HIGGS PRODUCTION AND DECAY

The Higgs boson can only be detected through its decay products. Being the SM Higgs a scalar particle, and with a rather small width ($\Gamma \sim 4\,MeV$), one can treat to a very good approximation the production and decay mechanisms as factorized.

Indeed we can write

$$d\sigma(I \to H \to F) = \frac{1}{2s} |\mathcal{M}(I \to H)|^2 \frac{1}{(s-m_H^2)^2 + \Gamma_H^2 m_H^2} |\mathcal{M}(H \to F)|^2 \, d\Phi_F$$

In the narrow width approximation we can write

$$\frac{1}{(s-m_H^2)^2 + \Gamma_H^2 m_H^2} \simeq \frac{\pi}{\Gamma_H m_H} \delta(s - m_H^2)$$

[the constant is fixed by imposing that the normalization is the same

$$\int_{-\infty}^{+\infty} \frac{1}{(s-m_H^2)^2 + \Gamma_H^2 m_H^2} = \frac{\pi}{\Gamma_H m_H}$$]

We can thus write

$$\sigma(I \to H \to F) = \frac{1}{2s} \int |\mathcal{M}(I \to H)|^2 \frac{\pi}{\Gamma_H m_H} |\mathcal{M}(H \to F)|^2 \delta(s - m_H^2) \, d\Phi_F$$

$$= \frac{1}{2s\,\Gamma_H} \int |\mathcal{M}(I \to H)|^2 \, 2\pi \, \delta(s - m_H^2) \int \frac{1}{2m_H} |\mathcal{M}(H \to F)|^2 d\Phi_F$$

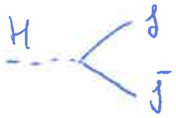
$$= \frac{1}{\Gamma_H} \sigma(I \to H) \, \Gamma(H \to F) = \sigma(I \to H) \, Br(H \to F)$$

where we have defined the BRANCHING RATIO $Br(H \to F) = \dfrac{\Gamma(H \to F)}{\Gamma_H}$

HIGGS BOSON DECAYS

As discussed at the beginning of this course, in the SM, once the Higgs boson mass is fixed, its profile is uniquely determined. The couplings to the gauge bosons and the fermions, in particular, are uniquely determined, and are proportional to their masses. As such, the Higgs decays into the heaviest SM particles allowed by kinematics. The case $m_H \simeq 125$ GeV is particularly interesting, because it will allow to study many different (and interesting) decay modes.

$H \to f\bar{f}$

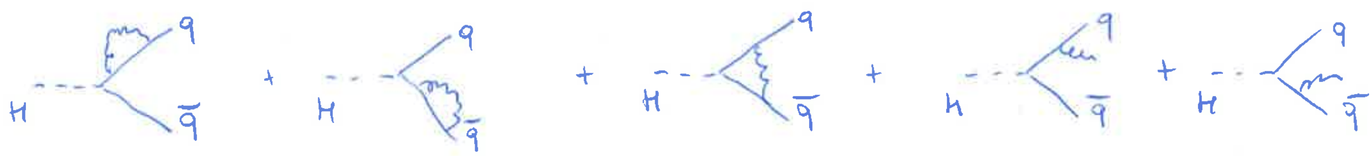


The decay into a fermion pair is a simple two-body decay and the corresponding width is

$$\Gamma = \frac{G_F}{4\sqrt{2}\pi}\, m_H\, m_f^2\, \beta_f^3 \qquad \beta_f = \sqrt{1-\frac{4m_f^2}{m_H^2}}$$

In the case in which the fermion is a quark this result has to be multiplied by a factor of 3. The quantity β_f is the fermion velocity, and it vanishes in the threshold limit. The suppression factor β_f^3 is typical of a scalar particle (for a pseudoscalar the width would be proportional to β). We now focus on $f=q$.

The QCD corrections involve as usual virtual and real diagrams



It is interesting to consider the limit $m_q \to 0$. According to the KLN theorem one expects that, at each order in perturbation theory, there are no collinear singularities in the UNRENORMALIZED decay rate. However the collinear singularity is reintroduced through the renormalization procedure. In a spontaneously broken gauge theory in which the fermion acquires mass from the Higgs mechanism, the mass of the fermion and its coupling to the Higgs boson cannot be renormalized independently [Braaten-Leveille, 1980]. The collinear singularity is thus reintroduced through the renormalization procedure.

One finds $$\Gamma_{NLO} \approx \Gamma_{LO}\left(1 + C_F \frac{\alpha_s}{\pi}\left(\frac{9}{4} + \frac{3}{2}\log\frac{m_q^2}{m_H^2}\right)\right)$$

The limit $m_q \to 0$ is still smooth (due to the overall Yukawa factor present in Γ_{LO}) but at small m_q the large logarithmic term can be important and can turn the width to negative values. However the large logarithm can be absorbed in the redefinition of the quark mass $m_q \to m_q(m_H)$. After that the effect of NLO corrections is to increase the LO rate by about 20%.

