Theoretical Aspects of Standard-Model Higgs-Boson Decays and Leptoproduction

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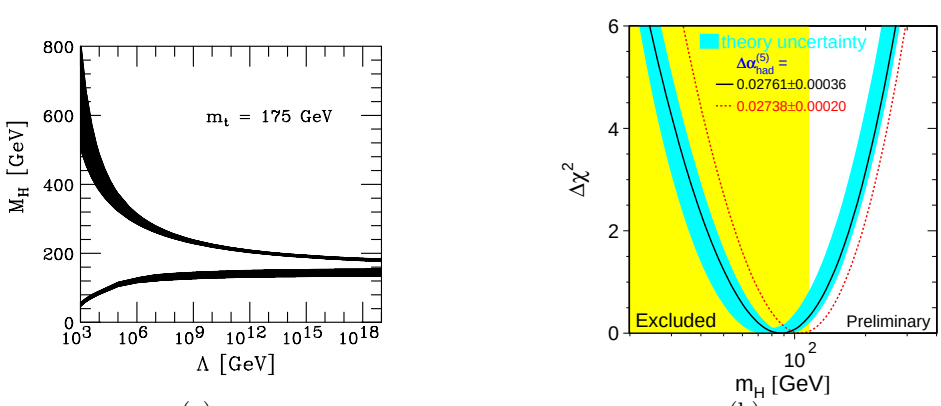
The Higgs boson is the missing link of the Standard Model of elementary particle physics. We review its decay properties and production mechanisms at a future e^+e^- linear collider and its e^-e^-, $e^\pm\gamma$, and $\gamma\gamma$ modes, with special emphasis on the influence of quantum corrections. We also discuss how its quantum numbers and couplings can be extracted from the study of appropriate final states.

I. INTRODUCTION

The $\mathrm{SU}(2)_I\times\mathrm{U}(1)_Y$ 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 $\mathrm{SU}(2)_I$-doublet scalar field Φ 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 V contains one mass and one self-coupling. Since the vacuum expectation value is fixed by the relation $v=2^{-1/4}G_F^{-1/2}\approx 246$ GeV, where G_F is Fermi's constant, there remains one free parameter in the Higgs sector, namely M_H. In fact, one has $V=\lambda H^2(v+H/2)^2+\cdots$, where $\lambda=M_H^2/(2v^2)$. The Higgs boson has the quantum numbers of the vacuum, namely electric charge $Q=0$, spin, parity, and charge conjugation $J^{PC}=0^{++}$. It has tree-level couplings to all massive particles with strengths that are determined by their masses, viz. $g_{ffH}=M_f/v$, $g_{VVH}=2M_V^2/v$, $g_{VVHH}=2M_V^2/v^2$, $g_{HHH}=6v\lambda$, and $g_{HHHH}=6\lambda$, where f denotes a generic fermion and $V=W,Z$. At a future e^+e^- linear collider (LC), an important experimental task will be to determine of the Higgs quantum numbers and couplings in order to distinguish between the minimal SM and possible extensions. In particular, the measurement of the Higgs self-couplings will allow one to directly test the Higgs mechanism.

Roughly speaking, the requirement that the running Higgs self-coupling $\lambda(\mu)$, where μ is the renormalization scale, stays finite (positive) for all values $\mu<\Lambda$, where Λ is the cutoff beyond which new physics operates, leads to the triviality upper bound (vacuum-stability lower bound) on M_H [1]. Assuming the SM to be valid up to the grand-unified-theory scale $\Lambda\approx 10^{16}$ GeV, one thus obtains $130\lesssim M_H\lesssim 185$ GeV [2] [see Fig. 1(a)]. This range comfortably lies between the lower bound on M_H from direct searches at CERN LEP2, 113 GeV, and the 95% confidence level upper bound from electroweak precision tests [3], 212 GeV, based on $\Delta\alpha_{\rm had}^{(5)}(M_Z^2)=0.02738\pm0.00020$ [4], and it is compatible with the 1σ range $76<M_H<181$ GeV resulting from the latter [3] [see Fig. 1(b)].



(a) (b)

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FIG. 1: (a) Triviality and vacuum-stability bounds on M_H [2] and (b) $\Delta\chi^2 = \chi^2 - \chi^2_{\rm min}$ as a function of M_H from fits to electroweak precision data from LEP taking into account the direct determinations of M_W and M_t [3].

It is interesting to consider a hypothetical scenario in which the Higgs boson is absent and to constrain the mass scale Λ of the new physics that would take its place. Using recent measurements of $\sin^2\theta^{\rm lept}_{\rm eff}$ and M_W [3], one finds that, in a class of theories characterized by simple conditions, the upper bound on Λ is close to or smaller than the upper bound on M_H, while in the complementary class Λ is not restricted by such considerations [5].

This review is organized as follows. In Sects. II and III, we discuss the decay properties of the Higgs boson and its main production mechanisms in e^+e^-, e^-e^-, $e^\pm\gamma$, and $\gamma\gamma$ collisions, emphasizing the influence of radiative corrections. In Sect. IV, we explain how to extract its quantum numbers and couplings from the study of final states. Sect. V contains our conclusions and a brief outlook.

