SELF-COUPLINGS OF ELECTROWEAK BOSONS:
THEORETICAL ASPECTS AND TESTS AT HADRON COLLIDERS

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ABSTRACT

The current theoretical understanding of anomalous gauge boson couplings is reviewed, and the direct measurement of these couplings in present and future hadron collider experiments is briefly discussed.

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1 Introduction

Although the electroweak Standard Model (SM) based on an \( SU(2) \times U(1) \) gauge theory has been very successful in describing contemporary high energy physics experiments, the three-vector-boson couplings predicted by this non-Abelian gauge theory remain largely untested experimentally. A direct measurement of these vector boson couplings is possible in present and future collider experiments, the three vector-boson couplings predicted successful in describing contemporary high energy physics.

Analogous to the introduction of arbitrary vector and axial vector couplings to fermions, the measurement of the \( W W V \), \( Z \gamma V \), \( W \gamma V \), \( Z \gamma, Z W \) vertices here presents a short overview of the theoretical aspects of vector boson self-couplings, and discuss their measurement in hadron collider experiments. Tests at LEP II and Linear Colliders are discussed in Refs. 1 and 2.

2 Theoretical Aspects

Analogous to the introduction of arbitrary vector and axial vector couplings \( g_v \) and \( g_A \) for the coupling of gauge bosons to fermions, the measurement of the WWV \( (V = W, Z) \) couplings can be made quantitative by introducing a more general WWV vertex. For our discussion of experimental sensitivities in Section 3 we shall use a parameterization in terms of the phenomenological effective Lagrangian (other, equivalent, parameterizations are possible):

\[
iL_{WWV} = g_{WWV} \left[ g_1^V (W^\dagger_\mu W^\nu V^\mu - W^\dagger_\mu V^\mu W^\nu) + \kappa_V W^\dagger_\mu W^\nu V^\mu + \frac{\lambda_V}{m_W^2} W^\dagger_\mu W^\nu V^\mu \right] . \tag{1}\]

Here the overall coupling constants are defined as \( g_{WWV} = e \) and \( g_{WWZ} = e \cot \theta_W \). \( W^\dagger_\mu W^\nu \) is defined as \( \partial_\mu W^\nu - \partial_\nu W^\mu \). \( V^\dagger_\mu V^\nu \) is defined as \( \partial_\mu V^\nu - \partial_\nu V^\mu \). Within the SM, at tree level, the couplings are given by \( g_1^Z = g_1^\gamma = \kappa_Z = \kappa_\gamma = 1 \), \( \lambda_Z = \lambda_\gamma = 0 \). \( g_1^V = 1 \) is fixed by electromagnetic gauge invariance. \( g_1^Z \), however, may well be different from its SM value 1 and appears at the same level as \( \kappa_\gamma \) or \( \kappa_Z \).

The effective Lagrangian of Eq. (1) parameterizes the most general Lorentz invariant and \( C \) and \( P \) conserving WWV vertex which can be observed in processes where the vector bosons couple to effectively massless fermions. Terms with higher derivatives are equivalent to a dependence of the couplings on the vector boson momenta and thus merely lead to a form-factor behaviour of these couplings. Analogous to the general WWV vertex it is possible to parameterize anomalous \( Z \gamma V, V = \gamma, Z \) couplings in terms of two coupling constants, \( h_3^V \) and \( h_4^V \) if CP invariance is imposed. All \( Z \gamma V \) couplings are \( C \) odd.

In the absence of a specific model of new physics, effective Lagrangian techniques are extremely useful. An effective Lagrangian parameterizes, in a model-independent way, the low-energy effects of the new physics to be found at higher energies. It is only necessary to specify the particle content and the symmetries of the low-energy theory. Although effective Lagrangians contain an infinite number of terms, they are organized in powers of \( 1/\Lambda \), where \( \Lambda \) is the scale of new physics. Thus, at energies which are much smaller than \( \Lambda \), only the first few terms of the effective Lagrangian are important. Since all experimental evidence is consistent with the existence of an \( SU(2) \times U(1) \) gauge symmetry it is natural to require the effective Lagrangian describing anomalous gauge boson couplings to possess this invariance. How this symmetry is realized depends on the particle content of the effective Lagrangian. If one includes a Higgs boson, the symmetry can be realized linearly, otherwise a nonlinear realization of the gauge symmetry is required.

