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## Implications for Supersymmetric Dark Matter Detection from Radiative $b$ Decays

Francesca M. Borzumati  
*II. Institut für Theoretische Physik  
Universität Hamburg, 22761 Hamburg, Germany*

Manuel Drees\*  
*Physics Department, University of Wisconsin, Madison, WI 53706, USA*

Mihoko M. Nojiri†  
*Theory Group, KEK, Oho 1-1, Tsukuba, Ibaraki 305, Japan*

### Abstract

We point out that combinations of parameters that predict large counting rates in experiments searching for supersymmetric dark matter often tend to predict a very large branching ratio for the inclusive decay  $b \rightarrow s\gamma$ . The recent experimental upper bound on this branching ratio, therefore, indicates that searches for supersymmetric dark matter might be even more difficult than previously anticipated.



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\*Heisenberg Fellow

†E-mail: NOJIRIN@JPNKEKVX

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The Lightest Supersymmetric Particle (LSP) is one of the most attractive particle physics candidates for the missing dark matter (DM) [1] in the Universe. In the simplest potentially realistic supersymmetric theory, the Minimal Supersymmetric Standard Model (MSSM) [2], the LSP is stable by virtue of a symmetry, the so-called  $R$ -parity. Calculations [3] have shown that if the LSP is the lightest of the four neutralino states present in this model, the relic density of LSPs left over from the Big Bang is in the desired range over a wide region of the supersymmetric parameter space. Very broadly this range can be defined by

$$0.025 \leq \Omega_{\text{LSP}} h^2 \leq 1, \quad (1)$$

where  $\Omega_{\text{LSP}}$  is the relic density in units of the closure density, and  $h$  is the Hubble constant in units of 100 km/(sec Mpc). Observations imply  $0.5 \leq h \leq 1$ , the lower range, perhaps, being favoured. The lower bound in (1) then follows from the requirement that there be enough relic LSPs to form the dark matter haloes of galaxies ( $\Omega_{\text{LSP}} \geq 0.1$ ). The upper bound is equivalent to the constraint that the Universe be at least 10 billion years old.

Unfortunately relic neutralinos are rather difficult to detect experimentally. Here we are interested in direct detection experiments [4], where one searches for the elastic scattering of an LSP off a target nucleus. The signal, provided by the energy deposited in the detector by the recoiling nucleus, has a rate proportional to the LSP–nucleus scattering cross section. Partly because of the Majorana nature of the LSP, this cross section is often quite small. It can be generally split into two parts [5], one due to spin–spin interactions and the other to scalar (spin-independent) interactions. For heavy target nuclei the spin-independent contribution usually dominates the spin dependent one [6], since it is enhanced by the square of the number of nucleons in the nucleus in question. This spin-independent interaction gets contributions from the exchange of the two neutral scalar Higgs bosons of the MSSM as well as from squark exchange. Unless squarks are quite close in mass to the LSP, the Higgs–exchange contribution usually dominates. We refer the reader to refs. [6, 7] for more details on LSP–nucleus interactions.

Thus, the LSP–nucleus scattering cross section depends, in general, on many parameters: the gaugino mass  $M_2$ , the higgsino mass  $\mu$  and ratio of Higgs vacuum expectation values  $\tan\beta$  entering the neutralino mass matrix\* [2]; the squark masses and mixings; and the masses and couplings of the Higgs bosons. At the tree level the Higgs sector of the MSSM [8] is completely specified in terms of two parameters, which we take to be  $\tan\beta$  and the mass  $m_P$  of the pseudoscalar Higgs boson. As it is well known, radiative corrections [9] to the mass of scalar Higgs bosons introduce also a dependence on the mass of the top–quark,  $m_t$ , as well as on the parameters describing the mass matrix for the scalar superpartners of the top–quark, or stop  $\tilde{t}$  (see below). In our analysis we include these corrections using the effective potential method [10]†. As for the slepton masses, needed for the calculation of the LSP relic density, we follow the conventional choice made in DM searches: we assume that the squared masses of all sfermions get the same soft supersymmetry breaking contribution  $m^2$  along the diagonal of their respective mass matrices. Our

\*We assume the usual unification relation between the  $U(1)$  gaugino mass  $M_1$  and the  $SU(2)$  gaugino mass  $M_2$ ,  $M_1 = (5/3)\tan^2\theta_W M_2 \simeq 0.5M_2$ .

