VECTOR RESONANCES FROM A STRONG ELECTROWEAK SECTOR
AT LINEAR COLLIDERS *)

R. Casalbuoni \textsuperscript{a,b)}, P. Chiappetta \textsuperscript{c)}, A. Deandrea \textsuperscript{d)}, S. De Curtis \textsuperscript{b)},
D. Dominici \textsuperscript{a,b)} and R. Gatto \textsuperscript{d)}

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a) Dipartimento di Fisica, Univ. di Firenze, I-50125 Firenze, Italy.
b) I.N.F.N., Sezione di Firenze, I-50125 Firenze, Italy.
c) CPT, CNRS, Luminy Case 907, F-13288, Marseille, France.
d) Dépt. de Phys. Théor., Univ. de Genève, CH-1211 Genève 4.
ABSTRACT

We explore the usefulness of very energetic linear $e^+e^-$ colliders in the $TeV$ range in studying an alternative scheme of electroweak symmetry breaking based on a strong interacting sector. The calculations are performed within the BESS model which contains new vector resonances. If the mass $M_V$ of the new boson multiplet lies not far from the maximum machine energy, or if it is lower, such a resonant contribution would be quite manifest. A result of our analysis is that also virtual effects are important. It appears that annihilation into a fermion pair in such machines, at the considered luminosities, would improve only marginally on existing limits if polarized beams are available and left-right asymmetries are measured. On the other hand, the process of $W$-pair production by $e^+e^-$ annihilation would allow for sensitive tests of the hypothesized strong sector, especially if the $W$ polarizations are reconstructed from their decay distributions, and the more so the higher the energy of the machine. An $e^+e^-$ collider with c.m. energy $\sqrt{s} = 500 \, GeV$ could improve the limits on the model for the range $500 < M_V(\, GeV) < 1000$ when $W$ polarization is not reconstructed. If $W$ polarizations are reconstructed, then the bounds improve for the entire expected range of $M_V$. These bounds become more stringent for larger energy of the collider. We have also studied the detectability of the new resonances through the fusion subprocesses, but this channel does not seem to be interesting even for a collider with a c.m. energy $\sqrt{s} = 2 \, TeV$. 
1. INTRODUCTION

Several laboratories are at present engaged in projects of $e^+e^-$ linear colliders with sufficient energy and luminosity to test for new physics. Among such activity we mention the work being done at SLAC, at KEK, at Novosibirsk and Serpukhov, at DESY and Darmstadt, at CERN with the CLIC project, at various laboratories participating to the TESLA concept, with several other groups and individuals contributing to such developments. Much interest has of course been devoted to the physics potentialities of high energy $e^+e^-$ colliders. Studies have been concentrated on a collider having c.m. energy up to 500 GeV, but at the same time possibilities of c.m. energies of 1 or 2 TeV have also been discussed. A recent workshop at Munich, Annecy, Hamburg was especially centered on $e^+e^-$ at 500 GeV c.m. energy [1], and further developments were discussed at a Conference in Finland [2]. In this paper we shall describe in some more detail the calculations we had presented in Hamburg [3] and Saariselkä [4] and extend our previous work to $e^+e^-$ colliders in the TeV energy range.

The sensitivity of future $e^+e^-$ and $pp$ colliders to the exploration of the electroweak symmetry breaking mechanism has been largely discussed, both in the context of the standard model or of its supersymmetric extension, and for alternative schemes where the Higgs field is composite and an underlying strong theory is responsible for the electroweak breaking. In this last scenario the longitudinal components of $W$ and $Z$ become strongly interacting in the TeV region. Possible discovery of this strong interaction is among the physics prospects of the future $pp$ colliders, LHC and SSC. Experiments at $e^+e^-$ linear colliders in the TeV region seem to be complementary, because the structure of the events is relatively simpler than at hadron colliders, and therefore offer the possibilities for more detailed studies.

We have considered the sensitivity of $e^+e^-$ linear colliders, for different options of total center of mass energies and luminosities, to a model which corresponds to a breaking of the electroweak symmetry due to a strongly interacting sector: breaking electroweak symmetry strongly (BESS) [5]. In BESS the electroweak symmetry breaking is obtained via a non-linear realization and no Higgs particles are present. New gauge bosons $V$ appear and they could be produced as real resonances if their mass is below the collider energy. Because of beamstrahlung and synchrotron radiation, in a high energy collider, one expects in this case to see dominant peaks below the maximum c.m. energy even without having to tune the beam energies. If instead the masses of the $V$ bosons are higher than the maximum c.m. energy, they would give rise to indirect effects in the $e^+e^- \rightarrow f^+f^-$ and $e^+e^- \rightarrow W^+W^-$ cross sections, which we discuss quantitatively below.

The BESS model appears as the simplest definite scheme against which to test for a possible strong electroweak sector. The vector resonances of the BESS model are bound states of a strongly interacting sector. In this sense they are similar to ordinary $\rho$ vector mesons, or to the techni-$\rho$ particle of technicolor theories. An effective lagrangian describing in an unified way mass terms and interactions of the standard electroweak
gauge bosons and these new vector resonances $V$ was derived in ref. [5] as a gauged non-linear $\sigma$-model. The BESS model contains as additional parameters the mass $M_V$ of the new bosons (which are a degenerate triplet of an additional $SU(2)$ gauge group), their gauge coupling $g''$ (which is assumed to be much larger than $g$, and $g'$), and a parameter $b$ specifying the direct coupling of $V$ to the fermions. The standard model (SM) is recovered in the limit $g'' \rightarrow \infty$, and $b = 0$. Mixings of the ordinary gauge bosons to the $V$'s are of the order $O(g/g'')$. Due to these mixings, $V$ bosons are coupled to fermions even for $b = 0$. Furthermore these couplings are still present in the $M_V \rightarrow \infty$ limit, and therefore the new gauge boson effects do not decouple in the large mass limit.

In this paper we will be mainly interested in the study of the effects due to vector resonances of the strongly interacting electroweak symmetry breaking sector. We will extend our previous analysis which was dedicated to a machine in the range of 500 GeV [3,4]. Our description of the vector resonances is rather general and after a convenient specification of the parameters is also apt to describe a standard techni-$\rho$ ($\rho_T$) state. The production of $\rho_T$ at $e^+e^-$ machines was analyzed by M.E. Peskin at Saariselkä [6] (see also references therein). Similar studies were also done by Iddir et al. [7]. They implement the strongly interacting regime by inserting in the Born amplitude for $e^+e^- \rightarrow W^+W^-$ strong final state corrections from $WW \rightarrow WW$. The corrections can be described by the Omnés function which is approximated through a rescaling of the Gounaris-Sakurai model [8]. These authors note that the most significant sign of a vector resonance in $W^+W^-$ would manifest itself in the backward hemisphere, since the transverse channel $W_TW_T$ dominates forward due to $\nu$-exchange in $t$-channel.

In his summary talk at Saariselkä J.L. Barklow [9] concludes that a $\rho_T$ state of $M_{\rho_T} = 1.7$ TeV would be visible at machines with c.m. energies greater or of the order of 1 TeV already for a luminosity of 50 $fb^{-1}$. K. Hikasa [10] has proposed to look for possible strong $WW$ scattering effects by using the interference between $W_LW_L$ and $W_TW_T$ in the $WW$ pair production process. These would manifest themselves in correlations between the azimuthal angles of the decay fermions from $W^+$ and $W^-$, specifically in the dependence of the distribution from the sinus of the sum of the two azimuthal angles for $W^+ \rightarrow f^+f^-$ decay (both decaying leptonically, or one leptonically, the other into $c\bar{s}$).

