BUHEP-97-27
UCB-PTH-97/44

September 19, 1997
hep-ph/9709297

Negative contributions to S in an effective field theory

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Abstract

We show that an effective field theory that includes non-standard couplings between the electroweak gauge bosons and the top and bottom quarks may yield negative contributions to both the S and T oblique radiative electroweak parameters. We find that that such an effective field theory provides a better fit to data than the standard model (the χ^2 per degree of freedom is half as large). We examine in some detail an illustrative model where the exchange of heavy scalars produces the correct type of non-standard couplings.

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1 Introduction

An important drawback of the standard model (SM) and its minimal supersymmetric extensions is that there are no theoretical constraints on the Yukawa couplings of the Higgs doublet, so that there are no clues about the quark and lepton spectrum. On the other hand, models which involve dynamical mechanisms for fermion mass generation, such as extended technicolor [1], typically induce corrections to precision electroweak observables that are in disagreement with data. Usually, the S and T oblique radiative electroweak parameters are too large [2], and the coupling of the Z gauge boson to b quarks is shifted so that the ratio of $Z \to b\bar{b}$ to $Z \to$ hadrons branching fractions is too small [5].

In this paper we show that, although shifts in oblique parameters and weak gauge couplings may individually be in disagreement with the precision electroweak data, their combination may lead to a much better fit than the SM. As we will demonstrate this possibility arises because shifts in weak gauge couplings can produce significant contributions to the oblique parameters. In Section 2 we discuss an effective field theory in which non-standard couplings of the electroweak gauge bosons to the third generation quarks may yield negative contributions to both S and T. In Section 3 we fit this effective field theory to the electroweak data. The couplings of the effective theory considered here can be produced, for example, by the exchange of heavy scalars in a technicolor model, as we show in Section 4. Our conclusions are presented in Section 5.

2 An effective field theory calculation

In extensions of the SM, in addition to the mass terms for the electroweak gauge bosons induced by the Higgs mechanism, there are new terms which can be generated in the effective Lagrangian below the scale of new physics. Two[1] phenomenologically important terms (of dimension 2 and 4 respectively) are [2, 3]:

$$\mathcal{L}_{\rm eff}^{\rm oblique} = -\alpha T^0 v^2 \, W^{3\mu} W^3_\mu - \frac{S^0}{16\pi} \, g g' \, B^{\mu\nu} W^3_{\mu\nu} \ , \tag{2.1}$$

where $W^3_{\mu\nu}$ and $B^{\mu\nu}$ are the gauge field strength tensors corresponding to the neutral gauge bosons of the $SU(2)_W \times U(1)_Y$ group that mix to produce the physical Z and the

[1]There is a third oblique correction parameter U, but it is generally much smaller than S and T in models without extra gauge bosons, and we will not consider it here.

photon, α is the electromagnetic coupling constant, g and g' are the weak and hypercharge gauge couplings, and v is the weak scale. These terms give tree-level contributions to the oblique radiative correction parameters S and T which can be defined [2] as:

$$S \equiv -\frac{8\pi}{M_Z^2}\left(\Pi_{3Y}(M_Z^2) - \Pi_{3Y}(0)\right) , \tag{2.2}$$

$$T \equiv \frac{4}{\alpha v^2}\left[\Pi_{11}(0) - \Pi_{33}(0)\right] , \tag{2.3}$$

where $\Pi_{jk}(q^2)$ are the vacuum polarizations of the W_j^μ and B^μ electroweak gauge fields due to non-SM physics, with the gauge couplings factored out[2]. Thus, in the effective theory we have tree-level and loop contributions to the oblique parameters:

$$S = S^0 + S^{\rm loop} , \tag{2.4}$$

$$T = T^0 + T^{\rm loop} . \tag{2.5}$$

New physics in the electroweak symmetry breaking sector can also induce changes in the interactions of quarks with electroweak gauge bosons [4, 5, 6]. In dynamical models of electroweak symmetry breaking, the new interactions are also responsible for generating masses for quarks and leptons, so we would expect to see the largest effects in the couplings of the top-bottom doublet. Below the scale of the new physics we can parameterize these effects with an effective Lagrangian which includes three new parameters $(\delta g_L, \delta g_R^{t,b})$:

$$\begin{aligned} \mathcal{L}_{\rm eff}^{\rm vertex} \;=\;& \delta g_L \left(g W_j^\mu - g'\delta_{3j} B^\mu\right) \overline{q}_L \gamma_\mu \sigma^j q_L \\ &+ \left(g W_3^\mu - g' B^\mu\right)\left[\delta g_R^t(\overline{t}_R\gamma_\mu t_R) + \delta g_R^b(\overline{b}_R\gamma_\mu b_R)\right] , \end{aligned} \tag{2.6}$$

where $q_L \equiv (t_L, b_L)$ is the left-handed $t-b$ quark doublet, and σ_j are the Pauli matrices. Such shifts in electroweak couplings produce tree-level effects in precision electroweak observables. In addition to these tree-level contributions (which come from integrating out physics above the effective field theory cutoff, Λ) there are loop corrections [7] from physics below the scale Λ which renormalize the coefficients in $\mathcal{L}_{\rm eff}$. If Λ is significantly larger than M_Z, then the leading logarithms from these one-loop corrections may be numerically important. For example in technicolor models, light pseudo-Nambu-Goldstone bosons can contribute to S and T at one-loop. In this section we will calculate the one-loop contributions to S and T from the shifts in the t and b gauge couplings.

[2] We are using the definition for hypercharge where $Y \equiv 2(Q - T_3)$.

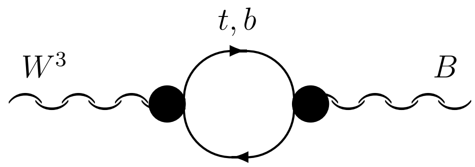

Figure 1: Contributions from t and b to Π_{3Y}. The $\bullet$ represents the effective couplings of the quarks to the gauge bosons.

It is convenient to use the following effective couplings:

$$\begin{aligned} \overline{g}_L^{t,b} &\equiv \pm\left(\frac{1}{2}+\delta g_L\right) , \\ \overline{Y}_L^{t,b} &\equiv Y_L \mp 2\delta g_L , \\ \overline{Y}_R^{t,b} &\equiv Y_R^{t,b} - 2\delta g_R^{t,b} , \end{aligned} \tag{2.7}$$

where $Y_L = 1/3$, $Y_R^t = 4/3$, and $Y_R^b = -2/3$ are the quark hypercharges. The one-loop contribution of the t quark to the vacuum polarization (see Fig. 1) is then:

$$\begin{aligned} \Pi_{3Y}^{(t)}(q^2) &= -\frac{N_c}{8\pi^2}\int_0^1 dx \left\{ 2\left(\overline{g}_L^t \overline{Y}_L^t + \delta g_R^t \overline{Y}_R^t\right)\left[x(1-x)q^2 - \frac{m_t^2}{2}\right]\right. \\ &\left. + \left(\overline{g}_L^t \overline{Y}_R^t + \delta g_R^t \overline{Y}_L^t\right) m_t^2\right\} \ln\left[\frac{\Lambda^2}{m_t^2 - x(1-x)q^2}\right] , \end{aligned} \tag{2.8}$$

where $N_c = 3$, m_t is the t mass, and a similar expression holds for the b quark contribution. Thus the one-loop result for the vacuum polarization is:

$$\begin{aligned} \Pi_{3Y}^{(t,b)}(M_Z^2) - \Pi_{3Y}^{(t,b)}(0) &\approx -\frac{N_c}{12\pi^2} M_Z^2 \left[\left(\overline{g}_L^t \overline{Y}_L^t + \delta g_R^t \overline{Y}_R^t\right) \ln\left(\frac{\Lambda}{m_t}\right)\right. \\ &\left. + \left(\overline{g}_L^b \overline{Y}_L^b + \delta g_R^b \overline{Y}_R^b\right) \ln\left(\frac{\Lambda}{M_Z}\right)\right] , \end{aligned} \tag{2.9}$$

where we kept only the leading-log terms (finite terms can be absorbed into the definition of S^0 and T^0 through the matching conditions at the scale Λ) and ignored terms suppressed by m_b^2/M_Z^2. To obtain the contribution to S we must subtract the SM contribution; the result to leading order in $\delta g_{L,R}$ is

$$\begin{aligned} S^{(t,b)} &\approx \frac{2N_c}{3\pi}\left\{\left[(Y_L - 1)\delta g_L + Y_R^t\, \delta g_R^t\right] \ln\left(\frac{\Lambda}{m_t}\right)\right. \\ &\left. + \left[-(Y_L+1)\delta g_L + Y_R^b\, \delta g_R^b\right] \ln\left(\frac{\Lambda}{M_Z}\right)\right\} . \end{aligned} \tag{2.10}$$

