Precision Electroweak Experiments and Heavy Physics: A Global Analysis

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(Received 12 September 1990)

The radiative corrections of heavy degenerate and nondegenerate electroweak multiplets do not decouple from low-energy experiments. Their effects can be parametrized in a general renormalization-scheme-independent way. For isospin-breaking h_V and chiral-breaking h_{AZ} and h_{AW}, derived from the Z and W self-energies, a global analysis of existing neutral-current and Z and W data yields $-1.2 < h_V < 0.9$, $-3.5 < h_{AZ} < 1.0$, $-4.4 < h_{AW} < 2.4$ (90% C.L.). The expectations for future experiments are described, and the implications for the top-quark and Higgs-boson masses, new generations, supersymmetry, and technicolor are discussed.

PACS numbers: 12.15.−y, 12.50.−d, 14.80.Er, 14.80.Gt

One of the consequences of electroweak symmetry breaking is that radiative corrections due to heavy particles with weak isospin do not decouple from low-energy interactions. Such processes thus constrain the possible types and masses of particles heavier than the Z. The generic effects of this "nondecoupling" have been known for some time.[1] The analysis of heavy-top-quark and -Higgs-boson corrections was presented in Ref. 2, while more recently important partial analyses have been applied to technicolor theories in the papers of Peskin and Takeuchi[3] and Golden and Randall.[4]

The constraints on heavy physics are actually better than previously believed, as pointed out by Marciano and Rosner, because of recent improvements in atomic parity-violation experiments.[5] We present a new global analysis of heavy physics, incorporating the best current data and parametrizing the nondecoupling in a general renormalization-scheme-invariant way, derived from the formalism of Kennedy and Lynn.[6] We represent the heavy physics in a general three-parameter form and fit to all three independently without constraints.

Our analysis includes the complete one-loop standard-model (SM) contributions with variable top-quark and Higgs-boson masses. The effects of nonstandard (NS) physics are considered here *only* insofar as they appear in the gauge-boson self-energies ("oblique" in Ref. 3), as these occur universally in all electroweak processes. Let J_A, J_Z, J_W be the electroweak gauge currents. Thus $J_A = eJ_Q$, $J_Z = (e/sc)(J_3 - s^2 J_q)$, and $J_W = (e/s) \times J_1$, where 1 and 3 refer to weak isospin. Defining the scalars Π as the transverse parts of the self-energies (proportional to $g_{\mu\nu}$) then

$$\Pi_{QQ} = \langle J_Q J_Q \rangle, \quad \Pi_{3Q} = \langle J_3 J_Q \rangle, \quad \Pi_{33} = \langle J_3 J_3 \rangle, \quad \Pi_{11} = \langle J_1 J_1 \rangle . \tag{1}$$

Ward identities require $\Pi_{QQ} = q^2 \Pi'_{QQ}$ and (ignoring the gauge-boson loops) $\Pi_{3Q} = q^2 \Pi'_{3Q}$. Nondecoupling of heavy physics occurs in three independent finite combinations of Π's. The first is due to the breaking of custodial $SU(2)_V$ isospin (ρ parameter):

$$\Delta_\rho(q^2) \equiv \Pi_{11}(q^2) - \Pi_{33}(q^2) . \tag{2}$$

A second and third arise from the breaking of chiral $SU(2)_A$ isospin. One occurs in the neutral current,

$$\Delta_3(q^2) = \Pi_{33}(0) + q^2 \Pi'_{3Q} - \Pi_{33} ; \tag{3}$$

and the other in the charged current,

$$\Delta_1(q^2) = \Pi_{11}(0) + q^2 \Pi'_{3Q} - \Pi_{11} . \tag{4}$$

$\Delta_\rho(q^2)$ receives contributions from any isomultiplet with mass splittings. $\Delta_3(q^2)$ and $\Delta_1(q^2)$ arise from chiral symmetry breaking, e.g., fermion masses or technicolor. Π is the proper self-energy, computable perturbatively to any order, or nonperturbatively.[7]