The BESS model at LEPII and at a collider energy of $\sqrt{s} = 600$ GeV has been studied in ref. [11].

Using the recent LEP1 data combined with UA2/CDF data and low energy electroweak data, we have first derived the present bounds on the BESS parameter space. We have studied how these limits can be improved for the following three options for the future machines: $\sqrt{s} = 500, 1000, 2000$ GeV, with integrated luminosity of 20, 80, 20 $fb^{-1}$ respectively (for the parameters of various designs see ref. [12]).

We have analyzed cross-sections and asymmetries for the channel $e^+e^- \rightarrow f^+f^-$ and $e^+e^- \rightarrow W^+W^-$. For the purposes of our calculation we have also assumed that it will be possible to separate $e^+e^- \rightarrow W_L^+W_L^-$, $e^+e^- \rightarrow W_T^+W_T^-$, and $e^+e^- \rightarrow W^+_T W^-_T$. 

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The distribution of the \( W \) decay angle in its c.m. frame depends indeed in very distinct ways from its helicity, being peaked forward (backward) with respect to the production direction for positive (negative) helicity or at 90\(^\circ\) for zero helicity. One may hope that by looking at the different \( W \) decay modes it will be possible to extract such a very useful information.

At very energetic linear colliders the pair production of \( W^+W^- \) through \( W^+W^- \), \( \gamma\gamma \), \( \gamma Z \), \( ZZ \) fusion have to be taken into account. Since the final states for the annihilation process, for the fusion subprocesses through \( W^+W^- \), or \( \gamma, Z \) are different, we treat them as three independent reactions and examine their sensitivities to the BESS model parameters.

One clear advantage of \( e^+e^- \) collisions with respect to \( pp \) is that all decay modes from a \( WW \) or \( ZZ \) pair may contribute to the observable signatures, and not only those which include at least a leptonic mode.

One has to stress that an important advantage of a linear collider is the possibility one can envisage of a large longitudinal beam polarization.

We find that the differential cross-sections to \( f \bar{f} \) would allow to improve the already existing limits on BESS from LEP1 and UA2/CDF if longitudinally polarized beams will be available, particularly by considering left-right asymmetries in hadron production.

A considerable improvement of the limits on BESS is obtained by the study of the differential \( WW \) cross section. As a matter of fact the cross section acquires in BESS a term which rises linearly with \( s \) plus a constant term which is the most relevant at the energies considered here, whereas in the SM the cross-section goes down as \( 1/s \).

This improvement becomes even more effective if the \( W \) polarization is reconstructed, because the new gauge bosons are mainly coupled to longitudinal \( W \)'s.

We will show that, if no deviations from SM expectations are seen at colliders in the \( TeV \) range, the parameter space left to the BESS model is very small. Direct search for \( V \) bosons at hadronic colliders has been extensively studied. In this case the discovery of the neutral resonance \( V^0 \) looks difficult. It has however been shown that LHC and SSC are ideal to pin down a charged resonance up to 2 \( TeV \) through the \( WZ \) channel. We see that very energetic \( e^+e^- \) linear colliders are complementary to hadron colliders.

\( W \)-fusion reactions can directly inform on \( WW \) scattering. They are expected at a very energetic \( e^+e^- \) collider. In principle, such a process could be interesting because it would allow to study a wide range of mass spectrum for the \( V \) resonance also if the energy of the machine is not adjustable downward. Note however that in such linear colliders beamstrahlung will already automatically provide for a spectrum of lower initial energies. Besides, as usual in fusion reactions, the c.m. energy available for the two colliding \( W \)'s is strongly lowered with respect to the original \( e^+e^- \) energy. This makes the exploration for possible effects of a high mass strong electroweak sector via \( W \)-fusion
less encouraging, also in view of the backgrounds from processes leading to the same final states and utilizing the full beam energy. In fact we find that also at a 2 TeV $e^+e^-$ energy the number of events is too low for detectability.

Already in 1983 Ginzburg et al. [13] had suggested the possibility of obtaining an energetic photon beam by collision of the electron bunches with a laser beam operating in the visible spectrum. This technique should allow keeping a high luminosity for the photon beam from the back-scattered laser and should allow for a photonic spectrum mostly concentrated at the highest energies, not much smaller than the electron energy. Such a behaviour is quite different from that of the expected beamstrahlung photons concentrated at the lower energies, and, even more so, from that of the bremsstrahlung photons. The possibility would allow for energetic photon-photon collisions and electron-photon collisions. For the purpose of the present paper, where we are interested in a possible strong electroweak sector, $\gamma\gamma$ collisions would appear of interest if resonant behaviours are present in states of zero angular momentum which can couple to two real photons. These may be the pseudogoldstone states as in general expected in these models [14] when one considers the terms due to the Adler-Bell-Jackiw anomaly. Also interesting will be the production of pairs or more of such pseudos from $\gamma\gamma$. As far as the vector-type resonances, such as in the simplest BESS considered here, they will not be prominent from an initial state of two real photons. Their contribution will be through rescattering between W’s in $\gamma\gamma$ and in $\gamma e$ collisions. In $\gamma e$ collisions, contributions will also appear through the modification of the leptonic vertices which are as prescribed in BESS. In both cases the additional terms coming from the BESS model are negligible.

The paper is organized as follows. We will recall in section 2 some basics of the BESS model and the present limits on parameter space. In section 3 the process $e^+e^- \to \bar{f}f$ is examined. The process $e^+e^- \to W_{L,T}^+W_{L,T}^-$ is discussed in section 4. Section 5 is devoted to the fusion subprocesses $e^+e^- \to W_{L,T}^+W_{L,T}^-e^+e^-$ and $e^+e^- \to W_{L,T}^+W_{L,T}^-\bar{\nu}\nu$. We conclude in section 6.

2. THE BESS MODEL

The increasing accumulation of precise measurements of the electroweak parameters gives the possibility of finding some hints for going beyond the standard model. Through mixing effects the contribution of vector resonances from a possible strong sector would affect masses and couplings of ordinary gauge bosons. Therefore precise measurements of the width of $Z$, its mass and forward backward asymmetries, performed at LEP, allow for restrictions on the unknown parameters of the BESS model which are: the mass $M_V$, the direct coupling to fermions $b$, and the gauge coupling constant $g''$ of the $V$ bosons.

BESS is a non-renormalizable theory; therefore, when radiative corrections are considered, a cut-off $\Lambda$ has to be introduced. This cut-off plays the role of the parameter $m_H$, Higgs mass, of the SM. The strategy we follow in our analysis is to consider for
BESS the same one-loop radiative corrections as in the SM and to interpret \( m_H \) as the cut-off.

Using the following LEP1 data averaged on the four LEP experiments, as communicated at the Dallas Conference [15]:

\[
M_Z = 91.187 \pm 0.007 \text{ GeV} \\
\Gamma_Z = 2492 \pm 7 \text{ MeV} \\
\Gamma_h = 1737.1 \pm 6.7 \text{ MeV} \\
\Gamma_\ell = 83.33 \pm 0.30 \text{ MeV} \\
A^p_{FB} = 0.0154 \pm 0.0027 \\
A^p_{pol} = 0.140 \pm 0.018 \\
A^b_{FB} = 0.098 \pm 0.012
\]

and

\[
\frac{M_W}{M_Z} = 0.8807 \pm 0.0031
\]

from CDF/UA2 [16], we can derive bounds on the BESS model, that we express as 90\% C.L. contours in the plane \((b, g/g''')\) for given \( M_V \) (see Fig. 1).

