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Constraints on Dynamical Symmetry Breaking Mechanisms from Electroweak Data

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IN MEMORIAM M.A.B. BÉG *

Abstract

Consistency of the Salam-Weinberg theory, including quantum corrections, with high precision data from LEP and elsewhere imposes non-trivial bounds on the parameters of this theory, in particular the top quark mass. We take stock of the available experimental information in the electroweak sector with the view of constraining possible additional interactions, such as present in dynamical symmetry breaking scenarios. Using the Peskin-Takeuchi isospin conserving, S, and -violating, T, parametrization of new physics contribution to vacuum polarization corrections, we show here that the full one family technicolor models are ruled out at the 95% C.L. from the LEP data and m_W-measurements alone. We stress the role of improved precision measurements of the W-boson mass and the decay width $\Gamma(Z \to b\bar{b})$ in the enhanced sensitivity gained on such interactions.

*To be published in the M.A.B. Bég Memorial Volume, World Scientific Publishers, Singapore (1991).

1 Introduction

The Standard Model (SM) of particle physics[1] is remarkably successful in describing all the present electroweak experimental measurements spanning an energy range from a few electron Volts (eV), as in the atomic parity-violation experiments,[2] to about 100 GeV, as is the case for experiments at LEP. This consistency is both quantitative and impressive,[3-5] and leaves little doubt that the underlying group structure of the standard model, namely $SU(2) \times U(1)$ with three fermion families, will survive as the correct effective theory at present energies, in a granderconceivably more attractive- framework.

Despite this stunning success of the standard model, two important ingredients of this model, namely the Higgs boson and the top quark, are still missing in direct experimental searches. Indirect hints for the existence of the top quark are, however, numerous and convincing. In particular, the isodoublet character of the beauty quark, established in e^+e^- annihilation experiments through the forward-backward charge asymmetry measurements is well established,[5] which could be interpreted as an existence proof for the isospin partner of the b-quark, namely the top quark. Likewise, measurement of the $B^0 - \bar{B}^0$ mixing with the mixing parameter $x_d \equiv (\Delta M/\Gamma) = 0.70 \pm 0.13$,[6] not only requires a SM top quark, it also pitches its mass beyond 100 GeV.[7] In the same vein, the consistency of the SM in the electroweak sector, judged at the one-loop quantum correction level, allows constraints on the top quark mass. Recent analyses predict a top quark with SM couplings in the range $m_t = 140 \pm 40$ GeV.[8] Since the direct searches of the top quark, with the SM couplings, also yield a lower bound of 90 GeV on its mass,[9] there is no objective reason to be apprehensive about the standard model in the top quark sector.

The experimental constraints on the Higgs boson mass, as well as direct tests of the underlying Higgs mechanism for the generation of masses in the SM, as incorporated by Weinberg and Salam,[1] are at present, on the other hand, less incisive. For example, a recent analysis by Ellis and Fogli[8] to constrain the SM Higgs boson mass from electroweak data yields 1.8 $GeV \leq m_H \leq 6$ TeV at 68 % C.L. A more stringent lower bound on m_H is provided by LEP experiments from direct searches of the Higgs Boson, namely $m_H \geq 48$ GeV.[4] The standard model relation: $\rho \equiv m_W/m_Z \cos\theta_W = 1$, including the radiative corrections from known physics, imposes serious constraints on new physics, in particular on the Higgs representations, as emphasized some time ago by Veltman.[10] The consist determination of this parameter in a number of independent electroweak measurements- quantitatively established to at least $O(1\ \%)$ level- requires the Higgs boson to be a weak isodoublet, *if* the spontaneous symmetry breaking is induced by the Higgs field(s). It is, however, known that the relation $\rho = 1$, at the tree level, also follows in alternative models of symmetry breaking, like the technicolor/hypercolor models (hereafter generically called TC models),[11] where this relation emerges as a consequence of a remaining $SU(2)$ symmetry (also called custodial symmetry) of the charged and neutral (would be) Goldstone boson couplings to the weak currents.[12,13] Including the quantum corrections both in SM and TC models, precision measurements of the quantity ρ in the ongoing electroweak experiments are expected to provide a distinction between the two scenarios.[14,15]

The issue of distinguishing a weakly interacting SM Higgs from a strongly interacting Higgs (and Higgs-like system) is a central one. Future high energy experiments in e^+e^- annihilation and pp collisions have the potential of searching for a Higgs boson with a mass of practically upto 1 TeV.[16] At these energies, it should also be possible to discriminate between the canonical Higgs mechanism and theories where the gauge symmetries are dynamically broken. For example, in TC models where the $SU(2) \times U(1)$ symmetry is broken by the $< F\bar{F} >$ condensates, with F a Technifermion, one has very often a rich spectrum of hadrons.[11] These additional hadrons are of two variety; the so-called Techni-hadrons, like Techni-rho, Techni-eta etc., which have a mass of the order of the technicolor scale, denoted generically here by Λ_{TC}, with $\Lambda_{TC} = O(1\ TeV)$, and the pseudo-goldstone bosons (PGB's), denoted in what follows by π''s, which are present in such theories and are light on the scale of Λ_{TC}. The PGB's have their masses protected by global symmetries and get masses only after gauge interactions are taken into account. Both of these varieties of new hadrons have to be discovered, in a scaled up sense of the QCD hadron spectrum, if TC scenarios are correct. The experimental verdict on the PGB's and Techni-hadrons is not yet in, and their searches in the expected mass scales will be set forth at ongoing and future high energy colliders (for searches of Techni-hadrons at the SSC and LHC, see, for example, Ref. 17).