II. DECAY PROPERTIES

At the tree level, the Higgs boson decays to pairs of massive fermions and gauge boson, the partial widths being

$$\begin{aligned}\Gamma\left(H\to f\bar{f}\right) &= \frac{g^2_{ffH}}{4\pi}\,\frac{N_c M_H}{2}\left(1-\frac{1}{r_f}\right)^{3/2},\\ \Gamma(H\to VV) &= \frac{g^2_{VVH}}{4\pi}\,\frac{3\delta_V}{8M_H}\left(1-\frac{4}{3}r_V+\frac{4}{3}r_V^2\right)\left(1-\frac{1}{r_V}\right)^{1/2},\end{aligned} \tag{1}$$

respectively, where $N_c = 1\,(3)$ for leptons (quarks), $\delta_{W,Z} = 2, 1$, and $r_i = M_H^2/(4M_i^2)$. If $1/4 < r_i < 1$ ($r_i < 1/4$), then one of the (both) final-state particles are forced to be off shell, so that one is dealing with three-particle (four-particle) decays [6]. The Higgs boson also couples to photons (gluons), through loops involving charged (coloured) massive particles, and one is led to consider the loop-induced decays $H\to Z\gamma$ [7], $H\to\gamma\gamma$ [8], $H\to gg$ [9], *etc.*

In order to match the high experimental precision to be achieved with a future e^+e^- LC, it is indispensable to take radiative corrections into account. A review of radiative corrections relevant for SM Higgs-boson phenomenology may be found in Refs. [10, 11]. At one loop, the electroweak corrections to $\Gamma\left(H\to f\bar{f}\right)$ [12, 13], $\Gamma(H\to VV)$ [12, 14, 15], and $\Gamma\left(H\to Zf\bar{f}\right)$ [10] and the QCD ones to $\Gamma(H\to q\bar{q})$ [16] are well established, including the dependence on all particle masses. Beyond one loop, only dominant classes of corrections were investigated, sometimes only in limiting cases. These include corrections enhanced by the strong-coupling constant α_s, the top Yukawa coupling g_{ttH}, and the Higgs self-coupling λ. Specifically, the two-loop QCD corrections were found for $\Gamma(H\to l^+l^-)$ [17], $\Gamma(H\to q\bar{q})$ $(q\neq t)$ [18, 19], $\Gamma\left(H\to t\bar{t}\right)$ [20], $\Gamma(H\to Z\gamma)$ [21], $\Gamma(H\to\gamma\gamma)$ [22], and $\Gamma(H\to gg)$ [19, 23]. Even three-loop QCD corrections were calculated, namely for $\Gamma(H\to q\bar{q})$ $(q\neq t)$ [24], $\Gamma(H\to\gamma\gamma)$ [25], and $\Gamma(H\to gg)$ [26]. In the last case, they are quite significant, the correction factor being [26]

$$K_{gg} = 1+\frac{215}{12}\,\frac{\alpha_s^{(5)}(M_H)}{\pi}+\left(\frac{\alpha_s^{(5)}(M_H)}{\pi}\right)^2\left(156.808-5.708\,\ln\frac{M_t^2}{M_H^2}\right), \tag{2}$$

which approximately amounts to $1+0.66+0.21$ for $M_H = 100$ GeV.

An efficient way of obtaining corrections leading in $X_t = g^2_{ttH}/(4\pi)^2$ to processes involving low-mass Higgs bosons is to construct an effective Lagrangian by integrating out the top quark. This may be conveniently achieved by means of a low-energy theorem [27], which relates the amplitudes of two processes which differ by the insertion of an external Higgs-boson line carrying zero four-momentum. A naïve version of it may be derived by observing the following two points: (i) the interactions of the Higgs boson with the massive particles in the SM emerge from their mass terms by substituting $M_i\to M_i(1+H/v)$; and (ii) a Higgs boson with zero four-momentum is represented by a constant field. This immediately implies that a zero-momentum Higgs boson may be attached to an amplitude, $\mathcal{M}(A\to B)$, by carrying out the operation

$$\lim_{p_H\to 0}\mathcal{M}(A\to B+H) = \frac{1}{v}\sum_i\frac{M_i\partial}{\partial M_i}\mathcal{M}(A\to B), \tag{3}$$

where i runs over all massive particles which are involved in the transition $A \to B$. This low-energy theorem comes with two caveats: (i) the differential operator in Eq. (3) does not act on the M_i appearing in coupling constants, since this would generate tree-level vertices involving the Higgs boson that do not exist in the SM; and (ii) Eq. (3) must be formulated for bare quantities if it is to be applied beyond the leading order.

In this way, the effective Lagrangian describing the l^+l^-H, W^+W^-H, and ZZH interactions is found to be

$$\mathcal{L}_{\rm eff} = \frac{H}{v}\left[-\sum_l m_l \bar{l}l(1+\delta_u) + 2M_W^2 W_\mu^- W^{+\mu}(1+\delta_{WWH}) + M_Z^2 Z_\mu Z^\mu (1+\delta_{ZZH})\right], \tag{4}$$

with [17, 28–30]

$$\begin{aligned}
\delta_u &= X_t\left\{\frac{7}{2} + 3\left[\frac{149}{8} - 6\zeta(2)\right]X_t - [3+2\zeta(2)]A - 56.703\,A^2\right\},\\
\delta_{WWH} &= X_t\left\{-\frac{5}{2} + \left[\frac{39}{8} - 18\zeta(2)\right]X_t + [9-2\zeta(2)]A + 27.041\,A^2\right\},\\
\delta_{ZZH} &= X_t\left\{-\frac{5}{2} - \left[\frac{177}{8} + 18\zeta(2)\right]X_t + [15-2\zeta(2)]A + 17.117\,A^2\right\},
\end{aligned} \tag{5}$$