This result is quite general given a weak doublet of fermions with masses $m_b \ll M_Z < m_t$. Using the values of the hypercharges, and discarding terms proportional to $\ln(m_t/M_Z)$, we get

$$S^{(t,b)} \approx -\frac{4}{3\pi}\left(3\,\delta g_L - 2\,\delta g_R^t + \delta g_R^b\right)\ln\left(\frac{\Lambda}{M_Z}\right) \ . \tag{2.11}$$

We see from this expression that the contribution to S can be of either sign (since the shifts in the couplings can be of either sign) and large (it is enhanced with respect to the finite contribution of a weak-doublet of fermions by the logarithm and by a color factor).

We now move on to the isospin breaking effects which contribute to T. There are no corrections to T from δg_L; the only large correction to T is due to δg_R^t (a similar effect is discussed in [8]):

$$\Pi_{11}^{(t,b)}(0) - \Pi_{33}^{(t,b)}(0) \approx \delta g_R^t\,\frac{N_c}{4\pi^2}\,m_t^2\,\ln\left(\frac{\Lambda}{m_t}\right) \ , \tag{2.12}$$

which gives

$$T^{(t,b)} \approx \delta g_R^t \frac{3m_t^2}{\pi^2 \alpha v^2}\ln\left(\frac{\Lambda}{m_t}\right) \ . \tag{2.13}$$

Again we see that the contribution can be of either sign. We also note that $\delta g_R^t < 0$ induces negative contributions for both S and T.

3 Comparison with Experiment

In order to assess the usefulness of this formalism we have performed a fit to the precision electroweak data using the standard techniques [6, 9] with the parameters S, T, δg_L, and δg_R^b (at present there is no direct precision measurement involving δg_R^t). The deviations from the SM predictions for the physical quantities used in the fit in terms of these parameters are given in the Appendix. The experimental values [10, 11] and SM predictions [12] as well as the best fit values are given in Table 1.

The best-fit values are

$$\begin{aligned} \delta g_L &= 0.004 \pm 0.013 \ , \\ \delta g_R^b &= 0.036 \pm 0.068 \ , \\ S &= -0.40 \pm 0.55 \ , \\ T &= -0.25 \pm 0.46 \ . \end{aligned} \tag{3.1}$$

These values give a very good fit to the data; the χ^2 (i.e. sum of the squares of deviations over one standard-deviation errors) per degree of freedom (df) is $\chi^2/\mathrm{df} = 0.7$, while for

Quantity	Experiment	SM	Fit
Γ_Z	2.4947 ± 0.0026	2.4925	2.4948
R_e	20.756 ± 0.057	20.717	20.789
R_μ	20.795 ± 0.039	20.717	20.789
R_τ	20.831 ± 0.054	20.717	20.789
σ_h	41.489 ± 0.055	41.492	41.444
R_b	0.2179 ± 0.0011	0.2156	0.2177
R_c	0.1720 ± 0.0056	0.1720	0.1716
A^e_{FB}	0.0161 ± 0.0025	0.0155	0.0170
A^μ_{FB}	0.0165 ± 0.0014	0.0155	0.0170
A^τ_{FB}	0.0204 ± 0.0018	0.0155	0.0170
$A_\tau(P_\tau)$	0.1401 ± 0.0067	0.1440	0.1503
$A_e(P_\tau)$	0.1382 ± 0.0076	0.1440	0.1503
A^b_{FB}	0.0985 ± 0.0022	0.1010	0.0985
A^c_{FB}	0.0734 ± 0.0048	0.0720	0.0755
A_{LR}	0.1551 ± 0.0040	0.1440	0.1503
M_W	80.38 ± 0.14	80.34	80.35
M_W/M_Z	0.8814 ± 0.0011	0.8810	0.8811
$g_L^2(\nu N \to \nu X)$	0.3003 ± 0.0039	0.3030	0.3024
$g_R^2(\nu N \to \nu X)$	0.0323 ± 0.0033	0.0300	0.0297
$g_{eA}(\nu e \to \nu e)$	-0.503 ± 0.018	–0.507	–0.506
$g_{eV}(\nu e \to \nu e)$	-0.025 ± 0.019	–0.037	–0.039
$Q_W(Cs)$	-72.11 ± 0.93	–72.88	–72.56
$R_{\mu\tau}$	0.9970 ± 0.0073	1.0000	1.0000

