Indirect Effects of New Resonances at Future Linear Colliders

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ABSTRACT

In this paper we consider a general $SU(2)_L \otimes SU(2)_R$ invariant Lagrangian describing scalar, vector and axial-vector resonances. By expanding the $WW$ and the $ZZ$ scattering amplitude up to the fourth order in the external momenta we can compare the parameters of our Lagrangian with the ones used in the effective chiral Lagrangian formalism. In the last approach there has been a recent study of the fusion processes at future $e^+e^-$ colliders at energies above 1 $TeV$. We use these results to put bounds on the parameter space of our model and to show that for the case of vector resonances the bounds obtained from the annihilation channel in fermion pairs are by far more restrictive, already at energies of the order of 500 $GeV$. 
1 Introduction

Future electron positron and proton proton colliders offer the possibility of testing the nature of spontaneous breaking of the electroweak symmetry, the most important open question not yet clarified by the high energy experiments done at LEP and Tevatron.

In particular if no light (supersymmetric) Higgs boson is found, the $WW$ physics, accessible at such accelerators, can explain which kind of mechanism is responsible for the dynamical breaking of the electroweak symmetry. Different models of $WW$ scattering have been suggested, mainly based on the chiral symmetry of the scalar sector of the Standard Model (SM) and on the effective Lagrangian approach, both with and without resonances. In fact if the $WW$ channel becomes strongly interacting new resonances are expected, and measuring the $WW$ amplitudes one can test the properties of these new resonances.

A model describing new vector, axial vector resonances, interacting with the standard gauge vector bosons was proposed in [1][2] (BESS model). In this paper this model is further generalized to include also a scalar resonance.

The production of the vector resonances at LHC was studied in a detailed way in a convenient region of the parameter space where the production cross section is sizeable [3][4]. Also we have studied the signature of this model at linear electron positron colliders with center of mass energy below the threshold for the production of these new resonances. In this low energy region we have studied indirect effects in the annihilation channel both into fermions and $WW$ pairs and in the fusion channel of the $WW$ and $ZZ$ rescattering [5].

In this paper the $WW$ and $ZZ$ scattering amplitudes are obtained by a general $SU(2)_L \times SU(2)_R$ invariant Lagrangian, which is derived by an effective Lagrangian describing scalar, vector, axial-vector resonances, in the low energy limit up to fourth order in the momentum expansion. This allows to find relations among the parameters of our model and the ones commonly used in the effective chiral Lagrangian formalism [6].

Concerning the limitations coming from future linear colliders, our analysis shows that the processes $e^+e^- \rightarrow W^+W^-\bar{\nu}\nu$ and $e^+e^- \rightarrow ZZ\bar{\nu}\nu$ are not the most important for testing strongly interacting $WW$ system in presence of new vector resonances, because the bounds on the relevant chiral parameters from the annihilation channel are more stringent.
2 The model

A model describing vector and axial-vector resonances interacting in a consistent way with the three Goldstone bosons necessary to give mass to the standard gauge vector bosons, \(W\) and \(Z\), was considered in ref. \[2\]. The idea is very simple and consists in requiring that also the new resonances acquire mass through the Higgs mechanism. In order to do that we need to introduce at least nine Goldstone boson fields. The other requirement is that after spontaneous symmetry breaking the vector resonances are triplets under the unbroken \(SU(2)\). The simplest way of realizing this situation is to assume beyond the standard global symmetry \(SU(2)_L \otimes SU(2)_R\), a local symmetry \(SU(2)\) for each new vector resonance. This symmetry group \(G \otimes H\) with

\[
G = [SU(2)_L \otimes SU(2)_R]_{\text{global}}, \quad H = [SU(2)_L \otimes SU(2)_R]_{\text{local}}, \quad (2.1)
\]

spontaneously breaks down to \(SU(2)\).

The physics below the scale of the possible strong interaction producing these resonances can be studied in terms of an effective Lagrangian. The methods to construct such a Lagrangian are the standard ones used to build up non-linear realizations (see ref. \[7\]). We will describe the Goldstone bosons by three independent \(SU(2)\) elements: \(L\), \(R\) and \(M\), whose transformation properties with respect to \(G \otimes H\) are the following

\[
L'(x) = g_L L(x) h_L(x), \quad R'(x) = g_R R(x) h_R(x),
\]

\[
M'(x) = h_R(x) M(x) h_L(x) \quad (2.2)
\]

where \(g_{L,R} \in G\) and \(h_{L,R} \in H\). Besides the invariance under \(G \otimes H\), we will also require an invariance under the following discrete left-right transformation, denoted by \(P\)

\[
L \leftrightarrow R, \quad M \leftrightarrow M^\dagger \quad (2.3)
\]

which ensures that the low-energy theory is parity conserving.