The natural choice of measurables occurs at low energies ($q^2=0$: α, G_F, ρ) and at the gauge-boson poles ($q^2 = M_Z^2, M_W^2$: $M_Z, M_W, \Gamma_Z, \Gamma_W, Z$ asymmetries). G_F here refers to the universal Fermi constant with the nonelectromagnetic vertex and box corrections specific to given initial and final states removed; e.g.,

$$G_F = G_\mu \left\{ 1 + \frac{\alpha(M_Z)}{2\pi s_\theta^2} \left[3 - \left(\frac{1}{2} c_\theta^2 - 3\frac{c_\theta^2}{s_\theta^2} \right) \ln c_\theta^2 \right] - \frac{\alpha(M_Z)}{8\pi} \left(5\frac{c_\theta^4}{s_\theta^4} - 3 \right) \ln c_\theta^2 \right\}^{-1} , \tag{5}$$

where $c_\theta^2 \equiv M_W^2/M_Z^2$, and $G_\mu = 1.16637 \times 10^{-5}$ GeV^{-2} is the standard muon decay constant.[6,8] Dimensionless heavy physics parameters are then h_V, h_{AZ}, h_{AW}:[9]

$$\begin{aligned} \alpha h_V &\equiv 4\sqrt{2} G_F \Delta_\rho(0) , \\ h_{AZ} &\equiv -16\pi \Delta_3(Z)/M_Z^2 , \\ h_{AW} &\equiv -16\pi \Delta_1(W)/M_W^2 . \end{aligned} \tag{6}$$

Our SM reference point is $m_t = M_H = M_Z$, and we display explicitly only deviations from this reference; i.e.,

$$h = h^{\rm NS} + \Delta h^{\rm SM} , \tag{7}$$

where $h^{\rm NS}$ represents the contributions of nonstandard physics, and keeping only the leading quadratic and logarithmic terms,

$$\Delta h_V^{SM} = \frac{3G_F M_Z^2}{8\sqrt{2}\alpha\pi^2}\left[\left(\frac{m_t}{M_Z}\right)^2 - 1\right] - \frac{3G_F}{8\sqrt{2}\alpha\pi^2}(M_Z^2 - M_W^2)\ln(M_H^2/M_Z^2), \quad (8a)$$

$$\Delta h_{AZ}^{SM} = -\frac{1}{6\pi}\ln(m_t^2/M_Z^2) + \frac{1}{12\pi}\ln(m_H^2/M_Z^2), \quad (8b)$$

$$\Delta h_{AW}^{SM} = +\frac{1}{3\pi}\ln(m_t^2/M_Z^2) + \frac{1}{12\pi}\ln(m_H^2/M_Z^2). \quad (8c)$$

Let M_{Z0} and M_{W0} be the gauge-boson masses with $h_V = h_{AZ} = h_{AW} \equiv 0$. Then the physical masses are

$$M_Z^2 = M_{Z0}^2\left[\frac{1-\alpha h_V}{1-4\sqrt{2}G_F M_{Z0}^2(h_{AZ}/16\pi)}\right], \quad (9a)$$

$$M_W^2 = M_{W0}^2\left[\frac{1}{1-4\sqrt{2}G_F M_{W0}^2(h_{AW}/16\pi)}\right]. \quad (9b)$$

Here and below, $\alpha \equiv \alpha(0) \simeq 1/137$. Next, define $\Gamma_{Z,W}$ as the gauge-boson widths renormalized by the pole residues:[3,6] $\Gamma \equiv \Gamma^0/\text{residue}$. Then

$$\Gamma_Z = G_F M_Z^3 \gamma_Z/(1-\alpha h_V), \quad (10a)$$

$$\Gamma_W = G_F M_W^3 \gamma_W, \quad (10b)$$

a convenient representation that incorporates all of the nonheavy SM physics into $\gamma_{Z,W}$ in a renormalization-scheme-invariant way. M_{Z0}, M_{W0}, and γ_Z are the canonical functions of the weak mixing angle s_R^2, which is taken to be $s_*^2(M_Z^2)$ in the scheme of Ref. 6 or $\sin^2\theta(M_Z^2)_{\overline{MS}}$ in that of Ref. 10. ($\overline{MS}$ denotes the modified minimal-subtraction scheme.) All ordinary (i.e., due to light physics) radiative corrections are incorporated into these expressions and into γ_W.[11]

Finally, the low-energy neutral-current matrix element is

$$\mathcal{M}_{NC} = \frac{4\sqrt{2}G_F}{1-\alpha h_V}(I_{3L} - s_R^2 Q)(I'_{3L} - s_R^2 Q'). \quad (11)$$

Equation (11) is schematic only. The ordinary (light-physics) radiative corrections, including the running of s_R^2 to $q^2=0$ and vertex and box corrections should be included in $\mathcal{M}_{NC}$ even though they are not displayed.

