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$b \rightarrow s\gamma$ and supersymmetry with large $\tan\beta$

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Abstract

We discuss the constraints on the supersymmetric parameter space from the decay mode $b \rightarrow s\gamma$ for large values of $\tan\beta$. We improve the theoretical prediction for the decay rate by summing very large radiative corrections to all orders in perturbation theory. This extends the validity of the perturbative calculation to the large $\tan\beta$ regime. This resummation of terms of order $\alpha_s^n \tan^{n+1}\beta$ is based on a recently proposed effective lagrangian for the Yukawa interaction of bottom quarks. Moreover, we identify an additional source of $\tan\beta$ -enhanced terms, which are of order $\alpha_s \tan\beta$ and involve the charged Higgs boson, and analyse their behaviour in higher orders of perturbation theory. After correcting the current expressions for this rare decay branching ratio at next-to-leading order, we obtain that, contrary to recent claims, the measured branching ratio of $b \rightarrow s\gamma$ constrains the supersymmetric parameter space in a relevant way, even if $\tan\beta$ is large.

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The measured branching ratio $\mathcal{BR}(b \rightarrow s\gamma)$ is known to provide a valuable constraint on the parameter space of the Minimal Supersymmetric Standard Model (MSSM). Increasing attention is devoted to scenarios of the MSSM with a large value of $\tan\beta$, the ratio of the two Higgs vacuum expectation values: this region of the parameter space is experimentally least constrained by the bounds coming from Higgs searches at the LEP experiments [1]. Theoretical interest in large $\tan\beta$ scenarios stems from GUT theories with bottom–top Yukawa unification, which require $\tan\beta = \mathcal{O}(50)$ [2–4]. Recently it has been claimed [6] that, if $\tan\beta$ is sufficiently large, the next-to-leading-order corrections to $\mathcal{BR}(b \rightarrow s\gamma)$ [7, 8] can wash out the constraints on the MSSM parameter space which arise from the leading-order calculation [10]. This claim is based on the fact that, at next-to-leading order, two large $\tan\beta$ -enhanced corrections occur: first, supersymmetric QCD (SQCD) corrections to the chargino-mediated transition lead to terms proportional to $\alpha_s \tan^2\beta$. Secondly, new $\tan\beta$ -enhanced contributions, absent at leading order, appear in SQCD corrections to the charged Higgs diagrams. The latter terms are of order $\alpha_s \tan\beta$. Therefore, for sufficiently large values of $\tan\beta$ the next-to-leading-order corrections may be of the same order as the leading-order ones, which involve one power of $\tan\beta$ less. Cancellations between the leading-order and the next-to-leading-order contributions of the supersymmetric particles to the decay amplitude may occur, depending on the specific values of the supersymmetry-breaking parameters. It was claimed in Ref. [6] that such cancellations do occur in the minimal supergravity model. In this letter, we shall analyse the dominant, $\tan\beta$ -enhanced, corrections to $\mathcal{BR}(b \rightarrow s\gamma)$. We shall explain how to resum large radiative corrections to this rare decay branching ratio to all orders in perturbation theory. After this, we shall reanalyse the modifications of the bounds on the minimal supergravity model parameter space which arise after radiative corrections are included. We conclude that the restriction on the sign of the Higgsino mass parameter μ , obtained at leading order [4], is robust under the inclusion of higher-order corrections.

In a previous publication [11] we have analysed the $\tan\beta$ -enhanced radiative corrections, which stem from the renormalization of the Yukawa coupling to down-type fermions. These corrections can efficiently be summed to all orders in perturbation theory and can be cast into an effective lagrangian:

$$\mathcal{L} = -h_d^{ij} \bar{d}_R^i H_1 Q_L^j - \delta h_d^{ij} \bar{d}_R^i H_2 Q_L^j - h_u^{ij} \bar{u}_R^i (i\tau_2 H_2^*) Q_L^j - \delta h_u^{ij} \bar{u}_R^i (i\tau_2 H_1^*) Q_L^j + \text{h.c.}, \quad (1)$$