†Since not only corrections growing as  $\log(m_t/m_c)$  are included, it is technically easier to present our results for fixed  $m_P$ , rather than for fixed mass of one of the scalar Higgs bosons.

main result is independent of this assumption. Finally, the specification of the neutralino mass matrix, due to gauge invariance, completely determines also the chargino sector.

Having fixed the (s)particle spectrum it is imperative to first check for consistency with experimental and theoretical constraints before we use this spectrum to predict LSP detection rates. In particular,  $M_2$ ,  $\mu$  and  $\tan\beta$  must be chosen such that charginos and neutralinos escape detection at LEP [11]. Similarly, searches for neutral Higgs bosons at LEP [11] constrain the parameters of the Higgs sector.

There is yet another constraint which has so far been ignored in estimates of LSP detection rates. The CLEO II collaboration has given [12] a 95% c.l. upper bound of  $5.4 \cdot 10^{-4}$  on the branching ratio for inclusive  $b \rightarrow s\gamma$  decays. Moreover, the observed exclusive channel  $B \rightarrow K^*\gamma$  translates into a (model dependent) lower bound on the inclusive decay of order  $10^{-4}$ . These bounds are relevant for LSP searches since within the MSSM the  $Br(b \rightarrow s\gamma)$  is determined [13] by the *same* parameters that determine LSP detection rates, i.e. the masses and mixings of squarks and charginos as well as the mass of the charged Higgs boson,  $m_{H^\pm}$ . This is related to  $m_P$  by†:

$$m_{H^\pm}^2 = m_P^2 + m_W^2, \quad (2)$$

where  $m_W \simeq 80$  GeV is the mass of the  $W$  bosons. In particular, a light charged Higgs gives a large positive contribution to the amplitude  $\mathcal{A}(b \rightarrow s\gamma)$ . Loops involving charginos (or, in general, gluinos) and squarks, in contrast, give contributions with either sign and decouple in the limit of large sparticle masses. Therefore, spectra of supersymmetric particles with rather light higgses when sparticles are taken to be heavy, although well suited for DM detection, tend to give results for the  $Br(b \rightarrow s\gamma)$  similar to the ones obtained in the two–higgs–doublet model (type II) [14, 15]. Thus, one may expect clashes with the experimental upper bound on this decay in regions of parameter space where the counting rates are at the highest values.

In order to quantify this statement we have to specify the amount of flavor mixing in the squark sector through which transitions from the third to the second generation of quarks, such as the decay  $b \rightarrow s\gamma$  can occur. As mentioned, mimicking as closely as possible the treatment of squark masses in previous treatments of LSP detection [6, 16], we assume that all sfermions have the same soft supersymmetry breaking mass. This implies that *no* contributions to the decay  $b \rightarrow s\gamma$  can come from loops mediated by neutral gauginos, gluinos or neutralinos. Flavor mixing in the quark sector, however, will introduce some mixing in the squark sector as well. Following ref. [13], we work in a quark basis in which current and mass eigenstates coincide for right handed quarks as well as left-handed down-type quarks. Flavor mixing, therefore, shows up only in the left–left sector of the  $6 \times 6$  mass matrix for  $u$ -type squarks,  $\tilde{u}$ :

$$\mathcal{M}_{\tilde{u}}^2 = \begin{pmatrix} \mathcal{M}_{\tilde{u}_L}^2 & \mathcal{M}_{\tilde{u}_{LR}}^2 \\ (\mathcal{M}_{\tilde{u}_{LR}}^2)^\dagger & \mathcal{M}_{\tilde{u}_R}^2 \end{pmatrix}. \quad (3)$$

†Eq.(2) holds at tree level. Radiative corrections to this relation are very small [10] unless one somewhat artificially allows  $\tan\beta < 1$ .

The  $3 \times 3$  left-left, right-right and left-right mixing submatrices  $\mathcal{M}_{\tilde{u}_L}^2$ ,  $\mathcal{M}_{\tilde{u}_R}^2$  and  $\mathcal{M}_{\tilde{u}_R}^2$  are given by:

$$\left(\mathcal{M}_{\tilde{u}_L}^2\right)_{ij} = (m^2 + 0.35m_Z^2 \cos 2\beta) \delta_{ij} + m_t^2 V_{3i}^* V_{3j}; \quad (4a)$$

$$\left(\mathcal{M}_{\tilde{u}_R}^2\right)_{ij} = (m^2 + 0.15m_Z^2 \cos 2\beta) \delta_{ij} + m_t^2 \delta_{i3} \delta_{j3}; \quad (4b)$$

$$\left(\mathcal{M}_{\tilde{u}_R}^2\right)_{ij} = -(A_t + \mu \cot \beta) m_t V_{3i}^* \delta_{j3}. \quad (4c)$$

