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A New Constraint on a Strongly Interacting Higgs Sector

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ABSTRACT

We show that an integral S over the spectral function of spin-l states of the Higgs sector is constrained by precision weak-interaction measurements. Current data excludes large technicolor models; experiments at LEP and SLC will soon provide more stringent limits on Higgs strong interactions.

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The most pressing question in the study of weak interactions is the nature of the Higgs sector. In the standard SU(2) x U(1) theoryo fweak interactions, some new particle or set of forces is needed to break the gauge symmetry. However, models of these Higgs particles ranging from a minimal doublet of scalar-fields to elaborate models with a rich spectrum are still consistent with experiment. We do not even know whether the new sector is weakly interacting, containing fundamental scalar fields and symmetry-breaking forces visible in perturbation theory, or whether it is strongly interacting, so that the symmetry-breaking results from nonperturbative phenomena.

Beyond the fact that the Higgs sector does break $SU(2)$ x U(1) in the required way, the main nontrivial piece of information that we have about its nature comes from the precision study of weak interaction parameters.[1] The relation[2] between the W and Z^0 masses and the value of the weak mixing angle $\sin^2\theta_w$

$$m_W^2/m_Z^2\cos^2\theta_w \cong 1\,, \tag{1}$$

to about 1%, implies that the Higgs sector has an approximate global $SU(2)$ 'custodial' symmetry.[3] Depending on the exact definition used to compute $\sin^2\theta_w$, the violation of the relation (1) measures the SU(2) symmetry violation in the Higgs sector.

In this letter, we will extend this conclusion to demonstrate that precision weak interaction experiments also constrain an ***isospin-symmetric*** observable of the Higgs sector. In essence, we will show that, by comparing weak interaction parameters, one can constrain not only the isospin asymmetry of this sector but also its total size. A longer and more complete version of this argument will be presented in ref. 4.

Our analysis will be based on the general formalism for weak interaction radiative corrections presented in refs. 5 and 6. This work begins from the observation that radiative corrections to weak-interaction processes involving light quarks and leptons due to new physics beyond the standard model appears dominantly through vacuum polarization amplitudes; in most models, the new physics does not significantly affect the vertex corrections or box diagrams. For example, modifications of the Higgs sector give vacuum polarization corrections of order α but vertex corrections are suppressed by an additional factor of $(m/m_W)^2$, where-m is the external fermion mass. The authors of ref. 5 called such corrections ***oblique*** and presented general formulae for the oblique radiative corrections to weak-interaction observables as combinations of vacuum polarization amplitudes. For specific cases, such as the effect of the heavy top quark, this formalism reproduced results well-known from the earlier literature.

We begin this analysis by writing the most important of these general formulae. Let us notate vacuum polarization amplitudes by using the subscripts 1,3 to denote the weak isospin currents $J^{\mu}_{1,3}$ and the subscript Q to denote the electric charged current. We write the Z^0 current as $(e/sc)\cdot(J_3 - s^2 J_Q)$, where s, c denote the sine and cosine of the weak mixing angle. The W boson self-energy is then written $i(e^2/s^2)\Pi_{11}(q^2)$; the Z^0–photon mixing amplitude is $i(e^2/sc)(\Pi_{3Q}(q^2) - s^2\Pi_{QQ}(q^2))$. The Ward identity requires $\Pi_{3Q}(0) = \Pi_{QQ}(0) = 0$.

As our basic set of weak-interaction observables, we choose the following:[7] α, G_F, m_W, m_Z, and the amplitudes $s^2_*(q^2)$, $\rho_*(q^2)$, Z_* defined by Kennedy and Lynn in ref. 6.[8] These last three amplitudes are defined as follows: We write the Z–f–$\overline{f}$ vertex as $(\text{const})\cdot(I^3 - s^2_*(q^2)Q)$. Then $s^2_*(q^2)$ is the ratio of the J_3 and J_Q terms in the Z^0 current and thus determines the weak interaction forward-backward and polarization asymmetries. For example, the polarization asymmetry for Z^0 production in e^+e^- annihilation is $A_{LR} \cong 8(1/4 - s^2_*(m_Z^2))$. The quantity ρ_* is a parameter of the effective Lagrangian of low-energy weak interactions:

$$\mathcal{L}_{eff} = \frac{4G_F}{\sqrt{2}}\left[J^{\mu}_{+}J^{\mu}_{-} + \rho_*(0)\left(J^{\mu}_3 - s^2_*(0)J^{\mu}_Q\right)^2\right]. \qquad (2)$$

