Higgs boson plus photon production at the LHC: a clean probe of the b-quark parton densities

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Abstract

Higgs boson production in association with a high $p_T$ photon at the CERN Large Hadron Collider is analyzed, in the framework of the MSSM model, for the heavier neutral Higgs bosons. The request of an additional photon in the exclusive Higgs boson final state selects $b$-quark pairs among the possible initial partonic states, since gluon-gluon initial states are not allowed by C-parity conservation. Hence, the measurement of cross sections for neutral Higgs boson plus photon production can provide a clean probe of the $b$-quark density in the proton as well as of the $b$-quark electromagnetic coupling. The suppression of the production rates by the $b$-quark electromagnetic coupling can be compensated by the enhanced Higgs boson Yukawa coupling to $b$’s in the large $\tan\beta$ regime. The Higgs boson decay into a tau-lepton pair is considered, and irreducible backgrounds with corresponding signal significances are evaluated.

The search for the Higgs boson, that is at the origin of the electroweak symmetry breaking and fermion mass generation, remains one of the main tasks of the Large Hadron Collider at CERN. Although the Higgs boson is still eluding any direct experimental test, electroweak precision tests requires, for data consistency in the Standard Model (SM), a mass of the Higgs boson close to the LEP2 experimental bound of 114 GeV[1].

A light Higgs-boson mass scenario is naturally predicted in the framework of the minimal supersymmetric extension of the SM (MSSM), with a mass range close to the present experimental bound. In the MSSM the Higgs boson sector contains two complex Higgs doublets, leading to two CP-even ($h, H$) and one CP-odd ($A$) neutral Higgs bosons, and a pair of charged Higgs bosons. At tree level, the Higgs sector of MSSM is fully specified by 2 parameters (e.g., $\tan\beta$, the ratio of the up- and down-Higgs doublet vacuum expectation values, and the mass of the pseudoscalar $m_A$),
while higher-order radiative corrections to the effective potential, and hence to the masses, depend on other relevant SUSY parameters. Large tan $\beta$ scenarios are favored in the MSSM, since they can easily provide both a viable Dark Matter candidate and large radiative corrections to the lightest Higgs boson mass. Other indications in this direction comes from the SUSY contribution to the anomalous magnetic moment of the muon $(g - 2)_{\mu}[2]$.

The $A/H$ boson Yukawa couplings to $b$-quarks and tau leptons are also enhanced by a large tan $\beta$. The CDF and D0 collaborations at Tevatron have recently analyzed the process $p\bar{p} \rightarrow \phi \rightarrow \tau^+\tau^-$, where $\phi = \{h, H, A\}$ [3, 4]. Although no evidence for a Higgs boson signal has been reported, the limits obtained for tan $\beta$ and $m_A$ by the CDF collaboration turn out to be weaker than expected, in the mass region $m_A \sim 120 - 160$ GeV. This could be a hint of a MSSM scenario with $m_A \sim 160$ GeV and tan $\beta \gtrsim 45$ [5].

In the large tan $\beta$ regime, the $b$-quark fusion process becomes one of the main mechanisms for heavy Higgs boson production at hadron colliders. While, in the SM, the $b\bar{b} \rightarrow h$ cross section at the LHC is at least two orders of magnitude smaller that the main gluon fusion cross section, in the MSSM the $b$ fusion rates are comparable to the $gg \rightarrow A/H$ ones, even at moderate values of tan $\beta$ [6]. In particular, for $m_{A,H} \gtrsim 150$ GeV and tan $\beta \gtrsim 5$, the bulk of the $A/H$ bosons produced at the LHC comes either from $gg \rightarrow A/H$ or from $b\bar{b} \rightarrow A/H$, the latter contributing dramatically to the Higgs discovery especially at large values of $m_{A,H}$ and tan $\beta$.

The $b$ fusion process is particularly sensitive not only to the Higgs coupling to the $b$ quark, but also to the actual composition of the proton in terms of the $b$-quark partons. At present, the $b$-quark parton density is derived perturbatively from the gluon parton density, and there is no direct measurement of it. In fact, the uncertainty in the gluon content of the proton propagates to the $b$ component prediction. In the SM, one is planning to measure the $b$-quark parton densities through processes containing one $b$-jet in the final state, such as $bg \rightarrow bZ/b\gamma$ [7]. The $bg \rightarrow bZ/b\gamma$ cross sections are linearly dependent on the $b$-quark densities. A measurement of the $b\bar{b} \rightarrow h$ production rate would be remarkably more sensitive, but there is no hope in the SM of separating the $b\bar{b} \rightarrow h$ signal from the $gg \rightarrow h$ production.

