ONE LOOP CALCULATION OF THE $\epsilon_3$ PARAMETER
WITHIN THE EXTENDED BESS MODEL

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ABSTRACT

The existence of a strongly interacting sector responsible for the electroweak symmetry breaking is assumed. As a consequence vector and axial-vector bound states may be formed. These resonances mix with the Standard Model gauge bosons and are of primary phenomenological importance for the LEP physics. The extended BESS model is an effective scheme based on the symmetry group $SU(8)_L \otimes SU(8)_R$, describing in a consistent way the interactions among the pseudo-Goldstone bosons, vector and axial-vector resonances and the standard gauge bosons. In a previous paper, the contribution from extended BESS to the electroweak oblique corrections was evaluated. However, only an estimate of the effects coming from mass and wave function renormalization of the new resonances, was given. Here we complete the evaluation by computing explicitly these effects. We confirm the previous result, that is, in spite of the great precision of the present LEP measurements, the extended BESS parameter space is not very much constrained.
1 Introduction

Precision measurements are now giving serious restrictions on the possibility of a new strong interacting sector being at the origin of the symmetry breaking in the electroweak theory. A natural consequence of such a strong sector is the occurrence of resonances in the TeV region. Among them, spin-1 resonances would be of particular interest as they would, already through mixing effects, affect the self-energies of the standard model gauge bosons.

An effective scheme (BESS model ref. [1]), introduced several years ago, describes a minimal version of such a scenario by introducing a triplet of vector resonances and no Goldstone bosons, being based on a non linear $SU(2)_L \otimes SU(2)_R$ chiral model.

If the strongly interacting sector gives rise to pseudo-Goldstone bosons, it has been shown [2] that their loop contribution to the vector boson self-energies is not negligible. For this reason, in ref. [3], an extended version of the BESS model, based on a chiral $SU(8)_L \otimes SU(8)_R$, was given.

The model contains 60 pseudo-Goldstone bosons in the spectrum and, as a further generalization, it describes also axial-vector resonances together with vector ones. The physics of this model, as far as the LEP experiments are concerned, is such that it gives corrections only to the self-energies of the standard electroweak gauge bosons. These corrections have been evaluated in ref. [3] at one-loop level by defining the theory with a cut-off $\Lambda$. In that calculation the contributions coming from mass and wave function renormalization of the new resonances were evaluated in an approximate way using a dispersive representation. Here we want to give a more complete analysis by including in the calculation all the one-loop diagrams contributing to self-energies of the new vector and axial-vector resonances.

2 Gauge boson self-energies within the extended BESS model

The extended BESS model effective Lagrangian, as introduced in ref. [3], describes vector and axial-vector gauge bosons interacting with the Goldstone bosons and the Standard Model (SM) vector bosons.

The effective Lagrangian is obtained by enlarging the symmetry group of the BESS model based on a non-linear $SU(2)_L \otimes SU(2)_R \sigma$ model as described in ref. [1], to contain a "hidden" $SU(8)_L \otimes SU(8)_R$ gauge group, by introducing the covariant derivatives with respect to the gauge fields associated to the "hidden" symmetry and with respect to the SM gauge fields, and by adding together all the independent invariant terms containing at most two derivatives.

The new resonances $V^A$ and $A^A$ ($A=1,...,63$) are described as Yang-Mills fields which acquire mass through the same symmetry breaking mechanism which gives mass to the ordinary gauge bosons.

Working in the unitary gauge for the $V$, $A$ and SM gauge bosons leaves us with 60 Goldstones as physical particles in the spectrum. We will assume that they acquire mass through radiative corrections.
An important property is that the mixing terms among the SM gauge bosons and the new vector and axial-vector resonances are the same as in the axial-vector extension of the $SU(2)$-BESS model \cite{4} plus an extra term in the neutral sector involving the singlet (under $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$) field $V_D$. The neutral Lagrangian which gives the mixing among the new resonances and the standard $SU(2)_L \otimes U(1)_Y$ gauge bosons is

$$\mathcal{L}^{\text{mix}} = \frac{g^2}{8} \left[ a(gW^3 - g'Y)^2 + b(gW^3 + g'Y - g''V^3)^2 \right. $$

$$+ b\left( \frac{2}{\sqrt{3}}g'Y - g''V_D \right)^2 + c(gW^3 - g'Y + g''A^3)^2 + dg''^2(A^3)^2 \right]$$ (2.1)

where $a$, $b$, $c$, $d$ are free parameters, $g$ and $g'$ are the gauge coupling constants of $W$ and $Y$ respectively, and $g''$ is the self coupling of the new gauge bosons. From eq. (2.1) we see that the mixing terms are proportional to $x \equiv g/g''$ so one can decouple the new resonances from the SM gauge bosons, by taking the large $g''$ limit (small $x$). The SM results are recovered by putting $x = 0$ and fixing the normalization $a + cd/(c + d) = 1$.

