THEORETICAL ASPECTS OF HIGGS PHYSICS†

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We review the present status of Higgs physics within the standard model and its extensions. First, we briefly summarize the current experimental exclusion limits from the direct searches with LEP1 and the Tevatron, and assess the discovery potential of LEP2. Then, we report the mass bounds resulting from global fits to the latest electroweak precision data, and compile various theoretical principles which lead to restrictions—or even determinations—of the Higgs-boson masses. Perturbative upper bounds are discussed in some detail. Finally, we survey the recent progress in the computation of higher-order radiative corrections to the $b\bar{b}$ decay width of the standard-model Higgs boson.

1 Introduction

The SU(2)$_L$×U(1)$_Y$ structure of the electroweak interactions has been consolidated by an enormous wealth of experimental data during the past three decades. The canonical way to generate masses for the fermions and intermediate bosons without violating this gauge symmetry in the Lagrangian is by the Higgs mechanism of spontaneous symmetry breaking. In the minimal standard model (SM), this is achieved by introducing one complex SU(2)$_L$-doublet scalar field $\Phi$ with $Y = 1$. The three massless Goldstone bosons which emerge via the electroweak symmetry breaking are eaten up to become the longitudinal degrees of freedom of the $W^\pm$ and $Z$ bosons, i.e., to generate their masses, while one CP-even Higgs scalar boson $H$ remains in the physical spectrum. The Higgs potential contains one mass and one self-coupling. Since the vacuum expectation value (VEV) is fixed by the relation $v = 2M_W/g$, where $g$ is the SU(2)$_L$ gauge coupling, there remains one free parameter in the Higgs sector, namely, $M_H$.

A phenomenologically interesting extension of the SM Higgs sector that keeps the electroweak $\rho$ parameter at unity in the Born approximation, is obtained by adding a second complex SU(2)$_L$-doublet scalar field $\Phi_2$ with $Y = -1$. This leads to the two-Higgs-doublet model (2HDM). After the three Goldstone bosons have been eliminated, there remain five physical Higgs scalars: the neutral CP-even $h^0$ and $H^0$ bosons, the neutral CP-odd $A^0$ boson, and the charged $H^\pm$ pair. The most general CP-conserving Higgs potential for the 2HDM contains three masses and four self-couplings. Subtracting the $M_H$ constraint on the two VEV’s $v_1$ and $v_2$, one is left with six free parameters, which are usually taken to be $m_{h^0}$, $m_{H^0}$, $m_{A^0}$, $m_{H^\pm}$, $\tan \beta = v_2/v_1$, and the weak mixing angle $\alpha$ which relates the weak and mass eigenstates of $h^0$ and $H^0$. According to the Glashow-Weinberg theorem, flavour-changing neutral currents may be avoided if all fermions with the same electric charge couple to the same Higgs doublet. In the 2HDM of type II, the up-type/down-type fermions couple to the Higgs doublet with $Y = \pm 1$.

The Higgs sector of the minimal supersymmetric extension of the SM (MSSM) consists of a 2HDM of type II. Since the Higgs self-couplings are then determined by the gauge couplings, there are only two free parameters in the MSSM Higgs sector, which are conveniently chosen to be $m_{h^0}$ and $\tan \beta$. However, a large number of new masses, couplings, and mixing angles connected with the supersymmetric partners enter via loop effects. In the supergravity-inspired MSSM, all these degrees of freedom may be related to just five parameters at the GUT scale: the Higgs mixing $\mu$, the sfermion mass $m_0$, the gaugino mass $m_{1/2}$, the trilinear sfermion-Higgs coupling $A$, and the bilinear Higgs coupling $B$. After the renormalization-group (RG) evolution down to the electroweak scale, the heavy sfermions are approximately degenerate, with mass $M_S$. The following three scenarios are often considered in the literature: (i) no mixing: $A = 0$, $|\mu| \ll M_S$; (ii) typical mix-
Table 1: 95% CL $M_H$ lower bounds (in GeV) from LEP1.

