Higgs boson interactions in supersymmetric theories with large tan β

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Abstract

We show that radiative corrections to the Higgs potential in supersymmetric theories with large tan β generically lead to large differences in the light Higgs boson decay branching fractions compared to those of the standard model Higgs boson. In contrast, the light Higgs boson production rates are largely unaffected. We identify Wh associated production followed by Higgs boson decays to photons or to leptons via WW^* as potential experimental probes of these theories.

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It is well known that supersymmetry requires two Higgs doublets to give masses to the up-type quarks and to the down-type quarks. Hence, we use the terminology "up-Higgs" and "down-Higgs" to indicate these two Higgs bosons, H_u and H_d respectively. The ratio of the vacuum expectation values,

$$\tan \beta \equiv \frac{\langle H_u^0 \rangle}{\langle H_d^0 \rangle}, \tag{1}$$

has both theoretical and experimental consequences.

Theoretically, a large value of $\tan \beta$ near m_t/m_b appears to be required in simple $b - \tau - t$ Yukawa unification theories such as supersymmetric $SO(10)$ [1]. Since the b-Yukawa and the t-Yukawa couplings are equal at the high scale they will be nearly equal at the low scales. This is because the bottom and top quarks are distinguished only by their weak hypercharge, and renormalization group evolution is dominated by large Yukawa and QCD interactions.

Furthermore, a large value of $\tan \beta \gtrsim 30$ is required for minimal gauge mediated models which solve the soft-CP problem [2, 3]. That is, CP violation effects are non-existent in the soft mass terms of a softly broken supersymmetric theory with gauge mediated supersymmetry breaking and with $B_\mu(M) = 0$ at the messenger scale M. If B_μ is zero at the messenger scale then arbitrary phases in the lagrangian can be rotated away, and CP violating effects do not get induced by renormalization group running. However, renormalization group running does induce a non-zero (CP conserving) B_μ term at the weak scale which is phenomenologically viable only if $\tan \beta$ is large [3, 4]. This can be seen readily from one of the conditions for proper electroweak symmetry breaking:

$$\sin \beta \cos \beta = \frac{B_\mu(M) + \Delta B_\mu}{m_A^2}. \tag{2}$$

Since $B_\mu(M) = 0$ and ΔB_μ is a small renormalization group induced mass, and m_A must be large enough to escape detection, the right-hand side of the equation is much less than 1. Therefore,

$$\sin \beta \cos \beta \ll 1 \implies \tan \beta \gg 1. \tag{3}$$

Unification of the third family Yukawa couplings and CP conservation in gauge mediated models both imply a strong preference for large $\tan \beta$. It is important to keep in mind that these preferences survive even when there is non-universality among the scalar masses in the theory. In supergravity, non-universality is even expected. Furthermore, large $\tan \beta$ is of course accessible in parameter space even if third family Yukawa unification is not realized,

or if the CP violating interactions are kept under control in a supersymmetric theory by some other means.

The mass matrix of the scalar Higgs bosons in the $\{H_d^0, H_u^0\}$ basis can be parametrized by,

$$M^2 = \begin{pmatrix} m_A^2 \sin^2\beta + m_Z^2 \cos^2\beta & -\sin\beta\cos\beta(m_A^2 + m_Z^2) \\ -\sin\beta\cos\beta(m_A^2 + m_Z^2) & m_A^2 \cos^2\beta + m_Z^2 \sin^2\beta \end{pmatrix} + \begin{pmatrix} A_{11} & A_{12} \\ A_{12} & A_{22} \end{pmatrix}, \quad (4)$$

where m_A is the physical pseudoscalar Higgs boson mass, and A_{ij} are radiative corrections induced by particle and sparticle loops. The dependences of A_{ij} on the sparticle masses and mixing angles can be found in several references [5]. For larger $\tan\beta$ and arbitrary mixing, a random scan over supersymmetric parameter space is enough to convince oneself that the off diagonal A_{12} can be as large as $\pm(40\,\text{GeV})^2$.

