Motivated by recent results from the LHC experiments, we analyze Higgs couplings in two Higgs doublet models with an approximate PQ symmetry. Models of this kind can naturally accommodate sizable modifications to Higgs decay patterns while leaving production at hadron colliders untouched. Nearby the decoupling limit, we integrate out the heavy doublet to obtain the effective couplings of the SM-like Higgs and express these couplings in a physically transparent way, keeping all orders in \((m_h/m_H)\) for small PQ breaking. Considering supersymmetric models, we show that the effects on the Higgs couplings are considerably constrained.

**Introduction.** It is with a light heart that we assume, as a working hypothesis, that the recent measurements presented by CMS and ATLAS \([1,2]\) hint to an Higgs boson of mass \(m_h \approx 125\) GeV. Once this first step is taken, it is only human to collect data and thoroughly inspect them for hints of new physics. Presently, the data are consistent with \(O(1)\) enhancements with respect to a Standard Model (SM) Higgs boson, for both gluon fusion (GF) and vector boson fusion (VBF) production channels, in the \(\gamma\gamma\) \([3,4]\), and possibly also in the ZZ and WW decay modes \([5,6]\).

If the \(\gamma\gamma\) rate (and potentially also ZZ andWW rates) turn out to be larger than in the SM, in both GF and VBF, then the multiple enhancements are more easily interpreted in terms of non-standard Higgs decay, rather than production. The simplest explanation being an \(O(1)\) reduction in the \(hbb\) coupling. However, a suppression in \(hbb\) must not be accompanied by a similar suppression in \(ht\). Otherwise, without fortuitous interference between different new physics effects, GF production would be reduced by a similar amount. The situation is then such that (i) the couplings of \(h\) to down-type quarks and to up-type quarks exhibit non-universal sensitivity to new physics, and (ii) the effect in the Higgs-top coupling is more pronounced than in the Higgs-bottom coupling. These requirements are met in two Higgs-doublet models (2HDMs) with an approximate PQ symmetry. 

In this class of models, some hierarchy, with one doublet parametrically heavier than the other, is motivated by experimental constraints. First, electroweak precision measurements limit the contribution of new physics to the \(\rho\) parameter. This implies that new physics at the TeV scale should approximately conserve the diagonal \(SU(2)_c\) custodial symmetry. Since \((H^-,A,H^+)\) transforms as a \(3\) of \(SU(2)_c\), the splitting between \(m_{H^\pm}\) and \(m_A\) is constrained in these models \([7,8]\). Furthermore the field \(H \pm iA\) is charged under PQ, so the splitting between \(m_{H^\pm}\) and \(m_A\) is of the order of the PQ breaking, that we are assuming is moderate. Second, the mass of the charged Higgs boson \(H^\pm\) is constrained by its contribution to the decay \(B \to X\gamma\) and recent calculations give the bound \([9]\)

\[
m_{H^\pm} > 295 \text{ GeV}
\]

at 95\%CL. Without accidental cancellations, this bound also applies to models with a richer particle spectrum, such as supersymmetry. As a result, it is natural to expect the whole doublet \((H^+,H+iA)\) to be parametrically heavier than \(m_h \approx 125\) GeV.

A mass hierarchy motivates an effective theory analysis of the 2HDM with only a single SM-like Higgs boson at the weak scale. In this paper we take on this analysis, adopting a somewhat different approach than in the previous literature. Focusing our attention to modified Higgs couplings, we present our results in a way that we find more transparently related to the symmetries of the theory from which the 2HDM and, eventually, the SM is assumed to descend. We show how these symmetries are reflected in observable Higgs signals and demonstrate the utility of our approach by easily deriving the modified Higgs couplings in several supersymmetric embeddings of the 2HDM.