| Collab. | Years     | $10^9 \, qq$ | $M_{H}^{\text{min}}$ |
|---------|-----------|--------------|---------------------|
| ALEPH   | 89–95     | 4.5          | 63.9                |
| DELPHI  | 91, 92, 94, 95 | 3.1          | 58.3                |
| L3      | 91–94     | 3.1          | 60.2                |
| OPAL    | 89–95     | 4.4          | 59.6                |

Table 2: 95% CL lower bounds on $m_{h^0}$, $m_{A^0}$, and $m_{H^\pm}$ (in GeV) from LEP1.

| Collab. | $m_{h^0}^{\text{min}}$ | $m_{A^0}^{\text{min}}$ | $m_{H^\pm}^{\text{min}}$ |
|---------|------------------------|-------------------------|---------------------------|
| ALEPH   | 45.5                   | 45.0                    | 41.7                      |
| DELPHI  | 45.4                   | 45.2                    | –                         |
| L3      | 42.0                   | 23.0                    | 42.8                      |
| OPAL    | 44.3                   | 23.5                    | 44.1                      |

The $A = -\mu = M_S$; and (iii) maximal mixing: $A = \sqrt{3}M_S$, $|\mu| \ll M_S$.

2 Direct Higgs searches

2.1 LEP1 and Tevatron

The SM Higgs boson was searched at LEP1 via Bjorken’s process $e^+e^- \rightarrow Hf\bar{f}$ by looking for a pair of acoplanar jets together with missing energy in the $H\nu\bar{\nu}$ channel, or for a hadronic event with an energetic lepton pair in the $Hl^+l^-$ ($l = e, \mu$) channel. The extraction by ALEPH of the 95% CL lower bound on $M_H$ is illustrated in Fig. 1.

The present LEP1 bounds are summarized in Table 1.

The $h^0$ and $A^0$ bosons of the 2HDM or MSSM could have been produced at LEP1 via Higgs-strahlung $e^+e^- \rightarrow h^0 f\bar{f}$, where the $h^0$ boson is radiated off the resonant $Z$ boson, associated production $e^+e^- \rightarrow h^0A^0$, and the Yukawa process $e^+e^- \rightarrow h^0/A^0f\bar{f}$ ($f = b, \tau$), where the $h^0$ and $A^0$ bosons couple to the fermion line. The $h^0Z^*$ and $h^0A^0$ channels are complementary, in the sense that their cross sections are proportional to $\sin^2(\beta - \alpha)$ and $\cos^2(\beta - \alpha)$, respectively. In the framework of the MSSM with $M_S = 1$ TeV and no squark mixing, DELPHI has excluded the shaded regions in the $(m_{h^0}, \tan \beta)$ plane of Fig. 2. The current 95% CL lower bounds on $m_{h^0}$ and $m_{A^0}$ in the MSSM obtained by the four LEP1 experiments are listed in Table 2.

In the more general context of the 2HDM, the $h^0$ ($A^0$) boson could have been missed in the $h^0Z^*$ and $h^0A^0$ channels, if $\sin^2(\beta - \alpha) \approx 0$ (1) and (or) $m_{h^0} + m_{A^0} > M_Z$. Then, the Yukawa process would be dominant. The MSSM scenario with $m_{A^0} \ll M_Z$ and $\tan \beta \gg 1$ was considered particularly interesting since it could par-
tially explain the $R_b$ anomaly, which was then still existing. ALEPH$^b$ searched for the $h^0$ and $A^0$ bosons of the 2HDM using the Yukawa process with $b\bar{b}b\bar{b}$, $\tau^+\tau^-b\bar{b}$, $\tau^+\tau^\mp\tau^-\tau^\mp$, and $\tau^+\tau^-X$ final states, where $X$ is a system with low charged multiplicity. The $b\bar{b}b\bar{b}$ channel turned out to be less efficient in ruling out the small-$m_{A^0}$ large-$\tan\beta$ scenario than anticipated$^5$. It is amusing to observe that, in the 2HDM with $\alpha \approx \beta$, the most stringent $m_{h^0}$ lower bound presently comes from HERA.$^6$

