VACUUM POLARIZATION EFFECTS OF NEW
PHYSICS ON ELECTROWEAK PROCESSES

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ABSTRACT

The effects of new physics on vacuum polarization corrections to electroweak processes can in general be parametrized in terms of six real constants, in the limit of neglecting terms which vanish for $\Lambda \rightarrow \infty$, with Λ being the scale of new physics. Three of these parameters can be reabsorbed in the definitions of α , G_F and m_Z , while the remaining three are observables. On the basis of simple models, it is shown that all three can be important and should be kept in a model-independent analysis of the data. This is equivalent to treating Δr , $\Delta\rho$ and $\bar{s}w^2$ (an effective $\sin^2\theta_W$ for on-shell Z couplings) as independent observables. By now available data allow a separate model-independent determination of these quantities. Implications for technicolour and supersymmetric models and anomalous γWW or ZWW couplings are discussed.

1 It is well known that the radiative corrections to electroweak processes involving light quarks and leptons as external states are sensitive to virtual effects of heavy particles. In general, heavy particles do not decouple, for example because of the longitudinal gauge boson modes whose couplings increase with masses. The upper limit on the top quark mass m_t derived from the presence [1] of large terms of order $G_F m_t^2$ in the electroweak corrections arises from this phenomenon and is at present its most interesting application. The effect of virtual heavy particles is typically concentrated in the so-called "oblique" corrections [2] arising from heavy loops in the gauge boson vacuum polarisation diagrams. Actually the top quark is exceptional in this respect, because large $G_F m_t^2$ terms are also present [3] in the $Z \rightarrow b\bar{b}$ vertex corrections. However, these vertex corrections are only there because the top is the main partner of the b quark in the weak charged current. Large virtual effects of heavy particles from new physics beyond the standard model are thus even more expected to be mainly contained in the oblique corrections.

The quadratic terms in m_t of order $G_F m_t^2$ only contribute to the first term in an expansion of the vacuum polarisation amplitudes $\Pi(q^2)$ in powers of q^2 , while the successive terms are down by relative powers of q^2/m_t^2 . In general, for new physics arising at a scale Λ , one can safely neglect higher terms in the expansion, as they are suppressed by increasing powers of q^2/Λ^2 , with $q^2 \leq m_Z^2$ for physics at LEP and at lower energy scales. If only the non-vanishing terms at $q^2=0$ are kept in the expansion, then one obtains the well-known result [2],[4] that the whole effect of oblique corrections can be described by a single quantity $\Delta\rho$. However these constant terms are not of order $G_F \Lambda^2$ because they must vanish in the limit of no gauge and isospin symmetry breaking (for example, in the t quark case, the relevant energy scale is the t-b mass splitting). In all sensible models the symmetry breaking scale is much less than Λ . As a consequence the next order terms in the expansion, which by dimension do not vanish in the large Λ limit (i.e. are typically of order $G_F m_W^2$) can be numerically important. For example it has recently been pointed out [5-8] that such terms can indeed be sizeable in technicolour theories [9] where a second quantity in addition to $\Delta\rho$ is introduced (denoted by αS in Ref. 5 while $\Delta\rho$ is called αT).

In this note we discuss the general structure of the oblique corrections under the assumption that only constant and linear

terms in q^2 are kept in the expansion of the vacuum polarisation amplitudes. As we shall see, in a model independent approach, the effect of oblique corrections for precision experiments at energies up to m_Z is to be described, at this level of accuracy, in terms of three independent parameters. One can conveniently choose Δr_W (relating m_W/m_Z to m_Z through $\alpha(m_Z)$ and G_F), $\Delta\rho$ (entering in the ratio of neutral to charged current four-fermion interactions) and $\Delta k'$ (describing the relative difference between the effective $\sin^2\theta_W$ for the Z couplings on-shell and the weak mixing angle $\sin^2\theta_W^0$ without weak radiative corrections). In the standard model Δr_W , $\Delta\rho$ and $\Delta k'$ are related, as far as the large terms of order $G_F m_t^2$ are concerned, by (see e.g. [10]):