**2HDM analysis.** Consider a type-II 2HDM with \(H_{1,2} \sim (1,2)_{\pm1/2}\), where only \(H_1\) couples directly to \(Q_Ld_R\) and only \(H_2\) couples directly to \(Q_Lu_R\) at high energies. Neglecting leptons for now, the Lagrangian is \([10,11]\)

\[
- \mathcal{L} = H_1^\dagger D^2 H_1 + H_2^\dagger D^2 H_2 + m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + \frac{\lambda_1}{2} |H_1|^4 + \frac{\lambda_2}{2} |H_2|^4 + \lambda_3 |H_1|^2 |H_2|^2 + \lambda_4 |H_1| \sigma_2 |H_2|^2 + \left\{ \frac{\lambda_5}{2} (H_1^\dagger H_2)^2 + (H_1^\dagger H_2) (m_{12}^2 + \lambda_6 |H_1|^2 + \lambda_7 |H_2|^2) + \lambda_8 |H_2|^2 + \lambda_9 |H_1|^2 \right\}.
\]

The parameters \(m_{12}^2, \lambda_6, \lambda_7\) and \(\lambda_8\) violate a \(U(1)_{PQ}\) under which \((H_1^\dagger H_2)\) has charge +1. A discrete \(Z_2\) subgroup of this \(U(1)_{PQ}\) controls the mixing between the two doublets. In this paper, we loosely refer to approximate \(Z_2\) as the PQ limit. Since the coupling \(\lambda_5\) is even under the...
$Z_2$, it does not need to be small for our analysis to apply and indeed we will treat it collectively with other $Z_2$-even couplings. We parameterize spontaneous symmetry breaking (SSB) in a unitary gauge with

$$H_1 = \left( \begin{array}{c} h^+ \\ h_{1+} \end{array} \right), \quad H_2 = \left( \begin{array}{c} 0 \\ h_{2} \end{array} \right), \quad \langle h_2 \rangle = v_2, \quad (3)$$

where $a, h_1, h_2$ and the VEV $v_2$ are real.

It is possible to diagonalize the Higgs mass matrix and express the couplings in terms of the rotation angle $\alpha$ connecting the interaction to the mass basis and of the ratio $\tan \beta$ between the VEVs of $H_2$ and $H_1$. This procedure gives $r_d \equiv \frac{v_d a}{m_a} = -(\sin \alpha / \cos \beta)$, $r_u \equiv \frac{v_u a}{m_u} = (\cos \alpha / \sin \beta)$ and $r_y \equiv \frac{v_y (\alpha, \beta)}{m_y} = \sin(\beta - \alpha)$. The trigonometric expressions for the $r_X$'s are useful as they provide the exact result and make apparent simple algebraic relations between them. They are less useful, however, if one looks for more insight into the underlying theory. Here, much in the spirit of [13], we abandon the exact but somewhat less revealing $\alpha - \beta$ formulation in favor of a perturbative expansion, keeping track of the couplings in Eq. (2) as we work out the solution.

Our scheme is useful if the doublet $H_1$ is heavier than $H_2$, so that around the scale $m_h$ only $H_2$ is accessible. With this framework in mind we will obtain an effective action for $h_2$ to order $(B/M_1^2)^3$, where

$$M_1^2 = m_1^2 + \frac{\lambda_{35} h_2^2}{2}, \quad B = m_{\tau_2}^2 + \frac{\lambda_7 h_2^2}{2} \quad (4)$$

with$^1 \lambda_{35} = \lambda_3 + \lambda_5$. We will not need to assume that $\lambda_{35} v^2 \ll m_1^2$. This will improve the accuracy of our results for a mild hierarchy $m_h \lesssim m_H$.

Before proceeding to integrate out the heavy fields in $H_1$, we note some simplifying properties of the Lagrangian. First, we assume that CP is conserved to a good approximation, and take all the potential couplings to be real. Under this assumption, scalars and pseudo scalars do not mix and we need only consider diagrams involving the two neutral scalars $h_1$ and $h_2$. Second, as defined in Eq. (2), $\lambda_4$ projects neutral onto charged states and vice versa. It does not enter in tree diagrams with no charged external Higgs fields and we can ignore it in what follows. Third, working to $O(B^3/M_1^6)$, we can ignore $\lambda_6$ and $\lambda_1$ that affect the results beginning at $O(B^5/M_1^6)$ and $O(B^4/M_1^6)$, respectively.

