The Higgs resonance shape in gluon fusion: Heavy Higgs effects

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Abstract

We study the influence of two–loop radiative corrections of enhanced electroweak strength on Higgs production at the LHC. We consider Higgs production by the gluon fusion mechanism, with the subsequent decay of the Higgs boson into a pair of Z bosons, and incorporate the resonance shape corrections up to order \( \left( \frac{g^2 m_H^2}{m_W^2} \right)^2 \). We take into account the full \( gg \to ZZ \) process and the \( q\bar{q} \to ZZ \) background, as well as the subsequent decay of the Z pair into leptons. We also discuss the theoretical uncertainty related to the use of the equivalence theorem in this process.
1 Introduction

The arguably most important task of the Large Hadron Collider is to provide evidence regarding the nature of the mechanism responsible for the spontaneous breaking of the electroweak symmetry. This is crucial for understanding the physics at the electroweak scale and may give hints about the dynamics underlying the standard model at higher energies.

The simplest model of spontaneous symmetry breaking, the Higgs sector of the minimal standard model, will be testable at the LHC up to an energy scale of \( \sim 1 \) TeV. This covers the mass range where the Higgs particle is believed to make sense as a fundamental field. This is related to the triviality of the Higgs sector, which is a nonperturbative problem. Lattice simulations of the \( \phi^4 \) theory seem to indicate an upper bound of the order of 650 GeV for \( m_H \), beyond which the Higgs mass would be larger than the cutoff scale. Therefore a strongly interacting Higgs sector would not exist. However, these results must be interpreted with caution. There are indications that they depend on the lattice regularization, and in fact the triviality bound may be considerably higher. Also the presence of a large Yukawa coupling may play a role because it relaxes the bound derived from the position of the Landau pole. Unfortunately, it is difficult to check this nonperturbatively because of notorious problems with treating fermions on the lattice in the presence of Yukawa interactions. One cannot exclude the possibility of a Higgs mass of the order of 1 TeV.

The Higgs boson is produced in hadron collisions via the gluon fusion and the vector boson fusion processes. With \( m_t \approx 180 \) GeV, the gluon fusion process dominates unless the Higgs is very massive, of the order of 1 TeV. Heavy Higgs bosons mainly decay into pairs of vector bosons. The decay channel \( H \rightarrow ZZ \rightarrow 2l^+l^- \) provides a clean signal if \( m_H \) is smaller than \( \sim 800 \) GeV. For a heavier Higgs, the channel \( H \rightarrow ZZ \rightarrow l^+l^-\nu\bar{\nu} \) must be included. The Higgs then appears as a Jacobian peak in the \( p_T \) distribution of the \( Z \) bosons. To approach the 1 TeV mass range, one needs to consider the \( H \rightarrow ZZ \rightarrow l^+l^- +2 \) jets and \( H \rightarrow W^+W^- \) channels, which are affected by large backgrounds.

The Higgs production mechanisms at hadron colliders were studied extensively at leading order. The NLO QCD corrections to the gluon fusion process were calculated recently, and indicate a substantial increase of the Higgs production. It is the purpose of this paper to study the NNLO corrections of enhanced electroweak strength to Higgs production via gluon fusion, with the Higgs subsequently decaying into a \( Z \) pair.
The large $m_H$ range is affected by potentially large electroweak radiative corrections due to the strong selfinteraction of the Higgs sector. For instance, the radiative corrections of increased electroweak strength to the Higgs decay into vector bosons are known at two–loop order [7], and are of the order of 24\% for $m_H \sim 900$ GeV. It is therefore legitimate to ask what is the effect of such corrections on the production and decay of heavy Higgs bosons.

A consequence of the Slavnov–Taylor identities of the theory is the equivalence theorem. This theorem relates in the high energy limit the amplitudes with external longitudinal vector bosons $V_L$ and the corresponding amplitudes with the $V_L$ replaced by Goldstone bosons [8]–[11]. The equivalence theorem is a useful tool for calculating leading contributions in the Higgs mass in a simple way, but within this context the resonance region should be treated carefully. The equivalence theorem holds order by order in perturbation theory, while in the resonance region a Dyson summation is needed which contains incomplete contributions from higher orders. If one fails to carry out the calculation consistently at the given order in the coupling constant, the result will differ from the corresponding leading $m_H$ contributions calculated in the full electroweak theory (see for instance ref. [12] and [14], and references therein).