where ζ is Riemann's zeta function, with value $\zeta(2) = \pi^2/6$, and $A = \alpha_s^{(6)}(M_t)/\pi$. Notice that δ_u is universal in the sense that it comprises just the renormalizations of the Higgs-boson wave function and vacuum expectation value. The analytic expressions of the $O(A^2)$ terms may be found in Ref. [29]. In $O\left(X_t^2\right)$, also the full M_b dependence is available [30]. From Eq. (4), one reads off that $\Gamma(H \to l^+l^-)$, $\Gamma(H \to W^+W^-)$, and $\Gamma(H \to ZZ)$ receive the correction factors $K_{ll} = (1+\delta_u)^2$, $K_{WW} = (1+\delta_{WWH})^2$, and $K_{ZZ} = (1+\delta_{ZZH})^2$, respectively. The $O(X_t)$, $O\left(X_t^2\right)$, and $O(X_tA)$ corrections to $\Gamma(H \to q\bar{q})$, where $q \neq b, t$, coincide with those for $\Gamma(H \to l^+l^-)$. The $O(X_tA^2)$ corrections to $\Gamma(H \to q\bar{q})$ were found in Ref. [31]. The effective-Lagrangian method in connection with the low-energy theorem was also employed to obtain the $O(X_tA)$ [32] and $O(X_tA^2)$ [31] corrections to $\Gamma(H \to b\bar{b})$, the $O(X_t)$ corrections to $\Gamma(H \to \gamma\gamma)$ [30], and the $O(X_t)$ [30, 31, 33] and $O(X_tA)$ [34] corrections to $\Gamma(H \to gg)$.

The expansion of $\Delta\rho = 1 - 1/\rho$, which measures the deviation of the electroweak ρ parameter from unity, analogous to Eq. (5) reads [35–37]

$$\Delta\rho = X_t\{3 + 3[19 - 12\zeta(2)]X_t - 2[1+2\zeta(2)]A - 43.782\,A^2\}. \tag{6}$$

The analytic expression of the $O(A^2)$ term and the full M_b dependence in $O\left(X_t^2\right)$ may be found in Refs. [36, 37], respectively. The $O\left(X_t^2\right)$ term exhibits a strong dependence on M_H, so that its value for $M_H = 0$ does not provide a useful approximation for realistic values of M_H [38]. Furthermore, subleading electroweak two-loop corrections, of $O\left(X_tG_FM_W^2\right)$, for $e\nu_\mu$ scattering and muon decay are not actually suppressed in magnitude against the $O\left(X_t^2\right)$ one [39]. Thus, the approximations by the $O\left(X_t^2\right)$ terms for $M_H = 0$ in Eq. (5) should be taken with a grain of salt. The coefficients of X_tA and X_tA^2 in Eqs. (5) and (6) are all negative and sizeable relative to the one of X_t. This is related to the use of the pole mass M_t. In fact, the convergence behaviour of these expansions may be considerably improved [29] by expressing them in terms of the scale-invariant $\overline{\rm MS}$ mass, $\mu_t = m_t(\mu_t)$, which is related to M_t by $\mu_t = M_t(1 - 4A/3 - 6.459\,A^2)$ [40].

At one loop in the conventional OS renormalization scheme, the production and decay rates of the Higgs boson exhibit singularities proportional to $(2M_V - M_H)^{-1/2}$ as M_H approaches $2M_V$ from below [15]. This problem is of phenomenological interest because the values $2M_W$ and $2M_Z$, corresponding to the W- and Z-boson pair production thresholds, lie within the M_H range favoured by the arguments presented in Sect. I. We recall that the OS mass M and total decay width Γ of an unstable boson are defined as

$$M^2 = M_0^2 + {\rm Re}\,A(M^2), \qquad M\Gamma = -Z\,{\rm Im}\,A(M^2), \qquad Z = \left[1 - {\rm Re}\,A'(M^2)\right]^{-1}, \tag{7}$$

where M_0 and $A(s)$ are the bare mass and unrenormalized self-energy, respectively, appearing in the propagator $i\left[s - M_0^2 - A(s)\right]^{-1}$. However, $A(s)$ possesses a branch point if s is at a threshold. If the threshold is due to a two-particle state with zero orbital angular momentum, then ${\rm Re}\,A'(s)$ diverges as $1/\beta$, where β is the relative velocity common to the two particles, as the threshold is approached from below [15, 41]. These threshold singularities are eliminated when the definitions of mass and total decay width are based on the complex-valued position $\bar{s}$ of the propagator's pole [41, 42], as [42]

$$\bar{s} = M_0^2 + A(\bar{s}) = m_2^2 - im_2\Gamma_2, \qquad m_2\Gamma_2 = -Z_2\,{\rm Im}\,A\left(m_2^2\right), \qquad Z_2 = \left[1 - \frac{{\rm Im}\,A\left(m_2^2\right) - {\rm Im}\,A(\bar{s})}{m_2\Gamma_2}\right]^{-1}. \tag{8}$$

