An alternative model for the electroweak symmetry breaking sector and its signature in future $e\gamma$ colliders.

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Abstract

We perform a preliminary study of the deviations from the Standard Model prediction for the cross section for the process $e\gamma \rightarrow \nu_e W\gamma$. We work in the context of a higgsless chiral lagrangian model that includes an extra vector resonance $V$ and an anomalous $\gamma W V$ coupling. We find that this cross section can provide interesting constraints on the free parameters of the model once it is measured in future $e\gamma$ colliders.

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1 Introduction

The Standard Model has been tested to a remarkable accuracy at LEP but there is still very little information about its electroweak symmetry breaking (EWSB) sector apart from a lower limit on the Higgs boson mass, $M_H > 63.5$ GeV \[^\text{1} \] . On the theoretical side, the existence of a Higgs boson by itself makes it difficult to understand the fine-tuning that protects its mass to be as large as the Planck mass. There are two possible ways to avoid this so-called hierarchy problem \[^\text{2} \]. One could either have a supersymmetric theory, in which quantum corrections to the Higgs mass are only logarithmic instead of quadratic, or one could have a strongly interacting underlying theory which has the EWSB sector as its low-energy effective theory. In spite of some new indirect evidence in favor of supersymmetric models \[^\text{3} \], here we will concentrate on the second approach, which has rich experimental consequences at the TeV scale that could be tested at future hadron colliders.

The presence of this strongly interacting sector would manifest itself primarily in the scattering of the longitudinal components of the weak gauge bosons $W_L, Z_L$, which at high energies behave as the pseudo Nambu-Goldstone bosons originating from the global symmetry breaking occurring in the EWSB sector, $SU(2)_L \times SU(2)_R \rightarrow SU(2)_V$. The residual $SU(2)_V$ symmetry is responsible for keeping $\rho = M_W/(M_Z \cos \theta) = 1$ at tree level. Therefore, $W_L$’s and $Z_L$’s behave as techni-pions of this underlying strong dynamics and we can describe their interactions by using chiral lagrangian techniques. If this scenario is correct, it is very plausible that resonances would also play a role in describing physical processes. In fact, amplitudes derived solely from the pion sector of a chiral lagrangian grow with energy and violate unitarity rather quickly; the presence of resonances tend to unitarize these amplitudes. At this point, the simplest possibility is that two types of resonances could appear: scalar and vector resonances. The usual Higgs sector is an example of the former choice; however, in QCD we know that vector resonances are more important in restoring unitarity, \textit{i.e.}, in saturating scattering amplitudes. Therefore, we should keep an open mind with respect to what type of resonance would show up in $W_L W_L$ scattering. Both types of resonances can be nicely described in terms of chiral lagrangians. The possibility of strong $W_L W_L$ interactions has been recently studied in detail at future hadron colliders \[^\text{4} \] and at future $e^+ e^-$ machines \[^\text{5} \].

In this work we will concentrate on a higgsless chiral lagrangian model, the so-called BESS model \[^\text{6} \], which incorporates an extra vector resonance $V$ analogous to a techni-rho . We extend this model by including new anomalous terms that give rise to $\gamma W V$ interactions. We explore the consequences of such an interaction at future $e\gamma$ colliders, obtained from a $\sqrt{s} = 1$ TeV $e^+ e^-$ machine via laser back-scattering off one
of the high energy electron (positron) beam. This $\gamma WV$ coupling is the analogue of the coupling responsible for $\rho \rightarrow \pi \gamma$ decays in QCD. In a previous work \cite{7} a similar model was used to estimate the contribution from this new resonance to $\gamma \gamma \rightarrow Z_L Z_L$ but the results were not encouraging due to a lack of a $s-$channel contribution. A similar $\gamma Z \omega_T$ coupling, where $\omega_T$ is the techni-omega has been used to study the possibility of detecting a $\omega_T$ in future hadron colliders \cite{8}.