$$\Delta r_W = (c^2 - s^2)/s^2 \Delta k' = -c^2/s^2 \Delta\rho \quad (1)$$

where $s^2 = 1 - c^2 = \sin^2\theta_W$. In general, however, these three parameters do not receive oblique contributions which are proportional to each other and we will produce examples of models where all three are linearly independent. The case where the oblique corrections can be described in terms of only two parameters [5-8] is quite plausible but not really general. The status of a separate independent determination of Δr_W , $\Delta\rho$ and $\Delta k'$ from the existing data is also presented. The impact of this determination on technicolour [9] or supersymmetric [11],[12] models is briefly illustrated.

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Given the fine structure constant α , the Fermi constant G_F as measured from μ -decay and the Z mass m_Z (which together with m_t and m_H are taken as the basic input parameters of the electroweak theory) the observables Δr_W , $\Delta\rho$ and $\Delta k'$ are defined in the following way. Δr_W is equivalent to m_W and is given by the relation

$$(1 - m_W^2/m_Z^2) m_W^2/m_Z^2 = \pi \alpha(m_Z) / (\sqrt{2} G_F m_Z^2 (1 - \Delta r_W)) \quad (2)$$

with $\alpha(m_Z) = \alpha/(1 - \Delta\alpha) = 1/128.8$ [10], so that the effect of the running of α (due to known physics) has been extracted out of $(1 - \Delta r) = (1 - \Delta\alpha)(1 - \Delta r_W)$. Note that

$$s_W^2 = 1 - m_W^2/m_Z^2 \quad (3)$$

is one commonly adopted definition of $\sin^2\theta_W$ [13]. In order to define $\Delta\rho$ and $\Delta k'$ we first introduce g_V and g_A , the vector and axial vector couplings of the on-shell Z to charged leptons (assuming $e-\mu-\tau$ universality), which are measured from the Z partial widths and the asymmetries at $q^2=m_Z^2$. The precise definitions of g_V and g_A are specified by the relations (taken as exact):

$$G(Z \rightarrow \bar{l}l) = G_F m_Z^3 (g_V^2 + g_A^2) / (6\pi\sqrt{2}) \quad (4)$$

$$A_{FB}^l(q^2 = m_Z^2) = 3g_V^2 g_A^2 / (g_V^2 + g_A^2)^2$$

In terms of g_V and g_A , $\Delta\rho$ and $\Delta k'$ are defined by:

$$g_A = -\sqrt{\rho}/2 = -(1 + \Delta\rho/2)/2 \quad (5)$$

$$g_V/g_A = 1 - 4\bar{s}_W^2 = 1 - 4(1 + \Delta k')s_0^2$$

Here $\bar{s}_W^2$ is the effective $\sin^2\theta_W$ for on-shell Z [2],[10],[14] while $s_0^2 = \sin^2\theta_W^0$ is defined by:

$$s_0^2 c_0^2 = \pi\alpha(m_Z)/(\sqrt{2}G_F m_Z^2) \quad (6)$$

(The reader should pay attention not to confuse s_W^2 , $\bar{s}_W^2$ and s_0^2).

Strictly speaking one should introduce a pair of $\Delta\rho$ and $\Delta k'$ for each flavour of fermions. In the standard model $\Delta\rho$ and $\Delta k'$ can be defined as universal quantities only if small terms, for example arising from box diagrams, vertex corrections, and imaginary parts are neglected [10]. We work in this approximation and consider oblique corrections which are larger than such terms.

Of course, parameters closely related to Δr_W , $\Delta\rho$ and $\Delta k'$ are commonly introduced in the literature [2],[13],[14] as auxiliary tools to approximately describe the radiative corrections in the standard model. The point here is that they are also necessary and sufficient to characterise in a general way the non-vanishing contributions (as Λ goes to infinity) to the electroweak parameters arising from vacuum polarisation corrections due to some new physics at the scale Λ .