Table 1: Experimental [10]-[12] and predicted values of electroweak observables for the SM for $\alpha_s(M_Z) = 0.115$. The SM values correspond to the best-fit values (with $m_t = 173$ GeV, $m_{\rm Higgs} = 300$ GeV) in [12], with $\alpha(M_Z) = 1/128.9$, and corrected for the change in $\alpha_s(M_Z)$.

the SM we find $\chi^2/\mathrm{df} = 1.50$. Thus we see that the SM (with α_s taken as an input rather than as an additional parameter to be fit to the electroweak data [9]) provides a relatively poor fit to the data [9]. Another way of stating this is that assuming the SM is correct, the probability of data giving a larger χ^2 is only 6%, while assuming the extended model is correct, the probability of finding a larger χ^2 is 78%.

For comparison, we also considered the case where only the S and T parameters are added to the SM, and found that the fit is almost as poor as in the case of the SM: $\chi^2/\mathrm{df} = 1.48$, corresponding to $S = -0.09 \pm 0.34$ and $T = 0.03 \pm 0.34$ (see also [13]).

Note that in the absence of isospin violation effects other than the one in eq. (2.13), the bound on T provides a tight constraint on δg_R^t:

$$\delta g_R^t < 0.02 \tag{3.2}$$

at the 95% confidence level (taking $\Lambda = 1$ TeV). It is amusing that measurements below the top quark threshold can produce such a tight bound on the top coupling.

4 A technicolor model with weak singlet scalars

In technicolor models, the quark and lepton masses may be produced by the exchange of gauge bosons [1], weak-doublet scalars [14], or weak-singlet scalars [15]. This latter alternative may explain certain features of the quark and lepton spectrum, and does not require complicated dynamics.

As an example application for the effective field theory calculation presented in Section 1, we examine in some detail the effects of a technicolor model with weak-singlet scalars on precision electroweak measurements. Because the main effects on electroweak observables are due to the sector responsible for electroweak symmetry breaking and the t and b masses, we will not be concerned with the mechanism which generates the other fermion masses. Also, we will assume the effects of the physics which keeps the scalars light (e.g. compositeness or supersymmetry) to be negligible at a scale of order 1 TeV (note that the supersymmetry breaking scale can be higher than in minimal supersymmetric extensions of the SM since fine-tuning is not an issue).

In addition to the SM fermions, consider one doublet of technifermions, P and N, and 3 scalars, ϕ, ω_t, ω_b, which transform under the $SU(N_{\mathrm{TC}}) \times SU(3)_{\mathrm{C}} \times SU(2)_{\mathrm{W}} \times U(1)_{\mathrm{Y}}$

gauge group (with $N_{\rm TC}$ even) as:

$$\Psi_R = \begin{pmatrix} P_R \\ N_R \end{pmatrix} \;:\; (N_{\rm TC},1,2)_0 \;, \quad P_L \;:\; (N_{\rm TC},1,1)_{+1} \;, \quad N_L \;:\; (N_{\rm TC},1,1)_{-1} \;,$$

$$\phi \;:\; (N_{\rm TC},\overline{3},1)_{-\frac{1}{3}} \;, \quad \omega_t \;:\; (N_{\rm TC},\overline{3},1)_{-\frac{7}{3}} \;, \quad \omega_b \;:\; (N_{\rm TC},\overline{3},1)_{\frac{5}{3}} \;. \tag{4.1}$$

The most general Yukawa interactions are contained in

$$\mathcal{L}_{\rm Y} = C_q \overline{\Psi}_R q_L \phi + C_t \overline{t}_R P_L \phi^\dagger + C_b \overline{b}_R N_L \phi^\dagger + C_{\omega_t} \overline{t}_R N_L \omega_t^\dagger + C_{\omega_b} \overline{b}_R P_L \omega_b^\dagger + {\rm h.c.} \;, \tag{4.2}$$

where the Yukawa coupling constants, $C_q, C_t, C_b, C_{\omega_t}, C_{\omega_b}$, are defined to be positive.