On the contrary, in the MSSM, the inclusive heavy Higgs production is predicted to be quite sensitive to the $b\bar{b} \rightarrow h$ contribution, although a non-negligible contamination from $gg \rightarrow h$ could make the extraction of the information on the $b$-quark densities more involved for large portions of the MSSM parameter space.

The aim of the present work is to study the exclusive process $pp \rightarrow \phi \gamma$, with $\phi = \{h, H, A\}$, at the LHC, where a high $p_T$ photon is required in the final state. The
Figure 1: Feynman diagrams for the process $b\bar{b} \rightarrow A/H \gamma$.

advantage to be gained from the photon is that at leading order the exclusive $\phi \gamma$ final state selects the $b\bar{b}$ initial state through the channel $b\bar{b} \rightarrow \phi \gamma$ shown in Figure 1. Indeed, the $gg \rightarrow \phi \gamma$ transition is forbidden at any order in perturbation theory by C-parity conservation. One can easily check, at one loop, that the box amplitude for $gg \rightarrow \phi \gamma$ exactly vanishes, even for off-shell external states. Hence, possible next-to-leading corrections coming from extra soft gluons emitted from the initial gluons vanish, too, while gluon emission from the box vanishes in the soft limit due to gauge invariance. Hence, the vanishing of the exclusive $\phi \gamma$ production from gluon fusion is stable under radiative corrections. Note that the process $gg \rightarrow \phi Z$ does not share the same property \cite{8}. Note also that when studying the $A/H \rightarrow \tau \tau$ signal with an additional $b$-tagging, there is in principle a large contamination arising from the $gb$ initial state. In addition, in this case, tagging a $b$-jet gives rise to double counting in the Monte Carlo simulation of events \cite{9}.

Of course, the presence of an additional photon in $b\bar{b} \rightarrow \phi \gamma$ suppresses the corresponding Higgs production rates, due to both the electromagnetic coupling $\alpha_{em}$ and the $b$-quark electric charge $Q_b = -1/3$. Nevertheless, we will see that the large-tan $\beta$ enhancement in the $A/H$ production can partly compensate this suppression, and leads to measurable rates.

In large-tan $\beta$ scenarios, the most promising signature for the process $pp \rightarrow \phi \gamma$ is the one arising from the decays $A/H \rightarrow \tau \tau$. Indeed, one has a branching ratio $BR(A/H \rightarrow \tau \tau) \simeq 10\%$, and, on the other hand, the irreducible background for the $\tau \tau \gamma$ final state has a purely electroweak origin, and will be shown to be well under control.

The $\tau \tau$ signature in Higgs boson production at the LHC has been extensively studied both in the SM Higgs boson production via vector boson fusion and in the inclusive MSSM A and H boson production \cite{10}. It turns out to be in general extremely promising, providing a discovery channel in the SM case for quite light Higgs-boson masses \cite{11}. Also, a few percent accuracy on the measurement of the $A/H$ masses in the MSSM is expected on the basis of the $\tau \tau$ invariant mass reconstruction \cite{12}.
Although the corresponding Higgs boson BR in $\tau\tau$ is in general one order of magnitude smaller than that of the dominant decay mode into $b\bar{b}$, the possibility of tagging $\tau$ leptons allows a drastic reduction of the background. On the other hand, despite the fact that undetected neutrinos are present among the decay products of the $\tau$’s, one can reconstruct the complete $\tau\tau$ invariant mass, whenever the two $\tau$ lepton are neither back-to-back nor collinear in the laboratory frame [13]. Requiring a large $p_T$ photon in the final state of the process $pp \to \phi \gamma$ naturally produces the required kinematical $\tau\tau$ configuration.

Note that the $A$ and $H$ bosons masses, as well as their couplings to fermions and decay widths, tend to be degenerate at large $\tan\beta$ and large $m_A$. As a consequence, the experimental separation of the two signals arising from $A$ and $H$ will be challenging. In the following analysis, we will then combine the statistical samples corresponding to the $A$ and $H$ resonances.

