EXPERIMENTAL SIGNATURES OF A PARITY VIOLATING ANOMALOUS COUPLING $g_5^Z$

G. Valencia

Department of Physics and Astronomy
Iowa State University
Ames, IA 50011

Abstract

I discuss the experimental signatures of a parity violating but CP conserving interaction in the symmetry breaking sector of the electroweak theory.
1 Introduction

The standard model of electroweak interactions has now been tested thoroughly in a number of experiments. The only sector that has not been tested directly is the electro-weak symmetry breaking (or Higgs) sector. It is very important to understand in detail the experimental signatures for the symmetry breaking sector. These vary from the direct search for new particles such as a Higgs boson, to the search for indirect manifestations of the existence of these new particles.

A convenient parameterization of these indirect effects of new particles at energies below threshold for their production is that of anomalous gauge boson couplings, the subject of this meeting. As discussed by Wudka [1], there are several ways in which these anomalous gauge boson couplings may be written in terms of a low energy effective Lagrangian.

I choose to study the case of a strongly interacting symmetry breaking sector in which there is no light Higgs boson, and therefore, use an effective Lagrangian with a non-linearly realized symmetry breaking. My motivation for this choice is simple: if there is a light Higgs boson we will find it directly and not through its contributions to anomalous couplings. I furthermore choose the “Gasser and Leutwyler”[2] construction of the effective Lagrangian because it makes the discussion of global symmetries transparent.

First I briefly review the formalism in order to establish the notation and discuss the possible size of the parity violating anomalous coupling from simple dimensional analysis. I then study the indirect bounds that can be placed on this coupling from its one-loop contribution to rare decays and partial $Z$ widths. Finally I discuss how to isolate the parity violating coupling in future high energy experiments.

2 Formalism

2.1 Effective Lagrangian

The starting point is the minimal standard model without a Higgs boson. This model can be written as the usual standard model, but replacing the scalar sector with the effective Lagrangian [3]:

$$\mathcal{L}^{(2)} = \frac{v^2}{4} \text{Tr} \left( D_{\mu} \Sigma \right) \left( D_{\mu} \Sigma \right).$$

(1)

The matrix $\Sigma \equiv \exp( iw \cdot \vec{r}/v)$, contains the would-be Goldstone bosons $w_i$ that give the $W$ and $Z$ their mass via the Higgs mechanism. Their interactions with the $SU(2)_L \times U(1)_Y$ gauge bosons follow from the covariant derivative:

$$D_{\mu} \Sigma = \partial_{\mu} \Sigma + \frac{i}{2} g W_{\mu}^i \tau^i - \frac{i}{2} g' B_{\mu} \Sigma \tau_3.$$  

(2)

The details of the physics that break electroweak symmetry determine the next-to-leading order effective Lagrangian. At energies small compared to $\Lambda$, it is sufficient
to consider those terms that are suppressed by $E^2/\Lambda^2$ with respect to Eq. [4]. There are three terms in this next to leading order effective Lagrangian that contribute to gauge boson self-energies at tree level (and thus to the LEP observables $\epsilon_{1,2,3}$ of Ref. [4]). For later reference, the one that respects the custodial symmetry is $L_{10}$.

There are several terms in the next to leading order effective Lagrangian that contribute to three gauge boson couplings at tree level. Only two of them respect the custodial symmetry, for later reference they are $L_{9L}, L_{9R}$.

Finally, there are several terms in the next to leading order effective Lagrangian that contribute to couplings with at least four gauge bosons. Two of these terms respect the custodial symmetry and for later reference they are $L_{1}, L_{2}$.