Let us assume that such new physics effects are concentrated in vacuum polarisation amplitudes, so that these terms must be well defined, gauge invariant and observable. In general we have:

$$\Pi_{ij,\mu\nu}(q^2) = -ig_{\mu\nu}[A_{ij} + q^2 F_{ij}(q^2)] + q_\mu q_\nu \text{ terms} \quad (7)$$

where i,j stand for W,Z,γ . We now assume that only $F_{ij}(0)$ for each i,j is retained in an expansion in q^2 . Taking into account that $\Pi_{\gamma\gamma}(0) = \Pi_{ZZ}(0) = 0$, one is left with a total of six independent constants: A_{WW} , F_{WW} , A_{ZZ} , F_{ZZ} , $F_{Z\gamma}$, where $A_{WW} = A_{WW}(0)$ etc. and the label $q^2 = 0$ has been omitted. These constants are taken as real numbers because there are no thresholds associated with the new physics at scales $q^2 \ll m_Z^2$ so much smaller than Λ^2 .

It is simple to derive the results [12]:

$$\delta\alpha/\alpha = -F_{\gamma\gamma} \quad (8)$$

$$\delta G_F/G_F = A_{WW}/m_W^2 \quad (9)$$

$$\delta m_Z^2/m_Z^2 = -A_{ZZ}/m_Z^2 - F_{ZZ} \quad (10)$$

where $\delta\alpha$, δG_F and δm_Z^2 indicate the shifts in the corresponding quantities due to the vacuum polarisation effects under discussion. These relations show that three combinations of α , G_F and m_Z . The remaining three combinations can be expressed as follows in terms of Δr_W , $\Delta\rho$ and $\Delta k'$:

$$\Delta\rho = \varepsilon_1 \quad (11)$$

$$\Delta k' = (-c^2 \varepsilon_1 + \varepsilon_3)/(c^2 - s^2) \quad (12)$$

$$\Delta r_W = -c^2/s^2 \varepsilon_1 + (c^2 - s^2)/s^2 \varepsilon_2 + 2 \varepsilon_3 \quad (13)$$

where

$$\varepsilon_1 = A_{ZZ}/m_Z^2 - A_{WW}/m_W^2 = (A_{33} - A_{WW})/m_W^2 \quad (14)$$

$$\varepsilon_2 = F_{WW} - F_{33} \quad (15)$$

$$\varepsilon_3 = F_{3\gamma}/s - F_{33} = c/s F_{30} \quad (16)$$

The indices 3,0 refer to $W_3 = s\gamma + cZ$ and $W_0 = c\gamma - sZ$, the $SU(2)$ completion of W^+ and W^- and the $U(1)$ gauge vector, respectively. Equivalently the three observables can be related to the quantities $\Delta_p(0)$ and the derivatives in $q^2=0$ of $\Delta_p(q^2)$ and of $\Delta_3(q^2)$ introduced in ref.[2]. Note that ϵ_1 and ϵ_2 are only different from zero if a difference between W^+ and W_3 is induced. For example a split isospin multiplet of heavy particles can contribute to ϵ_1 and ϵ_2 while an unsplit multiplet cannot. On the other hand, both split and unsplit multiplets can contribute to ϵ_3 .

3

As already mentioned, in the standard model, large corrections due to heavy top as well as logarithmic terms in the large Higgs mass limit appear in the vacuum polarization amplitudes. Precisely, the quadratic terms in m_t only contribute to ϵ_1 , leading to Eq. (1). The physical Higgs gives rise to logarithmic terms which are mainly in ϵ_1 with a small term in ϵ_3 as well. For large m_H , keeping only the leading logarithm, one has[15]:

$$\epsilon_1 = G_F m_W^2 / (2\pi^2 \sqrt{2}) \left(-3/4 s^2/c^2 \log(m_H^2/m_W^2) \right) \quad (17)$$

$$\epsilon_2 = 0 \quad (18)$$

$$\epsilon_3 = G_F m_W^2 / (2\pi^2 \sqrt{2}) \left(1/12 \log(m_H^2/m_W^2) \right) \quad (19)$$

The interesting observation of refs. [5-8] is that, while ϵ_2 and ϵ_3 are relatively small in the standard Higgs case, so that the Higgs contribution approximately follows the same pattern as the $G_F m_t^2$ terms, ϵ_3 can be large in technicolour theories (actually so large that many versions of the model are already excluded by present data, as will be clear below). In the notation of ref.5 we have

$$\epsilon_1 = \alpha T \quad \epsilon_3 = \alpha S / (4s^2) \quad (20)$$

In technicolour theories and similarly in chiral Lagrangian type models [5-8] no large contributions to ϵ_2 are expected.