By diagonalizing the mass matrix it turns out that the SM gauge boson masses get corrected by terms of the order of $(g/g'')^2$ (the photon remains massless) while the vector bosons (and the axial ones) are degenerate if one neglects the weak corrections, with their masses proportional to $g''^2$.

In ref. \cite{3} we have explicitly evaluated the interaction terms among the SM gauge bosons, the new resonances and the pseudo-Goldstone bosons (PGB). In the following application, we will need also the kinetic terms for the $V$ and $A$ resonances which are triplets under $SU(2)$. The explicit expression is given in ref. \cite{5}.

The BESS model parameter space is four-dimensional; we will choose, as free parameters, the masses of the vector and the axial-vector resonances $M_V$ and $M_A$, their gauge coupling constant $g''$ and $z = c/(c + d)$, which measures the ratio of the mixings $W - A$ and $W - V$.

The mixing described in (2.1) induces corrections to the self-energies of the SM. We define the scalar part of the SM vector boson self-energies through the relation

$$\Pi_{ij}^{\mu \nu}(p^2) = -i\Pi_{ij}(p^2)g^{\mu \nu} + p^{\mu}p^{\nu} \text{ terms}$$ (2.2)

where the indices $i$ and $j$ run over the ordinary gauge vector bosons.

In the neutral sector we choose to work with the SM fields ($\theta$ is the Weinberg angle)

$$W^3 = c_{\theta} Z + s_{\theta} \gamma$$ (2.3)

$$Y = -s_{\theta} Z + c_{\theta} \gamma$$ (2.4)

The corrections from the vector and axial-vector resonances and from the PGB of extended BESS are purely oblique and only affect the scalar gauge boson self-energy terms $\Pi_{WW}$, $\Pi_{33}$, $\Pi_{30}$, and $\Pi_{00}$.

It is then convenient to introduce the following combinations (see ref. \cite{3}):

$$\epsilon_1 = \frac{\Pi_{33}(0) - \Pi_{WW}(0)}{M_W^2}$$ (2.5)

$$\epsilon_2 = \frac{\Pi_{WW}(M_W^2) - \Pi_{WW}(0)}{M_W^2} - \frac{\Pi_{33}(M_Z^2) - \Pi_{33}(0)}{M_Z^2}$$ (2.6)

$$\epsilon_3 = \frac{c_{\theta} \Pi_{30}(M_Z^2) - \Pi_{30}(0)}{s_{\theta} M_Z^2}$$ (2.7)
In ref. [3] we have shown that the tree-level contribution to the \( \epsilon \) parameters coming from the mixing of the \( V \) and \( A \) bosons with the \( W \), \( Z \) and \( \gamma \) is the same we found for \( SU(2) \)-BESS model in ref. [5], or

\[
\epsilon_1 = 0 \quad \epsilon_2 \simeq 0 \quad \epsilon_3 \simeq x^2 (1 - z^2) \quad (2.8)
\]

where the last two results were obtained for large \( M_{V,A} \) \( (M_{V,A} \gg M_{W,Z}) \). The reason is that the new \( V_D \) boson mixes only with \( Y \) (see (2.1)) and, as a consequence, it does not affect \( \Pi_{30} \) and it contributes only to the definition of the electric charge and \( M_Z \). This means that our result can be extended to a \( SU(N)_L \otimes SU(N)_R \) model, in fact in these models the new fields associated to the diagonal generators will be mixed only with the hypercharge field \( Y \).