The future E821 experiment at BNL is expected to decrease the experimental error on $(g-2)_\mu$ by a factor of twenty or more and will thus lead to interesting $m_{h^0}$ and $m_{A^0}$ lower bounds via loop effects.$^7$

Searching for the $H^\pm$ bosons of the 2HDM via pair production $e^+e^- \rightarrow H^+H^-$, the LEP1 experiments$^8$ could almost exclude the entire $m_{H^\pm}$ range allowed by kinematics (see Table$^8$). CDF$^9$ has looked for $H^{\pm}$ bosons in the decay products of $t\bar{t}$ pairs, and has greatly increased the LEP1 lower bound on $m_{H^\pm}$ for $\tan\beta \gtrsim 60$ (see Fig. 3). If $\tan\beta \gg 1$, then the cascade $t \rightarrow H^+b \rightarrow \tau^+\nu_t\tau^-\bar{\nu}_t\bar{b}$ is more likely to happen than $t \rightarrow W^+b \rightarrow \tau^+\nu_t\tau^-\bar{\nu}_t\bar{b}$ and leads to a larger missing transverse energy.$^{10}$

2.2 LEP2

The theoretical and experimental aspects of Higgs phenomenology at LEP2 have recently been summarized in a comprehensive report.$^{11}$ Higgs-strahlung is the dominant production mechanism of the SM Higgs boson at LEP2. In Fig. 4, its cross section is shown as a function of $M_H$, for three values of CM energy. $^{12}$ Electroweak radiative corrections$^{13}$ are included here.

![Figure 3: Region in the ($\tan\beta, m_{H^\pm}$) plane excluded at 95% CL by CDF.](image)

![Figure 4: Cross section of Higgs-strahlung at LEP2 as a function of $M_H$.](image)

![Figure 5: $H\nu\bar{\nu}$ production via Higgs-strahlung, $W^+W^-$ fusion, and their interference in the threshold region at LEP2.](image)

The minimum luminosity $L_{\text{min}}$ needed per experiment for a combined 5$\sigma$ discovery or a 95% CL exclusion of the SM Higgs boson under realistic LEP2 conditions$^{14}$ is shown as a function of $M_H$ in Fig. 6. We see that, at $\sqrt{s} = 192$ GeV, $L_{\text{min}} = 150 \text{ pb}^{-1}$ is sufficient to discover the SM Higgs boson with $M_H \lesssim 95$ GeV. On the other hand, a 95 GeV Higgs boson can be excluded at
95% CL with $L_{\text{min}}$ as low as 33 pb$^{-1}$, while, with $L_{\text{min}} = 150 \text{ pb}^{-1}$, the $M_H$ range way up to 98 GeV may be excluded.

Figure 6: LEP2 exclusion and discovery limits for the SM Higgs boson.

The 5$\sigma$ discovery and 95% CL exclusion limits in the ($m_{h^0}, \tan \beta$) plane of the MSSM that may be reached at LEP2 with a luminosity of 150 pb$^{-1}$ per experiment may seen from Fig. 6. Here, $M_t = 175$ GeV, $M_S = 1$ TeV, and typical squark mixing as described in the Introduction are assumed. Comparing Figs. 2 and 6, we observe that LEP2 will be able to close the low-$\tan \beta$ region which was not covered by LEP1.

Figure 7: LEP2 exclusion and discovery limits in the ($m_{h^0}, \tan \beta$) plane.

### 3 Higgs mass bounds

#### 3.1 Global fits

Experimental precision tests of the standard electroweak theory are sensitive to the Higgs boson via quantum corrections. A global fit to the $e^+e^-$ data on the Z-boson observables from ALEPH, DELPHI, L3, OPAL, and SLD, the $p\bar{p}$ data on $M_W$ from CDF, D0, and UA2, the $\nu N$ data on $\sin^2\theta_W$ from CCFR, CDHS, and CHARM, and the Tevatron data on $M_t$ from CDF and D0 which were presented during this conference yields $M_H = 149^{+14}_{-82}$ GeV with $\chi^2_{\text{min}}$/d.o.f. = 19/14 $\approx$ 1.36. The resulting $M_H$ distribution of $\Delta \chi^2 = \chi^2 - \chi^2_{\text{min}}$ is shown in Fig. 8, where the shaded band represents the estimated theoretical error due to missing higher-order corrections. From Fig. 8, one may read off an 95% CL upper bound on $M_H$ at 550 GeV. The electroweak precision data have reached a quality which even allows for global fits within the MSSM.

