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Testing an anomalous top quark via precision electroweak measurements

Maurizio Frigeni

Dipartimento di Fisica, Università di Pisa

and

Riccardo Rattazzi¹

CERN, Geneva

The effect on electroweak corrections of anomalous interactions of the top quark with the Z^0 is studied. Such anomalous terms are expected to arise in the context of dynamical symmetry breaking. It is shown that, depending on m_t , the whole set of radiative-correction parameters can be relevant. This set consists of three observables coming from the vacuum polarization graphs plus the correction to the $Z^0 b\bar{b}$ vertex. The recent data coming from collider experiments are used to place bounds on the anomalous terms.

The same observables are evaluated in a simple renormalizable model in which the top quark mixes with an isosinglet up-type quark.

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¹ A. Della Riccia fellow

1. It is well known that precision measurements at LEP can open a window on physics above the Fermi scale. The most relevant effects due to virtual heavy particles come typically through corrections to vector boson propagators, the so-called “oblique” corrections [1]. An exception to this rule is represented by the top quark, since large $G_F m_t^2$ terms appear in the $Z\bar{b}b$ vertex corrections [2]. However, this is because b and t fit into the same weak doublet. Any possible anomalous behaviour of the top quark should show up both in oblique and vertex corrections.

The purpose of this letter is to study the impact on precision measurements of anomalous non-universal interactions of the top quark with the Z^0 . We shall deal with two models where such extra couplings arise in completely different ways: in one case they are induced by dynamical gauge symmetry breaking [3], while in the other they simply appear via the mixing of the top quark with a heavy isosinglet up-type quark [4].

Inspired by an analogy to low-energy chiral symmetry breaking, it has been argued in Ref. [3] that, if the electroweak symmetry breaking is dynamical, residual interactions between fermion and vector bosons should in general arise. One could also argue that such anomalous couplings depend somehow on the mass of the involved fermions. According to this picture one expects very small deviations for the light fermions (consistent with experimental observations). Conversely, large deviations are expected for the heavy top quark. These could give measurable effects in precision tests at LEP. In Ref. [5] the contribution to the ρ parameter and to the $Z\bar{b}b$ vertex were studied. The aim of this paper is to study the whole set of radiative-correction parameters, including ϵ_2 and ϵ_3 defined in Ref. [6]. A comparison is made with the latest LEP data, thus updating the analysis in Ref. [5].

Of course this model is only phenomenological and renormalizability is lost. In order to make a comparison with a renormalizable model, we consider the simplest possible case in which anomalous couplings of the top are present, consisting of the addition of a heavy isosinglet up-type quark. This choice is also of speculative interest. In fact, many grand-unified models require the addition of

isosinglet fermions. Furthermore, the addition of such an isosinglet characterizes many models which try to explain the heaviness of the top quark [7].

This paper is organized as follows. In Sec. 2 we recall the definition of the three parameters $\epsilon_{1,2,3}$ which determine the pattern of oblique radiative corrections in Z^0 physics [6]. In Sec. 3 we calculate these corrections in the dynamical-breaking model, and make a comparison with the recent experimental data coming from the four LEP experiments. In Sec. 4 we comment on the radiative effects in the case of the isosinglet. Finally, Sec. 5 is devoted to our conclusions.

2. The effect of the oblique corrections can be seen as a small modification of the tree level relations among the fundamental parameters of the electroweak theory. In general this effect can be described [6] with the introduction of three observables. One of these, $\Delta\rho$, determines the relative strength of charged and neutral current four-fermion interactions. The other two, Δr_W and $\Delta k'$, describe the relative difference among the three possible definitions of the sine of the Weinberg angle: s_W , s_o and $\bar{s}_W$. More precisely, Δr_W is defined by

$$s_W^2 c_W^2 = (1 - \frac{M_W^2}{M_Z^2}) \frac{M_W^2}{M_Z^2} = \frac{\pi\alpha(M_Z)}{\sqrt{2}G_F M_Z^2(1 - \Delta r_W)} = \frac{s_o^2 c_o^2}{1 - \Delta r_W} \quad (2.1)$$

where $\alpha(M_Z) = 128.8$ is the running fine structure constant evaluated at $s = M_Z^2$ and $G_F = 1.166 \times 10^{-5} [\text{GeV}]^{-2}$ is the Fermi constant as measured in μ -decay. The other two observables are given through the on-shell couplings g_V and g_A of the Z^0 to fermions

$$g_A = T_L^3(1 + \frac{\Delta\rho}{2}) \quad (2.2a)$$

$$\frac{g_V}{g_A} = 1 + 4Q_f \bar{s}_W^2 = 1 + 4Q_f(1 + \Delta k')s_o^2. \quad (2.2b)$$