It appears in R", the ratio of neutral to charged current cross sections in deep inelastic neutrino scattering. The quantity Z_* is the wavefunction renormalization of the Z^0 at its pole; it is measureable through the formula for the Z^0 width:

$$\Gamma_Z = Z_* \frac{\alpha_* m_Z}{6 s_*^2 c_*^2} \sum_f (I_f^3 - s_*^2 Q_f)^2 \cdot N_f \; , \tag{3}$$

where $c_*^2 = 1 - s_*^2$, starred parameters are evaluated at $q^2 = m_Z^2$, and N_f is the effective number of colors for the flavor f, including the QCD correction.

Since the Z^0 mass has now been measured with spectacular accuracy at LEP, it is most convenient to base-predictions in weak interaction theory on the measured values of α, G_F, m_Z. We find it convenient to define a weak mixing angle $\theta_w|_Z$ by

$$\sin 2\theta_w|_Z \equiv \left(\frac{4\pi\alpha_{*,0}(m_Z^2)}{\sqrt{2} G_F m_Z^2} \right)^{1/2} \; . \tag{4}$$

where $\alpha_{*,0}(m_Z^2)$ is the running electric charge, evaluated at the Z^0 mass, with the renormalization computed from known physics only. Then $\alpha_{*,0}^{-1} = 128.80 \pm 0.12$,[9] $m_Z = 91.172 \pm 0.032$ GeV[10] gives $\sin^2 \theta_w|_Z = 0.23147 \pm 0.00039$. The oblique corrections to the various weak interaction observables may be written as

$$m_W^2 - m_Z^2 \cos^2\theta_w|_Z$$

$$= -\left\{ \frac{e^2c^2}{s^2(c^2-s^2)} \left[\Pi_{33}(m_Z^2) - 2s^2\Pi_{3Q}(m_Z^2) - \frac{s^2}{c^2}\,\Pi_{11}((\frac{c^2-s^2}{c^2} \quad \Pi_{11}(m_W^2) \right] \right.$$

$$\left. + \frac{e^2s^2m_W^2}{c^2-s^2}\left[\Pi'_{QQ}(m_Z^2) - \Pi'_{QQ}(0)\right] \right\}$$

$$s_*^2(q^2) - \sin^2\theta_w|_Z$$

$$= \left\{ \frac{e^2}{c^2-s^2} \left[\frac{\Pi_{33}(m_Z^2) - 2s^2\Pi_{3Q}(m_Z^2) - \Pi_{11}(0)}{m_Z^2} - (c^2-s^2)\,\frac{\Pi_{3Q}(q^2)}{q^2} \right] \right.$$

$$\left. + \frac{e^2s^2}{c^2-s^2} \left[s^2\Pi'_{QQ}(m_Z^2) - c^2\Pi'_{QQ}(0) + (c^2-s^2)\,\Pi'_{QQ}(q^2) \right] \right\}$$

$$\rho_*(0) - 1 = -\frac{e^2}{s^2c^2m_Z^2}(\Pi_{33}(0) - \Pi_{11}(0))$$

$$Z_*(q^2) - 1 = \frac{e^2}{s^2c^2}\frac{d}{dq^2}\left(\Pi_{33} - 2s^2\Pi_{3Q} + s^4\Pi_{QQ}\right)\Big|_{q^2=m_Z^2}$$

$$- e^2\Pi'_{QQ}(0) - \frac{e^2(c^2-s^2)}{s^2c^2}\left(\Pi'_{3Q}(q^2) - s^2\Pi'_{QQ}(q^2)\right), \tag{5}$$

where $\Pi'(q^2)$ denotes $\Pi(q^2)/q^2$. We note that these formulae should be used only to compute the nonstandard corrections to these observables; the standard model radiative corrections are not generally oblique and do not fall into such a simple form. However, the corrections due to the Higgs sector are of the form of (5), and we may analyze the Higgs contribution by examining the right-hand sides of these relations.