Figure 8: $M_H$ dependence of $\Delta \chi^2 = \chi^2 - \chi^2_{\text{min}}$ resulting from a global fit to the latest experimental data.

#### 3.2 Theoretical principles

There have been various attempts to bound or even determine $M_H$ from first principles.
Naturalness

By requiring that the Higgs-boson mass countertern, i.e., the difference between the bare and renormalized masses, be devoid of quadratic ultraviolet divergences, Veltman derived the mass relation,

$$\sum_f N_f M_f^2 = \frac{3}{2} M_W^2 + \frac{3}{4} M_Z^2 + \frac{3}{4} M_H^2,$$

where $N_f = 1$ (3) for leptons (quarks). Naively inserting the known pole masses, this leads to $M_H \approx 164$ GeV if $M_t = 175$ GeV.

Noncommutative geometry

It is possible to construct the SM on the basis of the graded Lie algebra SU(2|1), which contains the SU(2) x U(1) V Lie algebra in its even part. The essential ingredient is an algebraic superconnection which incorporates both the gauge and Higgs fields and whose curvature automatically generates the spontaneously broken realization of the SM. This leads to a geometrical interpretation of the Higgs mechanism. As a by-product, one obtains the mass relation $M_H = \sqrt{2} M_W$. Taking this as an initial condition at the GUT scale and performing a one-loop RG analysis, one finds the physical Higgs-boson mass to be $M_H = 164$ GeV if $M_t = 175$ GeV.

Triviality and vacuum stability

Triviality and vacuum stability are features connected with the RG-improved effective potential $V_{\text{eff}} = -m^2(\mu)|\Phi|^2 + \lambda(\mu)|\Phi|^4/2$. The dependence of the quartic self-coupling $\lambda(\mu)$ on the renormalization scale $\mu$ is determined by the RG equations. Roughly speaking, the triviality upper bound (vacuum-stability lower bound) on $M_H$ follows from the requirement that $\lambda(\mu)$ stay finite (positive) for all $\mu < \Lambda$, where $\Lambda$ is the cutoff beyond which new physics operates. Assuming the SM to be valid up to the GUT scale $\Lambda = 10^{16}$ GeV, one thus obtains $130$ GeV $\lesssim M_H \lesssim 180$ GeV. In turn, should the SM Higgs boson be discovered at LEP2, new physics is expected to appear below $\Lambda = 10$ TeV.

Metastability

Depending on $M_H$ and $M_t$, $V_{\text{eff}}$ at finite (or zero) temperature may have a deep stable minimum at $\langle \Phi_{\text{min}} \rangle \gg G_F^{-1/2}$, i.e., the physical electroweak minimum may just be metastable. Then, an absolute lower bound on $M_H$ follows from the condition that the probability, normalized w.r.t. the expansion rate of the Universe, for the decay of the metastable vacuum by thermal fluctuations (or quantum tunneling) be negligibly small. The $M_H$ bounds thus obtained are somewhat below the usual stability bounds.

Multiple point criticality principle

It has been argued that a mild form of locality breaking in quantum gravity due to baby universes, which are expected to render coupling constants dynamical, leads to the realization of the so-called multiple-point criticality principle in Nature. According to this principle, Nature should choose coupling constants such that the vacuum can exist in degenerate phases, i.e., $V_{\text{eff}}(\Phi_{\text{min},1}) = V_{\text{eff}}(\Phi_{\text{min},2})$. In order that the dynamical mechanism be relevant, these authors also require a strong first-order phase transition between the two vacua, implemented by $\langle \Phi_{\text{min},2} \rangle \approx M_{\text{Planck}} = 2 \times 10^{19}$ GeV. Via the usual RG analysis, these two assumptions then lead to a simultaneous prediction of $M_t$ and $M_H$, namely, $M_t = (173 \pm 4)$ GeV and $M_H = (135 \pm 9)$ GeV.