Integrating out $h_1$ we obtain

$$- \mathcal{L}_{eff} = \frac{1}{2} y_{b} D_{\mu} h_{2} D^{\mu} h_{2} + \frac{1}{2} m_{h_{2}}^{2} h_{2}^{2} + \frac{\lambda_2}{8} \frac{h_{2}^{4}}{\sqrt{2}} + \frac{Y_{a} b_{b} h_{2} t}{\sqrt{2}}$$

$$- \frac{1}{2} B h_{2} D^{2} + \frac{1}{2} B h_{2} B h_{2} - \frac{Y_{b} b_{b} Q}{\sqrt{2}} \frac{1}{D^{2} + M_{1}^{2}} B h_{2}. \quad (5)$$

The interactions of the canonically normalized SM-like Higgs $h$ with the fermions and gauge bosons can be read off from (5), after accounting for wave function renormalization at $O(B^2/M_1^4)$. In particular, the bottom-Higgs Lagrangian is given by

$$Y_{b} \frac{b_{b}}{\sqrt{2}} \frac{1}{D^{2} + M_{1}^{2}} B \left( v_{2} + \left( 1 - \frac{f^2}{2} \right) h \right) \quad (6)$$

with

$$v_{2} = v \left( 1 - \frac{f^2}{2 v^2} \right), \quad v^2 = \frac{1}{\sqrt{2} 2 G_{F}} \cong (246 \text{ GeV})^{2}, \quad f = \left( \begin{array}{c} B h_{2} \end{array} \right), \quad f' = \frac{\partial f}{\partial v_{2}}. \quad (7)$$

Using (5) and (6) we obtain:

$$r_{b} = \frac{v_{b} h \bar{b}}{m_{b}} = \left( 1 - \frac{m_{b}^{2}}{M_{1}^{2}} \right) \left( 1 + \frac{\lambda_{7} v_{2}^{2}}{B} - \frac{\lambda_{35} v_{2}^{2}}{M_{1}^{2} - m_{h}^{2}} \right), \quad (8)$$

$$r_{t} = \frac{v_{b} h \bar{t}}{m_{t}} = 1 + \frac{B^{2}}{2 M_{1}^{2}} \left( 1 - \left( 1 + \frac{\lambda_{7} v_{2}^{2}}{B} - \frac{\lambda_{35} v_{2}^{2}}{M_{1}^{2}} \right)^{2} \right), \quad (9)$$

$$r_{v} = \frac{v_{b} v \bar{v}}{2 m_{v}^{2}} = 1 - \frac{B^{2}}{2 M_{1}^{2}} \left( \frac{v_{2}^{2}}{B} - \frac{\lambda_{35} v_{2}^{2}}{M_{1}^{2}} \right)^{2}. \quad (10)$$

The appearance of terms $(m_{h}^{2}/M_{1}^{2})$ in $r_{b}$ is due to the derivative operator in the effective vertex (6). The $\Box$ operator is replaced by $\Box \rightarrow -m_{h}^{2}$ when acting on an external Higgs particle and by $\Box \rightarrow 0$ when acting on the vacuum. Since $m_{h}$ corresponds to the physical mass, the $\Box$ operator automatically includes radiative corrections to $m_{h}$.

The deviation of $r_{b}$ from unity is parametrically small, beginning at $O(B^3/M_1^6)$. The deviation of $r_{t}$ scales similarly. In contrast, the deviation of $r_{b}$ does not scale with $(B/M_1^2)$. It can be parametrically $O(1)$ provided that either (i) $\lambda_{7} v_{2}^{2} \sim m_{t_{2}}^{2}$ or (ii) $\lambda_{35} v_{2}^{2} \sim m_{l_{2}}^{2}$.

The condition $\lambda_{7} v_{2}^{2} \sim m_{t_{2}}^{2}$ implies that the hard and soft breakings of the PQ are comparable at the scale of SSB. Note that it is perfectly possible to have $\lambda_{7} v_{2}^{2} \sim m_{t_{2}}^{2}$ and $m_{l_{2}}^{2} \ll m_{t}^{2}$. For instance, if the theory at some high scale has $m_{l_{2}}^{2} \sim 0$ but finite $\lambda_{7}$, we can expect $m_{l_{2}}^{2} \sim \lambda_{7} m_{l_{2}}^{2}/(4 \pi)^{2}$ at scales below $m_{l_{1}}$. In this case we can have an $O(1)$ correction to $r_{b}$ while the heavier doublet can be very heavy, $m_{H} \sim \text{TeV}$. This shows that $r_{b}$ is a sensitive probe for hard breaking of the PQ [13].

