Effects of an Anomalous *ZWW* Vertex on *Z* Decays

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We consider the effects of an anomalous quadrupole moment of the W on the Z^0 branching ratios and asymmetries at the CERN e^+e^- collider LEP. Stringent experimental constraints from the ρ parameter are evaded if $\lambda_\gamma = \lambda_Z \equiv \lambda$. We therefore choose a scheme based on a custodial global $\mathrm{SU}(2)_{\mathrm{weak}}$ symmetry, implemented with the W-dominance mechanism such that $\rho = 1$ is enforced. With the expected accuracy at LEP, we find that, for the scale of new physics $\Lambda = 1$ TeV, a limit $|\lambda| \leq 0.34$ can be obtained for this coupling. This is as good a limit as can be set at LEP 200.

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To date, theoretical considerations of the effects of new physics at the CERN e^+e^- collider LEP have generally concentrated on three cases.[1] First of all, the Z boson could decay to new particles. Second, the rates for certain rare decays could be considerably enhanced through new physics. Finally, the effects of new Z bosons could be seen by running off the Z peak (in general, the presence of an extra Z boson will lead to a mass shift for the standard Z boson, which could also be seen at the Z peak). However, there is a fourth possibility, which has hardly been considered, namely, that of looking for extra radiative corrections due solely to new physics. One might expect that these "anomalous radiative corrections" would be extremely small. This is true for those extensions of the standard model which are based on gauge theories and hence renormalizable (e.g., supersymmetry, left-right symmetric models, etc.). However, if the new physics gives rise to nonrenormalizable terms (i.e., if it is an effective theory at LEP energies), these terms will introduce divergences in loop calculations proportional to $\ln(\Lambda^2/M_W^2)$ or higher, where Λ is the scale of the new physics (compositeness). In such models there is the possibility of "large" anomalous radiative corrections which could be detected through high-precision measurements of the Z^0 parameters. It is this possibility which we address in this Letter.

We consider the radiative corrections at LEP due to an anomalous *ZWW* coupling, which can only appear in models not based on a local $\mathrm{SU}(2)\times\mathrm{U}(1)$ symmetry. If one views this coupling as arising from a nongauge model of weak interactions (as would be the case if the gauge bosons were composite), one question which must be addressed is the well-confirmed universality of the weak coupling, which we usually associate with a local gauge symmetry, while another issue which should be explained relates to the fact that $\rho = 1$. These should constitute the first requirements to impose on any model.

These properties can be recovered in nongauge models incorporating both a global SU(2) symmetry and the W-dominance mechanism. This approach[2] is borrowed from the composite ρ mesons system. One assigns the W bosons to the triplet representation of a global SU(2). Electromagnetic interactions are introduced via W^0-γ mixing and W dominance. In such a model, for the W self-interactions one obtains the same terms as in the standard model for the dimension-four operators. In particular, the anomalous electromagnetic and weak magnetic moments of the W are equal to unity. One can add a C- and P-conserving dimension-six operator, which is allowed by the custodial global SU(2) symmetry and W dominance:

$$\mathcal{L}_{WWV} = \mathcal{L}_{SM} + i g_{WWV} (\lambda_V / M_W^2) W^\dagger_{\mu\nu} W^\nu{}_\lambda V^{\lambda\mu}\,, \qquad (1)$$

where V stands for the photon or the Z^0, $g_{WW\gamma} = -e$, and $g_{WWZ} = -e\cot\theta_W$.

λ_V is used to define the anomalous quadrupole moment of the W. The global SU(2) symmetry automatically gives $\lambda_Z = \lambda_\gamma \equiv \lambda$. Note that if we had taken a more general Lagrangian in which, for instance, the relation $\lambda_Z = \lambda_\gamma \equiv \lambda$ did not hold, we would have had to consider the contribution of the *ZWW* anomalous terms to the *WW*, *ZZ*, and $Z\gamma$ two-point functions. The correction to the ρ parameter due to this contribution has been calculated[3] to be proportional to $(\lambda_\gamma - \lambda_Z)\Lambda^4$, where Λ is the scale of new physics. Thus, for $\Lambda \simeq 1$ TeV, the measurement of the ρ parameter will, in general, put very tight bounds on the anomalous couplings. However, these bounds can be evaded if $\lambda_Z = \lambda_\gamma$.