Perturbation-theory breakdown

An attractive way of constraining $M_H$ from above is to require that the Higgs sector be weakly interacting, so that perturbation theory is meaningful. The resulting bounds depend somewhat on the considered process and the precise definition of perturbation-theory breakdown, but they are independent of assumptions concerning the scale $\Lambda$ of new physics. At one and two loops, the leading high-$M_H$ corrections to physical observables related to Higgs-boson production or decay are of $\mathcal{O}(G_F M_H^2)$ and $\mathcal{O}(G_F^2 M_H^4)$, respectively. For $M_H$ increasing, the $\mathcal{O}(G_F^2 M_H^4)$ corrections eventually exceed the $\mathcal{O}(G_F M_H^2)$ ones in size, so that the perturbative expansions in $G_F M_H^2$ cease to usefully converge. The values $M_H^{\text{max}}$ where this happens may be used to define a perturbative upper bound on $M_H$. 

\[ 5 \]
As first examples, the Higgs decays to pairs of fermions and intermediate bosons have recently been studied through $\mathcal{O}(G_F^2 M_H^2)$. This task may be greatly simplified in the limit of interest, $M_H \gg 2M_Z$, through the use of the Goldstone-boson equivalence theorem. This theorem states that the leading high-$M_H$ electroweak contribution to a Feynman diagram may be calculated by replacing the intermediate bosons $W^{\pm}$ and $Z$ with the respective would-be Goldstone bosons $w^{\pm}$ and $z$ of the symmetry-breaking sector. In this limit, the gauge and Yukawa couplings may be neglected against the Higgs self-coupling. By the same token, the Goldstone bosons may be taken to be massless, and the fermion loops may be omitted. The resulting correction factors $K_f$ for $\Gamma (H \rightarrow f \bar{f})$ is independent of the fermion flavour $f$. Similarly, $\Gamma (H \rightarrow W^+ W^-)$ and $\Gamma (H \rightarrow ZZ)$ receive the same correction factor $K_V$. In the on-mass-shell renormalization scheme, the results read

$$K_f = 1 + \lambda \left( 13 - 2\pi \sqrt{3} \right) + \lambda^2 \left[ 12 - 169\pi \sqrt{3} + 170\zeta(2) - 252\zeta(3) + 12 \left( 13\pi + 19\sqrt{3} \right) \right. \times C_{\lambda} \left( \frac{\pi}{3} \right) \left. \right]$$

$$\approx 1 + 11.1\% \left( \frac{M_H}{\text{TeV}} \right)^2 - 8.9\% \left( \frac{M_H}{\text{TeV}} \right)^4,$$

$$K_V = 1 + \lambda \left( 19 - 6\pi \sqrt{3} - 10\zeta(2) \right) + 62.0 \lambda^2$$

$$\approx 1 + 14.6\% \left( \frac{M_H}{\text{TeV}} \right)^2 + 16.9\% \left( \frac{M_H}{\text{TeV}} \right)^4,$$

where $\lambda = (G_F M_H^2 / 16\pi^2 \sqrt{2})$, $\zeta$ is Riemann's zeta function, and $C_{\lambda}$ is Clausen's integral. The $\mathcal{O}(G_F M_H^2)$ terms in Eq. (2) have been known for a long time. $K_f$ and $K_V$ are displayed as functions of $M_H$ in Fig. 8, from where we may read off the $M_H^{\text{max}}$ values 1114 GeV and 930 GeV, respectively. The nonperturbative value of $K_V$ at $M_H = 727$ GeV may be extracted from a recent lattice simulation of elastic $\pi \pi$ scattering in the framework of the four-dimensional $O(4)$-symmetric nonlinear $\sigma$ model in the broken phase, where the $\sigma$ resonance was observed.