The comparison with BESS includes one-loop electroweak radiative corrections calculated with a cut-off of 1 TeV and for \( \alpha_s = 0.12 \). The bounds in Fig. 1 are almost independent of the mass of the new resonances \( V \) and of the choice of the cut-off. The region gets shifted by changing the values of \( \alpha_s, m_{top} \), becoming narrower for increasing \( m_{top} \). So the region of negative \( b \) values is practically excluded for ”light” top masses.

By including the low energy data coming from cesium parity violation experiments the bounds do not significantly change. In the context of BESS, LEP 200 is expected to increase only marginally the sensitivity over LEP1. The relevant modification will be brought by the more accurate determination of \( M_W \).

In order to compare with a standard techni-\( \rho \), one has to take \( b = 0 \), choose

\[
\frac{g}{g'''} = \sqrt{2} \frac{M_W}{M_{\rho T}}
\]

and identify \( M_{\rho T} \) with \( M_V \) [17].

3. \textbf{SENSITIVITY FROM } \( e^+e^- \rightarrow f^+f^- \)

In the discussion of high energy linear colliders, we start with the fermion channels. Our analysis is based on the following observables:

\[
\sigma^\mu, \quad R = \sigma^h/\sigma^\mu \\
A^e^+e^-\rightarrow\mu^+\mu^-, \quad A^e^+e^-\rightarrow\bar{b}b \\
A^e^+e^-\rightarrow\mu^+\mu^-, \quad A^e^+e^-\rightarrow\bar{b}b, \quad A^e^+e^-\rightarrow had
\]
where $A_{FB}$ and $A_{LR}$ are the forward-backward and left-right asymmetries, and $\sigma^{h(\mu)}$ is the total hadronic ($\mu^+\mu^-$) cross section. Assuming that one can measure the final $W$ polarization, we can add to our observables the longitudinal and transverse polarized $W$ differential cross sections and asymmetries.

To evaluate these observables within the BESS model, we recall that in the neutral sector the couplings of the fermions to the gauge bosons $Z$ and $V$ are $[5,18]$

$$e \left( v_Z^f + \gamma_5 a_Z^f \right) \gamma_\mu Z^\mu + e \left( v_V^f + \gamma_5 a_V^f \right) \gamma_\mu V^\mu$$

(3.2)

where $v_{Z,V}^f$ and $a_{Z,V}^f$ are the vector and axial-vector couplings given by

$$v_Z^f = \frac{1}{\sin 2\theta} \left( AT_3^L + 2BQ_{e.m.} \right)$$

$$a_Z^f = \frac{1}{\sin 2\theta} AT_3^L$$

$$v_V^f = \frac{1}{\sin 2\theta} \left( CT_3^L + 2DQ_{e.m.} \right)$$

$$a_V^f = \frac{1}{\sin 2\theta} CT_3^L$$

(3.3)

where

$$A = \frac{\cos \xi}{\cos \psi} (1 + b)^{-1} \left[ 1 + b \sin^2 \theta \left( 1 - \frac{\tan \xi \tan \theta \sin \psi}{\tan \xi \tan \theta \sin \psi} \right) \right]$$

$$B = -\sin^2 \theta \frac{\cos \xi}{\cos \psi} \left( 1 - \frac{\tan \xi \sin \psi}{\tan \theta} \right)$$

$$C = \frac{\sin \xi}{\cos \psi} (1 + b)^{-1} \left[ 1 + b \sin^2 \theta \left( 1 + \frac{\cot \xi \tan \theta \sin \psi}{\tan \xi \tan \theta \sin \psi} \right) \right]$$

$$D = -\sin^2 \theta \frac{\sin \xi}{\cos \psi} \left( 1 + \frac{\cot \xi \tan \psi}{\tan \theta} \right)$$

(3.4)

with $\sin \theta = g'/\sqrt{g'^2 + g''^2}$ and $e = g \sin \theta \cos \psi$. The mixing angles in the $M_V \gg M_W$ and large $g''$ limit are

$$\xi = -\frac{\cos 2\theta g}{\cos \theta g''}$$

$$\psi = 2 \sin \theta \frac{g}{g''}$$

(3.5)

The total cross section for the process $e^+e^- \rightarrow f^+f^-$ is given by (at tree level)

$$\sigma = \frac{\pi \alpha_{em}^2 s}{3} \sum_{h_f,h_e} |F(h_f,h_e)|^2$$

(3.6)

with $\alpha_{em} = e^2/(4\pi)$ and

$$F(h_f,h_e) = -\frac{1}{s} e_f + \frac{(v_Z^f + h_fa_Z^f)(v_Z + h_ea_Z)}{s - M_Z^2 + iM_Z\Gamma_Z} + \frac{(v_V^f + h_fa_V^f)(v_V + h_ea_V)}{s - M_V^2 + iM_V\Gamma_V}$$

(3.7)
where \( h_f, h_e = \pm 1 \) are the helicities of \( f \) and \( e \) respectively, \( e_f \) is the electric charge of \( f \) (in units of \( e_{\text{proton}} = 1 \)), \( v_{Z,V} = v_{Z,V}^e \) and \( a_{Z,V} = a_{Z,V}^e \), and \( \Gamma_{Z,V} \) are the widths of the neutral gauge bosons. The partial width of the \( V \) bosons corresponding to decays into fermion-antifermion and \( WW \) is given by

\[
\Gamma_V^h + 3(\Gamma_V^l + \Gamma_V') + \Gamma_V^W
\]

(3.8)

where \( \Gamma_V^h \) is the sum of all quark-antiquark widths (including top), with

\[
\Gamma_V^f = \frac{M_V \alpha_{\text{em}}}{3} \left( v_{V}^2 + a_{V}^2 \right)
\]

(3.9)

\[
\Gamma_V^W = \frac{M_V}{48} \alpha_{\text{em}} g_{VWW}^2 \left( 1 - 4 \frac{M_W^2}{M_V^2} \right)^{1/2} \left( \frac{M_V}{M_W} \right)^4 \times \left[ 1 + 20 \left( \frac{M_W}{M_V} \right)^2 + 12 \left( \frac{M_W}{M_V} \right)^4 \right]
\]

(3.10)

\[
g_{VWW} = \frac{\cos^2 \phi}{\tan \theta} \left( \frac{\sin \xi}{\cos \psi} - \tan \theta \tan \psi \cos \xi \right) + \frac{\cos \xi \sin^2 \phi g''}{2 \sin \theta} \frac{g'}{g}
\]

(3.11)

and, in the limit \( M_V \gg M_W \) and large \( g'' \)

\[
\phi = -\frac{g'}{g''}
\]

(3.12)

The forward-backward asymmetry in the present case is given by

\[
A_{FB}^{e^+e^-\rightarrow f^+f^-} = \frac{x}{1 + \frac{1}{3} x^2} \frac{(1 - P) \sum_{h_f,h_e} h_f h_e |F(h_f,h_e)|^2 + 2 P \sum_{h_f} h_f |F(h_f,1)|^2}{(1 - P) \sum_{h_f,h_e} |F(h_f,h_e)|^2 + 2 P \sum_{h_f} |F(h_f,1)|^2}
\]

(3.13)

where \( x \) is the detector acceptance \((x \leq 1)\), and \( P \) is the degree of longitudinal polarization of the electron beam. In the following analysis we will assume \( x = 1 \).