Models with large $\tan\beta$ and large m_A allow suppression of the $b\bar{b}$ decay mode resulting from cancellation of tree-level and off-diagonal terms in the Higgs mass matrix. There are two expansions that we would like to perform to illustrate this feature. The first is an expansion of the Higgs mass matrix about pure H_u^0 and H_d^0 *mass* eigenvalues. For these weak eigenstates to be mass eigenstates, mixing in the Higgs mass matrix is not allowed. At tree level this is impossible since in our convention M_{12}^2 is negative definite. However, when radiative corrections are introduced a cancellation can occur between the tree-level off-diagonal piece and the quantum corrections. For this to occur,

$$-\sin\beta\cos\beta(m_Z^2 + m_A^2) + A_{12} = 0 \quad (5)$$

is required. Note that if $\tan\beta$ is very large then the tree-level contribution gets small and the radiative correction term, A_{12}, can compete with it to yield $M_{12} = 0$ [6]. Solving for the mass eigenvalues we obtain,

$$m_A^2 = -m_Z^2 + \frac{A_{12}}{\sin\beta\cos\beta} \quad (6)$$

$$m_{H_d^0}^2 = m_Z^2 \cos 2\beta + A_{12} \tan\beta + A_{11} \quad (7)$$

$$m_{H_u^0}^2 = -m_Z^2 \cos 2\beta + \frac{A_{12}}{\tan\beta} + A_{22}. \quad (8)$$

To demonstrate the behavior of these eigenvalues, we plot their values in the limit that all soft squark masses are the same (m_{susy}) at the low scale. Squark mass degeneracy is broken
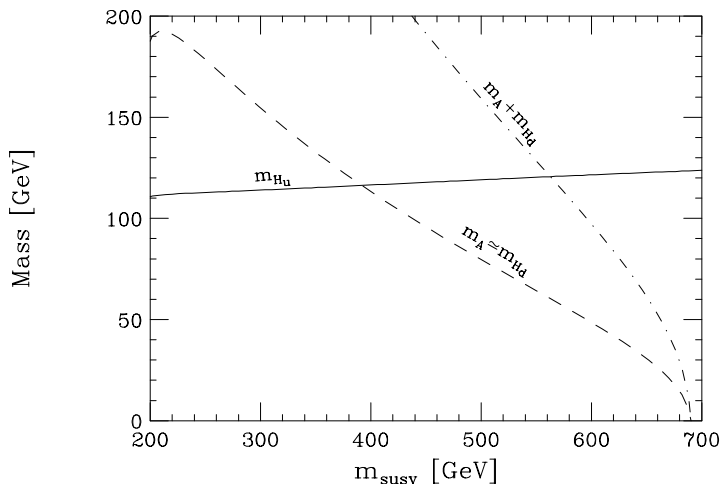
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Figure 1: The mass eigenvalues $m_A \simeq m_{H_d^0}$, $m_{H_u^0}$, and $m_A + m_{H_d^0}$ as a function of m_{susy}. $m_A + m_{H_d^0}$ is constrained to be greater than $\sim 142\,\text{GeV}$ from LEPII analyses.

by mixing introduced through a common A-term at the low scale and the μ term. Taking $\tan\beta = 40$ and $\mu = A = 300\,\text{GeV}$ we plot $m_A \simeq m_{H_d^0}$, $m_{H_u^0}$ and $m_A + m_{H_d^0}$ as a function of m_{susy}. We have plotted $m_A + m_{H_d^0}$ since it is constrained to be greater than $\sim 142\,\text{GeV}$ from LEPII analyses [7]. We see that H_d^0 and A have mass,

$$m_A^2 \simeq m_{H_d^0}^2 \simeq -m_Z^2 + A_{12}\tan\beta \tag{9}$$

which ranges widely with m_{susy} in this model. For A_{12} not too small m_A is naturally large for large $\tan\beta$. H_u^0 has a mass equal to,

$$m_h^2 = m_{H_u^0}^2 \simeq m_Z^2 + A_{22}, \tag{10}$$

which is logarithmically sensitive to m_{susy}.