The contribution of heavy particles to these observables via vacuum polarization graphs is quite simple. We can write the vacuum polarization tensor as

$$\Pi_{ij}^{\mu\nu} = -ig^{\mu\nu}(A_{ij} + q^2 F_{ij}(q^2)) \quad (2.3)$$

where i, j stand for W, Z, γ . For heavy particles we can approximate the functions $F_{ij}(q^2)$ with $F_{ij}(0) = F_{ij}$. Considering that $A_{\gamma\gamma} = A_{Z\gamma} = 0$, one is left

with 6 constants A_{WW} , F_{WW} , A_{ZZ} , F_{ZZ} , $F_{Z\gamma}$, $F_{\gamma\gamma}$. Three linear independent combinations of these amount to a simple redefinition of α , G_F and M_Z [8]. The combinations

$$\epsilon_1 = \frac{A_{ZZ}}{M_Z^2} - \frac{A_{WW}}{M_W^2} = \frac{A_{33} - A_{WW}}{M_W^2} \quad (2.4a)$$

$$\epsilon_2 = F_{WW} - F_{33} \quad (2.4b)$$

$$\epsilon_3 = \frac{F_{3\gamma}}{s_W} - F_{33} = \frac{c_W}{s_W} F_{03} \quad (2.4c)$$

affect the physical observables as

$$\Delta\rho = \epsilon_1 \quad (2.5a)$$

$$\Delta k' = \frac{-c_W^2 \epsilon_1 + \epsilon_3}{c_W^2 - s_W^2} \quad (2.5b)$$

$$\Delta r_W = -\frac{c_W^2}{s_W^2} \epsilon_1 + \left(\frac{c_W^2}{s_W^2} - 1\right) \epsilon_2 + 2\epsilon_3 . \quad (2.5c)$$

In the standard model the leading contribution is due to a heavy top quark, and resides in $\epsilon_1 \simeq 3G_F m_t^2 / 8\sqrt{2}\pi^2$. No large contribution arises in ϵ_2 and ϵ_3 . This picture is not essentially altered by supersymmetry [8]. On the other hand one expects large contributions to both ϵ_1 and ϵ_3 in technicolor models [9], while ϵ_2 is still less relevant.

A last comment on eq.(2.2) is in order. In general $\Delta\rho$ and $\Delta k'$ receive non-universal corrections due to light known particles and arising from box, vertex diagrams and imaginary parts of the propagators. In the standard model, these are small and reasonably well-known quantities, with the exception of the $Z\bar{b}b$ vertex correction. In this case the left-handed $Z\bar{b}b$ vertex receives a correction proportional to m_t^2 , so that (2.2) is modified by

$$g_A^b = T_L^3 \left[1 + \frac{1}{2} (\Delta\rho + \Delta\rho_{\text{vertex}}^b) \right] \quad (2.6a)$$

$$\frac{g_V^b}{g_A^b} = 1 + 4Q_f \bar{s}_W^2 = 1 + 4Q_f \left(1 + \Delta k' - \frac{1}{2} \Delta\rho_{\text{vertex}}^b \right) s_o^2 . \quad (2.2b)$$

In the standard model one has $\Delta\rho_{\text{vertex}}^b = -(\sqrt{2}G_F/4\pi^2)m_t^2$.

3. In Ref. [3] it has been shown that, if the symmetry breakdown in the electroweak theory is dynamical, residual interactions between fermions and Goldstone bosons should in general arise. These extra terms can be shown to give non-universal interactions between the $SU(2) \times U(1)$ gauge fields and the fermions. These interactions are the analogue of the extra couplings of the pions to nucleons which arise in QCD when $SU(2) \times SU(2)$ is dynamically broken. In QCD the extra couplings are $O(1)$ (for instance the $\pi \bar{N} N$ form factor at zero momentum is characterized by an extra piece $k = g_A - 1 \simeq 0.25$). This can be seen as a consequence of the fact that in QCD there is only one mass scale Λ_{QCD} . For the electroweak theory the situation is rather different, since there is a full hierarchy of mass scales in the fermion spectrum ranging from $m_e \simeq 0.5 \text{ MeV}$ to $m_t \gtrsim 80 \text{ GeV}$ [10]. One could argue [3, 5] that, if fermions participate in the symmetry breakdown in a non-trivial fashion, the structure of the mass spectrum may in some way be reflected in the extra couplings to gauge bosons. Indeed by assuming that the anomalous terms are of order $m_f \sqrt{G_F}$ for a fermion f , consistency is met with the experimental data on light fermions [3]. For the top quark, however, since it has not yet been detected, we cannot put direct bounds on its extra couplings. On the other hand, given the present lower bound on m_t and supposing $k_{\text{top}} \sim m_t \sqrt{G_F}$, one could expect large deviations for the top. It is then natural, in order to study the effects of these anomalous couplings, to neglect any term (*) apart from those affecting the $Z^0 \bar{t} t$ vertex. We can then write the effective $Z^0 \bar{t} t$ Lagrangian as

