Higher-order corrections to Higgs-boson decays*

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We review the status of radiative corrections to the decay rates of the Standard-Model Higgs boson in the mass range accessible at LEP 2. Specifically, we consider corrections of $\mathcal{O}(\alpha)$, $\mathcal{O}(\alpha_G m_t^2)$, and $\mathcal{O}(G_F^2 M_H^4)$ to the fermionic decay rates and two-loop QCD corrections to the hadronic decay rates.

1. Introduction

One of the great puzzles of elementary particle physics today is whether nature makes use of the Higgs mechanism of spontaneous symmetry breaking to generate the observed particle masses. The Higgs boson, $H$, is the missing link sought to verify this concept in the Standard Model. Many of the properties of the Higgs boson are fixed, e.g., its couplings to the gauge bosons, $g_{VH} = \frac{25/4}{G_F^{3/2}} M_V^2$ ($V = W, Z$), and fermions, $g_{fH} = \frac{2^{1/4} G_F^{1/2}}{m_f}$, and the vacuum expectation value, $v = 2^{-1/4} G_F^{-1/2} \approx 246$ GeV. However, its mass, $M_H$, and its self-couplings, which depend on $M_H$, are essentially unspecified.

The failure of experiments at LEP 1 and SLC to detect the decay $Z \rightarrow f\bar{f}H$ has ruled out the mass range $M_H \leq 63.8$ GeV at the 95% confidence level [1]. At the other extreme, unitarity arguments in intermediate-boson scattering at high energies [2] and considerations concerning the range of validity of perturbation theory [3] establish an upper bound on $M_H$ at $(8\pi \sqrt{3}/3G_F)^{1/2} \approx 1$ TeV in a weakly interacting Standard Model.

The Higgs-boson discovery potential of LEP 1 and SLC is almost exhausted [4]. Prior to the advent of the LHC, the Higgs-boson search will be restricted to the lower mass range. With LEP 2 it should be possible to find a Higgs boson with $M_H \leq 100$ GeV when high energy and luminosity can be achieved [5]. A possible 4-TeV upgrade of the Tevatron might cover the $M_H$ range up to 120 GeV or so [6]. At an $e^+e^-$ linear collider operating at 300 GeV, 50 fb$^{-1}$ luminosity and a $b$-tagging efficiency of 50% would be sufficient to detect a Higgs boson with $M_H \leq 150$ GeV in the $\mu^+\mu^- b\bar{b}$ channel [7].

Below the onset of the $W^+W^-$ threshold, the Standard-Model Higgs boson is relatively long-lived, with $\Gamma_H < 100$ MeV, so that, to a good approximation, its production and decay processes may be treated independently. The low-mass Higgs boson, with $M_H \leq M_Z$, decays with more than 99% probability into a fermion pair [8]. With $M_H$ increasing, the $W^+W^-$ mode, with at least one $W$ boson being off shell, gradually gains importance. Its branching fraction surpasses that of the $\tau^+\tau^-$ mode at $M_H \approx 115$ GeV and that of the $b\bar{b}$ mode at $M_H \approx 135$ GeV [9]. In the near future, however, Higgs-boson searches will rely mostly on the $f\bar{f}$ modes.

Quantum corrections to Higgs-boson phenomenology have received much attention in the literature; for a review, see Ref. [10]. The experimental relevance of radiative corrections to the $f\bar{f}$ branching fractions of the Higgs boson has been emphasized recently in the context of a study dedicated to LEP 2 [11]. Techniques for the measurement of these branching fractions at a $\sqrt{s} = 500$ GeV $e^+e^-$ linear collider have been elaborated in Ref. [12].

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2. Electroweak corrections to $\Gamma (H \rightarrow f \bar{f})$

In the Born approximation, the $ff\bar{f}$ partial widths of the Higgs boson are given by [11]

$$\Gamma_0 (H \rightarrow f \bar{f}) = \frac{N_c G_F M_H m_f^2}{4\pi \sqrt{2}} \left( 1 - \frac{4m_f^2}{M_H^2} \right)^{3/2}, \quad (1)$$

where $N_c = 1$ (3) for lepton (quark) flavours.