For the renormalization due to a heavy top quark, it is well known that the dominant contribution to (5) comes from the difference

$$\Pi_{11} - \Pi_{33} = \frac{3\alpha}{16\pi\sin^2\theta_w}\frac{m_t^2}{m_W^2}, \tag{6}$$

which is momentum independent to a sufficient approximation. This renormalization affects m_W, s_*^2, ρ_* in a way which can be easily inferred. Any other momentum-independent, isospin-asymmetric contribution to the Π's will have the same general phenomenology. This difference cannot be larger than 1% or so-without having a grave effect on weak interaction phenomenology.

We now propose another simplification of (5) which is appropriate to conventional technicolor models.[11,12] In these models, the Higgs sector is a copy of the usual strong interactions scaled up to TeV energies. In such models, the technicolor interactions conserve conventional isospin and also parity, and we may divide the vacuum polarization amplitudes of weak isospin currents into contributions from correlators of vector and axial-vector currents: $\Pi_{11} = \Pi_{33} = (\Pi_{VV} + \Pi_{AA})/4$; $\Pi_{3Q} = \Pi_{VV}/2$. If we expand the Π's in powers of q^2 and ignore q^4 and above (making a relative error of m_W^2/M_T^2, where M is a technicolor mass scale),

$$\Pi_{VV}(q^2) = q^2\Pi'_{VV}(0) \quad \Pi_{AA}(q^2) = \Pi_{AA}(0) + q^2\Pi'_{AA}(0)\,. \tag{7}$$

Inserting (7) into (5), we find that m_W, s_*^2, and Z_* receive corrections proportional to $(\Pi'_{VV}(0) - \Pi'_{AA}(0))$. In a large technicolor model, as we will see below, this contribution can be as large as that of the heavy top quark. The idea that isospin-conserving contributions of technicolor can lead to large radiative corrections has been suggested independently by Golden and Randall.[13]

Let us now formalize these considerations as follows: Define the parameters S and T by[14]

$$\begin{aligned}\alpha S &= -e^2\frac{d}{dq^2}\left(\Pi_{VV}(q^2) - \Pi_{AA}(q^2)\right)\Big|_{q^2=0}\\ &= 4e^2\frac{d}{dq^2}\left(\Pi_{33}(q^2) - \Pi_{3Q}(q^2)\right)\Big|_{q^2=0} \qquad (8)\\ \alpha T &= \frac{e^2}{s^2m_W^2}\left(\Pi_{11}(0) - \Pi_{33}(0)\right)\,.\end{aligned}$$

Then, if we ignore terms of order q^2 in the isospin-violating pieces of the Π's and terms of order q^4 in the isospin-symmetric terms, we find

$$\begin{aligned} m_W^2 - m_Z^2 \cos^2\theta_w|_Z &= m_W^2 \frac{\alpha}{c^2 - s^2}\left[c^2 T - f s\right] \\ s_*^2(q^2) - \sin^2\theta_w|_Z &= \frac{\alpha}{c^2 - s^2}\left[-s^2 c^2 T + \frac{1}{4} S\right] \\ \rho_*(0) - 1 &= \alpha T \\ Z_* - 1 &= \frac{\alpha}{4 s^2 c^2} S \ . \end{aligned} \tag{9}$$

In principle, comparison of weak-interaction measurements can restrict S and T independently. We note, though that S and T are generally both positive, so that they tend to compensate one another in any single observable.

Before discussing the phenomenological constraints on S and T, let us compute the values of S and T in technicolor models. We begin with S. The combination of vacuum polarization amplitudes which appears in the definition of S obeys the dispersion relation

$$\Pi_{VV}(q^2) - \Pi_{AA}(q^2) = -\frac{q^2}{12\pi}\int \frac{ds}{\pi} \frac{\left[R_V(s) - R_A(s)\right]}{s - q^2} - f_\pi^2 \tag{10}$$

