]. The hidden local symmetry formalism is more restrictive on the allowed free-parameters (κ_n's etc) due to the additional assumption about that new local group $\mathcal{H}$.

5) In the R_ξ-gauge of hidden group $\mathcal{H}$, $\widetilde{\mathcal{T}}$ can be generally defined as $\widetilde{\mathcal{T}} = \mathcal{M}^\dagger\tau^3\mathcal{M}$ with the unitary matrix $\mathcal{M}$ representing the additional new Goldstone bosons associating with the breaking of $\mathcal{H}$. Under $\mathcal{H}$, $\mathcal{M}$ transforms as $\mathcal{M} \Rightarrow \mathcal{M}' = h_R^\dagger \mathcal{M} h_L$ and thus $\widetilde{\mathcal{T}} \Rightarrow \widetilde{\mathcal{T}}' = h_L^\dagger \widetilde{\mathcal{T}} h_L$, with $h_{L(R)} \in SU(2)'_{L(R)}$. In the unitary gauge of $\mathcal{H}$, $\mathcal{M} = 1$, so that $\widetilde{\mathcal{T}} = \tau^3$.

After eliminating the V and A fields in the heavy mass expansion, we obtain the leading terms (of no explicit B_μ-dependence):

$$\begin{aligned}\widehat{\mathcal{L}}_{\rm eff}^{VA} \;=\; & \frac{(\eta^2-1)^2 + 16\tilde{\eta}^2}{8\tilde{g}^2}\left[({\rm Tr}\mathcal{V}_\mu\mathcal{V}_\nu)^2 - ({\rm Tr}\mathcal{V}_\mu\mathcal{V}^\mu)^2\right] + \\ & \frac{\tilde{\eta}\left(4(3-\eta^2)\tilde{\eta} + (1-\eta^2)\eta\right)}{8\tilde{g}^2}\left[{\rm Tr}\mathcal{V}_\mu^2\,({\rm Tr}\mathcal{T}\mathcal{V}_\nu)^2 - {\rm Tr}\,(\mathcal{V}_\mu\mathcal{V}_\nu)\,{\rm Tr}\,(\mathcal{T}\mathcal{V}^\mu)\,{\rm Tr}\,(\mathcal{T}\mathcal{V}^\nu)\right] + O\left(\frac{1}{M_{V,A}^4}\right)\end{aligned} \tag{13}$$

which contributes to ℓ_n as follows:

$$\left\{\begin{array}{l} \ell_4 = \ell_4^v + \ell_4^a \\ \ell_5 = \ell_5^v + \ell_5^a \\ \ell_6 = \ell_6^v + \ell_6^a \\ \ell_7 = \ell_7^v + \ell_7^a \\ \ell_{10} = \ell_{10}^v + \ell_{10}^a \end{array}\right. \qquad \left\{\begin{array}{l} \ell_4^v = -\ell_5^v = 1/[2\sqrt{2}\tilde{g}v\Lambda^{-1}]^2 > 0 \\ \ell_4^a = -\ell_5^a = \left[\eta^2(\eta^2-2) + 16\tilde{\eta}^2\right]/[2\sqrt{2}\tilde{g}v\Lambda^{-1}]^2 \\ \ell_6^v = \ell_7^v = 0 \\ \ell_6^a = -\ell_7^a = -\tilde{\eta}\left[4(3-\eta^2)\tilde{\eta} + (1-\eta^2)\eta\right]/[2\sqrt{2}\tilde{g}v\Lambda^{-1}]^2 \\ \ell_{10}^v = \ell_{10}^a = 0 \end{array}\right. \tag{14}$$

where

$$\eta = \frac{4\kappa_2}{4\kappa_2 + \kappa_3}\,, \quad \tilde{\eta} = \frac{2\kappa_2 + 4\tilde{\kappa}_2}{(4\kappa_2 + \kappa_3) + 2(4\tilde{\kappa}_2 + \tilde{\kappa}_3)} - \frac{2\kappa_2}{4\kappa_2 + \kappa_3}\,, \tag{15}$$