Related to the large coupling domain is also the problem of unitarity violation when the Higgs mass exceeds certain limits. Unitarity breaking effects are always present within perturbation theory and are formally of higher order in the coupling constant. Numerically, they may become important if the coupling constant is large, and raise questions about the relevance of calculations of physical processes.

A few approaches were proposed in the literature to deal with such problems. To mention only two of them, ref. [13] proposes the use of a momentum dependent width of the Higgs in the resonant propagator instead of a constant width in order to reduce the size of the unitarity violations in longitudinal vector boson scattering. Ref. [15] notes that this procedure reproduces the correct result in the resonance region but has a bad behavior at higher energies, where the nonresonant propagator should be used. Therefore it proposes a summation of the full tree level Goldstone boson scattering amplitude rather than a Dyson summation of only the Higgs selfenergy in order to interpolate between the two regions.

These approaches amount in fact to changes in the incomplete contributions of higher order in the coupling constant. One should see them as leading order calculations. They are not complete at one–loop order, and their use is largely motivated by the fact that they lead to a solution which
has the desired numerical behaviour in the specific process considered. For instance, using the momentum dependent width in the Higgs resonant propagator gives a numerically well behaved amplitude in ref. [13], but would lead to an unwanted result in the case of fermion–fermion scattering. Indeed, the one–loop momentum dependent selfenergy used in ref. [13] would result in a shift in the position of the Higgs resonance in fermion scattering at leading order in $m_H$, which in fact is absent in the full one–loop radiative corrections [16]. In this case the constant width is obviously the better choice.

A systematic treatment would instead require a complete calculation of radiative corrections to the given process. This ensures both that the equivalence theorem is valid up to the desired order, and that the unitarity violating contributions are of a higher order in the coupling constant compared to the tree amplitude, which is the best one can do within perturbation theory.

This paper is devoted to the study of the two–loop radiative corrections of enhanced electroweak strength to the $gg \rightarrow ZZ$ process. By using the equivalence theorem, we derive in the following section the two–loop corrections to the shape of the Higgs resonance in the related process $t\bar{t} \rightarrow H \rightarrow zz$. We then study the influence of these corrections on Higgs searches by means of a Monte–Carlo simulation of the full $gg \rightarrow ZZ \rightarrow 4f$ process and of the $q\bar{q} \rightarrow ZZ \rightarrow 4f$ background at LHC.

## 2 Resonance shape corrections

In this section we derive the NNLO radiative corrections to the scattering process $t\bar{t} \rightarrow H \rightarrow zz$, where $z$ are Goldstone bosons. We are interested only in radiative corrections of enhanced electroweak strength, that is, corrections which grow like $m_H^2$ at one–loop order and like $m_H^4$ at two–loop.

In the following, we denote by $\lambda$ the selfcoupling of the Higgs field,

$$\lambda = \frac{g^2 m_H^2}{16\pi^2 m_W^2},$$

where $g^2 = 4\sqrt{2} m_W^2 G_F$, $G_F = 1.16637 \cdot 10^{-5} \text{GeV}^{-2}$, and $m_W = 80.22 \text{GeV}$.

The radiative corrections of leading order in $m_H$ to the $t\bar{t} \rightarrow H \rightarrow zz$ process are depicted in fig. 1. They consist of the radiative correction to the Yukawa coupling $V_{Ht\bar{t}}$, which is momentum independent [18, 13], the momentum dependent correction to the $Hzz$ vertex $V_{Hzz}(s)$, and the
correction to the Higgs propagator $\Pi(s)$. In the resonance region one needs to perform an all–order Dyson summation of the proper selfenergy. Far from the peak, this resonant propagator will differ from the nonresummed two–loop selfenergy insertion by contributions of order $\lambda^3$.