The next to leading order effective Lagrangian that respects the custodial symmetry is then:

$$L^{(4)} = \frac{v^2}{\Lambda^2} \left\{ L_1 \left[ \text{Tr} \left( D^\mu \Sigma D^\nu \Sigma \right) \right]^2 + L_2 \text{Tr} \left( D^\mu \Sigma D^\nu \Sigma \right) \text{Tr} \left( D^\mu \Sigma D^\nu \Sigma \right) 
- ig L_{9L} \text{Tr} \left( W^{\mu\nu} D^\rho \Sigma D^\sigma \Sigma \right) \right. 
- ig' L_{9R} \text{Tr} \left( B^{\mu\nu} D^\rho \Sigma D^\sigma \Sigma \right) 
+ gg' L_{10} \text{Tr} \left( \Sigma B^{\mu\nu} \Sigma W^{\mu\nu} \right) \right\}. \tag{3}$$

There are many more terms that break the custodial symmetry, but only one that violates parity while conserving $\mathcal{CP}$. This term gives rise to three and four gauge boson couplings and is the subject of this talk.

The motivation for considering this term is, of course, that we should explore all possibilities for the symmetry breaking sector. In theories where the electroweak symmetry breaking sector conserves parity, like the minimal standard model or most technicolor theories, this term is expected to be very small.

The parity violating and $\mathcal{CP}$ conserving effective Lagrangian at order $1/\Lambda^2$ is

$$L^{(4)}_{p.v.} = \frac{v^2}{\Lambda^2} g \hat{\alpha} \epsilon^{\alpha\beta\mu\nu} \text{Tr} \left( \tau_3 \Sigma D_{\mu} \Sigma \right) \text{Tr} \left( W_{\alpha\beta} D_{\nu} \Sigma \Sigma \right) \tag{4}$$

where $W_{\mu\nu}$ is the $SU(2)$ field strength tensor. In terms of $W_{\mu} \equiv W^i_{\mu} \tau_i$, it is given by:

$$W_{\mu\nu} = \frac{1}{2} \left( \partial_{\mu} W_{\nu} - \partial_{\nu} W_{\mu} + i \frac{g}{2} [W_{\mu}, W_{\nu}] \right). \tag{5}$$

It is easy to see that this is the only term that violates parity and conserves $\mathcal{CP}$ to order $1/\Lambda^2$.

In unitary gauge, the effects of the Lagrangian Eq. [4], are very simple. There is a three gauge boson interaction:

$$L^{(3)} = -\frac{\hat{\alpha} g^3 v^2}{\Lambda^2 c_0} \epsilon^{\alpha\beta\mu\nu} \left( W_{\mu}^- \partial_{\alpha} W_{\beta}^+ - W_{\beta}^+ \partial_{\alpha} W_{\mu}^- \right) Z_{\mu}. \tag{6}$$

2
which generates a $Z(g) \rightarrow W^+(p^+)W^-(p^-)$ coupling. In the notation of Ref. [3] we have the correspondence:

$$g_5^Z = \hat{\alpha} \frac{g^2 v^2}{c_\theta^2} \Lambda^2 = \frac{4M_Z^2}{\Lambda^2} \hat{\alpha}. \hspace{1cm} (7)$$

There is also a four gauge boson interaction required by electromagnetic gauge invariance:

$$L^{(4)} = \frac{2\hat{\alpha}g^4v^2s_\theta}{\Lambda^2c_\theta} \epsilon^{\alpha\beta\mu\nu} W^\alpha_{\alpha} W^\beta_{\beta} Z_{\mu} A_{\nu}. \hspace{1cm} (8)$$

This interaction contributes to the processes we discuss and must be considered simultaneously with that of Eq. 6. The Feynman rules for this interaction were written down in Ref. [6].

### 2.2 Natural size of $g_5^Z$ and Unitarity

Within the minimal standard model, the operator Eq. 4 is generated at one-loop by the splitting between top-quark and bottom-quark masses. In the limit $m_t \gg m_W$, and setting $m_b = 0$ one finds [4]:

$$\left( \frac{v^2}{\Lambda^2 \hat{\alpha}} \right)_{top} = \frac{N_c}{128\pi^2} \left( 1 - \frac{8}{3}s_\theta^2 \right) \approx 10^{-3} \hspace{1cm} (9)$$