and $\Lambda = \min\{M_V, M_A\}$. At the leading order, $\{M_V, M_A\} \simeq \left\{\tilde{g}v\sqrt{\kappa_1},\ \tilde{g}v\sqrt{\kappa_2 + \kappa_3/4}\right\}$, after ignoring the SM gauge couplings g and g'. In (14), the factor $1/[\tilde{g}v\Lambda^{-1}]^2 \simeq \kappa_1(\Lambda/M_V)^2 = O(\kappa_1)$ and all $SU(2)_c$-breaking terms depend on $\tilde{\eta}$. We see that the $SU(2)_c$-symmetric contribution from the axial-vector boson interactions to $\ell_4^a = -\ell_5^a$ becomes negative for $|\eta| < \sqrt{2}$, while the summed contribution $\ell_4 = -\ell_5 = [(\eta^2-1)^2 + 16\tilde{\eta}^2]/[2\sqrt{2}\tilde{g}v\Lambda^{-1}]^2 \geq 0$. The deviation of η and/or $\tilde{\eta}$ from $\eta(\tilde{\eta}) = 0$ represents the *non-QCD-like* EWSB dynamics.

• **Heavy Doublet Fermions** Consider a simple model for one flavor heavy chiral fermions which form a left-handed weak doublet $(U_L, D_L)^T$ and right-handed singlets $\{U_R, D_R\}$, and joins a new strong $SU(N)$ gauge group in its fundamental representation. Their small mass-splitting breaks the $SU(2)_c$ and is characterized by the parameter $\omega = 1 - (M_U/M_D)^2$. The anomaly-cancellation is ensured by assigning the $\{U, D\}$ electric charges as $\{+\frac{1}{2}, -\frac{1}{2}\}$. By taking $\{U, D\}$ as the source of the EWSB, the W, Z masses can be generated by heavy fermion loops. The new contributions to the quartic gauge couplings of W/Z come from the *non-resonant* $\{U, D\}$ box-diagrams. The leading results in the $1/M_{U,D}$ and ω expansions are summarized as follows:

$$\ell_4^f = -2\ell_5^f = \left(\frac{\Lambda}{4\pi v}\right)^2 \frac{N}{12} > 0\,; \quad \ell_6^f = -\ell_7^f = -\left(\frac{\Lambda}{4\pi v}\right)^2 \frac{7N}{240}\omega^2\,, \quad \ell_{10} = 0\,; \tag{16}$$

where $\Lambda = \min\{M_U, M_D\}$.

3. Testing the Quartic Gauge Couplings via WWZ/ZZZ-Productions

While the LHC will give the first direct test on these new quartic gauge couplings (QGCs), the large backgrounds limit its sensitivity to the parameter-ℓ_n's and cutting off the backgrounds significantly reduces the event rate. As shown in Ref. [16], even for the direct resonance production in the TeV regime only around 10 signal events were predicted for $W^\pm W^\pm$ channels at the LHC with a 100 fb^{-1} annual luminosity after imposing necessary cuts in the gold-plated modes (by pure leptonic decays). The corresponding study at the TeV $e^\pm e^-$ NLC opens a much more exciting possibility [17].

The present study focuses on how to make further precision tests at the NLC for bounding these QGCs via WWZ/ZZZ-production [20] (cf. Sec. 3)[6] and their interplay with the WW-fusion [21] (cf. Sec. 4), which is much cleaner than the LHC so that the final state W/Z's can be detected via the dijet mode and with large branching ratios. Due to the limited calorimeter energy resolution, the misidentification probability of W versus Z and the rejection of certain fraction of diboson events should be considered [17]. Inclusion of the leptonic decay of Z to e^-e^+ and $\mu^-\mu^+$ is also useful. To avoid the potential fusion backgrounds $e^-e^+ \to eeZZ, eeWW$ in studying the WWZ/ZZZ-production, we only add the $\mu^-\mu^+$ channel for the Z-decay. Including these we find the detection efficiencies for ZZZ and WWZ final states are about 16.8% and 18.4%, respectively. The signal diagrams only contain the s-channel Z-boson so that the relevant QGCs come from $ZZZZ$ and $WWZZ$ vertices. It turns out that $e^-e^+ \to WWZ$ has huge backgrounds due to the t-channel ν_e or e-ν_e exchange, and the kinematic cuts alone help very little. However, we find that such type of backgrounds involve the left-handed W-e-ν coupling and thus can be very effectively suppressed by using the right(left)-hand polarized $e^-(e^+)$ beam. The highest sensitivity is reached by maximally polarizing *both* e^- and e^+ beams.