The QCD corrections are now known to $O(\alpha_s^4)$ and can be written as [Baikov et al 2005]

$$\Gamma(H\to q\bar{q}) = \frac{3 G_F m_H}{4\sqrt{2}\pi}\, m_q^2(m_H)\, \tilde{R}$$

where

$$\tilde{R} = 1 + 5.6666\,\frac{\alpha_s}{\pi} + \left(35.94 - 1.359\, n_F\right)\left(\frac{\alpha_s}{\pi}\right)^2 + \left(164.14 - 25.77\, n_F + 0.259\, n_F^2\right)\left(\alpha_s/\pi\right)^3$$
$$+ \left(39.34 - 220.9\, n_F + 9.685\, n_F^2 - 0.0205\, n_F^3\right)\left(\frac{\alpha_s}{\pi}\right)^4$$

$H \to WW, ZZ$

The width for the Higgs decay into on-shell vector bosons is

$$\Gamma(H\to ZZ) = \frac{G_F m_H^3}{16\sqrt{2}\pi}\sqrt{1-x_Z}\left(1 - x_Z + \frac{3}{4}x_Z^2\right) \qquad x_V = \frac{4m_V^2}{m_H^2}$$

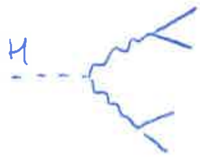
$$\Gamma(H\to WW) = \frac{G_F m_H^3}{8\sqrt{2}\pi}\sqrt{1-x_W}\left(1 - x_W + \frac{3}{4}x_W^2\right)$$

The leading term as $x_V \to 0$ (large Higgs mass) goes like m_H^3 $\Rightarrow$ that's why a heavy Higgs boson becomes obese! This is a consequence of the fact that the decay is dominated by longitudinally polarized gauge bosons.

One finds $$\frac{\Gamma(H\to V_T V_T)}{\Gamma(H\to V_L V_L)} = \frac{\frac{1}{2}x_V^2}{(1-x_V/2)^2}$$

At high m_H, when the phase space factors can be ignored, one simply has $\Gamma_{WW}\sim 2\Gamma_{ZZ}$

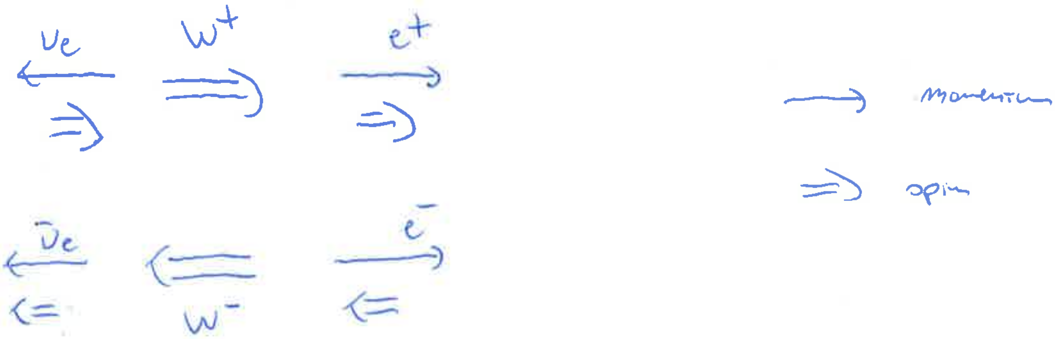
In the region $m_H\sim 125$ GeV we are clearly below the WW and ZZ threshold and one has to consider the full 4 fermion decay. NLO QCD+EW corrections to this decay process



have been computed, and are included in the PROPHECY4f program. The decays $H\to WW^*\to \ell\nu\ell\nu$ and $H\to ZZ^*\to 4\ell$ are characterized by a definite pattern of spin and angular correlations that allow to discriminate the spin/CP properties of the Higgs boson.

For the moment we consider the $H\to WW\to \ell\nu\ell\nu$ decay and assume $m_H\sim 2m_W$.

In this case the scalar nature of the Higgs, together with the V−A decay structure of the W bosons allow us to conclude that the two charged leptons must be very close in angle



as a consequence, a boost invariant quantity which turns out to be very sensitive to this behaviour is the $\Delta\phi_{ll}$: azimuthal separation of the charged leptons in the transverse plane. The potential of spin correlations in the $H\to WW\to \ell\nu\ell\nu$ decay mode was already pointed out by C. Nelson [1988]

The first realistic study using $\Delta\phi_{ll}$ was carried out by Dittmar-Dreiner [1996]

Loop-induced decays

The Higgs boson does not directly couple to $\gamma\gamma$ and gg. However the $H\gamma\gamma$, Hgg and also $HZ\gamma$ vertices are generated at quantum level.

