Leading Electroweak Correction to Higgs Boson Production at Proton Colliders

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At proton colliders, Higgs particles are dominantly produced in the gluon-gluon fusion mechanism. The Higgs-boson–gluon coupling is mediated by heavy quark loops, and the process can serve to count the number of heavy strongly interacting particles whose masses are generated by the Higgs mechanism. We present the two-loop leading electroweak radiative correction to this coupling, which is quadratically proportional to the heavy quark masses. It turns out that this correction is well under control across the physically interesting quark mass ranges.

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The fundamental particles, quarks, leptons, and gauge bosons, acquire their masses through the Higgs mechanism [1]. This mechanism requires the existence of at least one weak isodoublet scalar field, the self-interaction of which leads to a nonzero field strength in the ground state, inducing the spontaneous breaking of the SU(2) × U(1) electroweak symmetry down to the U(1) electromagnetic symmetry [2]. Among the four initial degrees of freedom, three Goldstone bosons will be absorbed to build up the longitudinal polarization states of the massive $W^\pm$ and Z bosons, and one degree of freedom will be left over, corresponding to a physical scalar particle, the Higgs boson.

The discovery of this particle is the most crucial test of the standard model and the search for it will be one of the most important missions of future high-energy colliders [3,4]. Unfortunately, in the standard model, the Higgs boson mass M_H is essentially a free parameter. The only information available is the lower limit $M_H > 63.8$ GeV [5] established from the negative results of the Higgs boson search in Z boson decays at the CERN e^+e^- collider LEP; this limit can be raised up to ~80 GeV in the second phase of LEP. However, from the requirement of vacuum stability and from the assumption that the standard model can be continued up to the grand unification scale, the Higgs boson mass could well be expected [6] in the window $80 < M_H < 180$ GeV, which is generally referred to as the intermediate mass range.

The dominant process for producing Higgs particles at proton colliders is the gluon-gluon fusion mechanism [7] $gg \to H$. The Hgg amplitude is built up by heavy quark triangular loops, Fig. 1. In the minimal standard model with three generations of fermions, the only significant contribution is the top quark. Since the quarks couple to the Higgs bosons proportionally to their masses, the loop particles will not decouple from the amplitude when they are much heavier than the Higgs boson. This coupling is therefore very interesting since it is sensitive to scales far beyond the Higgs boson mass and can be used as a possible "microscope" for new strongly interacting particles whose masses are generated by the Higgs mechanism. For instance, a fourth generation of fermions, the existence of which is still allowed by present experimental data [8] if the associated neutrino is heavy enough, would increase the $gg \to H$ production rate by an order of magnitude (see below).

To lowest order, the $gg \to H$ parton cross section can be expressed in terms of a form factor derived from the quark triangle diagram in Fig. 1,

$$\sigma^{\rm LO}(gg \to H) = \frac{G_F \alpha_S^2}{288\sqrt{2}\,\pi} \left| \sum_Q F_Q(\tau_Q) \right|^2 , \qquad (1)$$

with the form factor

$$F_Q(\tau_Q) = \frac{3}{2}\tau_Q^{-1}\left[1 + \left(1 - \tau_Q^{-1}\right)\arcsin^2\sqrt{\tau_Q}\,\right] \qquad (2)$$

approaching unity for quark masses slightly above half the Higgs boson mass, justifying the approximation of working in the limit $\tau_Q = M_H^2/4m_Q^2 \to 0$ already for $\tau_Q < 1$. In this limit, the two additional quarks of a fourth generation will give the same contribution to the amplitude, Eq. (2), as the top quark enhancing the production rate, Eq. (1), by a factor of 9. The Higgs boson production cross section for proton colliders is found by integrating the parton cross section, Eq. (1), over the gluons luminosity.

Because the precise knowledge of the $gg \to H$ production cross section is mandatory, quantum corrections must be included. The QCD corrections have been evaluated in Ref. [9] and found to be rather large, increasing the

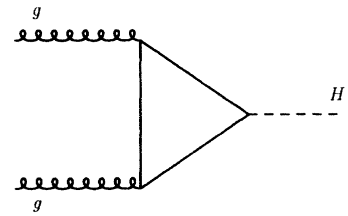

FIG. 1. The loop-mediated Higgs-boson–gluon–gluon coupling at lowest order.

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production rate by more than 50%. The next important radiative correction to the Hgg coupling, which is proportional to the square of the masses of the heavy quarks in the loop and is therefore potentially very large, is the two-loop $O(G_F m_Q^2)$ electroweak correction. In this Letter, we present the result for this leading correction. We will work in the limit $m_Q \to \infty$ since, as mentioned previously, this is a very good approximation for Higgs boson masses smaller than half the quark mass; this should hold at least for Higgs bosons in the intermediate mass range.