The crucial roles of the beam polarization and the higher collider energy for the WWZ-production are demonstrated in Fig. 1a, where $\pm 1\sigma$ exclusion contours for ℓ_4-ℓ_5 are displayed at $\sqrt{s}$ = 0.5, 0.8, 1.0 and 1.6 TeV, respectively. The beam polarization has much less impact on the ZZZ mode, due to the almost axial-vector type e-Z-e coupling. Including the same polarizations as in the case of the WWZ mode, we find about $10-20\%$ improvements on the bounds from the ZZZ-production. Assuming the two beam polarizations (90% e^- and 65% e^+), we summarize the final $\pm 1\sigma$ bounds for both ZZZ and WWZ channels and their combined 90% C.L. contours for 0.5 TeV with $\int \mathcal{L} = 50$ fb^{-1} in Fig. 1b (representing the *first direct probe* at the LC) and for 1.6 TeV with $\int \mathcal{L} = 200$ fb^{-1} in Fig. 1c (representing the *best* sensitivity gained from the final stage of the LC with energy around 1.5/1.6 TeV). We see that, at the 90% C.L. level, the bounds on ℓ_4-ℓ_5 at 0.5 TeV are within $O(10-20)$, while at 1.6 TeV they sensitively reach $O(1)$. The ellipses for the WWZ final state in ℓ_4-ℓ_5 plane are identical to those in ℓ_6-ℓ_7 plane, while the bands for the ZZZ final state in ℓ_6-ℓ_7 plane become tighter due to a factor of 2 enhancement from the 4Z-interaction vertex. ℓ_{10} only contributes to ZZZ final state and can be probed at the similar level. The new physics cutoff is chosen as $\Lambda = 2$ TeV in our plots and the numerical results for other values of Λ can be obtained by simple scaling. Finally, we have further performed a parallel analysis to Fig. 2b-c for the situation without e^+ beam polarization (with e^- polarization the same as before). For a two-parameter ($\ell_{4,5}$) study, the 90% C.L. results are compared as folllows:

$$\begin{array}{lll} \text{at 0.5 TeV}: & -12\ (-18) \le \ell_4 \le 21\ (27), & -17\ (-22) \le \ell_5 \le 9.5\ (15); \\ \text{at 1.6 TeV}: & -0.50\ (-0.67) \le \ell_4 \le 1.5\ (1.7), & -1.3\ (-1.5) \le \ell_5 \le 0.36\ (0.58); \end{array} \qquad (17)$$

where the numbers in the parentheses denote the bounds from polarizing the e^--beam alone. The comparison in (17) shows that without e^+-beam polarization, the sensitivity will decrease by about $15\% - 60\%$. Therefore, making use of the possible e^+-beam polarization with a degree around 65% will certainly be beneficial. In the above, the total rates are used to derive the numerical bounds. We have further studied the possible improvements by including different characteristic distributions, but no significant increase of the sensitivity is found for the above processes.

4. Interplay of WWZ/ZZZ-Production and WW-Fusion

To probe the QGCs (1), we know [9] that the WW-fusion amplitudes have the highest E-power dependence in the TeV regime while the s-channel signals of the WWZ/ZZZ-production lose an enhancement factor of $(E/v)^2$ relative to that of the fusion processes. When the collider energy is reduced by half (from 1.6 TeV down to 800 GeV), the sensitivity of the WW-fusion decreases by about a factor of 20 or more [21]. We therefore expect that $ee \to WWZ, ZZZ$ become more important at the earlier phase of the NLC and will be competitive with and complementary to fusions for the later stages of the NLC around $0.8 \sim 1$ TeV [20]. The following analysis reveals that even at the 1.5/1.6 TeV, $e^+e^- \to WWZ$ plays a crucial role in achieving a clean five-parameter analysis.

To completely determine all the QGCs, we need at least five independent processes. From WW-fusions alone, we can have

[6] The WWZ/ZZZ-production in the SM was studied in Ref. [18], and later some analyses on including the anomalous couplings have also appeared [19], but only for the case with unpolarized $e^\mp$ beams.