The second condition, $\lambda_{35} v_{2}^{2} \sim m_{l_{2}}^{2}$, implies that a sizable part of the mass of $H_{1}$ is driven by SSB. This is the relevant condition for models in which hard breaking of the PQ is absent or small (like e.g. the MSSM). In this case, a discernible deviation of $r_{b}$ from unity implies a light second doublet with $m_{H} \sim v$. The corrections $\sim (m_{h}^{2}/M_{1}^{2})$ coming from the derivative expansion can then be relevant; note that these terms correct $r_{b}$ with a definite positive sign. In the interesting case where soft

---

$^1$ Compared with the basis of [10], $(B/M_1^2) \sim 1/\tan \beta$ and our $\lambda_{35}$ equals their $\lambda_{35}$. 

---

2
PQ breaking is also small, $B \ll m_\phi^2$, we can expand $r_b$ in $(B/M_1^2)$. In that case, we can replace $M_1^2 \to m_H^2$, valid to $\mathcal{O}(B^2/M_1^4)$, in Eq. (8), obtaining

$$r_b \approx \left(1 - \frac{m_h^2}{m_H^2}\right)^{-1} \left[1 - \left(1 - \frac{m_h^2}{m_H^2}\right) \frac{\lambda_{35} v^2}{m_H^2}\right].$$

(9)

Eq. (9) is correct to all orders in $(m_h^2/m_H^2)$ and $(\lambda_{35} v^2/m_H^2)$ and to $\mathcal{O}(B^2/M_1^4)$.

Modifying the Higgs decays to bottom quarks affects the other search channels by changing the total width. A 125 GeV Higgs in the SM has $BR(h \to b\bar{b}) \approx 56\%$ so, for instance, the diphoton signal will be

$$\frac{\sigma \times BR(h \to \gamma\gamma)}{\sigma \times BR(h \to \gamma\gamma)_{SM}} \approx 1 + 0.56 (r_b^2 - 1).$$

(10)

The effect on the $ZZ, WW$ final states is similar. In Fig. 1 we plot the diphoton enhancement from Eqs. (9)-(10). The maximal enhancement is about a factor of two and is obtained for $m_H^2 = m_h^2 + \lambda_{35} v^2$. This means that, in the context of a 2HDM with an approximate PQ and for order one couplings $\lambda_{35} \sim 1$, taking the recent best fit ATLAS and CMS results [3, 4] at face value implies a light second doublet $m_H \sim 300$ GeV. Note that in Eq. (10) we neglected the charged Higgs loop contribution to the coupling $h\gamma\gamma$. In the appendix we show that this contribution is indeed negligible for the range of $m_H$ that is consistent with $b \to X_s\gamma$.

Finally, using Eqs. (5) we give a quick prescription for computing the correction to $r_b$ in models with a type-II 2HDM at low energies.

1. If the theory contains hard breaking of the PQ via $\lambda_7$, then significant deviation is possible even for $m_H \sim \text{TeV}$ in which case the leading effect is [10, 13]

$$r_b \approx 1 + \frac{2}{1 + 2m_{12}^2/(\lambda_7 v^2)}. \quad \quad \quad \quad (11)$$

2. If there is little or no hard breaking of the PQ, $\lambda_7 v^2 \ll m_{12}^2$, then a modified $r_b$ requires a light second doublet. When soft PQ breaking is also small, $B \ll m_h^2$, Eq. (9) resums all powers of $(m_h^2/m_H^2)$ and $(\lambda_{35} v^2/m_H^2)$.

So far we have neglected the Higgs coupling to leptons, but those can be added in a straightforward manner. If the doublet $H_1$ that couples to the down quarks couples also to the leptons, then $r_l = r_b$ and the change to the total width is amplified by a small factor $1 + (m_\tau/m_b^2)/3 \sim 1.1$.

**Supersymmetric examples.** We now examine supersymmetric extensions of the SM with a 2HDM effective theory near the weak scale and extract the modifications to Higgs observables.