The vector and axial-vector resonances are introduced as linear combinations of the gauge particles associated to the local group \(H\). The most general \(G \otimes H \otimes P\) invariant Lagrangian is given by \[2\]

\[
\mathcal{L}_R = \mathcal{L}_G + \mathcal{L}_{\text{kin}} \quad (2.4)
\]
where
\[ \mathcal{L}_G = -\frac{v^2}{4} f(L_\mu, R_\mu) \] (2.5)

with
\[ f(L_\mu, R_\mu) = aI_1 + bI_2 + cI_3 + dI_4 \] (2.6)

\[ I_1 = tr[(V_0 - V_1 - V_2)^2], \quad I_2 = tr[(V_0 + V_2)^2] \]
\[ I_3 = tr[(V_0 - V_2)^2], \quad I_4 = tr[V_1^2] \] (2.7)

and
\[ V_0^\mu = L^\dagger D^\mu L \]
\[ V_1^\mu = M^\dagger D^\mu M \]
\[ V_2^\mu = M^\dagger(R^\dagger D^\mu R)M \] (2.8)

The parameters \( a, b, c, d \) are not independent and we will fix later a condition among them, in such a way that \( v^2 = 1/(\sqrt{2}G_F) \). The covariant derivatives are defined by
\[ D_\mu L = \partial_\mu L - LL_\mu \]
\[ D_\mu R = \partial_\mu R - RR_\mu \]
\[ D_\mu M = \partial_\mu M - ML_\mu + R_\mu M \] (2.9)

where \( L_\mu \) and \( R_\mu \) are the gauge fields associated to the local symmetry group \( H \). The quantities \( V_i^\mu \) \((i = 0, 1, 2)\) are, by construction, invariant under the global symmetry \( G \) and covariant under the gauge group \( H \)
\[ (V_i^\mu)' = h_L^\dagger V_i^\mu h_L \] (2.10)

Their transformation properties under the parity operation, \( P \), are:
\[ (V_0 \pm V_2) \rightarrow \pm M(V_0 \pm V_2)M^\dagger, \quad V_1 \rightarrow -MV_1M^\dagger \] (2.11)

Out of the \( V_i^\mu \) one can build six independent quadratic invariants, which reduce to the four \( I_i \) listed above, when parity is enforced. The kinetic part will be written in the form
\[ \mathcal{L}_{kin} = \frac{1}{g'^2}tr[F_{\mu\nu}(L)]^2 + \frac{1}{g''^2}tr[F_{\mu\nu}(R)]^2 \] (2.12)
where $g''$ is the gauge coupling constant for the gauge fields $L_\mu$ and $R_\mu$,

$$F_{\mu\nu}(L) = \partial_\mu L_\nu - \partial_\nu L_\mu + [L_\mu, L_\nu]$$  \hspace{1cm} (2.13)

and analogously for $R_\mu$.

We will also consider a scalar field invariant under the group $G \otimes H \otimes P$. We will not be interested in the self-couplings of this field $S$, and we will consider only interaction terms with the previous fields at most linear in $S$. The relevant effective Lagrangian is then

$$\mathcal{L} = \frac{1}{2} \partial_\mu S \partial^\mu S - \frac{1}{2} m^2 S^2 - \frac{v\kappa}{2} S f(L_\mu, R_\mu) + \mathcal{L}_R + \cdots$$  \hspace{1cm} (2.14)

We can observe that in the case of a Higgs particle one has $\kappa = 1$.