In addition to the self-energy corrections arising from the mixing, one has loop corrections. In ref. [3] we have implemented the calculation of the \( \epsilon \) parameters at one-loop level. Being BESS a non renormalizable model, the loop integrals were evaluated using a cut-off \( \Lambda \). In Fig. 1 we list the graphs contributing to \( \epsilon_3 \) at one-loop level. In ref. [3] the calculation was performed neglecting the contributions to \( \epsilon_3 \) coming from the self-energies of the new resonances (which are schematically represented in Fig. 1 by the first two graphs) and we presented a simple estimate of these contributions through dispersive integrals. In particular the first two graphs in Fig. 1 represent the propagators for the \( V^3 \) and \( A^3 \) gauge bosons, corrected and opportunely renormalized at one-loop level, which must be inserted, via mixing, in the \( \Pi_{30} \) function. Here we will give the complete one-loop result for \( \epsilon_3 \).

### 3 One-loop evaluation of \( \epsilon_3 \)

The graphs contributing to the \( V^3 \) and \( A^3 \) self-energies are listed in Fig. 2. They have been evaluated using the technique of dimensional regularization. We have not included the tadpoles because they give a constant (as function of \( p^2 \)) contribution to the self-energies and so they can be disposed of by mass renormalization. Furthermore, being \( \epsilon_3 \) an isospin symmetric observable, we have assumed a common non-vanishing mass \( M_\pi \) for all the PGB.

The first contribution to the \( A_3 \) self-energy, given by the graph in Fig. 2a, is

\[
\Pi_{A_3}^{(a)}(p^2) = 15 \frac{g_{V,A}^2}{16\pi^2} \left[ \left( \frac{3}{4} - \frac{1}{4} \frac{M_\pi^2}{M_V^2} + \frac{p^2}{12M_V^2} \right) \left( \ln \frac{\Lambda^2}{M_V M_\pi} - \gamma \right) + \frac{M_\pi^2 - M_V^2}{2p^2} \left( 1 + \frac{1}{12p^2 M_V^2} - \frac{M_\pi^2 - M_V^2}{4M_V^2} \right) \ln \frac{M_\pi^2}{M_V^2} \right]
\]

\[
+ \left( \frac{1}{3} - \frac{5}{3} \frac{M_\pi^2}{M_V^2} + \frac{2}{9} \frac{p^2}{M_V^2} - \frac{1}{6M_V^2} \frac{(M_\pi^2 - M_V^2)^2}{p^2} \right) \beta(M_V^2, M_\pi^2, p^2) \ln l(M_V, M_\pi, p^2)
\]

\[
+ \frac{17}{12} - \frac{7}{12} \frac{M_\pi^2}{M_V^2} + \frac{2}{9} \frac{p^2}{M_V^2} + \frac{(M_\pi^2 - M_V^2)^2}{12M_V^4 p^2} \right] \quad (3.1)
\]
where the factor 15 gives the number of the pairs \( V\pi \) circulating in the loop (we are working in the unitary gauge where the \( SU(2) \) triplet \( \pi^a \) is eaten up by \( W^\pm, Z \)) and

\[
\beta(M_1^2, M_2^2, p^2) = \frac{\sqrt{(M_1^2 - M_2^2)^2 - 2(M_1^2 + M_2^2)p^2 + p^4}}{2p^2}
\]

(3.2)

\[
l(M_1, M_2, p^2) = \frac{s_+(M_1, M_2, p^2) - s_-(M_1, M_2, p^2)}{s_+(M_1, M_2, p^2) + s_-(M_1, M_2, p^2)}
\]

(3.3)

\[
s_{\pm}(M_1, M_2, p^2) = \sqrt{(M_1 \pm M_2)^2 - p^2}
\]

(3.4)

\( \gamma \) is the Euler’s constant (\( \simeq 0.577 \)), and \( g_{VA}(\pi) \) is a trilinear coupling given in ref. [3]

\[
g_{VA}(\pi) = \frac{v}{4} \varepsilon^{2} \frac{x^2}{r_V} \left( \frac{r_V}{r_A} - 1 \right)
\]

(3.5)

with \( v \simeq 246 \text{ GeV} \) and \( r_{VA} \simeq M_W^2/M_V^2 \) in the large \( g'' \) limit.