2.1. One-loop approximation

The full one-loop electroweak corrections to Eq. (1) are now well established [12,13]. They consist of an electromagnetic and a weak part, which are separately finite and gauge independent. They may be written as an overall factor,

$$\left( 1 + \frac{\alpha}{\pi} Q_f^2 \Delta_{\text{em}} + \Delta_{\text{weak}} \right).$$

For $M_H \rightarrow 2m_f$, $\Delta_{\text{em}}$ exhibits a threshold singularity with a positive sign,

$$\Delta_{\text{em}} = 3\zeta(2) \left( \frac{1}{\beta} + \beta \right) - 1 + \mathcal{O} (\beta^2 \ln \beta), \quad (2)$$

where $\beta = \sqrt{1 - 4m_f^2/M_H^2}$ is the fermion velocity in the centre-of-mass frame. This divergence is tamed by the overall threshold factor $\beta^3$ of Eq. (1). For $M_H \gg 2m_f$, $\Delta_{\text{em}}$ develops a large logarithm,

$$\Delta_{\text{em}} = -\frac{3}{2} \ln \frac{M_H^2}{m_f^2} + \frac{9}{4} + \mathcal{O} \left( \frac{m_f^2}{M_H^2} \ln \frac{M_H^2}{m_f^2} \right). \quad (3)$$

For $M_H \ll 2M_W$, the weak part is well approximated by [13]

$$\Delta_{\text{weak}} = \frac{G_F}{8\pi \sqrt{2}} \left\{ K_f m_t^2 + M^2_W \left( 3s_w^2 \ln c_w^2 - 5 \right) 
+ M_Z^2 \left[ \frac{1}{2} - 3 \left( 1 - 4s_w^2|Q_f|^2 \right) \right] \right\}, \quad (4)$$

where $c_w^2 = 1 - s_w^2 = M_W^2/M_Z^2$, $K_b = 1$, and $K_f = 7$ for all other flavours, except for top. The $t\bar{t}$ mode will not be probed experimentally anytime soon and we shall not be concerned with it in the remainder of this presentation. Equation (4) has been obtained by putting $M_H = m_f = 0$ ($f \neq t$) in the expression for the full one-loop weak correction. It provides a very good approximation for $f = \tau$ up to $M_H \approx 135$ GeV and for $f = b$ up to $M_H \approx 70$ GeV, the relative deviation from the full weak correction being less than 15% in each case. From Eq. (4) it is evident that the dominant effect is due to virtual top quarks. In the case $f \neq b$, the $m_t$ dependence is carried solely by the renormalizations of the wave function and the vacuum expectation value of the Higgs field and is thus flavour independent. These corrections are of the same nature as those considered in Ref. [14]. For $f = b$, there are additional $m_t$ dependent contributions from the $bbH$ vertex correction and the $b$-quark wave-function renormalization. Incidentally, they cancel almost completely the universal $m_t$ dependence. It is amusing to observe that a similar situation has been encountered in the context of the $Z \rightarrow f \bar{f}$ decays [15,16]. In summary, the universal virtual-top-quark term will constitute the most important part of the weak one-loop corrections to Higgs-boson decays in the near future. In Sect. 2.2, we shall present the two-loop gluon correction to this term.

It is interesting to consider the high-$M_H$ limit of $\Delta_{\text{weak}}$. The leading term is of $\mathcal{O} (G_F M_H^2)$ and flavour independent, and arises from loops in which only physical and unphysical Higgs bosons circulate. It reads [17]

$$\Delta_{\text{weak}} = \frac{G_F M_H^2}{8\pi \sqrt{2}} \left( \frac{13}{2} - \pi \sqrt{3} \right) \approx 11.1\% \left( \frac{M_H}{1\text{ TeV}} \right)^2. \quad (5)$$

It may be extracted conveniently in the framework of the Higgs-Goldstone scalar theory. As we shall see in Sect. 2.3, this formalism carries over to $\mathcal{O} (G_F^2 M_H^4)$.

2.2. Two-loop $\mathcal{O} (\alpha_s G_F m_f^2)$ corrections

The universal $\mathcal{O} (G_F m_f^2)$ term of $\Delta_{\text{weak}}$ resides inside the combination