### 3 Low-energy limit

We will discuss now the low-energy limit of the previous Lagrangian, by keeping terms up to the fourth order in the derivatives. Let us start with the scalar field part. At the lowest order we neglect the kinetic term. Then the equation of motion gives

$$S = -\frac{v\kappa}{2m^2} f(L_\mu, R_\mu)$$  \hspace{1cm} (3.1)

We see that $S$ is at least of the second order in the derivatives, therefore the kinetic term does not contribute at the lowest order, and substituting in eq. (2.14) we get

$$\mathcal{L} = \frac{v^2 \kappa^2}{8m^2} [f(L_\mu, R_\mu)]^2 + \mathcal{L}_R + \cdots$$  \hspace{1cm} (3.2)

To discuss the same limit for the vector and axial-vector resonances is convenient to choose the gauge $R(x) = M(x) = 1$. This can be reached by the gauge transformation $h_R(x) = R^{-1}(x)$, $h_L(x) = M^{-1}(x)R^{-1}(x)$. By defining

$$\omega_\mu = L^I \partial_\mu L$$  \hspace{1cm} (3.3)

we get

$$V_{0\mu} = \omega_\mu - L_\mu, \quad V_{1\mu} = -L_\mu + R_\mu, \quad V_{2\mu} = -R_\mu$$  \hspace{1cm} (3.4)
If we neglect again the kinetic term for $L_\mu$ and $R_\mu$, we can solve easily the equations of motion for the fields at the leading order in the derivatives. It is convenient to put

$$z = \frac{c}{c + d} \quad (3.5)$$

Then

$$R_\mu = \frac{1}{2}(1 - z)\omega + R_\mu^{(3)}, \quad L_\mu = \frac{1}{2}(1 + z)\omega + L_\mu^{(3)} \quad (3.6)$$

where $R_\mu^{(3)}$ and $L_\mu^{(3)}$ are of the third order in the derivatives. By substituting into $f(L_\mu, R_\mu)$ defined in (2.6) we find

$$f(L_\mu, R_\mu) = (a + \frac{cd}{c + d})\text{Tr}(\omega^2) + \cdots \quad (3.7)$$

where the dots denote terms which are at least of the sixth order in the derivatives. In order to recover the non-linear $\sigma$-model describing the standard breaking $SU(2)_L \otimes SU(2)_R$ to $SU(2)$ we require a relation among the parameters of the Lagrangian (2.4), namely

$$a + \frac{cd}{c + d} = 1 \quad (3.8)$$

By using the following property, which holds for $2 \times 2$ matrices of the form $\vec{A} \cdot \vec{\sigma}$, with $\vec{\sigma}$ the Pauli matrices,

$$\text{Tr}(ABCD) = \text{Tr}(AB)\text{Tr}(CD) - \text{Tr}(AC)\text{Tr}(BD) + \text{Tr}(AD)\text{Tr}(BC) \quad (3.9)$$

we get

$$(\text{Tr}(\omega^2))^2 = \text{Tr}(\omega_\mu\omega^\mu\omega_\nu\omega^\nu) \quad (3.10)$$

and therefore

$$f^2(L_\mu, R_\mu) = \text{Tr}(\omega_\mu\omega^\mu\omega_\nu\omega^\nu) \quad (3.11)$$

From the observation that $\omega_\mu$ is a 1-form with zero curvature, that is satisfying

$$\partial_\mu\omega_\nu - \partial_\nu\omega_\mu + [\omega_\mu, \omega_\nu] = 0 \quad (3.12)$$

and using eq. (3.6) and eq. (2.13), we get

$$F_{\mu\nu}(L_\mu) = F_{\mu\nu}(R_\mu) = -\frac{1}{4}(1 - z^2)[\omega_\mu, \omega_\nu] + \cdots \quad (3.13)$$
Finally we find

\[
\mathcal{L} = -\frac{v^2}{4}Tr(\omega^2) + \frac{1}{4g''^2}(1 - z^2)^2Tr[\omega_\mu\omega_\nu\omega_\mu\omega_\nu]
+ \left(\frac{v^2\kappa^2}{8m^2} - \frac{1}{4g''^2}(1 - z^2)^2\right)Tr[\omega_\mu\omega_\nu\omega_\mu\omega_\nu] + \cdots
\]  

(3.14)

This expression can be compared with the general expression for the ungauged chiral Lagrangian [6] at the same order in the derivatives (we will do our calculations by using the equivalence theorem [8]). We see that at the fourth order, only the terms \(L_4\) and \(L_5\) are different from zero (we are using the standard notations of ref. [6]), therefore we get the following expressions for the chiral parameters \(\alpha_4\) and \(\alpha_5\)

\[
\alpha_4 = \frac{(1 - z^2)^2}{4g''^2}, \quad \alpha_4 + \alpha_5 = \frac{\kappa v^2}{8m^2}
\]  

(3.15)

Let us notice that this result holds for the model given by the Lagrangian (2.14) where we have explicitly considered only terms with at most two derivatives. Higher order terms in the derivative expansion add contributions and the relations (3.15) change.

