Relaxing $b \to s\gamma$ Constraints on the Supersymmetric Particle Mass Spectrum

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Abstract

We consider the radiative decay $b \to s\gamma$ in a supersymmetric extension of the standard model of particle interactions, where the $b$-quark mass is entirely radiative in origin. This is accomplished by the presence of nonholomorphic soft supersymmetry breaking terms in the Lagrangian. As a result, the contributions to the $b \to s\gamma$ amplitude from the charged Higgs boson and the charginos/neutralinos are suppressed by $1/\tan^2 \beta$ and $\mathcal{O}(\alpha/\alpha_s)$ respectively, allowing these particles to be lighter than in the usual supersymmetric model. Their radiatively generated couplings differ from the usual tree-level ones and change the collider phenomenology drastically. We also study how this scenario may be embedded into a larger framework, such as supersymmetric SU(5) grand unification.

1. Introduction. The minimal supersymmetric standard model (MSSM) [1] is one of the most popular extensions of the standard model (SM). In recent years, both theorists and experimentalists have devoted enormous amounts of time to study its predictions. While superpartner masses are expected to be below 1 TeV in some scenarios, there is actually a lot of uncertainty regarding the soft supersymmetry (SUSY) breaking sector of the theory. In the most general case, the MSSM contains more than one hundred free parameters. Nevertheless, present collider data as well as low-energy experiments are starting to place nontrivial constraints on the supersymmetric particle mass spectrum.

One of the processes known to put stringent constraints on new physics is the radiative decay $b \to s\gamma$ with a branching ratio experimentally determined to be in the range [2]

$$2 \times 10^{-4} < \text{BR}(B \to X_s \gamma) < 4.5 \times 10^{-4}. \quad (1)$$

This result agrees with the SM prediction. On the other hand, in the MSSM framework, the decay $b \to s\gamma$ receives large additional contributions from charged Higgs-boson, chargino, neu-
The leading-order quantum-chromodynamics (LO QCD) corrections to $b \to s\gamma$ are known in the MSSM for arbitrary flavour structures [3,4], while the next-to-leading-order (NLO) analyses have been performed only for specific scenarios [5]. The charged Higgs-boson contribution always adds constructively to that of the SM (i.e. the W-boson contribution). The magnitudes of chargino and neutralino contributions depend strongly on $\tan \beta$. For large values of $\tan \beta$, the chargino contribution becomes the dominant one. In that case, its sign is determined by that of the $\mu$ parameter and can add constructively or destructively. Despite its strong interaction, the gluino contribution is generally the subdominant one. Its sign depends on the details of the flavour structure of the model [4], e.g. in SUSY SU(5) with a radiatively driven mass spectrum, it adds constructively to the chargino contribution [6].

An important issue for understanding the $b \to s\gamma$ constraints on SUSY models, recently reemphasized in Ref. [7], is the proper inclusion of loop corrections to the $b$-quark Yukawa coupling [8] when calculating the $b \to s\gamma$ rate. These are enhanced for large $\tan \beta$ and their sign is also determined by the $\mu$ parameter. This implies a strong correlation between the values of the $b$-quark Yukawa coupling, SUSY model parameters, and the $b \to s\gamma$ constraints.

The gauge-coupling unification in the MSSM strongly suggests that there is a grand unified theory (GUT) above the unification scale $M_{\text{GUT}} \sim 2 \times 10^{16}$ GeV, such as SU(5) or SO(10) which allows tau-bottom or tau-bottom-top Yukawa coupling unification respectively. However, successful Yukawa unification [9] is achieved only for one sign of the $\mu$ parameter, $\text{sign}(\mu) = -$ in our convention, and for very large values of $\tan \beta$, say $\sim 30-50$ for SU(5) and $\sim 50$ for SO(10) [10]. For $\text{sign}(\mu) = -$, all dominant contributions to $b \to s\gamma$ add constructively, implying thus very strong constraints on these SUSY scenarios. Typically, the SUSY mass scale must exceed 1 TeV to be consistent with the bound given by Eq.(1) [10]. The charged Higgs-boson mass is forced to be large in this case, typically above the reach of the Tevatron at Fermilab as well as that of the Large Hadron Collider (LHC) at CERN. In more general scenarios, $\text{sign}(\mu) = +$ allows cancellations between different terms. However for large $\tan \beta$, the cancellation can happen only for a restricted part of the parameter space and stringent constraints may still be in force. In conclusion, the $b \to s\gamma$ bound of Eq.(1) implies strong constraints on MSSM parameters in general, and on GUT scenarios in particular.