The symbols  $V_{ij}$  indicate here elements of the CKM mixing matrix,  $\mu$  is the mass parameter entering the neutralino mass matrix, and  $A_t$  is a soft supersymmetry breaking parameter of order  $m$ . When writing eqs.(4) we have neglected all Yukawa couplings except for the top quark. Similarly, the left-right mixing in (4c) is significant only for the third generation of squarks.

We are now in a position to discuss quantitatively the correlation between the relic LSP detection rate and the  $Br(b \rightarrow s\gamma)$ . For definiteness we focus on a detector consisting of isotopically pure  $^{76}\text{Ge}$  since such a device is now under construction. The next round of experiments is expected to reach a sensitivity of about 0.1 events/(kg-day) which improves on the current best limits [17] by about a factor of 100.

We show in figs.1 the LSP counting rate in such a detector (solid lines) as well as the branching ratio for  $b \rightarrow s\gamma$  (dashed lines) as function of various parameters of the MSSM and for fixed top-quark mass,  $m_t = 175$  GeV. We give results for the case of a heavy LSP ( $M_2 = 500$  GeV and  $\mu = 400$  GeV, giving  $m_{\text{LSP}} \simeq 200$  GeV) and the case of a much lighter one ( $M_2 = 100$  GeV,  $\mu = -100$  GeV, giving  $m_{\text{LSP}} \simeq 50$  GeV). We fix the remaining supersymmetric parameters to be, in general,  $A_t = 0$ ,  $\tan\beta = 2$ ,  $m_P = 100$  GeV, and choose  $m$  to be respectively  $m = 500$  GeV and  $m = 200$  GeV in the case of the heavy and light LSP. We then deviate from these points in the supersymmetric parameter space by varying  $m_P$  (fig. 1a),  $\tan\beta$  (1b),  $m$  (1c) or  $A_t$  (1d).

We have chosen  $\mu > M_1$  so that the LSP is predominantly a gaugino; this is necessary [3] to satisfy the lower bound on  $\Omega_{\text{LSP}} h^2$  in (1). In both cases, however, the higgsino components of the LSP are still quite substantial, leading to sizable couplings of the LSP to Higgs bosons. The LSP detection rate can therefore be quite large *if* Higgs bosons are light.

This is illustrated in fig. 1a, where the dependence on  $m_P$  is shown. If  $M_2$ ,  $m$  and  $\mu$  are large (upper pair of curves),  $m_P$  can be chosen as small as 50 GeV without violating direct Higgs search limits. This would lead to an LSP counting rate of 0.2 events/(kg-day), in principle observable in the next round of direct detection experiments. Nevertheless, demanding that  $Br(b \rightarrow s\gamma)$  be  $\leq 5.4 \cdot 10^{-4}$  implies  $m_P \geq 170$  GeV and a counting rate of less than 0.02 events/(kg-day), a value too small to be detectable in the near future. Notice, however, that no error in the theoretical determination of  $Br(b \rightarrow s\gamma)$  has been assumed, as yet. This discussion is postponed to a later point of this paper.

For our second choice of parameters (lower pair of curves) the lower bound on  $m_P$  is set by Higgs searches. In this case, which is characterized by rather small squark and

chargino masses, there is a sizable contribution to the decay  $b \rightarrow s\gamma$  from chargino-squark loops which interferes *destructively* with the contributions from  $W$  and  $H^\pm$  loops. The Higgs contributions decouple in the limit of large  $m_P$ . As a result this scenario gives for  $m_P > 300$  GeV, values of  $Br(b \rightarrow s\gamma)$  *below* the Standard Model (SM) prediction of  $3.2 \cdot 10^{-4}$  [18]. It should be noted also that for fixed  $m_P$  this scenario gives a significantly smaller counting rate than the case with heavy LSP, even though a lighter LSP means a larger LSP flux (the mass density of relic neutralinos in the vicinity of the solar system is assumed to be fixed) and less suppression due to nuclear form factors [16]. The reason is that negative values of  $\mu$  (in the convention of ref. [8], which we follow throughout) always imply less gaugino-higgsino mixing in the neutralino sector, and hence smaller couplings of the LSP to Higgs bosons.

