Low Energy Constraints and Anomalous Triple Gauge Boson Couplings

Dieter Zeppenfeld

Department of Physics, University of Wisconsin, Madison, WI 53706, USA

ABSTRACT

Low energy (1-loop) constraints on anomalous triple gauge boson vertices (TGV’s) are revisited and compared to the sensitivity achievable at LEP II and at future linear e\(^+\)e\(^-\) colliders. The analysis is performed within the framework of an effective Lagrangian of gauge invariant dimension six operators with the gauge bosons and a single Higgs doublet field as the low energy degrees of freedom. The low energy data do not directly bound TGV’s but they provide strong constraints on models which lead to anomalous gauge boson interactions in addition to other low energy effects.

1. Introduction

Over the last four years e\(^+\)e\(^-\) collision experiments at LEP and at the SLAC linear collider have beautifully confirmed the predictions the Standard Model (SM). At present experiment and theory generally agree at the 1% level or better in the determination of the vector boson couplings to the various fermions, which may rightly be considered a confirmation of the gauge boson nature of the W and the Z. Nevertheless the most direct consequence of the SU(2) \(\times\) U(1) gauge symmetry, the nonabelian self-couplings of the W, Z, and photon, remain poorly measured to date. Even if the underlying theory is SU(2) \(\times\) U(1) invariant, novel strong interactions in the gauge boson–Higgs sector may lead to anomalous WWZ and WW\(\gamma\) couplings.

One of the major reasons for raising the energy of the LEP collider above the W-pair threshold is the systematic study and measurement of these triple gauge boson vertices (TGV’s) via the process e\(^+\)e\(^-\) \(\rightarrow\) W\(^+\)W\(^-\) W pair production together with measurements of the single W production cross section at a future linear e\(^+\)e\(^-\) or e\(\gamma\) collider will provide us with an excellent measurement of the three vector boson couplings. One can quantify the sensitivity of all these experiments by parameterizing the most general WWV \((V = Z, \gamma)\) vertex in terms of an effective Lagrangian \(\mathcal{L}_{eff}^{WWV}\). Considering \(C\) and \(P\) even couplings only, it takes the form:

\[
\mathcal{L}_{eff}^{WWV} = g_{WWV} \left( g_V^V (W_{\mu
u}^\dagger W^\mu - W^{\dagger \mu} W_{\mu
u}) V_{\nu} + \kappa_V W_{\mu}^\dagger W_{\nu} V_{\mu\nu} + \frac{\lambda_V}{m_W^2} W_{\mu}^\dagger W_{\mu} W_{\nu} V_{\mu\nu} \right).
\]

(1)

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Here the overall coupling constants are defined as $g_{WW\gamma} = e$ and $g_{WWZ} = e \cot \theta_W$. Within the SM the couplings are given by $g_1^2 = g_1^2 = \kappa_Z = \kappa_\gamma = 1$, $\lambda_Z = \lambda_\gamma = 0$.

While the value of $g_1^2$ is fixed by electromagnetic gauge invariance (it is just the electric charge of the $W^+$) the other couplings may well deviate from their SM values and need to be determined experimentally. At LEP II one expects a sensitivity to deviations from the SM predictions of $\Delta \kappa \approx \Delta \lambda \approx 0.1 ... 0.2$ while the future $e^+e^-$ linear colliders will push the precision of these measurements to the 1\% level or below.

2. Effective Lagrangians

The question arises whether the present high precision measurements at LEP and at lower energies already give comparable constraints via 1-loop corrections to $S$-matrix elements which involve the $WWV$ vertices. Many such investigations have been performed in the past, usually, however, in a framework which introduces the deviations from the SM in such a way as to violate $SU(2) \times U(1)$ gauge-invariance. As a result the 1-loop contributions from anomalous $WWV$ interactions to oblique parameters like $\delta \rho$ or the $S, T, U$ parameters of Peskin and Takeuchi turn out to be quadratically or even quartically divergent. This in turn has lead to very stringent bounds from existing low-energy data.