4 Pion-Pion scattering amplitude

In this Section we will study the scattering amplitude for the \(W^+_LW^-_L \rightarrow W^+_LW^-_L\) and \(W^+_LW^-_L \rightarrow Z_LZ_L\), in the case of the Lagrangian considered before, i.e. for the case of a model with vector, axial vector and scalar particles. We use the unitary gauge for the \(L_\mu\) and \(R_\mu\) fields, characterized by the following expressions for fields \(L, R, M\):

\[
L = \exp(\frac{i\pi}{v})\exp(\frac{i}{v}(\lambda - z\pi))
\]  

(4.1)

\[
R = \exp(-\frac{i\pi}{v})\exp(\frac{i}{v}(\rho + z\pi))
\]  

(4.2)

\[
M = \exp(-\frac{i}{v}(\rho + z\pi))\exp(\frac{i}{v}(\lambda - z\pi))
\]  

(4.3)
where
\[ \pi = \vec{\pi} \cdot \frac{\vec{\sigma}}{2}, \quad \lambda = \vec{\lambda} \cdot \frac{\vec{\sigma}}{2}, \quad \rho = \vec{\rho} \cdot \frac{\vec{\sigma}}{2} \] (4.4)

It is not difficult to show that (apart from normalization constants) \( \vec{\lambda} \) and \( \vec{\rho} \) are the Goldstone bosons associated to \( L_\mu \) and \( R_\mu \) respectively, whereas \( \vec{\pi} \) are the Goldstone bosons necessary to give mass to \( W \) and \( Z \). The unitary gauge is then defined by taking \( \vec{\lambda} = \vec{\rho} = 0 \) in the previous expressions, that is
\[ L = \exp(\frac{i}{v}(1 - z)\pi) = L^i, \quad M = \exp(-\frac{2i}{v}z\pi) \] (4.5)

Using this gauge it is easy to evaluate the mass eigenstates which are \( V_\mu = (R_\mu + L_\mu)/2 \) and \( A_\mu = (R_\mu - L_\mu)/2 \) with masses given respectively by
\[ M^2_{V_i} = \frac{b v^2}{4} g''^2; \quad M^2_A = \frac{(c + d) v^2}{4} g''^2 \] (4.6)

In this gauge, from the \( \pi \pi \) scattering amplitudes, by using the equivalence theorem, we get
\[ M(W^+_L W^-_L \rightarrow Z_L Z_L) = A(s, t, u) \] (4.7)
\[ M(W^+_L W^-_L \rightarrow W^+_L W^-_L) = A(s, t, u) + A(t, s, u) \] (4.8)

where
\[ A(s, t, u) = \left(1 - \frac{3}{4} \beta\right) \frac{s}{v^2} - \frac{1}{4} \frac{M^2_V}{v^2} \beta \left(\frac{s - u}{t - M^2_V} + \frac{s - t}{u - M^2_V}\right) - \frac{\kappa^2}{v^2} \frac{s^2}{s - m^2} \] (4.9)

and
\[ \beta = 4 \frac{M^2_V}{g''^2 v^2} (1 - z^2)^2 \] (4.10)

We can now expand the amplitude \([10]\) up to the fourth power in the momenta getting the following expression
\[ A(s, t, u) = \frac{s}{v^2} + \frac{(1 - z^2)^2}{g''^2 v^4} \left(-s^2 + 2st + 2t^2\right) + \frac{\kappa^2}{m^2 v^2} s^2 \] (4.11)

Comparing with the expression obtained by Kilian \([9]\) from the chiral Lagrangian \([6]\)
\[ A(s, t, u) = \frac{s}{v^2} + 4\alpha_4 \frac{t^2 + u^2}{v^4} + 8\alpha_5 \frac{s^2}{v^4} \]
\[ = \frac{s}{v^2} + 4\alpha_4 \frac{-s^2 + 2st + 2t^2}{v^4} + 8(\alpha_4 + \alpha_5) \frac{s^2}{v^4} \] (4.12)
we get again the relation (3.13) correlating the chiral parameters $\alpha_4$ and $\alpha_5$ with the ones of the model presented here.