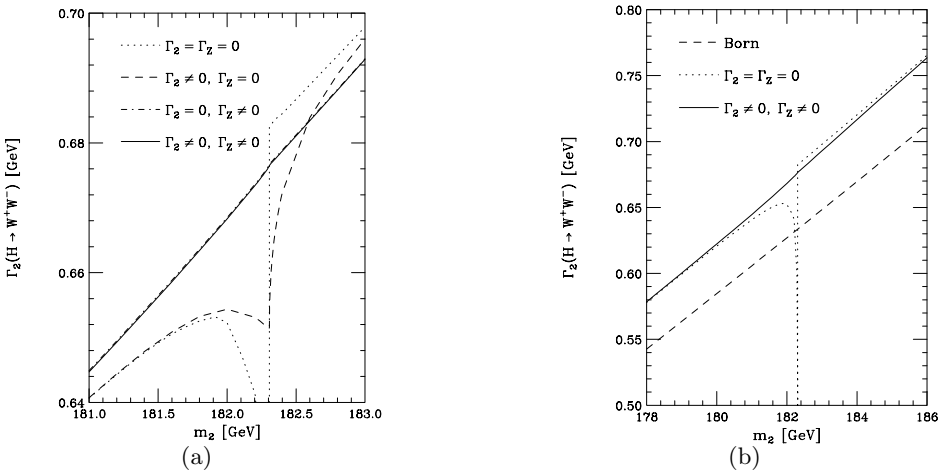


FIG. 2: $\Gamma(H \to W^+W^-)$ as a function of M_H [42]. (a) The threshold singularity at $M_H = 2M_Z$ in the OS scheme (dotted line) is regularized by adopting the pole scheme (dashed line), allowing for the Z-boson width to be finite (dot-dashed line), or both (solid line). (b) The one-loop results in the OS scheme (dotted line) and in the pole scheme with $\Gamma_Z \neq 0$ (solid line) are compared with the tree-level result (dashed line).

This is illustrated in Figs. 2 (a) and (b) for $\Gamma(H \to W^+W^-)$ in the vicinity of $M_H = 2M_Z$.

It is fair to say that radiative corrections for Higgs-boson decays have been explored to a similar degree as those for Z-boson decays. Unfortunately, this does not necessarily lead to similarly precise theoretical predictions. In fact, the errors on the latter are dominated by parametric uncertainties, mainly by those in $\alpha_s^{(5)}(M_Z)$ and the quark masses (see Fig. 3) [43, 44].

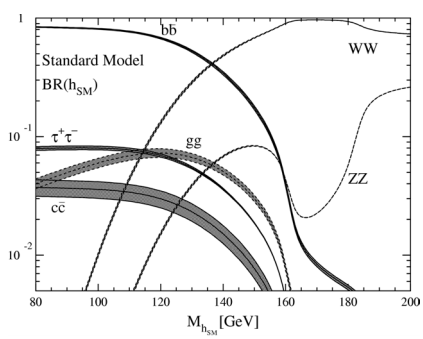


FIG. 3: Higgs-boson decay branching fractions, including theoretical uncertainties, as functions of M_H [44].

III. PRODUCTION IN e^+e^- COLLISIONS

The dominant mechanisms of Higgs-boson production in e^+e^- collisions are Higgs-strahlung and W^+W^- fusion, which, at the tree level, proceed through the Feynman diagrams depicted in Fig. 4(a). The cross section of ZZ fusion, $e^+e^- \to e^+e^-H$, is approximately one order of magnitude smaller than the one of W^+W^- fusion, because of weaker couplings. The total cross section of Higgs-strahlung reads

$$\sigma(e^+e^- \to ZH) = \frac{g_{ZZH}^2}{4\pi} \frac{G_F\left(v_e^2 + a_e^2\right)}{96\sqrt{2}} \frac{\sqrt{\lambda}\left(\lambda + 12sM_Z^2\right)}{s^2 D}, \tag{9}$$

where $v_f = 2I_f - 4s_w^2 Q_f$ and $a_f = 2I_f$ are the $Zf\bar{f}$ vector and axial-vector couplings, respectively, $\sqrt{s}$ is the centre-of-mass energy, $\lambda = \left[s - (M_Z + M_H)^2\right]\left[s - (M_Z - M_H)^2\right]$, and $D = \left(s - M_Z^2\right)^2 + M_Z^2\Gamma_Z^2$. Here, I_f is the third component of weak isospin of the left-handed component of f, Q_f is the electric charge of f, and

$c_w^2 = 1 - s_w^2 = M_W^2/M_Z^2$. The one of VV fusion may expressed as a one-dimensional integral [46]. They are both shown in Fig. 4(b) as functions of M_H for $\sqrt{s} = 350$, 500, and 800 GeV [45].

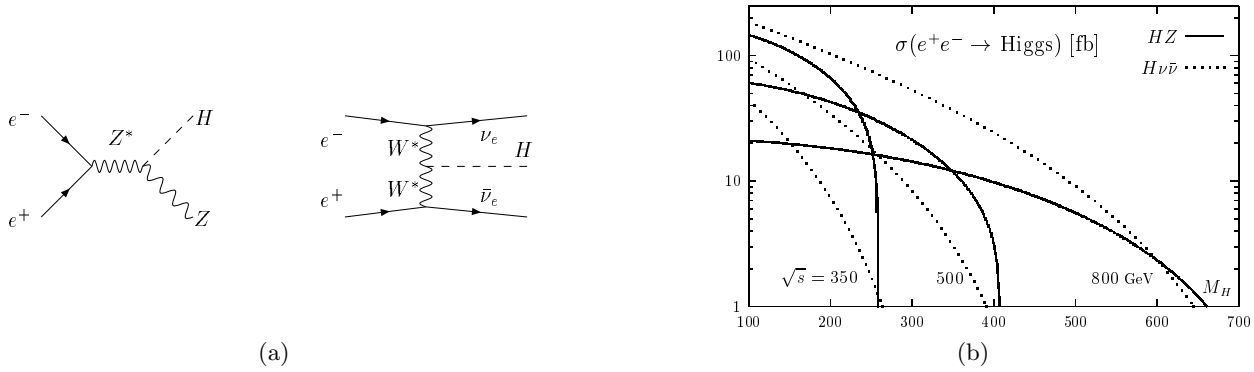


FIG. 4: (a) Feynman diagrams for Higgs-strahlung and W^+W^- fusion [45]. (b) Total cross sections of Higgs-strahlung and W^+W^- fusion as functions of M_H [45].

