Investigating the light-gluino scenario through $b \to s\gamma$

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We investigate the impact of a light gluino, which might have escaped detection at colliders, on inclusive radiative B decays mediated through penguinlike diagrams. We find that the viability of the scenario depends largely on the magnitude of the flavor-violating c parameter and on the charged Higgs boson mass. Some previously allowed regions of parameter space are now disfavored.

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There has been continued speculation that a light gluino ($\sim$ 2–5 GeV) has escaped detection at the colliders [1]. This assertion had also been fueled by the observation that a light, colored, neutral fermion improves the agreement between low- and high-energy α_s measurements; the light gluino is a strong candidate to satisfy such a requirement. This possibility has been looked into by a number of experiments [2], but it is still very much open, crying out for verification. It is noteworthy that the direct search limits on squark masses from the Collider Detector at Fermilab (CDF) Collaboration [3] are evaded in the presence of a light gluino; the squarks need, in principle, to be heavier than only $M_Z/2$, from nonobservation at the CERN e^+e^- collider LEP. However, it has been pointed out [4] that the precision LEP measurements disfavor squarks below 60 GeV associated with such a light gluino. Of late, a particularly interesting gateway to examine various varieties of new physics, including this speculative light-gluino scenario, has been provided by the inclusive B-decay studies, setting a limit $B(b \to s\gamma) < 5.4 \times 10^{-4}$ at 95% C.L. [5]. It has already been pointed out [6,7] that this rare decay has a strong influence on restricting the parameter space of supersymmetry (SUSY). This motivates us to examine in this paper the present status of the light gluino through this "microscope." SUSY contributions to the rare decay $b \to s\gamma$ have been examined in the literature [8] earlier. On top of these investigations we adopt a timely specialization to the recently reheated issue of a light gluino, following the improved experimental measurement.

The branching ratio of $b \to s\gamma$ is given in units of the semileptonic b-decay branching ratio as

$$\frac{B(b\to s\gamma)}{B(b\to ce\overline{\nu})} = \frac{6\alpha}{\pi\rho\lambda}\left|\frac{K_{tb}K_{ts}^*}{K_{bc}}\right|^2 \times\left[\eta^{16/23}A_\gamma+\tfrac{8}{3}(\eta^{14/23}-\eta^{16/23})A_g+C\right]^2, \quad (1)$$

where $\eta = \alpha_s(M_Z)/\alpha_s(m_b)$, $\rho = (1 - 8r^2 + 8r^6 - r^8 - 24r^4 \ln r)$ with $r = m_c/m_b$, $\lambda = 1 - 1.61\alpha_s(m_b)/\pi$, and C is a coefficient from a complete calculation of the leading-logarithmic QCD corrections; K is the standard Cabibbo-Kobayashi-Maskawa matrix. It may be noted that the m_b^5 dependence in the partial decay widths of the b quark cancels out in Eq. (1). An $O(m_s^2/m_b^2)$ part in the branching ratio is neglected. We take $B(b \to ce\overline{\nu}) = 0.107$. A_γ and A_g are the coefficients of the effective operators for bs-photon and bs-gluon interactions [9] following from

$$\mathcal{L}_{\text{eff}} = \sqrt{\frac{G_F^2}{8\pi^3}}K_{tb}K_{ts}^*\ \bar{s}\ \sigma^{\mu\nu} \times\left[\sqrt{\alpha}A_\gamma F_{\mu\nu}+\sqrt{\alpha_s}A_gT_aG^a_{\mu\nu}\right](m_bP_R+m_sP_L)\ b. \quad (2)$$

The contributions to A_γ and A_g from W bosons, charged Higgs bosons, and gauginos are listed in [7].

The core of the interaction under our investigation is contained in a particular subset of SUSY induced by the quark-squark-gluino Lagrangian. For the sake of making this note self-contained, we extract, in what follows, the essence of the formalism of our earlier work [4, 10]. The quark-squark-gluino Lagrangian is given by

$$\mathcal{L}_{q\tilde{q}\tilde{g}} = i\sqrt{2}\,g_s\tilde{q}_i^{\dagger a}\overline{\tilde{g}}_\alpha(\lambda_\alpha/2)_{ab} \times\left[\Gamma_L^{ip}\frac{1-\gamma_5}{2}+\Gamma_R^{ip}\frac{1+\gamma_5}{2}\right]q_p^b, \quad (3)$$

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where, for three generations of quarks $p = 1-3$, $i = 1-6$ (for each quark flavor there are two squark states), the color indices $a, b = 1-3$ and $\alpha = 1-8$. The (6×3) matrices Γ_L and Γ_R are determined by the quark and squark mass matrices shown below.