Assuming the ϕ techniscalar is sufficiently heavy to be integrated out, this results in a t mass

$$m_t \approx \frac{C_q C_t}{M_\phi^2} \pi v^3 \left(\frac{3}{N_{\rm TC}}\right)^{1/2} \;. \tag{4.3}$$

The only four-fermion operators induced by techniscalar exchange that will induce shifts in the electroweak couplings of the t and b are

$$\mathcal{L}_{\rm eff}^{\rm 4F} = -\left[\frac{C_t^2}{2M_\phi^2}(\overline{P}_L\gamma^\mu P_L) + \frac{C_{\omega_t}^2}{2M_{\omega_t}^2}(\overline{N}_L\gamma^\mu N_L)\right](\overline{t}_R\gamma_\mu t_R) - \frac{C_{\omega_b}^2}{2M_{\omega_b}^2}(\overline{P}_L\gamma^\mu P_L)(\overline{b}_R\gamma_\mu b_R)$$
$$-\frac{C_q^2}{4M_\phi^2}\left[\left(\overline{\Psi}_R\gamma^\mu\sigma_j\Psi_R\right)(\overline{q}_L\gamma_\mu\sigma_j q_L) + \left(\overline{\Psi}_R\gamma^\mu\Psi_R\right)(\overline{q}_L\gamma_\mu q_L)\right] \;. \tag{4.4}$$

These operators have an effect on the Z couplings to $t\overline{t}$ and $b\overline{b}$ which can be evaluated using an effective Lagrangian approach [5]:

$$\begin{aligned}
\overline{\Psi}_R\gamma^\mu\sigma_j\Psi_R &= -i\frac{v^2}{2}{\rm Tr}\left(D^\mu\,\Sigma\sigma_j\Sigma^\dagger\right) \;, \\
\overline{P}_L\gamma^\mu P_L &= i\frac{v^2}{2}{\rm Tr}\left(\Sigma^\dagger\frac{\sigma_3+1}{2}D^\mu\Sigma\right) \;, \\
\overline{N}_L\gamma^\mu N_L &= i\frac{v^2}{2}{\rm Tr}\left(\Sigma^\dagger\frac{-\sigma_3+1}{2}D^\mu\Sigma\right) \;.
\end{aligned} \tag{4.5}$$

Since Σ transforms as $C\Sigma W^\dagger$ under $SU(2)_C \times SU(2)_W$ (where $SU(2)_C$ is the custodial symmetry which has $U(1)_Y$ as subgroup), the covariant derivative is given by

$$D^\mu\Sigma = \partial^\mu\Sigma + ig\Sigma\frac{\sigma^k}{2}W_k^\mu - ig'\frac{\sigma^3}{2}B^\mu\Sigma \;. \tag{4.6}$$

With these expressions for operators, a comparison of eqs. (4.4) and (2.6) yields the couplings induced by scalar exchange [we eliminate M_ϕ using eq. (4.3)]:

$$\begin{aligned}
\delta g_L &= -\frac{C_q}{C_t}\frac{m_t}{8\pi v}\left(\frac{N_{\rm TC}}{3}\right)^{1/2} , \\
\delta g_R^t &= \frac{C_t}{C_q}\frac{m_t}{8\pi v}\left(\frac{N_{\rm TC}}{3}\right)^{1/2} - \frac{C_{\omega_t}^2 v^2}{8M_{\omega_t}^2} , \\
\delta g_R^b &= \frac{C_{\omega_b}^2 v^2}{8M_{\omega_b}^2} .
\end{aligned} \tag{4.7}$$