While the effective Lagrangian $L_{WWV}$ of Eq.(1) is general enough for a discussion of weak boson pair production, low energy observables are affected at the 1-loop level not only by the TGV’s. One also expects contributions from other new interactions which are induced by the new physics simultaneously with anomalous values of $g_1$, $\kappa$, or $\lambda$. In order to take such effects into account while avoiding an inflation of free parameters, some simplifying assumptions are needed.

Given the excellent agreement of the measured fermion couplings with the SM gauge theory predictions, I shall assume in the following that

i) The $W$, $Z$, and photon are indeed the gauge bosons of a spontaneously broken $SU(2) \times U(1)$ local symmetry.

ii) New contributions to the gauge boson–fermion couplings can be neglected.

iii) The low energy effects of the new interactions which are responsible for anomalous $WWV$ couplings are described by an effective Lagrangian with the $SU(2) \times U(1)$ gauge fields and the Higgs doublet field as the low energy degrees of freedom:

$$L_{\text{eff}} = \sum_i \frac{f_i}{\Lambda^2} O_i + \sum_i \frac{f^{(8)}_{i}}{\Lambda^4} O^{(8)}_i + \ldots .$$

Here the scale $\Lambda$ may be identified with the typical mass of new particles associated with the new physics. Because of assumptions i) and ii) only gauge invariant operators $O_i$ are allowed which can be constructed out of the Higgs field $\Phi$, covariant
derivatives of the Higgs field, \( D_\mu \Phi \), and the field strength tensors \( W_{\mu\nu} \) and \( B_{\mu\nu} \) of the \( W \) and the \( B \) gauge fields:

\[
[D_\mu, D_\nu] = \hat{B}_{\mu\nu} + \hat{W}_{\mu\nu} = i \frac{g'}{2} B_{\mu\nu} + i g \frac{\sigma^a}{2} W^a_{\mu\nu}.
\] (3)

The use of a linear realization of the electroweak symmetry breaking in terms of the Higgs doublet field \( \Phi \) allows to discuss Higgs mass effects in the following. It is general enough, though, since nonlinear realizations of the symmetry breaking sector can be simulated by the \( m_H \rightarrow \Lambda \) limit. As has been emphasized by Burgess and London, the gauge invariance assumption does not really provide any constraints on e.g. anomalous TGV’s induced by \( \mathcal{L}_{\text{eff}} \), since the phenomenological Lagrangian \( \mathcal{L}_{WWV} \) can be regarded as the unitary gauge version of an explicitly \( SU(2) \times U(1) \) invariant effective Lagrangian. Constraints arise when making one additional assumption:

iv) The effective Lagrangian may be truncated at the dimension six level, \( i.e. \) corrections of order \( m_W^2/\Lambda^2 \) or \( v^2/\Lambda^2 \) can be neglected in the low energy effects.

This last assumption, while limiting the applicability of the subsequent analysis somewhat, is general enough to elucidate the generic problems of low energy constraints on the \( WWV \) couplings, as we shall see later.

A complete list of \( SU(2) \times U(1) \) invariant dimension six operators has been given in Ref. 12 and has by now been employed in the analysis of \( WWV \) couplings by many authors. Using the SM equations of motion for the Higgs doublet field and identifying operators which only differ by a total derivative, 11 independent operators can be constructed at the dimension six level. Of these only 9 contribute to four-fermion amplitudes up to 1-loop:

\[
\mathcal{L}_{\text{eff}} = \sum_{i=1}^{9} \frac{f_i}{\Lambda^2} \mathcal{O}_i = \frac{1}{\Lambda^2} \left( f_{\Phi,1} (D_\mu \Phi)^\dagger \Phi^i (D^\mu \Phi) + f_{BW} \Phi^i \hat{B}_{\mu\nu} \hat{W}^{\mu\nu} \Phi 
+ f_{DW} Tr([D_\mu, \hat{W}_{\nu\rho}] [D^\mu, \hat{W}^{\nu\rho}]) - f_{DB} \frac{g^2}{2} (\partial_\mu B_{\nu\rho})(\partial^\mu B^{\nu\rho})
+ f_B (D_\mu \Phi)^\dagger \hat{B}^{\mu\nu} (D_\nu \Phi) + f_W (D_\mu \Phi)^\dagger \hat{W}^{\mu\nu} (D_\nu \Phi) + f_{WWW} Tr[\hat{W}_{\mu\nu} \hat{W}^{\mu\nu} \hat{W}_\rho] 
+ f_{WW} \Phi^i \hat{W}_{\mu\nu} \hat{W}^{\mu\nu} \Phi + f_{BB} \Phi^i \hat{B}_{\mu\nu} \hat{B}^{\mu\nu} \Phi \right),
\] (4)