At leading order, the diagram of fig. 1 reads simply

$$\frac{1}{s - m_H^2 + i m_H \Gamma_H}$$

up to the tree level $H\bar{t}t$ and $Hzz$ couplings, which are constants independent of $s$. $\Gamma_H$ in eq. 2 is the tree level Higgs width. This expression is correct at leading order, i.e., $\mathcal{O}(\lambda^{-1})$ on the resonance and $\mathcal{O}(1)$ far from the resonance.

The radiative corrections change this expression into:

$$V_{H\bar{t}t} \frac{1}{s - m_H^2 + i m_H \Pi(s)} V_{Hzz}(s).$$

We are primarily interested in radiative corrections in the resonance region. We therefore perform a momentum expansion around $s = m_H^2$:

$$\Pi(s) = \Pi(m_H^2) + \frac{s - m_H^2}{m_H^2} \Pi' + \frac{(s - m_H^2)^2}{m_H^4} \Pi'' + \ldots$$

$$V_{Hzz}(s) = V_{Hzz}(m_H^2) + \frac{s - m_H^2}{m_H^2} V'_{Hzz} + \frac{(s - m_H^2)^2}{m_H^4} V''_{Hzz} + \ldots.$$ (4)

For treating the radiative corrections consistently as an expansion in the coupling constant $\lambda$, one notes that in the resonance region the quantity $s - m_H^2$ is of the order of $m_H \Gamma_H$, that is, $\mathcal{O}(\lambda)$. This must be taken into account together with the loop expansion of the quantities $V_{H\bar{t}t}$, $\Pi(m_H^2)$, $V_{Hzz}(m_H^2)$, $\Pi'$, $V'_{Hzz}$, $\ldots$. To have the complete NNLO amplitude, one needs to keep the following contributions:

$$V_{H\bar{t}t} = 1 + V^{(1-\text{loop})}_{H\bar{t}t} + V^{(2-\text{loop})}_{H\bar{t}t} + \mathcal{O}(\lambda^3)$$

$$V_{Hzz}(s) = 1 + V^{(1-\text{loop})}_{Hzz}(m_H^2) + V^{(2-\text{loop})}_{Hzz}(m_H^2) + \frac{s - m_H^2}{m_H^2} V'_{Hzz} + \mathcal{O}(\lambda^3)$$

$$\Pi(s) = \Pi^{(1-\text{loop})}(m_H^2) + \Pi^{(2-\text{loop})}(m_H^2) + \Pi^{(3-\text{loop})}(m_H^2) + \frac{s - m_H^2}{m_H^2} \Pi' + \mathcal{O}(\lambda^4).$$ (5)
The correction to the Yukawa coupling reads \[18\] \[19\]:

\[
V_{Hti}^{(1\text{-}loop)} = \left( \frac{13}{16} - \frac{\pi \sqrt{3}}{8} \right) \lambda = .132325 \lambda \\
V_{Hti}^{(2\text{-}loop)} = (-.26387 \pm 1.3 \cdot 10^{-4}) \lambda^2.
\] (6)

The correction to the $H_{zz}$ vertex was calculated in ref. \[6\]:

\[
V_{Hzz}^{(1\text{-}loop)}(m_H^2) = \left[ \frac{19}{16} + \frac{5 \pi^2}{48} - \frac{3 \sqrt{3} \pi}{8} + i \pi \left( \frac{\log 2}{4} - \frac{5}{8} \right) \right] \lambda \\
= (.17505951 - i 1.4190989) \lambda \]

\[
V_{Hzz}^{(2\text{-}loop)}(m_H^2) = - \left[ (.53673 \pm 4.1 \cdot 10^{-4}) + i (.32811 \pm 3.1 \cdot 10^{-4}) \right] \lambda^2.
\] (7)

By evaluating the one–loop diagrams which contribute to the $H \rightarrow zz$ decay, on finds also:

\[
V_{Hzz}^{(1\text{-}loop)} = \left[ 1 + \frac{\sqrt{3} \pi}{12} - \frac{5 \pi^2}{48} + i \pi \left( \frac{1}{8} \frac{\log 2}{4} \right) \right] \lambda \\
= (.42536605 - i .15169744) \lambda.
\] (8)

Let us now turn to the correction to the Higgs propagator.