Based on these examples, one could imagine that ϵ_2 could be neglected in general, on the grounds that it requires isospin violations in the derivatives F_{WW} and F_{33} (see eq.15). For $\epsilon_2 = 0$, the following relation among the physical observables holds:

$$c^2/s^2 \Delta_p + 1/(c^2-s^2) \Delta_{rw} - 2\Delta K' = \epsilon_2/s^2 = 0 \quad (21)$$

Although it is clearly more difficult to introduce large isospin violations in the derivatives (as the examples of m_t , m_H and the effective chiral Lagrangian models [5-8] illustrate), we think that in general all the effects of the F_{ij} 's, which contribute to both ϵ_2 and ϵ_3 , cannot be neglected. In fact the F_{ij} 's, being dimensionless, are not necessarily suppressed by a large mass. We now explicitly produce a phenomenological model, where both ϵ_2 and ϵ_3 can be large.

4

We consider the possibility of anomalous γWW and ZWW couplings [16], [17] by adding to the standard model Lagrangian the non-renormalizable couplings

$$\Delta L = iM_W^2 (g_\gamma \lambda_\gamma F_{\mu\nu} \gamma + g_Z \lambda_Z F_{\mu\nu} Z) W^{+\mu} W^{-\nu} \quad (22)$$

where $g_\gamma = e$, $g_Z = e c/s$ and $F_{\mu\nu} \gamma$, $F_{\mu\nu} Z$, $W_{\mu\nu}$ are the non-Abelian field strengths associated to the γ , Z and W bosons respectively

$$F_{\mu\nu} \gamma = \partial_\mu A_\nu - \partial_\nu A_\mu + ig_\gamma W_\mu^+ W_\nu^- - (\mu \leftrightarrow \nu) \quad (23)$$

$$F_{\mu\nu} Z = \partial_\mu Z_\nu - \partial_\nu Z_\mu + ig_Z W_\mu^+ W_\nu^- - (\mu \leftrightarrow \nu) \quad (24)$$

$$W_{\mu\nu}^{\pm} = (\partial_\mu - ig_\gamma A_\mu - ig_Z Z_\mu) W_\nu^{\pm} - (\mu \leftrightarrow \nu) \quad (25)$$

These couplings would contribute to the magnetic and electric quadrupole moment vertices the extra terms: ($i = \gamma, Z$)

$$\Delta\mu_i = -\lambda_i g / 2M_W \quad (26)$$

$$\Delta Q_i = -\lambda_i g / M_W^2 \quad (27)$$

We concentrate on these couplings, and not on the anomalous γWW and ZWW couplings of dimension 4, since the λ 's contribute, at the level of the leading divergences, to vacuum polarization amplitudes and not to vertices. Retaining only the well defined leading divergences, cut off at a scale Λ , one obtains for the relevant vacuum polarisation functions [16]:

$$\Delta\Pi_{WW} = g^2 / (32\pi^2) q^2 \{ \Lambda^4 / m_W^4 (s^2 \lambda_\gamma^2 + c^2 \lambda_Z^2) + 12 \Lambda^2 / m_W^2 (s^2 \lambda_\gamma + c^2 \lambda_Z) \} \quad (28)$$

$$\Delta\Pi_{ZZ}=g^2/(32\pi^2)q^2c^2\{\Lambda^4/m_W^4\lambda_Z^2+12\Lambda^2/m_W^2\lambda_Z\} \quad (29)$$

$$\Delta\Pi_{\gamma\gamma}=g^2/(32\pi^2)q^2s^2\{\Lambda^4/m_W^4\lambda_\gamma^2+12\Lambda^2/m_W^2\lambda_\gamma\} \quad (30)$$

$$\Delta\Pi_{Z\gamma}=g^2/(32\pi^2)q^2cs\{\Lambda^4/m_W^4\lambda_\gamma\lambda_Z+6\Lambda^2/m_W^2(\lambda_\gamma+\lambda_Z)\} \quad (31)$$