In particular, we will study the sensitivity of the cross section for the process $e\gamma \rightarrow \nu_e W\gamma$ to a free parameter $k$ that parametrizes this new coupling. The advantages of studying such a process is that it is independent of the usual Standard Model Higgs boson and the new vector particle with the anomalous coupling would appear as a $s-$channel resonance.

This paper is organized as follows: in the next section we review the salient features of the BESS model; in section 3 we extend the model by introducing anomalous couplings terms responsible for the $\gamma WV$ coupling; in section 4 we compute the cross section for the process $e\gamma \rightarrow \nu_e W\gamma$ in the presence of the new anomalous coupling and we conclude in section 5.

2 Brief review of the BESS model

The BESS model \cite{6} is based on the so-called hidden symmetry approach \cite{9} that successfully describes the low energy QCD phenomenology involving $\rho$’s and $\pi$’s. Applied to the electroweak sector, this hidden symmetry approach describes the interactions of a new vector resonance $V$ with the particles of a higgsless Standard Model. In this section we will summarize its relevant results.

Let us begin by writing down a global $SU(2)_L \times SU(2)_R$ non-linear $\sigma$ model which describes the low-energy interactions of the ligh degrees of freedom (the Nambu-Goldstone bosons $w^a(a = 1, 2, 3)$) of a strongly interacting fundamental gauge theory:

$$\mathcal{L} = \frac{v^2}{4} Tr[\partial_\mu U \partial^\mu U^\dagger],$$

where

$$U = \exp[2i w/v] \quad , \quad w = w^a \tau^a / 2 \quad , \quad v = 246 \text{ GeV}. \quad (2)$$

The above lagrangian is invariant under the transformation $U \rightarrow LUR^\dagger$, where $L, R \in SU(2)$.

The simplest BESS model is obtained by introducing a vector resonance as a gauge boson of an additional local $SU(2)_V$ symmetry. Furthermore, we will gauge the $SU(2)_L \times SU(2)_R$ symmetry to introduce the usual gauge fields. To this end we
re-parametrize $U$ as:

$$U(x) = \xi_L^i(x) \xi_R(x),$$

(3)

where under a local $SU(2)_L \times SU(2)_R \times SU(2)_V$ transformation we have:

$$\xi_L \to h \xi_L L^\dagger, \quad \xi_R \to h \xi_R R^\dagger, \quad h \in SU(2)_V.$$  

(4)

The fields $\xi$ can be represented in the following form:

$$\xi_L = \exp[i\sigma/v] \exp[-iw/v], \quad \xi_R = \exp[i\sigma/v] \exp[iw/v],$$

(5)

where $\sigma = \sigma^a r^a/2$ is sometimes called a compensator field.

Under a parity transformation in the strongly interacting sector one has:

$$\sigma \leftrightarrow \sigma, \quad w \leftrightarrow -w, \quad \xi_L \leftrightarrow \xi_R, \quad U \leftrightarrow U^\dagger, \quad \vec{x} \leftrightarrow -\vec{x}.$$  

(6)

The gauge fields $\tilde{V}_\mu, l_\mu$ and $r_\mu$ transform as usual:

$$\tilde{V}_\mu \to h \tilde{V}_\mu h^\dagger + \frac{i}{g_V} (\partial_\mu h) h^\dagger, \quad l_\mu \to L l_\mu L^\dagger + \frac{i}{g_L} (\partial_\mu L) L^\dagger, \quad r_\mu \to R r_\mu R^\dagger + \frac{i}{g_R} (\partial_\mu R) R^\dagger.$$  

(7)

The covariant derivatives are given by:

$$D_\mu \xi_L = \partial_\mu \xi_L - ig_V \tilde{V}_\mu \xi_L + ig_L \xi_L l_\mu, \quad D_\mu \xi_R = \partial_\mu \xi_R - ig_V \tilde{V}_\mu \xi_R + ig_R \xi_R r_\mu,$$

(8)

which transform as:

$$D_\mu \xi_L \to h D_\mu \xi_L L^\dagger, \quad D_\mu \xi_R \to h D_\mu \xi_R R^\dagger.$$  