While direct searches of Higgses/PGB's are continuing, it is imperative to look for indirect signatures of dynamical symmetry breaking through precision electroweak experiments. In this context, the vacuum-polarization amplitudes (so-called oblique corrections) could be advantageously used to search for, or set limits on, such additional interactions.[18] The prerequisite for meaningful constraints is the availability of high precision electroweak data, which, thanks to LEP, are now at hand. This data have been used to set limit on TC models,[19-24] supersymmetry, and higher dimensional Higgs representations.[23] In addition, possible anomalous couplings of the top quark, which may emerge in specific realizations of models, where the $SU(2) \times U(1)$ symmetry is dynamically broken only by the top quark-antiquark condensate,[25] have also been entertained in literature.[26,27] Electroweak data, in particular from LEP, constrain such couplings but we shall not discuss such models here.

In this report, dedicated to the memory of our colleague, M.A.B. Bég, who had contributed very significantly to the early development of the subject of dynamical symmetry breaking,[28-30] we present an analysis of the recent electroweak data with

the view of constraining effects which are typical of such symmetry breaking mechanisms. For the TC models, on which we concentrate here, we present an analysis, based on the updated data from LEP and the present measurements of m_W, making use of the Peskin-Takeuchi isospin-conserving, S, and -breaking, T, parametrization of new effects in the vacuum polarization amplitudes.[19] In doing so, we have included the data from the atomic parity-violating experiments (in Cesium)[2,3] also, which, as emphasized by Marciano and Rosner,[20] is very efficient in constraining the variable S.

The results of our analysis can be summarized as follows. We first incorporate the one-loop radiative corrections in the standard model, using the so-called $\overline{MS}$ -scheme.[31,32] The consistency of this framework with the present electroweak measurements (in particular the LEP data and m_W) allows us to set upper and lower bounds on m_t, which are in line with the ealier such analyses, namely that $m_t < 200\ GeV$ at 90% C.L. Having done this, we follow the Peskin-Takeuchi analysis[19] involving the variable S and T and determine that the present data exclude $S < 0.8$ at 90% C.L., and $S < 1.0$ at 95% C.L . These numbers correspond to an assumed top quark mass of 150 GeV and a Higgs mass of 100 GeV. The dependence of S on m_t and M_H is rather weak. Simple TC models, which are scaled-up versions of QCD, predict:[19,20]

$$S = N_{TG} \times\ 2.1 + 0.4(N_{TC} - 4) \tag{1}$$

for a full technicolor generation. Thus, our analysis rules out a full technigeneration model at 95% C.L. We emphasize that the role of data from Cesium parity-violation is not crucial in reaching this conclusion, though this piece of data is still significant in constraining S. Similar estimates yield:

$$S = N_{TG} \times\ 0.4 + 0.08(N_{TC} - 4) \tag{2}$$

for N_{TC} number of technicolors and N_{TG} number of technigeneration for a doublet of techniquarks. Our analysis suggests that TC models with more than one doublet of techniquarks are also ruled out at 90% C.L. We briefly discuss the importance of improving the precision on the W-boson mass determination, as well as on the decay width, $\Gamma(Z \to b\bar{b})$, in future LEP experiments, in addition to the improvement of the Cesium data, which has already been emphasized in literature.[20]

This report is organized as follows. In Section 2, we review the general analysis of new physics contribution to various electroweak observables. Section 3 describes the (S,T)-analysis. This formalism is applied to the present electroweak data and the resulting contours are displayed in Figs. 1–3. Section 4 contains concluding remarks about the TC models.

2 General analysis of new contributions to oblique corrections

In this section we study the modifications induced by physics associated with a scale Λ much larger than the Fermi scale on the predictions for the various electroweak observables. We begin by assuming that the SM is basically correct; we then take fully into account the electroweak SM corrections using the modified minimal subtracted scheme ($\overline{MS}$).[31,32] In this scheme the weak mixing angle defined at a scale $\mu = m_Z$ is related to the fine structure constant, α, the Fermi constant, G_μ and the Z mass, m_Z, through[33]

$$\hat{s}^2 \hat{c}^2 = \frac{A^2}{m_Z^2 (1 - \Delta\hat{r})} \tag{3}$$

where $\hat{s}^2 \equiv \sin^2\hat{\theta}_W(m_Z)$, $\hat{c}^2 \equiv 1 - \hat{s}^2$, $A = (\pi\,\alpha/\sqrt{2}\,G_\mu)^{\frac{1}{2}}$ and

$$\Delta\hat{r} = Re\left[\frac{A_{ZZ}(m_Z^2)}{m_Z^2} - \frac{A_{WW}(0)}{m_W^2} - \frac{2\delta e}{e}\right]_{\overline{MS}} + \cdots \tag{4}$$