As for the Higgs-strahlung process, the electromagnetic [46, 47] and weak [46, 48] corrections are fully known at one loop. The latter is shown in Fig. 5 as a function of M_H for $\sqrt{s} = 192$ GeV, 500 GeV, 1 TeV, and 2 TeV.

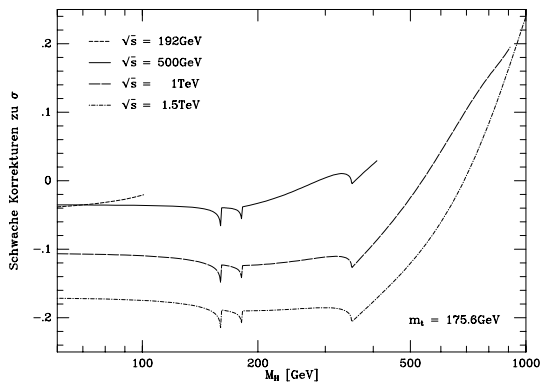


FIG. 5: Weak correction to $\sigma(e^+e^- \to ZH)$ as a function of M_H [46].

The electroweak corrections for VV fusion, a $2 \to 3$ process, are not yet available. However, the leading effects can be conveniently included as follows. The bulk of the initial-state bremsstrahlung can be taken into account in the so-called leading logarithmic approximation provided by the structure-function method, by convoluting the tree-level cross section with a radiator function, which is known through $O(\alpha^2)$ and can be further improved by soft-photon exponentiation [49]. The residual dominant corrections of fermionic origin can be incorporated in a systematic and convenient fashion by invoking the so-called improved Born approximation (IBA) [50]. These are contained in $\Delta\rho$ and $\Delta\alpha = 1 - \alpha/\overline{\alpha}$, which parameterizes the running of Sommerfeld's fine-structure constant from its value α defined in Thomson scattering to its value $\overline{\alpha}$ measured at the Z-boson scale. The recipe is as follows. Starting from the Born formula expressed in terms of c_w, s_w, and α, one substitutes $\alpha \to \overline{\alpha} = \alpha/(1 - \Delta\alpha)$ and $c_w^2 \to \overline{c}_w^2 = 1 - \overline{s}_w^2 = c_w^2(1 - \Delta\rho)$. One then eliminates $\overline{\alpha}$ in favour of G_F by exploiting the relation $(\sqrt{2}/\pi)G_F = \overline{\alpha}/\left(\overline{s}_w^2 M_W^2\right) = \overline{\alpha}/\left(\overline{c}_w^2 \overline{s}_w^2 M_Z^2\right)(1 - \Delta\rho)$, which correctly accounts for the leading fermionic corrections. Finally, one includes the corrections enhanced by X_t that are generated by Eq. (4). One thus obtains the correction factors

$$K_{ffH} = 1 + 2\delta_{ZZH} + 2\left[1 - 4c_w^2\left(\frac{Q_e v_e}{v_e^2 + a_e^2} + \frac{Q_f v_f}{v_f^2 + a_f^2}\right)\right]\Delta\rho \approx K_{eeH}, \qquad K_{\nu_e\nu_e H} = 1 + 2\delta_{WWH}, \tag{10}$$

for Higgs-strahlung, W^+W^- fusion, and ZZ fusion, respectively [29]. The interference of the scattering amplitudes for $\nu_e\bar{\nu}_e H$ and e^+e^-H production by Higgs-strahlung with those for W^+W^- and ZZ fusion, respectively, is negligible for $\sqrt{s} > M_Z + M_H$ [51]. It is important to keep in mind that the IBA is only reliable if $\sqrt{s}, M_H \ll 2M_t$.

It may be possible to operate a future e^+e^- LC in e^-e^-, $e^\pm\gamma$, or $\gamma\gamma$ modes. In e^-e^- collisions, Higgs bosons will be mainly produced via ZZ fusion, $e^-e^- \to e^-e^-H$ [52]. Its cross section emerges from the one of $e^+e^- \to e^+e^-H$ by crossing symmetry, as explained in Ref. [46], and it has a size very similar to the latter. The dominant Higgs-boson production mechanisms in $e^\pm\gamma$ collisions include the processes $e^\pm\gamma \to \nu_e W^\pm H$ [53, 54], $e^\pm\gamma \to e^\pm ZH$ [53], and $e^\pm\gamma \to e^\pm\gamma\gamma \to e^\pm H$ [55], which proceeds via charged-fermion and W-boson loops. In $\gamma\gamma$ collisions, Higgs bosons will be chiefly created through $\gamma\gamma$ fusion, $\gamma\gamma \to H$ [56], which is mediated by the same types of loops. Cutting open the W-boson loops leads to the process $\gamma\gamma \to W^+W^-H$ [57], which benefits from the huge cross section of the parent process $\gamma\gamma \to W^+W^-$. The process $\gamma\gamma \to t\bar{t}H$ [58] is sensitive to the top Yukawa coupling g_{ttH}, but it suffers from phase-space suppression.