The second contribution to \( \Pi_A \) is
given by the graph in Fig. 2b, is

\[
\Pi_A^{(0)}(p^2) = 16 \frac{g''^2}{16 \pi^2} \left[ a(M_A^2, M_V^2, p^2) + b(M_A^2, M_V^2, p^2) \left( \ln \frac{\Lambda^2}{M_A^2 M_V^2} - \gamma \right) \right.
\]

\[
\left. + c(M_A^2, M_V^2, p^2) \beta(M_A^2, M_V^2, p^2) \ln l(M_A, M_V, p^2) \right]
\]

(3.6)

\[
+ d(M_A^2, M_V^2, p^2) \ln \left( \frac{M_V^2}{M_A^2} \right)
\]

(3.7)

where, again the factor 16 is a multiplicity factor, and

\[
a(x_1, x_2, t) = \frac{1}{72x_1x_2t} \left( -405x_1x_2(x_1 + x_2)t + 90(x_1^3 + x_2^3)t + 6f(x_1, x_2, t) - 452x_1x_2t^2 - 182(x_1^2 + x_2^2)t^2 + 16t^4 + 70(x_1 + x_2)t^4 \right)
\]

(3.8)

\[
b(x_1, x_2, t) = \frac{1}{12x_1x_2} \left( 9(x_1^3 + x_2^3) - 36x_1x_2(x_1 + x_2) - 17(x_1^2 + x_2^2)t \right.
\]

\[
- 50x_1x_2t + 7(x_1 + x_2)t^2 + t^3 \right)
\]

(3.9)

\[
c(x_1, x_2, t) = \frac{1}{6x_1x_2} \left( -f(x_1, x_2, t) - 8(x_1^3 + x_2^3)t + 32x_1x_2(x_1 + x_2)t + 18(x_1^2 + x_2^2)t^2 \right.
\]

\[
+ 32x_1x_2t^2 - 8(x_1 + x_2)t^3 - t^4 \right)
\]

(3.10)

\[
d(x_1, x_2, t) = \frac{x_1 - x_2}{24x_1x_2t^2} \left( f(x_1, x_2, t) + 7(x_1^3 + y_1^3)t - 43x_1x_2(x_1 + x_2)t - 17(x_1^2 + x_2^2)t^2 \right.
\]

\[
- 53x_1x_2t^2 + 9(x_1 + x_2)t^3 \right)
\]

(3.11)

The contribution to the \( V^3 \) self-energy coming from the graph in Fig. 2c can be evaluated from \( \Pi_A^{(0)} \) with the substitution \( M_V \rightarrow M_A \).

From the graph in Fig. 2d we get

\[
\Pi_V^{(d)}(p^2) = 15 \frac{g''^2}{8 \pi^2} \left[ \left( \frac{p^2}{6} - M_\pi^2 \right) \left( \ln \frac{\Lambda^2}{M_\pi^2} - \gamma \right) \right.
\]

\[
+ \frac{8}{3} p^2 \alpha^3(M_\pi^2, p^2) \arctan \left( \frac{1}{2\alpha(M_\pi^2, p^2)} - \frac{7}{3} M_\pi^2 + \frac{4}{9} p^2 \right) \right]
\]

(3.12)
where, again, the factor 15 is a multiplicity factor

\[ \alpha(M^2, p^2) = \sqrt{\frac{M^2}{p^2} - \frac{1}{4}} \tag{3.13} \]

and the trilinear coupling \( g_{V\pi\pi} \) (as given in ref. [3]) is

\[ g_{V\pi\pi} = \frac{g'' x^2}{4 r_V} (1 - z^2) \tag{3.14} \]

Finally, the contributions corresponding to the graphs in Figs. 2e and 2f can be evaluated directly from \( \Pi^{(b)}_{\pi} \) with the substitutions \( M_A \to M_V \) and \( M_V \to M_A \) respectively.

We now introduce counterterms in order to get canonical propagator at the mass of the resonance

\[ \Pi^\text{ren}_i(p^2) = \Pi_i(p^2) + p^2(Z_{3i} - 1) + Z_i \tag{3.15} \]

with \( i = A, V \). The renormalization conditions are

\[
\begin{align*}
\left. \Re \Pi^\text{ren}_i(p^2) \right|_{p^2=M_i^2} &= 0 \\
\lim_{p^2\to M_i^2} \frac{p^2 \big/ M_i^2}{(p^2 - M_i^2 + \Re \Pi^\text{ren}_i(p^2))} &= 1
\end{align*}
\]