The observables of Eqs. (9)-(11) depend on h_V, h_{AZ}, and h_{AW} in a simple way, provided M_{Z0}, M_{W0}, γ_Z, γ_W, and $\mathcal{M}_{NC}$ are written in terms of s_R^2. We leave s_R^2 as a free parameter to be determined in a simultaneous fit of h_V, h_{AZ}, h_{AW}, and s_R^2. An alternative would be to express the observables in terms of s_Z^2, which is the apparent value of the weak mixing angle calculated from M_Z assuming $h_{V,A} \equiv 0$ (e.g., $s_Z^2 = 0.2333 \pm 0.0004$ in the $\overline{MS}$ scheme, $s_Z^2 = 0.2330 \pm 0.0004$ in the $*$ scheme). From Eq. (9a),

$$s_R^2 = s_Z^2 - \frac{s_R^2 c_R^2}{1-2s_R^2}\left[\alpha h_V - 4\sqrt{2}G_F M_{Z0}^2 \frac{h_{AZ}}{16\pi}\right]. \quad (12)$$

One then fits with $h_{V,A}$ as free parameters.[5] Although the apparent dependence of Eqs. (9)-(11) on $h_{V,A}$ would be different, the two approaches are equivalent.

For the data set, we take (a) the low-energy neutral-current results of Ref. 2; (b) the CERN LEP Z-pole measurements ($M_Z = 91.172 \pm 0.031$ GeV, $\Gamma_Z = 2.498 \pm 0.020$ GeV, $\Gamma_{l^+l^-} = 84.0 \pm 0.9$ MeV, and $R = \Gamma_{had}/\Gamma_{l^+l^-} = 20.89 \pm 0.27$, where $\Gamma_{l^+l^-}$ is the average width for $Z \to e^+e^-$, $\mu^+\mu^-$, $\tau^+\tau^-$, and Γ_{had} is the total width to hadrons);[12] (c) the weak charge $Q_W = -71.04 \pm 1.58(\text{expt}) \pm 0.88(\text{th})$ for atomic parity violation in cesium, based on the data of Noecker *et al.*[13] and the new atomic calculation of Blundell, Johnson, and Sapirstein;[14] (d) the combined $M_W/M_Z = 0.880 \pm 0.004$ ($M_W = 80.2 \pm 0.3$ GeV) from UA2 and CDF;[15] and (e) the preliminary CHARM-II measurement of $\sigma(\nu_\mu e)/\sigma(\bar{\nu}_\mu e)$,[16] yielding

$$\sin^2\theta_W^0 = 0.240 \pm 0.009(\text{stat}) \pm 0.008(\text{syst})$$

in the tree-level formula.

Figure 1 displays the sensitivity of various experimental results to h_V, h_{AW}, and h_{AZ}. In particular, Q_W^{Cs} combined with M_Z strongly restricts h_{AZ}, as was emphasized in Ref. 5. Similarly, deep-inelastic neutrino scattering and M_Z, or Γ_Z and $\Gamma_{l^+l^-}$ combined with M_Z, each separately yield a correlated region in the (h_V,h_{AZ}) plane. The combination of M_W/M_Z with M_Z does also, if we assume $h_{AW} = h_{AZ} \equiv h_A$. The combination of M_W, M_Z, and the other quantities allow a simultaneous determination of h_V, h_{AW}, h_{AZ}, and s_R^2. One obtains the following:

$$\begin{aligned} &h_V = -0.1 \pm 0.8, \quad h_{AZ} = -1.1 \pm 1.7, \\ &h_{AW} = -0.8 \pm 2.6, \\ &\sin^2\theta(M_Z^2)_{\overline{MS}} = 0.230 \pm 0.004, \\ &s_*^2(M_Z^2) = 0.230 \pm 0.004. \end{aligned} \quad (13)$$