The first four operators, \( \mathcal{O}_{\Phi,1}, \mathcal{O}_{BW}, \mathcal{O}_{DW}, \) and \( \mathcal{O}_{DB} \), affect the gauge boson two-point functions at tree level and as a result the coefficients of these four operators are severely constrained by present low energy data.

Of the remaining five, \( \mathcal{O}_{WWW}, \mathcal{O}_W \) and \( \mathcal{O}_B \) give rise to non-standard triple gauge boson couplings. Their presence in the effective Lagrangian leads to deviations of the \( WWV \) couplings from the SM, namely:

\[
\kappa_\gamma = 1 + (f_B + f_W) \frac{m_W^2}{2\Lambda^2}, \quad \kappa_Z = 1 + \left( f_W - s^2(f_B + f_W) \right) \frac{m_Z^2}{2\Lambda^2},
\] (5a)
\[ g_1^2 = 1 + f_W \frac{m_Z^2}{2\Lambda^2} = \kappa_Z + \frac{s^2}{c^2}(\kappa_\gamma - 1), \quad (5b) \]
\[ \lambda_\gamma = \lambda_Z = \frac{3m_W^2g^2}{2\Lambda^2} f_{WW} = \lambda, \quad (5c) \]

with \( s = \sin \theta_W \). As mentioned earlier, the correlations between different anomalous \( WWV \) couplings exhibited in the last two equations are due to the truncation of the effective Lagrangian at the dimension six level and do not hold any longer when dimension eight operators are included.\(^{14}\) In addition to anomalous triple boson vertices the operators of Eq. (5) provide anomalous gauge boson Higgs couplings, and these additional interactions play a very important role in canceling the divergencies which have plagued earlier loop analyses.

### 3. Low Energy Observables

The effective Lagrangian of Eq. (4) can now be used to calculate loop corrections involving anomalous \( WWV \) couplings. A range of observables have been considered in the past. Classical examples are \((g - 2)_\mu\)\(^{20}\) and the \( b \to s\gamma \) decay rate.\(^{21}\) I will not discuss these processes in detail here because of two reasons: i) the oblique corrections give more stringent constraints on the 11 operators considered above and ii) a complete discussion along the lines given below for the oblique corrections necessitates the introduction of additional operators which involve fermions.\(^{22}\) In the case of the anomalous magnetic moment of the muon this is e.g. the magnetic moment operator

\[ \mathcal{O}_{g-2} = \frac{m_\mu}{v} \bar{L} \sigma^{\mu\nu} (a_W \tilde{W}^{\mu\nu} + a_B \tilde{B}^{\mu\nu}) \Phi R, \quad (6) \]

where \( L \) denotes the left-handed \((\nu_\mu, \mu)\) doublet field. These direct tree level contributions need to be considered together with the loop corrections involving \( WWV \) couplings in order to perform a model-independent analysis of the low energy bounds and for absorbing the divergencies of the loop integration.\(^{20}\)