There are some proposals to satisfy the $b \to s\gamma$ constraints which still allow light sparticle masses accessible at future colliders. In Ref. [11], it has been pointed out in the context of SO(10) GUT that with $m, A \gg M_{1/2}$, where $m$ and $A$ denote generically common masses of matter multiplets and soft $A$ terms respectively, and $M_{1/2}$ is the common gaugino mass, there is a part of the parameter space for which the $b \to s\gamma$ rate is in the range given by Eq.(1). In that case, the scalar superpartner masses are very large and suppress the new SUSY contributions to $b \to s\gamma$ whereas the wino masses can still be light. This scenario implies that the only discoverable SUSY particles are light charginos and neutralinos. The new Higgs bosons are heavy and the collider phenomenology discussed in Ref. [12] is not allowed.

The authors of Ref. [4,7] argue instead that there might be new flavour violation present in the squark sector of the model which modifies and enhances the gluino contribution which then cancels the other large SUSY contributions. This scenario requires the introduction of new
unknown flavour physics in general SUSY models and is not realized in GUTs [6]. Also, the cancellation of two large terms cannot be considered a natural solution.

The purpose of this letter is to show that the dominant SUSY contributions to $b \to s \gamma$ may be naturally suppressed if the SUSY radiative corrections to the $b$-quark Yukawa coupling are large. We consider the limit of vanishingly small down-quark Yukawa couplings so that the corresponding quark masses are generated radiatively [13,14]. Making the usual assumption that the trilinear $A$ terms are proportional to Yukawa couplings, we are forced to introduce the nonholomorphic $A'$ terms [14–17] to give a correct mass to the $b$-quark. After all, the soft $A'$ terms should be included in the complete SUSY Lagrangian on general grounds. We find that in this scenario, the charged Higgs-boson and the dominant chargino contributions to $b \to s \gamma$ rate are suppressed by $1/\tan^2 \beta$ and $\mathcal{O}(\alpha/\alpha_s)$ respectively. Therefore, the soft, radiatively induced couplings of these particles change the Tevatron and LHC phenomenology considerably by reducing their production rates. This scenario can also be embedded into a GUT framework, such as SUSY SU(5), removing the stringent constraint on the MSSM parameter space coming from Yukawa unification. We discuss the sparticle mass spectrum in that case.

2. Proposed model and the radiative decay $b \to s \gamma$. In the following we work with the usual particle content of the MSSM [1]. However, we assume that the Yukawa coupling matrix in the $(d, s, b)$ quark sector is vanishing, i.e. $f_{d_{ij}} = 0$. In this case, the relevant MSSM superpotential $W$ for quark and Higgs left-chiral superfields is

$$W = Q_i(f_{u_{ij}})U_j^c H_2 - \mu H_1 H_2. \quad (2)$$

Further we make the usual assumption that the trilinear soft SUSY breaking terms $a_{ij}$ have the same structure as the Yukawa coupling matrices: $a_{ij} = A \cdot f_{ij}$. Thus, neglecting leptons, the most general soft SUSY breaking terms in the MSSM are given by

$$- \mathcal{L}_{\text{soft}} = \tilde{Q}^c_i (m_{\tilde{Q}}^2)_{ij} \tilde{Q}_j + \tilde{U}^c_i (m_{\tilde{U}}^2)_{ij} \tilde{U}_j + \tilde{D}^c_i (m_{\tilde{D}}^2)_{ij} \tilde{D}_j + m_{\tilde{H}_1}^2 H_1^1 H_1 + m_{\tilde{H}_2}^2 H_2^1 H_2$$

$$+ (\tilde{Q}_i (A_u \cdot f_{u_{ij}}) \tilde{U}_j^c H_2 - \tilde{Q}_i (A_d' \cdot f_{u_{ij}}) \tilde{D}_j^c (i \tau_2 H_2^*) + B_H H_1 H_2 + \frac{1}{2} M_1 \tilde{B} \tilde{B} + \frac{1}{2} M_2 \tilde{W}^a \tilde{W}^a + \frac{1}{2} M_3 \tilde{g}^a \tilde{g}^a + h.c.), \quad (3)$$