(9)

We finally define the “building blocks” $\alpha_{L\mu}$ and $\alpha_{R\mu}$ to construct a chiral invariant lagrangian by:

$$\alpha_{L\mu} = D_\mu \xi_L \xi_L^\dagger, \quad \alpha_{R\mu} = D_\mu \xi_R \xi_R^\dagger,$$

(10)

which transform under local $SU(2)_L \times SU(2)_R \times SU(2)_V$ as:

$$\alpha_{L\mu} \to h \alpha_{L\mu} h^\dagger, \quad \alpha_{R\mu} \to h \alpha_{R\mu} h^\dagger.$$  

(11)

Under parity we have $\alpha_{L0} \leftrightarrow -\alpha_{R0}$ and $\alpha_{Li} \leftrightarrow \alpha_{Ri}.$

Finally, the most general lowest order lagrangian which respects the parity-like symmetry $L \leftrightarrow R$ is given by:

$$\mathcal{L} = -\frac{v^2}{4} Tr[(\alpha_{L\mu} - \alpha_{R\mu})^2] - a\frac{v^2}{4} Tr[(\alpha_{L\mu} + \alpha_{R\mu})^2] + \mathcal{L}_{\text{kin}}$$

(12)
where
\[ \mathcal{L}_{\text{kin}} = -\frac{1}{2} Tr[(V_{\mu \nu})^2] - \frac{1}{2} Tr[(L_{\mu \nu})^2] - \frac{1}{2} Tr[(R_{\mu \nu})^2] \] (13)
are the kinetic terms involving the usual non-abelian field-strength for the vector resonance and gauge fields.

In unitary gauge \((w = \sigma = 0, \xi_L = \xi_R = 1)\) and specializing the gauge fields to the Standard Model fields the \(\alpha_{\mu}'s\) reduce to :
\[ \alpha_{L\mu} = -igV_{\tilde{V}_\mu} + igW_\mu, \quad \alpha_{R\mu} = -igV_{\tilde{V}_\mu} + ig'B_\mu \tau^3/2, \] (14)
and the usual BESS model is given by the lagrangian :
\[ \mathcal{L}_{\text{BESS}} = -\frac{v^2}{4} \left\{ Tr[(ig\tilde{W}_\mu - ig'B_\mu \tau^3/2)^2] + aTr[(ig\tilde{W}_\mu + ig'B_\mu \tau^3/2 - 2igV_{\tilde{V}_\mu})^2] \right\} + \mathcal{L}_{\text{kin}} \] (15)

Notice that here we prefer to keep a canonical kinetic term for the \(\tilde{V}_\mu\) field and for this reason our coupling constant \(g_v\) differs from the coupling constant \(g''\) used in references [6] : \(g_v = g''/2\). Since it is not relevant in what follows, we assume that there is no direct coupling between the new vector resonance \(\tilde{V}_\mu\) and fermions. This is equivalent to set the parameter \(b\) defined in reference [6] to zero.

The fields appearing in Eq. [15] are unphysical and one obtains the physical fields \(A_\mu, Z_\mu, W_\mu\) and \(V_\mu\) by diagonalizing the mass matrix in the neutral and charged sector. In the charged sector one has :
\[ \tilde{W}^\pm = W^\pm \cos \phi + V^\pm \sin \phi \] (16)
\[ \tilde{V}^\pm = -W^\pm \sin \phi + V^\pm \cos \phi. \] (17)

In the neutral sector the situation is slightly more complicated :
\[ \tilde{W}_3 = A_1 \cos \theta + A_2 \sin \theta \] (18)
\[ B = -A_1 \sin \theta + A_2 \cos \theta \] (19)
\[ \tilde{V}_3 = V \cos \xi \cos \psi - Z \sin \xi \cos \psi + A \sin \psi \] (20)

where
\[ A_1 = Z \cos \xi + V \sin \xi \] (21)
\[ A_2 = A \cos \psi + Z \sin \psi \sin \xi - V \sin \psi \cos \xi \] (22)