$$
\begin{array}{lll}
\text{Full process :} & \text{Sub} - \text{process :} & \text{Relevant parameter :} \\
e^-e^+ \to \nu\bar{\nu}W^-W^+\,, & (W^-W^+ \to W^-W^+), & (\ell_{4,5})\,, \\
e^-e^- \to \nu\bar{\nu}W^-W^-\,, & (W^-W^- \to W^-W^-), & (\ell_{4,5})\,; \\
e^-e^+ \to \nu\bar{\nu}ZZ\,, & (W^-W^+ \to ZZ), & (\ell_{4,5};\ \ell_{6,7})\,, \\
e^-e^+ \to e^{\pm}\nu W^{\mp}Z\,, & (W^{\mp}Z \to W^{\mp}Z), & (\ell_{4,5};\ \ell_{6,7})\,, \\
e^-e^+ \to e^-e^+ZZ\,, & (ZZ \to ZZ), & ([\ell_4+\ell_5]+2[\ell_6+\ell_7+\ell_{10}])\,.
\end{array} \tag{18}
$$

We see that for a complete determination the e^-e^- mode is necessary in opening the W^-W^- channel. The first two processes in (18) provide a clean test on $\ell_{4,5}$, and by including the third and fourth reactions $\ell_{6,7}$ can be further disentangled, and finally the fifth channel provides the unique probe on ℓ_{10}. Though this scheme is complete in principle, the realistic situation is much more involved. The small e-e-Z coupling suppresses the total rates of the last two channels (especially the fifth). Furthermore, the WZ-channel has large γ-induced $eeWW$ background in which one e is lost in the beam-pipe and one W misidentified as Z. A cut on the missing $p_\perp(\nu)$ is imposed to specially suppress this background. Even though, the final sensitivity still turns out to be less useful in constraining the ℓ_6-ℓ_7 space (cf. Fig. 2a below) [21]. To sensitively bound $\{\ell_6, \ell_7\}$ (especially ℓ_6) well below $O(1)$, we propose to use the triple gauge boson production mechanism $e^-e^+ \to WWZ$. Fig. 2a demonstrates the interplay of WW-fusion and WWZ-production for discriminating the $SU(2)_c$-breaking parameters ℓ_6-ℓ_7 at $\sqrt{s} = 1.6$ TeV and with an annual luminosity of 200 fb^{-1}. To constrain ℓ_{10}, both ZZZ and $eeZZ$ channels are available. Assuming that $\ell_{4,5;6,7}$ are constrained by the processes mentioned above, we set their values to be zero (the reference point) for simplicity and define the statistic significance $S = |\mathcal{N} - \mathcal{N}_0|/\sqrt{\mathcal{N}_0}$ which is a function of ℓ_{10}. (Here $\mathcal{N}$ is the total event-number while $\mathcal{N}_0$ is the number at $\ell_{10} = 0$.) As shown in Fig. 2b, at 1.6 TeV, the sensitivity of $e^-e^+ \to eeZZ$ for probing ℓ_{10} is better than that of $e^-e^+ \to ZZZ$.

In summary, the first direct probe on these QGCs will come from the early phase of the LC at 500 GeV, where the WW-fusion processes are not useful. The two mechanisms become more competitive and complementary at energies $\sqrt{s} \sim 0.8 - 1$ TeV. At a later stage of the LC, $\sqrt{s} = 1.6$ TeV, the 90% C.L. one-parameter bounds from the fusion processes become very sensitive, for $\Lambda = 2$ TeV:

$$
\begin{array}{c}
-0.13 \le \ell_4 \le 0.10\,, \quad -0.08 \le \ell_5 \le 0.06\,; \\
-0.22 \le \ell_6 \le 0.22\,, \quad -0.12 \le \ell_7 \le 0.10\,, \quad -0.21 \le \ell_{10} \le 0.21\,;
\end{array} \tag{19}
$$

obtained for $\int\mathcal{L} = 200$ fb^{-1} with a 90% (65%) polarized $e^-(e^+)$ beam. The bounds on $\ell_{4,5}$ are about a factor of $3 \sim 6$ stronger than that from WWZ/ZZZ-modes (cf. Table 1); while the bounds on $\ell_{6,7,10}$ are comparable. For a complete

multi-parameter analysis, the WWZ-channel is crucial for determining ℓ_6-ℓ_7 even at a 1.6 TeV LC.

Table 1: Combined 90% C.L. bounds on ℓ_{4-10} from WWZ/ZZZ-production. For simplicity, we set one parameter to be nonzero at a time. The bound on ℓ_{10} comes from ZZZ-channel alone.