The corresponding upper limits on the h's are $h_V < 0.9$ (1.2), $h_{AZ} < 1.0$ (1.6), $h_{AW} < 2.4$ (3.3) at 90% (95%) C.L.[17] The results are hardly changed if, as in Ref. 3, we constrain $h_{AW} = h_{AZ} \equiv h_A$, for which $h_A < 0.7$ (1.2). Similarly, $h_V > -1.2$, $h_{AZ} > -3.5$, $h_{AW} > -4.4$ at 90% C.L. These results can be regarded as direct constraints on new heavy physics h^{NS} for the reference value $m_t = M_H = M_Z$. For arbitrary m_t,M_H, they apply to $h^{NS} + \Delta h^{SM}$, where the Δh^{SM} are given in (8). For $m_t \gg M_Z$ the dominant modification is the positive $(m_t/M_Z)^2$ contribution to Δh_V^{SM}, so that the upper limit on h_V^{NS} is more stringent: e.g., $h_V^{NS} < -0.5$ (-0.2) for $m_t = 200$ GeV, $M_H = M_Z$.

In Fig. 2 we display the sensitivity of expected future measurements of M_W, A_{LR}, and Q_W. It is clear that these future measurements will significantly improve the determination of h_{AZ}, an improvement mainly due to the polarization asymmetry A_{LR}. In fact, A_{LR} alone, combined with present data, would improve the 1σ uncertainty in h_{AZ} in Eq. (13) to approximately ± 0.5. The

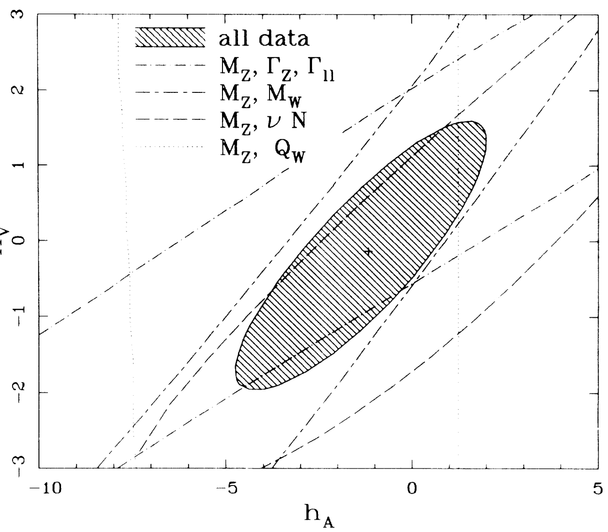

FIG. 1. 90%-C.L. regions in (h_A, h_V) from simultaneous fits of (Q_W, M_Z), $(\nu N, M_Z)$, $(\Gamma_Z, \Gamma_{l^+l^-}, M_Z)$, (M_W, M_Z), and *(all data)*. The first three fits are to $h_A \equiv h_{AZ}$, h_V, and s_R^2, and do not constrain h_{AW}. The last two are for the special case $h_{AW} = h_{AZ}$.

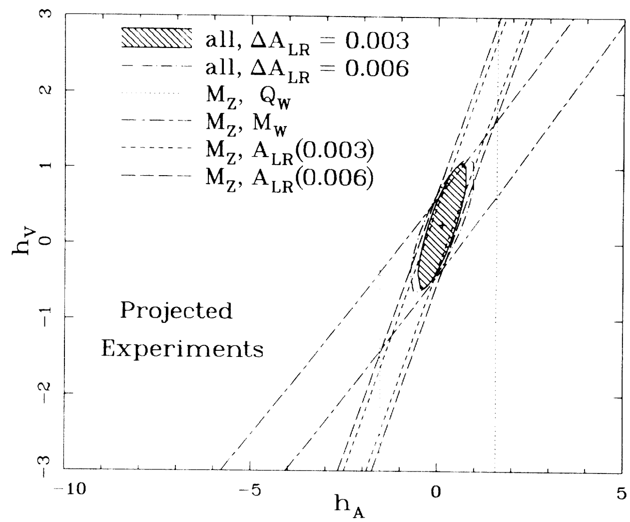

FIG. 2. 90%-C.L. regions in (h_A, h_V) from projected future experiments: (Q_W, M_Z), (M_W, M_Z), (A_{LR}, M_Z), and *(all data)*, where A_{LR} is the polarization asymmetry of $e^+e^- \to \mu^+\mu^-$ at the Z pole. $\Delta Q_W/Q_W = 1\%$, $\Delta M_W = 100$ MeV, and $\Delta A_{LR} = 0.003$ (0.006) for LEP (SLAC Linear Collider) are assumed.