The mixing angles introduced above are given by :
\[ \tan \theta = g'/g, \quad \tan \psi = \frac{g' \cos \theta}{g_v} = \frac{g \sin \theta}{g_v} \] (23)
and in the limit \( g_V \gg g, g' \):

\[
\phi = -\frac{g}{2g_V}, \quad \xi = -\frac{g^2 - g'^2}{g_V\sqrt{g^2 + g'^2}}.
\]

(24)

The physical photon field is massless (\( M_A = 0 \)), and in the limit \( g_V \gg g, g' \) the masses of the other vector bosons are given by:

\[
M_W^2 = \frac{v^2}{4g^2} \left( 1 - \left( \frac{g}{2g_V} \right)^2 \right)
\]

(25)

\[
M_Z^2 = \frac{v^2}{4}(g^2 + g'^2) \left( 1 - \frac{(g^2 - g'^2)^2}{(g^2 + g'^2)4g_V^2} \right)
\]

(26)

\[
M_V^2 = ag^2_v v^2
\]

(27)

The Standard Model couplings are also modified due to mixing. For instance, the electric charge in this case is given by:

\[
e = g \sin \theta \cos \psi.
\]

There are two free parameters in this lagrangian, namely, \( g_V \) and \( a \). In the limit \( g_V \to \infty \), one recover the mass relations and couplings of the Standard Model. Precision measurements from LEP put constraints on the parameters of the model [10].

The width of the vector resonance (for \( M_V \gg M_W \)) is given by:

\[
\Gamma_V = \frac{aM_V^3}{192\pi v^2}
\]

(28)

and we can substitute \((M_V, \Gamma_V)\) for \((g_V, a)\) as free parameters.

In this work we are interested in a possible \( \gamma W V \) vertex and in principle one might think that such a vertex could occur in the kinetic terms after substituting for the physical fields. However, an explicit computation shows that no such interaction appears. In fact, one can show that if such an interaction resulted from the kinetic terms then electromagnetic gauge invariance would be violated in the process \( \gamma W \to \gamma W \). Therefore, we are led to expand the BESS model by introducing extra terms that induce new “anomalous” interactions. In the next section we discuss the motivation and describe this anomalous BESS model.

### 3 The anomalous BESS model

The parity-like operation \( L \leftrightarrow R \) is not a symmetry of the underlying theory. It corresponds to a symmetry of the theory under \( w(\vec{x}, t) \to -w(\vec{x}, t) \), which forbids transitions between even and odd numbers of the pseudo-Goldstone bosons \( w \). However, parity conservation implies in the symmetry \( w(\vec{x}, t) \to -w(-\vec{x}, t) \) and it is possible to
write down parity-conserving terms in the lagrangian that violate the $L \leftrightarrow R$ symmetry \[11\]. In QCD, these terms describe processes like $\rho, \omega \rightarrow \pi \gamma$.

In order to write the possible terms down we go back to the general case and define the following modification of the field strength for the external gauge fields:

$$\hat{L}_{\mu\nu} = \xi_L L_{\mu\nu} \xi_L^\dagger, \quad \hat{R}_{\mu\nu} = \xi_R R_{\mu\nu} \xi_R^\dagger$$

so that under $SU(2)_L \times SU(2)_R \times SU(2)_V$ we have $V_{\mu\nu}, \hat{L}_{\mu\nu}$ and $\hat{R}_{\mu\nu}$ transforming in the same manner as the “building blocks” $\alpha_{\mu}$’s. Under parity we have:

$$V_{00} \leftrightarrow V_{00}, \quad V_{0i} \leftrightarrow -V_{0i}, \quad V_{ij} \leftrightarrow V_{ij}$$

$$\hat{L}_{00} \leftrightarrow \hat{R}_{00}, \quad \hat{L}_{0i} \leftrightarrow -\hat{R}_{0i}, \quad \hat{L}_{ij} \leftrightarrow \hat{R}_{ij}.$$ 