The study of the $\mu$ dependence of Higgs-boson observables in the $\overline{\text{MS}}$ renormalization scheme provides another aspect of perturbation-theory breakdown in the Higgs sector. If perturbation theory is valid, the $\mu$ dependence should be reduced as higher-order corrections are included. While this empirical rule is satisfied for $\Gamma (H \rightarrow W^+ W^-)$ at $M_H = 400$ GeV, it is clearly violated at $M_H = 700$ GeV as may be seen from Fig. 9.

### 4 Radiative corrections to $\Gamma (H \rightarrow b \bar{b})$

The SM Higgs boson with intermediate mass $M_H \lesssim 2M_W$ decays dominantly to $b \bar{b}$ pairs. The radiative corrections to the partial width of this decay involve two very different scales, namely, $M_H$ and $M_t$. It is therefore convenient to treat this process in the framework of a $n_f = 5$ effective Yukawa Lagrangian, i.e., to integrate out the top quark. This leads to a RG-improved formulation,
which provides a natural separation of the $n_f = 5$ QCD corrections at scale $\mu = M_H$ and the top-quark-induced $n_f = 6$ corrections at scale $\mu = M_t$. One thus obtains the following structure:

$$
\Gamma_{bb} = \Gamma^\text{Born}_{bb} \left[ (1 + \Delta_b^\text{QED}) (1 + \Delta_b^\text{weak} \big|_{x_t=0}) \times (1 + \Delta_b^\text{QCD}) (1 + \Delta_b^\text{t}) + \Xi_b^\text{QCD} \Xi_b^\text{t} \right].
$$

If the Born formula

$$
\Gamma^\text{Born}_{bb} = \frac{3G_FM_H m_b^2}{4\pi} \left( 1 - \frac{4m_b^2}{M_H^2} \right)^{3/2}
$$

is written with the QED and QCD $\overline{\text{MS}}$ mass evaluated with $n_f = 5$ quark flavours at scale $\mu = M_H$, $m_b^5(M_H)$, then $\Delta_b^\text{QED}$ and $\Delta_b^\text{QCD}$ are finite for $m_b = 0$ and read

$$
\Delta_b^\text{QED} = \frac{17}{3} \alpha + a_5^2 \left( \frac{8851}{144} - \frac{47}{6} \zeta(2) - \frac{97}{6} \zeta(3) \right)
+ a_3^2 \left( \frac{34873057}{46656} - \frac{10225}{54} \zeta(2)
- \frac{80095}{216} \zeta(3) - \frac{25}{6} \zeta(4) + \frac{1945}{36} \zeta(5) \right)
= 5.66667 a_5 + 29.14671 a_5^2 + 41.75760 a_5^3,
$$

where $a_5 = \alpha_s^5(M_H)/\pi$. By means of scale optimization according to the principles of fastest apparent convergence (FAC) or minimal sensitivity (PMS), the coefficient of $a_5^2$ may be estimated to be $-981$. In the limit $M_H \ll 2 M_W$, the weak correction, with the $\mathcal{O}(G_F M_H^2)$ term stripped off, reads

$$
\Delta_b^{\text{weak}} \big|_{x_t=0} = \frac{G_F M_H^2}{8\pi^2 \sqrt{2}} \left( \frac{1}{6} - 7 \frac{c_w^2}{3} - 16 \frac{c_w^4}{3} \right)
+ \frac{3}{2} \frac{c_w^2}{s_w} \ln c_w,
$$

where $c_w^2 = 1 - s_w^2 = M_W^2/M_Z^2$. The leading contributions due to the $b\bar{b}$ and $b\bar{b}g$ cuts of the double-triangle diagrams where the top quark propagates in one of the triangles are contained in $\Xi_b^\text{QCD}$

$$
\Xi_b^\text{QCD} = a_5 \left( -\frac{76}{3} + 8\zeta(2) - \frac{4}{3} \ln^2 \frac{m_b^2}{M_H^2} \right).
$$