where τ_2 is the usual two by two Pauli matrix, $Q_L = (u, d)_L$, and a gauge-invariant contraction of weak and colour indices has been implicitly assumed. In the above, we have ignored small $SU(2)_L$ breaking effects. The dominant contributions to the couplings δh_d and δh_u are induced via SQCD corrections. Assuming that the right- and left-handed soft-supersymmetry-breaking mass parameters are generation-independent, they are proportional to the couplings h_d and h_u :

$$\delta h_d^{\text{SQCD}} = h_d \frac{2\alpha_s}{3\pi} M_{\tilde{g}} \mu I(m_{\tilde{b}_L}, m_{\tilde{b}_R}, M_{\tilde{g}}),$$

$$\delta h_u^{\text{SQCD}} = h_u \frac{2\alpha_s}{3\pi} M_{\tilde{g}} \mu I(m_{\tilde{t}_L}, m_{\tilde{t}_R}, M_{\tilde{g}}), \quad (2)$$

where $M_{\tilde{g}}$ is the gluino mass and $m_{\tilde{b}_{L,R}}, m_{\tilde{t}_{L,R}}$ are the left- and right-handed mass parameters of the down- and up-squarks respectively. The dependence of this loop integral on its parameters is given by

$$I(a, b, c) = \frac{1}{(a^2 - b^2)(b^2 - c^2)(a^2 - c^2)} \left(a^2 b^2 \log \frac{a^2}{b^2} + b^2 c^2 \log \frac{b^2}{c^2} + c^2 a^2 \log \frac{c^2}{a^2} \right). \quad (3)$$

Once SU(2) breaking effects are included, these left- and right-handed squark mass parameters should be replaced by the squark mass eigenvalues, $m_{\tilde{t}_{1,2}}, m_{\tilde{b}_{1,2}}$ (or eventually $m_{\tilde{s}_L}$ in the charged Higgs boson vertex corrections involving the left-handed strange quark)¹. Also, the appropriate CKM angles, which distinguish the charged and neutral Higgs couplings, should be included. A precise description of the $\tan\beta$ -enhanced couplings, proportional to h_b , of the Higgs particles to top and bottom quarks was discussed in detail in Ref. [11]. In our application to $b \rightarrow s\gamma$ we need the coupling of the charged Higgs boson to the right-handed bottom and left-handed top quarks, for which, ignoring small CKM angle effects, an all-order resummation of the large $\tan\beta$ -enhanced corrections is achieved by replacing the tree-level relation between the coupling h_b in Eq. (1) and the bottom mass by

$$h_b = \frac{g}{\sqrt{2} M_W \cos\beta} \frac{\overline{m}_b(Q)}{1 + \Delta m_b^{\text{SQCD}}}. \quad (4)$$

Here g is the SU(2) gauge coupling, M_W is the W-boson mass and \overline{m}_b is the bottom mass renormalized in a mass-independent renormalization scheme like $\overline{\text{MS}}$ [12] (as we explained in Ref. [11], the above relation can be simply understood in terms of the effective lagrangian of Eq. (1)). The $\tilde{t}bH^+$ vertex is renormalized at the scale Q , which enters \overline{m}_b in Eq. (4). When applied to $b \rightarrow s\gamma$ the scale Q equals the scale μ_W , at which top and H^+ are integrated out. Following Eqs. (1) and (2), the resummed $\tan\beta$ -enhanced SQCD corrections are contained in [3, 4]

$$\Delta m_b^{\text{SQCD}} = \frac{2\alpha_s}{3\pi} M_{\tilde{g}} \mu \tan\beta I(m_{\tilde{b}_1}, m_{\tilde{b}_2}, M_{\tilde{g}}), \quad (5)$$

where α_s should be evaluated at a scale of the order of the masses entering I . If these masses differ so much that they induce large logarithms, the scale of α_s is set by the largest of these masses. The necessity of including these large corrections, proportional to Δm_b^{SQCD} , in the computation of $\mathcal{BR}(b \rightarrow s\gamma)$ at large values of $\tan\beta$, was first emphasized in Ref. [5].

¹The vertex corrections to the charged and neutral Higgs will involve the superpartners of the corresponding quarks appearing in the external legs.