Here A_{ij} refers to the ij self-energy, $\frac{2\delta e}{e}$ contains the vacuum polarization correction to the photon propagator, the subscript $\overline{MS}$ denotes both the $\overline{MS}$ renormalization and the choice $\mu = m_Z$ for the 't Hooft mass scale, and the dots here and henceforth represent box and vertex contributions suitable for the process under consideration (in this case muon decay). Any new physics contribution will affect the determination of $\hat{s}^2$ from the Z mass through its effects in the various self-energies yelding a deviation from the SM value, $\hat{s}_0^2$, of

$$\hat{s}^2 = \hat{s}_0^2 + \frac{\hat{s}_0^2\hat{c}_0^2}{\hat{c}_0^2 - \hat{s}_0^2}\delta(\Delta\hat{r}) \equiv \hat{s}_0^2 + \delta\hat{s}^2 \tag{5}$$

(hereafter the subscript 0 appended to a quantity refers to its value in the SM).

In table 1 we list the experimental values of the various electroweak observables we use in our analysis together with their currently best determined values. In the following we discuss the dependence of these observables on new contributions. We begin with the W mass. To predict the W mass one uses the relation[37]

$$m_W^2 = \frac{A^2}{\hat{s}^2 (1 - \Delta\hat{r}_W)} \tag{6}$$

where

$$\Delta\hat{r}_W = Re\left[\frac{A_{WW}(m_W^2) - A_{WW}(0)}{m_W^2} - \frac{2\delta e}{e}\right]_{\overline{MS}} + \cdots \tag{7}$$

then assuming additional contributions to the self-energies we get

$$m_W = m_{W_0} - \frac{m_{W_0}}{2}\frac{\delta\hat{s}^2}{\hat{s}_0^2} + \frac{m_{W_0}}{2}\delta(\Delta\hat{r}_W) \tag{8}$$

Observable	Experimental value	Ref.
W mass	80.14 ± 0.31 (GeV)	9
$\Gamma(Z \to \text{all})$	2.485 ± 0.009 (GeV)	34
$\Gamma(Z \to \text{hadrons})$	1.740 ± 0.009 (GeV)	34
$\Gamma(Z \to b\bar{b})$	0.385 ± 0.023 (GeV)	36
$\sin^2\bar{\theta}_W$	0.2319 ± 0.0028	35
ρ_{eff}	0.998 ± 0.005	35
$Q_W(^{133}_{55}Cs)$	$-71.04 \pm 1.58 \pm 0.88$	2

Table 1: Electroweak observables used in our analysis and their current experimental values.

with

$$m_{W_0}^2 = \frac{A^2}{\hat{s}_0^2 (1 - \Delta\hat{r}_{W_0})}.$$

where $\Delta\hat{r}_{W0}$ is defined by (7) and takes into account only the SM contribution.

The modifications induced in the partial and total widths of the Z^0, whose present experimental values are given in table 1, can be summarized in the discussion of the partial decay width $Z^0 \to f\bar{f}$, where f stands for a generic fermion. We have in the $\overline{MS}$ framework[38] *

$$\Gamma_f = N_c^f \frac{\hat{\alpha}}{\hat{s}^2\hat{c}^2}\frac{m_Z\bar{\rho}_f(m_Z^2)}{48}\left[1 + \left(1 - 4|Q_f|\hat{k}_f\hat{s}^2\right)^2\right] \tag{9}$$

where N_c^f is the color factor, Q_f the electric charge of the fermion f,

$$\bar{\rho}_f(q^2) = 1 + \frac{\overline{A}_{ZZ}(q^2)}{q^2 - m_Z^2} + \cdots \tag{10a}$$

$$\hat{k}_f(q^2) = 1 - \frac{\hat{c}}{\hat{s}}\frac{A_{\gamma Z}^{\overline{MS}}(q^2)}{q^2} + \cdots \tag{10b}$$

with

$$\overline{A}_{ZZ}(q^2) = \left[A_{ZZ}(q^2) - Re\,A_{ZZ}(m_Z^2) - i\frac{q^2}{m_Z^2}Im\,A_{ZZ}(m_Z^2)\right]_{\overline{MS}}$$

$$A_{\gamma Z}^{\overline{MS}}(q^2) = \left[A_{\gamma Z}(q^2) - A_{\gamma Z}(0)\right]_{\overline{MS}}$$

*In Ref. 38 $\bar{\rho}_f$ was denoted $\bar{\rho}_{ff}$ and it is understood that we are considering only the real part of $\hat{k}$ and $\bar{\rho}_f$ neglecting $O(\alpha^2)$ contributions.

and $\hat{e}^2 = 4\pi\hat{\alpha}$ is the $\overline{MS}$ coupling at $\mu = m_Z$ as defined in Ref. 32. It is worth reminding that the dependence of $\overline{\rho}_f(q^2)$ and $\hat{k}_f(q^2)$ on the fermion f under consideration is due to vertex diagrams and "hidden" in Eqs. (10a), (10b) in the dots. In the case $Z^0 \to b\bar{b}$ these contributions involve the top quark and for large value of m_t are numerically not negligible.