where $R_V(s)$, $R_A(s)$ are the analogues of R(s), the cross section for e^+e^- annihilation to hadrons in units of the point cross section, with the electromagnetic current replaced by the vector and axial vector isospin currents. Asymptotically, both $R_V(s)$ and $R_A(s)$ tend to the sums of the squares of the isospins of technifermions. In an asymptotically free gauge theory, one expects that the leading two terms of (10) vanish for large q^2; this leads to Weinberg's first and second spectral function sum rules relating R_V and R_A.[15] For our application, we need the $q^2 \to 0$ limit of (10), so that S is expressed as a zeroth Weinberg sum rule:

$$S = \frac{1}{3\pi}\int \frac{ds}{s}\left\{\left[R_V(s) - R_A(s)\right] - \frac{1}{4}\left[1 - (1 - \frac{m_H^2}{s})^3 \theta(s - m_H^2)\right]\right\} . \tag{11}$$

The last term of (11) is the contribution of the standard model Higgs boson sector which the technicolor theory replaces; m_H is the mass of the standard-model Higgs.

Note that, as $s \to 0$, $R_V(s) \to 1/4$ (since we consider the chiral limit where the pions are massless) and $R_A(s) \to 0$. Thus, subtraction of the standard-model contribution regularizes the infrared divergence in the first term of (11).

To evaluate (11), we need the vector and axial vector spectral functions for the technicolor theory. These we evaluated in the following way: We fit the data on the cross section for e^+e^- annihilation to hadrons in the familiar strong interactions and extracted its isovector part, in the form of the ρ resonance, an enhancement around 1400 MeV, and high-energy continuum. We then represented the axial-vector spectral function by the $a_1(1260)$ resonance and continuum, fitting the height of the resonance and the approach to the continuum by using the Weinberg sum rules. Having thus obtained $R_V(s)$ and $R_A(s)$ in the usual strong interactions, we obtained these quantities for a technicolor theory with 3 colors and 2 flavors by multiplying energies by (F_π/f_π), where f_π = 93 MeV, F_π = 250 GeV. To obtain the spectral functions for other values of the number of technicolors N_C and flavors N_F, we used the scaling laws appropriate to gauge theories at large N_C: The area under a resonance should scale as N_C, while the width varies as N_F/N_C. These simple scaling considerations should give results to about 20% accuracy. For the masses of pseudo-Goldstone bosons, we used the estimates given in refs. 16, 17. To evaluate S, we took m_H = 1 TeV. Using this procedure, we found

$$S = \begin{cases} 0.4 + 0.08(N_C - 4) & 1 \text{ doublet} \\ 2.1 + 0.4(N_C - 4) & 1 \text{ generation} \end{cases} , \qquad (12)$$

where the first line of (12) refers to the minimal model with one technifermion doublet and the second line to a model whose fermions have the quantum numbers of (ν, e, u, d). The value of S is roughly proportional to the total number of weak doublets. The full details of our estimate will be presented in ref. 4. For 1 generation of technifermions, this value of S gives a shift Δm_W = -500 MeV, in agreement with other estimates of the technicolor radiative correction.[5,13] This is a very large correction, and one is tempted to conclude on this basis that such a model is excluded. However, this depends on the value of T, to which we now turn.

Our estimate of T in technicolor models is much less reliable. In ref. 18, this quantity was estimated by evaluating the vacuum polarization diagrams using noninteracting technifermions with scale-dependent masses $\Sigma(k^2)$. This gives

$$T = \frac{1}{s^2 m_W^2}\frac{N_C}{16\pi}\int_0^\infty dk^2 \frac{k^4}{(k^2+\Sigma^2)^4}(\Sigma_U^2-\Sigma_D^2)^2 . \tag{13}$$

The value of T depends on the isospin splitting of the dynamical masses, which must appear if the model is to predict a large mass difference for the b and t quarks. But the integral is dominated by the contribution from the technicolor mass scale, exactly the region where the estimate is unreliable. If we crudely estimate $(\Sigma_U^2 - \Sigma_D^2) = 2m_t\Sigma$, independent of momentum, we find

$$T = \left(\frac{3}{16\pi s^2}\frac{m_t^2}{m_W^2}\right)\cdot\left[\frac{4N_C}{9}\right] . \tag{14}$$

The factor in parentheses is the standard top quark contribution, which should be added to the technicolor contribution (14). Note that the factor in brackets makes the technicolor contribution larger by roughly a factor $(N_C/2)$;[19] thus, the expectation for T in technicolor models is dangerously large. Even a more careful estimate of (13) with scale-dependent Σ's, given in ref. 18, show that it is very difficult to obtain $T < 5$ in models with $m_t > 100$ GeV. On the other hand, we should emphasize again that (13) itself is highly unreliable. A more sophisticated approach, or special tailoring of the technicolor dynamics, could well yield a smaller value of T.