For comparison, in this same limit one obtains $\frac{v^2}{\Lambda^2} L_1 \sim -3 \times 10^{-4}$ and $\frac{v^2}{\Lambda^2} L_2 \sim 6 \times 10^{-4}$ [4]. We see that in this limit (in which the custodial symmetry is violated “maximally”), the parity violating coupling is of the same size as other anomalous couplings. Of course, this limit is not allowed by the size of $\Delta \rho$. Taking the scale $\Lambda$ to be a few TeV ($2$ TeV for definiteness), one expects that in theories where there is no custodial symmetry and $\rho \approx 1$ accidentally, $\hat{\alpha}$ can be of order one. On the other hand, in theories with a custodial symmetry, one expects $\hat{\alpha}$ to be at most as large as $\Delta \rho$.

Our effective Lagrangian formalism breaks down at some scale $\Lambda \leq 3$ TeV, and this manifests itself in amplitudes that grow with energy and violate unitarity at some scale related to $\Lambda$. By studying the high energy behavior of longitudinal vector boson scattering one finds that the effective Lagrangian description breaks down between 1 and 2 TeV. For our numerical estimates we will work with energies up to 2 TeV. We thus turn the question around and ask how large can $\hat{\alpha}$ be so that all scattering amplitudes remain below their unitarity bound at energies up to 2 TeV. The answer is $|\hat{\alpha}| < 5$. This means that the bounds that can be placed at high energy experiments on this coupling will only be meaningful if they are better than $|\hat{\alpha}| < 5$. For bounds placed at lower energy machines such as LEP, the bad high energy behavior shows up in the need for counterterms to the one-loop calculations. Since we do not know what those counterterms are, the bounds obtained will be connected to “naturalness” assumptions.
Present Bounds

In this section we study the bounds that already exist on $g_5^Z$. They follow from considering the one-loop effects of the operator Eq. 4 in the coupling of a $Z$ boson to fermions. These observables do not single out the effects of $g_5^Z$, they are sensitive to most of the anomalous couplings.

3.1 Rare $K$- and $B$-meson decays

These rare decays receive contributions from the parity violating effective Lagrangian Eq. 4 at the one-loop level. One-loop amplitudes with one vertex from the $\mathcal{O}(1/\Lambda^2)$ effective Lagrangian are $\mathcal{O}(1/\Lambda^4)$. A complete study thus requires the next to next to leading order counterterms, as well as two loop contributions from the leading order effective Lagrangian. It is clear that there are several contributions to these decays that occur at the same order as the one-loop contribution from $g_5^Z$ and that they could cancel: we assume that they do not.

In unitary gauge, the $g_5^Z$ coupling affects this decays by modifying the “$Z$-penguin” diagram as discussed in Ref. [6, 8]. The result is dominated by top-quark intermediate states and is finite due to a GIM cancellation. With $x_t = m_t^2/m_W^2$ and defining

$$W(x_t) \equiv \frac{3}{4} x_t \left( \frac{1}{1 - x_t} + \frac{x_t \log x_t}{(1 - x_t)^2} \right) \tag{10}$$

one can write down the result in terms of the notation of Ref. [9], by replacing:

$$Y(x_t) \rightarrow \hat{Y}(x_t) = Y(x_t) + g_5^Z c_2 W(x_t)$$

$$Y(x_t) = \frac{x_t}{8} \left( \frac{x_t - 4}{x_t - 1} + \frac{3x_t}{(x_t - 1)^2} \log x_t \right) \tag{11}$$

One finds for example:

$$\Gamma(B_s \rightarrow \mu^+ \mu^-) = \frac{G_F^2}{\pi} \left( \frac{\alpha}{4\pi s_0^2} \right)^2 F_B^2 m_{\mu}^2 m_B |V_{tb}V_{ts}^*|^2 \hat{Y}(x_t)^2. \tag{12}$$