To consider the dominant chargino–squark contributions, we can write an effective lagrangian analogous to Eq. (1), by ignoring the corrections $\delta h_{u,d}$ and replacing the charged Higgs and one of the two quarks by their superpartners. Thereafter, we can proceed in exactly the same way: by replacing the Yukawa coupling h_b in the stop–bottom–chargino coupling with Eq. (4) we encounter all terms of order $\alpha_s^n \tan^n \beta$ for $n = 0, 1, 2, 3, \dots$. Since the leading-order chargino–stop contribution grows linearly with $\tan \beta$, this resummation leads to contributions to $\mathcal{BR}(b \rightarrow s\gamma)$ of order $\alpha_s^n \tan^{n+1} \beta$. Similar to the charged Higgs case, one can therefore implement the resummation into an existing leading-order calculation by simply replacing h_b by its expression given in Eq. (4). We have done this with the Wilson coefficients C_7 and C_8 , which contain the supersymmetric terms relevant to the $b \rightarrow s\gamma$ amplitude. After expanding Eq. (4) to first order in α_s we have indeed reproduced the dominant terms of the known next-to-leading-order result of [7]: these terms of order $\alpha_s \tan^2 \beta$ stem from the chargino–squark contributions and equal $-\Delta m_b^{\text{SQCD}}$ times the corresponding leading-order piece involving h_b . While the authors of [6] have correctly connected the $\alpha_s \tan^2 \beta$ term to the sbottom mixing, they have erroneously claimed the absence of terms of order $\alpha_s^2 \tan^3 \beta$. Since the sbottom mixing terms proliferate into the Yukawa counterterm through Δm_b^{SQCD} , they iteratively show up in any order of perturbation theory, as explained in detail in [11]. Expanding Eq. (4) to first order in α_s also reproduces one term of order $\alpha_s \tan \beta$ associated with SQCD corrections to the $H^+ \bar{t}_L b_R$ vertex in the charged Higgs diagram.

Interestingly enough, as has been shown in Ref. [7], there is an additional source of $\tan \beta$ -enhanced corrections in the charged Higgs diagrams: while the tree-level $H^+ \bar{t}_R s_L$ vertex is suppressed by $1/\tan \beta$, this vertex suppression is lifted at the one-loop level, so that the next-to-leading-order charged-Higgs contribution to $\mathcal{BR}(b \rightarrow s\gamma)$ is $\tan \beta$ -enhanced with respect to the leading-order one. This feature originates from the loop-induced flavour-violating couplings δh_u^{ij} described in Eq. (1), which involve a right-handed top quark, a left-handed strange quark, and H_1^+ , the charged component of the doublet H_1 . The disappearance of the $\tan \beta$ suppression is due to the fact that, at large values of $\tan \beta$, the physical charged Higgs can be approximately identified with H_1^+ , while its component on H_2^+ is suppressed by $1/\tan \beta$. The absence of a $\tan \beta$ suppression in this loop-induced coupling of the charged Higgs leads to a term of order $\alpha_s \tan \beta$ in the charged Higgs diagram, where the factor of $\tan \beta$ comes from the bottom-quark Yukawa coupling in Eq. (4).²

Schematically, the resummation of the $\tan \beta$ -enhanced corrections to the chargino and charged Higgs contributions can be summarized as follows: the $\tan \beta$ -enhanced chargino contributions

²While the term involving δh_u^{ij} is important for the $H^+ \bar{t} s$ vertex in the effective lagrangian in Eq. (1), it is subdominant in the effective $H^+ \bar{t} b$ coupling. Therefore it has not been considered in Ref. [11], which deals with Higgs and top decays.

to $\mathcal{BR}(b \rightarrow s\gamma)$ are

$$\mathcal{BR}(b \rightarrow s\gamma)|_{\chi^\pm} \propto \mu A_t \tan\beta f(m_{\tilde{t}_1}, m_{\tilde{t}_2}, m_{\tilde{\chi}^+}) \frac{m_b}{v(1 + \Delta m_b)}, \quad (6)$$

where all dominant higher-order contributions are included through Δm_b , and f is the loop integral appearing at one loop. The relevant charged-Higgs contributions to $\mathcal{BR}(b \rightarrow s\gamma)$ in the large $\tan\beta$ regime are

$$\mathcal{BR}(b \rightarrow s\gamma)|_{H^\pm} \propto \frac{m_b (h_t \cos\beta - \delta h_t \sin\beta)}{v \cos\beta (1 + \Delta m_b)} g(m_{H^\pm}, m_t), \quad (7)$$