The H_u^0 mass eigenvalue is always light (between m_Z and $135\,\text{GeV}$). By assumption it cannot decay into $b\bar{b}$ or $\tau^+\tau^-$, and is only allowed to decay into $c\bar{c}$, gg, $\gamma\gamma$, WW^*, $Z^{(*)}\gamma$, and $ZZ^{(*)}$. Since H_u^0 carries all the burden of electroweak symmetry breaking, it couples with full electroweak strength to the vector bosons, and the up-type quarks. Therefore, H_u^0 acts like the standard model Higgs in this limit in every way except it does not couple to down quarks or leptons[8]. The decay branching fractions [9] for this state are shown in Fig. 2. According to Eq. 10, the shaded region is theoretically not allowed in the high $\tan\beta$ limit.
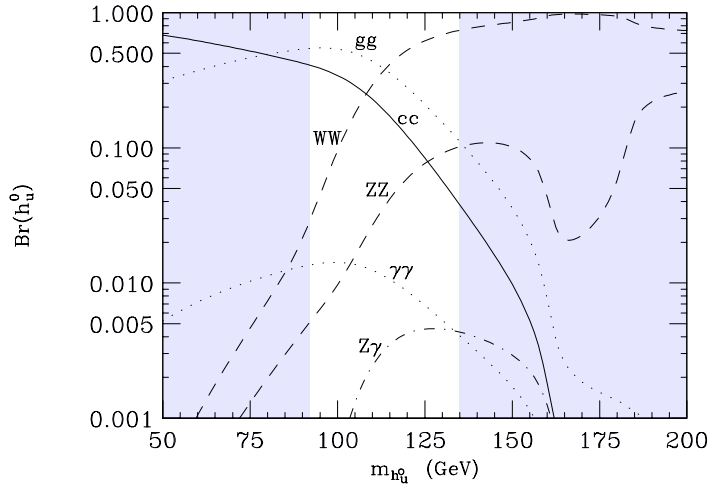
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Figure 2: Decay branching fractions of the H_u^0 mass eigenstate. The shaded region is theoretically not allowed with high $\tan\beta$.

Only in the region $m_Z \lesssim m_{H_u^0} \lesssim 135\,\mathrm{GeV}$ does one expect to find the H_u^0 at a high energy collider.

The above discussion may appear unsatisfactory since it seems to imply that a special finetuning of m_A is needed for interesting effects to occur. To address this concern, we expand in a second limit, the large m_A limit, in order to justify our claim that large deviations occur generically in large $\tan\beta$ and large m_A.

For a moment, let us ignore the radiative corrections and expand about large m_A with $A_{ij} = 0$. The lightest Higgs boson coupling to quarks and vector bosons compared to the standard model Higgs boson coupling to quarks and vector bosons is,

$$\frac{(h\bar{d}d)_{susy}}{(h\bar{d}d)_{sm}} = 1 + 2\epsilon_A + 2\epsilon_A^2 + \cdots \tag{11}$$

$$\frac{(h\bar{u}u)_{susy}}{(h\bar{u}u)_{sm}} = 1 - 2\epsilon_A\epsilon_\beta^2 - 4\epsilon_A^2\epsilon_\beta^2 + \cdots \tag{12}$$

$$\frac{(hVV)_{susy}}{(hVV)_{sm}} = 1 - 2\epsilon_A^2\epsilon_\beta^2 + \cdots \tag{13}$$

where

$$\epsilon_A \equiv m_Z^2/m_A^2 \quad \text{and} \quad \epsilon_\beta \equiv 1/\tan\beta. \tag{14}$$

There are several points to notice in the above expressions. First, the coupling to up-type fermions decouples to the standard model value much faster than the coupling to the

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down-type fermions by a factor of $1/\tan^2\beta$. Second, the vector boson couplings to the light Higgs boson decouples even more rapidly to the standard model result. Therefore, even at tree-level, we expect the coupling to down-type fermions to deviate from the standard model much more than the couplings to the up-type fermions and vector bosons.

Now, if we add radiative corrections to the expansion we get an even more interesting result. Keeping the first tree-level correction and the leading radiative correction terms,

$$\frac{(h\bar{d}d)_{susy}}{(h\bar{d}d)_{sm}} \;=\; 1 + 2\epsilon_A - \frac{A_{12}}{m_Z^2}\frac{\epsilon_A}{\epsilon_\beta} + \cdots \tag{15}$$

$$\frac{(h\bar{u}u)_{susy}}{(h\bar{u}u)_{sm}} \;=\; 1 - 2\epsilon_A\epsilon_\beta^2 - \frac{1}{2}\frac{A_{12}^2}{m_Z^4}\epsilon_A^2 + \frac{A_{12}}{m_Z^2}\epsilon_A\epsilon_\beta + \cdots \tag{16}$$

$$\frac{(hVV)_{susy}}{(hVV)_{sm}} \;=\; 1 - 2\epsilon_A^2\epsilon_\beta^2 - \frac{1}{2}\frac{A_{12}^2}{m_Z^4}\epsilon_A^2 + \cdots . \tag{17}$$