If the \( SU(2) \times U(1) \) symmetry is realized linearly and the analysis is restricted to operators of dimension 6, three operators give rise to anomalous triple gauge boson couplings:

\[
\mathcal{L}_{eff} = \frac{1}{\Lambda^2} [f_B \mathcal{O}_B + f_W \mathcal{O}_W + f_{WWW} \mathcal{O}_{WWW}] . \tag{2}\]

The explicit form of the three operators \( \mathcal{O}_B, \mathcal{O}_W \), and \( \mathcal{O}_{WWW} \) can be found in Ref. 3. If \( f_B = f_W \), the number of independent WWV couplings is reduced to two. Choosing \( \kappa_\gamma \) and \( \lambda_\gamma \) as independent parameters, the WWZ couplings are then given by ("HISZ scenario"):

\[
\Delta g_1^Z = \frac{1}{2 \cos^2 \theta_W} \Delta \kappa_\gamma , \tag{3}\]
\[ \Delta \kappa_Z = \frac{1}{2} (1 - \tan^2 \theta_W) \Delta \kappa_\gamma, \quad (4) \]
\[ \lambda_Z = \lambda_\gamma, \quad (5) \]
where \( \theta_W \) is the weak mixing angle, \( \Delta \kappa_V = \kappa_V - 1 \) and \( \Delta g_i^Z = g_i^Z - 1 \). However, it should be noted that these relations are modified when operators of dimension 8 or higher are included.

Within the nonlinear realization scenario, there are two operators which contribute to anomalous couplings at lowest order. They correspond to \( \mathcal{O}_B \) and \( \mathcal{O}_W \) in Eq. (3). In both, the linear and non-linear realization scenarios, non-standard three vector boson couplings are of \( \mathcal{O}(m_W^2/\Lambda^2) \). If the energy scale of the new physics is \( \sim 1 \) TeV, anomalous gauge boson couplings are expected to be no longer than \( \mathcal{O}(10^{-2}) \). Current high precision experiments still allow anomalous couplings of \( \mathcal{O}(1) \).

Tree level unitarity uniquely restricts the WWV and ZZV couplings to their SM gauge theory values at asymptotically high energies. This implies that either any deviation of \( g_Y^i \), \( \kappa_Y \), \( \lambda_Y \) and \( h_Y^i \), \( i = 3, 4 \) has to be described by a form factor, or that the effective Lagrangian describing the anomalous vector boson self-interactions breaks down at very high energies, \( \sqrt{s} \). The functional behaviour of the form factors depends on the details of the underlying new physics. Effective Lagrangian techniques are of little help here because the low energy expansion which leads to the effective Lagrangian exactly breaks down where the form factor effects become important. Therefore, ad hoc assumptions have to be made. Here, we assume a behaviour similar to the nucleon form factor
\[ \Delta \kappa_V (\hat{s}) = \frac{\Delta \kappa_V^0}{(1 + \hat{s}/\Lambda_{FF}^2)^n}, \quad (6) \]
and similarly for the other couplings. \( \Delta_{FF} \) in Eq. (3) is the form factor scale which is a function of the scale of new physics, \( \Lambda \). We shall assume that \( n = 2 \) for WWV couplings, and \( n = 3 \) (\( n = 4 \)) for \( h_Y^i \) (\( h_Y^V \)). Since anomalous couplings are probed over a large \( \hat{s} \) range in di-boson production at hadron colliders, it is mandatory to take form factor effects into account in these processes in order to avoid unphysically large cross sections at high energies.

Some of the features of anomalous couplings, namely form factors and the necessity to consider the full S-matrix elements can nicely be illustrated by some very non-anomalous physics, namely fermion loop corrections within the SM. At the same time the problem of how to implement finite W width effects while maintaining gauge invariance when dealing with processes involving vector boson self-interactions can be addressed.