There are similar contributions to the decays $K_L \rightarrow \mu^+ \mu^-$ and $K^+ \rightarrow \pi^+ \nu \bar{\nu}$. Because $K_L \rightarrow \mu^+ \mu^-$ is dominated by long distance physics, it can only be used to place a “theoretical” bound on $g_5^Z$ by requiring the new contribution to be less than the standard model short distance contribution [8]. This results in

$$g_5^Z < \mathcal{O}(1). \tag{13}$$

If the rates for the short distance dominated processes $B_s \rightarrow \mu^+ \mu^-$ or $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ are measured to within factors of two, the same bound Eq. 13 will be obtained. To improve this bound would require a precision measurement of the rate, combined with detailed knowledge of all the standard model parameters (CKM angles, top quark mass, and decay constants) [8].
3.2 Partial Z widths at LEP

High precision measurements at the Z pole at LEP combined with polarized forward backward asymmetries at SLC put stringent limits on any new physics beyond the standard model. These measurements are now sufficiently precise to limit the one-loop contribution of anomalous three gauge boson couplings to the Z pole observables.

The bounds arise because at the one-loop level Eq. 3 modifies the $Z f \bar{f}$ couplings. Because the operator modifies the gauge boson self-couplings, its one-loop effects on the Z couplings to fermions affect both the flavor diagonal and the flavor changing vertices considered in the previous section. It turns out that the flavor diagonal vertices provide better constraints due to the extraordinary precision of the LEP measurements.

The flavor diagonal calculation is different from the flavor changing calculation in that the one-loop effects of $g_5^Z$ are now divergent. Nevertheless, from the effective field theory perspective this is not significant. In both cases, there are other contributions to the physical processes from other non-renormalizable interactions between fermions and gauge bosons. In the previous section we adopted the point of view that those other interactions did not cancel the contributions of $g_5^Z$ to rare decays. In this section we adopt the point of view that the renormalization of such couplings removes any divergence from physical amplitudes. As is usual in effective field theory calculations, we estimate the size of the $g_5^Z$ contribution to the physical amplitudes from the leading non-analytic terms that go like log($\mu$).

In addition to the direct contribution of $g_5^Z$ to the $Z f \bar{f}$ vertex we must consider indirect effects due to renormalization. In particular, the operator of Eq. 3 also modifies the $W^\pm \to \ell^\pm \nu$ coupling, contributing in this way to muon decay and thus introducing a renormalization of $G_F$. In terms of the input parameters: $G_F$ as measured in muon decay, $\alpha_s(M_W^2) \approx 1/128.8$ and the physical $Z$ mass, and using a $s^2_\theta$ defined by the relation:

$$s^2_Z c^2_Z \equiv \frac{\pi \alpha_s}{\sqrt{2} G_F M_Z^2},$$

we find:

$$\frac{\delta \Gamma_j^5}{\Gamma_j^{(0)}} = \frac{3\alpha}{2\pi} g_5^Z \log\left(\frac{\mu}{M_W}\right) \left[ \frac{2L_f}{L_f^2 + R_f^2} s^2_\theta + \left(1 + 2R_f(L_f + R_f) \frac{c^2_\theta}{L_f^2 + R_f^2} \right) \frac{c^2_\theta}{s^2_\theta - c^2_\theta} \right].$$

(15)

Where the shifts in the partial decay widths of the Z are defined by:

$$\Gamma(Z \to f \bar{f}) = \Gamma_j^{SM} + \delta \Gamma_j^5 \equiv \Gamma_j^{SM} \left(1 + \frac{\delta \Gamma_j^5}{\Gamma_j^{SM}}\right).$$

(16)

To place bounds on $g_5^Z$ (and other couplings) we compare the standard model predictions, $\Gamma_j^{SM}$, including the one loop QED and QCD radiative corrections with the most recent results from LEP. We use the theory numbers of Langacker [12].