$$\mathcal{L}_{Z\bar{t}t} = \frac{g}{2c_W} \bar{t} \gamma_\mu \{ P_L (1 + k_L) + P_R k_R - \frac{8}{3} s_W^2 \} t Z^\mu \quad (3.1)$$

where $P_{L,R}$ are chirality projectors and $k_{L,R}$ parametrize the extra interaction. The Lagrangian with a modified vertex (3.1) is non-renormalizable. As a matter of fact, radiative corrections coming from (3.1) bring a physically meaningful cut-off dependence. However, since eq.(3.1) is deduced from a gauge-invariant non-linear chiral Lagrangian by going to the unitary gauge [3], a gauge-invariant

(*) For example a simple form $\sim \sqrt{m_i m_j G_F}$, for the anomalous couplings involving two fermions i, j and a vector boson, is consistent with our choice.

cut-off must be chosen. It is then useful to make the calculation in dimensional regularization and interpret $1/\epsilon$ as $\ln \Lambda$ [11]. Λ represents the typical scale of the interactions which break the gauge group.

The expressions of $\epsilon_{1,2,3}$ can be readily found as functions of k_L and k_R

$$\epsilon_1 = \frac{3G_F m_t^2}{8\sqrt{2}\pi^2} \left\{ 1 + 2[2(k_R - k_L) - (k_R - k_L)^2] \ln \frac{\Lambda^2}{m_t^2} \right\} \quad (3.2a)$$

$$\epsilon_2 = -\frac{G_F M_W^2}{\sqrt{2}\pi^2} \left\{ \frac{k_L}{2} + \frac{1}{4}(k_L^2 + k_R^2) \right\} \ln \frac{\Lambda^2}{m_t^2} \quad (3.2b)$$

$$\epsilon_3 = -\frac{G_F M_W^2}{\sqrt{2}\pi^2} \left\{ \frac{k_L}{6} - \frac{k_R}{3} + \frac{1}{4}(k_L^2 + k_R^2) \right\} \ln \frac{\Lambda^2}{m_t^2} \quad (3.2c)$$

In deriving these expressions, only terms proportional to $1/\epsilon$ were retained. Finite anomalous contributions must be discarded in this cut-off dependent model. The expression for ϵ_1 was already given in Ref. [5], and corresponds to the standard contribution with the addition of terms proportional to the axial anomalous coupling $k_A = k_L - k_R$. Notice that for $k_A = 2$ formula (3.2a) goes back to the standard model result. In fact $k_A = 2$ corresponds to a simple change of sign of the standard axial current interaction. A relevant phenomenological feature of eq.(3.2a) is that the anomalous contribution can take negative values. This means that larger values of m_t than in the standard case can be compatible with the experimental bound on $\Delta\rho$.

The standard contribution due to the top quark has not been added to ϵ_2 and ϵ_3 in eq.(3.2bc) since it is not relevant compared to the present experimental precision. Notice that, ϵ_2 and ϵ_3 being derivative quantities, they receive contributions from both the vector k_V and axial k_A combinations of couplings. Moreover, they depend only logarithmically on m_t , so that the results we get from their comparison with the experimental data will be less affected by the (unknown) value of the top mass.

We recall the expression for $\Delta\rho_{\text{vertex}}^b$ obtained from this model [5], inclusive of the standard model leading contribution

$$\Delta\rho_{\text{vertex}}^b = -\frac{\sqrt{2}G_F m_t^2}{4\pi^2} \left\{ 1 + (k_R/2 - 2k_L) \ln \frac{\Lambda^2}{m_t^2} \right\} \quad (3.2d)$$

Notice that the $O(m_t^2)$ contributions to the vertex and to ϵ_1 are no longer proportional to each other (in the standard case $\Delta\rho_{\text{vertex}}^b = -4/3\Delta\rho$).