combination of A_{LR}, M_W, and Q_W will lower the error in h_V to about ± 0.4, while not essentially changing the uncertainty in h_{AZ}.

Within the context of the standard model $(h^{\rm NS}=0)$ the upper limit on h_V reproduces the limits $m_t < 172$ GeV (190 GeV) at 90% (95%) C.L. for $M_H = M_Z$, or $m_t < 196$ GeV (212 GeV) for $M_H = 1$ TeV. There are no useful constraints on M_H.

Our analysis places important restrictions on new physics even with the loosest conditions, $m_t = M_H = M_Z$. (All limits are quoted at 90% C.L.) Each member of an isomultiplet of degenerate heavy fermions contributes to h_{AZ}:

$$h_{AZ} = (L_3 - R_3)^2 N_C / 3\pi , \tag{14}$$

where L_3, R_3 are the left- and right-handed couplings to the isocurrent J_3, and $N_C = 1$ (3) for leptons (quarks) is the color factor. For ordinary fermions, $L_3 = I_{3L}$ and $R_3 = 0$; i.e., $1/6\pi$ for a lepton doublet and $2/3\pi$ for a complete generation. The upper limit on h_{AZ} then allows no more than four degenerate heavy generations. A split fermion doublet contributes

$$h_{AZ} = \frac{N_C}{6\pi}\left[1 - 4\bar{Q}\ln\left(\frac{m_u}{m_d}\right)\right], \tag{15}$$

for $m_u \geq m_d \gg M_Z$, where $\bar{Q} = \frac{1}{2}(Q_u + Q_d)$. Mass splittings are constrained by h_V. The general contribution to h_V from a split doublet of scalars or fermions is[1]

$$\alpha h_V = \frac{N_C G_F}{8\sqrt{2}\pi^2}\left[m_u^2 + m_d^2 - \frac{2m_u^2 m_d^2}{m_u^2 - m_d^2}\ln\left(\frac{m_u^2}{m_d^2}\right)\right]. \tag{16}$$

Let $\alpha h_V \equiv N_C G_F \Delta m^2 / 8\sqrt{2}\pi^2$. Then $\Delta m^2 < [(250\ {\rm GeV})/N_C^{1/2}]^2$.

Supersymmetry[18] (SUSY) and technicolor[19] (TC) theories are also constrained by our analysis. The only significant source of nondecoupled radiative corrections in SUSY arises from the top/bottom squark doublet:

$$\alpha h_V \cong \frac{3G_F}{8\sqrt{2}\pi^2} m_t^2 \times \begin{cases} 1, & m_{3/2} < m_t , \\ O(m_t/m_{3/2}), & m_{3/2} > m_t , \end{cases} \tag{17}$$

where $m_{3/2}$ is the soft SUSY-breaking scale. If $m_{3/2} \gg m_t$, the stop decouples; for $m_{3/2} < m_t$, the SM upper limit on the top-quark mass is reduced by $\sqrt{2}$; i.e., $m_t < 122$ GeV. The gauginos and Higgsinos decouple from the h's in the heavy limit as they become degenerate and their gauge couplings become pure vector [Eq. (14)].[18] Simple TC theories have been presented in Refs. 3 and 4, using "scaled-up" QCD analyses based on sum rules and chiral perturbation theory. In Ref. 3,

$$h_{AZ} \simeq N_{\rm TG} \times \begin{cases} 0.4 + 0.08(N_{\rm TC} - 4), \\ 2.1 + 0.4(N_{\rm TC} - 4), \end{cases} \tag{18}$$