Note that all $\Delta\Pi_{ij}$ are proportional to q^2 , so that $\Delta A_{ij} = 0$ and $\Delta F_{ij} = \Delta\Pi_{ij}/q^2$. By means of eqs.(28-31), it is a simple matter of algebra to obtain the corresponding contributions to $\epsilon_1, \epsilon_2, \epsilon_3$ (denoted by $\Delta\epsilon_i$)

$$\Delta\epsilon_1 = 0 \quad (32)$$

$$\Delta\epsilon_2 = -g^2/(32\pi^2)\Lambda^4/m_W^4s^2c^2(\lambda_\gamma\lambda_Z)^2 \quad (33)$$

$$\Delta\epsilon_3 = -g^2/(32\pi^2)\{\Lambda^4/m_W^4c^2(\lambda_\gamma\lambda_Z) (s^2\lambda_\gamma+c^2\lambda_Z) + 6\Lambda^2/m_W^2c^2(\lambda_\gamma\lambda_Z)\} \quad (34)$$

We see that in this phenomenological model $\Delta\epsilon_1=\Delta\rho=0$ whereas both $\Delta\epsilon_2$ and $\Delta\epsilon_3$ are different from zero. A similar term is not found in the effective chiral Lagrangian approach of refs.[6],[8] because there it is assumed that the only breaking of the custodial SU(2) is induced by the U(1) coupling g' . Note in fact that all the contributions to the physical observables are proportional to $\lambda_\gamma\lambda_Z$ which vanishes for a W_3WW coupling. A strong bound on $|\lambda_\gamma\lambda_Z|$, although Λ dependent, is therefore already implied by available experimental data. For $\lambda_\gamma = \lambda_Z = \lambda$, only $\log\Lambda$ terms would survive in vacuum polarization amplitudes and in vertices [18] as well. As a consequence, in this case, the significance of precision tests on λ at LEP1 and in low energy experiments is weakened relative to production experiments at hadron colliders and LEP2 [19].

This phenomenological model shows that ϵ_2 can also be significant, so that the presence of all the 3 terms $\epsilon_1, \epsilon_2, \epsilon_3$ cannot be excluded in a model independent way. This strengthens the point that $\Delta r_w, \Delta\rho$ and $\Delta k'$ should be treated as independent parameters.

It is interesting to study the values that can be attributed to $\Delta r_w, \Delta\rho$ and $\Delta k'$ from present data. Δr_w is directly obtained from Eq. (2) by using the value of m_W/m_Z measured by CDF and UA2 [20]:

$$s_W^2 = 1 - m_W^2/m_Z^2 = 0.2265 \pm 0.006 \quad (35)$$

One obtains

$$\Delta r_w = -0.015 \pm 0.018 \quad (36)$$

a direct information on $\Delta\rho$ comes from the value of g_A for charged leptons measured at LEP. The combined value of g_A from the four LEP experiments is given by [21]:

$$g_A = -0.501 \pm 0.0025 \quad (37)$$

From Eq. (5) one obtains

$$\Delta\rho = 0.004 \pm 0.010 \quad (38)$$

Similarly, the combined value [21] for g_V is $g_V = -0.045 \pm 0.006$, which implies

$$g_V/g_A = 1 - 4(1 + \Delta k')s_0^2 = 0.090 \pm 0.012 \quad (39)$$

Solving for $\bar{s}_W^2 = (1 + \Delta k')s_0^2$, one finds $\bar{s}_W^2 = 0.2275 \pm 0.003$. By comparison with the value of s_0^2 defined in Eq. (6) and determined by $m_Z = 91.177 \text{ GeV}$ [21] to be $s_0^2 = 0.2313$, one gets