In this case the up-type fermions and vector boson couplings to the light Higgs boson still decouples rapidly to the standard model result. In large $\tan\beta$ models, $A_{12} \sim m_Z^2$ for much of the parameter space. Further, unless m_A is very heavy, we expect ϵ_A to be less than or comparable to ϵ_β for large $\tan\beta$. In such models there is an $\mathcal{O}(1)$ correction to the down-type fermion couplings to the light Higgs boson. In particular, we see that if $A_{12} \simeq m_Z^2\epsilon_A/\epsilon_\beta$ then $h = H_u^0$ and the coupling to down-type fermions may be shut off completely. This is the same result we concluded from the previous expansion above. Here, however, we have shown that the coupling to down-type fermions is generically altered by $\mathcal{O}(1)$ corrections in either direction depending on the sign of A_{12}, but the coupling to up-type fermions and vector bosons are standard model-like.

At high $\tan\beta$, the pseudoscalar state A^0 may be produced copiously at high energy colliders, and significant constraints obtained [10]. Here we focus on the CP-even scalar states, since one, which we are calling h, is guaranteed to be light (below $\lesssim 135\,\mathrm{GeV}$). The main production mechanisms for h at high energy hadron colliders are $qq' \to Wh$, $gg \to h$, and $qq \to t\bar{t}h$. At electron-positron colliders the production modes are $e^+e^- \to hZ$, $e^+e^- \to t\bar{t}h$, and $\gamma\gamma \to h$. Since the h couplings to the top-quark and vector bosons are very close to the standard model values, these cross-sections are not expected to deviate much from the Standard Model. This is good news to not have to worry about complicated mixing between non-standard production and decay. Only the branching fractions are altered in high

$\tan\beta$ models[*].

For a Higgs boson mass near $120\,\mathrm{GeV}$ and below, the main decay modes are $h \to b\bar{b}$ and $h \to \tau\tau$. If these two were the only decay modes then deviations in the down-type fermion couplings would just cancel in the branching fraction and only the standard model result would be seen. However, there are other important decay modes of the Higgs boson, most notably the $h \to \gamma\gamma$ decay mode. In the standard model its branching fraction is $\sim 1.4 \times 10^{-3}$ in the mass range $100\,\mathrm{GeV} < m_h < 150\,\mathrm{GeV}$. In supersymmetry, this branching fraction can be altered by three factors. First, supersymmetry sparticle loops involving squarks and charged Higgs bosons can change the $h \to \gamma\gamma$ amplitude. Given current limits on sparticle masses and some fortuitously large loop integrals for the W boson compared to charged scalars, this effect is small over almost all of the parameter space [11]. Second, changes in the $h\bar{t}t$ and hWW couplings can change the $h \to \gamma\gamma$ amplitude. However, we argued above that this does not occur in large $\tan\beta$ models. And third, large corrections to $h\bar{b}b$ and $h\tau\tau$ couplings can change the denominator in $B(h \to \gamma\gamma)$. This effect is large in the high $\tan\beta$ models.

Since , $(h \to WW^*)$ is unaffected in the large $\tan\beta$ region also, its branching ratio is determined only by the changes wrought by deviations in the down-type fermion couplings. In Fig. 3 we plot the deviation in the $h \to VV$ ($V = \gamma, W^{(*)}, Z^{(*)}$) branching fraction compared to the standard model branching fraction as a function of $(hbb)_{susy}/(hbb)_{sm}$. We also plot the deviation of $B(h \to b\bar{b})$ for reference. The plot was made with $m_h = 120\,\mathrm{GeV}$. Since the Higgs boson decays predominantly into down-type fermions at this mass, the branching fraction into vector bosons is most sensitive to deviations in the hbb coupling. The sensitivity (insensitivity) increases even more for $B(h \to VV)$ $(B(h \to b\bar{b}))$ for lighter Higgs masses.