The left-right asymmetry is given by

\[
A_{LR}^{e^+e^-\rightarrow f^+f^-} = P \frac{\sum_{h_f,h_e} h_e |F(h_f,h_e)|^2}{\sum_{h_f,h_e} |F(h_f,h_e)|^2}
\]

(3.14)

The notations are the same as for the forward-backward asymmetry.

In the following numerical analysis, following the existing studies of 500 GeV \( e^+e^- \) linear colliders [3,19], we assume a relative systematic error in luminosity of \( \delta L/L = 1\% \) and \( \delta \epsilon_{\text{had}}/\epsilon_{\text{had}} = 1\% \) (which is perhaps an optimistic choice due to the problems connected with the \( b \)-jet reconstruction), \( \delta \epsilon_{\mu}/\epsilon_{\mu} = 0.5\% \), where \( \epsilon_{\text{had},\mu} \) denote the selection efficiencies. We shall also assume the same systematic errors for the 1 and 2 TeV machines. Finally we have considered an integrated luminosity \( L = 20 \, fb^{-1} \) for \( \sqrt{s} = 500 \, GeV \), \( L = 80 \, fb^{-1} \) for \( \sqrt{s} = 1 \, TeV \) and \( L = 20 \, fb^{-1} \) for \( \sqrt{s} = 2 \, TeV \). These
integrated luminosities correspond to about one year ($10^7$ sec.) of running. One should of course take into account beamstrahlung effects. However for two body final states, as we consider here, the practical effect is a reduction of the luminosity. This means that with the assumed nominal luminosity one has to run for a correspondingly longer period.

In the case we cannot reach the mass of the resonance we can get restrictions on the parameter space by combining the observables of eq. (3.1). Throughout this paper we assume $m_{t_{top}} = 150$ GeV and $\Lambda = 1$ TeV. Our results are shown in Figs. 2 and 3. Unfortunately the most sensitive observables are the left-right asymmetries, which means that, if the beams are not polarized, one does practically get no advantage over LEP1 from this channel.

The contours shown in Fig. 2 correspond to the regions which are allowed at 90% C.L. in the plane $(b, g/g'')$, assuming that the BESS deviations for the observables of eq. (3.1) from the SM predictions are within the experimental errors. The results are obtained assuming a longitudinal polarization of the electron $P_e = 0.5$ (solid line) and $P_e = 0$ (dashed line). In Fig. 2 we assume $\sqrt{s} = 500$ GeV and $M_V = 600$ GeV. As it is clear there is no big improvement with respect to the already existing bounds from LEP1. Increasing the energy of the machine does not drastically change the results. We have also explored the sensitivity with respect to $M_V$, by choosing $b = 0$ and $P_e = 0.5$. For instance, in Fig. 3, we show the allowed region in the plane $(M_V, g/g'')$ to be compared with the already existing limit $g/g'' < 0.06$ from LEP1 at $b = 0$ (see Fig. 1). Here we consider both the case of $\sqrt{s} = 500$ GeV (solid line) and $\sqrt{s} = 1000$ GeV (dashed line). The bounds improve only for $M_V$ close to the value of the energy of the machine.

All these conclusions become much more negative if one assumes a higher systematical error for the hadron selection efficiency.

4. SENSITIVITY FROM $e^+e^- \rightarrow W^+W^-$

In this Section we will consider the $WW$ channel, which is expected to be more sensitive, at high energy, than the $f\bar{f}$ channel to effects coming from a strongly interacting electroweak symmetry breaking sector. In the case of a vector resonance this is simply due to the strong coupling between the longitudinal $W$ bosons and the resonance. Furthermore this interaction, in general, destroys the fine cancellation among the $\gamma-Z$ exchange and the neutrino contribution occurring in the SM. This effect gives rise, in the case of the BESS model, to a differential cross-section increasing linearly with $s$. However, by using the explicit expression given below one can show that the leading term in $s$ is suppressed by a factor $(g/g'')^4$. Therefore the effective deviation at the energies considered here is given only by the constant term, which is of the order $(g/g'')^2$.

We shall consider the $WW$ channel, for one $W$ decaying leptonically and the other
hadronically. The main reason for choosing this decay channel is to get a clear signal to reconstruct the polarization of the W’s [20]. Let us consider first the following observables:

\[
\frac{d\sigma}{d\cos \theta}(e^+e^- \rightarrow W^+W^-) = \left( \frac{d\sigma}{d\cos \theta}(P_e = +P) - \frac{d\sigma}{d\cos \theta}(P_e = -P) \right) / \frac{d\sigma}{d\cos \theta}
\]

where \( \theta \) is the \( e^+e^- \) center of mass scattering angle. In the \( e^+e^- \) center of mass frame the angular distribution \( d\sigma/d\cos \theta \) and the left-right asymmetry read [11]

\[
\frac{d\sigma}{d\cos \theta} = \frac{2\pi \alpha_{em}^2}{\sqrt{s}} \left\{ 2a_W^4 \left[ \frac{4}{M_W^2} + p^2 \sin^2 \theta \left( \frac{1}{M_W^4} + \frac{4}{t^2} \right) \right] + G_1 p^2 \left[ \frac{4s}{M_W^2} + \left( 3 + \frac{sp^2}{M_W^4} \right) \sin^2 \theta \right] + G_1' \left[ 8 \left( 1 + \frac{M_W^2}{t} \right) + 16 \frac{p^2}{M_W^4} \right. \right. \\
\left. \left. + \frac{p^2}{s} \sin^2 \theta \left( \frac{s^2}{M_W^4} - 2 \frac{s}{M_W^2} - 4 \frac{s}{t} \right) \right] \right\}
\]

and

\[
A_{LR}(\cos \theta) = -P \frac{2\pi \alpha_{em}^2}{\sqrt{s}} \left\{ 2a_W^4 \left[ \frac{4}{M_W^2} + p^2 \sin^2 \theta \left( \frac{1}{M_W^4} + \frac{4}{t^2} \right) \right] + G_2 p^2 \left[ \frac{4s}{M_W^2} + \left( 3 + \frac{sp^2}{M_W^4} \right) \sin^2 \theta \right] + G_1' \left[ 8 \left( 1 + \frac{M_W^2}{t} \right) + 16 \frac{p^2}{M_W^4} \\
\left. \left. + \frac{p^2}{s} \sin^2 \theta \left( \frac{s^2}{M_W^4} - 2 \frac{s}{M_W^2} - 4 \frac{s}{t} \right) \right] \right\} \right\} / \frac{d\sigma}{d\cos \theta}
\]

where

\[
p = \frac{1}{2} \sqrt{s(1 - 4M_W^2 / s)^{1/2}}
\]

\[
t = M_W^2 - \frac{1}{2} s[1 - \cos \theta(1 - 4M_W^2 / s)^{1/2}]
\]