At the order in which we are working, the quantity $\Pi(s)$ is the imaginary part of the proper Higgs selfenergy. The real part of the Higgs proper selfenergy at $s = m_H^2$ and its first derivative with respect to $s$ are absorbed into the Higgs mass counterterm and wave function renormalization, respectively. The real part of the selfenergy may contribute an imaginary piece to eq. 5 only at $O(\lambda^4)$, through the quantity $\Pi''(2\text{-}loop)$. Therefore $\Pi(m_H^2)$ is the total width of the Higgs including the appropriate radiative corrections.

At tree level, the main decay channels of a heavy Higgs are given by:

\[
\Gamma_{H \rightarrow W^+W^-}^{\text{(tree)}} = \frac{g^2 m_H^3}{64 \pi m_W^2} \left[ 1 - 4 \frac{m_W^2}{m_H^2} \right] \times \left[ 1 - 4 \frac{m_W^2}{m_H^2} + 12 \frac{m_W^4}{m_H^4} \right]^{1/2}
\]
\[ \Gamma_{H \rightarrow Z^0 Z^0}^{(\text{tree})} = \frac{g^2}{128\pi} m_H^3 \left( 1 - 4 \frac{m_Z^2}{m_H^2} \right)^{1/2} \times \left[ 1 - 4 \frac{m_Z^2}{m_H^2} + 12 \frac{m_{W^2}}{m_H^2} \right] \]

\[ \Gamma_{H \rightarrow t\bar{t}}^{(\text{tree})} = \frac{3g^2}{32\pi} \frac{m_H m_{t}}{m_W^2} \left[ 1 - 4 \frac{m_t^2}{m_H^2} \right]^{3/2} . \] (9)

The radiative corrections to the decay widths of eqns. 9 are known at two–loop level in the heavy Higgs approximation \([15, 19, 7]\):

\[ \Gamma_{H \rightarrow W^+ W^–, Z^0 Z^0} = \Gamma_{H \rightarrow W^+ W^–, Z^0 Z^0}^{(\text{tree})} \times \left[ 1 + \lambda \left( \frac{19}{8} + \frac{5\pi^2}{24} - \frac{3\sqrt{3}\pi}{4} \right) + \lambda^2 \left( .97103 \pm 8.2 \cdot 10^{-4} \right) \right] \]

\[ = \Gamma_{H \rightarrow W^+ W^–, Z^0 Z^0}^{(\text{tree})} \times \left[ 1 + .350119 \lambda + \left( .97103 \pm 8.2 \cdot 10^{-4} \right) \lambda^2 \right] \]

\[ \Gamma_{H \rightarrow t\bar{t}} = \Gamma_{H \rightarrow t\bar{t}}^{(\text{tree})} \times \left[ 1 + \lambda \left( \frac{13}{8} - \frac{\pi \sqrt{3}}{4} \right) - \lambda^2 \left( .51023 \pm 2.5 \cdot 10^{-4} \right) \right] \]

\[ = \Gamma_{H \rightarrow t\bar{t}}^{(\text{tree})} \times \left[ 1 + .264650 \lambda - \left( .51023 \pm 2.5 \cdot 10^{-4} \right) \lambda^2 \right] . \] (10)

Note that in eqns. 10 some incomplete subleading contributions are present in the radiative corrections, as discussed in ref. [7]. They appear if one multiplies the full tree level width by the radiative correction factor calculated in the leading \(m_H\) approximation. These terms are of the same order in the coupling constant as the theoretical uncertainty related to the use of the equivalence theorem while calculating radiative corrections. It is thus not possible to decide unambiguously whether it is better to keep them or to drop them without calculating the complete subleading contributions explicitly. Numerically, this ambiguity is at 1% level at most. As such it can be safely neglected.