Since all anomalous propagator effects cancel in this model, the only nonstandard radiative correction of interest is due to the diagram of Fig. 1. The standard-

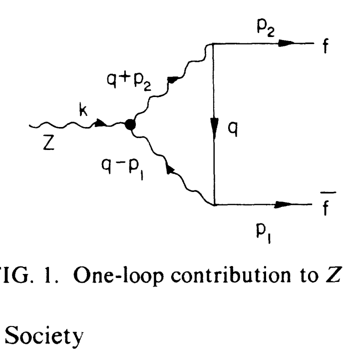

FIG. 1. One-loop contribution to $Z \to f\bar{f}$.

model ZWW vertex in this loop gives a very small contribution. We therefore need only consider the anomalous ZWW coupling of Eq. (1). This vertex will contribute, through the loop diagram, to the $Z^0 \to f\bar{f}$ rates and asymmetries. Before proceeding with the calculation we note first that, since the anomalous coupling is constructed out of field strengths, the $Zf\bar{f}$ vertex must be transverse in the Z field without imposing that f or $\bar{f}$ are on shell. Because the W's couple to left-handed fermions, the induced (CP-conserving) coupling must take the form

$$\bar{u}(p_2)[iA\sigma_{\mu\nu}k^\nu + B(k^2\gamma_\mu - k_\mu \not{k})(1-\gamma_5)]v(p_1)\,. \quad (2)$$

The A term breaks chirality and is negligible if one ignores fermion masses. For a real Z, the k_μ term disappears, so that the contribution to the vertex has the form $M_Z^2\gamma_\mu(1-\gamma_5)$. Of course, this term would vanish for a real photon. Simply from the structure of the coupling we can expect that this anomalous vertex will yield a fairly large contribution in high-energy processes (due to the k^2 term). These contributions will be larger than for the muon $(g-2)$, since for the muon only the chirality-breaking term contributes and is suppressed by the muon mass. We also note that the vertex satisfies $k_1^\alpha\Gamma_{\mu\alpha\beta} = k_2^\beta\Gamma_{\mu\alpha\beta} = 0$, where k_1 and k_2 are the momenta of the W's, so that the longitudinal part of the W propagators does not contribute in the loop, which simplifies the calculation somewhat.

Neglecting the mass of the final-state fermions, the amplitude for the process of Fig. 1 can be written

$$\mathcal{M}_{\text{loop}} = ie\cot\theta_W \frac{\lambda}{M_W^2}\frac{g^2}{4}\bar{u}(p_2)(\not{\epsilon}\gamma^\alpha\not{p}_2 - \not{p}_1\gamma^\alpha\not{\epsilon})(1-\gamma_5)v(p_1)I_\alpha\,, \quad (3)$$

where

$$I_\alpha = \int \frac{d^4q}{(2\pi)^4}\frac{q_\alpha q^2}{(q^2-m_{f'}^2)[(q-p_1)^2-M_W^2][(q+p_2)^2-M_W^2]}\,. \quad (4)$$

We have ignored finite terms since the divergent piece I_α will give the dominant contribution for this process. To evaluate the linearly divergent integral of Eq. (4), we introduce a regulator $-\Lambda^2/(q^2-\Lambda^2)$ and obtain

$$I_\alpha = \frac{i}{32\pi^2}\ln\left(\frac{\Lambda^2}{M_W^2}\right)(p_1-p_2)_\alpha + \text{finite}\,. \quad (5)$$

We should remark that this result does not depend on the regularization procedure since we keep only the leading (logarithmic) divergence, which is regulator independent; ambiguity arises only in the finite pieces.