(30)

The extra terms invariant under $SU(2)_L \times SU(2)_R \times SU(2)_V \times$ parity are given by \[11\]:

$$\mathcal{L}_{\text{anom}} = \sum_{i=1}^{4} \kappa_i \mathcal{L}_i$$

(31)

where

$$\mathcal{L}_1 = \epsilon^{\mu\nu\alpha\beta} \text{Tr} [\alpha_L \mu \alpha_L \nu \alpha_L \alpha_R \beta - \alpha_R \mu \alpha_R \nu \alpha_R \alpha_L \beta]$$

(32)

$$\mathcal{L}_2 = \epsilon^{\mu\nu\alpha\beta} \text{Tr} [\alpha_L \mu \alpha_R \nu \alpha_L \alpha_R \beta]$$

(33)

$$\mathcal{L}_3 = i \epsilon^{\mu\nu\alpha\beta} \text{Tr} [V_{\mu\nu}(\alpha_L \alpha_R \beta - \alpha_R \alpha_L \beta)]$$

(34)

$$\mathcal{L}_4 = i \epsilon^{\mu\nu\alpha\beta} \text{Tr} [\hat{L}_{\mu\nu} \alpha_L \alpha_R \beta - \hat{R}_{\mu\nu} \alpha_R \alpha_L \beta]$$

(35)

These terms are formally of order $O(p^6)$ compared to the $O(p^4)$ terms of Eq. [12] but higher order terms respecting the $L \leftrightarrow R$ symmetry would not generate the interactions we are interested in.

In unitary gauge, $\mathcal{L}_1$ and $\mathcal{L}_2$ generate only 4–boson couplings which are of higher order in the coupling constant. However, no new $\gamma\gamma W^+ W^-$ coupling is present because of the antisymmetric tensor. The relevant terms in $\mathcal{L}_3$ are naively supressed by a factor of $g_g V$ with respect to $\mathcal{L}_4$ and therefore here we will concentrate on $\mathcal{L}_4$.

The relevant terms in $\mathcal{L}_4$ are given by, in unitary gauge:

$$\mathcal{L}_4 = 2i \epsilon^{\mu\nu\alpha\beta} \text{Tr} [\partial_\mu \bar{W}_\nu (-ig_V \bar{V}_\alpha + ig \bar{W}_\alpha) (-ig_V \bar{V}_\beta + ig' B_\beta \tau_3/2) - \partial_\nu B_\mu \tau_3/2 (-ig_V \bar{V}_\alpha + ig' B_\alpha \tau_3/2) (-ig_V \bar{V}_\beta + ig \bar{W}_\beta)] + 4$–gauge couplings$$(36)

It is now straightforward but tedious to use the relations Eqs.[16–24] in order to derive the $\gamma WV$ interaction lagrangian:

$$\mathcal{L}_{\gamma WV} = -\frac{i}{6} g (\frac{g}{g_V})^2 (\sin \theta - \cos \theta) \cos \psi \epsilon^{\mu\nu\alpha\beta} \partial_\mu A_\nu (V_\alpha^+ W^-_\beta - V^-_\alpha W^+_\beta)$$

(37)
Notice that the overall coupling is very small, being of the order of $O(g^3/g^2_\gamma)$. However, this term is to be multiplied by a free parameter $\kappa_4$ and we introduce an effective coupling $k$ defined by:

$$L_{\gamma W} = i \ k \ e^{\mu\alpha\beta} \partial_\mu A_\nu \ (V^+_{\alpha} W^-_{\beta} - V^-_{\alpha} W^+_{\beta}).$$

The same type of coupling was derived in Ref. [8] for the $\omega_T Z \gamma$ interaction, where $\omega_T$ is the technicolor analogue of the $\omega$-meson. We would like to point out that there is no theoretical reason to expect large values $|k| = O(1)$. In fact, $|k| \simeq 10^{-2}$ in QCD but in this paper we adopted the approach of deriving what type of limits on $|k|$ could be achieved in future $e\gamma$ colliders.