Our 90% confidence level interval for the allowed values of $g_5^Z$ is shown in Table [10]. To place this result in perspective, it is instructive to compare them with results for the couplings that respect the custodial symmetry in Eq. 3 [13]:

5
Table 1: 90% confidence level intervals for $g^Z_5$ from different LEP observables.

| Coupling (Λ = 2 TeV) | 90% confidence level interval |
|----------------------|-----------------------------|
| $L_{10}(M_Z)_{new}$  | (-0.46,0.77)                |
| $L_{9L}$             | (-22,16)                    |
| $L_{9R}$             | (-77,94)                    |
| $L_1 + 5/2L_2$       | (-28,26)                    |
| $\hat{\alpha}$      | (-9,5)                      |
| $g^Z_5$              | (-0.07,0.04)                |

From Table 1 we see that the best limits are placed on the coupling that contributes to the $Z$ self-energy at tree level, $L_{10}$. The bounds on the other couplings are obtained by taking only one of them to be non-zero at a time, and they are all comparable. A deviation in the partial $Z$ widths from their standard model value could not be attributed to a single coupling. In order to isolate the effects of $g^Z_5$ we consider in the next two sections other observables that single out the parity violating operator.

## 4 Future Bounds

In this section we discuss the most promising reactions to place bounds on $g^Z_5$ in future colliders. These bounds arise from considering observables that single out the coupling $g^Z_5$ making through its parity violating nature.

### 4.1 Forward-backward asymmetry in $e^+_Le^-_R \rightarrow W^+W^-$

In this section we study the effect of the parity violating operator Eq. 4 on the process $e^+e^- \rightarrow W^+W^-$. This process receives contributions from $s$-channel $\gamma$ and $Z$ exchange diagrams and from a $t$-channel neutrino exchange diagram. The latter contributes only to $e^+_Le^-_R \rightarrow W^-\bar{W}^+$. The differential cross-section for right-handed electrons is found to be [3, 4]:

$$\frac{d\sigma_{TT}}{d(\cos \theta)}|_{e^-_R} = \frac{\pi \alpha^2 \beta^3}{s} \frac{m^4_Z}{(s - m^2_Z)^2} \sin^2 \theta$$

$$\frac{d\sigma_{LL}}{d(\cos \theta)}|_{e^-_R} = \frac{\pi \alpha^2 \beta^3}{32s} \frac{c^4_\theta}{(s - m^2_Z)^2} (5 + \beta^2)^2 \sin^2 \theta$$

$$\frac{d\sigma_{TL}}{d(\cos \theta)}|_{e^-_R} = \frac{\pi \alpha^2 \beta^3}{s} \frac{m^2_Z}{c^2_\theta (s - m^2_Z)^2} \left(1 + \cos^2 \theta + 2\beta \frac{s}{m^2_Z g^Z_5 \cos \theta}\right)$$  \hspace{1cm} (17)

where $\beta^2 = 1 - 4m^2_W/s$. Other anomalous couplings do not contribute to the forward backward asymmetry in $e^+_Le^-_R \rightarrow W^-\bar{W}^+$ and they are not considered here.

As can be seen from Eq. 17, there is a term in $\sigma_{TL}$ that is linear in $\cos \theta$ (the scattering angle in the center of mass). This term gives rise to a forward-backward asymmetry. Although there is a similar term in the differential cross-section for
\( e_L^− e_R^+ \rightarrow W^+ W^- \), in that case one also has a \( t \)-channel neutrino exchange diagram that gives rise to a very large forward-backward asymmetry within the minimal standard model. Thus, if we want to isolate the \( g_5^Z \) term, it is very important to have right-handed electrons. Since the cross-section for left-handed electrons is several orders of magnitude larger than that for right-handed electrons, it presents a formidable background.

We find [6] that the largest sensitivity to \( g_5^Z \) occurs in the forward-backward asymmetry at high center of mass energies. This sensitivity decreases dramatically if there is any contamination of left handed electrons as shown in Figure 1, where we present the forward-backward asymmetry for an \( e^+ e^- \) collider with \( \sqrt{s} = 500 \) GeV. This figure shows the great sensitivity of the observable to the coupling \( g_5^Z \).