IV. QUANTUM NUMBERS AND COUPLINGS FROM FINAL STATES

The spin, parity, and charge-conjugation quantum numbers J^{PC} of Higgs bosons can be determined at a future e^+e^- LC in a model-independent way. The observation of the decay or fusion processes $H \rightleftharpoons \gamma\gamma$ would rule out $J = 1$ by the Landau-Yang theorem and, at the same time, fix C to be positive [59].

The angular distribution of $e^+e^- \to ZH$ depends on J and P. The SM Higgs boson is a 0^{++} state, and its couplings to two Z bosons is proportional to $\vec{\epsilon}\cdot\vec{\epsilon'}$ in the laboratory frame, where $\vec{\epsilon}$ and $\vec{\epsilon'}$ are the polarization three-vectors of the Z bosons. In order to distinguish the SM Higgs boson from a CP-odd 0^{-+} state A, or a CP-violating mixture of the two, which will be generically denoted by Φ, one may consider a $ZZ\Phi$ coupling of the form [60]

$$\Gamma_{ZZ\Phi} = i g_{ZZH}\left(g^{\mu\mu'} + i\frac{\eta}{M_Z^2}\epsilon^{\mu\mu'\nu\nu'}p_\nu p'_{\nu'}\right), \tag{11}$$

where p, p' and μ, μ' are the incoming four-momenta and Lorentz indices of the two Z bosons, respectively, and η is a dimensionless factor. In the case $\eta = 0$, we recover the SM Higgs boson, while the absence of the first term in Eq. (11) corresponds to an A boson. In the laboratory frame, the ZZA coupling is proportional to $(\vec{\epsilon}\times\vec{\epsilon'})\cdot(\vec{p}-\vec{p'})$. The CP-odd case is realized in the minimal supersymmetric extension of the SM and in two-Higgs-doublet models (2HDM) without CP violation, in which the ZZA couplings are induced at the level of quantum loops. However, in a more general scenario, η need not be loop suppressed, and it is useful to allow for η to be arbitrary in the experimental data analysis. In a general 2HDM, the three neutral Higgs bosons correspond to arbitrary mixtures of CP eigenstates, and their production and decay processes exhibit CP violation. The differential cross section of $e^+e^- \to Z\Phi$ that results from the coupling of Eq. (11) reads

$$\begin{aligned}\frac{d\sigma(e^+e^-\to Z\Phi)}{d\cos\theta} &= \frac{g_{ZZH}^2}{4\pi}\,\frac{G_F\left(v_e^2+a_e^2\right)}{16\sqrt{2}}\,\frac{M_Z^2\sqrt{\lambda}}{sD}\left[1+\frac{\lambda}{8sM_Z^2}\sin^2\theta+\eta\frac{2v_ea_e}{v_e^2+a_e^2}\,\frac{\sqrt{\lambda}}{M_Z^2}\cos\theta\right.\\ &\quad\left.+\,\eta^2\frac{\lambda}{8M_Z^4}(1+\cos^2\theta)\right],\end{aligned} \tag{12}$$

where θ is the polar angle of the Z boson w.r.t. to the beam axis in the laboratory frame. Thus, the angular distribution of $e^+e^- \to ZA$, namely $(1/\sigma)d\sigma(e^+e^- \to ZA)/d\cos\theta = (3/8)(1+\cos^2\theta)$, is very distinct from the SM one, which is $(1/\sigma)d\sigma(e^+e^- \to ZH)/d\cos\theta \approx (3/4)\sin^2\theta$ for $\sqrt{s} \gg M_Z$ [59]. The presence of the interference term (linear in η) in Eq. (12), would generate a forward-backward asymmetry, which would be a clear signal for CP violation. Another discriminator between the CP-even and CP-odd cases is provided by the threshold behaviour of the cross section, which is proportional to $\beta = \sqrt{\lambda}/s$ and β^3, respectively [59]. In the most general situation, where the particle produced in association with the Z boson corresponds to a J^P state, the threshold behaviour is $\sigma(e^+e^- \to Z\Phi) \propto \beta^n$, where $n = 1$ for $J^P = 0^+, 1^+, 2^+$, $n = 3$ for $J^P = 0^-, 1^-, 2^-$, $n = 2J-3$ for $J^P = 3^-, 4^+, 5^-, \ldots$, and $n = 2J-1$ for $J^P = 3^+, 4^-, 5^+, \ldots$. We conclude that the observation of a threshold behaviour linear in β would rule out the assignments $J^P = 0^-, 1^-, 2^-, 3^\pm, 4^\pm, \ldots$.

The angular distribution of $e^+e^- \to Z\Phi$ can also be exploited to establish the $J = 0$ nature of the Higgs bosons. To this end, it should be compared with the one of $e^+e^- \to ZZ$, which exhibits a distinctly different angular momentum structure. Owing to the electron exchange in the t-channel, the $e^+e^- \to ZZ$ amplitude is built up by many partial waves, which peak in the forward and backward directions. In Fig. 6, the angular distributions of ZH, ZA, and ZZ production are shown for $\sqrt{s} = 500$ GeV, assuming a Higgs-boson mass of 120 GeV.