where $N_{\rm TC}$ and $N_{\rm TG}$ are the number of technicolors and technigenerations; and the first and second lines refer to doublets of techniquarks and full technigenerations, respectively. Even a single full TG is ruled out at 90% C.L., while $N_{\rm TG} < 2.5$ for doublets of techniquarks. The constraint on h_{AZ} clearly favors theories with the fewest possible number of technicolors and technigenerations. As stressed in Refs. 3 and 4, however, such scaled-up QCD models are not viable, while the radiative corrections from realistic ("walking") technicolor theories have not been reliably calculated.[19] The isospin breaking in TC is usually stated in terms of the pseudo-Goldstone decay constants:

$$\alpha h_V = (F_{\pm}^2 - F_0^2)/F_0^2 < 0.007 \,, \tag{19}$$

where $F_0 = 246$ GeV.

Our formulation can easily be extended to allow higher-dimensional Higgs representations with symmetry-breaking vacuum expectation values (VEVs). Define

$$\rho_0 \equiv \frac{\sum_i (L_i^2 - L_{3i}^2 + L_{3i}) |\langle \phi_i \rangle|^2}{\sum_i 2 L_{3i}^2 |\langle \phi_i \rangle|^2} \,, \tag{20}$$

where L_i is the weak isospin of the VEV $\langle \phi_i \rangle$. Then the expressions for Γ_Z and $\mathcal{M}_{\rm NC}$ in Eqs. (10a) and (11) are multiplied by ρ_0, while the right-hand side of Eq. (9a) is divided by ρ_0. Clearly, ρ_0 and h_V always occur in the combination $\rho_0/(1-\alpha h_V)$, so that the limits on h_V translate into $0.991 < \rho_0/(1-\alpha h_V) < 1.007$. This universal combination of ρ_0 and h_V means that the effects of $\rho_0 \neq 1$ and $h_V \neq 0$ are not separable with gauge-boson self-energy measurements alone. Even with $\rho_0 \neq 1$, however, in the special case of large top-quark mass, there is still an upper limit of $m_t < 350$ GeV (90% C.L.) that arises from the logarithmic top-quark-mass dependence of $h_{AZ,W}$ [Eqs. (8b) and (8c)] and of the vertex corrections to the $Z \to b\bar{b}$ partial width.[20]

In closing, we stress that the formalism of Eqs. (1)-(4),(6), (9)-(11) is specific to the SU(2)×U(1) gauge group and cannot be applied to the effect of new gauge bosons (e.g., a Z') without fundamental modifications. In models with extra neutral gauge bosons, the Z-Z' mixing lowers the predicted value of M_Z and mimics a negative value for h_{AZ}, an attractive possibility given the central value of h_{AZ} obtained in Eq. (13). Nevertheless, the effect of a Z' in Γ_Z and $\mathcal{M}_{NC}$ cannot be so described, because of the Z-Z' mixing in the Z couplings and the Z' exchange. Analysis of new-gauge-boson interactions requires additional independent parameters in the global fit, the effects of which cannot be subsumed into the h's.

This work was supported by the Aspen Center for Physics and by the Department of Energy Contract No. DE-AC02-76-ERO-3071.

Note added.—A new analysis by Altarelli and Barbieri of electroweak radiative corrections due to heavy physics has appeared recently, with a treatment similar to that given above.[21]

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[11]Equations (8a), (8b), and (8c) are derived using the scheme of Ref. 6. In the $\overline{\rm MS}$ scheme of Ref. 10, the equivalent h's in Eqs. (9a) and (9b) have a different logarithmic dependence on m_t: $(2/3\pi)(\frac{3}{4} + 2s_R^2 - 16 s_R^2 c_R^2)$ (instead of $-1/6\pi$) in Eq. (8b) and $(2/\pi)(-8s_R^2/9 + \frac{1}{2})$ (instead of $+1/3\pi$) in Eq. (8c), where $c_R^2 = 1 - s_R^2$. This difference arises from the different dependences of $\sin^2\theta(M_Z^2)_{\overline{\rm MS}}$ and $s_*^2(M_Z^2)$ on the top-quark threshold.

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[17]Most types of new physics yield positive contributions to the h's. If one imposes the extra restriction $h_{V,A} > 0$, then one obtains the weaker limits $h_V < 1.3$, $h_{AZ} < 2.1$, $h_{AW} < 3.7$ at 90% C.L.

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