The would-be mass singularity proportional to $\ln^2(m_b^2/M_H^2)$ cancels if the $b\bar{b}(g)$ and $gg$ decay channels are combined. The leading top-quark-induced corrections are concentrated in

$$
\Delta_b^t = a_6 \left( \frac{5}{9} + \frac{2}{3} L \right) + x_t \left[ 1 + a_6 \left( -\frac{4}{3} - 4\zeta(2) \right) 
+ a_6^2 \left( -59.1626 + \frac{2}{3} L \right) \right],
\Xi_b^t = a_6 \left( -\frac{1}{12} - \frac{1}{12} x_t \right),
$$

where $\mu_t = m_t^6(\mu_t)$, $a_6 = \alpha_s^6(\mu_t)/\pi$, $L = \ln(\mu_t^2/M_H^2)$, and $x_t = (G_F \mu_t^2/8\pi^2 \sqrt{2})$. The $L$-dependent terms may be resummed by exploiting the RG-invariance of the energy-momentum tensor.

The QCD correction to $\Gamma(H \to gg)$ includes a contribution due to the $b\bar{b}g$ final state, which may also be interpreted as a $\mathcal{O}(a_s^2 M_H^2/m_t^2)$ correction to $\Gamma(H \to b\bar{b})$.

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**References**

1. S.L. Glashow and S. Weinberg, *Phys. Rev.* D **15**, 1958 (1977).
2. D. Buskulic et al. (ALEPH Collaboration), Pa13–026, *Phys. Lett.* B **384**, 427 (1996); M. Acciarri et al. (L3 Collaboration), Pa11–016, *Phys. Lett.* B **385**, 454 (1996).
3. C. Martínez Rivero et al. (DELPHI Collaboration), Pa07–095; P. Igo-Kemenes et al. (OPAL Collaboration), Pa13–004.
4. P. Janot, *Nucl. Phys.* B (Proc. Suppl.) **38**, 264 (1995); O. Adriani et al. (L3 Collaboration), *Z. Phys.* C **57**, 355 (1993).
5. J. Kalinowski and M. Krawczyk, *Phys. Lett.* B **361**, 66 (1995); *Acta Phys. Pol.* B **27**, 961 (1996).
6. J.D. Wells, C. Kolda, and G.L. Kane, *Phys. Lett.* B **338**, 219 (1994); *Phys. Rev. Lett.* **76**, 869 (1996).
7. ALEPH Collaboration, Pa13–026.
8. A.C. Bawa and M. Krawczyk, *Phys. Lett.* B **357**, 637 (1995).
9. L. Roberts, in these proceedings.
10. M. Krawczyk and J. Zochowski, Report Nos. IFT 15/96 and hep-ph/9608321 (August 1996).

11. D. Decamp et al. (ALEPH Collaboration), Phys. Rep. 216, 253 (1992); B. Lindemann and A. Sopczak, L3 Note Nos. 1966 and hep-ph/9607327 (June 1996); G. Alexander et al. (OPAL Collaboration), Phys. Lett. B 370, 174 (1996).

12. F. Abe et al. (CDF Collaboration), Phys. Rev. D 54, 735 (1996); Pa13–025.

13. M. Carena, P.M. Zerwas (conveners) et al., in Physics at LEP2, edited by G. Altarelli, T. Sjöstrand, and F. Zwirner, CERN Yellow Report No. 96–01, Vol. 1, p. 351.

14. E. Gross, B.A. Kniehl, and G. Wolf, Z. Phys. C 63, 417 (1994); 66, 321(E) (1995).

15. J. Fleischer and F. Jegerlehner, Nucl. Phys. B 216, 469 (1983); F.A. Berends and R. Kleiss, Nucl. Phys. B 260, 32 (1985); B.A. Kniehl, Z. Phys. C 55, 605 (1992); A. Denner, J. Kühnbeck, R. Mertig, and M. Böhm, Z. Phys. C 56, 261 (1992).

16. W. Kilian, M. Krämer, and P.M. Zerwas, Phys. Lett. B 373, 135 (1996).

17. J. Alcaraz et al. (LEP Electroweak Working Group and SLD Heavy Flavor Group), CERN Report No. LEPEWWG/96–02 (30 July 1996).

18. W. de Boer, A. Dabelstein, W. Hollik, W. Mösle, and U. Schwickerath, Report Nos. Pa11–001, IEKP–KA/96–08, KA–TP–18–96, and hep-ph/9609208 (August 1996).