New physics will modify the SM prediction for Γ_f according to

$$\Gamma_f = \Gamma_{f_0}(1+\delta\rho_f) - N_c^f \frac{\hat{\alpha}_0}{\hat{s}_0^2\hat{c}_0^2}\frac{m_Z\overline{\rho}_{f_0}(m_Z^2)}{6}\left(|Q_f|\hat{k}_{f_0}\hat{s}_0^2 - 4Q_f^2\hat{k}_{f_0}^2\hat{s}_0^4\right)\left(\frac{\delta\hat{s}^2}{\hat{s}_0^2} + \frac{\delta\hat{k}_f}{\hat{k}_{f_0}}\right) \quad (11)$$

where $\delta\rho_f$ is given by [†]

$$\delta\rho_f = \frac{\delta\overline{\rho}_f(m_Z^2)}{\overline{\rho}_{f_0}(m_Z^2)} - \frac{\hat{c}_0^2 - \hat{s}_0^2}{\hat{s}_0^2\hat{c}_0^2}\delta\hat{s}^2 + \frac{\delta\hat{\alpha}}{\hat{\alpha}_0} \quad (12)$$

with $\delta\overline{\rho}_f(m_Z^2)$, $\delta\hat{k}_f$ $\delta\hat{\alpha}$ being the shifts in the parameters $\overline{\rho}_f$, $\hat{k}_f$ and $\hat{\alpha}$ induced by new physics.

We discuss now the parameters ρ_{eff} and $\sin^2\overline{\theta}_W$. In the framework of an improved Born approximation one defines the leptonic partial width, Γ_l, and forward-backward asymmetry, A_{FB}^l, through the expressions

$$\Gamma_l = \frac{G_\mu m_Z^3}{6\pi\sqrt{2}}\left(\overline{v}_l^2 + \overline{a}_l^2\right) \quad (13)$$

$$A_{FB}^l = 3\frac{\overline{v}_e\,\overline{a}_e}{\overline{v}_e^2 + \overline{a}_e^2}\frac{\overline{v}_l\,\overline{a}_l}{\overline{v}_l^2 + \overline{a}_l^2} \quad (14)$$

with

$$\overline{v}_l = -\frac{1}{2}\sqrt{\rho_{eff}}(1 - 4\sin^2\overline{\theta}_W) \qquad \overline{a}_l = -\frac{1}{2}\sqrt{\rho_{eff}}.$$

Comparing Eqs. (13) and (9) we have,

$$\rho_{eff} = \frac{\hat{\alpha}\,\pi}{\hat{s}^2\hat{c}^2}\frac{\overline{\rho}_l(m_Z^2)}{\sqrt{2}\,G_\mu m_Z^2} \quad (15)$$

$$\sin^2\overline{\theta}_W = \hat{k}_l\hat{s}^2 \quad (16)$$

and consequently

$$\rho_{eff} = \rho_{eff_0} + \frac{\hat{\alpha}_0\,\pi}{\hat{s}_0^2\hat{c}_0^2}\frac{\overline{\rho}_{l_0}(m_Z^2)}{\sqrt{2}\,G_\mu m_Z^2}\delta\rho_l \equiv \rho_{eff_0} + \delta\rho_{eff} \quad (17)$$

$$\sin^2\overline{\theta}_W = \sin^2\overline{\theta}_{W_0} + \hat{k}_{l_0}\delta\hat{s}^2 + \delta\hat{k}\hat{s}_0^2 \equiv \sin^2\overline{\theta}_{W_0} + \delta\sin^2\overline{\theta}_W \quad (18)$$

[†] Do not confuse $\delta\overline{\rho}_f$ with $\delta\rho_f$!

In atomic parity-violation one has for the so-called weak charge, Q_W,

$$Q_W(Z,A) = \rho_{pv}\left[2Z - A - 4Z\kappa_{pv}(0)\hat{s}^2\right] \quad (19)$$

where

$$\rho_{pv} = 1 + \frac{A_{WW}(0)}{m_W^2} - \frac{A_{ZZ}(0)}{m_Z^2} + \cdots \quad (20)$$

and $\kappa_{pv}(0)$ is given by an expression similar to Eq. (10b) (for the Cesium atom $Z = 55$, $A = 133$). The additional contributions can be easily derived, giving

$$Q_W(Z,A) = Q_{W_0}(Z,A) + \frac{\delta\rho_{pv}}{\rho_{pv_0}}Q_{W_0}(Z,A) - 4Z\rho_{pv_0}(\delta\kappa_{pv}\hat{s}_0^2 + \kappa_{pv_0}\delta\hat{s}^2) \quad (21)$$

3 (S,T) formulation of new oblique corrections

It is well known that heavy particles, in general, need not decouple and could affect electroweak processes presently under experimental investigation through virtual effects in the oblique corrections.