With these results, one can identify in eq. 5:

\[ \Pi^{(1-\text{loop})}(m_H^2) + \Pi^{(2-\text{loop})}(m_H^2) + \Pi^{(3-\text{loop})}(m_H^2) = \Gamma_{H \rightarrow W^+ W^–} + \Gamma_{H \rightarrow Z^0 Z^0} + \Gamma_{H \rightarrow t\bar{t}} + \Gamma_{H \rightarrow 4w, wz, wz, 4z} \] (11)

with \(\Gamma_{H \rightarrow W^+ W^–}, \Gamma_{H \rightarrow Z^0 Z^0}\) and \(\Gamma_{H \rightarrow t\bar{t}}\) given by eqns. 10.
The term $\Gamma_{H \rightarrow 4w,wwzz,4z}$ originates from the three–loop cut diagrams shown in fig. 2. All other cut diagrams of three–loop selfenergy diagrams are already contained in $\Gamma_{H \rightarrow WW+W^+,Z^0Z^0}^{(2–loop)}$.

Including or dropping $\Gamma_{H \rightarrow t\bar{t}}$ in eq. 11 is in principle irrelevant at the order in which we are working because this is a contribution of higher order in the top quark mass. In practice this is not a large effect because the $t\bar{t}$ branching ratio is small and because of the partial cancellation between the one–loop and the two–loop corrections to this partial width.

The quantity $\Pi'_{\text{(2–loop)}}$, which also contributes to eq. 5, was calculated in ref. [16]:

$$\Pi'_{\text{(2–loop)}} = m_H \lambda^2 \frac{3\pi}{4} \left( 1 + \frac{\pi \sqrt{3}}{12} - \frac{5\pi^2}{48} \right) = 1.002245142 m_H \lambda^2 . \quad (12)$$

To evaluate $\Gamma_{H \rightarrow 4w,wwzz,4z}$, let us introduce the following notations:

$$\Gamma_{H \rightarrow 4z} = \varphi_{4z} \lambda^2 \Gamma_0 ,$$

$$\Gamma_{H \rightarrow 2z2w} = \varphi_{2z2w} \lambda^2 \Gamma_0 ,$$

$$\Gamma_{H \rightarrow 4w} = \varphi_{4w} \lambda^2 \Gamma_0 , \quad \Gamma_0 = \frac{3\pi}{8} m_H \lambda . \quad (13)$$

We calculate the dimensionless factors $\varphi$ by means of a Monte–Carlo integration over the four–body phase space of the cut diagrams of fig. 2.

The diagrams themselves were calculated through the equivalence theorem, therefore only the leading $m_H$ terms are kept at the level of the amplitude. However, when integrating over the phase space, we allow for a finite mass of the vector bosons by keeping the momenta of the outgoing Goldstone bosons at an invariant mass equal to the vector boson masses, $m_Z$ and $m_W$. This procedure is similar to the way we multiplied in eqns. 10 the full tree level width by the loop corrections obtained via the equivalence theorem. In both cases the result contains incomplete subleading terms from the phase space integration (eq. 10 contains also contributions from the transversal degrees of freedom of the vector bosons). This is motivated by the assumption that the phase space factor is numerically the main subleading effect. To check the reliability of the equivalence theorem calculation of this process, we calculated the decay $H \rightarrow 4Z^0$ in the electroweak
However, the equivalence theorem and the exact $H$ at three–loop level, where diagrams with four massless particles on the cut lines are produced in pairs. In this case the equivalence theorem works better than the exact massless particles on the cut lines because the Goldstone bosons are always loop calculations. At two–loop level there are only cut diagrams with two massless particles because the Goldstone bosons are always produced in pairs. In this case the equivalence theorem works better than the exact massless particles on the cut lines.