The determination of the J^{PC} quantum numbers of the Higgs bosons can be refined by taking the angular distributions of their decay products into account. The $J = 0$ property manifests itself in the complete absence

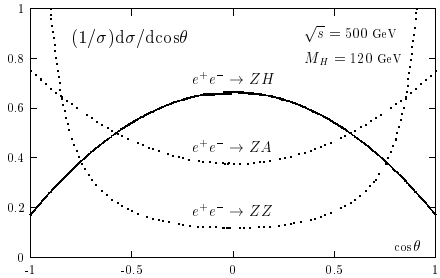


FIG. 6: $\sigma(e^+e^- \to ZH)$, $\sigma(e^+e^- \to ZA)$, and $\sigma(e^+e^- \to ZZ)$ as functions of $\cos\theta$ [59].

of angular correlations between the initial- and final-state particles. The criteria to distinguish between CP-even and CP-odd Higgs bosons or mixtures thereof include the polarization of the vector bosons $V = W, Z$ in the decay $\Phi \to VV$, the distribution in the mass M_* of the virtual boson V^* in the decay $\Phi \to VV^*$, and characteristic features of the angular distribution of the decay $\Phi \to V^*V^* \to \left(f\bar{f}\right)\left(f'\bar{f}'\right)$ [59, 60].

In the effective-Lagrangian approach, the $ZZ\Phi$ coupling in Eq. (11) is not the most general one [62–64]. In fact, the first term may come with a fudge factor different from unity, and there may be two more independent CP-even terms. Similarly, there may be an effective $Z\gamma\Phi$ coupling, involving two CP-even and one CP-odd terms. The most general effective $ZV\Phi$ interaction Lagrangian reads [62, 63]

$$\mathcal{L}_{\text{eff}} = \frac{g_{ZZH}}{2}(1+a_Z)HZ_\mu Z^\mu + \frac{g_{ZZH}}{M_Z^2}\sum_{V=Z,\gamma}\left[b_V HZ_{\mu\nu}V^{\mu\nu} + c_V\left(\partial_\mu HZ_\nu - \partial_\nu HZ_\mu\right)V^{\mu\nu} + \tilde{b}_V HZ_{\mu\nu}\tilde{V}^{\mu\nu}\right], \quad (13)$$

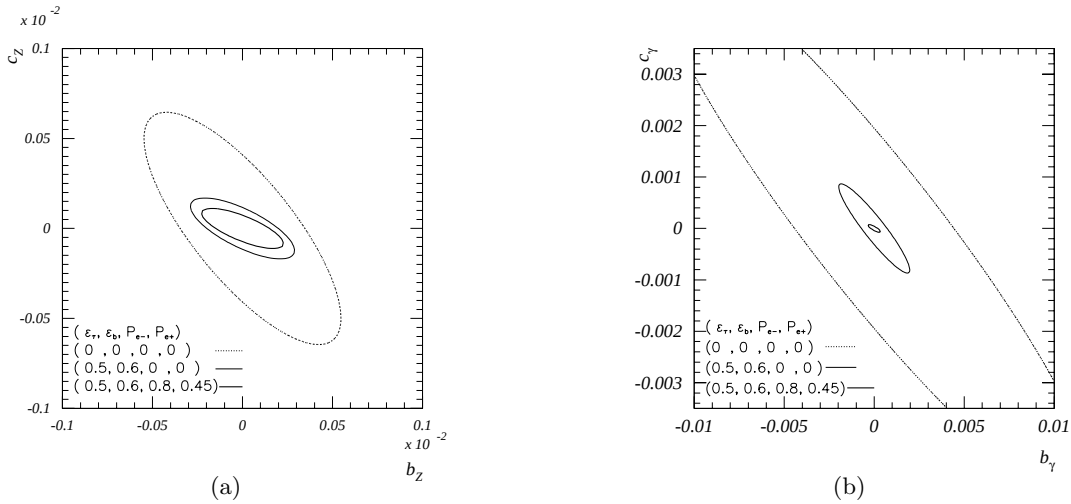
where $V_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu$ and $\tilde{V}_{\mu\nu} = \epsilon_{\mu\nu\alpha\beta}V^{\alpha\beta}$. Here, we have neglected the scalar components of the vector bosons, by putting $\partial_\mu Z^\mu = \partial_\mu V^\mu = 0$. The couplings a_Z, b_Z, c_Z, b_γ, and c_γ are CP-even, while $\tilde{b}_Z$ and $\tilde{b}_\gamma$ are CP-odd.

With sufficiently high luminosity, it should be possible to determine, by means of the optimal-observable method [65, 66], most of these couplings from the angular distribution of $e^+e^- \to Z\Phi \to \left(f\bar{f}\right)\Phi$. The achievable bounds can be significantly improved by measuring the tau-lepton helicities, identifying the bottom-hadron charges, polarizing the electron and positron beams, and running at two different values of $\sqrt{s}$ [63]. The results for energy $\sqrt{s} = 500$ GeV, luminosity $L = 300$ fb^{-1}, efficiencies $\epsilon_\tau = 50\%$ and $\epsilon_b = 60\%$, and polarizations $P_{e^-} = \pm 80\%$ and $P_{e^+} = \mp 45\%$ are visualized in Figs. 7(a) and (b). Here, the couplings are assumed to be real, and a_Z is fixed. In order to also determine a_Z, one needs to perform the experiment at two different values of $\sqrt{s}$. We observe that the $ZZ\Phi$ couplings are generally well constrained, even for $\epsilon_\tau = \epsilon_b = P_{e^-} = P_{e^+} = 0$, while the $Z\gamma\Phi$ couplings are not. The constraints on the latter may be significantly improved by the above-named options, especially by beam polarization.