The high $\tan\beta$ gauge mediated models with $B_\mu = 0$ at the messenger scale provide a good specific illustration of the expected deviations. It is conservative since the value of μ and m_A tends to be higher in these high $\tan\beta$ models than supergravity mediated theories. If we take the messenger scale to be equal to $10^4\Lambda$, where $\Lambda \equiv F_S/S$ (see [3] for discussions of gauge mediation parameters), then we can self-consistently calculate the value of $\tan\beta$ as a function of Λ which ensures that $B_\mu(M) = 0$. We find that with chargino mass between

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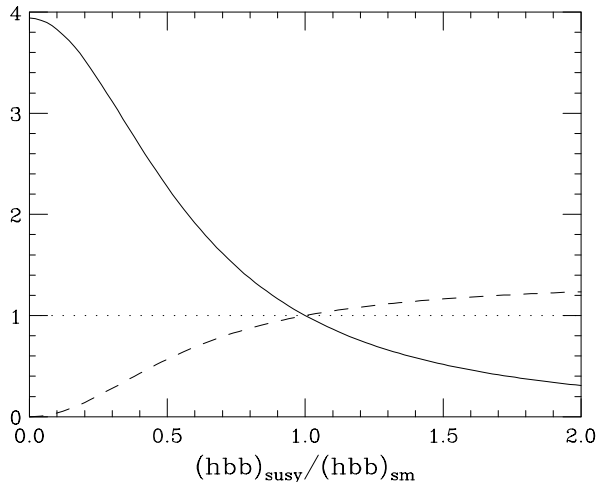[*]Note that $\mu^+\mu^- \to h$ would be altered at a muon collider.

Figure 3: $B(h \to VV)_{susy}/B(h \to VV)_{sm}$ (solid) and $B(h \to b\bar{b})_{susy}/B(h \to b\bar{b})_{sm}$ (dash) as a function of the variation in the $h\bar{b}b$ coupling for $m_h = 120\,\mathrm{GeV}$. A dotted line at the value 1 corresponds to the standard model reference.

$100\,\mathrm{GeV}$ and $200\,\mathrm{GeV}$,

$$105\,\mathrm{GeV} \lesssim m_h \lesssim 115\,\mathrm{GeV} \tag{18}$$

$$1.06 \lesssim \frac{B(h \to b\bar{b})_{susy}}{B(h \to b\bar{b})_{sm}} \lesssim 1.03 \tag{19}$$

$$0.55 \lesssim \frac{B(h \to VV)_{susy}}{B(h \to VV)_{sm}} \lesssim 0.85. \tag{20}$$

As expected the deviation in the Higgs to vector bosons branching fraction is significant. We have also scanned parameter space in supergravity theories with non-universal scalar masses and found $\pm\mathcal{O}(1)$ deviations in the $h \to VV$ branching fractions over much of parameter space.

An important question is whether or not experiment can see the effects of these deviations in the Higgs branching fractions to vector bosons. It is probably hopeless to try to see even 50% effects in the $\gamma\gamma$ rate from $gg \to h$ fusion. The QCD corrections to this rate are nearly a factor of two by themselves with large uncertainties [13]. A more fruitful approach may be to search for Wh associated production with the W decaying leptonically and the Higgs boson decaying to $\gamma\gamma$ [14] or $WW^* \to l\nu l\nu$ [12, 6]. The total rates for these production mechanisms in the standard model at the LHC with $\sqrt{s} = 14\,\mathrm{TeV}$ is given in Fig. 4. The advantages to these modes is the smaller background and the ability to cancel some systematics by taking ratios of this with some other electroweak process, for example, $\sigma(WZ)$.
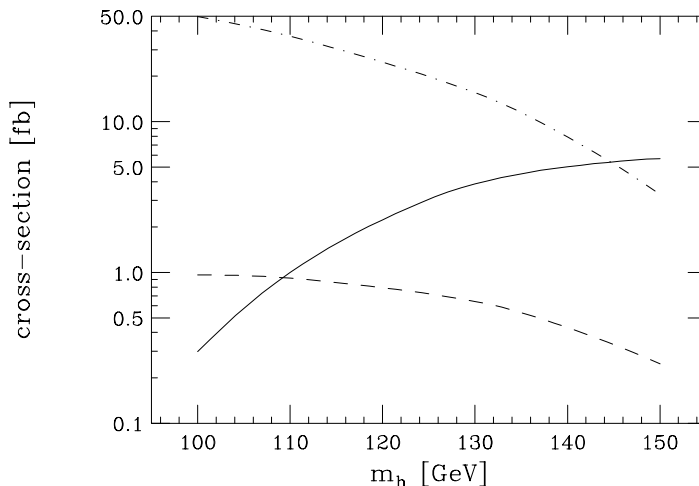
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Figure 4: $\sigma(Wh \to 3l)$ (solid), $\sigma(Wh \to l\gamma\gamma)$ (dash), and $\sigma(Wh \to l\tau\tau)$ (dot-dash) in the standard model at the LHC with 14 TeV center-of-mass energy.