$\sqrt{s}$ (TeV)	0.5	0.8	1.0	1.6
$\int\mathcal{L}$ (fb^{-1})	50	100	100	200
WWZ/ZZZ Bounds (at 90%C.L.)	$-9.5 \le \ell_4 \le 11.7$	$-2.7 \le \ell_4 \le 3.2$	$-1.7 \le \ell_4 \le 2.0$	$-0.50 \le \ell_4 \le 0.58$
	$-9.8 \le \ell_5 \le 8.9$	$-3.1 \le \ell_5 \le 2.3$	$-1.9 \le \ell_5 \le 1.4$	$-0.54 \le \ell_5 \le 0.36$
	$-5.0 \le \ell_6 \le 5.8$	$-1.5 \le \ell_6 \le 1.6$	$-0.95 \le \ell_6 \le 1.0$	$-0.28 \le \ell_6 \le 0.28$
	$-5.0 \le \ell_7 \le 5.7$	$-1.5 \le \ell_7 \le 1.5$	$-0.95 \le \ell_7 \le 0.92$	$-0.28 \le \ell_7 \le 0.26$
	$-4.3 \le \ell_{10} \le 5.2$	$-1.4 \le \ell_{10} \le 1.4$	$-0.83 \le \ell_{10} \le 0.88$	$-0.26 \le \ell_{10} \le 0.26$
Range of $\lvert\ell_n\rvert$	$\le O(4 \sim 10)$	$\le O(1 \sim 3)$	$\le O(0.8 \sim 2)$	$\le O(0.3 \sim 0.6)$

5. Concluding Remarks

Despite the constantly increasing evidence in supporting the Standard Model (SM) over the past 30 years, we particle physicists have been struggling in search for " *New Physics Beyond the SM* " so far [1]. Among the numerous ways for going beyond the SM, the Higgs boson hypothesis [22] stands out. Though the direct lower Higgs-mass-bound is gradually pushed up [1], the unitarity and triviality forbid it to go beyond the TeV scale, at which we are facing an exciting strong electroweak symmetry breaking (EWSB) dynamics. Below the new heavy resonance, we have to *first* probe the EWSB parameters formulated by means of the electroweak chiral Lagrangian (EWCL), among which the quartic gauge interactions penetrate the pure Goldstone dynamics. After analyzing the different patterns of these quartic couplings in connection with typical underlying resonance/non-resonance models, we perform a model-independent study on constraining them via WWZ/ZZZ-production and WW-fusion at the next generation $e^{\pm}e^-$ linear colliders (NLC). The main focus is then put onto the *interplay* of these two production mechanisms for achieving a complete probe of the EWSB mechanism. The important roles of both polarized e^- and e^+ beams are revealed and analyzed.

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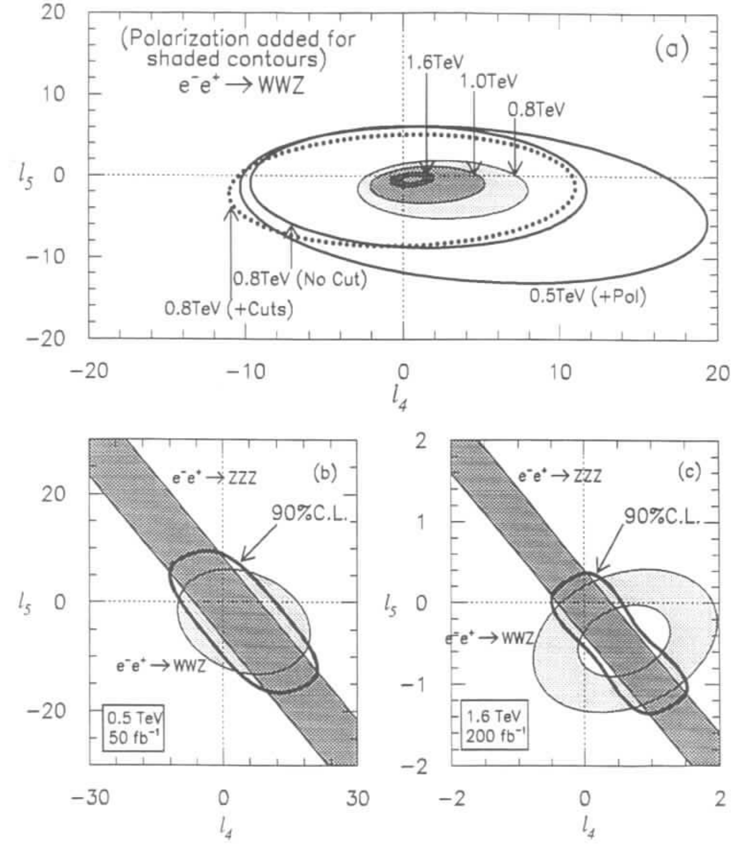
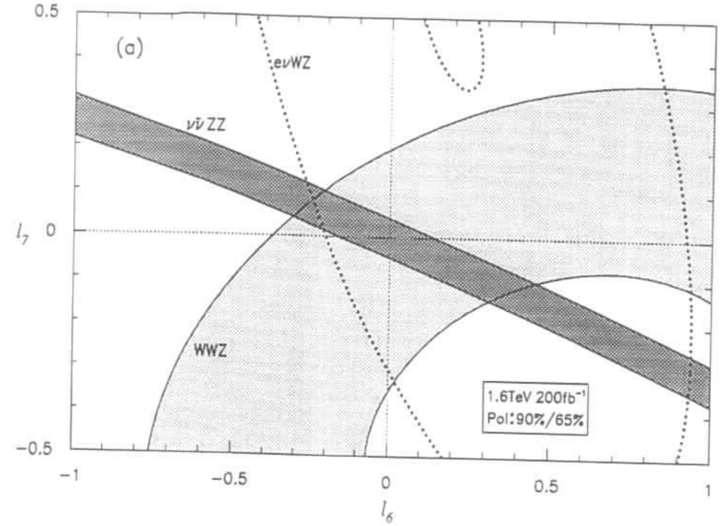