Once g_{ZZH} has been pinned down, the top Yukawa coupling g_{ttH} can be extracted by studying the process $e^+e^- \to t\bar{t}H$ [67]. The QCD correction to its cross section can be of either sign, depending on $\sqrt{s}$, and reach a magnitude of several ten percent [68, 69]. This may be seen from Fig. 8(a), where the Born and QCD-corrected cross sections are shown as functions of M_H for $\sqrt{s} = 500$ GeV, 1 TeV, and 2 TeV. Anomalous top Yukawa couplings may be extracted from the angular distribution of $e^+e^- \to t\bar{t}\Phi$ with the help of the optimal-observable method [66]. The analysis of double Higgs-strahlung, $e^+e^- \to ZHH$, and W^+W^- double-Higgs fusion, $e^+e^- \to \bar{\nu}_e\nu_e HH$, offers the possibility to extract the trilinear Higgs self-coupling g_{HHH} [70]. The cross sections of these two processes are relatively modest, but they can be enhanced by factors 2 and 4, respectively, by using beam polarization. They are shown as functions of M_H for $\sqrt{s} = 500$ GeV, 1 TeV, and 1.6 TeV in Fig. 8(b). The sensitivity to g_{HHH} is strongest close to the production thresholds.

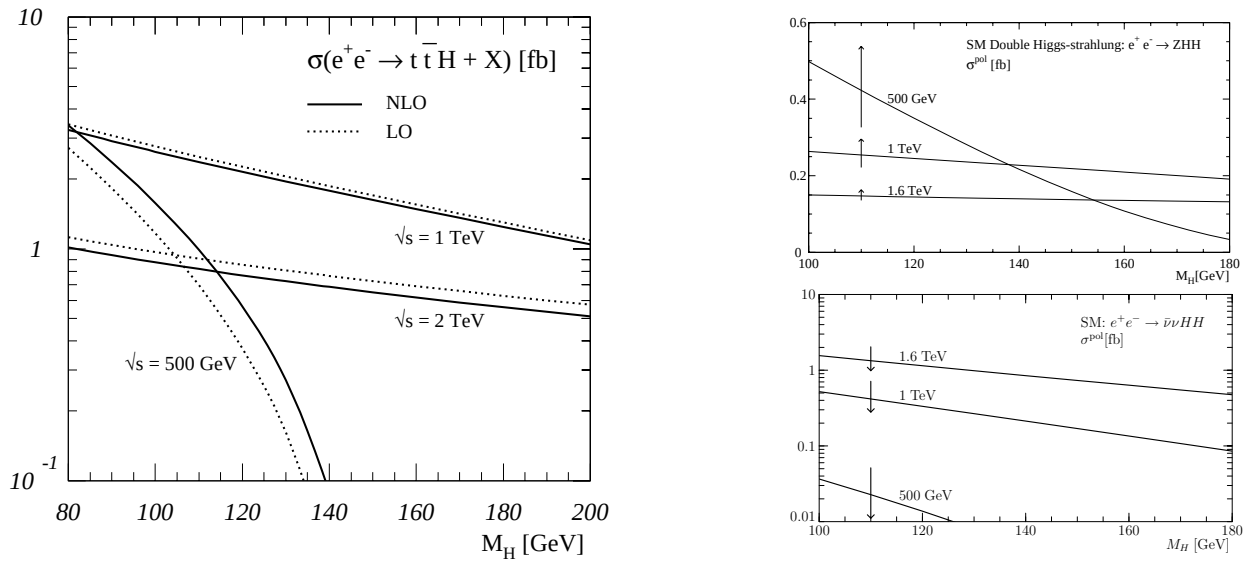
V. CONCLUSIONS AND OUTLOOK

We reviewed theoretical results that are relevant for the phenomenology of the SM Higgs boson at a future e^+e^- LC, putting special emphasis on radiative corrections to its partial decay widths and production cross sections, and on the logistics of extracting its quantum numbers and couplings from the analysis of appropriate final states. It is fair to say that theoretical predictions for partial decay widths and production cross sections are generally in good shape. However, the precision on the partial decay widths is limited by parametric



(a) (b)

FIG. 7: Contours of $\chi^2 = 1$ in the (a) (b_Z, c_Z) and (b) (b_γ, c_γ) planes for 300 fb^{-1} of data at $\sqrt{s} = 500$ GeV [63]. In each case, the other degrees of freedom are integrated out.



(a) (b)

FIG. 8: (a) $\sigma(e^+e^- \to t\bar{t}H)$ with and without QCD corrections as a function of M_H [69]. (b) $\sigma(e^+e^- \to ZHH)$ and $\sigma(e^+e^- \to \nu_e\bar{\nu}_e HH)$ as functions of M_H [70]. The vertical arrows indicate the shifts in cross section induced by the variation of g_{HHH} by $\pm 50\%$.

uncertainties, mainly by those in $\alpha_s^{(5)}(M_Z)$ and the quark masses. The strategies for the determination of the Higgs profile are also well elaborated.

The list of urgent tasks left to be done includes the calculation of the full $O(\alpha)$ corrections for important $2 \to 3$ processes, such as W^+W^- fusion, ZZ fusion, and $t\bar{t}H$ associated production, and the inclusion of background processes and detector simulation.

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