These problems appear as well in the analysis of oblique corrections involving the operators which were discussed in the previous section, and my discussion here closely follows the one in Ref. 16. A vast amount of experimental data can be understood as the measurement of 4-fermion \( S \)-matrix elements. This includes the recent LEP data, neutrino scattering experiments, atomic parity violation, \( \mu \)-decay, and the \( W \)-mass measurement at hadron colliders. Since the data are now sensitive to electroweak loop-corrections, these SM corrections must be considered at the same time as the new physics contributions. After correcting for SM box contributions and non-universal vertex corrections (in particular the top mass dependence of the \( Zb\bar{b} \)-vertex) the remaining SM contributions as well as all divergent new physics contributions can be parameterized in a simple way. The four-fermion amplitudes for massless external fermions are given by

\[ \mathcal{M}(p_1, p_2, p_3, p_4) = I(q^2) J_\mu(p_1, p_2) J^\mu(p_3, p_4). \quad (7) \]
Here the $J_\mu$ only depend on the wave functions of the external fermions and the helicity dependent $I(q^2)$ are given by

$$I_{CC}(q^2) = \frac{\bar{g}_W^2(q^2)/2}{q^2 - m_W^2 + im_W \Gamma_W}$$

(8)

for CC amplitudes of left-handed fermions, while NC amplitudes may be written as

$$I_{NC}(q^2) = \frac{\bar{e}^2(q^2)}{q^2} Q_{f_i} Q_{f_3} + \frac{\bar{g}_Z^2(q^2)}{q^2 - m_Z^2 + i m_Z \Gamma_Z} \left( T_{3 f_i}^f - \bar{s}^2(q^2) Q_{f_i} \right) \left( T_{3 f_3}^f - \bar{s}^2(q^2) Q_{f_3} \right).$$

(9)

$Q_{f_i}$ denotes the electric charge and $T_{3 f_i}$ the third component of the weak isospin of fermion $f_i$.

The free parameters, which need to be determined by experiment, are the four form-factors $\bar{e}^2(q^2)$, $\bar{g}_W^2(q^2)$, $\bar{g}_Z^2(q^2)$, and $\bar{s}^2(q^2)$ and the $W$ and $Z$ mass. Three measurements are needed to define the parameters of the SM, and these may be chosen as $m_Z$, the Fermi constant $G_F \propto \bar{g}_W^2(0)/m_W^2$, and $\alpha = \bar{e}^2(0)/4\pi$. Only after these values have been fixed can the remaining data be used to place constraints on new physics contributions.

An analysis of the available data has recently been performed by Hagiwara et al. For $m_t = 140$ GeV the LEP and SLC data can be summarized in terms of

$$\bar{g}_Z^2(m_Z^2) = 0.5524 \pm 0.0017, \quad \bar{s}^2(m_Z^2) = 0.2319 \pm 0.0011.$$  \hspace{1cm} (10)

A slight top mass dependence of the extracted results is negligible compared to the errors. In a similar fashion the low-energy data on neutrino scattering and atomic parity violation determine the same form-factors at zero momentum transfer:

$$\bar{g}_Z^2(0) = 0.5462 \pm 0.0035, \quad \bar{s}^2(0) = 0.2359 \pm 0.0048.$$  \hspace{1cm} (11)

Finally, the $W$-mass measurement at hadron colliders together with the input value of $G_F$ can be translated into a measurement of $\bar{g}_W^2(0)$:

$$\bar{g}_W^2(0) = 0.4217 \pm 0.0027.$$  \hspace{1cm} (12)

These five measurements are closely related to other formulations of the oblique corrections, like the $S,T,$ and $U$ parameters of Peskin and Takeuchi. S and T, for example, are given by

$$S = \frac{4 \bar{s}^2(m_Z^2) \bar{e}^2(m_Z^2)}{\alpha(m_Z^2)_{SM}} - \frac{16\pi}{\bar{g}_Z^2(0)};$$

(13a)

$$1 - \alpha T = \frac{1}{\bar{\rho}} \frac{\bar{g}_W^2(0)}{\bar{g}_Z^2(0)} \frac{m_Z^2}{m_W^2}.$$  \hspace{1cm} (13b)
The new feature here is the inclusion of the $q^2$ dependence of the form-factors. Indeed, new physics contributions like the operators $\mathcal{O}_{DW}$ or $\mathcal{O}_{DB}$ do lead to a non-trivial $q^2$ dependence of the form-factors in Eq. (9), and the more general analysis is needed to constrain these operators. Low energy bounds are obtained by fitting

\begin{align}
S &= S_{SM}(m_t, m_H) + \Delta S, \\
T &= T_{SM}(m_t, m_H) + \Delta T\text{ etc.}
\end{align}

(14a)

(14b)

to the data. Here the SM contributions ($S_{SM}$ etc.) introduce a significant dependence on the as yet unknown values of the Higgs and the top quarks masses.