Note that up to color and electroweak charges factors, the quark contribution to the $HZ\gamma$ and $H\gamma\gamma$ couplings is the same as the one for the Hgg coupling. At future e^+e^- colliders, the $H\gamma\gamma$ amplitude can be precisely measured in the process $\gamma\gamma \to H$, the high-energy photons being generated by Compton backscattering of laser light [10]; this amplitude is also important since the $\gamma\gamma$ decay of the Higgs boson is the most promising detection channel of this particle at hadron colliders. The leading $O(G_F m_Q^2)$ correction to the Hgg amplitude presented here will be the same for the $H\gamma\gamma$ and $HZ\gamma$ amplitudes.

The technique that we use to calculate the two-loop $O(G_F m_Q^2)$ correction to the Hgg coupling has been known for some time [11,12]. Writing the basic Higgs-boson–quark Lagrangian as

$$\mathcal{L}(HQ\bar{Q}) = -\left(\sqrt{2}\,G_F\right)^{1/2} m_Q^0 H Q_0 \bar{Q}_0 , \qquad (3)$$

the Hgg coupling at small momentum can be derived from the condition that the matrix element $\langle gg|\theta_\mu^\mu|0\rangle$ of the trace of the energy-momentum tensor

$$\theta_\mu^\mu = (1 - \delta_2) m_Q^0 Q\bar{Q} + \frac{1}{2}\frac{\beta(\alpha_S)}{g_S} G_{\mu\nu}G^{\mu\nu} \qquad (4)$$

vanishes in the low-energy limit. Here, $G_{\mu\nu}$ is the gluon field strength tensor, $\alpha_S = g_S^2/4\pi$ with g_S the strong coupling constant, and $\beta(\alpha_S)$ is the QCD β function to which a quark contributes by an amount

$$\frac{\beta(\alpha_S)}{g_S} = \frac{\alpha_S}{6\pi}[1 + \delta_1], \qquad (5)$$

where the term δ_1 denotes the higher-order contribution. To evaluate this contribution at $O(\alpha_S G_F m_Q^2)$, one needs to consider the two-loop diagrams shown in Fig. 2 and the corresponding counterterms. In renormalizable gauges, the virtual scalar bosons exchanged in the loops correspond either to the Higgs boson or to the neutral and charged Goldstone bosons. Note that in the amplitude for a quark of a given flavor, the virtual exchange of the charged Goldstone boson will introduce the weak isospin partner of this quark.

The term δ_2 in Eq. (4) arises from a subtlty in the use of the low-energy theorem [12]: In renormalizing the $HQ\bar{Q}$ interaction, Eq. (3), the counterterm for the Higgs-boson–quark Yukawa coupling is not the $HQ\bar{Q}$ vertex with a subtraction at zero momentum transfer, $\Gamma_{HQ\bar{Q}}(q^2 = 0)$ (which is implicitly used in the low-energy theorem), but

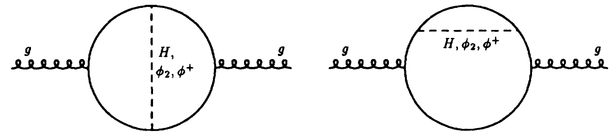

FIG. 2. Generic diagrams contributing to the QCD β function at $O(\alpha_S G_F m_Q^2)$.

rather is determined by the counterterms for the quark mass δm_Q and quark wave function Z_2^Q. This has to be corrected for, and one then has

$$\delta_2 = \left(Z_2^Q - 1\right) - \frac{\delta m_Q}{m_Q} + \Gamma_{HQ\bar{Q}}(q^2 = 0). \qquad (6)$$

Finally, one needs to include the renormalization of the Higgs boson wave function and the vacuum expectation value of the Higgs field. This is achieved by multiplying the one-loop generated Hgg coupling by a factor $1 + \delta_3$ where, in terms of the W and H boson vacuum polarization functions at zero momentum transfer, δ_3 reads

$$\delta_3 = -\frac{1}{2}\left[\frac{\Pi_{WW}(0)}{M_W^2} + \frac{\partial \Pi_{HH}(M_H^2)}{\partial M_H^2}\bigg|_{M_H^2 \to 0}\right]. \qquad (7)$$

The complete $O(G_F m_Q^2)$ correction to the effective Hgg coupling will be then given by

$$\mathcal{L}(Hgg) = \left(\sqrt{2}\,G_F\right)^{1/2} \frac{\alpha_S}{12\pi} H G_{\mu\nu}G^{\mu\nu}(1 + \delta), \qquad (8)$$

with

$$\delta = \delta_1 + \delta_2 + \delta_3, \qquad (9)$$

and the corrected $gg \to H$ cross section at this order will read

$$\sigma(gg \to H) = \sigma^{\rm LO}(gg \to H)[1 + 2\delta]. \qquad (10)$$

Using dimensional regularization, we have evaluated the contribution of the diagrams shown in Fig. 2 as well as those of the various one-loop self-energy and vertex functions which enter the counterterms in δ_1 and the terms δ_2 and δ_3, in the case of a weak isodoublet of heavy quarks (U, D) with masses $m_U \neq m_D$. The calculation has been performed in the on-shell scheme which is usually used in the electroweak theory [13]. In this scheme, the quark masses correspond to the poles of the quark propagators.