$$\Delta k' = -0.017 \pm 0.013 \quad (40)$$

Incidentally, had we defined, rather than $\Delta k'$ via eq.5, the more usual quantity Δk from

$$g_V/g_A = 1 - 4(1 + \Delta k)s_W^2, \quad (41)$$

the measurement of g_V/g_A would have given a relation between Δk and Δr_w . We then prefer to work with

$$\Delta k' = \Delta k + c^2/(c^2 - s^2)\Delta r_w \quad (42)$$

in order to obtain directly $\Delta k'$ from gv/gA . Notice also that $\Delta k'$, unlike Δk , is independent of ε_2 .

Important additional information is obtained from neutrino hadron scattering. The ratios $R_V = \sigma_{NC}^{vN} / \sigma_{CC}^{vN}$ and $R_{\bar{V}} = \sigma_{NC}^{\bar{v}N} / \sigma_{CC}^{\bar{v}N}$ of neutral to charged currents in deep inelastic neutrino and antineutrino scattering are given by [22],[23]:

$$R_V = \rho^2(q^2) [1/2 - \bar{s}w^2(q^2) + 5/9(1+r)\bar{s}w^4(q^2)] \quad (43)$$

$$R_{\bar{V}} = \rho^2(q^2) [1/2 - \bar{s}w^2(q^2) + 5/9(1+\bar{r})\bar{s}w^4(q^2)] \quad (44)$$

where r ($\bar{r} = 1/\bar{r}$) is also measured. The experimental values of $R_V, \bar{V}$ and r , corrected for the non-isoscalarity of the target, QED effects, weak boxes and vertices, but no oblique corrections, are given in Ref. [23].

After correcting for the small effect due to the change of scale from the typical q^2 of the neutrino scattering experiments up to m_Z^2 [23],[24] one can express the neutrino results as an ellipse in the $\Delta k'$ - Δp plane as plotted in Fig. 1. This figure, which also shows the constraints from the measurements of g_A [Eq. (37)] and gv/gA [Eq. (39)], summarizes, together with Eq. (36), the main present information on $\Delta r_w, \Delta p$ and $\Delta k'$. Note that the determination of Δp obtained from g_A , given in Eq. (38), is considerably improved by adding the information from R_V and $R_{\bar{V}}$ as can be seen from Fig.1. The resulting value is:

$$\Delta p = 0.001 \pm 0.007 \quad (45)$$

We also find $\varepsilon_2 = 0.002 \pm 0.019$ and $\varepsilon_3 = -0.008 \pm 0.012$.

It is of interest to know the foreseen "ultimate" precision in the determination of these parameters at LEP. From the expected uncertainty in the m_W measurement at LEP2 [19], $\sigma(m_W) = \pm 100 \text{ MeV}$, one has correspondingly $\sigma(\Delta r_w) = \pm 0.005$. The following uncertainties in the measurements of the τ -polarization, left-right and $b\bar{b}$ forward-backward asymmetries are foreseen [26]: $\sigma(A_\tau) = \pm 0.011$, $\sigma(A_{LR}) = \pm 0.003$, $\sigma(A_{FB}(b\bar{b})) = \pm 0.005$. They respectively correspond to the following uncertainties in $\Delta k'$: $\sigma(\Delta k') = \pm 0.007$, ± 0.0016 , ± 0.005 . Finally, from about a factor of 2 reduction in the errors on the Z-widths, the error on Δp can go in Eq. (40) down to $\sigma(\Delta p) = \pm 0.007$.

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To distinguish different physical models, it may be useful to try to isolate the measured values of the different ε_i -terms defined in Section 2.. Independently of ε_2 , the values of ε_1 and ε_3 can be read off from Fig. 1, also showing the standard model predictions [25] for different values of m_t and m_H . As noticed, since Δr_w depends on ε_2 as well, the constraint coming from the m_W measurement cannot be represented in the $\Delta k'$ - Δp or ε_1 - ε_3 planes. Conversely, the relative smallness of ε_2 could be seen by the consistency of the present constraint from m_W/m_Z , Eq. (36), with the other measurements, if it were drawn in Fig. 1 for $\varepsilon_2 = 0$.