$$\delta = -\frac{\Pi_{WW}(0)}{M_W^2} - \Re \Pi_{HH} (M_H^2), \quad (6)$$

where $\Pi_{WW}$ and $\Pi_{HH}$ are the unrenormalized self-energies of the $W$ and Higgs bosons, respectively. The same is true of its QCD correction.
For $M_H < 2m_t$ and $m_b = 0$, the one-loop term reads \[13\]
\[
\delta_0 = N_c G_F m^2 \left( \frac{1}{2} \right) \left( \frac{1}{2} - r^2 \right) \arcsin \sqrt{r}
\]
\[
- \frac{1}{4} \frac{1}{2} r^2 \right], \tag{7}
\]
where $r = (m_H^2/4m^2)$. In the same approximation, the two-loop term may be written as \[13\]
\[
\delta_1 = \frac{N_c C_F}{2} \frac{G_F m^2}{\alpha_s} \left( \frac{2}{3} \frac{19}{4} r^2 \right) - \text{Re} H_1^1(r),
\]
where $H_1^1$ has an expression in terms of dilogarithms and trilogarithms. In the heavy-quark limit ($r \ll 1$), one has \[13\]
\[
H_1^1(r) = 6\zeta(3) + 3\zeta(2) - \frac{13}{4} + \frac{122}{135} r + \mathcal{O}(r^2). \tag{9}
\]
Combining Eqs. (7) and retaining only the leading high-$m_t$ terms, one finds the QCD-corrected coefficients $K_f$ for $f \neq b$, \[13\]
\[
K_f = 7 - 2 \left( \frac{1}{3} + \frac{3}{\pi} \right) \alpha_s \approx 7 - 4.004 \alpha_s. \tag{10}
\]
We recover the notion that, in Electroweak Physics, the one-loop $\mathcal{O}(G_F m^2)$ terms get screened by their QCD corrections. The QCD correction to the shift in $\Gamma (H \rightarrow f \bar{f})$ induced by a pair of quarks with arbitrary masses may be found in Ref. [13].

### 2.3. Two-loop $\mathcal{O}(G_F^2 M_H^2)$ corrections

We shall now outline the derivation of the $\mathcal{O}(G_F M_H^2)$ and $\mathcal{O}(G_F^2 M_H^4)$ corrections to $\Gamma (H \rightarrow f \bar{f})$ \[14\]. In the limit of $M_H \gg M_W$, power counting and inspection of coupling constants (in the 't Hooft-Feynman gauge) reveal that we may concentrate on those one- and two-loop diagrams which involve only the physical Higgs boson and the longitudinal polarization states of the intermediate bosons. The fermion mass and wave-function renormalizations do not receive contributions in the orders considered. Even though the process $H \rightarrow f \bar{f}$ does not involve external gauge bosons, the fact that the corrections of interest arise solely from diagrams containing only virtual Higgs and gauge bosons allows us to simplify the calculation enormously by using the Goldstone-boson equivalence theorem \[21\]. In particular, the dominant contributions for $M_H \gg M_W$ can be calculated by replacing the propagators of the gauge bosons by the propagators of the corresponding scalar Goldstone bosons, and setting the gauge coupling and the gauge-boson masses to zero \[21\].

Once the equivalence theorem has been invoked, all the information necessary for the calculation may be obtained from the Lagrangian of the Higgs or symmetry-breaking sector of the Standard Model. This characterizes the kinematics and interactions of the Higgs boson, $H$, and the Goldstone bosons, $w^\pm$ and $z$. The latter remain massless and satisfy a residual SO(3) symmetry. The boson masses and wave functions are renormalized according to the usual on-mass-shell procedure. Furthermore, we fix the physical vacuum expectation value and quartic coupling by $v = 2^{-1/4}G_F^{1/2}$ and $\lambda = G_F M_H^2 / \sqrt{2}$, respectively. Straightforward algebra shows that the Higgs-fermion Yukawa coupling receives a multiplicative correction of the form $(Z_H/Z_w)^{1/2}$, where $Z_H$ and $Z_w$ are the wave-function renormalizations of $H$ and $(w^\pm, z)$. Thus, the leading high-$M_H$ corrections to the fermionic Higgs-boson decay rates can be included by multiplying Eq. (1) with the overall factor $Z_H/Z_w$, independently of the fermions involved.

The wave-function renormalizations $Z_H$ and $Z_w$ are defined as \[14\]
\[
Z_w^{-1} = 1 - \frac{d}{dp^2} \Pi_w^0 (p^2) \bigg|_{p^2 = 0},
\]
\[
Z_H^{-1} = 1 - \frac{d}{dp^2} \text{Re} \Pi_H^0 (p^2) \bigg|_{p^2 = M_H^2}, \tag{11}
\]
where $\Pi_w^0$ and $\Pi_H^0$ are the self-energy functions for the bare fields. Using dimensional regularization, we find that $Z^{-1}$ can be written in factored form, \[14\]
\[
Z^{-1}_\sigma = \left( 1 + a_\sigma \lambda \xi + b_\sigma \lambda^2 \xi^2 \right) \left( 1 + \frac{3}{\epsilon} \lambda^2 \xi^2 \right) + \mathcal{O} \left( \lambda^4 \xi^3 \right), \quad \sigma = w, H, \tag{12}
\]
so rapidly with