19. M. Veltman, Acta Phys. Pol. B 12, 437 (1981).

20. R. Häußling, N.A. Papadopoulos, and F. Scheck, Phys. Lett. B 260, 125 (1991); R. Coquereaux, R. Häußling, N.A. Papadopoulos, and F. Scheck, Int. J. Mod. Phys. A 7, 2809 (1992).

21. Y. Okumura, Report Nos. CHUBU9607 and hep-ph/9608208 (October 1996).

22. N. Cabibbo, L. Maiani, G. Parisi, and R. Petronzio, Nucl. Phys. B 158, 295 (1979); M. Lindner, Z. Phys. C 31, 295 (1986).

23. G. Altarelli and G. Isidori, Phys. Lett. B 337, 141 (1994); J.A. Casas, J.R. Espinosa, and M. Quirós, Phys. Lett. B 342, 171 (1995), Phys. Lett. B 382, 374 (1996).

24. J.R. Espinosa and M. Quirós, Phys. Lett. B 353, 257 (1995).

25. C.D. Froggatt and H.B. Nielsen, Report Nos. GUTPA–96–7–2, Pa08–012, and hep-ph/9607302 (July 1996); see also R. Rodenberg, Pa12–001 (July 1996).

26. M. Veltman, Phys. Lett. B 70, 253 (1977).

27. L. Durand, B.A. Kniehl, and K. Riesselmann, Phys. Rev. Lett. 72, 2534 (1994); 74, 1699(E) (1995); A. Ghinculov, Phys. Lett. B 337, 137 (1994); 346, 426(E) (1995); V. Borodulin and G. Jikia, Report No. Freiburg–THEP 96/19 and hep-ph/9609447 (September 1996).

28. A. Ghinculov, Nucl. Phys. B 455, 21 (1995); A. Frink, B.A. Kniehl, D. Kreimer, and K. Riesselmann, Pa07–081, Phys. Rev. D 54, 4548 (1996).

29. J.M. Cornwall, D.N. Levin, and G. Tiktopoulos, Phys. Rev. D 10, 1145 (1974); 11, 972(E) (1975); C.E. Vayonakis, Lett. Nuovo Cimento 17, 383 (1976).

30. M. Veltman, Acta Phys. Pol. B 8, 475 (1977); W.J. Marciano and S.S.D. Willenbrock, Phys. Rev. D 37, 2509 (1988).

31. M. Göckeler, H.A. Kastrup, J. Westphalen, and F. Zimmermann, Nucl. Phys. B 425, 413 (1994).

32. U. Nierste and K. Riesselmann, Phys. Rev. D 53, 6638 (1996).

33. K.G. Chetyrkin, B.A. Kniehl, and M. Steinhauser, Report Nos. MPI/PhT/96–65 and hep-ph/9610456 (July 1996).

34. E. Braaten and J.P. Leveille, Phys. Rev. D 22, 715 (1980); S.G. Gorishny, A.L. Kataev, S.A. Larin, and L.R. Surguladze, Mod. Phys. Lett. A 5, 2703 (1990); K.G. Chetyrkin, Report Nos. MPI/PhT/96–61 and hep-ph/9608315 (August 1996), Phys. Lett. B (in press).

35. K.G. Chetyrkin, B.A. Kniehl, and A. Sirlin, Report No. MPI/PhT/96–62.

36. B.A. Kniehl, Nucl. Phys. B 376, 3 (1992).

37. S.A. Larin, T. van Ritbergen, and J.A.M. Vermaseren, Phys. Lett. B 362, 134 (1995); K.G. Chetyrkin and A. Kwiatkowski, Nucl. Phys. B 461, 3 (1996).

38. T. Inami, T. Kubota, and Y. Okada, Z. Phys. C 18, 69 (1983); A. Djouadi, M. Spira, and P.M. Zerwas, Phys. Lett. B 264, 440 (1991).

39. A. Djouadi, M. Spira, and P.M. Zerwas, Z. Phys. C 70, 427 (1996).