We have then specialized to two particular cases of physical relevance: (i) $m_U \gg m_D$ which corresponds to the approximate contribution of the top-bottom weak isodoublet since $m_t \sim 174$ GeV [14] is much larger than $m_b \sim 5$ GeV, and (ii) $m_U \sim m_D$ which corresponds to the contribution of an additional generation of fermions since in this case the mass splitting between the members of the extra weak isodoublet is highly constrained by electroweak precision measurements [8]. The lengthy results in the general case $m_U \neq m_D$ as well as the tedious details of the calculation will be given elsewhere [14]; in this short Letter we will simply present our final results in the two special cases of interest.

In the minimal standard model with three fermion families, the $O(\alpha_S G_F m_t^2)$ contribution to the top quark loop amplitude in the limit $m_t \gg m_b$ is given by

$$\delta = +\frac{G_F\sqrt{2}}{32\pi^2} m_t^2 . \quad (11)$$

Because of a large cancellation among the various δ_i contributions (in units of δ/m_t^2 one has $\delta_1 = -12$, $\delta_2 = +6$, and $\delta_3 = 7$), the total correction is very small for a value $m_t \sim 200$ GeV, which can be viewed as a conservative upper bound on the top quark mass [15], it amounts to a mere (positive contribution of) 0.2%. Therefore, contrary to the QCD corrections which have been found to be very large [9], the leading electroweak correction to the top quark loop mediated Higgs-boson–gluon coupling turns out to be very small. Note that the correction is free of infrared similarities for $m_b \to 0$, as required by the Kinoshita-Lee-Nauenberg theorem [16].

In the case of a fourth family of heavy quarks with degenerate masses, $m_U = m_D = m_Q$, the $O(G_F m_Q^2)$ correction to one of the quarks amplitude is given by

$$\delta = -\frac{G_F\sqrt{2}}{8\pi^2} m_Q^2 . \quad (12)$$

This negative correction will therefore screen the value of the one-loop generated Hgg coupling. However, the correction is rather small since for realistic values of the quark masses [17], $m_Q < 500$ GeV (obtained from the requirement that in the scattering of heavy quarks, weak interactions do not become strong and perturbation theory is still reliable), it does not exceed the 5% level. It is only for quark masses larger than ~ 2 TeV, for which perturbation theory breaks down already at the tree level [17], that the radiative correction will exceed the one-loop result. Therefore, the $O(G_F m_Q^2)$ correction to the Hgg amplitude is well under control for quark masses in the range interesting for perturbation theory, and the counting of new heavy quarks via the Hgg coupling will not be jeopardized by these radiative corrections.

Note that the calculation in the equal mass case has been first performed in Ref. [18]. However, only the irreducible contribution δ_1 (including quark mass, wave function, and vertex renormalizations with a subtraction at zero momentum transfer for the Higgs-boson–quark vertex) has been evaluated: the proper renormalization of the Higgs-boson–quark Yukawa coupling and the renormalizations of the Higgs boson wave function and vacuum expectation value have been omitted. As a consequence, the result of Ref. [18] is a factor of 3 larger compared to our result.

Note also that in the previous equation only the contribution of the heavy quarks of the fourth generation has been taken into account. Additional contributions will be induced by the extra weak isodoublet of leptons (with a right-handed component for the heavy neutrino, for the mass of the latter particle to be generated through the standard Higgs mechanism) via the renormalization of the Higgs boson wave function and the one of the vacuum expectation value of the Higgs field. If one assumes that the masses of the heavy leptons are approximately equal to those of the quarks, the total contribution of the weak isodoublets of quarks and leptons to the coefficient δ will be smaller by a factor of 3 than in Eq. (12).

Finally, we observe that in this equal mass case, the quark mass renormalization does not contribute to the amplitude in the limit $m_Q \to \infty$, and, therefore, the result for the correction δ is independent on the scheme in which the quark mass is defined. This can be understood by recalling that in this limit, the quark contribution to the one-loop amplitude decouples in the sense that there is no more dependence on the quark mass.

In conclusion, we have presented the two-loop leading $O(G_F m_Q^2)$ electroweak radiative correction to the Higgs-boson–gluon coupling. This coupling is very interesting since it is sensitive to sales far beyond the Higgs boson mass. In the case of the minimal standard model with only three fermion families, the correction to the heavy top quark contribution is very small: less than 0.2% for a top quark mass smaller than 200 GeV. If the standard model is extended to include a fourth generation of heavy fermions, the corrections to the additional quark loop amplitudes are well under control across the physically interesting quark mass ranges for perturbation theory, since in this case they do not exceed the 5% level.

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