The analytical expression for the tree-level differential cross section $\sigma_\phi \equiv \sigma(b\bar{b} \to \phi \gamma)$, for $\phi = \{h, H, A\}$, is given by

$$
\frac{d\tilde{\sigma}_\phi}{d\hat{t}}(\hat{s}, \hat{t}) = \frac{\alpha_{em} Q^2_\phi \lambda^2_\phi}{4N_c(1-4r_b)} \left\{ \frac{F^\phi_1(\hat{s})}{(\hat{t}-m^2_b)(\hat{u}-m^2_b)} + F^\phi_2(\hat{s}) \left( \frac{1}{(\hat{t}-m^2_b)^2} + \frac{1}{(\hat{u}-m^2_b)^2} \right) \right\},
$$

(1)

where

$$
F^\phi_1(\hat{s}) = (1-r_\phi)^2 + 2(r_\phi - r_b \xi_\phi)(1-2r_b), \quad F^\phi_2(\hat{s}) = -2r_b(r_\phi - r_b \xi_\phi),
$$

(2)

$N_c$ is the color number, $r_i = m_i^2/\hat{s}$ ($i = \phi, b$), $\xi_{h,H} = 4$, and $\xi_A = 0$. The Yukawa couplings $\lambda_\phi$ are given by $\lambda_h = \lambda^{SM}(-\sin \alpha/\cos \beta)$, $\lambda_H = \lambda^{SM}(\cos \alpha/\cos \beta)$, and $\lambda_A = \lambda^{SM} \tan \beta$, being $\lambda^{SM} = m_b/v$, with $v$ the electroweak vacuum expectation value. The Mandelstam variables are defined as $\hat{t} = (p_b - p_\gamma)^2$, $\hat{u} = (p_b + p_\gamma)^2$, $\hat{s} = (p_b + p_\gamma)^2$. Finally, since (at tree-level) $\sin 2\alpha = -\sin 2\beta m^2_H + m_b^2)/(m^2_H - m_b^2)$, at large $\tan \beta$ the cross sections for the $H$ and $A$ boson production are much larger than the $h$ boson cross section.

In the $m_b \to 0$ limit, in Eq. (1) one has $F^\phi_1(\hat{s}) \to (1 + r^2_\phi)$, $F^\phi_2(\hat{s}) \to 0$, and the scalar and pseudoscalar boson cross sections have the same expression in terms of Yukawa couplings and masses, due to the chiral symmetry restoration.

We also take into account the leading-order QCD corrections to the Yukawa couplings, by using the running $b$-quark mass $m_b(\mu)$, in the Yukawa coupling, evaluated at a scale $\mu \sim m_\phi$. For $m^\text{pole}_b = 4.6$ GeV, one has $m_b(\mu) \simeq 2.81$ and 2.56 GeV, for $\mu \simeq 150$ and 500 GeV, respectively. These results have been obtained using the HDECAY program [14], also used to calculate the masses and widths of the Higgs bosons.
In order to obtain the inclusive cross sections for the process \( pp \rightarrow \phi \gamma \), we impose a minimal cut \( p_T^{\text{cut}} > 30 \text{ GeV} \) on the \( \gamma \) transverse momentum. The latter also regularizes soft and collinear singularities arising in the \( m_b \rightarrow 0 \) limit.

The \( pp \) cross section \( \sigma_{\phi \gamma} \), after integrating over \( p_T^2 > p_T^{\text{cut}} \), is then

\[
\sigma_{\phi \gamma} = \int_{x_1^c}^1 dx_1 \int_{x_1/x_2}^1 dx_2 \left[ f_b(x_1) f_b(x_2) + \text{c.c.} \right] \tilde{\sigma}_{\phi \gamma}(\hat{s}),
\]

where \( \hat{s} = x_1 x_2 S \), \( x_1^c \equiv \left( p_T^{\text{cut}} + \sqrt{(p_T^{\text{cut}})^2 + m_\phi^2} \right)^2 / S \), \( \sqrt{S} \) is the \( pp \) center-of-mass energy, and \( f_b(x) \) is the \( b \)-quark parton density in the proton. In Eq. 3, the inclusive partonic cross section \( \tilde{\sigma}_{\phi \gamma}(\hat{s}) \) can be obtained by integrating Eq. 2:

\[
\tilde{\sigma}_{\phi \gamma}(\hat{s}) = \frac{c_{\text{em}} Q_\phi^2 \lambda_\phi^2}{2 N_c \hat{s}} \left\{ \frac{1 + r_\phi^2}{1 - r_\phi^2} \log \left( \frac{1 + \sqrt{\Delta}}{1 - \sqrt{\Delta}} \right) - 2 \frac{m_b^2}{(p_T^{\text{cut}})^2} \sqrt{\Delta} r_\phi (1 - r_\phi) \right\},
\]