The contribution of heavy particles in the triangular loops does not decouple (the conditions of the DECOUPLING THEOREM are not satisfied since the couplings with the Higgs grow with the mass)

$\Rightarrow$ this implies that these decays are very interesting as they are sensitive to scales well above m_H

- $H \to \gamma\gamma$

The $H \to \gamma\gamma$ decay is sensitive to both colored and colorless particles in the loop. The corresponding width can be written as

$$\Gamma(H \to \gamma\gamma) = \frac{G_F \alpha^2 m_H^3}{128\sqrt{2}\pi^3} \left| \sum_q 2 N_c Q_q^2 A_{1/2}(\tau_q) + A_1(\tau_W) \right|^2$$

where we have considered only the contribution of the heavy quarks and of the W boson. The function $A_{1/2}(\tau_q)$ is the same encountered in $gg \to H$ and has the form

$$A_{1/2}(\tau_Q) = \tau_q \left(1 + (1-\tau_q) f(\tau_q)\right) \qquad \tau_q = \frac{4 m_q^2}{m_H^2}$$

$$f(\tau_q) = \begin{cases} \arcsin^2 \sqrt{\frac{1}{\tau_Q}} & \tau_Q \geq 1 \\ -\frac{1}{4}\left(\log \frac{1+\sqrt{1-\tau_Q}}{1-\sqrt{1-\tau_Q}} - i\pi\right)^2 & \tau_Q < 1 \end{cases}$$

The function $A_1(\tau_W)$ is instead

$$A_1(\tau_W) = -\left[2 + 3\tau_W\left(1 + (2-\tau_W) f(\tau_W)\right)\right] \qquad \tau_W = \frac{4m_W^2}{m_H^2}$$

In the limit $\tau_W \to \infty$ we have $\quad A_1(\tau_W) \to -7$

For $m_H = 125$ GeV, $m_W = 80.385$ GeV we have $\quad A_1(\tau_W) \approx -8.32$

⇒ not too different from the $\tau_W \to \infty$ limit!

Since $A_{1/2}(\tau_t) \approx 1.83$ we have

$$\left|2 N_c Q_t^2 A_{1/2}(\tau_t) + A_1(\tau_W)\right| \approx \left|1.83 - 8.32\right|$$

⇒ the top quark and the W boson interfere destructively and the W contribution is about 4.5 times larger than the top contribution.

In the limit $m_H \to 0$ the $H \to \gamma\gamma$ amplitude can be computed by using an effective lagrangian obtained by integrating out the top and W fields. To obtain such effective lagrangian we can use the low-energy theorem used for the Hgg interaction. We have to evaluate the contribution of charged massive particles to the on-shell photon self-energy.

The result can be written with the effective lagrangian

$$\mathcal{L}_{\gamma\gamma} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} \left[1 + \Pi^t_{\gamma\gamma}(0) + \Pi^W_{\gamma\gamma}(0)\right]$$

where $F_{\mu\nu}$ is the electromagnetic tensor and

$$\Pi^t_{\gamma\gamma}(0) = N_c Q_t^2 \frac{\alpha}{3\pi} \left(\frac{4\pi\mu^2}{m_t^2}\right)^{\epsilon} \frac{\Gamma(1+\epsilon)}{\epsilon}$$

$$\Pi^W_{\gamma\gamma}(0) = -\frac{\alpha}{4\pi}\left(\frac{4\pi\mu^2}{m_W^2}\right)^{\epsilon}\Gamma(1+\epsilon)\left[\frac{7}{\epsilon}+\frac{2}{3}+O(\epsilon)\right]$$