To investigate these effects we assume a renormalization prescription that eliminates any new physics from $A_{\gamma\gamma}$ and $A_{\gamma Z}$. In this way only the W and Z self-energy will be affected by additional contributions. Following Peskin and Takeuchi,[19] we divide these contributions into an isospin-breaking part, T, defined by[20]

$$\frac{A_{WW}^{new}(0)}{m_W^2} - \frac{A_{ZZ}^{new}(0)}{m_Z^2} = \alpha(m_Z)T \quad (22)$$

and an isospin-conserving part, S, such as

$$\left.\frac{A_{WW}^{new}(m_W^2) - A_{WW}^{new}(0)}{m_W^2}\right|_{\overline{MS}} = \frac{\alpha(m_Z)}{4\hat{s}_0^2}S \quad (23a)$$

$$\left.\frac{A_{ZZ}^{new}(m_Z^2) - A_{ZZ}^{new}(0)}{m_Z^2}\right|_{\overline{MS}} = \frac{\alpha(m_Z)}{4\hat{s}_0^2\hat{c}_0^2}S \quad (23b)$$

where $\alpha(m_Z) = 1/127.8$ is the running α at the scale m_Z [*]. It should be noticed that our definition of the S and T parameters differs slightly from that of Ref. 20, where these parameters are defined for a fixed value of $\hat{s}_0^2$, namely $\hat{s}_0^2 = 0.2323$, corresponding to the top mass $m_t = 140$ GeV. Here, instead, $\hat{s}_0^2$ is assumed to be the value of the weak angle as calculated in the SM. This value depends on the top and Higgs masses, and this dependence is included in our analysis.

[*] $\alpha(m_Z)$ and $\hat{\alpha}$ differs in the treatment of the top-quark loops. In fact $\hat{\alpha}$ contains contributions from the top quark even when it is more massive than m_Z. In the parametrization of additional contributions we prefer to use the more conventional $\alpha(m_Z)$.

As pointed out in the literature[20,23,24] the above parametrization (Eqs. (22), (23)) is not really the most general one for new effects in the vacuum polarization amplitudes. In general, one has to allow two different parameters, S_W and S_Z, in Eq. (23a) and (23b), respectively. The use of a single parameter, S, in these equations implicitly assumes that there is no large isospin violation in the derivatives of the self-energies. As shown in Ref. 24, one can explicitly construct models where this condition is not satisfied yielding a large contribution both to S_Z and to $S_Z - S_W$ (in Ref. 24 the corresponding parameters are denoted by ϵ_3 and ϵ_2 respectively). In technicolor theories, however, an underlying symmetry suppresses $S_W - S_Z$. Therefore, as in Ref. 19 we consider only the isospin-conserving contribution $S \simeq S_W \simeq S_Z$.

According to Eq. (5) a non-zero S and/or T contribution will affect $\hat{s}^2$, adding to the predicted SM value $\hat{s}_0^2$ the quantity

$$\delta\hat{s}^2 = \frac{\alpha(m_Z)}{4(\hat{c}_0^2 - \hat{s}_0^2)}S - \frac{\alpha(m_Z)\,\hat{s}_0^2\,\hat{c}_0^2}{\hat{c}_0^2 - \hat{s}_0^2}T. \tag{24}$$

Noticing that

$$\delta(\Delta\hat{r}_W) = \frac{\alpha(m_Z)}{4\,\hat{s}_0^2}S$$

we have for the W mass (Cf. Eq. (8))

$$m_W = m_{W_0}\left[1 - \frac{1}{2}\frac{\alpha(m_Z)}{\hat{c}_0^2 - \hat{s}_0^2}\left(\frac{S}{2} - \hat{c}_0^2\,T\right)\right]$$

We now come to the discussion of the Z partial widths. From Eq. (10b) and our renormalization prescription it follows that $\delta\hat{\alpha} = \delta\hat{k}_f = 0$. Concerning $\delta\overline{\rho}_f(m_Z^2)$ we have, neglecting terms $O(m_Z^2/\Lambda^2)$,

$$\delta\overline{\rho}_f(m_Z^2) = \frac{\alpha(m_Z)}{4\hat{c}_0^2\hat{s}_0^2}S \tag{25}$$