It is our purpose to study how one can constrain the possible values of $k$ by computing deviations from the Standard Model predictions. In this paper we’ll focus on the particular process $e\gamma \to \nu_e W\gamma$, attainable at future $e\gamma$ colliders and in the next section we compute the cross section for this process in the anomalous BESS model.

4 A possible signature for the anomalous BESS model in $e\gamma$ colliders

We are interested in the process $e\gamma \to \nu_e W\gamma$, depicted in Fig. 1. The exact Standard Model cross section for this process including realistic cuts was computed in Ref. [12]. We’ll assume that the initial electron beam has a monochromatic energy of 500 GeV and that the initial photon has an energy spectrum originated from laser back-scattering off a 500 GeV electron (positron) beam. In this preliminary study we’ll use the effective-$W$ approximation [13] to estimate the cross section and we won’t try to implement any realistic acceptance cuts. One can straightforwardly improve upon our results by computing an exact cross section with appropriate cuts for a more detailed study. Since we work in unitary gauge we don’t need to invoke the equivalence theorem to derive our results.

The cross section can be written as:

$$\frac{d\sigma(s)}{dM_{\gamma W}} = \frac{2M_{\gamma W}}{s} \sum_{i=L,T} \left( \frac{d\mathcal{L}}{d\tau} \right)_{\gamma W_i} \hat{\sigma}_i(\tau s),$$

where $M_{\gamma W}$ is the $\gamma - W$ invariant mass, $\hat{\sigma}_i$ is the subprocess $\gamma W \to \gamma W$ cross section for a given initial $W$ helicity $i$, $\tau = M^2_{\gamma W}/s$ and the so-called luminosity function $d\mathcal{L}/d\tau$ is given by:

$$\left( \frac{d\mathcal{L}}{d\tau} \right)_{\gamma W_i} = \int_{\tau}^{x_m} \frac{dx}{x} f_{e\to\gamma}(x) f_{e\to W_i}(\tau/x).$$
The functions $f_{e\rightarrow \gamma}(x)$ and $f_{e\rightarrow W_i}(x)$ are respectively the energy spectrum of the laser back-scattered photon and the probability of finding a $W$ of polarization $i$ carrying an energy fraction $x$ of the parent electron.

The energy spectrum of a back-scattered photon as a function of the energy fraction $x$ of the initial electron’s energy is given by [14] :

$$f_{e\rightarrow \gamma}(x) = \frac{N(x, \xi)}{D(\xi)} \tag{41}$$

where $\xi$ is defined in terms of the electron mass $m_e$, the electron initial energy $E$ and the energy of the photon in the laser beam $\omega_0$ :

$$\xi = \frac{4E\omega_0}{m_e^2} \tag{42}$$

and

$$N(x, \xi) = 1 - x + \frac{1}{1 - x} - \frac{4x}{\xi(1 - x)} + \frac{4x^2}{\xi^2(1 - x)^2} \tag{43}$$

$$D(\xi) = \int_0^{x_m} dx \ N(x, \xi)$$

$$= \left[ 1 - \frac{4}{\xi} - \frac{8}{\xi^2} \right] \ln(1 + \xi) + \left( \frac{1}{2} + \frac{8}{\xi} - \frac{1}{2(1 + \xi)^2} \right) \tag{44}$$

where $x_m = \frac{\xi}{1 + \xi}$ is the maximum energy fraction carried by the scattered photon. If $x_m \geq 0.828$ other processes start to compete with the photon back-scattering reaction degrading the photon beam. In our case, with a laser energy $\omega_0 = 1.17$ eV we have $x_m = 0.90$. Therefore, here we’ll take $x_m = 0.828$.