where the $A'_d$ term is the nonholomorphic soft term and it does not cause quadratic divergences. Therefore, it has been emphasized in Ref. [15], and more recently in Ref. [16], that these terms should be included in the MSSM to study its low-energy phenomenology; their omission cannot be justified in the general context. In the framework of high-energy physics, the nonholomorphic terms are generated, and not necessarily suppressed, in scenarios of spontaneous SUSY breaking such as those being mediated by supergravity [17]. In strongly coupled supersymmetric gauge theories, the nonholomorphic soft terms occur naturally [18]. Without understanding the real origin of SUSY breaking, we should keep the $A'$ terms in the general soft Lagrangian of the MSSM, such as Eq.(3). Radiatively induced quark masses from the $A'$ terms have been recently studied in Ref. [14].
The \((u, c, t)\) quark masses come directly at tree level from the hard Yukawa interaction in the superpotential Eq. (2). However, the \((d, s, b)\) quark masses must be generated radiatively \([13, 14]\). Taking only the dominant gluino-mediated contribution and neglecting intergenerational mixings, the one-loop soft bottom quark mass is given by

\[
m_b = -\frac{2\alpha_s}{3\pi} m_{\tilde{g}} m_t A'_b I(m_{\tilde{b}_1}^2, m_{\tilde{t}_2}^2, m_{\tilde{g}}^2). \tag{4}
\]

Here the necessary chirality violation is due to the gluino mass insertion and the term \(m_t A'_b\) is the off-diagonal entry of the sbottom mass matrix. The loop function \(I(m_1^2, m_2^2, m_3^2)\) is given by \([8]\)

\[
I(m_1^2, m_2^2, m_3^2) = -\frac{m_2^2m_2^2\ln(m_1^2/m_2^2) + m_3^2m_3^2\ln(m_2^2/m_3^2) + m_1^2m_1^2\ln(m_3^2/m_1^2)}{(m_1^2 - m_2^2)(m_2^2 - m_3^2)(m_3^2 - m_1^2)}. \tag{5}
\]

Obtaining the correct bottom quark mass via Eq. (4) implies stringent constraints on the model parameters.

As the tree-level bottom Yukawa term is missing in the superpotential Eq. (2), the \(\bar{t}_L b_R H^+\) coupling is induced radiatively by the diagram depicted in Fig. 1. The induced soft Lagrangian at the one-loop level can be expressed as

\[
\mathcal{L}_{\text{Yukawa}}^{\text{rad}} = -\frac{m_b}{v \tan \beta} \bar{t} [r_1 P_L + r_2 P_R] b H^+ + h.c., \tag{6}
\]

where the form factors \(r_1\) and \(r_2\) are given by

\[
r_1 = \frac{\sin 2\theta_i \sin 2\theta_b}{4I(m_{\tilde{b}_1}^2, m_{\tilde{b}_2}^2, m_{\tilde{g}}^2)} \left[ C_0(0, 0, m_{H}^2; m_{\tilde{t}_1}^2, m_{\tilde{g}}^2, m_{\tilde{b}_1}^2) + C_0(0, 0, m_{H}^2; m_{\tilde{t}_2}^2, m_{\tilde{g}}^2, m_{\tilde{b}_2}^2) \\
- C_0(0, 0, m_{H}^2; m_{\tilde{t}_1}^2, m_{\tilde{g}}^2, m_{\tilde{b}_2}^2) - C_0(0, 0, m_{H}^2; m_{\tilde{t}_2}^2, m_{\tilde{g}}^2, m_{\tilde{b}_1}^2) \right], \tag{7}
\]

\[
r_2 = \frac{1}{I(m_{\tilde{b}_1}^2, m_{\tilde{b}_2}^2, m_{\tilde{g}}^2)} \left[ \cos^2 \theta_i \sin^2 \theta_b C_0(0, 0, m_{H}^2; m_{\tilde{t}_1}^2, m_{\tilde{g}}^2, m_{\tilde{b}_1}^2) + \sin^2 \theta_i \cos^2 \theta_b C_0(0, 0, m_{H}^2; m_{\tilde{t}_2}^2, m_{\tilde{g}}^2, m_{\tilde{b}_2}^2) \right].
\]
Fig. 2. The scatter plot for the form factors $r_1$ and $r_2$ in the GUT scenario.