Using Eqs. (24) and (25) one obtains to $0(\alpha)$

$$\delta\rho_f = \alpha(m_Z)T \tag{26}$$

then

$$\Gamma_f = \Gamma_{f_0} + \alpha(m_Z)\,T\left[\Gamma_{f_0} + N_c^f\frac{\hat{\alpha}_0}{\hat{c}_0^2 - \hat{s}_0^2}\frac{m_Z}{6}\overline{\rho}_{f_0}(m_Z^2)\left(|Q_f|\hat{k}_{f_0} - 4Q_f^2\hat{k}_{f_0}^2\hat{s}_0^2\right)\right]$$
$$- \frac{\alpha(m_Z)\,S}{4(\hat{c}_0^2 - \hat{s}_0^2)}N_c^f\frac{\hat{\alpha}_0}{\hat{s}_0^2\,\hat{c}_0^2}\frac{m_Z}{6}\overline{\rho}_{f_0}(m_Z^2)\left(|Q_f|\hat{k}_{f_0} - 4Q_f^2\hat{k}_{f_0}^2\hat{s}_0^2\right) \tag{27}$$

It is easy to see using Eqs. (17), (24) and (26) that

$$\delta\rho_{eff} = \frac{\hat{\alpha}_0\,\pi}{\hat{s}_0^2\hat{c}_0^2}\frac{\overline{\rho}_{f_0}(m_Z^2)}{\sqrt{2}\,G_\mu\,m_Z^2}\alpha(m_Z)T \tag{28}$$

$$\delta\sin^2\overline{\theta}_W = \hat{k}_{l_0}\left[\frac{\alpha(m_Z)}{4(\hat{c}_0^2 - \hat{s}_0^2)}S - \frac{\hat{s}_0^2\hat{c}_0^2}{\hat{c}_0^2 - \hat{s}_0^2}\alpha(m_Z)\,T\right] \tag{29}$$

For the weak charge of Cesium we have, instead, $\delta\kappa_{pv} = 0$ and

$$\delta Q_W(Z,A) = \alpha(m_Z)\,T\left[Q_{W_0} + 4Z\,\rho_{pv_0}\kappa_{pv_0}\frac{\hat{s}_0^2\hat{c}_0^2}{\hat{c}_0^2 - \hat{s}_0^2}\right] - \frac{Z\,\rho_{pv_0}\kappa_{pv_0}}{\hat{c}_0^2 - \hat{s}_0^2}\alpha(m_Z)\,S \tag{30}$$

In Fig. 1 we plot the allowed contours in the S and T plane (90 % C.L.) obtained by using LEP data and m_W from table 1. We have used $m_H = 100$ GeV and the values of the top quark mass are indicated on the figure. The SM point, $S = T = 0$, is well outside the allowed ellipse for $m_t = 200$ GeV showing that such a heavy top seems to be ruled out in the SM – a result already obtained in the latest general analyses of the electroweak data.[8] From the same figure it is clear that the case of a very heavy top ($m_t > 180$ GeV) is allowed only if there are compensating effects coming from new physics contributing *negatively* to T. It should be remarked that while in general additional corrections tend to give a positive contribution to T, that is equivalent to a positive contribution to the ρ-parameter[10] or a negative contribution to Δr,[39] there are specific examples where negative contributions to T are obtained for definite regions of the parameter space. Examples of this sort are a 4-th generation with Majorana neutrino,[40] and the 2-Higgs doublet model in a non-supersymmetric scenario.[41]

Fig. 2 shows the 95 % C.L. allowed domain in the (S,T) plane for $m_t = 150$ GeV and $m_H = 100$ GeV, using LEP data and m_W, and including the atomic-parity violation result. The 90% C.L. (S,T) contour using all the data in table 1 is also shown in this figure. From Fig. 2, we get $S < 1.4$ (at 95% C.L.) using LEP plus m_W data alone, with the constraint being more stringent if one uses in addition the Cesium data, giving $S < 1.0$ (95% C.L.). We remark that LEP plus m_W data alone rule out simple technicolor models with one full Technigeneration such as the one considered in Ref. 19, that predict $S \simeq 2$. As already pointed out in literature[20] the Cesium data is quite efficient in cutting the positive S region. This is still the case with the improved LEP data used here (as compared to the ones used in Ref. 20). At 90% C.L. including the Cesium data we find $S < 0.8$. This would then also rule out TC models with more than one Techni-doublet, as discussed in the introduction.

In Fig. 3 a comparison is made between the present capability of constraining S and T using LEP plus Cesium data, and what is foreseeable with a precision on the W mass of 80 MeV ($\Delta m_W/m_W \sim 10^{-3}$) and an error on the $Z \to b\overline{b}$ reduced to half