Table 1: The value of the factors $\varphi$ calculated via the equivalence theorem ($\varphi^{(ET)}$) and in the electroweak theory, as a function of the Higgs mass. For $\varphi^{(ET)}$, the momenta of the outgoing Goldstone bosons have invariant masses $m_Z$ and $m_W$. $\varphi^{(Z_L^0)}_{4z}$ corresponds to the decay $H \rightarrow 4Z_L^0$ calculated in the electroweak theory in the approximation $\epsilon_L(p_\mu) \approx p_\mu/m_Z$. $\varphi^{(Z^0)}_{4z}$ is the exact result, summed over the polarization states of the $Z$ bosons.

\[
\begin{array}{|c|c|c|c|c|c|}
\hline
m_H[GeV] & \varphi^{(ET)}_{4w} & \varphi^{(ET)}_{2z2w} & \varphi^{(ET)}_{4z} & \varphi^{(Z_L^0)}_{4z} & \varphi^{(Z^0)}_{4z} \\
\hline
600 & 4.14 \cdot 10^{-5} & 1.88 \cdot 10^{-5} & 8.97 \cdot 10^{-6} & 9.07 \cdot 10^{-6} & 1.69 \cdot 10^{-6} \\
700 & 6.42 \cdot 10^{-5} & 3.34 \cdot 10^{-5} & 1.58 \cdot 10^{-6} & 1.67 \cdot 10^{-6} & 4.67 \cdot 10^{-6} \\
800 & 8.47 \cdot 10^{-5} & 4.81 \cdot 10^{-5} & 2.22 \cdot 10^{-6} & 2.38 \cdot 10^{-6} & 8.74 \cdot 10^{-6} \\
900 & 1.02 \cdot 10^{-4} & 6.24 \cdot 10^{-6} & 2.79 \cdot 10^{-6} & 3.01 \cdot 10^{-6} & 1.33 \cdot 10^{-5} \\
1000 & 1.18 \cdot 10^{-4} & 7.53 \cdot 10^{-5} & 3.25 \cdot 10^{-6} & 3.51 \cdot 10^{-6} & 1.78 \cdot 10^{-5} \\
1100 & 1.31 \cdot 10^{-4} & 8.70 \cdot 10^{-5} & 3.66 \cdot 10^{-6} & 3.93 \cdot 10^{-6} & 2.22 \cdot 10^{-5} \\
1200 & 1.41 \cdot 10^{-4} & 9.65 \cdot 10^{-5} & 4.01 \cdot 10^{-6} & 4.29 \cdot 10^{-6} & 2.62 \cdot 10^{-5} \\
1500 & 1.64 \cdot 10^{-4} & 1.20 \cdot 10^{-4} & 4.73 \cdot 10^{-6} & 5.00 \cdot 10^{-6} & 3.59 \cdot 10^{-5} \\
2000 & 1.87 \cdot 10^{-4} & 1.45 \cdot 10^{-4} & 5.41 \cdot 10^{-6} & 5.61 \cdot 10^{-6} & 4.61 \cdot 10^{-5} \\
3000 & 2.05 \cdot 10^{-4} & 1.66 \cdot 10^{-4} & 5.93 \cdot 10^{-6} & 6.04 \cdot 10^{-6} & 5.51 \cdot 10^{-5} \\
5000 & 2.16 \cdot 10^{-4} & 1.80 \cdot 10^{-4} & 6.24 \cdot 10^{-6} & 6.29 \cdot 10^{-6} & 6.08 \cdot 10^{-5} \\
\infty & 2.23 \cdot 10^{-4} & 1.89 \cdot 10^{-4} & 6.41 \cdot 10^{-6} & 6.41 \cdot 10^{-6} & 6.41 \cdot 10^{-5} \\
\hline
\end{array}
\]

theory. We calculated the decay $H \rightarrow 4Z_L^0$ in the approximation that the longitudinal polarization vectors of the $Z$ are given by $\epsilon_L(p_\mu) \approx p_\mu/m_Z$, as well as the complete $H \rightarrow 4Z^0$ decay.

The results are given in table 1. The limit $m_W = m_Z = 0$, in which the equivalence theorem is exact, corresponds to $m_H \rightarrow \infty$. One can see that the equivalence theorem with the phase-space factors included approximates well the electroweak decay width in the approximation $\epsilon_L(p_\mu) \approx p_\mu/m_Z$. However, the equivalence theorem and the exact $H \rightarrow 4Z^0$ result start to agree only for $m_H$ of the order of 2.5—3 TeV. The equivalence theorem is not a good approximation for $m_H \sim 1$ TeV because $m_H$ is not large enough compared to the four–particle threshold at $\sim 360$ GeV.