The quantity \( a_W \) is the \( \nu eW \) coupling

\[
a_W = \frac{1}{2\sqrt{2}\sin \theta}(1 + b)^{-1} \left( \cos \phi - \frac{b}{2g} \sin \phi \right)
\]

and

\[
G_1 = \left( \frac{e_e}{s} \right)^2 + (v_Z^2 + a_Z^2) g_{ZWW} \frac{1}{(s - M_Z^2)^2 + M_Z^4 \Gamma_Z^2} + 2 \frac{e_e}{s} v_Z g_{ZWW} \frac{s - M_Z^2}{(s - M_Z^2)^2 + M_Z^4 \Gamma_Z^2} \\
+ (v_V^2 + a_V^2) g_{VWW} \frac{1}{(s - M_V^2)^2 + M_V^4 \Gamma_V^2} + 2 \frac{e_e}{s} v_V g_{VWW} \frac{s - M_V^2}{(s - M_V^2)^2 + M_V^4 \Gamma_V^2} \\
+ 2(v_Zv_V + a_Za_V) g_{ZWWgVWW} \frac{(s - M_Z^2)(s - M_V^2)}{[(s - M_Z^2)^2 + M_Z^4 \Gamma_Z^2][(s - M_V^2)^2 + M_V^4 \Gamma_V^2]}
\]

(4.6)
\[ G_1' = a_W^2 \left[ e_e \frac{e_e}{s} \right. \]
\[ + g_{ZW} (v_2 + a_z) \frac{s - M_Z^2}{(s - M_Z^2)^2 + M_Z^2 \Gamma_Z^2} \]
\[ + g_{VV} (v_V + a_V) \frac{s - M_V^2}{(s - M_V^2)^2 + M_V^2 \Gamma_V^2} \]
\[ G_2 = 2 \left( e_e a_z g_{ZW} \frac{s - M_Z^2}{(s - M_Z^2)^2 + M_Z^2 \Gamma_Z^2} + e_e a_V g_{VV} \frac{s - M_V^2}{(s - M_V^2)^2 + M_V^2 \Gamma_V^2} \right. \]
\[ + \frac{1}{16} a_z v_z g_{ZWW} \left( s - M_Z^2 \right) + a_V v_V g_{VVW} \left( s - M_V^2 \right) \]
\[ \left. + (a_z v_2 + v_z a_V) g_{ZW} g_{VV} \frac{(s - M_Z^2)(s - M_V^2) + M_Z \Gamma_Z M_V \Gamma_V}{((s - M_Z^2)^2 + M_Z^2 \Gamma_Z^2)((s - M_V^2)^2 + M_V^2 \Gamma_V^2)} \right) \]

where
\[ g_{ZW} = \frac{\cos^2 \phi}{\tan \theta} \left( \frac{\cos \xi}{\cos \psi} + \tan \theta \tan \psi \sin \xi \right) \]

and \( g_{VVW} \) as given in eq. (3.11).

Assuming that the final \( W \) polarization can be reconstructed by using the \( W \) decay distributions, it is convenient to examine the cross sections for \( W_L W_L, W_T W_L \), and \( W_T W_T \). One has [21]

\[ \frac{d \sigma_{LL}}{d \cos \theta} = \frac{2 \pi \alpha_{em}^2 p}{\sqrt{s}} \left\{ \frac{a_W^4}{8} \frac{1}{M_W^4 \ell^2} \left[ s^3 (1 + \cos^2 \theta) - 4M_W^4 (3s + 4M_W^2) \right. \right. \]
\[ - 4(s + 2M_W^2) p \sqrt{s s \cos \theta} \sin^2 \theta \]
\[ + \frac{1}{16} G_1 \frac{1}{M_W^4} \sin^2 \theta (s^3 - 12sM_W^4 - 16M_W^6) \]
\[ + G_1' \sin^2 \theta \frac{1}{2t} [p s \sqrt{s \cos \theta} \frac{1}{2M_W^4} (s + 2M_W^2) \]
\[ - \frac{1}{4M_W^4} (s^3 - 12sM_W^4 - 16M_W^6) \right\} \]

\[ \frac{d \sigma_{TL}}{d \cos \theta} = \frac{2 \pi \alpha_{em}^2 p}{\sqrt{s}} \left\{ \frac{a_W^4}{8} \frac{1}{M_W^4 \ell^2} \left[ s^2 (1 + \cos^2 \theta) + 4M_W^4 (1 + \cos^2 \theta) \right. \right. \]
\[ - 4(4p^2 + s \cos^2 \theta) p \sqrt{s \cos \theta} \cos \theta + 2s (s - 6M_W^2) \cos^2 \theta - 4sM_W^2 \]
\[ + 2G_1 s \frac{p^2}{M_W^4} (1 + \cos^2 \theta) \]
\[ + 2G_1' \frac{p s \sqrt{s}}{t M_W^4} \left[ \cos \theta (4p^2 + s \cos^2 \theta) - 2p \sqrt{s} (1 + \cos^2 \theta) \right] \}

\[ \frac{d \sigma_{TT}}{d \cos \theta} = \frac{2 \pi \alpha_{em}^2 p}{\sqrt{s}} \left\{ \frac{2a_W^4}{8} \frac{1}{M_W^4 \ell^2} \left[ s (1 + \cos^2 \theta) - 2M_W^2 - 2p \sqrt{s} \cos \theta \right] \sin^2 \theta \right. \]
\[ + 2G_1 p^2 \sin^2 \theta + G_1' \frac{\sin^2 \theta}{2t} \left[ 4p \sqrt{s} \cos \theta - 8p^2 \right] \}

\]
The left-right asymmetries for longitudinal and/or transverse polarized \( W \) can be easily obtained as in eq. (4.3) by substituting \( G_1 \) by \( G_2 \) in eqs. (4.10), (4.11), (4.12), and dividing by the corresponding differential cross section. In the case of the \( WW \) channel we assume \( \delta B/B = 0.005 \) [22], where \( B \) denotes the product of the branching ratio for \( W \to \text{hadrons} \) and that for \( W \to \text{leptons} \), and we assume 1\% for the acceptance. The assumed value is \( B = 0.29 \).

Figs. 4, 5 show the deviations of the BESS model with respect to the SM with values of the BESS parameters \( b = 0 \) and \( g/g'' = 0.05 \), for the two cases \( \sqrt{s} = 1 \) \( \text{TeV} \), \( M_V = 1.5 \) \( \text{TeV} \), and \( \sqrt{s} = 2 \) \( \text{TeV} \), \( M_V = 2.5 \) \( \text{TeV} \) respectively. Furthermore we plot both the unpolarized and the longitudinally polarized differential cross-section for decay of one \( W \) leptonically and the other hadronically. The systematic errors are the same as discussed before and the statistical errors are evaluated assuming integrated luminosities of \( 80 \) \( \text{fb}^{-1} \) and \( 20 \) \( \text{fb}^{-1} \) for machines at \( \sqrt{s} = 1 \) \( \text{TeV} \) and \( \sqrt{s} = 2 \) \( \text{TeV} \) respectively. We observe that, as already noticed [7, 9], the bigger deviations are away from the forward region. In the longitudinal channel the deviations are much bigger and concentrated in the central region.

To discuss the restrictions on the parameter space for masses of the resonance a little higher than the available energy we have here taken into account the experimental efficiency. We have assumed an overall detection efficiency of 10\% including the branching ratio \( B = 0.29 \) and the loss of luminosity from beamstrahlung [20]. This gives an effective branching ratio of about 0.1. We have again considered an integrated luminosity of \( 20 \) \( \text{fb}^{-1} \), \( 80 \) \( \text{fb}^{-1} \) and \( 20 \) \( \text{fb}^{-1} \) for the three cases \( \sqrt{s} = 0.5, 1 \) and \( 2 \) \( \text{TeV} \).