FIGURE 1. Probing ℓ_4-ℓ_5 via WWZ and ZZZ production processes. The roles of the polarization and the higher collider energy for $e^-e^+ \to WWZ$ are shown by the $\pm 1\sigma$ exclusion contours in (a). The integrated luminosities used here are 50 fb^{-1} (at 500 GeV), 100 fb^{-1} (at 800 GeV) and 200 fb^{-1} (at 1.0 and 1.6 TeV). In (b) and (c), the $\pm 1\sigma$ contours are displayed for ZZZ/WWZ final states at $\sqrt{s}$ =0.5 and 1.6 TeV respectively, with two beam polarizations (90% e^- and 65% e^+); the thick solid lines present the combined bounds at 90% C.L.



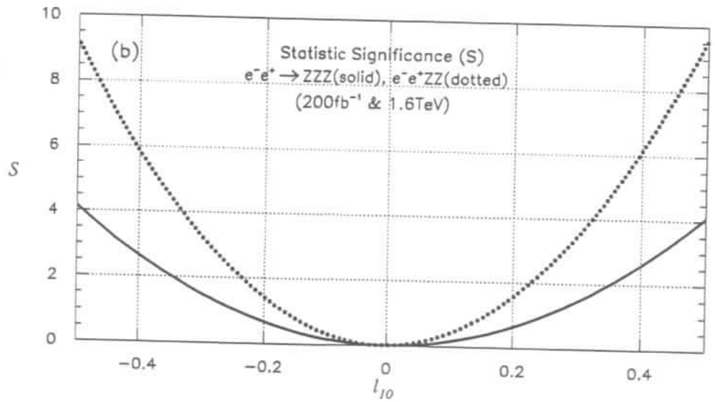


FIGURE 2. Interplay of the WW-fusion and WWZ/ZZZ-production for discriminating ℓ_6-ℓ_7 and ℓ_{10} at $\sqrt{s}$ =1.6 TeV with $\int \mathcal{L}$ =200 fb^{-1}: (a). $\pm 1\sigma$ exclusion contours for $e^-e^+ \to \nu\bar{\nu}ZZ$, $e^+\nu W^-Z/e^-\bar{\nu}W^+Z$, and $e^-e^+ \to WWZ$ with polarizations (90% e^- and 65% e^+). (b). Statistic significance versus ℓ_{10} for $e^-e^+ \to ZZZ$, e^-e^+ZZ (with unpolarized $e^\mp$ beams).