Let us consider \( W \gamma \to \ell \nu \gamma \) production at hadron colliders. Replacing the W propagator factors \( 1/(q^2 - m_W^2) \) by a naive Breit-Wigner form, \( 1/(q^2 - m_W^2 + i m_W \Gamma_W) \), where \( \Gamma_W \) is the W width, will disturb the gauge cancellations between the individual Feynman graphs and thus lead to an amplitude which is not electromagnetically gauge invariant. Finite width effects are properly included by resumming the imaginary part of the W vacuum polarization. In the unitary gauge and for \( q^2 > 0 \) the W propagator is thus given by
\[ D_W^{\mu \nu}(q) = \frac{-i}{q^2 - m_W^2 + i q^2 \gamma_W} \left( g^{\mu \nu} - \frac{q^\mu q^\nu}{m_W^2} (1 + i \gamma_W) \right), \quad (7) \]
with \( \gamma_W = \Gamma_W/m_W \). Note that the W propagator has received a \( q^2 \) dependent effective width which actually would vanish in the space-like region. A gauge invariant expression for \( q\bar{q} \to \ell \nu \gamma \) is obtained by attaching the final state photon in all possible ways to all charged particle propagators in the Feynman graphs, including the charged fermions inside the W vacuum polarization loops. As a result, for the process considered here, the lowest order \( WW \gamma \) vertex function \( \Gamma_0^\alpha \beta \mu \) is replaced by
\[ \Gamma_0^\alpha \beta \mu = \Gamma_0^\alpha \beta \mu (1 + i \gamma_W). \quad (8) \]

The modification of the lowest order \( WW \gamma \) vertex in Eq. (8) looks like the introduction of anomalous couplings \( \gamma \gamma = 1 + i \gamma_W \) and one may thus worry that the full amplitude will violate unitarity at large center of mass energies \( \sqrt{s} \). While indeed the vertex is modified, this modification is compensated by the effective \( \hat{s} \)-dependent width in the propagator. This interplay of propagator and vertex corrections illustrates the remarks made above. The leading one-loop contributions, namely the imaginary parts of \( WW \gamma \) vertex and inverse W propagator, lead to a change of the S-matrix element for W production which can be parameterized in terms of form factors by
\[ g_1^V (\hat{s}) = \kappa_V (\hat{s}) = 1 - \frac{i \Gamma_W m_W}{\hat{s} - m_W^2 + i \gamma_W \hat{s}}, \quad (9) \]
and the form factor scale is set by the W mass.

3 Tests at Hadron Colliders

The signals of anomalous gauge boson couplings in di-boson production at hadron colliders (\( q\bar{q} \to W^+ W^- \), \( W \gamma \), \( Z \gamma \), \( WZ \)) can be understood from the high energy behaviour of the anomalous contributions to the helicity amplitudes. Terms proportional to non-standard couplings increase with energy like \( (\sqrt{s}/m_W)^m \), where \( m \) depends on the coupling and the process considered, and \( \hat{s} \) is the di-boson invariant mass squared. For large values of \( \sqrt{s} \), the non-standard contributions to the helicity amplitudes would dominate, and would suffice to explain differential distributions of the photon and the W/Z decay products. Due to the fact that anomalous couplings only contribute via s-channel W, Z or
The current limits on anomalous WWV and ZγV couplings (h_V^Z = h_V^γ = 0 in the SM at tree level) from Tevatron experiments are summarized in Table 1. For the Main Injector Era, integrated luminosities of the order of 1 fb^{-1} are envisioned, and through further upgrades of the Tevatron accelerator complex, an additional factor 10 in luminosity may be gained (TeV33). The substantial increase in integrated luminosity will make it possible to test the WWV and ZγV vertices with much greater precision than in current experiments. In Fig. 2 we compare the limits on WWV couplings expected from e^+e^- → W^+W^- → ℓνjj and the various di-boson production processes in hadron collisions in the HISZ scenario [see Eqs. (3) - (5)] for the envisioned energies and integrated luminosities of the Tevatron and LEP II. Similar bounds are obtained if different relations between the anomalous couplings are assumed. The bounds expected from Tevatron and LEP II data for Δκ are quite similar (≈ 0.2), whereas the Tevatron enjoys a clear advantage in constraining λ_5 (|λ_5^0| < 0.06), if correlations between the two couplings are taken into account. It should be noted, however, that the strategies to extract information on vector boson self-interactions at the two machines are very different. At the Tevatron one exploits the strong increase of the anomalous contributions to the helicity amplitudes with energy to derive limits. At LEP II, on the other hand, information is extracted from the angular distributions of the final state fermions. Data from the Tevatron and LEP II thus yield complementary information on the nature of the WWV couplings.