The maximum possible sensitivity to a deviation in the cross-section can be estimated by assuming a perfect detector with no background and no theoretical uncertainty in the calculation of the standard model rate. The sensitivity to deviations in the branching fraction are then given by $\Delta B/B \geq 1/\sqrt{\sigma I}$ where I is the time integrated luminosity, and σ is the specific cross-section we are investigating (e.g., $\sigma(l\gamma\gamma)$). For 100 fb^{-1} this sensitivity is *at best* 10% at the 1σ level. At 1000 fb^{-1}, which would require several years of high-luminosity LHC running, the sensitivity is at best 3% at the 1σ level. The ATLAS and CMS technical design reports [15] provide estimates on signal efficiencies and background rates [15] (both real and fake) to this process for the ATLAS and CMS detectors that will be installed at LHC. An updated study at Snowmass 96 estimated that with LHC data and NLC data ($\sqrt{s} = 500$ GeV and 200 fb^{-1}) one could determine $B(h \to \gamma\gamma)$ to $\lesssim 16\%$ accuracy [16] in the range $80 \lesssim m_h \lesssim 130$ GeV.

Such a measurement of $B(h \to \gamma\gamma)$ would be advantageous in piecing together large $\tan\beta$ supersymmetric theories. This is even true in the gauge mediated model discussed above. Gauge mediation models do not allow much mixing in the top squark sector because the A-terms are zero at the messenger scale. Therefore, the Higgs masses in this scenario were rather light, and the off-diagonal A_{12} entry in the Higgs mass matrix was not as large as it generically is in supergravity models. In the case of large mixing, the *hbb* coupling can be more volatile than we found in gauge mediation models and the Higgs mass can be larger.

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In Fig. 4 we see that the $l\gamma\gamma$ rate falls off rather rapidly when m_h is above 120 GeV, and the $Wh \to 3l$ rate increases rapidly. For Higgs masses above 120 GeV the most useful probe of Higgs coupling deviations may very well be the $3l$ signal [12, 6]. Here, the background could be as large as 1 fb [6]. With the detailed background cuts applied also to the signal we estimate that for a 130 GeV Higgs boson the deviation in the $3l$ signal at best could be measured to within 20% at 1σ with 100 fb^{-1}. With 10^4 fb^{-1} the rate could be measured to perhaps better than 8%. Therefore, it is extremely important, and even necessary to study the $3l$ signal in detail at the LHC as a probe of non-standard couplings to the light Higgs boson. Interestingly, the $3l$ signal from supersymmetric chargino and neutralino production is highly suppressed in the large $\tan\beta$ region [17]. The $Wh \to 3l$ signal therefore dominates the $3l$ signal from superpartner production.

It should be noted in conclusion that the gauge mediation model which we used to calculate these results is correlated with deviations in the $b \to s\gamma$ observable [4]. Substantial deviations in $b \to s\gamma$ are generically expected in all large $\tan\beta$ models. However, given the expected supersymmetry contribution to this amplitude, the expected experimental data, and the uncertainties in QCD corrections, it is likely that deviations will show up as anomalies at the few σ level at B-factories, and so conclusive evidence for new physics would not be ensured. A suppression of $l\gamma\gamma$ events at LHC for the given mass range would then be helpful to support the claim of supersymmetric theories with high $\tan\beta$. Furthermore, there are large regions of parameter space in the MSSM not necessarily tied to any minimal model which predicts small deviations in $B(b \to s\gamma)$ but $\mathcal{O}(1)$ corrections to the $h\bar{b}b$ coupling. These are best investigated by the Higgs boson observables discussed above.

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