The technifermion contribution to S is estimated to be [2]

$$S^0 \approx 0.1 \ N_{\rm TC} , \tag{4.8}$$

so that, for[3] $N_{\rm TC} = 4$ and $\Lambda \approx 1$ TeV, the prediction of this model is

$$S = S^0 + S^{(t,b)} \approx 0.4 - 1.02\left(3\,\delta g_L - 2\,\delta g_R^t + \delta g_R^b\right) . \tag{4.9}$$

In addition to the direct isospin violation (2.13), there are "indirect" contributions to T from the technifermion mass spectrum which can be only roughly estimated [16]:

$$T^0 \sim \frac{N_{\rm TC}}{16\pi^2 \alpha v^2}\left(\Sigma_P(0) - \Sigma_N(0)\right)^2 , \tag{4.10}$$

where $\Sigma_P(q^2)$ and $\Sigma_N(q^2)$ are the technifermion self-energies. In this model, the origin of the indirect isospin violation is the difference between the C_t and C_b Yukawa couplings which accounts for the $t-b$ mass splitting. The four-fermion operator responsible for m_t,

$$\frac{C_t C_q}{M_\phi^2}\left(\overline{\Psi}_R q_L\right)\left(\overline{t}_R P_L\right) , \tag{4.11}$$

gives a one-loop correction to Σ_P which is quadratic divergent:

$$\Sigma_P(0) - \Sigma_N(0) = -\frac{3}{16\pi^3}\frac{m_t^2}{v^3}\Lambda'^2 , \tag{4.12}$$

where Λ' is expected to be of order 1 TeV, and generically different than Λ. Thus, the result is extremely sensitive to Λ',

$$T^0 \approx 0.51\left(\frac{N_{\rm TC}}{4}\right)\left(\frac{\Lambda'}{1{\rm TeV}}\right)^4 , \tag{4.13}$$

[3] It may be possible to have $N_{\rm TC} = 2$ if the scalars have a significant effect on the vacuum alignment [17], this issue has not been studied in detail in the literature.

and it is not possible to evaluate precisely this isospin breaking effect. For $N_{\rm TC} = 4$, the final expression for T is

$$T = T^0 + T^{(t,b)} \approx 0.51 \left(\frac{\Lambda'}{1\text{TeV}}\right)^4 + 34.0\, \delta g_R^t \ . \tag{4.14}$$

It is interesting that all the contributions from ω_t and ω_b exchange decrease S and T, as can be seen from eqs. (4.7), (4.9) and (4.14). For a range of values of the parameters $C_{t,q,\omega_t,\omega_b}$, M_{ω_t,ω_b} and Λ', the predictions for δg_L, δg_R^b, S and T, are within the correct ball-park (the best-fit values of eq. (3.1)). For example, for $C_t/C_q = 2.5$, $C_{\omega_t}(0.5\,\text{TeV})/M_{\omega_t} = 2$, $C_{\omega_b}(0.5\,\text{TeV})/M_{\omega_b} = 1.7$ and $\Lambda' = 1.2$ TeV, we find

$$\delta g_L \approx -0.013 \ , \ \delta g_R^b \approx 0.088 \ , \ S \approx 0.27 \ , \ T \approx -0.32 \ . \tag{4.15}$$

We emphasize, though, that given the sensitivity of T^0 to Λ', and the fact that the value of Λ' is fixed by the dynamics and not known precisely, the prediction for T is only potentially successful.

5 Conclusions

As we have seen in Section 3, if the SM were the correct theory of nature, then the probability of producing such a poor fit to electroweak measurements is only 6%. Therefore, it is worthwhile to look for extensions or alternatives to the SM which provide better fits to the existing data. The effective theory discussed in this paper greatly improves the fit (the χ^2/df is half as large) by including four additional parameters: two in the gauge boson self-energies and two in the gauge boson couplings to third generation quarks.

It is interesting that this effective theory can arise below the weak scale from a technicolor model which offers insight into the pattern of quark and lepton masses (for example, the mass hierarchy between the three generations may be induced by the masses of the exchanged scalars [15]). Given that compositeness and/or supersymmetry seem to be required to keep the scalars light, this technicolor model can only be considered as an effective theory valid up to roughly 1 TeV. However, to decide whether this model is a better candidate than the SM, one would need improved methods of computing the isospin breaking effects.