Because of the much higher energies accessible at the Tevatron and the steep increase of the anomalous contributions to the helicity amplitudes with energy, Tevatron experiments will be able to place significantly better bounds on the ZγV couplings (|h_{30}^V| < 0.024, |h_{40}^V| < 0.0013 for Λ_{FF} = 1.5 TeV and \int \mathcal{L} dt = 10 fb^{-1}) than LEP II (|h_{30,40}^V| < 0.5). The Tevatron limits, however, do depend non-negligibly on the form factor scale assumed.

Since terms proportional to the non-standard WWV and ZγV couplings in the di-boson production amplitudes grow quickly with energy, one expects that experiments at the LHC are able to improve significantly the bounds which can be obtained at the Tevatron. The effect of QCD corrections, which are very large in di-boson production processes at the LHC, and the tt, ttγ and ttZ backgrounds can be reduced by imposing a jet veto. Table 2 compares the sensitivities which can be achieved in di-boson production at the LHC with 100 fb^{-1} for Λ_{FF} = 3 TeV. If the integrated luminosity is reduced by a factor 10, the bounds listed in Table 2 are weakened by about a factor 2. The limits depend non-negligibly on

![Figure 1: The differential cross section for the transverse momentum of the Z boson in pp → W^+Z at the Tevatron in the SM case (solid line) and for various anomalous WWZ couplings. A form factor scale of Λ_{FF} = 1 TeV has been assumed.](image-url)
We have discussed the theoretical aspects of factor 2, weakens the limits by a factor five to ten e-couplings in the HISZ scenario [see Eqs. (3) – (5)] from WWV. Figure 2: Comparison of the expected sensitivities on anomalous "HISZ scenario" [see Eqs. (1) – (3)] from e+e− → WW+WW− → ℓνjj at LEP II and various processes at the Tevatron.

4 Conclusions

We have discussed the theoretical aspects of WWV and ZγV couplings and their measurement at the Tevatron and LHC. These couplings are defined through a phenomenological effective Lagrangian, analogously to the general vector and axial vector couplings, gV and gA, for the coupling of gauge bosons to fermions. The major goal of these measurements will be the confirmation of the SM predictions. If the energy scale of the new physics responsible for the non-standard gauge boson couplings is ~ 1 TeV, these couplings are expected to be no larger than O(10−2).

Present data from di-boson production at the Tevatron yield bounds typically in the range of 0.5 – 2.0. Within the next 10 years, experiments conducted at the Tevatron and at LEP II are expected to confirm the SM WWV (ZγV) couplings at the 10% (1%) level. At the LHC one expects to probe anomalous WWV couplings with a precision of O(10−4 – 10−3). The limits on the ZγV couplings are very sensitive to the value of ΛFF. For ΛFF ≥ 3 TeV, the bounds which can be achieved are of O(10−3) for h3/2, and of O(10−5) for h1/2.

Table 2: Expected 95% CL limits on anomalous WWV, V = γ, Z, and ZV couplings from experiments at the LHC with 100 fb−1 and ΛFF = 3 TeV (ℓ1,2 = e, µ). Only one of the independent couplings is assumed to deviate from the SM at a time. A jet veto and standard lepton identification cuts are imposed to reduce backgrounds.

| channel | limit |
|---------|-------|
| Wγ → ℓνγ | −0.080 < Δκγ < 0.080 |
| −0.0057 < λγ < 0.0057 |
| WZ → ℓ1ν1ℓ2ν2 | −0.0060 < Δκγ < 0.0097 |
| −0.0053 < λγ < 0.0067 |
| HIZS scenario | −0.064 < Δκγ < 0.107 |
| −0.0076 < λγ < 0.0075 |
| WW → ℓ1ν1ℓ2ν2 | −0.025 < Δκγ < 0.047 |
| −0.0079 < λγ < 0.0078 |
| WWγ → ℓ1ν1ℓ2ν2 | −0.018 < Δκγ < 0.027 |
| −0.0111 < λγ < 0.0084 |
| Zγ → e+e−γ | −0.0013 < h3/2 < 0.0013 |
| −6.8 · 10−6 < h1/2 < 6.8 · 10−6 |

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