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure1.png}
\caption{Forward-backward asymmetry for the process \( e^+ e^- \rightarrow W^+ W^- \) for \( \sqrt{s} = 500 \) GeV. The different curves from upper most to lowest correspond to a fraction of right handed electrons in the beam of 0\%, 90\%, 95\%, 99\% and 100\%.
}
\end{figure}

Unfortunately it also shows how this sensitivity is lost if there is even a small fraction of left handed electrons.

The total cross section is also sensitive to the value of \( g_5^Z \); however, a deviation in the total cross section from the standard model value would not single out the \( g_5^Z \) coupling.

### 4.2 High energy \( e^- \gamma \rightarrow \nu W^- Z \)

In this section we explore the possibility of observing the effects of the parity violating operator Eq. 4 via the anomalous four-gauge-boson coupling that it generates. We thus turn our attention to high energy vector-boson fusion experiments. Given the form of the four vector-boson interaction, Eq. 8, we look at processes involving one
photon and one Z. There are several possibilities, for example Zγ production in high energy e⁺e⁻ or pp colliders. This process, however, suffers from large standard model backgrounds. We consider instead a high energy e⁻γ collider where we can cleanly identify the process e⁻γ → νW⁻Z, and where we can also consider a polarized photon if need be.

To understand the physics, we first use the equivalence theorem to compute the the polarized cross sections for Wγ → wz, \( \sigma(\lambda^W, \lambda^γ) \) [6]:

\[
\begin{align*}
\sigma_{+−} &= \sigma_{−+} = \frac{\alpha^2}{s_0^2} \frac{1}{3s} \\
\sigma_{++} &= \sigma_{−−} = \frac{\alpha^2}{s_0^2} \frac{1}{3s} \left( |g_5^Z|^2 c_0^4 \frac{s^2}{m_W^4} \right) \\
\sigma_{L+} &= \sigma_{L−} = \frac{\alpha^2}{s_0^2} \frac{1}{3s} \left( |g_5^Z|^2 c_0^4 \frac{s^3}{4 m_W^6} \right)
\end{align*}
\]

(18)

From Eq. (18) we see that the \( g_5^Z \) term does not interfere with the lowest order term. This means that we can only construct observables sensitive to \( g_5^Z \) that are parity even and can thus be generated by other anomalous couplings. However, it is possible that the cross section is more sensitive to the \( |g_5^Z|^2 \) term than to those terms proportional to \( L_{9L}, L_{9R} \) or \( L_{10} \) in very high energy machines. The reason is that the \( |g_5^Z|^2 \) term is the only one that contributes to the amplitude where all three vector-bosons are longitudinally polarized (this is the source of \( \sigma_{L±} \) in Eq. (18)) and we expect these terms of “enhanced electroweak strength” to dominate at high energies. This is indeed the case, as shown by a numerical simulation [13].

To demonstrate the significant sensitivity of this process to \( \hat{\alpha} \), we consider a 2 TeV e⁺e⁻ collider (the e⁻γ differential cross section is folded with the energy spectrum of the back scattered photon). We make use of the relative enhancement of \( \hat{\alpha} \) at higher energies to isolate this coupling with a set of cuts like:

\[
|\cos \theta_V| < 0.8, \quad p_T(WZ) > 30 \text{ GeV}, \quad M(WZ) > 0.5 \text{ TeV}.
\]

(19)

Here the \( p_T(WZ) \) cut is optimized to suppress reducible backgrounds from other sources. The numerical results support our earlier conclusion. The interference of \( g_5^Z \) with the lowest order term (which vanishes in the effective W approximation) is very small, so it is not possible to single out the \( g_5^Z \) term through a parity violating observable. On the other hand, by isolating the high invariant mass region for the WZ pair, we significantly enhance the contribution of \( g_5^Z \) with respect to other couplings as shown in Figure 2. In that Figure we show a 3σ significance resulting from the anomalous couplings \( \hat{\alpha}, L_{9L}, \) and \( L_{9R} \), at \( \sqrt{s_{ee}} = 2 \text{ TeV} \) for the cuts of Eq. (19) as a function of the integrated luminosity. We see that the coefficient \( \hat{\alpha} \) can be probed here to a level less than \( (\Lambda/2 \text{ TeV})^2 \).
5 Conclusions