It should be stressed that the class of high energy theories which may lead to the effective theory discussed in Sections 2 and 3 is potentially large. For example, the analysis presented in Section 4 can be extended to any strong dynamics which gives rise to four-fermion interactions. Of course, the sizes and signs of the contributions to the S and T parameters are model dependent.

Acknowledgments

We would like to thank Sekhar Chivukula for helpful discussions and comments on the manuscript. *This work was supported in part by the National Science Foundation under grants PHY-9057173, and PHY-95-14797; and by the Department of Energy under grant DE-FG02-91ER40676.*

Appendix

In this Appendix we list the deviations from the SM predictions for the electroweak data in terms of δg_L, δg_R^b, S and T.

$$\Gamma_Z = (\Gamma_Z)_{SM}\left(1 - 3.8\times10^{-3}S + 0.010\,T - 0.707\,\delta g_L + 0.128\,\delta g_R^b\right) \tag{A.1}$$

$$R_{e,\mu,\tau} = (R_{e,\mu,\tau})_{SM}\left(1 - 2.9\times10^{-3}S + 2.0\times10^{-3}T - 1.01\,\delta g_L + 0.183\,\delta g_R^b\right) \tag{A.2}$$

$$\sigma_h = (\sigma_h)_{SM}\left(1 + 2.2\times10^{-4}S - 1.6\times10^{-4}T + 0.404\,\delta g_L - 0.073\,\delta g_R^b\right) \tag{A.3}$$

$$R_b = (R_b)_{SM}\left(1 + 6.6\times10^{-4}S - 4.0\times10^{-4}T - 3.56\,\delta g_L + 0.645\,\delta g_R^b\right) \tag{A.4}$$

$$R_c = (R_c)_{SM}\left(1 - 1.3\times10^{-3}S + 10.0\times10^{-4}T + 1.01\,\delta g_L - 0.183\,\delta g_R^b\right) \tag{A.5}$$

$$A_{FB}^{e,\mu,\tau} = (A_{FB}^{e,\mu,\tau})_{SM} - 6.8\times10^{-3}S + 4.8\times10^{-3}T \tag{A.6}$$

$$A_\tau(P_\tau) = (A_\tau(P_\tau))_{SM} - 0.028\,S + 0.020\,T \tag{A.7}$$

$$A_e(P_\tau) = (A_e(P_\tau))_{SM} - 0.028\,S + 0.020\,T \tag{A.8}$$

$$A_{FB}^b = \left(A_{FB}^b\right)_{SM} - 0.020\,S + 0.014\,T - 0.035\,\delta g_L - 0.191\,\delta g_R^b \tag{A.9}$$

$$A_{FB}^c = (A_{FB}^c)_{SM} - 0.016\,S + 0.011\,T \tag{A.10}$$

$$A_{LR} = (A_{LR})_{SM} - 0.028\,S + 0.020\,T \tag{A.11}$$

$$M_W = (M_W)_{SM}\left(1 - 3.6\times10^{-3}S + 5.5\times10^{-3}T\right) \tag{A.12}$$

$$M_W/M_Z = (M_W/M_Z)_{SM}\left(1 - 3.6\times10^{-3}S + 5.5\times10^{-3}T\right) \tag{A.13}$$

$$g_L^2(\nu N\to\nu X) = \left(g_L^2(\nu N\to\nu X)\right)_{SM} - 2.7\times10^{-3}S + 6.6\times10^{-3}T \tag{A.14}$$

$$g_R^2(\nu N \to \nu X) = \left(g_R^2(\nu N \to \nu X)\right)_{SM} + 9.4 \times 10^{-4} S - 1.9 \times 10^{-4} T \tag{A.15}$$

$$g_{eA}(\nu e \to \nu e) = (g_{eA}(\nu e \to \nu e))_{SM} - 3.9 \times 10^{-3} T \tag{A.16}$$

$$g_{eV}(\nu e \to \nu e) = (g_{eV}(\nu e \to \nu e))_{SM} + 7.2 \times 10^{-3} S - 5.4 \times 10^{-3} T \tag{A.17}$$

$$Q_W(Cs) = (Q_W(Cs))_{SM} - 0.796\, S - 0.011\, T \tag{A.18}$$

$$R_{\mu\tau} = (R_{\mu\tau})_{SM} \tag{A.19}$$

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