As an illustration of the impact of these considerations in constraining possible extensions of the standard model, consider first the case of technicolour theories [9]. As already mentioned, the parameter ε_3 is expected [5-8] to be significantly larger in technicolour theories than in the standard model. More precisely, the following estimate is given in Ref.[5] for the additional contribution to ε_3 (N_C is the number of technicolours):

$$\begin{aligned} \Delta \varepsilon_3 = & 0.0032 + 0.0006(N_C-4) \\ & 0.017 + 0.0032(N_C-4) \end{aligned} \quad (46)$$

for the minimal models containing one technicolour doublet (first line) or one whole technicolour generation with the same $SU(2)_L \times U(1)$ quantum numbers of (v, e, u, d) (second line). The value of $\Delta \varepsilon_3$ increases in a roughly proportional way to the total number of weak doublets. The importance of the recent data on gv/gA or $\Delta k'$ in constraining $\Delta \varepsilon_3$ from above is apparent from Fig. 1. Even more apparent is the potential impact of the foreseen precision on the various asymmetry measurements. Notice that the bound that one can obtain on ε_3 is independent from the value of ε_2 .

In the minimal supersymmetric standard model [11] all significant corrections are indeed concentrated in the vacuum polarisation amplitudes if all s-partners are heavy enough, i.e. not accessible to LEP2. In this case, for whatever spectrum in the two-doublet Higgs sector, with good approximation, the radiative corrections reproduce [12] the values obtained in the standard model for a relatively light Higgs (i.e. $m_H < m_Z$). The only

exception would be the case of a large contribution of the $\bar{t}-\bar{b}$ doublet to ε_1 [12],[27], which could be realised for a particularly light s-top.

7

The sizeable effects of virtual heavy particles are usually concentrated in the vacuum polarizations of the gauge bosons. As mentioned in Section 1, this is not however a general rule. Large corrections may also appear in the vertex functions, as exemplified by the heavy top effects in the $Zb\bar{b}$ vertex. In turn this calls for a way to describe these effects as well.

Under the assumption of consistently also neglecting in vertex functions all terms of order q/Λ , where Λ is the hypothetical new scale, the physical observables introduced in Section 2 are sufficient to describe all possible effects in a fixed fermionic channel. The difference due to the vertex corrections with respect to the oblique effects lies in the fact that the form factors g_V and g_A may become non-universal. In general they may depend on the fermion under consideration, so that one is led to define the parameters $\Delta\rho^f$ and Δk^f from $\Gamma(Z \rightarrow f\bar{f})$ and the corresponding asymmetries, generalising Eqs. (3)-(6). For the b-quark, from the top vertex corrections only, one has

$$\begin{aligned}\Delta\rho^b(\text{ top vertex-corrections}) &= -(\sqrt{2}G_F/4\pi^2)m_t^2 \\ \Delta k^b(\text{ top vertex-corrections}) &= (\sqrt{2}G_F/8\pi^2)m_t^2\end{aligned}\quad (47)$$

If no other large vertex corrections occur, these contributions should be added to the oblique effects to obtain the overall $\Delta\rho^b$ and Δk^b .

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Figure Caption

Constraints on $\Delta k'$ and $\Delta \rho$ (at $q^2 = m_Z^2$) imposed by the measurements of g_A (vertical band), g_V/g_A (horizontal band) and neutrino or antineutrino deep inelastic scattering (the area inside the ellipse [23]). The shaded area denotes the $1-\sigma$ intersection. The region corresponding to the prediction of the standard model, including complete radiative corrections [25], for $m_Z = 91.177$ GeV, $m_t = 90-190$ GeV and $m_H = 40-1000$ GeV is shown. The measured value of $\Delta k'$ is slightly smaller than required for a perfect agreement with the standard model. The ϵ_3 axis is also shown. Note that the standard model predicts a slightly positive value of ϵ_3 .