As m_t grows, and at fixed values of the k 's, ϵ_1 and (or) $\Delta\rho_{\text{vertex}}$ become the most relevant effect. However, for m_t not very large ($m_t \lesssim 200 \text{ GeV}$) also the effect of $\epsilon_{2,3}$ is relevant. When making a comparison with the experimental data all the observables must then be taken into account. In what follows we compare our formulae to the present data coming from precision measurements in the four LEP experiments. The data on the leptonic partial widths of the Z^0 and on the leptonic F-B-asymmetries yield for the the effective Z^0 couplings [12](*)

$$\begin{aligned} g_A &= \frac{1}{2} + \frac{1}{4}\Delta\rho = 0.4998 \pm 0.0017 \\ \frac{g_V}{g_A} &= 0.074 \pm 0.012 \end{aligned} \tag{3.3}$$

where a mean has been taken over the four LEP experiments [14]. The error in (3.3) is mainly statistical. Using the value $M_Z = 91.177 \text{ GeV}$, one has from eq.(2.1) $s_o^2 = 0.2313$, and then from eq. (2.2)

$$\Delta k' = 0.001 \pm 0.013 . \tag{3.4}$$

In addition, the value of M_W/M_Z measured by CDF and UA2 [14] yields

$$\Delta r_W = -0.015 \pm 0.018 . \tag{3.5}$$

A direct application of eq.(2.5) gives the following bounds for $\epsilon_{1,2,3}$

$$\begin{aligned} \epsilon_1 &= -0.0008 \pm 0.0068 \\ \epsilon_2 &= -0.008 \pm 0.011 \\ \epsilon_3 &= 0.000 \pm 0.009 \end{aligned} \tag{3.6}$$

(*) The same data can be found in Ref. [13]. In that Ref., however, due to a different procedure for combining statistical and systematic errors, the quoted value for g_A is: $g_A = 0.4998 \pm 0.0012$

where the errors on $\Delta\rho$, $\Delta k'$, Δr_W have been propagated quadratically. A comparison between eq.(3.2) and (3.6) bounds k_L and k_R .^(†)

An additional bound is provided by the experimental value of the ratio between hadronic and leptonic width [13]

$$\frac{\Gamma_h}{\Gamma_l} = 20.94 \pm 0.12 \quad (3.7)$$

whose expression in terms of our observables is given by

$$R = \frac{\Gamma_h}{\Gamma_l} = 6 \frac{(5 - \frac{28}{3}s_o^2 + \frac{88}{9}s_o^4)}{1 + (1 - 4s_o^2)^2} \left(1 + \frac{\alpha_s}{\pi}\right) (1 + 0.26\epsilon_1 - 0.34\epsilon_3 + 0.25\Delta\rho_{\text{vertex}}^b) \quad (3.8)$$

where the numerical coefficients were computed using $s_o^2 = 0.23$, and α_s is the strong coupling constant.

We show in figures 1 and 2 which part of the (k_L, k_R) plane is allowed by the above data, for different values of m_t and the cut-off Λ . In doing that we doubled the errors given above for the experimental quantities. Thus we may roughly consider our bounds to be at the 2-sigma level (part of the error is systematic). In order to take into account the uncertainties related to QCD corrections, we also considered $\alpha_s = 0.12 \pm 0.01$ [16].

The bounds coming from $\Delta\rho$ and R were already shown in [5]. There, however, older LEP data were used and ϵ_2 , ϵ_3 were not taken into account. This is surely correct for large enough values of the top mass. However, for moderately large values of m_t (≤ 200 GeV), we expect the contributions from ϵ_2 and ϵ_3 to be of the same order. As a matter of fact we see in fig.(1) ($m_t = 150$ GeV) that the use of all the observables implies stronger bounds than the simple use of R and ϵ_1 . For $m_t = 120$ GeV, see fig.(2), the bound from R becomes irrelevant. This is due to the fact that, for negative values $k_L \sim k_R < 0$, the contribution

(†) It should be noticed that, in the interplay between the experimental data (3.3)-(3.5) and eq.(3.6) some corrections due to known effects (such as those due to box and vertex diagrams) should be taken into account. However, a direct evaluation of these terms by using the formulae in Ref. [15] shows that they are not larger than the errors quoted in eq.(3.6).

from $\Delta\rho_{\text{vertex}}^b$ in eq.(3.8) is compensated by that of ϵ_3 . At $m_t = 250$ GeV the combined bounds from ϵ_1 and R are more restrictive. We do not report the graph in this case. By choosing $\Lambda = 3$ TeV, the allowed region reduces to a strip with $-0.4 < k_L + k_R < 1.2$, $-0.06 < k_R - k_L < -0.02$.