In supersymmetry, holomorphy of the superpotential requires a second Higgs doublet in order to couple the Higgs sector to both up- and down-type quarks. Identifying $H_2 = i\sigma_2 H_1^\dagger$, $H_2 = H_2$, the tree level quartic couplings of the MSSM are

$$\lambda_1 = \lambda_2 = \frac{g^2 + g'^2}{4}, \quad \lambda_3 = -\frac{g^2 + g'^2}{4}, \quad \lambda_4 = \frac{g^2}{2}, \quad \lambda_5 = \lambda_6 = \lambda_7 = 0.$$  \quad (12)

The coupling $\lambda_{35} = \lambda_3 + \lambda_5 \approx -0.14$ is negative and so tends to increase $h\bar{b}b$. With $\lambda_{35}$ fixed and assuming $m_h = 125$ GeV, Eq. (9) tells us that, neglecting loop corrections to the bottom Yukawa, the value of $r_b$ depends only on $m_H$ with little sensitivity to the details of the supersymmetric spectrum. This result is in good agreement with FeynHiggs [14]. The MSSM prediction for $r_b$ in this case translates to the vertical dashed line in Fig. 1.

In some corners of the MSSM parameter space, our analysis ceases to give the leading result due to loop effects outside of the Higgs sector [10, 13, 17]. The bottom Yukawa is corrected by stop-higgsino and sbottom-gaugino loops. The first contribution scales as $(\mu A_t/m_{\text{soft}}^2)/(4\pi^2)$, and can dominate the coupling for $(\mu A_t/m_{\text{soft}}^2) \sim \text{few}$. The second contribution scales as $(Y_{tB}\mu m_{\tau}/m_{\text{soft}}^2)/(4\pi^2)$ and can dominate for $(Y_{tB}\mu m_{\tau}/m_{\text{soft}}^2) \gtrsim (4\pi^2)(B/m_H^2) \sim (4\pi^2)/\tan \beta$. In addition, light sfermions can affect the $h\gamma\gamma$ (and potentially $hGG$) vertex. In particular, very light staus could give an enhancement if $(M_2^2/B)$ (or $\tan \beta$) is so large that the tau Yukawa becomes $\mathcal{O}(1)$ [15]. This defines a lower limit for $(B/M_1^2)$ where we can neglect bottom and tau loop corrections: $(B/M_1^2) \gtrsim (\sqrt{2}m_e/v) \sim 1/40$. 

**FIG. 1:** Contours of $\sigma \times BR(h \to \gamma\gamma)/(\text{SM})$ vs. $\lambda_{35}$ and $m_H$, for $m_h = 125$ GeV and $\lambda_\tau = 0$. (Recall $\lambda_{35} = \lambda_3 + \lambda_5$.) The MSSM prediction, neglecting loop corrections to the bottom Yukawa (but effectively including corrections to the Higgs potential), is shown by the dashed line.
It is interesting to ask whether simple extensions of the MSSM that accommodate a large Higgs mass in a more natural way can also reduce $hbb$. We briefly examine three such models, the NMSSM, the BMSSM and a $U(1)_X$. Considering a $Z_3$ version of the NMSSM, we write the superpotential
\[
W = \lambda S H_u H_d + \frac{\kappa}{3} S^3.
\]  
(13)

We assume that $S$ is somewhat heavy so that at low energy the theory can be described by the 2HDM. The coupling $\lambda_3$ is then given by
\[
\lambda_3 = -\frac{g^2 + g'^2}{4} + |\lambda|^2 \approx -0.14 + 0.5 \left| \frac{\lambda}{0.7} \right|^2
\]  
(14)

and is larger than in the MSSM. Still, as can be seen from Fig. 4, for this effect alone to achieve $r_b < 1$, $\lambda > 0.7$ is required, above the limit of perturbative unification $[18]$. Adding to Eq. (13) the supersymmetric mass terms $\mu H_u H_d$ and $\frac{\lambda}{4} S^2$ with $\Lambda \gg \mu$ produces the BMSSM $[7]$. The spurion $\epsilon_1 = (\lambda^2 \mu^4 / \Lambda)$ carries $PQ$ charge -1 and induces $\lambda_3 = \lambda_3^{\text{MSSM}} - 2 \epsilon_1$. For $(B/M_t^2) \lesssim 1/10$, $\epsilon_1$ could decrease $r_b$, while making a negligible correction to $m_h$ $[13]$.