Here we summarize the bounds that can be placed on the coupling $g_5^Z$ from all the processes discussed in this talk. We also compare them with the natural size expected for $g_5^Z$. We see from Table 2 that the current bound at LEP is an order of magnitude better than the bounds from rare decays. An indication of how precise the LEP measurements are is the fact that the LEP bound can only be improved by one order of magnitude in a 2 TeV $e^+e^-$ collider. In principle, $g_5^Z$ can be bound precisely by studying the forward backward asymmetry in $e_L^+e_R^- \rightarrow W^+W^-$, however, it is not clear that it will ever be possible to achieve the high degree of polarization that would be required. From the numbers in Table 2 we conclude that an observation of a non-zero value for $g_5^Z$ would be very strong evidence against a custodial symmetry in the electroweak symmetry breaking sector.

| Process | Bound on $|g_5^Z|$ |
|---------|----------------------|
| Rare Decays | $\mathcal{O}(1)$ |
| Partial $Z$ widths | $5 \times 10^{-2}$ |
| $A_{fb}(e_L^+e_R^- \rightarrow W^+W^-)$ | Potentially very good, needs $P \sim 100\%$ |
| $e^-\gamma \rightarrow \nu W^-Z$ | $5 \times 10^{-3}$ (in a 2 TeV $e^+e^-$ collider) |
| Natural size | $10^{-4}$ (with custodial symmetry) |
| | $10^{-2}$ (without custodial symmetry) |

Figure 2: $3\sigma$ sensitivity of an $e^+e^-$ collider at $\sqrt{s_{ee}} = 2$ TeV (operating in the $e^-\gamma$ mode) to $\hat{\alpha}, L_{9L}$ and $L_{9R}$ with the cuts Eq. [11]. The curves are shown as a function of integrated luminosity, and we set $\Lambda = 2$ TeV.
6 Acknowledgements

This work was supported in part by a DOE OJI award under contract number DEFG0292ER40730. I am grateful to S. Dawson for many pleasant collaborations on the matters of this talk. I also thank my collaborators T. Han and K. Cheung.

References

[1] J. Wudka, these proceedings.

[2] Following the work of Gasser and Leutwyler for the case of chiral symmetry breaking in QCD, J. Gasser and H. Leutwyler, *Ann. of Phys.* **158** 142 (1984).

[3] T. Appelquist and C. Bernard, *Phys. Rev.* **D22** 200 (1980); A. Longhitano, *Nucl. Phys.* **B188** 118 (1981).

[4] G. Altarelli, R. Barbieri and F. Caravaglions, *Nucl. Phys.* **B405** 3 (1993).

[5] K. Hagiwara, *et. al., Nucl. Phys.* **B282** 253 (1987).

[6] S. Dawson and G. Valencia, *Phys. Rev.* **D49** 2188 (1994).

[7] S. Dawson and G. Valencia, *Nucl. Phys.* **B348** 23 (1991).

[8] X.-G. He, *Phys. Lett.* **B319** 327 (1993).

[9] G. Buchalla and A. Buras, *Nucl. Phys.* **B400** 225 (1993).

[10] S. Dawson and G. Valencia, *Phys. Lett.* **B333B** 207 (1994); *erratum Phys. Lett.* **341B** 452 (1995).

[11] M. Peskin, *Theory of Precision Electroweak Measurements*, Lectures presented at the 17th SLAC Summer Institute, Stanford, CA (1989).

[12] P. Langacker, UPR-0624T; and private communication.

[13] The notation and normalization is that of S. Dawson and G. Valencia, [hep-ph 9410364](http://arxiv.org/abs/hep-ph/9410364) to appear in *Nucl. Phys.* **B**.

[14] C. Ahn, *et. al., Nucl. Phys.* **B309** 221 (1988).

[15] K. Cheung, *et. al., Phys. Rev.* **D51** 5 (1995).