(3.16)

where we have taken the real part of \( \Pi^\text{ren}_i \) because the vacuum polarization tensor develops an imaginary part above threshold for the possible decay processes of \( A^3 \) and \( V^3 \). By substituting (3.15) in (3.16) we get

\[ Z_{3i} - 1 = -\frac{d}{dp^2} \left. \Re \Pi_i(p^2) \right|_{p^2=M_i^2} \tag{3.17} \]

\[ Z_i = -\Re \Pi_i(M_i^2) - M_i^2(Z_{3i} - 1) \tag{3.18} \]

We are now able to compute the first two contributions to \( \Pi_{30} \) shown in Fig. 1. In fact, one has only to substitute the expressions for \( \Pi^\text{ren}_A(p^2) \) and \( \Pi^\text{ren}_V(p^2) \) in the following relation (see ref. [3])

\[ \Pi^\text{self}_{30}(p^2) = -x^2 \frac{8\theta}{c_B} p^2 \left( \frac{M_V^2}{p^2 - M_V^2 + \Pi^\text{ren}_V(p^2)} - z^2 \frac{M_A^2}{p^2 - M_A^2 + \Pi^\text{ren}_A(p^2)} \right) \tag{3.19} \]

and then, to use it in eq. (2.7) for the calculation of \( \epsilon_3^\text{self} \). The value of \( \epsilon_3^\text{self} \) predicted by the model is real because \( p^2 = M_Z^2 \) is below threshold. Notice that by considering \( \Pi^\text{ren}_A(p^2) = \Pi^\text{ren}_V(p^2) = 0 \) in (3.19), one recovers, for large \( M_{V,A} \), the tree level value of \( \epsilon_3 \) as given in eq. (2.8).

The contribution to \( \epsilon_3 \) of the remaining graphs of Fig. 1 was calculated in ref. [3]:

\[ \epsilon_3^{\text{loop}} \simeq \frac{g''^2}{16\pi^2} 5 \left( \log \frac{M_A^2}{M_Z^2} - \gamma \right) \left[ 2 - \frac{1}{2} \frac{x^4}{r_A} (1 - z^2)^2 \right. \\
- \frac{x^2}{r_V} \left( 1 - z^2 \right) \left( 1 - z^2 \right) \left( 1 - \frac{r_V}{r_A} \right) - z^2 \frac{x^2}{r_A} \left( 1 - \frac{r_A}{r_V} \right)^2 \right] \\
- \frac{x^2}{r_V} \left( 1 - z^2 \right) \left( 1 - z^2 \right) \left( 1 - \frac{r_V}{r_A} \right) \left( A(M_A^2, M_A^2) + 1 \right) \\
- \frac{x^2}{r_A} \left( 1 - \frac{r_V}{r_A} \right)^2 \left( A(M_Z^2, M_A^2) + 1 \right) \right\} \tag{3.20} \]
where

\[ A(M_\pi^2, M^2) = \frac{M^6 + 9M^4M_\pi^2}{(M^2 - M_\pi^2)^3} \log \frac{M_\pi^2}{M^2} \]

\[ + \frac{1}{6(M^2 - M_\pi^2)^3}(M_\pi^6 - 27M_\pi^4M^2 - 9M_\pi^2M^4 + 35M^6) \]  \hspace{1cm} (3.21)

Furthermore, the numerical analysis has to include electroweak radiative corrections \( (\epsilon_3^{\text{rad}}) \). Our assumption, physically plausible in the presence of new degrees of freedom related to a larger scale, is to take the usual one-loop radiative corrections of the SM with the identification of the Higgs mass as the cut-off \( \Lambda \) which regularizes BESS at high momenta. Finally we have

\[ \epsilon_3 = \epsilon_3^{\text{rad}} + \epsilon_3^{\text{self}} + \epsilon_3^{\text{loop}} \]  \hspace{1cm} (3.22)

4 Numerical results

We now want to evaluate numerically the \( \epsilon_3 \) parameter as predicted by the extended BESS model. In order to reduce the parameter space, we will assume, as we did in ref. [3], the validity of the Weinberg sum rules (WSR) [7]. In this way we can eliminate two parameters, for example \( M_A \) and \( g'' \) in favour of \( M_V \) and \( z \). We get [3]

\[ M_A^2 = \frac{M_V^2}{z} \quad g'' = \frac{2M_V}{v} \sqrt{1 - z} \]  \hspace{1cm} (4.1)

with \( 0 < z < 1 \). Then, the free parameters of our analysis are: \( M_V \), \( z \), \( M_\pi \) and \( \Lambda \).