For a collider at \( \sqrt{s} = 500 \) \( \text{GeV} \) the results are illustrated in Fig. 6. The contours have been obtained by taking 18 bins in the angular region restricted by \( |\cos \theta| < 0.95 \). This figure illustrates the 90\% C.L. allowed regions for \( M_V = 600 \) \( \text{GeV} \) obtained by considering the unpolarized \( WW \) differential cross-section (dotted line), the \( W_L W_L \) cross section (dashed line), and the combination of the left-right asymmetry with all the differential cross-sections for the different final \( W \) polarizations (solid line). We see that already at the level of the unpolarized cross-section we get important restrictions with respect to LEP1. In Fig. 7 we have examined the possibility that the total number of events is reduced of a factor 0.5 from the losses due to the reconstruction of the polarization of the \( W \)’s. We see that, least to say, this does not reduce the efficiency of the machine with respect to LEP1 in restricting the BESS parameter space.

For colliders with \( \sqrt{s} = 1, 2 \) \( \text{TeV} \) and for \( M_V = 1.2 \) and \( 2.5 \) \( \text{TeV} \) respectively, the allowed region, combining all the observables, reduces in practice to a line and the analysis is better discussed in the plane \( (M_V, g/g'') \), as we shall see later on. Therefore, even the unpolarized \( WW \) differential cross section measurements can improve the bounds, as shown in Fig. 8 for \( \sqrt{s} = 2 \) \( \text{TeV} \) and \( M_V = 2.5 \) \( \text{TeV} \).

Finally in Figs. 9-11 we assume \( b = 0 \), that is no direct coupling of the \( V \) to the fermions, and we analyze the allowed region in the plane \( (M_V, g/g'') \). We work in the same hypotheses as before. Recalling that the bound on \( g/g'' \) from LEP1 is
about 0.06 for $b = 0$, we see that if the $W$ polarization is not observed the machine at $\sqrt{s} = 500 \text{ GeV}$ improves the LEP1 limit up to $M_V \approx 1 \text{ TeV}$. If the $W$ polarization is measured the LEP1 limit is improved in all $M_V$ range. For colliders at $\sqrt{s} = 1, 2 \text{ TeV}$ we get a big improvement on the bound for all the situations.

5. SENSITIVITY FROM FUSION SUBPROCESSES

Another mechanism to produce $W^+W^-$ pairs is the fusion of a pair of ordinary gauge bosons, each being initially emitted from an electron or a positron. In the so-called effective-W approximation the initial $W, Z, \gamma$ are assumed to be real and the cross section for producing a $W^+W^-$ pair is obtained by a convolution of the fusion subprocess with the luminosities of the initial $W, Z, \gamma$ inside the electrons and positrons. There are two fusion subprocesses which contribute to produce $W^+W^-$ pairs. The first one is $e^+e^- \rightarrow W_{L,T}^0W_{L,T}^0e^+e^-$. It is mediated by $W^\pm$ and $V^\pm$ exchanges in the $t$ and $u$ channels. The second fusion subprocess we consider is \( e^+e^- \rightarrow W^+_{L,T}W^-_{L,T}\nu\bar{\nu} \). It is mediated by $\gamma, Z$ and $V^0$ exchanges in the $s$ and $t$ channels. Both processes get a contribution from the gauge boson quadrilinear couplings.

In principle, the fusion processes are interesting because they allow to study a wide range of mass spectrum for the $V$ resonance from one given $e^+e^-$ c.m. energy.

In the $e^+e^-$ center-of-mass frame the invariant mass distribution $d\sigma/dM_{WW}$ reads

\[
\frac{d\sigma}{dM_{WW}} = \frac{1}{4\pi s} \frac{1}{M_{WW}^2} \sum_{i,j} \sum_{l_1,l_2} \int_{(p_T^2)_{\text{min}}}^{M_{WW}^2/4} dp_T^2 \int_{\log \tau}^{-\log \sqrt{\tau}} dy \ f_i^{l_1}(\sqrt{\tau}e^y) f_j^{l_2}(\sqrt{\tau}e^{-y}) \\
\times \frac{1}{p} \frac{1}{\sqrt{M_{WW}^2 - 4p_T^2}} |M(V_i^{l_1}V_j^{l_2} \rightarrow W_{i_1}^{\pm}W_{i_2}^{\pm})|^2
\]

where $p_T$ is the transverse momentum of the outgoing $W$, $\tau = M_{WW}^2/s$, $p$ and $p'$ are the absolute values of the three momenta for incoming and outgoing pairs of vector bosons: $p = (E_1^2 - M_1^2)^{1/2} = (E_2^2 - M_2^2)^{1/2}$ and $p' = (\sqrt{M_{WW}^2/2})(1 - 4M_1^2/M_{WW}^2)^{1/2}$ with $E_i$ the fraction of the electron (or positron) energy of the vector boson $V_i$ with mass $M_i$ and helicity $l_i$. The structure functions $f$ appearing in the previous formula are

\[
\begin{align*}
  f^+(x) &= \frac{\alpha_{\text{em}}}{4\pi} \frac{(v + a)^2}{x} + (1 - x)^2 \frac{(v - a)^2}{x} \log \frac{s}{M^2} \\
  f^-(x) &= \frac{\alpha_{\text{em}}}{4\pi} \frac{(v - a)^2}{x} + (1 - x)^2 \frac{(v + a)^2}{x} \log \frac{s}{M^2} \\
  f^0(x) &= \frac{\alpha_{\text{em}}}{\pi} \left( v^2 + a^2 \right) \frac{1 - x}{x}
\end{align*}
\]

and represent the probability of having inside the electron a vector boson of mass $M$ with fraction $x$ of the electron energy. In eq. (5.2) $v$ and $a$ are the vector and axial-vector couplings of the gauge bosons to fermions. More precisely for the photon $v = -1$. 

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and \(a = 0\), for \(Z v = v_Z^f\) and \(a = a_Z^f\) as given in eq. (3.3), and for \(W v = a = a_W\) as given in eq. (4.5). The quadrilinear couplings are

\[
g_{WWWW} = \frac{e^2}{\sin^2 \theta \cos^2 \psi} \frac{\cos^4 \phi}{4} + \frac{g''^2}{4} \sin^4 \phi
\]

\[
g_{WZZ} = \frac{e^2}{\cos^2 \phi} g_Z W W + \frac{g''^2}{4} \sin^2 \phi \sin^2 \xi \cos^2 \psi
\]

\[
g_{WW \gamma \gamma} = e^2 \cos^2 \phi + \frac{g''^2}{4} \sin^2 \phi \sin^2 \psi
\]

\[
g_{WWZ \gamma} = e^2 g_{ZW} W - \frac{g''^2}{4} \sin^2 \phi \sin \xi \sin \psi \cos \psi
\]

with \(g_{ZW} W\) as given in eq. (4.9) and the mixing angles \(\xi, \psi\) and \(\phi\) in eqs. (3.5) and (3.12).