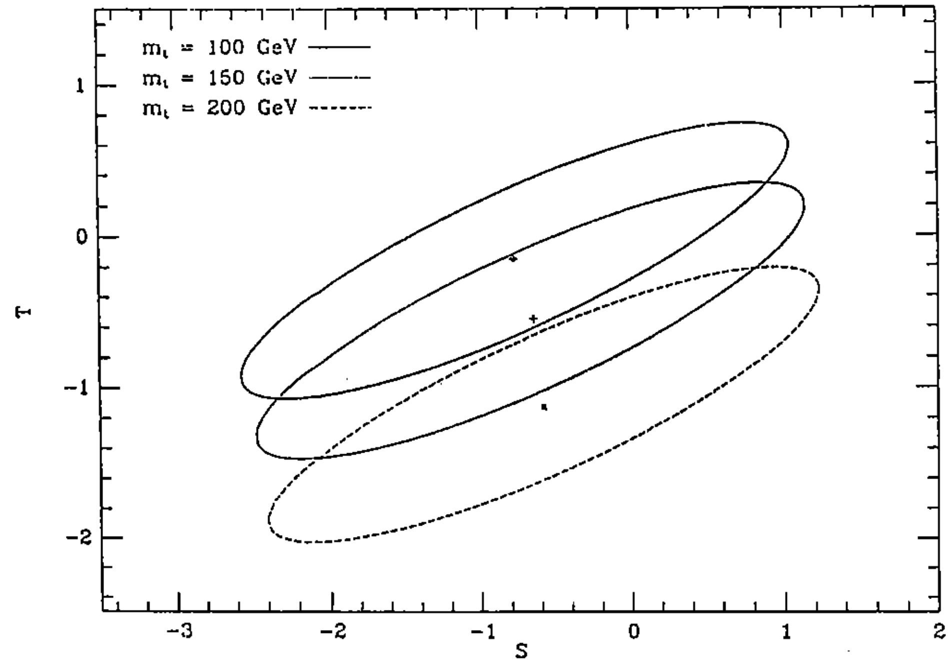

Figure 1: Allowed region in the (S,T) plane at 90% C.L., using LEP data and m_W from table 1. The value of the Higgs mass is fixed at $m_H = 100$ GeV.

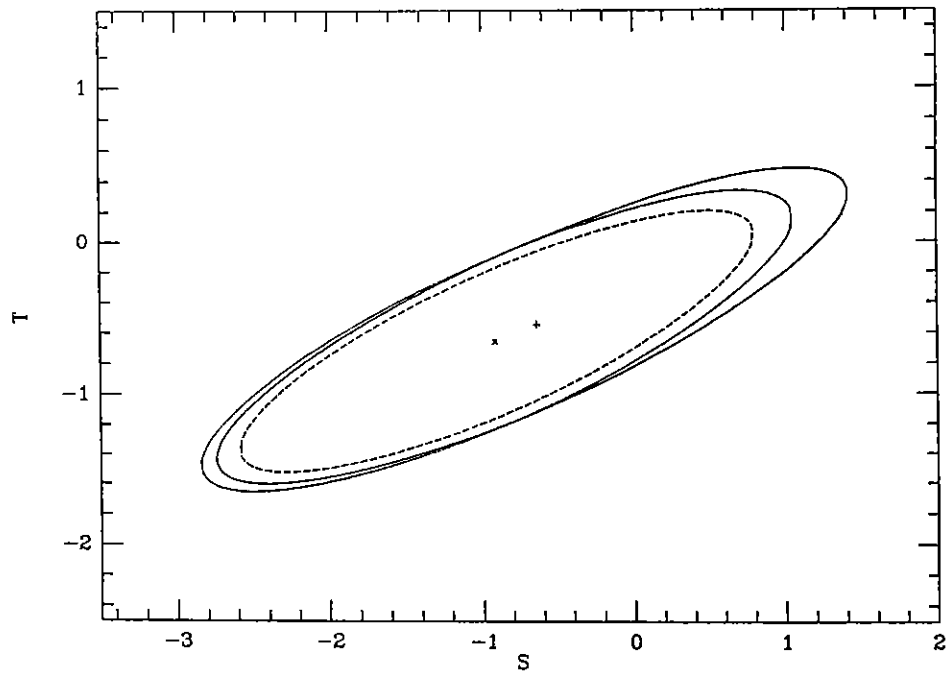

Figure 2: Allowed regions in the (S,T) plane; at 95% C.L. using LEP data and m_W (solid line), and including atomic-parity violation data (dotted line). The dashed curve corresponds to 90% C.L. countor using all the data of table 1. The curves are drawn for a top mass $m_t = 150$ GeV and a Higgs mass $m_H = 100$ GeV.