The importance of this term increases changes drastically when the two-loop term is in-

and discarding terms beyond $\mathcal{O}\left(\hat{\lambda}^2\right)$, we obtain the alternative representation

$$
\frac{Z_H}{Z_w} = 1 + \left(a_w - a_H\right)\hat{\lambda} + \left(b_w - b_H - a_wa_H + a_h^2\right)\hat{\lambda}^2 + \ldots
$$

where $a_w = 1$, $a_H = 2\pi \sqrt{3} - 12 \approx -1.117$, $b_w \approx -24.769$, and $b_H \approx 265.764$ \cite{13}. Equation (13), which naturally emerges from our formalism, automatically resums one-particle-reducible Higgs-boson self-energy diagrams in a way that conforms with the standard procedure in $Z$-boson physics; see, e.g., Ref. \cite{16}. However, we have no control of terms beyond $\mathcal{O}\left(\hat{\lambda}^2\right)$, and are not aware of any physical organizing principle analogous to that provided at high energies by the renormalization group, which would allow us to select an optimum resummation scheme. Expanding Eq. (13) and discarding terms beyond $\mathcal{O}\left(\hat{\lambda}^2\right)$, we obtain the alternative representation

$$
\frac{Z_H}{Z_w} = 1 + \left(a_w - a_H\right)\hat{\lambda} + \left(b_w - b_H - a_wa_H + a_h^2\right)\hat{\lambda}^2 + \ldots
$$

$$
\approx 1 + 11.1\% \left(\frac{M_H}{1\text{TeV}}\right)^2 - 78.6\% \left(\frac{M_H}{1\text{TeV}}\right)^4,
$$

which extends Eq. (3) to two loops.

We are now in a position to explore the implications of our results. In Fig. 1, we show the leading electroweak corrections to $\Gamma(H \to f\bar{f})$ in the one- and two-loop approximations with and without resummation of one-particle-reducible higher-order terms plotted as functions of $M_H$. We shall concentrate first on the expanded expression in Eq. (13). While the $\mathcal{O}(G_F M_H^2)$ term (upper dotted line) gives a modest increase in the rates, by 11.1% at $M_H = 1$ TeV, the situation changes drastically when the two-loop term is included. The importance of this term increases so rapidly with $M_H$ that it already cancels the one-loop term completely for $M_H = 375$ GeV. By $M_H = 530$ GeV, it has twice the magnitude of the one-loop term, and the sum of one- and two-loop corrections (lower dotted line) reaches the same magnitude as the one-loop corrections, but with a reversed sign. The perturbation series for $\Gamma(H \to f\bar{f})$ in powers of $\lambda$ or $G_F M_H^2$ clearly ceases to converge usefully, if at all, for values of $M_H$ beyond about 400 GeV. A Higgs boson with a mass larger than 400 GeV effectively becomes a strongly interacting particle in the electroweak processes which contribute to the correction, a very surprising result. Conversely, $M_H$ must not exceed approximately 400 GeV if the standard electroweak perturbation theory is to be predictive for the decays $H \to f\bar{f}$. Note that, for $M_H \gtrsim 400$ GeV, one cannot use the usual unitarization schemes invoked in studies of $W_L^\pm$, $Z_L$, $H$ scattering to restore the predictiveness for the Higgs-boson width, as no unitarity violation occurs.

One might expect to improve the perturbative result in the upper $M_H$ range somewhat by resumming the one-particle-reducible contributions to the Higgs-boson wave-function renormalization according to Eq. (13). This leads to an insignificant increase of the one-loop correction (upper solid line), while the negative effect of the two-loop correction is appreciably lessened (lower solid line), i.e., the ratio of two- to one-loop corrections is rendered more favourable theoretically. However, in the mass range below $M_H = 600$ GeV, this effect is too feeble to change our conclusions concerning the breakdown of perturbation theory. Moreover, the resummation of reducible one-loop terms in the perturbation series does not yield a proper estimate for the size of the two-loop terms, so that there is no reason to favour this approach in the present problem.

3. QCD corrections to $\Gamma(H \to 2j)$

Both $H \to q\bar{q}$ and $H \to gg$ lead to two-jet final states. These processes are relevant phenomenologically for the search for the low- and intermediate-mass Higgs boson at $e^+e^-$ colliders, while they are very difficult to discern from background reactions at hadron colliders. Both decay
channels receive significant QCD corrections.