This increases the importance of the parameter S. Since S is a simple integral over the zeroth-order spectrum of technicolor, we believe that the estimate we have given above is valid within a wide variety of technicolor models. Let us now address the question of whether S can be constrained independently of the value of T.

If we use the dependence (9) to. compute the weak interaction radiative correction due to new physics, any single measurement selects a line (or, rather, a band) in the ***S-T*** plane. By combining measurements, we may determine the region in the ***S-T*** plane that they collectively allow. In this analysis, we take the standard model calculations at the fixed values m_t = 150 GeV, m_H = 1000 GeV based on the measured values of α, G_F, and m_Z = 91.17 GeV, as reference values.[20] We do not intend to systematically reanalyze the whole corpus of weak interaction data; rather we select the few best-measured observables. These are listed, together with their current and reference values, in Table 1. The first three quantities in this table—R^ν, m_W/m_Z, and Γ_Z give the bands in the ***S-T*** plane shown in Fig. l(a). Unfortunately, the three curves are almost parallel, so that they do not give a strong constraint on S independently of the value of ***T***. From (9), it is clear that a direct measurement of s_*^2 would be a powerful way to constrain S, since the effect of S relative to ***T*** is largest in this quantity. Unfortunately, there seems to be no highly accurate direct measurement of s_*^2. The two best current measurements are listed in the fourth and fifth lines of Table 1; these produce the rather wide bands shown in Fig. l(b). Combining the five measurements, we find the likelihood contours shown in Fig. l(c). Fig. l(d) compares these measurements to the standard model with a varying top quark mass, and to two technicolor models using the estimate (14) for ***T***. We re-emphasize that, in the technicolor models only the computation of S is trustworthy. The likelihood analysis yields a 90% confidence upper limit on S: S $<$ 3.6 This would exclude models with two generations of technifermions-for which we expect S $\cong$ 4 for N_C = 4-but not yet models with a single generation.

The next two years of data-taking at LEP and SLC will produce much more accurate direct measurements of $s_*^2(m_Z^2)$, through the measurement of the polarization asymmetry A_{LR} and the forward-backward asymmetry to $b\bar{b}$ at the Z^0. Fig. l(d) shows the band allowed by a polarization asymmetry measurement with 10^5 Z^0s and 40% polarization; such a measurement would either verify or strongly constrain the possibility of a new strong interaction sector associated with the Higgs bosons.

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Table 1. Measurements used in to constrain S and T

0 bservable	Measured Value	(ref.)	Standard Model	(ref.)
R_o^ν	0.3095 ± 0.004	21	0.3115	22
m_W/m_Z	0.8791 ± 0.0036	23	0.8786	24
Γ_Z	2.540 ± 0.026	10	2.501	25
$\sin^2\theta_w^0(\nu e)$	0.233 ± 0.014	26	0.231	27
$A_{FB}^b(Z^0)$	0.133 ± 0.104	28	0.065	29
$A_{LR}(Z^0)$	?		0.129	29

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FIGURE CAPTIONS

1) Allowed region of the ***S,T*** plane: (a) bands in the ***S,T*** plane allowed by the measured values of ***R”***, m_W/m_Z, Γ_Z, within 1 σ errors; (b) bands allowed by the measurements of νe scattering and A^b_{FB}, within 1 a; (c) likelihood contours based on the five measurements, corresponding to 68% and 90% probability (the dotted curves include only the measurements in (a)); (d) comparison of these contours with the predictions of two technicolor models, and with the band allowed by an anticipated measurement of A_{LR}. All four figures show the standard model prediction for ***T*** and ***S*** as a function of the top quark mass, with stars at m_t = 50, 100, 150, 200, 250 GeV; the last figure shows the analogous curves for technicolor models with one doublet and one generation, respectively, N_C = 4, and ***T*** estimated using (14).