\[
\begin{align*}
&\cos^2 \theta_1 \cos^2 \theta_\tilde{b} C_0(0, 0, m_H^2, m_{\tilde{t}_1}^2, m_{\tilde{g}}^2, m_{\tilde{b}_2}^2) \\
&+ \sin^2 \theta_1 \sin^2 \theta_\tilde{b} C_0(0, 0, m_H^2, m_{\tilde{t}_2}^2, m_{\tilde{g}}^2, m_{\tilde{b}_2}^2) \end{align*}
\]

Here the three-point functions $C_0(0, 0, m_H^2; m_{\tilde{t}_1}^2, m_{\tilde{b}_2}^2, m_{\tilde{b}_1}^2)$ are defined in [14] and we do not present them here. In the limit $m_H \to 0$, which is a good approximation if $m_H \ll m_{\tilde{g}}, m_{\tilde{b}}, m_{\tilde{t}}$ the three-point functions become $C_0(0, 0, 0; m_{\tilde{b}_1}^2, m_{\tilde{b}_2}^2) \to I(m_{\tilde{t}_1}^2, m_{\tilde{g}}^2, m_{\tilde{b}_1}^2)$.

As with the usual hard tree-level coupling which we have omitted, the Lagrangian Eq.(6) is also proportional to $m_b$. However, the most important feature to notice is that the coupling is now suppressed by $\tan \beta$ and not enhanced by it. This is because the nonholomorphic $A'$ term couples to $H_2$ rather than to $H_1$ as in the usual case. The induced couplings are momentum-dependent and are characterized by the form factors $r_1$ and $r_2$. If the couplings arise form the tree-level superpotential, then to lowest order, $r_1 = 0$ and $r_2 = 1$. In our scenario, the deviation from these values is expected to be small, say a few percent. Assuming that all stop and sbottom masses are degenerate as a first approximation, $m_{\tilde{b}_1} \approx m_{\tilde{b}_2} \approx m_{\tilde{t}_1} \approx m_{\tilde{t}_2}$, we have $r_1 \to 0$ independently of the stop and sbottom mixing angles. If the charged Higgs boson is lighter than the coloured sparticles, then we also have $r_2 \to 1$ in this case. To show the realistic values of the form factors, we plot in Fig. 2 the values of $r_1$ against $r_2$ for the scenario (to be considered in detail in the next Section) where the SUSY mass spectrum is generated via renormalization-group running, assuming GUT conditions. Thus the interaction term proportional to $r_1$ in Eq.(6) is just a small correction to the hard Yukawa term arising from the superpotential Eq.(2) and we neglect it in further discussion. The second term in Eq.(6) mimics the missing hard term in the superpotential but is suppressed by $1/\tan^2 \beta$ compared to it.

For quarks which obtain their masses radiatively, their higgsino couplings are also generated radiatively. The details for the neutral higgsinos are given in Ref. [14]; similar arguments apply.
also for the charged higgsinos. Without going into details, the most important feature of the radiatively induced couplings of the right-handed $b$-quark to the higgsino $\tilde{H}$ and squark is that it is induced by loops involving binos. Thus the radiatively induced couplings for $b_R$ are always suppressed by $\mathcal{O}(\alpha/\alpha_s)$ compared to the hard couplings. We show below that this suppression factor also applies to the dominant chargino and neutralino contributions to the $b \to s\gamma$ amplitude.

The effective Hamiltonian for $b \to s\gamma$ in SUSY models can be expressed in two terms:

$$H_{\text{eff}} = H_{\text{eff}}^{CKM} + H_{\text{eff}}^{\tilde{g}},$$

where

$$H_{\text{eff}}^{CKM} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^{*} \sum_i C_i(\mu) \mathcal{O}_i(\mu)$$

contains the SM as well as the charged Higgs-boson, chargino, and neutralino contributions with the same flavour structure as in the SM. The explicit formulas including LO QCD corrections can be found in Ref. [3]. The gluino contribution $H_{\text{eff}}^{\tilde{g}}$ may exhibit in addition the new flavour violation present in general SUSY models but absent in our scenario; details including LO QCD corrections can be found in Ref. [4]. The decay width of $b \to s\gamma$ can be written as

$$\Gamma(b \to s\gamma) = \frac{m_b^5 G_F^2 |V_{tb}V_{ts}^{*}|^2 \alpha}{32\pi^4} |C_{7}^{\text{eff}}|^2,$$