Using the most recent results from LEP experiments combined with the low energy weak data [8]

\[ \epsilon_3 = (1.3 \pm 3.1) \times 10^{-3} \]  \hspace{1cm} (4.2)

we can derive the bounds on the parameter space of the extended BESS model coming from the estimate of the corrections to \( \epsilon_3 \).

Concerning the electroweak radiative corrections, following ref. [8], we use

\[ \epsilon_3^{\text{rad}} = 0.00696 \]  \hspace{1cm} (4.3)

corresponding to \( m_{\text{top}} = 150 \, \text{GeV} \) and \( M_H = \Lambda = 1.5 \, \text{TeV} \).

In Fig. 3 we give the 90\% C.L. allowed region in the plane \( (z, M_V) \) for two values of \( M_\pi/\Lambda \) and \( \Lambda = 1.5 \, \text{TeV} \). The effects of including the contribution coming from the \( V \) and \( A \) self-energies are generally small and go in the direction of enlarging the allowed region. This result is in accordance with the observations done in ref. [3] that the \( V \) and \( A \) self-energy contributions are generally negative and decrease the value of \( \epsilon_3^{\text{tree}} \) of the tree level. The corrections due to the self-energies are more sizeable in the region close to \( z = 1 \). In fact, by using the WSR, the expression of \( \epsilon_3^{\text{loop}} \) in eq. (3.20) vanishes in the limit \( z \to 1 \). The reason for this is twofold. First, the contributions from the three graphs with PGB loops (see Fig. 1) cancel among themselves, second, the WSR imply that for \( z = 1 \) also \( r_V = r_A \) corresponding to a complete degeneration among \( V \) and \( A \) resonances (notice also that \( \epsilon_3^{\text{tree}} = 0 \) for \( z = 1 \)). This latter property is not true anymore for the
complete expression of $\epsilon_3$ since the $V$ and $A$ resonances get different corrections from the self-energies, also at $z = 1$.

In Fig. 4 we give the 90% C.L. allowed region in the plane $(M_V, M_\pi)$ for two values of $z$ and $\Lambda = 1.5 \, TeV$. We notice that for $M_\pi > 100 \, GeV$ the bounds do not depend very much on the mass of the PGB and that the allowed region shrinks for increasing values of $z$.

Finally, for completeness, in Fig. 5 we also give the 90% C.L. allowed region in the plane $(z, M_\pi)$ for two values of $M_V$ and $\Lambda = 1.5 \, TeV$.

From Figs. 3 and 4 we see that there is an absolute lower bound on $M_V$, independent on $z$ and $M_\pi$, of about $900 \, GeV$.

In conclusion the calculation we have done here shows that the approximation used in ref. [3] for the evaluation of the self-energies of the new resonances was essentially correct. However it should be noticed that a sizeable difference arises in the region close to $z = 1$ where, as explained before, the different one-loop contributions for the vector and the axial-vector resonances forbid the cancellation arising from the other diagrams.

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Figure Captions

**Fig. 1** Graphs contributing to $\epsilon_3$ at one-loop level within the extended BESS model.

**Fig. 2** $A^3$ and $V^3$ self-energy contributions.

**Fig. 3** 90% C.L. allowed regions in the plane $(z, M_V)$ from $\epsilon_3$ for $\Lambda = 1.5$ $TeV$. The solid (dashed) line is the lower bound coming from the total one-loop effect for $M_\pi/\Lambda = 0.10$ ($M_\pi/\Lambda = 0.35$).

**Fig. 4** 90% C.L. allowed regions in the plane $(M_V, M_\pi)$ from $\epsilon_3$ for $\Lambda = 1.5$ $TeV$. The solid (dashed) line is the bound coming from the total one-loop effect for $z = 0.2$ ($z = 0.8$). The allowed regions lie on the right of the curves.

**Fig. 5** 90% C.L. allowed regions in the plane $(z, M_\pi)$ from $\epsilon_3$ for $\Lambda = 1.5$ $TeV$. The solid (dashed) line is the bound coming from the total one-loop effect for $M_V = 1$ $TeV$ ($M_V = 1.5$ $TeV$). The allowed regions lie on the left of the curves.
This figure "fig1-1.png" is available in "png" format from:

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