where we have left the $\cos\beta$ and $\sin\beta$ factors associated with their sources, and g is the loop-integral appearing at the one-loop level. In the above δh_t proceeds from the flavour-violating coupling δh_u in Eq. (1), where, as we already explained, the squark masses should be replaced by the superpartners of the left-handed strange squark and of the right-handed top squark. Since the left-handed and right-handed top squarks mix in a non-trivial way, the loop integral $I(a, b, c)$ should be replaced by the sum of two-loop integrals for the two stop mass eigenstates with appropriate projection factors $\cos^2\theta_{\tilde{t}}$ and $\sin^2\theta_{\tilde{t}}$, respectively [7],

$$\delta h_t = h_t \frac{2\alpha_s}{3\pi} \mu M_{\tilde{g}} \left(\cos^2\theta_{\tilde{t}} I(m_{\tilde{s}_L}, m_{\tilde{t}_2}, M_{\tilde{g}}) + \sin^2\theta_{\tilde{t}} I(m_{\tilde{s}_L}, m_{\tilde{t}_1}, M_{\tilde{g}}) \right). \quad (8)$$

Observe that, as we first explained in Ref. [11] and follows from Eqs. (6) and (7), the corrections to the bottom-quark Yukawa coupling originate from mass counterterms and lead to terms of order $\tan^n\beta$ at n th-order of perturbation theory. The one-loop induced $H^+ \tilde{t}_R s_L$ vertex, instead, is of order $\tan^0\beta$, and higher-order corrections to it cannot induce any $\tan\beta$ -enhanced term. The above guidelines are useful to extend the validity of the calculation presented in Ref. [7] to large values of $\tan\beta$.

It should be stressed that, within the conventions of Ref. [7], the sign of the one-loop contribution, proportional to $\mu H_2(u_i, x_j)$ in the notation of Appendix A of Ref. [7], to the couplings of the charged Goldstone and Higgs fields to the right-handed top and left-handed down-like quark should be the ones that arise from Eq. (1), which are opposite to the ones stated in Ref. [7].³ A change of sign in these couplings leads to an inversion in the sign of those $\tan\beta$ -enhanced charged-Higgs contributions to $b \rightarrow s\gamma$ not related with Δm_b (the sign of δh_t in Eq. (7)). Once the correct sign is used, we find that the higher-order $\tan\beta$ -enhanced corrections to the charged-Higgs and chargino contributions to $b \rightarrow s\gamma$ lead to an enhancement (suppression) of these contributions for negative (positive) values of $\mu M_{\tilde{g}}$.

³The authors of Ref. [7] have independently detected these sign errors. They posted a revised version of Ref. [7] to the hep-ph archive during completion of this article.

At large values of $\tan\beta$, the leading-order chargino contributions to the amplitude of the decay rate $b \rightarrow s\gamma$ are proportional to $A_t\mu$. In supergravity models, the sign of A_t is opposite to the one of the gaugino masses. This sign relation holds unless the boundary values of A_t at the high-energy input scale is one order of magnitude larger than the gaugino soft-supersymmetry-breaking mass parameters [2, 13]. Since the charged Higgs-top diagram leads to a contribution to the amplitude of the decay $b \rightarrow s\gamma$ of the same sign as the SM contribution, and the relative sign of the chargino contribution is governed by the sign of $A_t\mu$, negative values of $A_t\mu$ (or equivalently, defining the gaugino masses as positive, positive values of μ) are necessary in order to render values of $\mathcal{BR}(b \rightarrow s\gamma)$, in agreement with the ones observed experimentally [4, 10]. The present study shows that for negative values of μ the next-to-leading-order corrections to the charged Higgs and the chargino-stop contributions further enhance the $b \rightarrow s\gamma$ decay amplitude. Therefore, it follows that even after considering higher-order effects, positive values of μ are necessary in order to obtain correct values for $\mathcal{BR}(b \rightarrow s\gamma)$ within minimal supergravity models, for which the sign of A_t at low energies tends to be negative [14].