where \( \Delta = 1 - 4 \left( p_T^{\text{cut}} \right)^2 / \left[ \hat{s} (1 - r_\phi^2) \right] \). In Eq. 2, we retained the leading contributions in \( m_b \), that are of order \( \mathcal{O}[m_b^2 / (p_T^{\text{cut}})^2] \). They arise by integrating the terms \( m_b^2 / (t - m_b^2) \) and \( m_b^2 / (u - m_b^2) \) in the differential cross section in Eq. 1, and have an impact of the order of a few permil for \( p_T^{\text{cut}} \approx 30 \text{ GeV} \). Note that, in the inclusive cross section, \( m_b^2 / (t - m_b^2) \) and \( m_b^2 / (u - m_b^2) \) terms give also finite contributions that survive the \( m_b \rightarrow 0 \) limit after integration. This mass discontinuity is due to a chirality-flip mechanism induced by collinear photons [15].

The inclusive cross section for \( pp \rightarrow A \gamma \) versus \( m_A \) and \( \tan \beta \) is shown in Figure 2 for \( p_T^{\text{cut}} \approx 30 \text{ GeV} \). We use the parton density set CTEQ6L1 [16], with the factorization scale \( \mu_F \approx m_A / 2 \). Cross sections of a few hundreds fb’s are obtained for a relatively light \( m_A \) and large \( \tan \beta \) values. Assuming \( p_T^{\text{cut}} = 30, 40, 50, 60 \text{ GeV} \), for \( m_A = 150 \) (500) GeV and \( \tan \beta = 40 \), one gets \( \sigma = 194, 116, 73.6, 49.0 \) (6.12, 4.36, 3.26, 2.50) fb, respectively. Similar results hold for the heavy scalar Higgs boson process \( pp \rightarrow H \gamma \), for \( m_H \approx m_A \) and same \( \tan \beta \).

Note that the corresponding \( pp \rightarrow A \gamma \) cross section at Tevatron turns out to be 2.5 fb, for \( p_T^{\text{cut}} \gtrsim 20 \text{ GeV} \), if \( m_A = 150 \text{ GeV} \) and \( \tan \beta \approx 50 \).

We now present the results of a first analysis of the \( \tau \tau \) signature corresponding to the process \( pp \rightarrow \phi \gamma \rightarrow \tau^+ \tau^- \gamma \) in the MSSM at the LHC. It aims to separate the \( b\bar{b} \rightarrow \phi \gamma \) signal from the main irreducible background corresponding to the \( \tau \) pair production via an off-shell \( Z/\gamma \) associated to a high \( p_T \) photon. The spurious signal coming from the channels \( gg, b\bar{b} \rightarrow \phi \rightarrow \tau^+ \tau^- \rightarrow \tau \tau \gamma \), mediated by an off-shell \( \tau \) that radiates a photon, has also to be kept under control, since only the resonant \( \phi \rightarrow \tau \tau \) production associated to a high \( p_T \) photon can disentangle the initial \( b\bar{b} \) partonic state.
Figure 2: Total cross section (in pb) for the process $pp \to A \gamma$ at the LHC, for $p_T > 30$ GeV, and several values of $\tan \beta = 10, 20, 30, 40, 50$ (from lower to higher curve).

A much more detailed presentation than can be fit into this paper will be given in [17].

The $pp \to \tau \tau \gamma$ cross sections for signal and SM backgrounds have been calculated by using the full tree-level matrix element through CompHep [18], and also cross checked via our analytical results. Apart from the $\tau$ pair production via an off-shell $Z/\gamma$ associated to a high $p_T$ photon, a further irreducible SM backgrounds is given by the channel $pp \to W^+W^-\gamma$, when $W^\pm \to \tau^\pm\nu$. We checked that the latter contribution to the $\tau\tau\gamma$ signature is negligible, even including the other leptonic $W$ decays that can fake taus decays. Indeed, we find $\sigma(pp \to W^+W^-\gamma) \times Br(W \to l\nu)^2 \simeq 14$ fb, for $p_T > \sim 30$ GeV, which will be further reduced by cuts on the $\tau\tau$ invariant mass, etc.

Due to the mass and coupling degeneracy discussed above, the signal corresponding to $b\bar{b} \to A \gamma \to \tau \tau \gamma$ and $b\bar{b} \to H \gamma \to \tau \tau \gamma$ will be added coherently, increasing the total production rate by a factor two.