where $|C_{7}^{\text{eff}}|^2 = |C_7 + C_7^\prime|^2 + |C_7^\prime|^2$. Here $C_7$ stands for the total contribution from the effective Hamiltonian Eq.(10), while $C_7^\prime$ and $C_7^\prime$ arise from gluino loops. The present experimental result Eq.(1) implies the allowed range:

$$0.25 < |C_{7}^{\text{eff}}| < 0.375.$$  

Let us now consider different SUSY contributions to $b \to s\gamma$ in our scenario. The charged
Higgs-boson contribution is induced by two different chiral structures: the SM-like one with the chirality flip in the external $b$-quark line which is induced by the hard top Yukawa interactions, and the one with the chirality flip in the internal $t$-quark line which is induced by the radiative coupling Eq.(6). The latter diagram, presented in Fig. 3, is actually a two-loop diagram. A rigorous calculation would require that it be performed in two loops. However, since the coloured sparticles are expected to be much heavier than the Higgs bosons, we can integrate them out at their mass scale and assume with a good accuracy that the $H^+$ contribution is given by the one-loop diagram in Fig. 3 with the radiative coupling Eq.(6). In that case, the latter contribution to $C_7$, which is not suppressed by $\tan^2 \beta$ if the bottom Yukawa couplings are hard, becomes suppressed by $1/\tan^2 \beta$ implying that the full $H^+$ contribution to $C_7$ is suppressed by the same factor. Explicitly,

$$C_7^{H^+}(m_W) = \frac{1}{2} \frac{x_{tH}}{\tan^2 \beta} \left[ \frac{2}{3} f_1(x_{tH}) + f_2(x_{tH}) + r_2 \left( \frac{2}{3} f_3(x_{tH}) + f_4(x_{tH}) \right) \right], \quad (13)$$

where $x_{tH} = M_{H^+}^2 / m_t^2$ and the Inami-Lim type functions $f_i(x_{tH})$ can be found in Ref. [3]. This equation implies an important result: the mass bounds on the charged Higgs boson coming from the measurement of $b \to s\gamma$ can be relaxed and $H^+$ can be light.

The chargino and neutralino contributions to $C_7$ follow the same two chiral patterns discussed above. In this case the chirality flip may occur in the internal chargino/neutralino line implying the large enhancement factors $m_{\tilde{\chi}} / m_b$. These enhanced terms, induced by the higgsinos, dominate the chargino/neutralino contributions. However, as we argued before, the relevant higgsino couplings are suppressed by $\mathcal{O}(\alpha/\alpha_s)$ compared to the case of hard bottom Yukawa couplings. The exact calculation of $C_7^{\tilde{\chi}^+ \tilde{\chi}^0}$ involves two loops and is beyond the aim of this letter. In our numerical examples below, we take the known MSSM expressions for $C_7^{\tilde{\chi}^+ \tilde{\chi}^0}$ from Ref. [3] and suppress the dominant chirality flipping terms by $\alpha/\alpha_s$.

Because the $b$-quark mass is generated radiatively by the gluino loop, the gluino-mediated contribution to $b \to s\gamma$ is expected to be sizable. However, in the absence of new large flavour violation beyond the Cabibbo-Kobayashi-Maskawa (CKM) matrix, the gluino contribution is always subdominant compared with that of the SM or the charged Higgs-boson and chargino ones. Nevertheless, in our scenario, the gluino contribution may become the largest SUSY contribution to $b \to s\gamma$.

To conclude this Section, in the scenario of radiatively induced $(d, s, b)$ quark masses in the MSSM, the $b \to s\gamma$ rate does not impose serious constraints on the charged Higgs-boson, chargino, and neutralino masses. This allows the possibility of their production at colliders but with drastically modified couplings, implying new phenomenology at experiments.

3. Unification framework: Embedding the MSSM with the radiative $(d, s, b)$ quark masses into a unification framework is a well-motivated and appealing possibility. This can naturally happen in SUSY SU(5) GUT because the up and down Yukawa couplings are not unified in this model. This automatically solves all the constraints on the model parameters (see, e.g., Ref. [10]) coming from the prediction of $b \to \tau$ Yukawa unification; they are simply vanishing. In addition,
the constraints from $b \rightarrow s \gamma$ are practically removed, as shown in the previous Section.