Flavor violation stems from the fact that the quark and squark mass matrices are not diagonal in the same basis. The (6×6) $\tilde{d}$ mass squared matrix (in a basis in which the d-quark mass matrix is diagonal) is

$$M_{\tilde{d}}^2 = \begin{pmatrix} m_{0L}^2 I + \hat{M}_d^2 + cK^\dagger \hat{M}_u^2 K & Am_{3/2}\hat{M}_d \\ Am_{3/2}\hat{M}_d & m_{0R}^2 I + \hat{M}_d^2 \end{pmatrix}, \tag{4}$$

where m_{0L} and m_{0R} are flavor-blind supersymmetry-breaking parameters for the left- and right-type squarks, respectively. (For the sake of simplification we have taken $m_{0L} = m_{0R} = m_0$ for numerical purposes, which does not materially affect the conclusion of the paper.) Here, $\hat{M}_u$ and $\hat{M}_d$ are diagonal up- and down-quark mass matrices, respectively. The c term corresponds to a quantum mass correction for a d-type left squark driven by Higgsino exchange. It may be noted that c is the most crucial parameter, originating from an electroweak one-loop effect, which triggers flavor-violating interactions such as $b \to s\gamma$. In specific models c can be estimated by the renormalization group (RG) equations of the quark and squark mass parameters. In our analysis c is a phenomenological input. The off-diagonal block in Eq. (4) corresponds to left-right squark mixings and is proportional to the d-type quark mass matrix. Γ_L and Γ_R in Eq. (3) are

$$\Gamma_L = \tilde{U}^\dagger \begin{pmatrix} I \\ 0 \end{pmatrix}, \quad \Gamma_R = \tilde{U}^\dagger \begin{pmatrix} 0 \\ I \end{pmatrix}; \tag{5}$$

$\tilde{U}$ is the matrix that diagonalizes $M_{\tilde{d}}^2$; $m_{3/2}$ stands for the gravitino mass, and I is the (3×3) identity matrix. It should be mentioned that although the above mass matrix is of the texture that follows from $N = 1$ supergravity, a mild extension of the minimal supersymmetric standard model (MSSM) keeps the general structure unaltered.

When the c-induced SUSY interaction is turned on, A_γ and A_g in Eq. (2) pick up terms in addition to those given in [7]. Their modified expressions, denoted by A'_γ and A'_g, respectively, are given by

$$\begin{aligned} A'_\gamma &= A_\gamma + \frac{4}{9}\frac{\alpha_s(M_Z)}{\alpha}\sin^2\theta_W M_W^2\ S_\gamma, \\ A'_g &= A_g + \frac{\alpha_s(M_Z)}{6\alpha}\sin^2\theta_W M_W^2\ S_g. \end{aligned} \tag{6}$$

Although we compute with the complete set of parameters we present in the following the expressions of S_γ and S_g in the simplified case when $A = 0$:

$$S_\gamma = C_{11} + C_{21} \tag{7}$$

and

$$S_g = (C_{11} + C_{21}) + 9(\tilde{C}_{11} + \tilde{C}_{21}), \tag{8}$$

where the C and $\tilde{C}$ functions are the three-point integrals [11], the arguments of which are the three external and the three internal masses of the relevant penguins. Generically, the C functions correspond to the case when a photon (or a gluon) couples to the internal squark lines in the penguin diagrams, while the $\tilde{C}$ functions refer to the situation when a gluon is emitted from an internal gluino line. The C and $\tilde{C}$ functions in Eqs. (7) and (8) represent their final forms after the super-Glashow-Iliopoulos-Maiani (GIM) subtraction [generically, $C \equiv C(m_{\tilde{b}}^2) - C(m_{\tilde{d}}^2)$ and $\tilde{C} \equiv \tilde{C}(m_{\tilde{b}}^2) - \tilde{C}(m_{\tilde{d}}^2)$]. Both C and $\tilde{C}$ are proportional to cm_t^2, the mass splitting between $\tilde{b}_L$ and any of the remaining d-type squarks, controlling the rate of flavor violation. (In the actual calculation, the GIM subtraction is done numerically.) To evaluate the three-point functions we use the code developed in [12] and employed subsequently in [4,10]. We also cross-check our calculation by performing a systematic expansion in powers of the ratios of the masses of the light and heavy particles. The approximate expressions of S_γ and S_g used in Eq. (6), which agree within 1% with those in Eqs. (7) and (8), are shown below ($x = m_{\tilde{g}}^2/m_0^2$, where $m_{\tilde{g}}$ is the mass of the gluino):

$$S_\gamma = \frac{cm_t^2}{6m_0^4}\Big[(x-1)^{-4}(1-8x-17x^2) + 6(x-1)^{-5}x^2(x+3)\ln x\Big] \tag{9}$$

and

$$S_g = \frac{cm_t^2}{6m_0^4}\Big[(x-1)^{-4}(x^2+172x+19) + 6x(x-1)^{-5}(x^2-15x-18)\ln x\Big]. \tag{10}$$

Before we discuss our results, a few comments are in order. As mentioned at the outset, one of the major motivations for the resurrection of the light gluino scenario is that the running of α_s has a pronounced dependence on the presence of a light color octet fermion, which affects Eq. (1) sensitively. For the coefficient C, determined by the leading log QCD corrections, we use [13]

$$C = \sum_{i=1}^{8} h_i \eta^{a_i}, \tag{11}$$

where

$$a_i = \tfrac{14}{23}, \tfrac{16}{23}, \tfrac{6}{23}, -\tfrac{12}{23}, 0.4086, -0.4230, -0.8994, 0.1456, \tag{12}$$

$$h_i = \tfrac{626126}{272277}, -\tfrac{56281}{51730}, -\tfrac{3}{7}, -\tfrac{1}{14}, -0.6494, -0.0380, -0.0186, -0.0057\ . \tag{13}$$

In particular, we find that the usual W-exchange contribution to the branching ratio in Eq. (1) with $\eta = \alpha_s(M_Z)/\alpha_s(m_b) = 0.681$ [14] is reduced by 27% compared to the SM case where $\eta = 0.548$.

The results of our analysis are presented in Figs. 1 and 2. To appreciate the effect of the light gluino in the context of the full theory of SUSY, we have included the contributions of the charged Higgs bosons and the gauginos. Results are presented for three different values of the parameter c. The broken line corresponds to choosing $c = 0$, i.e., no contribution from the gluino sector at all. The SM contribution for $m_t = 180$ GeV is also shown as the dotted line (the m_t dependence of the branching ratio is rather mild). $m_{\tilde{g}}$ is set to 3 GeV in our analysis.

Since $c < 0$ is preferred in the MSSM, the lightest of the $\tilde{d}$-type squarks, dominantly $\tilde{b}_L$, has a mass $\simeq \sqrt{m_0^2 + cm_t^2}$ (for $A = 0$). Thus for a given choice of m_0 and for a fixed m_t, the maximum magnitude of c is restricted by the LEP bound $\sqrt{m_0^2 + cm_t^2} \geq 45$ GeV. For $m_t = 180$ GeV and $m_0 = 60$ GeV, this requires $|c| \leq 0.05$. In Fig. 1(a) we present the results for $m_{H^+} = 100$ GeV and the Higgs boson mixing parameter $\mu = 90$ GeV. This choice is motivated from the fact that the LEP constraint on the chargino masses when imposed in the light gluino scenario restricts $\tan\beta \sim 1$ and the maximum allowed value of μ to 90 GeV. To illustrate the variation with respect to μ we display in Fig. 1(b) the results for $\mu = 0$, although this corresponds to an unrealistic limit where a massless axion is present at the tree level. In Figs. 2(a) and 2(b) we exhibit the findings for m_{H^+} chosen to be 400 GeV.