Next, consider a gauge $U(1)_X$ extension under which the Higgs fields are charged, $q_{H_u} = q_{H_d} \equiv q_H$. The scalar potential receives a correction
\[
V = V_{\text{MSSM}} + \frac{g_X^2 q_h^2}{2} \left( |H_2|^2 + |H_1|^2 \right)^2
\]  
(15)

Leading to $\lambda_3 \rightarrow \lambda_3 + g_X^2 q_h^2$. The modification to the $hbb$ coupling in this example is in fact limited by the Higgs mass: since $\delta m_h^2 = \frac{\epsilon_1}{2} \approx 0.12$. We therefore do not expect large deviations from the MSSM tree level predictions.

Conclusions. We have shown that a 2HDM close to the decoupling limit with an approximate $PQ$ symmetry (or a $Z_2$ subgroup of it) can produce $O(1)$ deviations in the lightest Higgs couplings to SM particles. Integrating out the heavier Higgs, we presented the effective couplings of the light SM-like Higgs in a physically transparent way. Keeping derivative operators in the expansion allowed us to include automatically radiative corrections to the lightest Higgs mass, important in the limit of small $PQ$ breaking but a mild hierarchy $m_H \sim m_h$.

Our results are applicable to any type-II 2HDM not far from decoupling. In particular, considering the MSSM and some of its extensions, our analysis elucidates why it is hard to enhance $h \rightarrow \gamma\gamma$. If the current experimental hints are confirmed as the statistics increases, these results will allow to set bounds on the heavy Higgs mass $m_H$ in large regions of these models parameter space.

Assuming a 2HDM with an approximate $PQ$ and natural couplings, taking the recent best fit ATLAS and CMS results $[3, 4]$ at face value implies a not too-heavy second doublet $m_H \sim 300$ GeV. Further support for this possibility should come from better measurements of the SM-like Higgs decay patterns, where the generic type-II 2HDM predicts similar enhancements in gluon fusion and vector boson fusion production for $\gamma\gamma$, $ZZ$ and $WW$ final states, with correspondingly decreased $h \rightarrow bb$.

Acknowledgments. We thank Janie Chan, David E. Kaplan, Rouven Essig, JiJi Fan and Neal Weiner for useful discussions. We are grateful to Nima Arkani-Hamed for many insightful discussions and for comments on the manuscript. KB is supported by DOE grant DE-FG02-90ER40542.

Appendix: Charged Higgs contribution to $h\gamma\gamma$. The decay $h \rightarrow \gamma\gamma$ is mediated by a dimension five Lagrangian, that we parametrize by
\[
\mathcal{L}_\gamma = -\frac{2\pi \alpha v c}{\Lambda^2} h F_{\mu\nu} F^{\mu\nu} - \frac{2\pi \alpha v c}{\Lambda^2} h F_{\mu\nu} F^{\mu\nu}.
\]  
(16)

In the absence of CP-violation, $\bar{c}_\gamma = 0$. The contribution of the charged Higgs loop is given by
\[
\frac{c_\gamma}{\Lambda^2} = -\frac{\lambda H^+ H^+}{v} \frac{f(\tau)}{24 m^2 h^+}, \quad \tau = \frac{m_h^2}{4 m_{h^+}^2}
\]  
(17)

with
\[
f(r) = 3 \arcsin^2 \left( \sqrt{r} \right) - r \approx 1 + 0.6 r,
\]  
(18)

where the last approximation is valid to one percent accuracy for $m_h^2/4m_{h^+}^2 < 0.2$ or, in case of $m_h = 125$ GeV, for $m_{h^+} > 140$ GeV. The coupling $(\lambda H^+ H^+/v)$ can be obtained from Eq. (2),
\[
\frac{\lambda H^+ H^+}{v} = \lambda_{34}
\]  
(19)

with $\lambda_{34} = \lambda_3 + \lambda_2$ to $O(B^2/M_t^4)$. In the MSSM, for example, $\lambda_{34} = (g^2 - g'^2)/4 \approx 0.07$.