Anyway, in a phenomenological model like the one at hand, one should take into account all the available information. It is well conceivable, for instance, that there exist small additional couplings not related to the top quark and changing the above picture. One example is an anomalous $Z\bar{b}b$ vertex, which could give an extra contribution to R , naively of the same order as that coming from the top in eq(3.2d), if one assumes $k_b \sim m_b\sqrt{G_F}$. On the other hand, according to the same pattern, very small additional effects are expected for the light fermions e and μ . This observation increases the importance of ϵ_3 , since it is bounded through purely leptonic quantities ($\Delta k'$ in eq. (3.4)).

The use of ϵ_2 and ϵ_3 is relevant also in view of the foreseen precision in the determination of $\Delta k'$ and Δr_W . (It should be noticed that, conversely, the precision on $\Delta\rho$ is not expected to increase noticeably.) By combining the above quantities the ϵ_1 contribution can be cancelled

$$(c^2 - s^2)\Delta k' - s^2\Delta r_W = (c^2 - s^2)(\epsilon_3 - \epsilon_2) = \epsilon_- . \quad (3.9)$$

The ϵ_- observable is proportional to the vector combination $k_L + k_R$, since it is determined by $F_{3\gamma}$ in eq. (2.4) (in our model there are no anomalous contributions to F_{WW}). The measure of the left-right asymmetry with an error $\sigma(A_{LR}) = 0.003$ [6] allows a determination of $\Delta k'$ with $\sigma(\Delta k') = 0.0016$. The measure of M_W at LEP II [17] with $\sigma(M_W) = 100$ MeV corresponds to $\sigma(\Delta r_W) = 0.005$. Thus the resulting error on ϵ_- will allow us to bound $k_L + k_R$ at the 2-sigma level as $E < k_L + k_R < E + 1.2$, where E depends on the experimental value of ϵ_- ($\Lambda = 3$ TeV and $m_t = 150$ GeV have been assumed). This bound is quantitatively similar to the one that can already be obtained for $m_t = 250$ GeV by using the data on R ; it is, however, interesting since it is rather independent on m_t .

4. In this section we examine briefly the effects on precision measurements due to the presence of an isosinglet up quark.

Indicating by g'_L and g'_R the two components of this fermion and supposing that it mixes essentially with the top [4], the Yukawa interactions are given by

$$\mathcal{L} = h\bar{q}'_L H t'_R + m_1\bar{g}'_L g'_R + m_2\bar{g}'_L t'_R \quad (4.1)$$

where H is the Higgs field, m_1 and m_2 are bare mass terms and h is a Yukawa coupling. The fields q'_L and t'_R are the top-bottom left gauge doublet and the top right singlet. The mass eigenstates t_L and g_L arising from the mass matrix after S.S.B are related to t'_L and g'_L via a mixing angle ξ_L (*). Due to this mixing, the coupling of the top t_L to the Z^0 is non-standard. In addition there is a flavour-violating vertex $Z^0 t g$ [4].

The net contribution to $\Delta\rho$ is then [4]

$$\Delta\rho = \frac{3G_F}{8\sqrt{2}\pi^2} \left\{ \cos^4 \xi_L m_t^2 + \sin^4 \xi_L m_g^2 + 2 \sin^2 \xi_L \cos^2 \xi_L \frac{m_t^2 m_g^2}{m_g^2 - m_t^2} \ln \frac{m_g^2}{m_t^2} \right\} \quad (4.2)$$

which, in contrast to eq.(3.2), is always positive. In addition eq.(4.2) always strengthens the bound on m_t implied by the standard expression. In fact, it can be easily verified that the expression in curly brackets in (4.2) is larger than $\min[m_t^2, m_g^2]$, for any value of ξ_L . This property is shared by the more general formula for $\Delta\rho$ given in Ref. [4], which also takes into account a possible mixing with the lighter up quarks.