This is to be contrasted to the situation encountered in one– and two–loop calculations. At two–loop level there are only cut diagrams with two massless particles on the cut lines because the Goldstone bosons are always produced in pairs. In this case the equivalence theorem works better than at three–loop level, where diagrams with four massless particles on the cut
lines are allowed, and the threshold in the full electroweak theory is higher.

Generally, for given masses of the vector bosons and for fixed external momenta, one expects the equivalence theorem to become a progressively bad approximation as the number of loops increases. The subleading contributions become increasingly important because of the presence of multi-particle cuts.

In this calculation, the uncertainty related to subleading contributions in the four Goldstone decay is numerically negligible. Compared to the two Goldstone channel, the four Goldstone decay is strongly suppressed by phase space.

We have now all ingredients needed to evaluate eq. 3 at NNLO. At lowest order, the expression which we obtain is equivalent to eq. 2. At the same time, it incorporates the full radiative corrections of $\mathcal{O}(\lambda)$ and $\mathcal{O}(\lambda^2)$. It also contains some incomplete subleading contributions.

In fig. 3 we compare the LO expression to the NNLO calculation.

Although the purpose of this calculation was primarily to obtain the radiative corrections to the shape of the resonance in a consistent way, one notes that the amplitude is well behaved far from the resonance as well.

Compared to the corrections to the Higgs width, the resonance shape corrections of fig. 3 are relatively small. This is because of the partial compensation of the effect of the radiative corrections to $\Pi(s)$ and to $V_{Hzz}(s)$. For the same reason, the shift of the resonance peak towards lower momentum is smaller than the shift in fermionic scattering [16].

Finally, let us notice that the NNLO is the lowest order where nontrivial corrections to the resonance shape occur. At NLO there is only a correction to the Higgs width because the momentum dependence of the quantities $\Pi$ and $V_{Hzz}$ is of higher order in $\lambda$.

### 3 Gluon fusion at hadron colliders

The diagrams which contribute to the process $gg \to ZZ$ are shown in fig. 4. Among these, only the Higgs production diagram receives radiative corrections of enhanced electroweak strength.

The leading $m_H$ radiative corrections to the Higgs production diagram of fig. 4 a) are identical to the corrections to the process $t\bar{t} \to H \to zz$, which were discussed in the previous section, by the equivalence theorem. This is true in the limit of large top mass as well. In this limit the top quark does not decouple, and the $gg \to H$ subgraph results in the effective
interaction $\mathcal{L}_{ggH}^{\text{eff}} = g \alpha_s/(24 \pi m_W) H C_{\mu\nu}^a G^a_{\mu\nu}$, which receives the same radiative corrections at leading order in $m_H$ as the Yukawa coupling.

We incorporated the radiative corrections of the previous section in a Monte–Carlo event generator to calculate the $gg \rightarrow ZZ$ process and the $q\bar{q} \rightarrow ZZ$ background at LHC. The subsequent decay of the $Z$ pair into fermions is included in the narrow width approximation by using the density matrix formalism. Details of the calculation can be found in ref. \cite{20, 21}.

We take $\sqrt{s} = 14.5$ TeV for the CM energy of the LHC, and $m_t = 180$ GeV for the mass of the top quark. We use the MHRS parton distribution functions \cite{22}, with $\Lambda = 190$ MeV, evolved to the scale $Q^2 = s/4$, and $\alpha_S = 12\pi/(23 \log(Q^2/\Lambda^2))$ corresponding to five flavours. As a rough simulation of the detector geometry, we impose a rapidity cut on the outgoing leptons of $|y_l| < 2.5$, and request that the leptons have a transverse momentum larger than 20 GeV. The cross sections shown in the following correspond to four muons in the final state. The branching ratio should be multiplied by a factor 4 for muons or electrons in the final state, and 24 for $l^+l^-\nu\bar{\nu}$ (with these cuts, the cross sections for neutrinos in the final state will be underestimated to some extent). One can consider an integrated luminosity of $10^2 fb^{-1}$ to interpret the results.