From here, by applying the low-energy theorem we find

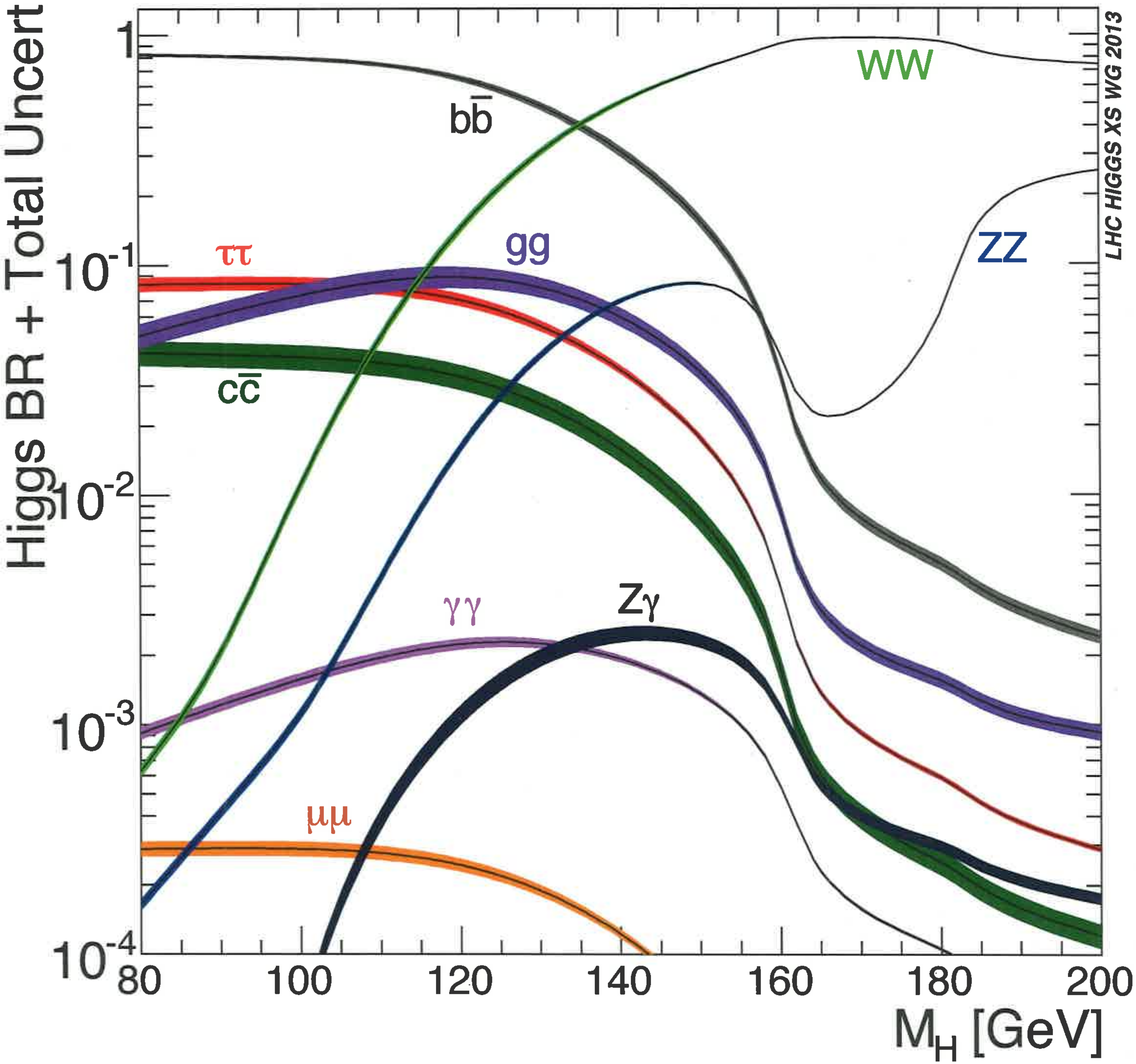
$$\mathcal{L}_{H\gamma\gamma} = \frac{\alpha}{2\pi} F^{\mu\nu}F_{\mu\nu}\,\frac{H}{v}\left(N_c\frac{Q_t^2}{3}-\frac{7}{4}\right)$$

which is valid for $m_H \ll m_t, m_W$. This expression is in agreement with the $m_H \to 0$ limit of the result discussed before.

The NLO corrections to the $H\to\gamma\gamma$ width consist only of VIRTUAL CORRECTIONS ($H\to\gamma\gamma+g$ is forbidden by color conservation). For $m_H = 125$ GeV their effect is completely cancelled by ELECTROWEAK CORRECTIONS. The QCD corrections have been evaluated up to 3-loop order but the complete effect (QCD+EW) is well below 1%. The $H\to\gamma\gamma$ decay mode is a discovery mode at the LHC due to the very clean signature. However the corresponding branching ratio is extremely small. ($BR(H\to\gamma\gamma)\sim 10^{-3}$).

$H\to Z\gamma$

The $H\to Z\gamma$ decay mode is less important than $\gamma\gamma$. It provides a comparable rate with respect to $\gamma\gamma$ but pays the additional suppression of the $Z\to e^+e^-$ decay mode, and will thus be important only at high luminosity. The NLO QCD corrections are known and are below 1%.


Higgs BR + Total Uncert
1
10^{-1}
10^{-2}
10^{-3}
10^{-4}
80
100
120
140
160
180
200
M_H [GeV]
$b\bar{b}$
WW
ZZ
$\tau\tau$
gg
$c\bar{c}$
$\gamma\gamma$
$Z\gamma$
$\mu\mu$
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