It is seen from Fig. 1(a) that choosing $c = -0.05$, or even $c = -0.03$, the squark-gluino contribution dominates over the rest for $m_0 < 100$ GeV. The figure corresponds to the situation when there is no left-right squark mixing, i.e., $A = 0$. Under these circumstances, if one

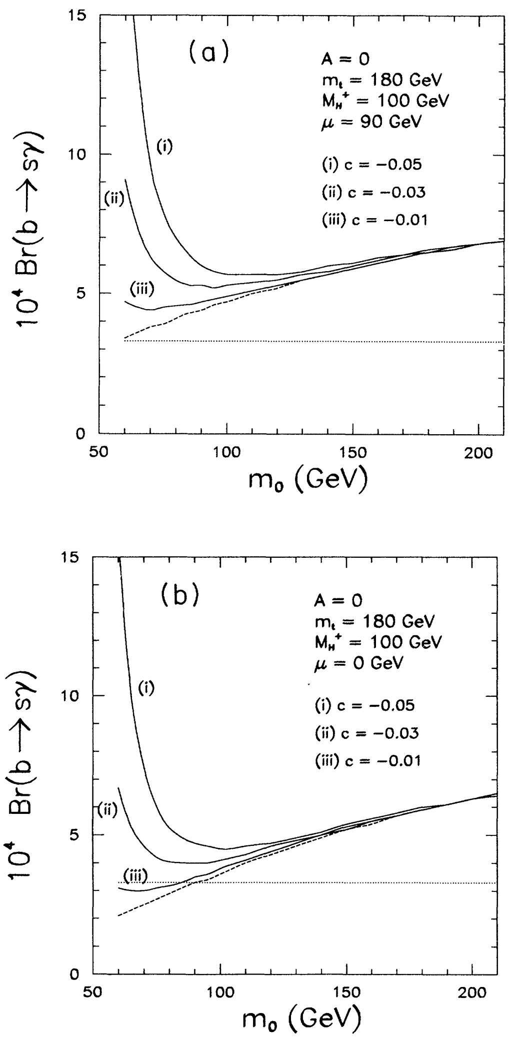

FIG. 1. The branching ratio for the process $b \rightarrow s\gamma$ as a function of the average squark mass (m_0) for $m_{H^+} = 100$ GeV and different values of the flavor-violation parameter c (solid lines). Also shown are the branching ratio with no contribution from the gluino sector (broken line) and from the standard model alone (dotted line). For (a) $\mu = 90$ GeV and for (b) $\mu = 0$ GeV.

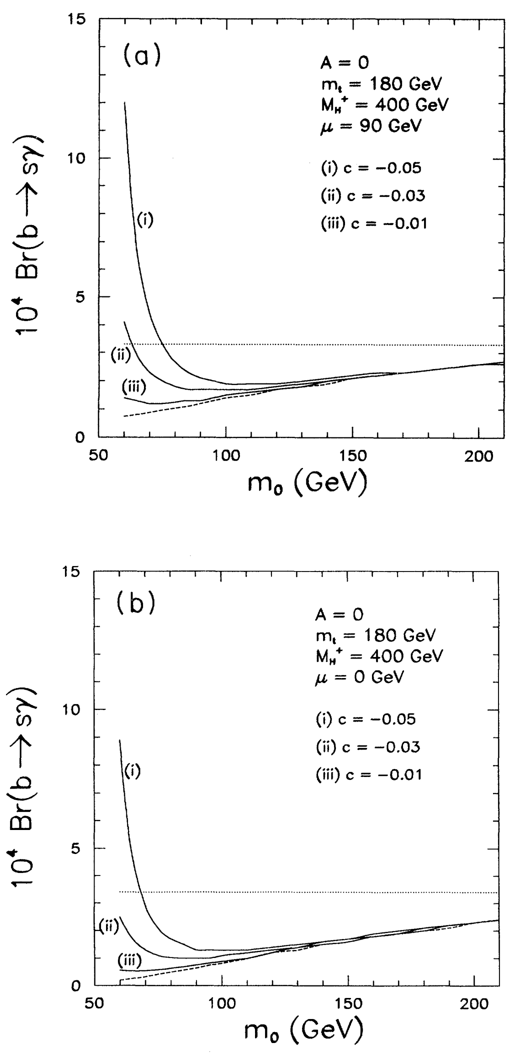

FIG. 2. Same as in Fig. 1 except that $m_{H^+} = 400$ GeV.