The four operators $\mathcal{O}_{DW}, \mathcal{O}_{DB}, \mathcal{O}_{BW}$, and $\mathcal{O}_{\Phi,1}$, contribute already at tree level,

\begin{align}
\Delta \delta \rho &= \alpha \Delta T = -\frac{v^2}{2\Lambda^2} f_{\Phi,1}, \\
\Delta S &= -32\pi s m_W^2 \frac{m_H^2}{\Lambda^2} (f_{DW} + f_{DB}) - 4\pi v^2 \frac{f_{BW}}{\Lambda^2},
\end{align}

(15a)

(15b)

with similar results for the other form-factors. Fitting these to the five data points one obtains measurements of the coefficients of the operators in the effective Lagrangian,

\begin{align}
f_{DW}/\Lambda^2 &= (0.56 \pm 0.79) \text{ TeV}^{-2}, \\
f_{DB}/\Lambda^2 &= (-8.0 \pm 11.9) \text{ TeV}^{-2}, \\
f_{BW}/\Lambda^2 &= (1.9 \pm 2.9) \text{ TeV}^{-2}, \\
f_{\Phi,1}/\Lambda^2 &= (0.11 \pm 0.20) \text{ TeV}^{-2},
\end{align}

(16a)

(16b)

(16c)

(16d)

for $m_H = 200$ GeV and $m_t = 140$ GeV. While the central values depend on the choice of $m_t$ and $m_H$, the quoted errors are unaffected. There are strong correlations between the coefficients of the dimension six operators, however, in particular between $f_{DB}, f_{BW}$ and $f_{\Phi,1}$.

While the contributions of these four operators are already constrained at the tree level, the remaining five, which include the anomalous $WWV$ couplings, only contribute at the 1-loop level to the oblique correction parameters. Contributions to the four-fermion amplitudes arise via the corrections to the gauge boson self-energies and also to the gauge boson–fermion vertices. In fact both need to be included to preserve gauge invariance. The complete calculation of the logarithmically enhanced contributions was performed in Ref. 16, partial results can be found in Refs. 13–15.

At intermediate steps of the calculation one still encounters quadratic divergencies due to the insertion of dimension six operators in the loops. These are all absorbed, however, into the renormalization of the SM parameters $m_Z, G_F$ and $\alpha$. All remaining logarithmically divergent terms are found to be renormalizations of the four operators which already contributed at tree level. Neglecting all terms which are not logarithmically enhanced, the leading effects are given by replacing $f_{DW}$ etc. in Eq. (15) by the
After renormalization of the SM parameters and of the operators which contribute at
tree level, the remaining corrections to four-fermion amplitudes are finite. The log \( \Lambda^2/\mu^2 \) terms in Eq. (17) describe mixing of the operators between the new physics scale \( \Lambda \) and the weak boson mass scale \( \mu = m_W \).

The quadratic divergencies observed in earlier work are cancelled by Higgs contributions to the vacuum polarization of the \( W \) and the \( Z \): \( SU(2) \times U(1) \) gauge invariance and the use of a linear realization for the symmetry breaking sector relates TGV’s to anomalous Higgs-gauge boson interactions. Gauge invariance guarantees the cancellation of all quadratic divergencies between gauge boson and Higgs contributions. A trace of the quadratic divergencies is preserved in the \( m_H^2 \) terms in the results of Eqs. (17): the Higgs graphs give rise to \(-\Lambda^2 + m_H^2 \log \Lambda \) terms. By including Higgs exchange we have therefore replaced quadratic divergencies by \( m_H^2 \) terms. In the limit \( m_H \to \Lambda \) the quadratic divergencies are recovered.