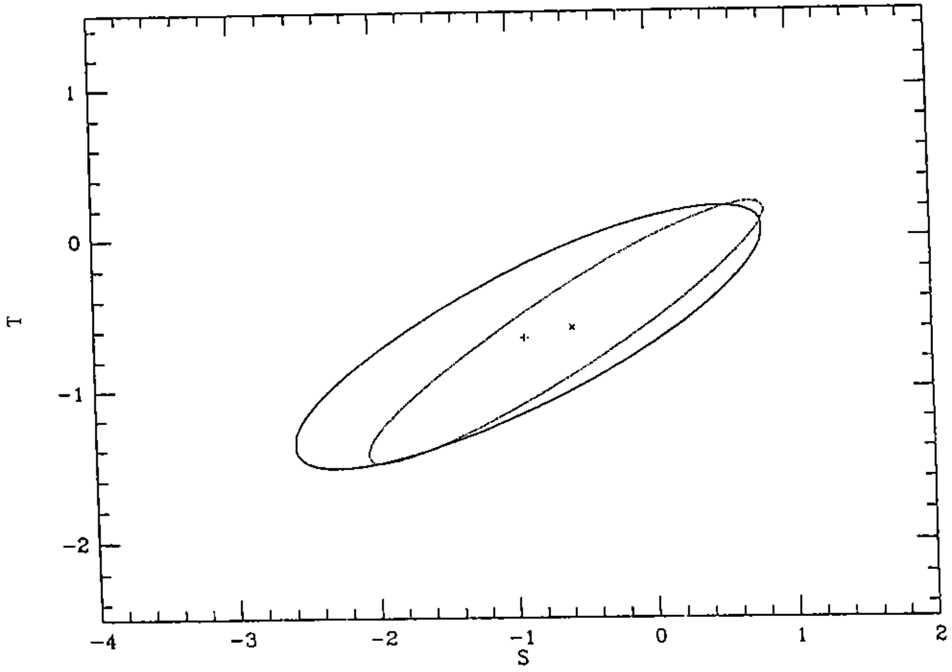

Figure 3: Allowed region at 90% C.L. using the present data from table 1 (solid line), and assuming an error of 80 MeV on the W mass and 12 MeV on the decay width $\Gamma(Z \to b\bar{b})$ (dotted line). We use $m_t = 150$ GeV and $m_H = 100$ GeV.

the present experimental error, i.e. 12 MeV. It is evident that these improvements can shrink the allowed region by roughly a factor 2. Clearly, this figure is illustrative only. One has to repeat the analysis with new measurements to get a precise bound on S.

4 Concluding Remarks

The analysis presented in sections 2 & 3 above is quite general, since most (though by no means all) new physics effects would show themselves through the non-zero S- and T-values. We have already mentioned the simplifying assumption of isospin invariance on the derivatives of the W- and Z- self-energies. This isospin breaking is much below the present level of experimental sensitivity in SM, it is also expected to be small in TC models which are scaled–up version of QCD. We leave the question of how to construct a natural TC model, in which the present electroweak constraints can be circumvented, to the standard bearers of such models!

Having discussed the constraints on the strongly interacting Higgs sector from precision electroweak measurements, perhaps a few general remarks on TC models concerning the production of PGB's are not out of place here. The PGB's are expected to have masses of $O(\alpha\Lambda_{TC}) \sim O(10\ GeV)$.[11] Concerning neutrals, it is known that the production mechanisms and decay properties of the SM Higgs and the neutral PGB's of the dynamical symmetry broken scenario are sufficiently different to be confused with each other.[29,30,42] These tests hinge primarily on the parity of the SM Higgs (O^+) and PGB's (O^-). The PGB's are expected to be negative-parity states in almost all realistic TC models. This can be motivated by analogies with the pattern of chiral symmetry breaking observed in QCD.[11] The different parities split the couplings of the SM Higgs and the PGB's with the gauge bosons. In particular, search strategies for the standard Higgs boson at LEP, which are based on the SM HZZ- coupling, become innocuous for the bounds on the π'^0, since the corresponding couplings $\pi'^0 ZZ$, being anomalous, occur only at the one-loop level (so-called triangle diagram), and hence are completely negligible.[30,42] The charged PGB's, $\pi'^{\pm}$ are, however, easy to produce, since in almost all models their couplings to the photon and the Z-boson are determined by their electroweak charges. They have measureable cross-sections and consequently have been searched for in e^+e^- annihilation experiments at PEP/PETRA/TRISTAN and at LEP. The signatures of the charged PGB's and charged Higgs bosons, $H^{\pm}$, which are present in extensions of the standard model with more than one Higgs doublet, are very similar. No such charged PGB or Higgs boson has todate been seen, and the present bound on the masses of $\pi'^{\pm}$ and $H^{\pm}$ is: $m_{\pi'^{\pm}}, m_{H^{\pm}} \geq 35 GeV$.[5]

While the neutral PGB's are admittedly difficult to produce and detect in ongoing experiments, the continued absence of the charged PGB's in the anticipated range of their mass, on the other hand, is curious. As already stated, the natural mass of the PGB's is $O(\alpha\Lambda_{TC})$, and the mass predictions can be vitiated only by introducing adhoc mass terms, which would take away the main appeal of TC models as viable candidates of mass generation. It is, therefore, a debateable proposition, whether the present negative searches of PGB's constitute a circumstantial evidence against the correctness of their mass estimates and/or against the TC models themselves, where such objects are naturally present. It should also be pointed out that the problem of fermion mass generation in the TC models (the so-called Extended Technicolor Models ETC), without running afoul with the flavour changing neutral current transitions[43] is still very much present. Attepmts at circumventing this by the so-called "walking" (instead of the usual running) coupling constants,[44] are, in our opinion, at best technical devices. The ETC models remain very much baroque constructions. In the ultimate analysis, only experiments are the jury. We look forward to the next generation of electroweak experiments at LEP and elsewhere to guide us in our search of mechanism(s) of spantaneous symmetry breaking.

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