uses the CLEO bound $B(b \to s\gamma) < 5.4 \times 10^{-4}$, the light gluino is completely disfavored, no matter what the squark mass is. It ought to be stressed that such a choice of c is in good consonance with the predictions from the RG evolution of the squark masses [15]. The case for $c = -0.01$ is close to the present experimental bound. If μ is smaller [Fig. 1(b)], the overall results turn out to be smaller and only for $c = -0.05$ can the experimental bound put some restrictions. As seen from Figs. 2(a) and 2(b), a larger choice of m_{H^+} (400 GeV in this case) diminishes the branching ratio significantly enough to practically remove all constraints. Other parameters remaining the same, choosing $A = 3$ decreases the effect very slightly, at most by $\sim 2\%$, which justifies our choice of setting $A = 0$ in all the figures for the sake of simplicity. Under this situation, i.e., in the absence of left-right squark mixings, we have also checked that for Figs. 1(a) and 2(a), changing the sign of μ does not alter the result. If one deviates from the MSSM and assumes a positive value for c, the gluino-induced effect becomes less prominent as a result of its destructive interference with the other sectors, and no significant bound could be set at all. It may be noted that varying $m_{\tilde{g}}$ in the range $(1 - 5)$ GeV has no numerical impact within the scale of the figure.

In conclusion, we have studied the process $b \to s\gamma$ in the context of the light gluino scenario. Some previously allowed regions of parameter space are now disfavored. For example, for $c = -0.05$, $m_t = 180$ GeV, and $M_{H^+} =$ 100 GeV, the light-gluino window is virtually closed for arbitrary choices of the squark masses. Needless to say, the sign and the magnitude of c, for which there is a significant freedom, as well as the mass of the charged Higgs boson m_{H^+} have crucial roles to play in drawing such a conclusion.

The consequence of a light gluino in the MSSM, in the context of unification of gauge and Yukawa couplings, has been shown [16] to pose a very tight restriction on the allowed values of $\alpha_s(M_Z)$, keeping it consistent, nevertheless, with the prediction at LEP. Additionally, if one demands the radiative breakdown of electroweak symmetry in the MSSM (irrespective of the criterion of unification), a light gluino is difficult to be accommodated [17]. This analysis, which probes a rather direct contribution of a light gluino, concludes that the window is still open, albeit with a smaller region of allowed parameter space. Further investigation and more accurate experimental measurements are therefore called for before any final verdict can be drawn on this issue.

Note added. Very recently CLEO has announced their first measurement [18] of the inclusive $b \to s\gamma$ decay with a branching ratio $(2.32 \pm 0.51 \pm 0.29 \pm 0.32) \times 10^{-4}$ setting a 95% upper (lower) limit of $3.66(0.98) \times 10^{-4}$ to the branching ratio. As a result of the reduction of the upper limit of the branching ratio from the existing bound of 5.4×10^{-4}, the parameter space of c and m_0 gets slightly more constrained than before. For example, in Fig. 1(a) even a choice of $c = -0.01$ cannot accommodate a light gluino, while in Fig. 1(b) a choice of $c = -0.03$ finds it difficult for given values of other parameters relevant for those figures. On the other hand, attention should be paid on theoretical uncertainties arising, for example, from the scale dependence of QCD corrections [a choice of $\alpha_s(m_b/2)$ to $\alpha_s(2m_b)$ in the definition of η in Eq. (1) changes the SM prediction of the decay rate from $\sim 4 \times 10^{-4}$ to $\sim 2 \times 10^{-4}$ for $m_t = 180$ GeV [19]] or from next-to-leading-order QCD corrections. These uncertainties can easily plague the above conclusions. Of late, a part of the next-to-leading-order QCD corrections has been performed with a consequence of reducing the QCD enhancement in the SM by $\sim 15\%$ [20]. Our conclusion still remains that the observed decay rate of inclusive $b \to s\gamma$ has to experience an improved statistics and reduced systematic errors and the theoretical uncertainties have to be resolved further, before anything concrete can be concluded in this issue.

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