Contrary to our results, by using the expressions given in Ref. [7], one would obtain an incorrect strong suppression (and even a change of sign at sufficiently large values of $\tan\beta$) of the charged-Higgs contributions for negative values of μ . Under these conditions, the authors of Ref. [6] incorrectly found that, for negative values of μ , acceptable values of $\mathcal{BR}(b \rightarrow s\gamma)$ may be obtained. Analogously, for positive values of μ they found an incorrect enhancement of the charged-Higgs contribution, rendering it difficult to find acceptable values for the rare bottom quark decay rate under study.

We give here a simple recipe, based on Eqs. (7) and (6), for implementing the all-order resummation of $\tan\beta$ -enhanced contributions into existing computations for the $b \rightarrow s\gamma$ amplitude: to this end multiply first in Eq. (4) of [7] the terms proportional to $1/\cos\beta$ with $1/(1 + \Delta m_b^{\text{SQCD}})$. Second, multiply the term proportional to A_d in Eq. (53) of first reference in [8], which contains the contribution of the charged Higgs, with $1/(1 + \Delta m_b^{\text{SQCD}})$. Third, add a term of the form $F_i^{(2)}(x) \cdot (1 - 1/(1 + \Delta m_b^{\text{SQCD}}))$ in Eq. (28) of first reference in [8], for $i = 7, 8$. Fourth, delete the term proportional to $\mu \tan\beta H_2(x_1, x_2)$ in H_d , U_d and $\Delta_d^{(2)}$, which are defined in appendices A.2.–A.4. of [7]. $\Delta_b^{(2)}$ enters the Wilson coefficients $C_{7,8}$ through $G_{7,8}^{\chi:2}$ in Eq. (13) and H_b , U_b enter these coefficients through Eqs. (25), (26) of [7]. Multiply the remaining $\tan\beta$ -enhanced terms, proportional to $\mu \tan\beta H_2(u_i, x_2)$, in the charged-Goldstone and charged-Higgs boson contributions by $1/(1 + \Delta m_b^{\text{SQCD}})$. The presence of these additional $\tan\beta$ -enhanced terms in the charged Goldstone diagram ensures the proper decoupling of the supersymmetric corrections in the limit of large supersymmetric particle masses.

To demonstrate which is the effect of the improvement in the theoretical prediction for $\mathcal{BR}(b \rightarrow s\gamma)$ developed in this work, we compare in Figs. 1 and 2 our result, including a resummation of the dominant $\tan\beta$ -enhanced radiative corrections to all orders of perturbation

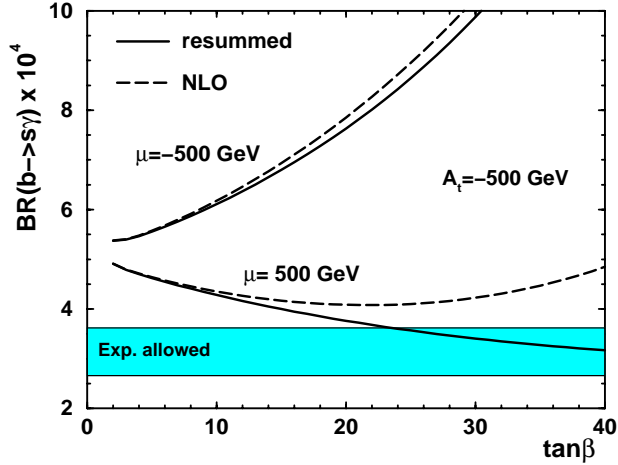


Figure 1: Comparison of the result for $\mathcal{BR}(b \rightarrow s\gamma)$, as obtained in this work, as a function of $\tan\beta$, with the NLO expression of [7], where we have replaced the appropriate signs in the next-to-leading-order corrections, as explained in the text. The charged-Higgs boson mass is 200 GeV and the light stop mass is 250 GeV. The values of μ and A_t are indicated in the plot while M_2 , the gluino, heavy-stop and down-squark masses are set at 800 GeV.

theory, with the next-to-leading-order result of Ref. [7], with the appropriate sign corrections explained above, for typical values of the supersymmetric parameters. In Fig. 1, a negative value of $A_t = -500$ GeV at low energies is chosen, as is usually the case in supergravity models. For the experimental measurement of the $b \rightarrow s\gamma$ branching ratio, we use the combined result of CLEO [15] and ALEPH [16], $\mathcal{BR}(b \rightarrow s\gamma) = (3.14 \pm 0.48) \times 10^{-4}$. Fig. 2 shows the complementary case of positive $A_t = 500$ GeV. Notice that, for $A_t > 0$, the behaviour of the MSSM result never differs crucially from the SM prediction: there is always some cancellation between competing terms, either between the charged Higgs and the chargino contributions (for values of $\mu < 0$) or between the LO and the $\tan\beta$ -enhanced NLO corrections (for values of $\mu > 0$).