The amplitudes of the vector bosons scattering processes within the BESS model are:

\[
M(W_1^+ W_2^- \rightarrow W_3^+ W_4^-) = -ie^2 \frac{f_s}{s} - ie^2 g_{ZW} W \frac{f_s}{s - M_Z^2}
\]

\[
- ie^2 g_{WW} W \frac{f_s}{s - M_W^2 + i \Gamma_W M_W} + (s \rightarrow t)
\]

\[
+ ig_{WW} W [2(\epsilon_1 \cdot \epsilon_2)(\epsilon_2^* \cdot \epsilon_4^*) - (\epsilon_1 \cdot \epsilon_2)(\epsilon_3^* \cdot \epsilon_4^*) - (\epsilon_1 \cdot \epsilon_3^*)(\epsilon_2 \cdot \epsilon_4^*)]
\]

\[
M(\gamma_1 \gamma_2 \rightarrow W_3^+ W_4^-) = -ie^2 \frac{h_t}{t - M_W^2} + (t \rightarrow u)
\]

\[
- ig_{WW} W [2(\epsilon_1 \cdot \epsilon_2)(\epsilon_3^* \cdot \epsilon_4^*) - (\epsilon_1 \cdot \epsilon_3^*)(\epsilon_2 \cdot \epsilon_4^*) - (\epsilon_1 \cdot \epsilon_4^*)(\epsilon_2 \cdot \epsilon_3^*)]
\]

\[
M(\gamma_1 Z_2 \rightarrow W_3^+ W_4^-) = -ie^2 g_{ZW} W \frac{h_t}{t - M_Z^2} + (t \rightarrow u)
\]

\[
- ig_{WW} Z [2(\epsilon_1 \cdot \epsilon_2)(\epsilon_3^* \cdot \epsilon_4^*) - (\epsilon_1 \cdot \epsilon_3^*)(\epsilon_2 \cdot \epsilon_4^*) - (\epsilon_1 \cdot \epsilon_4^*)(\epsilon_2 \cdot \epsilon_3^*)]
\]

\[
M(Z_1 Z_2 \rightarrow W_3^+ W_4^-) = -ie^2 g_{ZW} W \frac{h_t}{t - M_W^2}
\]

\[
- ie^2 g_{WW} W \frac{h_t}{t - M_V^2 + i \Gamma_V M_V} + (t \rightarrow u)
\]

\[
- ig_{WW} Z [2(\epsilon_1 \cdot \epsilon_2)(\epsilon_3^* \cdot \epsilon_4^*) - (\epsilon_1 \cdot \epsilon_3^*)(\epsilon_2 \cdot \epsilon_4^*) - (\epsilon_1 \cdot \epsilon_4^*)(\epsilon_2 \cdot \epsilon_3^*)]
\]

where

\[
f_s = [(\epsilon_1 \cdot \epsilon_2)(p_2 - p_1)\lambda - 2(\epsilon_1 \cdot p_2)\epsilon_2 \lambda + 2(\epsilon_2 \cdot p_1)\epsilon_1 \lambda]
\]

\[
\cdot [(\epsilon_3^* \cdot \epsilon_4^*)(p_3 - p_4)\lambda - 2(\epsilon_4^* \cdot p_3)\epsilon_3 \lambda + 2(\epsilon_3^* \cdot p_4)\epsilon_4 \lambda]
\]
and $f_t$ can be deduced from $f_s$ with the substitution $p_2 \leftrightarrow -p_3$ and $\epsilon_2 \leftrightarrow \epsilon_3^*$, while

\[
\begin{align*}
  h_t &= [2(\epsilon_1 \cdot p_3)\epsilon_3^* - (\epsilon_1 - \epsilon_3^*)(p_3 + p_1)\lambda + 2(p_1 \cdot \epsilon_3^*)\epsilon_1\lambda] \\
  &\quad \cdot [2(\epsilon_2 \cdot p_4)\epsilon_4^* - (\epsilon_2 - \epsilon_4^*)(p_4 + p_2)\lambda + 2(p_2 \cdot \epsilon_4^*)\epsilon_2\lambda] \\
  &\quad + (\epsilon_1 \cdot \epsilon_3^*)(p_1^2 - p_3^2)(\epsilon_2 \cdot \epsilon_4^*)(p_2^2 - p_4^2)/M_W^2
\end{align*}
\]

and $h_u$ can be deduced from $h_t$ with the substitution $p_3 \leftrightarrow p_4$ and $\epsilon_3^* \leftrightarrow \epsilon_4^*$.

In eqs. (5.4-7) $g_{ZWW}$ and $g_{VVW}$ are given in eq. (4.9) and (3.11) respectively and $\Gamma_V$ is the width of the $V$ resonance (we have not explicitly inserted the widths of the standard gauge bosons because they are irrelevant for this calculation).

For comparison, the amplitudes within the SM can be obtained from eqs. (5.4-8) by taking all the trilinear and quadrilinear vector boson couplings in the limit $g'' \to \infty$ and $b \to 0$ and adding the contribution due to the Higgs boson exchange (see ref. [23]). We are not considering here all the non-annihilation graphs contributing to the processes $e^+e^- \to W^+W^-e^+e^-$ and $e^+e^- \to W^+W^-\nu\bar{\nu}$ (at the order $\alpha_{em}^4$) in which the final $W$'s are emitted from the electron (positron) legs. This because we expect their contribution to lie mostly in a kinematical region different from the one we are interested in ($p_T \sim M_W$).

The result of our analysis is that there are not significative differences between the SM and the BESS model differential cross section in the case of the process $e^+e^- \to W^+W^-e^+e^-$. This is due, first of all, to the absence of the $s$ channel exchange of the $V$ resonance, secondly, to the dominance of the $\gamma\gamma$ fusion contribution, and to the fact that in the BESS model the couplings of the photon to the fermions and to $W^+W^-$ are the same as in the SM.

Concerning the process $e^+e^- \to W^+W^-\nu\bar{\nu}$, we have evaluated the differential cross sections $d\sigma/dM_{WW}$ both for the SM with $M_H = 100$ GeV and for the BESS model. The only channel which turns out to be useful is the one corresponding to longitudinally polarized final $W$'s. The results are illustrated in Figs. 12, 13 for two different choices of the BESS parameters. In Fig. 12 we compare $d\sigma/M_{WW}(LL)$ for the SM (dashed line) and for the BESS model (solid line) for $\sqrt{s} = 1.5$ TeV, $b = 0.01$, $g'' = 13$ and $M_V = 1$ TeV. In Fig. 13 and $\sqrt{s} = 2$ TeV, the mass of the resonance is $M_V = 1.5$ TeV and the other BESS parameters are the same as in Fig. 12. In both cases we are not applying any cuts except for $(p_T)_{min} = 10$ GeV. For simplicity we have only introduced the partial width of eq. (3.8).

Although theoretically there is a clear difference between the curves of the two models, the experimental situation is quite different.

Let us first consider the case of $\sqrt{s} = 1.5$ TeV. By integrating the differential cross section for $500 < M_{WW}(GeV) < 1500$ and considering an integrated luminosity of $80 \, fb^{-1}$ we obtain 127 $W_L$ pairs for the SM and 158 for the BESS model (with $M_V = 1$ TeV, $g'' = 13$ and $b = 0.01$) corresponding to a statistical significance of 2.75.
This result is quite discouraging considering the fact that we have not included the branching ratio. The situation does not significantly improve by considering a wider resonance, when varying the BESS parameters in the region allowed by the present bounds (see Fig. 1).

By increasing the energy to $2\, TeV$ and by considering an integrated luminosity of $20\, fb^{-1}$, the situation gets even worse. By integrating the differential cross section for $1000 < M_{WW}(GeV) < 2000$ we get 7 $W_L$ pairs for the SM and 13 for the BESS model (with $M_V = 1.5\, TeV$, $g'' = 13$ and $b = 0.01$) corresponding to a statistical significance of 2.27.

The channel corresponding to transverse-longitudinal final $W$’s leads to a very small bump in the region of the resonance above the SM backgrounds, which is not observable.