In counting the expected number of events for a given integrated luminosity, we assume quite conservatively a detection efficiency for the tau pair of $\epsilon_{\tau\tau} \simeq 0.2$. This takes into account both the leptonic $\tau \to \ell \nu, \nu \ell$ (35%) and the one-prong hadronic $\tau \to h\nu$ (50%) signatures for each of the two taus, assuming ID efficiency $\sim 90\%$ and $\sim 25\%$, respectively (giving a detection efficiency $\epsilon_{\tau} \simeq 0.44$ on a single tau) [19]. Note that the doubly hadronic decays contributes only by about 0.016 to $\epsilon_{\tau\tau} \simeq 0.2$ in this case.

In order to optimize the suppression of any $\tau\tau\gamma$ final state not coming from the resonant $b\bar{b} \to \phi \gamma \to \tau\tau\gamma$, we impose the following kinematical cuts:
Figure 3: Differential cross sections $d\sigma/dm_{\tau\tau}$ at the LHC for the process $b\bar{b} \rightarrow \tau\tau\gamma$, at $\tan \beta = 50$ and $m_A = 150$ GeV (left), 500 GeV (right), after sequential application of the cuts indicated (with $\Delta R \equiv \Delta R_{ij}$). Cuts on pseudo-rapidities and transverse momenta are detailed in the text.

- $p_T^\tau > 30, 30, 40, 50$ GeV, for $m_A = 150, 200, 300, 500$ GeV, respectively;
- $0.9 < m_{\tau\tau} < 1.1 m_A$ on the $\tau\tau$ invariant mass;
- $p_T^{\gamma \pm} > 20$ GeV, $|\eta_i| < 2.5$, $|\eta_{\gamma \pm}| < 2.5$, where $|\eta_i|$ is the $i$-particle pseudo-rapidity;
- $\Delta R_{\gamma \tau \pm} > 0.7$, $\Delta R_{\tau\tau} > 0.7$, $\Delta \phi_{\tau\tau} < 2.9$, where $\Delta R_{ij} = \sqrt{(\Delta \eta_{ij})^2 + (\Delta \phi_{ij})^2}$, and $\Delta \eta_{ij}$ and $\Delta \phi_{ij}$ stand for the difference in pseudo-rapidity and azimuthal angle of the particle $i$ and $j$, respectively.

The $\Delta R_{ij}$ cuts guarantee the needed spatial isolation to identify particles experimentally. On the other hand, both the $\Delta R_{\gamma \tau \pm} > 0.7$ and the $\Delta \phi_{\tau\tau} < 2.9$ cuts suppress the kinematical configurations where the photon is almost collinear to one of the $\tau$. The latter are typical of photons emitted off taus, like in $gg, b\bar{b} \rightarrow \phi \rightarrow \tau^* \tau \rightarrow \tau\tau\gamma$, that we want to deplete. Adding the cut $\Delta \phi_{\tau\tau} < 2.9$ after imposing $\Delta R_{\gamma \tau \pm} > 0.7$ has a modest impact at $m_A \sim 150$ GeV, but is quite effective at larger masses (see Figure 3). Imposing the $\tau\tau$ acollinearity condition $\Delta \phi_{\tau\tau} < 2.9$ also guarantees the reconstruction of the complete invariant mass for the $\tau\tau$ system despite the presence of invisible neutrinos in the $\tau$ decay products [20].

Figure 3 shows, for the $b$-fusion process $b\bar{b} \rightarrow \tau\tau\gamma$ mediated by $A$ and $H$, the effect of the cuts on $\Delta R_{\gamma \tau}, \Delta R_{\tau\tau}$, and $\Delta \phi_{\tau\tau}$ in dramatically reducing (especially at large $m_{A,H}$) the contribution of $\gamma$ emission off taus, that shifts $m_{\tau\tau}$ from the $m_{A,H}$ resonance. After applying all cuts, one gets a substantial purity of the $b\bar{b} \rightarrow \phi \gamma$ signal. In particular, the $\gamma$ radiation off b’s gives $83\%$ ($66\%$) of the observed rate, for $m_A = 150$ ($500$) GeV, at $\tan \beta = 50$. 