Including the tree-level contribution, the amplitude for $Z \to f\bar{f}$ is then

$$\mathcal{M} = \frac{g}{2\cos\theta_W}\bar{u}(p_2)\gamma_\mu[c_L'(1-\gamma_5) + c_R(1+\gamma_5)]v(p_1)\,, \quad (6)$$

where c_L and c_R are the usual standard-model couplings,

$$c_L = I_{3L} - Q\sin^2\theta_W\,, \qquad c_R = -Q\sin^2\theta_W\,, \quad (7)$$

and $c_L' = c_L + \delta$, with

$$\delta = \frac{\lambda\alpha}{8\pi\sin^2\theta_W}\ln\left(\frac{\Lambda^2}{M_W^2}\right)\,. \quad (8)$$

[In deriving Eq. (8), we have used the relation $\cos\theta_W = M_W/M_Z$.] It is important to notice that δ is flavor independent and that it only contributes to the left-handed part of the Z^0 current. One can observe shifts in partial widths and asymmetries at the Z^0 pole. For instance, the shift in the partial width to fermions is

$$\Delta\Gamma_f/\Gamma_f = 2\delta c_L^f/[(c_L^f)^2 + (c_R^f)^2]\,. \quad (9)$$

In order to get a reasonable limit on λ from such measurements, one has to keep in mind that the standard-model radiative corrections also induce shifts in partial widths and asymmetries of the same order of magnitude as the anomalous vertex contribution. These corrections depend largely on the t-quark mass since they come mainly from the t-quark loop correction to the Z propagator. To look unambiguously at the effect of the anomalous vertex one has to consider quantities which are free of dependence on the top-quark mass. One such quantity is the partial width to b quarks,[4] in which there are cancellations between the Z^0 propagator correction and the standard vertex diagram with a top quark in the loop. The limit obtained on λ from this observable is rather weak, as shown in Table I. It has been shown[5] that by forming appropriate combinations of observables (twiddled quantities), it is possible to get rid of all uncertainties due to the top-quark mass, thereby singling out the effects of new physics. The list and definitions of these observables are given in Table I together with the limits on λ that can be obtained from each observable, assuming the expected accuracies at LEP as given in Ref. 5. The best limit comes from a measurement of n which is basically the ratio of the hadronic and muonic decay widths. When the scale $\Lambda = 1$ TeV, this limit is $|\lambda| < 0.34$. The other observables listed in Table I, the quantities x, y, z, and w, involve different combinations of asymmetries normalized in such a way that the standard-model prediction is zero. The shift due to the

TABLE I. Limits on $|\lambda|$ (69% C.L.) from various observables at LEP for $\Lambda = 1$ TeV. For w, $s_z = 1 - c_z = 0.5\{1 - [1 - 4(38.65\ \text{GeV})^2/M_Z^2]^{1/2}\}$.

Observable	Expected LEP accuracy	Limit on $\|\lambda\|$
$\Delta\Gamma_{b\bar{b}}/\Gamma_{b\bar{b}}$	0.05	1.7
n: $\frac{3}{59}\Gamma_{\text{had}}/\Gamma_{\mu\mu} - [9/2\alpha(M_Z)]\Gamma_{\mu\mu}/M_Z - \frac{1}{2}$	0.0125	0.34
x: $A_{FB}^{\text{pol}(b\bar{b})} - \frac{1}{15}A_{LR}^{\text{had}} - \frac{9}{13}$	0.02	10.0
y: $A_{FB}^{\text{pol}(c\bar{c})} - \frac{9}{25}A_{LR}^{\text{had}} - \frac{9}{20}$	0.02	1.1
z: $[9/\alpha(M_Z)]\Gamma_{ee}/M_Z - \frac{1}{3}A_{LR}^{\text{had}} - 1$	0.02	1.0
w: $(M_W^2/M_Z^2)/c_z - \frac{1}{2}A_{LR}^{\text{had}} - 4s_z$	0.008	0.63

anomalous vertex is not as large for these observables, all of which require polarization.