The invariant mass spectrum of the $Z$ pair is shown in fig. 5 for different values of the Higgs mass. The background refers to $m_H \rightarrow \infty$. The transverse momentum distribution of the $Z$ bosons is shown in fig. 6. One notices that the way the radiative corrections influence the shape of the Higgs resonance differs from the effect shown in fig. 3 for the process $t\bar{t} \rightarrow H \rightarrow zz$. In particular, the cross section on top of the resonance is slightly increased. This is due to interference effects with the nonresonant diagrams in fig. 4.

4 Conclusions

We studied the effects of the radiative corrections of enhanced electroweak strength on Higgs production by gluon fusion at hadron colliders. The full correction to the Higgs resonance shape at next-to-next-to-leading order was derived. This was then incorporated in a Monte–Carlo simulation of the processes $gg \rightarrow ZZ$ and $q\bar{q} \rightarrow ZZ$ at LHC.

Compared to the similar radiative corrections to the $H \rightarrow V_LV_L$ decay, the Higgs resonance shape corrections turn out to be relatively small. This is due to cancellations between different contributions. Other than for the $t\bar{t} \rightarrow H \rightarrow zz$ scattering, in the gluon fusion process $gg \rightarrow H \rightarrow ZZ$ the
radiative corrections result in an increase of 10—20% of the Higgs signal due to interference effects with the nonresonant diagrams.

We did not address the question of the other Higgs production mechanism, the vector boson fusion. The cross section of this process becomes more important as the mass of the Higgs boson increases. For the vector boson scattering process, one–loop results are available \[23\] in the large \(m_H\) limit. Results at two–loop order in the large momentum limit also exist \[24\]. A complete two–loop analysis would be difficult to perform even via the equivalence theorem because one needs to evaluate two–loop box diagrams at finite external momenta. This problem deserves further investigation.

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Figure captions

Fig.1 The radiative corrections of enhanced electroweak strength to the process $t\bar{t} \rightarrow H \rightarrow zz$.

Fig.2 Cut diagrams which contribute to the imaginary part of the Higgs selfenergy at three–loop order. These contributions are not included in the two–loop $H \rightarrow zz, w^+w^−$ decay width.

Fig.3 Radiative corrections to the Higgs resonance shape in the process $t\bar{t} \rightarrow H \rightarrow zz$ for $m_H = 850$ GeV. The solid line is the squared absolute value of the leading order (eq. 2). The thin line corresponds to the NNLO expression (eq. 3).

Fig.4 Feynman diagrams for $Z^0$ pair production by gluon fusion.

Fig.5 Invariant mass distribution of the $Z^0$ pairs at LHC. The processes considered are $gg \rightarrow ZZ \rightarrow 2(\mu^+\mu^-)$ and $q\bar{q} \rightarrow ZZ \rightarrow 2(\mu^+\mu^-)$. We take $\sqrt{s} = 14.5$ TeV, and for the outgoing muons we request $p_T > 20$ GeV and $|y_l| < 2.5$. The solid line is the NNLO cross section, the dashed line is the tree level cross section, and the dotted line is the background (no Higgs production diagram). a) shows the total cross section, and b) shows the Higgs signal, with the background subtracted.

Fig.6 Transverse momentum distribution of the $Z^0$ bosons at LHC. The processes considered are $gg \rightarrow ZZ \rightarrow 2(\mu^+\mu^-)$ and $q\bar{q} \rightarrow ZZ \rightarrow 2(\mu^+\mu^-)$. We take $\sqrt{s} = 14.5$ TeV, and for the outgoing muons we request $p_T > 20$ GeV and $|y_l| < 2.5$. The solid line is the NNLO cross section, the dashed line is the tree level cross section, and the dotted line is the background (no Higgs production diagram). a) shows the total cross section, and b) shows the Higgs signal, with the background subtracted.