For $f_{e\rightarrow W_i}(x)$ we use [13] :

$$f_{e\rightarrow W_i}(x) = \begin{cases} \frac{\epsilon}{8\pi \sin^2 \vartheta} \frac{(x^2 + 2(1-x))}{x} \ln \left( \frac{4E^2}{M_W^2} \right) & i = T \\ \frac{\alpha}{4\pi \sin^2 \vartheta} \frac{1-x}{x} & i = L \end{cases} \tag{45}$$

We now compute the point cross section that includes the diagrams shown in Fig. 2, where the notation is defined. The amplitude is given by :

$$\mathcal{M} = \mathcal{M}_s + \mathcal{M}_t + \mathcal{M}_c \tag{46}$$

where the subscripts $s, t$ and $c$ stand for the $s$–channel, $t$–chanel and contact term contributions respectively.

We find :

$$\mathcal{M}_s = \varepsilon_{1\mu} \varepsilon_{2\nu} \varepsilon_{3\alpha}^{*} \varepsilon_{4\beta}^{*} \left( \frac{i\epsilon^2}{(p_1 + p_2)^2 - M_W^2} \right)$$
\[
\left[(p_1 - p_2)^\lambda g^{\mu\nu} + 2p_2^\mu g^{\nu\lambda} - 2p_1^\nu g^{\mu\lambda}\right] g_{\lambda\sigma} \left[(p_3 - p_4)^\sigma g^{\alpha\beta} + 2p_4^\alpha g^{\beta\sigma} - 2p_3^\beta g^{\alpha\sigma}\right]
\]

\[
\frac{-ie^2 M_W^2}{(p_1 + p_2)^2 - M_W^2} g^{\mu\nu} g^{\alpha\beta}
\]

\[
\left(\frac{ik}{(p_1 + p_2)^2 - M_V^2 + iM_V\Gamma_V}\right) \epsilon^{\lambda\mu\sigma\nu} \epsilon^{\nu\alpha\sigma'\beta} \left[p_{1\lambda} p_{3\lambda} \left[g_{\sigma\sigma'} - \frac{Q_{\sigma}Q_{\sigma'}}{M_V^2}\right]\right]
\]

\[
\mathcal{M}_t = \mathcal{M}_s (p_2 \leftrightarrow -p_4, \varepsilon_2 \leftrightarrow \varepsilon_4)
\]

\[
\mathcal{M}_c = -ie^2 \varepsilon_1 \varepsilon_2^* \varepsilon_3 \varepsilon_4^* \left[2g^{\mu\alpha} g^{\nu\beta} - g^{\mu\nu} g^{\sigma\alpha} - g^{\mu\sigma} g^{\nu\alpha}\right]
\]

In order to compute the subprocess cross section \(\hat{\sigma}_i(s)\) we square the amplitude Eq\[46\], summing over the helicities of all particles but the initial W and integrate it over all values of the scattering angle.

Before showing our results, let’s recall that our free parameters are \((M_V, \Gamma_V, k)\). In the limit \(k \to 0\) one recovers the Standard Model prediction. Here we’ll concentrate on two sets of masses and widths: \((M_V = 500 \text{ GeV}, \Gamma_V = 50 \text{ GeV})\) and \((M_V = 800 \text{ GeV}, \Gamma_V = 100 \text{ GeV})\), which are chosen such that there is no unitarity violation in \(W_L W_L\) scattering [15]. For each set of \((M_V, \Gamma_V)\) we vary the values of \(|k|\) from 0 to 0.3. Our results for the subprocess cross section are displayed in Figs. 3 and 4. Notice that the new resonance couples much more strongly to the longitudinally polarized \(W\)’s, as expected from their Nambu-Goldstone boson origin.

In Fig. 5 we present our results for the differential cross section Eq. [39]. These plots follow simply from convoluting the subprocess cross section with a rapidly falling luminosity function. In Table 1 we show the the cross section in femtobarns integrated around the resonance peak \((M_V - \Gamma_V < M_W < M_V + \Gamma_V)\). For an \(e^+e^-\) collider with a total yearly luminosity of 100 fb\(^{-1}\) [16] one expects to be able to put firm constraints on this particular model.