Adding this correction to the SM W and top loop contributions gives
\[
\Gamma(h \rightarrow \gamma\gamma) = \Gamma(h \rightarrow \gamma\gamma)_{\text{SM}} + \frac{\Gamma^W + \Gamma^t (1 - \frac{3}{4} \Gamma^t)}{2} - \frac{4 \pi^2 r^2}{\Lambda^2},
\]  
(20)

with the loop function $\Gamma^W = \Gamma^W + \frac{3}{4} (1 - \frac{3}{4}) I^t$ taken from $[19]$. For $m_h = 125$ GeV we find $\Gamma^W = -2.09, \Gamma_t = 0.34, \Gamma^t = -1.645$. Unless the charged Higgs is very light, or the relevant quartic couplings are large, the contribution to the $h\gamma\gamma$ coupling makes a negligibly small correction to the SM terms:
\[
\Gamma(h \rightarrow \gamma\gamma)_{\text{SM}} \approx 1.27 r_V - 0.27 r_t - 0.05 \left( \frac{m_{h^+}^2}{m_{h^+}} \right) \left( \frac{m_{h^+}}{350 \text{ GeV}} \right)^2.
\]  

*Electronic address: kblum@ias.edu
1 Electronic address: raffaele.dagnolo@sns.it

[1] ATLAS Collaboration, arXiv:1202.1408 [hep-ex].
[2] CMS Collaboration, arXiv:1202.1488 [hep-ex].
[3] ATLAS Collaboration, arXiv:1202.1414 [hep-ex].
[4] CMS Collaboration, arXiv:1202.1487 [hep-ex].
[5] ATLAS Collaboration, arXiv:1202.1415 [hep-ex]. G. Aad et al. [ATLAS Collaboration], arXiv:1112.2577 [hep-ex].
[6] CMS Collaboration, arXiv:1202.1489 [hep-ex]. CMS Collaboration, arXiv:1202.1416 [hep-ex].
[7] M. Dine, N. Seiberg and S. Thomas, Phys. Rev. D 76, 095004 (2007) [arXiv:0707.0005 [hep-ph]].
[8] J. Mrazek, A. Pomarol, R. Rattazzi, M. Redi, J. Serra and A. Wulzer, Nucl. Phys. B 853, 1 (2011) [arXiv:1105.5403 [hep-ph]].
[9] M. Misiak, H. M. Asatrian, K. Bieri, M. Czakon, A. Czarnecki, T. Ewerth, A. Ferroglia and P. Gambino et al., Phys. Rev. Lett. 98, 022002 (2007) [hep-ph/0609232].
[10] J. F. Gunion and H. E. Haber, Phys. Rev. D 67 (2003) 075019 [hep-ph/0207010].
[11] G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, M. Sher and J. P. Silva, arXiv:1106.0034 [hep-ph].
[12] P. M. Ferreira, R. Santos, M. Sher and J. P. Silva, arXiv:1112.3277 [hep-ph].

[13] L. Randall, JHEP 0802, 084 (2008) [arXiv:0711.4360 [hep-ph]].
[14] S. Heinemeyer, W. Hollik and G. Weiglein, Comput. Phys. Commun. 124, 76 (2000) [hep-ph/9812320];
S. Heinemeyer, W. Hollik and G. Weiglein, Eur. Phys. J. C 9, 343 (1999) [hep-ph/9812472]; G. Degrassi, S. Heinemeyer, W. Hollik, P. Slavich and G. Weiglein, Eur. Phys. J. C 28, 133 (2003) [hep-ph/0212020]; M. Frank, T. Hahn, S. Heinemeyer, W. Hollik, H. Rzehak and G. Weiglein, JHEP 0702, 047 (2007) [hep-ph/0611326].
[15] M. Carena, S. Gori, N. R. Shah and C. E. M. Wagner, arXiv:1112.3336 [hep-ph].
[16] G. L. Kane, G. D. Kribs, S. P. Martin and J. D. Wells, Phys. Rev. D 53, 213 (1996) [hep-ph/9508265].
[17] M. S. Carena, J. R. Ellis, A. Pilaftsis and C. E. M. Wagner, Nucl. Phys. B 586, 92 (2000) [hep-ph/0003180].
[18] L. J. Hall, D. Pinner and J. T. Ruderman, arXiv:1112.2703 [hep-ph].
[19] A. V. Manohar and M. B. Wise, Phys. Lett. B 636, 107 (2006) [hep-ph/0601212].
[20] A. Djouadi, Phys. Rept. 457, 1 (2008) [hep-ph/0503172];
A. Djouadi, Phys. Rept. 459, 1 (2008) [hep-ph/0503173].