Finally, we have also considered the fusion process in the charged channel $e^+e^- \rightarrow W_{L,T}^+Z_{L,T}\bar{\nu}e^-$. However, by comparing with the previous process $e^+e^- \rightarrow W_{L,T}^+W_{L,T}^-\bar{\nu}$, we can observe that, in this case, the SM cross section is bigger, being dominated by the $\gamma W \rightarrow WZ$ fusion process, while the BESS effect $WZ \rightarrow V \rightarrow WZ$ is of the same order of magnitude. So we expect a worse signal to background ratio.

6. CONCLUSIONS

We have shown the usefulness of very energetic $e^+e^-$ linear colliders in exploring an alternative scheme of electroweak symmetry breaking based on the existence of vector resonances from strong sector.

Such a study is complementary to those performed in the case of the LHC (SSC) proton colliders. In fact $pp$ colliders give a valuable opportunity to study the $V^\pm$ resonances through the $W^\pm Z$ decay [24], while the $V^0 \rightarrow W^+W^-$ channel is difficult to study in $pp$ due to background problems, and $V^0 \rightarrow l^+l^-$ has a very low rate [25]. On the contrary $e^+e^-$ colliders give the possibility of detecting new neutral vector bosons. This can be relevant in order to distinguish among the various models with additional vector bosons. For example, in left-right symmetric models the charged and the neutral vector resonance masses are split while in the BESS model they are degenerate (neglecting the electroweak corrections).

Our investigation shows that, even if the mass of the neutral resonance is higher than the c.m. energy of the collider, the process of $W$ pair production will allow to put very severe restrictions on the parameter space of the BESS model, especially so if the $W$ polarizations can be reconstructed from their decay distributions.

If no deviation from the SM prediction is found, already at $\sqrt{s} = 500\, GeV$ and integrated luminosity $L = 20\, fb^{-1}$, the BESS model parameters $g''$ and $b$ can be severely restricted, while for $M_V$ we find significant improvement with respect to LEP1 when the final $W$ polarization is reconstructed. With higher energy colliders ($\sqrt{s} = 1\, TeV$,
and \( L = 80 \text{ fb}^{-1} \) and \( \sqrt{s} = 2 \text{ TeV} \) and \( L = 20 \text{ fb}^{-1} \) the parameter space contracts and at \( b = 0 \) we get an upper bound on \( g/g'' \) of the order of 0.02, for any given value of \( M_V \).

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FIGURE CAPTION

Fig. 1 90% C.L. contours in the plane \((b, g/g'')\) for \(M_V = 1000 \text{ GeV}\), from LEP1 and CDF/UA2 data (we assume \(\alpha_s = 0.12\), and \(\Lambda = 1000 \text{ GeV}\)) for \(m_{\text{top}} = 120 \text{ GeV}\) (dashed line), \(m_{\text{top}} = 150 \text{ GeV}\) (solid line), and \(m_{\text{top}} = 180 \text{ GeV}\) (dotted line). The allowed regions are the internal ones. For increasing \(\alpha_s\) the allowed regions get shifted to the left.

Fig. 2 90% C.L. contours in the plane \((b, g/g'')\) for \(\sqrt{s} = 500 \text{ GeV}\) and \(M_V = 600 \text{ GeV}\) from the fermion channel (we assume \(m_{\text{top}} = 150 \text{ GeV}\), \(\alpha_s = 0.12\), \(\Lambda = 1000 \text{ GeV}\)). The solid line corresponds to polarization \(P_e = 0.5\) while the dashed line is for unpolarized electron beams. The allowed regions are the internal ones.

Fig. 3 90% C.L. contours in the plane \((M_V, g/g'')\) for \(b = 0\), from the fermion channel (we assume \(m_{\text{top}} = 150 \text{ GeV}\), \(\alpha_s = 0.12\), \(\Lambda = 1000 \text{ GeV}\), and polarization \(P_e = 0.5\)) for different choices of \(\sqrt{s} (\text{GeV})\): 500 (solid line), 1000 (dashed line). The lines give the upper bounds on \(g/g''\).

Fig. 4 Unpolarized and longitudinally polarized differential cross-section \(d\sigma/d\cos \theta \text{ (in pb)}\) in the WW channel, for one W decaying leptonically and the other hadronically, versus \(\cos \theta\) for the SM (dash line) and BESS model (solid line) corresponding to \(\sqrt{s} = 1000 \text{ GeV}\), \(M_V = 1500 \text{ GeV}\), \(b = 0\), and \(g/g'' = 0.05\). The error bars are the total errors and correspond to one standard deviation. The lower curves are for the \(W_LW_L\) channel.

Fig. 5 Same of Fig. 4 for \(\sqrt{s} = 2000 \text{ GeV}\) and \(M_V = 2500 \text{ GeV}\).

Fig. 6 90% C.L. contours in the plane \((b, g/g'')\) for \(\sqrt{s} = 500 \text{ GeV}\) and \(M_V = 600 \text{ GeV}\) from the unpolarized WW differential cross section (dotted line), from the \(W_LW_L\) differential cross section (dashed line) and from all the differential cross sections for \(W_LW_L, W TW_L, W TW_T\) combined with the WW left-right asymmetries (solid line). The allowed regions are the internal ones.

Fig. 7 90% C.L. contour in the plane \((b, g/g'')\) for \(\sqrt{s} = 500 \text{ GeV}\) and \(M_V = 600 \text{ GeV}\) from all the differential cross sections for \(W_LW_L, W TW_L, W TW_T\) combined with the WW left-right asymmetries, with the total number of events reduced by a factor 0.5 to account for possible losses due to the reconstruction of the polarization of the W’s. The allowed region is the internal ones.

Fig. 8 90% C.L. contour in the plane \((b, g/g'')\) for \(\sqrt{s} = 2000 \text{ GeV}\), an integrated luminosity of \(20 \text{ fb}^{-1}\) and \(M_V = 2500 \text{ GeV}\) from the unpolarized WW differential cross section. The allowed region is the internal ones.

Fig. 9 90% C.L. contours in the plane \((M_V, g/g'')\) for \(\sqrt{s} = 500 \text{ GeV}\), \(L = 20 \text{ fb}^{-1}\) and \(b = 0\). The solid line corresponds to the bound from the unpolarized WW
differential cross section, the dashed line to the bound from the longitudinally polarized $W_L W_L$ differential cross section, the dotted line to the bound from all the polarized differential cross sections $W_L W_L$, $W_T W_L$, $W_T W_T$ combined with the $WW$ left-right asymmetries. The lines give the upper bounds on $g/g''$.

Fig. 10 Same as Fig. 9 for $\sqrt{s} = 1000$ GeV and $L = 80$ $fb^{-1}$. Here the dashed and the dotted lines are almost coincident.

Fig. 11 Same as Fig. 10 for $\sqrt{s} = 2000$ GeV and $L = 20$ $fb^{-1}$. Here the dashed and the dotted lines are coincident.

Fig. 12 Longitudinally polarized differential cross-section $d\sigma/dM_{WW}(e^+ e^- \rightarrow W_L^+ W_L^- \nu\bar{\nu})$ (in $fb/GeV$) versus $M_{WW}$ for the SM (dash line) and BESS model (solid line) corresponding to $\sqrt{s} = 1500$ GeV, $M_V = 1000$ GeV, $b = 0.01$, and $g'' = 13$.

Fig. 13 Same of Fig. 12 for $\sqrt{s} = 2000$ GeV, $M_V = 1500$ GeV, $b = 0.01$, and $g'' = 13$. 