We have also calculated the cross sections for \((M_V = 1000 \text{ GeV}, \Gamma_V = 200 \text{ GeV})\) but due to the rapidly falling luminosity function we find that these values are out of reach for a 1 TeV \(e^+e^-\) collider.

### 5 Conclusion

We have performed a preliminary study of the deviations from Standard Model predictions of the cross section for the process \(e^+e^- \to \nu_e \gamma W\) in the context of a higgsless chiral lagrangian model that includes a new vector resonance \(V\) and an anomalous \(\gamma W V\) interaction. The model has as free parameters the mass \(M_V\), the width \(\Gamma_V\) and the anomalous coupling \(k\). No exhaustive search for the discovery region in the \((M_V, \Gamma_V, k)\) parameter space was intended and no realistic cuts and experimental efficiencies were
Our results suggest that if such a particle $V$ exists with a mass in the range $500-800$ GeV with an anomalous coupling strength $|k| \gtrsim 0.2$ one would expect an enhancement of $\mathcal{O}(50-100\%)$ over the Standard Model result for the cross section $e\gamma \rightarrow \nu e\gamma W$ integrated around the new particle mass. Even with a yearly luminosity of $10\text{ fb}^{-1}$ our results imply a number of events of the order of $\mathcal{O}(400-2000)$ per year, which should not be difficult to find in the reasonable clean environment of an $e\gamma$ collider. The situation improves for larger values of $|k|$ and a more realistic analyses may be worthwhile pursuing, where more definite constraints on the parameter space of the model can be set once such a cross section is measured. For $|k| \lesssim 0.1$, this process loses its sensitive to the presence of the new resonance and one would not be able to distinguish between the model studied here and the Standard Model result.

An interesting question is the consequence of this model for future hadron colliders, in particular in the mode $pp \rightarrow \gamma(Z)WX$. Work along these lines is now in progress.

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Figure Captions

Figure 1 : The process $e\gamma \rightarrow \nu_e W\gamma$.

Figure 2 : Feynman diagrams for the subprocess $\gamma W \rightarrow \gamma W$ cross section.

Figure 3a : Point cross section in fentobarns for the process $W_L\gamma \rightarrow W\gamma$ as a function of the $\gamma W$ center-of-mass energy with $M_V = 500$ GeV and $\Gamma_V = 50$ GeV. Solid line: $|k| = 0$ (Standard Model); dot-dashed line:$|k| = 0.1$ ; dotted line: $|k| = 0.2$; dashed line: $|k| = 0.3$.

Figure 3b : Same as Figure 3a but for the process $W_T\gamma \rightarrow W\gamma$.

Figure 4a : Same as Figure 3a but for $M_V = 800$ GeV and $\Gamma_V = 100$ GeV.

Figure 4b : Same as Figure 3b but for $M_V = 800$ GeV and $\Gamma_V = 100$ GeV.

Figure 5a : Differential cross sections in fb/GeV for the process $e\gamma \rightarrow \nu_e W\gamma$ as a function of the final state $W\gamma$ invariant mass with $M_V = 500$ GeV and $\Gamma_V = 50$ GeV. Solid line: $|k| = 0$ (Standard Model); dot-dashed line:$|k| = 0.1$ ; dotted line: $|k| = 0.2$; dashed line: $|k| = 0.3$.

Figure 5b : Same as Figure 5a but with $M_V = 800$ GeV and $\Gamma_V = 100$ GeV.
| $|k|$ | (500 GeV, 50 GeV) | (800 GeV, 100 GeV) |
|-----|------------------|-------------------|
| 0   | 128.1            | 28.29             |
| 0.1 | 132.8            | 29.12             |
| 0.2 | 202.4            | 41.19             |
| 0.3 | 502.8            | 93.32             |
| 0.4 | 1,311            | 233.4             |

Table 1: Integrated cross section around $M_V - \Gamma_V < M_{\gamma W} < M_V + \Gamma_V$ in fentobarns for different values of $|k|$ and $(M_V, \Gamma_V)$.
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