References

1. B.C. Allanach et al, *Report of the Working Group on 'Searches'*, hep-ph/9708250; A. Sopczak, *Searches for Higgs Bosons at LEP2*, hep-ph/9712283; and references therein.
2. M.S. Chanowitz, Phys. Rev. Lett. (1998), in press, and LBNL-40877 (hep-ph/9710308).
3. T. Appelquist and J. Carazzone, Phys. Rev. **D11** (1975) 2856.
4. For a nice review, H. Georgi, Ann. Rev. Nucl. & Part. Sci. **43** (1994) 209.
5. M.E. Peskin, Prog. Theor. Phys. Suppl. **123** (1996) 507; H.E. Haber, hep-ph/9703381; G.L. Kane, hep-ph/9709318.
6. E. Accomando et al, (ECFA/DESY LC Physics Working Group), Phys. Rep. (1998), in press, and DESY-97-100 (hep-ph/9705442); H. Murayama and M.E. Peskin, Ann. Rev. Nucl. & Part. Sci. **46** (1996) 533.
7. T. Appelquist and C. Bernard, Phys. Rev. **D22** (1980) 200; A.C. Longhitano, Nucl. Phys. **B188** (1981) 118; T. Appelquist and G.-H. Wu, Phys. Rev. **D48** (1993) 3235; and references therein.
8. S. Weinberg, Physica **96A** (1979) 327.
9. H.-J. He, Y.-P. Kuang, and C.-P. Yuan, in these proceedings, hep-ph/9708402 and DESY-97-141; Phys. Rev. **D55** (1997) 3038, Mod. Phys. Lett. **A11** (1996) 3061; Phys. Lett. **B382** (1996) 149; for an updated comprehensive review, hep-ph/9704276 and DESY-97-056, Lectures in the Proceedings of the CCAST (World Laboratory) Workshop on *Physics at TeV Energy Scale*, Vol. 72, pp.119-234.
10. J. Goldstone, Nuovo Cim. **19** (1961) 154.
11. B.W. Lee, C. Quigg and H. Thacker, Phys. Rev. **D16** (1977) 1519; M.S. Chanowitz and M.K. Gaillard, Nucl. Phys. **B261** (1985) 379.
12. H.-J. He and W.B. Kilgore, Phys. Rev. **D55** (1997) 1515; H.-J. He, Y.-P. Kuang, C.-P. Yuan, Phys. Rev. **D51** (1995) 6463; H.-J. He, Y.-P. Kuang, and X. Li, Phys. Rev. Lett. **69** (1992) 2619; Phys. Rev. **D49** (1994) 4842; Phys. Lett. **B329** (1994) 278; and references therein.
13. M. Peskin, Talk at the Snowmass Conference (June, 1996), and the working group summary report, hep-ph/9704217.
14. D. Dominici, in these proceedings; R. Casalbuoni, et al, Phys. Rev. **D53** (1996) 5201; and references therein.
15. C.G. Callan, S. Coleman, J. Wess, B. Zumino, Phys. Rev. **177** (1969) 2247.
16. J. Bagger, V. Barger, K. Cheung, J. Gunion, T. Han, G.A. Ladinsky, R. Rosenfeld, and C.-P. Yuan, Phys. Rev. **D49** (1994) 1246; **D52** (1995) 3878.
17. V. Barger, K. Cheung, T. Han, and R.J.N. Phillips, Phys. Rev. **D52** (1995) 3815.
18. V. Barger and T. Han, Phys. Lett. **B212** (1988) 117; V. Barger, T. Han, and R.J.N. Phillips, Phys. Rev. **D39** (1989) 146.
19. G. Belanger and F. Boudjema, Phys. Lett. **B288** (1992) 201; S. Dawson, A. Likhoded, G. Valencia, and O. Yuschenko, hep-ph/9610299,

and private communications.
20. T. Han, H.-J. He, and C.-P. Yuan, Phys. Lett. **B422** (1998) 294 and hep-ph/9711429.
21. H.-J. He, DESY-97-037, in the proceedings of " *The Higgs Puzzle* ", pp.207-217, Ringberg, Munich, Germany, December 8-13, 1996, Ed. B. Kniehl, World Scientific Pub.; E. Boos, H.-J. He, W. Kilian, A. Pukhov, C.-P. Yuan, and P.M. Zerwas, Phys. Rev. **D57** (1998) 1553 and hep-ph/9708310.
22. P.W. Higgs, Phys. Lett. **12** (1964) 132; Phys. Rev. Lett. **13** (1964) 508; F. Englert and R. Brout, Phys. Rev. Lett. **13** (1964) 321; G.S. Guralnik, C.R. Hagen, and T.W.B. Kibble, Phys. Rev. Lett. **13** (1964) 585.