In Fig. 3 we confront the resummed theoretical prediction for the $b \rightarrow s\gamma$ branching ratio with the experimental result to find the excluded region –from the curves, to the left– in the (M_{H^+}, M_{χ^+}) plane. Curves are shown for $\tan\beta = 10, 20$, and $m_{\tilde{t}_2} = 100, 300$ GeV, after a scan in the remaining parameters, as described in the caption.

In conclusion, we have presented in this article a method allowing to extend the validity of the NLO computation of the decay rate $b \rightarrow s\gamma$ to the large- $\tan\beta$ regime, which is based on a resummation to all orders of the dominant $\tan\beta$ -enhanced radiative corrections. We have proved that, contrary to recent claims, the next-to-leading-order corrections preserve the

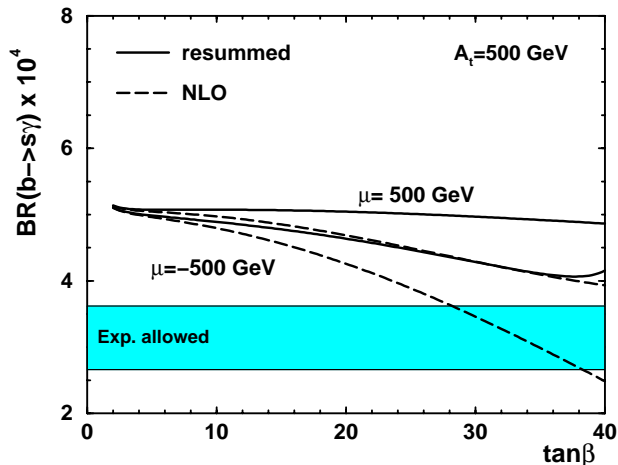


Figure 2: As in Fig. 1, but for positive $A_t = 500 \text{ GeV}$.

basic features of the leading-order results in constraining positive values of μA_t within minimal supergravity models, for which the low-energy values of A_t tend to be negative.

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Note added

After completion of this work similar results have been presented in [17]. While our paper focuses on the all-order resummation of $\tan\beta$ -enhanced SQCD corrections and the clarification of the findings of [6], the authors of [17] have discussed dominant two-loop contributions arising from $\tan\beta$ -enhanced SQCD and supersymmetric electroweak corrections and of large logarithms which arise for heavy superpartner masses.

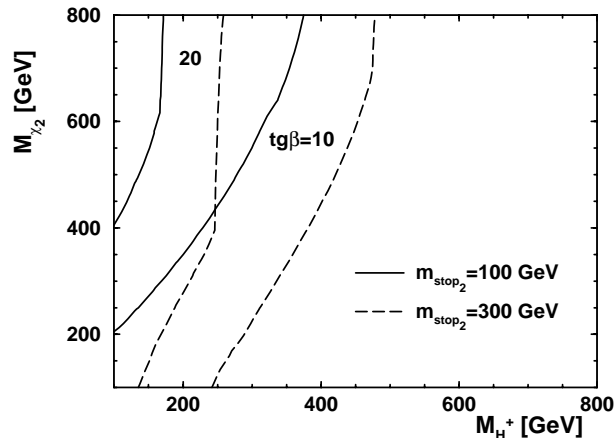


Figure 3: Combined bound on the charged Higgs and chargino masses, for various values of the mass of the lightest stop, $m_{\tilde{t}_2}$, and $\tan\beta$. The excluded region corresponds to light (heavy) charged Higgs (chargino) masses. We have scanned for $m_{\tilde{t}_2} < m_{\tilde{t}_1} \leq 1$ TeV, $m_{\tilde{\chi}_2^+} < m_{\tilde{\chi}_1^+} \leq 1$ TeV and $|A_t| \leq 500$ GeV, the rest of SUSY masses have been set at 1 TeV.

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