Before considering some particular cases, let us discuss the limitations on $\alpha_4$ and $\alpha_5$ coming from the partial wave unitarity conditions on the amplitudes (4.7), (4.8). The most restrictive bounds come from $s$-wave. By requiring $|a_0| \leq 1$, for $\sqrt{s} = 1.6$ TeV, we get the allowed region delimited from the solid lines showed in Fig. 1.

It is interesting to consider some particular case. We start considering the vector case (BESS model [1]), because here we can compare with the analysis done at the $e^+e^-$ colliders. We recall that the restrictions on the parameter space coming from the annihilation process $e^+e^- \rightarrow W^+W^-$, have been already considered in ref. [5]. Therefore we will be able to compare (at least for vector resonances) the sensitivity of the two processes. For the BESS case there are no axial-vector and scalar particles, that is we have to
put \( z = \kappa = 0 \). We get

\[
\alpha_4 = \frac{1}{4g''^2}, \quad \alpha_4 + \alpha_5 = 0 \quad (4.13)
\]

This means that the BESS model lies on the line \( \alpha_5 = -\alpha_4 \) with \( \alpha_4 \geq 0 \) and from Fig. 1 we can see that the unitarity bound is approximately \( \alpha_4 \leq 0.01 \) which corresponds to \( g'' \geq 5 \) (this is for \( \sqrt{s} = 1.6 \text{ TeV} \)).

From the Kilian \[9\] analysis made for a collider with polarized beams, \( \sqrt{s} = 1600 \text{ GeV} \) and an integrated luminosity of 200 fb\(^{-1}\) we find (for \( \chi^2 = 3 \), that is a confidence level of 91.7\% for a gaussian distribution)

\[
\frac{1}{4g''^2} \leq 0.0028 \quad (4.14)
\]

that is

\[
g'' \geq 9.45 \quad (4.15)
\]

or

\[
\frac{g}{g''} \leq 0.07 \quad (4.16)
\]

Following the analysis in ref. \[9\], for a collider with unpolarized beams, \( \sqrt{s} = 500 \text{ GeV} \) and an integrated luminosity of 20 fb\(^{-1}\) we find (again for \( \chi^2 = 3 \))

\[
\frac{g}{g''} \leq 0.038 \quad (4.17)
\]

We see that the indirect bounds we get on the parameter space from the annihilation process are by far more restrictive than the ones for the fusion. The situation is, of course, much different in the scalar channel, given the fact that this does not contribute to the annihilation process in a collider \( e^+e^- \).

The case of vector and axial-vector resonances is not too different, but in this case, from the analysis of ref. \[9\] one gets bounds only for the combination \( g''/(1 - z^2) \). The bound will be the same as given in (4.13), that is

\[
\frac{g''}{1 - z^2} \geq 9.45 \quad (4.18)
\]

An interesting observation is that the low-energy data give informations on a different combination of \( g'' \) and \( z \). In fact the low-energy deviations from the
SM, in the actual case, can be entirely expressed in terms of the observable $\epsilon_3$, and this is given by

$$\epsilon_3 = \left(\frac{g}{g'}\right)^2 \left(1 - z^2\right)$$

(4.19)

Notice that in the case $z = 1$, the so called degenerate BESS model (vector and axial-vector resonances degenerate in mass and in their couplings to the corresponding currents) [4], we loose any bounds; in fact in the scattering amplitude only the term corresponding to the low-energy theorems [12] survives. This follows from the fact that in the degenerate BESS the vector and axial-vector resonances decouple from the Goldstone bosons.

Finally for the case in which only the scalar particle is present (or one has also vector and axial-vector resonances degenerate), one gets from [4] a bound on the combination $\frac{\kappa^2 v^2}{m^2}$ ($\kappa > 0$)

$$\frac{\kappa^2 v^2}{8 m^2} \leq +0.0012$$

(4.20)

suggesting that for a coupling of order 1, one can test up to scalar masses of order of 2.55 TeV.

5 Conclusions

The BESS model provides an effective parametrization for strongly interacting $WW$ system in presence of new resonances. Below the resonance threshold, the relevant scattering amplitudes within the BESS model can be compared with the results from effective chiral Lagrangians. Our analysis shows that the fusion processes at future $e^+e^-$ linear colliders are not the most important for testing strongly interacting $WW$ system in presence of new vector resonances. In fact the bounds on the BESS model parameters from the annihilation channel, which we have studied in several previous papers, are most stringent and can be already obtained from colliders with a center of mass energy of 500 GeV.

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