It is interesting to see whether $\Delta\rho_{\text{vertex}}^b$ evaluated in this model turns out proportional to $\Delta\rho$. The calculation follows the same lines as in the standard case. The leading contributions, proportional to the squared fermion masses, arise from the triangle graphs in which a pseudo-Goldstone boson is propagated. The result mimics the standard one

$$\Delta\rho_{\text{vertex}}^b = -\frac{4}{3}\Delta\rho \quad (4.3)$$

where $\Delta\rho$ is given in eq.(4.2). Thus the experimental measurement of $\Gamma_{\text{had}}/\Gamma_{\text{lep}}$ cannot be used to get an additional bound on the parameters of this model.

(*) Another mixing angle ξ_R appears for the right-handed components. The value of this parameter is fixed by ξ_L and by the mass eigenvalues m_t , m_g .

The contribution to ϵ_2 and ϵ_3 is

$$\begin{aligned}\epsilon_2 &= \frac{G_F M_W^2}{4\sqrt{2}\pi^2} \left\{ \ln \frac{m_t^2}{M_Z^2} - \frac{1}{2} + \sin^2 \xi_L \ln \frac{m_g^2}{m_t^2} - \cos^2 \xi_L \sin^2 \xi_L H(m_t^2/m_g^2) \right\} \\ \epsilon_3 &= \frac{G_F M_W^2}{12\sqrt{2}\pi^2} \left\{ \ln \frac{m_t^2}{M_Z^2} + \frac{1}{6} + \sin^2 \xi_L \ln \frac{m_g^2}{m_t^2} - 3 \cos^2 \xi_L \sin^2 \xi_L H(m_t^2/m_g^2) \right\}\end{aligned}\quad (4.4)$$

where the function $H(x)$ is given by

$$H(x) = \frac{1}{(x-1)^2} \left\{ 4x + \frac{x+1}{x-1} (1+x^2-4x) \ln x \right\} - \frac{5}{3} \quad (4.5)$$

In analogy with the standard model case, the terms in eq.(4.4) are very small.

5. We have studied the electroweak radiative effects of anomalous top quark couplings to the Z^0 boson. Such anomalous terms are expected to arise in the context of dynamical gauge symmetry breaking. It turns out that, for moderately large values of the top quark mass $m_t \lesssim 150$ GeV, the whole set of observables, consisting of ϵ_1 , ϵ_2 , ϵ_3 and $\Delta\rho_{\text{vertex}}^b$, must be considered. In this range of m_t the bounds on the anomalous couplings k_L and k_R coming from the experimental data are essentially determined by $\epsilon_{1,2,3}$. These bounds are roughly given by $-2 < k_V < 2$, $-0.15 < k_A < 0.05$. The foreseen measurements with polarized beams and at LEP II will permit to reduce the allowed range in k_V of about a factor 3 ($\Delta k_V \sim 1$). For larger values of the top mass ($m_t \gtrsim 200$ GeV), the bounds on k_L and k_R given by the experimental values of $\Delta\rho$ and $R = \Gamma_{\text{had}}/\Gamma_{\text{lep}}$ are more significant than those coming from ϵ_2 and ϵ_3 . It must however be noticed that the case of a very heavy top quark ($m_t \gtrsim 250$ GeV), even if allowed in principle by equations (3.2a) and (3.2b), is not very plausible since, in order to be consistent with the experimental data, a considerable degree of fine tuning must be exercised on the anomalous k_L and k_R .

We have also studied the whole set of radiative observables in a simple renormalizable model in which the top quark is mixed with an isosinglet quark. It turns out that the pattern is analogous to the standard one in the improved Born approximation. The contribution to $\Delta\rho$ and the leading contribution to the $Z^0 \bar{b}b$ vertex correction are proportional to the same expression. The parameters ϵ_2 and ϵ_3 remain irrelevant.

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Figure captions

Fig. 1: Regions in the (k_L, k_R) plane allowed by present-day LEP data, when compared to the computed values of ϵ_1 (two dotted stripes centred respectively around $k_A = 0$ and $k_A = 2$), R (region delimited by the dashed lines), ϵ_2 (internal part of the lower standing circle) and ϵ_3 (internal part of the upper standing circle). The values $m_{top} = 150$ GeV and $\Lambda = 1$ TeV have been assumed for the top quark mass and the cut-off scale respectively.

Fig. 2: The same as Fig. 1, for $m_{top} = 120$ GeV and $\Lambda = 1$ TeV.

Fig. 1

$m_{\text{top}} = 120 \text{ GeV}, \quad \Lambda = 1 \text{ TeV}$

Fig. 2