3.1. Next-to-next-to-leading order corrections to $\Gamma (H \to q\bar{q})$

The one-loop QCD correction to $\Gamma (H \to q\bar{q})$ emerges from one-loop QED correction, discussed in Sect. 2.1, by substituting $\alpha_s C_F$, where $C_F = (N_c^2 - 1)/(2N_c) = 4/3$, for $\alpha Q^2$. This corresponds to the on-shell scheme, where the quark pole mass, $M_q$, is used as a basic parameter. From Eq. (3) it is apparent that, for $M_q \ll M_H/2$, large logarithmic corrections occur. In general, they are of the form $(\alpha_s/\pi)^n \ln^n(M_H^2/M_q^2)$, with $n \geq m$. Appealing to the renormalization-group equation, these logarithms may be absorbed completely into the running quark mass, $m_q(\mu)$, evaluated at scale $\mu = M_H$. In this way, these logarithms are resummed to all orders and the perturbation expansion converges more rapidly. This observation gives support to the notion that the $q\bar{q}H$ Yukawa couplings are controlled by the running quark masses.

The values of $M_q$ may be estimated from QCD sum rules. To obtain $m_q(M_H)$, one proceeds in two steps. Firstly, one evaluates $m_q(M_q)$ from

$$\frac{M_q}{m_q(M_q)} = 1 + C_F \frac{\alpha_s(M_q)}{\pi} + K_q \left(\frac{\alpha_s(M_q)}{\pi}\right)^2,$$

with

$$K_q \approx 16.11 - 1.04 \sum_{i=1}^{N_F-1} \left(1 - \frac{M_i}{M_q}\right),$$

where the sum extents over all quark flavours with $M_i < M_q$. Specifically, $K_c = 13.3$ and $K_b = 12.4$. Secondly, one determines $m_q(M_H)$ via the scaling law

$$m_q(M_H) = m_q(M_q) \frac{c_q(\alpha_s(M_H)/\pi)}{c_q(\alpha_s(M_q)/\pi)},$$

where

$$c_\circ(x) = \left(\frac{25}{6} x^2\right)^{12/25} (1 + 1.014x + 1.389x^2),$$

$$c_0(x) = \left(\frac{23}{6} x^2\right)^{12/23} (1 + 1.175x + 1.501x^2).$$

For $q \neq t$, the QCD corrections to $\Gamma (H \to q\bar{q})$ are known up to $O(\alpha_s^3)$. In the $\overline{MS}$ scheme, the result is

$$\Gamma (H \to q\bar{q}) = \frac{3G_F M_H m_q^2}{4\pi \sqrt{2}} \left[1 - 4 \left(\frac{m_q^2}{M_H^2}\right)^{3/2} + C_F \frac{\alpha_s}{\pi} \left(\frac{17}{4} - 30 \frac{m_q^2}{M_H^2}\right) + K_2 \left(\frac{\alpha_s}{\pi}\right)^2\right],$$

where $K_2 \approx 35.9399 - 1.3586N_F$ with $N_F$ being the number of quark flavours active at scale $M_H$, and it is understood that $\alpha_s$ and $m_q$ are to be evaluated at $\mu = M_H$. We note in passing that Eq. (15) may be translated into the on-shell scheme by using the above relation between $M_q$ and $m_q(M_H)$. The difference between these two evaluations is extremely small, which indicates that the residual uncertainty due to the lack of knowledge of the $O(\alpha_s^2 m_q^2/M_H^2)$ and $O(\alpha_s^3)$ terms is likely to be inconsequential for practical purposes.

3.2. Next-to-leading order corrections to $\Gamma (H \to gg)$

The hadronic width of the Higgs boson receives contributions also from the $H \to gg$ channel, which is mediated by massive-quark triangles, and related higher-order processes. The respective partial width is well approximated by

$$\Gamma (H \to gg) = \Gamma (H \to gg) \left[1 + \frac{\alpha_s}{\pi} \left(\frac{95}{4} - \frac{7}{6}N_F\right)\right],$$

where $N_F$ is the number of quark flavours active at scale $M_H$, and

$$\Gamma (H \to gg) = \frac{\alpha_s^2 G_F M_H^4}{36 \pi^2 \sqrt{2}} \left(1 + \frac{7}{60} M_H^2\right).$$

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FIGURE CAPTIONS

Figure 1. Universal electroweak correction factor for $\Gamma (H \to f \bar{f})$ to $O(G_F M_H^2)$ and $O(G_F^2 M_H^4)$ as a function of $M_H$. In each order, the expanded result, Eq. (14), is compared with the calculation where the one-particle-reducible Higgs-boson self-energy diagrams are resummed, Eq. (13).