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### Table 1: Signal cross section $\sigma_S$ for the $b$-fusion contribution to the process $pp \rightarrow \tau^+\tau^-\gamma$ at the LHC, and corresponding significance $S = n(S)/\sqrt{n(S) + n(B)}$ versus $m_A$ (in GeV) and $\tan\beta$ ($n(S)$ and $n(B)$ stand for the number of events for signal and irreducible background, respectively). An integrated luminosity of $\mathcal{L} = 100$ fb$^{-1}$ and a tau-pair reconstruction efficiency $\epsilon_{\tau\tau} \simeq 0.2$ are assumed in the $S$ evaluation. Kinematical cuts are defined in the text.

| $\tan\beta$ | $m_A$ (GeV) | $\sigma_S$ (fb) | $S$ | $\sigma_S$ (fb) | $S$ | $\sigma_S$ (fb) | $S$ | $\sigma_S$ (fb) | $S$ |
|-------------|-------------|-----------------|-----|-----------------|-----|-----------------|-----|-----------------|-----|
| 20          | 150         | 5.58            | 7.3 | 12.5            | 13  | 22.1            | 19  | 34.5            | 24  |
|             | 200         | 3.00            | 5.3 | 6.81            | 9.5 | 12.3            | 14  | 19.9            | 18  |
|             | 300         | 0.727           | 2.4 | 1.67            | 4.5 | 3.08            | 6.7 | 5.03            | 9.1 |
|             | 500         | 0.0981          | 0.72| 0.238           | 1.5 | 0.456           | 2.4 | 0.768           | 3.4 |

The $b$-fusion cross sections $\sigma_S$ (mediated by both $A$ and $H$ bosons), contributing to the $pp \rightarrow \tau\tau\gamma$ rate after the above cuts, are reported in Table 1. For the same choice of cuts, the $Z^*/\gamma^*$ background cross sections are $\sigma_B = 6.10$, 3.44, 1.12, 0.270 fb, for $m_A = 150$, 200, 300, 500 GeV, respectively. Values for the significance $S = n(S)/\sqrt{n(S) + n(B)} > 5$ (where $n(S)$ and $n(B)$ stand for the number of events for signal and irreducible background, respectively) are obtained for $m_A \lesssim 300$ GeV and $\tan\beta \gtrsim 30$, with an integrated luminosity of $\mathcal{L} = 100$ fb$^{-1}$, and the tau-pair detection efficiency $\epsilon_{\tau\tau} \simeq 0.2$.

In order to get a feeling of the sensitivity of the signal rates to the present estimate of the error in the $b$-quark parton density, we computed the signal rates for several PDF sets among the ones present in LHAPDF [21], and found a variation of at most about 20%. However, one should keep in mind that the actual uncertainty on the PDF could be quite larger than the one given by the present LHAPDF set [22].

Another source of theoretical uncertainty in our calculation comes from the SUSY corrections to the b-Yukawa coupling that we are not taking into account, and which are proportional to $1/(1 + \Delta_b)^2$. The term $\Delta_b$ depends also on further SUSY parameters (see [23] for more details). However, although these corrections turn out to be in general quite large, they cancel out when the $\phi \rightarrow \tau\tau$ decay is considered, leaving only a residual dependence $\approx 1/[(1 + \Delta_b)^2 + 9]$, just as is the case for the corresponding process without a photon [23].

Clearly, how well one will be able to determine both the $b$-quark parton density and the overall coupling $(\tan\beta)^2$ will crucially depend on the actual experimental precision on the measurement of the $b\bar{b} \rightarrow \phi\gamma$ cross section. We believe that both a
more sharp theoretical prediction, i.e. the inclusion of higher order QCD corrections, and a full experimental simulation of the signatures corresponding to the signal and background, will be needed to assess the actual potential of this process.

Presently, studies have been carried out in order to estimate the uncertainty on the measurement of $\tan \beta$ derived from the $b\bar{b}\phi$ production (see [24]). This uncertainty is expected to be quite large at the moment. Assuming a comparable uncertainty for the value of the $\tan \beta$ coupling in the process $b\bar{b} \rightarrow \phi \gamma$ would prevent the extraction of the $b$-quark density in a precise way. However, one could think of new strategies aimed to optimize the measurements of the relevant quantities coming from different processes. For example, one could consider the ratio of the $b\bar{b} \rightarrow \phi \gamma$ cross section over the $b\bar{b}\phi$ production cross section where the dominant $(\tan \beta)^2$ dependence (and uncertainty) would automatically drop off. One could then obtain an observable with higher sensitivity to the $b$-quark parton density.

In conclusion, we believe that the process $b\bar{b} \rightarrow \phi \gamma$, possibly combined with other observables, could help in the determination of both the $b$-quark parton density and the $b$-quark Yukawa coupling. In order to establish its actual potential, further theoretical and experimental studies will be needed.

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