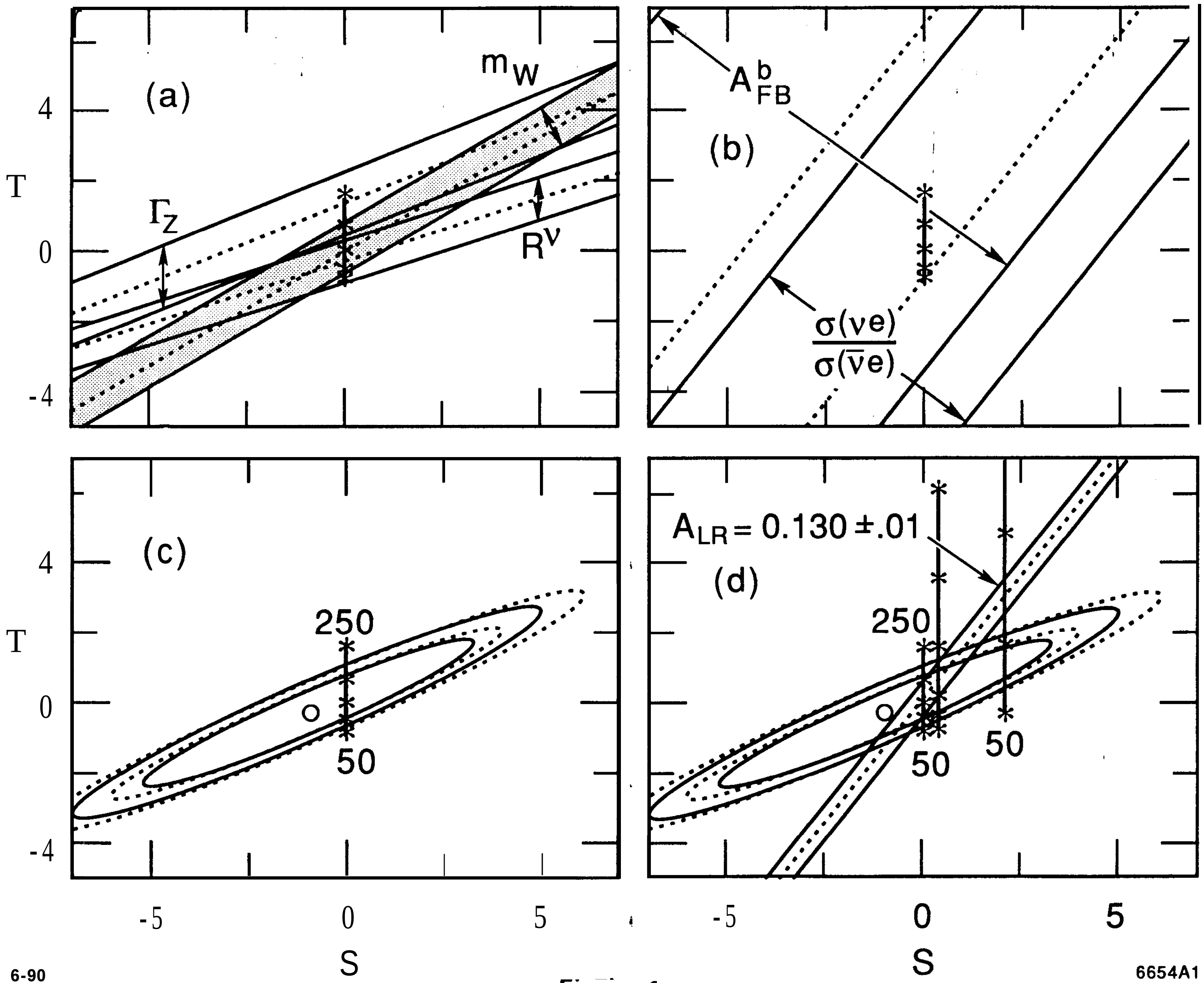
(a)
(b)
(c)
(d)
T
S
4
0
-4
-5
5
m_W
Γ_Z
R^ν
A_{FB}^b
$\frac{\sigma(\nu e)}{\sigma(\bar{\nu} e)}$
$A_{LR} = 0.130 \pm .01$
250
50
6-90
6654A1

Fig. 1