4. Low Energy Bounds on Anomalous \( WWV \) Couplings?

We have seen that all divergent 1-loop contributions involving TGV’s are just renormalizations of the coefficients of some other, independent, operators. Hence these \( \Lambda^2 \) or log \( \Lambda \) terms cannot be used for a direct measurement of the \( WWV \) couplings without making assumptions on the absence of cancellations between tree level and 1-loop contributions. Even including the finite corrections involving \( WWV \) couplings the five data points of Eqs. (10–12) are barely sufficient to limit the four tree level coefficients \( f_{DW}/\Lambda^2, f_{DB}/\Lambda^2, f_{BW}/\Lambda^2, \) and \( f_{\phi,1}/\Lambda^2 \) in addition to the SM top quark and Higgs boson mass dependences. Without additional assumptions the \( WWV \) couplings remain unconstrained by the present low-energy data.

One may assume, for example, that \( f_{BW} \) vanishes at the scale \( \Lambda = 1 \text{ TeV} \) and that the main contribution to \( f_{BW}' \) arises at 1-loop from \( \lambda = 3 g^2 f_{WWW} m_W^2/2\Lambda^2 \). Eq. (16c) then translates into

\[ \lambda = 0.89 \pm 1.35 \, , \]  
(18)
a constraint which is comparable to present hadron collider bounds.\textsuperscript{2}

The more traditional analysis of low energy bounds assumes that only one coupling differs from its SM value, hence that no cancellations occur between the contributions from different couplings. In the context of our 11 dimension-six operators this corresponds to considering the three cases $f_B \neq 0$, $f_W \neq 0$, and $f_{WWW} \neq 0$ with the coefficients of the remaining 10 operators vanishing. Choosing $m_t = 140$ GeV and $m_H = 200$ GeV one finds\textsuperscript{4}

\begin{align}
\kappa_\gamma &= 1 + f_B \frac{m_W^2}{2 \Lambda^2} = 1.04 \pm 0.06 \quad \text{for } f_B \neq 0 , \\
\kappa_\gamma &= 1 + f_W \frac{m_W^2}{2 \Lambda^2} = 1.01 \pm 0.09 \quad \text{for } f_W \neq 0 , \\
\lambda_\gamma &= \frac{3 m_W^2 g^2}{2 \Lambda^2} f_{WWW} = 0.03 \pm 0.16 \quad \text{for } f_{WWW} \neq 0 .
\end{align}

A more stringent assumption has been proposed by De Rújula et al.\textsuperscript{13} There are no obvious symmetries which distinguish the tree level operators $O_{BW}$, $O_{DW}$, $O_{DB}$, and $O_{\Phi,1}$ from the remaining seven. For a generic model of the underlying dynamics one may hence expect e.g. $|f_B + f_W| \approx |f_{BW}|$ which with the result of Eq. (16c) implies $|\kappa_\gamma - 1| = |f_B + f_W| \frac{m_W^2}{2 \Lambda^2} < 0.02$ at ”90% CL”, a value too small to be observable in $W^+W^-$ production at LEP, but still in the interesting range for future linear colliders.

As Einhorn et al.\textsuperscript{17} have argued, one should perhaps not expect anomalous couplings which are larger than these most stringent bounds. In extensions of the SM anomalous $WWV$ couplings arise from heavy particle loops (of mass $M$) with three external gauge bosons attached. Because of the universal factor $1/16\pi^2$ for loop integrals one should expect

\begin{equation}
\frac{f_i}{\Lambda^2} = \frac{1}{16\pi^2} \frac{c_f}{M^2} ,
\end{equation}

and hence $|f_i| m_W^2/\Lambda^2 < 10^{-3}$ even for masses as low as $M = 250$ GeV, unless the counting factor $c_f$ is substantially larger than unity, e.g. due to higher isospin multiplets or because of large multiplicities of the heavy particles.

The last two arguments indicate the difficulty of constructing realistic models which would predict large anomalous $WWV$ couplings. One must clearly state, however, that there is no proof that large anomalous couplings are ruled out, and naturalness arguments may well prove erroneous. The eleven dimension-six operators of Section II are independent and must therefore be constrained individually by experiment. For the $WWV$ couplings, $W^+W^-$ production and single $W$ production at future linear colliders are the ideal way to achieve this goal.
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