Nevertheless the mass spectrum in such a version of the MSSM is stringently constrained by the renormalization-group running of the model parameters, the requirement of radiative electroweak symmetry breaking, and most importantly, the requirement of generating a correct mass to the $b$ quark via Eq.(4). We start the running of gauge and top Yukawa couplings at $m_t$ using the two-loop MSSM renormalization group equations [19]. The bottom and tau Yukawa couplings are taken to be vanishing. We identify the unification scale $M_{GUT} = 2 \cdot 10^{16}$ GeV by the meeting of $g_1$ and $g_2$. At that scale we generate randomly the free parameters of the model: the common gaugino mass $M_{1/2}$, common squark mass $m_0$, common Higgs mass $M_{H_0}$, common $A$ parameter $A_0$, common $A'$ parameter $A'_0$, tan $\beta$, and $\text{sign}(\mu)$ in the following ranges:

$$100 \text{ GeV} < M_{1/2} < 1000 \text{ GeV},$$
$$100 \text{ GeV} < m_0 < 1000 \text{ GeV},$$
$$100 \text{ GeV} < M_{H_0} < 1000 \text{ GeV},$$
$$-2000 \text{ GeV} < A_0 < 1000 \text{ GeV},$$
$$-2000 \text{ GeV} < A'_0 < 1000 \text{ GeV},$$
$$3 < \tan \beta < 60.$$  \hspace{1cm} (14)

Note that tan $\beta$ is now a free parameter and is not constrained by $b - \tau$ Yukawa unification. With these initial values we run the model parameters to the weak scale, assume radiative symmetry-breaking conditions and calculate the sparticle and Higgs-boson mass matrices there. The renormalization-group equations for nonholomorphic terms can be found in [16]. We require that the branching ratio of $b \rightarrow s \gamma$ is in the allowed range (1) and that the radiative $b$-quark mass Eq.(4) is in the range $2.8 < m_b(m_Z) < 3.2$. The scatter plots are almost independent of the sign of the $\mu$ parameter, our results are presented for $\text{sign}(\mu) = -$.

In Fig. 4 we present the scatter plots of the allowed values of the charged Higgs-boson mass $M_{H^+}$ against the lightest chargino mass $M_{\tilde{\chi}^+}$, and against the lightest bottom squark mass $m_{\tilde{b}_1}$. Notice that in this scenario $H^+$ is bounded to be rather heavy, $M_{H^+} \gtrsim 400 \text{ GeV}$. This comes from the requirement of radiatively induced electroweak symmetry breaking. Because the $b$-quark Yukawa coupling is vanishing, the difference between the Higgs-boson mass parameters $M_{H_1}$ and $M_{H_2}$ is maximized. As their difference determines $m_A$, the charged Higgs boson is naturally quite heavy. However, the charginos can be light and be discovered at future colliders.

It is interesting to see the SUSY contribution to $b \rightarrow s \gamma$ in this scenario. In Fig. 5 we plot the total value of $|C_7^{Hf}|$ and the gluino contribution $C_7^g$ (see Eq.(11) for explanation) against the charged Higgs-boson mass $M_{H^+}$. The deviation from the LO SM value $C_7^{SM} = -0.29$ is small and, as follows from the figures, dominated by the gluino contribution. Thus in this scenario, the MSSM mass spectrum is not constrained by $b \rightarrow s \gamma$.

4. **Conclusions.** We have studied the decay $b \rightarrow s \gamma$ in the MSSM in the case the $(d, s, b)$ quark masses are generated radiatively. The soft radiatively generated $b_R$ couplings to the charged Higgs boson and higgsino are suppressed by $1/\tan^2 \beta$ and $\mathcal{O}(\alpha/\alpha_s)$ respectively. The dominant
Fig. 4. Scatter plots of the allowed values of the charged Higgs-boson mass $M_{H^+}$ against the lightest chargino mass $M_{\tilde{\chi}^+_1}$, and against the lightest bottom squark mass $m_{\tilde{b}_1}$.

Fig. 5. Scatter plots of the total $|C_7^{\gamma jj}|$ and the dominant gluino contribution $C_7^{\tilde{g}}$ against the charged Higgs-boson mass $M_{H^+}$.

Contributions of these particles to $b \to s\gamma$ are suppressed by the same factors allowing the existence of light $H^+$ and $\tilde{\chi}^+$. Their production and decay processes at future colliders are changed drastically.
If this scenario is realized in the framework of GUTs, then the constraints from $b - \tau$ Yukawa unification as well as from $b \rightarrow s\gamma$ are removed. Nevertheless, the lightest sparticles in that case are binos and winos.

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