The limit we get for λ should be compared with other limits obtained from loop contributions to various low-energy processes and from direct measurements of the ZWW or γWW vertex. In Ref. 3 a comprehensive survey of the constraints on general CP-conserving anomalous couplings was done. The most restrictive limit due to indirect effects in low-energy data comes from the contribution to the W and Z mass, i.e., the ρ parameter. As mentioned earlier, this constraint does not apply in the type of model we are considering, that with a global SU(2) symmetry. In fact, in this model, the constraints from νN scattering, the polarized electron-deuterium asymmetry, and the cross section and forward-backward asymmetry in $e^+e^- \to \mu^+\mu^-$ all disappear. The only low-energy process which gives a limit is the measurement of $g-2$ of the muon,[6] since just the photon vertex contributes (λ_γ). However, the scale-independent bound is rather weak, $-2 \le \lambda_\gamma \le 6$. Finally, one can impose unitarity constraints on the anomalous vertex from $f\bar{f}' \to WW, WZ, W\gamma$ assuming that the anomalous couplings have a form factor that goes to zero suddenly at a scale Λ where new physics takes over.[7] The limit obtained is $|\lambda| < 0.6$.

The cutoff Λ used to regularize the integral introduces just a logarithmic divergence, which means that the bounds on λ we extract are only mildly dependent on Λ, the scale of new physics. Thus, measurements at LEP will essentially give limits directly on the coupling λ, which can be contrasted with those obtained from direct searches (Table II). From Table II, we see that the limits on λ obtained from searches for anomalous radiative corrections at LEP are better than those from the DESY ep collider HERA, and competitive with Fermilab Tevatron limits. It is possible to search for anomalous ZWW couplings in another way at LEP, using the process $Z \to q\bar{q}e\nu$ through two virtual W's. However, one needs an optimistic number of events (5×10^7), and even then the limits obtained are far weaker than those using radiative corrections. Surprisingly, the limit on λ obtained at LEP is as good as or better than that from LEP 200—a measurement of the total cross section for $e^+e^- \to W^+W^-$, with an accuracy of 10% (5%), yields $|\lambda| < 0.5$ ($|\lambda| < 0.3$). Of course, our limit on λ would be improved at multi-TeV machines such as the CERN Large Hadron Collider (LHC) and the Superconducting Super Collider (SSC) where the 10^{-2} level can be reached.

TABLE II. Expected limits on $|\lambda|$ from other colliders. Unless otherwise indicated, all limits are at 69% C.L. For LEP 200, the limits (a) and (b) correspond to a 10% and 5% measurement of the total cross section, respectively.

Collider	Process	Limit on $\|\lambda\|$
HERA[a]	$ep \to eWX$	<0.9
Tevatron[b]	$p\bar{p} \to W\gamma$	<0.25–0.3
LEP[c]	$Z \to q\bar{q}e\nu$	$\lesssim 5$ (5×10^7 Z's)
LEP 200[d]	$e^+e^- \to W^+W^-$	<0.5 (a)
		<0.3 (b)
LHC[e]	$q\bar{q} \to W\gamma$	<0.054 (99.99% C.L.)
SSC[e]	$q\bar{q} \to W\gamma$	<0.039 (99.99% C.L.)

[a]Reference 8.
[b]Reference 9.
[c]Reference 10.
[d]References 3 and 11.
[e]Reference 12.

To conclude, we have calculated the new radiative corrections in $Z^0 \to f\bar{f}$ due to an anomalous ZWW quadrupole coupling (λ). At LEP, a limit of $\lambda < 0.34$ can be obtained for $\Lambda = 1$ TeV, where Λ is the scale of new physics. This limit, which would be a significant improvement over the present constraints, is better than that which can be obtained at HERA, and is competitive with Tevatron limits. The limit on λ is also as good as or better than that which can be obtained in direct searches at LEP 200. Only the very-high-energy hadron machines LHC and SSC could improve on these limits.

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