In fig. 1b we show the  $\tan\beta$  dependence for our two choices of the LSP mass. Here we have increased  $m_P$  to 150 GeV, and chosen  $\mu = -300$  GeV in the light LSP case, in order to ensure  $\Omega_{\text{LSP}} h^2 \geq 0.025$  for a sizable range of  $\tan\beta$ . We have terminated the curve for  $M_2 = 100$  GeV at  $\tan\beta = 33$  since larger values give a too small LSP relic density. The lower bound on  $\tan\beta$  is given by the requirement that the Higgs boson escapes detection at LEP [11]. We see that the  $\tan\beta$  dependence of the LSP counting rate is quite similar in both cases. At first, the rate decreases with increasing  $\tan\beta$  since the mass  $m_{h^0}$  of the light neutral Higgs  $h^0$  increases. This mass, however, remains essentially constant for  $\tan\beta \geq 5$ . At the same time, the coupling of the heavier neutral Higgs  $H^0$  to down-type quarks increases roughly as  $\tan\beta$ . In the case of the light LSP, the coupling of the LSP to  $H^0$  is less suppressed than the coupling to  $h^0$ . In addition, in this scenario, squark exchange contributions are not entirely negligible (recall that it is for  $m = 200$  GeV). As a result, the counting rate grows faster with  $\tan\beta$  in the case of the light LSP than in the case of the heavy one.

The difference between the two cases is much more dramatic for the  $Br(b \rightarrow s\gamma)$ . The contribution from  $H^\pm$  loops becomes independent of  $\tan\beta$  roughly for  $\tan\beta > 3$ . In the scenario with heavy LSP and even heavier charginos and squarks, the contribution from sparticle loops is always small compared to the SM and Higgs contributions, leading to an overall very weak dependence of the branching ratio on  $\tan\beta$ . Notice that the entire curve lies above the experimental upper bound of  $5.4 \cdot 10^{-4}$ . In contrast, in the case with a light LSP and comparatively light charginos and squarks, sparticle loops do contribute significantly. The observed strong  $\tan\beta$  dependence of this contribution is due to the fact that the chargino- $b$ - $\bar{t}$  interaction contains [8, 13] a term proportional to the bottom Yukawa coupling, which grows as  $\tan\beta$  for large  $\tan\beta$ . For the given values of  $M_2$  and  $m$  but positive  $\mu$  the contribution from this term would be positive, leading to predictions for  $Br(b \rightarrow s\gamma)$  rapidly growing and quickly exceeding the experimental upper bound. For the given case of negative  $\mu$  this contribution interferes destructively with the  $W$  and  $H^\pm$  loops, and one eventually gets into conflict with the experimental lower bound. As already mentioned earlier, it is difficult to translate the measured [12]  $Br(B \rightarrow K^* \gamma) = 4.5 \cdot 10^{-5}$  into a stringent lower bound on  $Br(b \rightarrow s\gamma)$ . The very conservative requirement of a lower bound of  $10^{-4}$  for the inclusive branching ratio gives the upper limit  $\tan\beta < 25$  and hence a counting rate of less than 0.23 events/(kg-day) in the case with light LSP.

The dependence on the parameter  $m$  is shown in fig. 1c. The LSP counting rate falls

monotonically with increasing  $m$  since, due to the above mentioned radiative corrections to the Higgs sector [9, 10],  $m_{h^0}$  increases with  $m$ . In the light LSP scenario the lower bound on  $m$  comes from the Higgs search limits, whereas in the heavy LSP case this bound is set by the requirement that the lightest squark (essentially a  $\tilde{t}$ -squark) be heavier than the lightest neutralino. A charged LSP, in fact, would result in too large an abundance of exotic isotopes [19] such as the one with a squark bound to a nucleus. At the lower end of  $m$  squark exchange contributions to LSP-nucleus scattering are quite important, including the  $\mathcal{O}(m_{\tilde{t}}^{-4})$  terms discussed in ref.[7]; this explains the rapid decrease of the expected counting rate with increasing  $m$  in this region.

In the two cases of light and heavy LSP, the  $Br(b \rightarrow s\gamma)$  increases with increasing  $m$ , since in both cases chargino contributions interfere destructively with  $W$  and  $H^\pm$  contributions. The  $m$  dependence is much stronger for the light LSP due to the presence of lighter charginos and hence potentially larger contributions from sparticle loops. In the limit of large squark masses these contributions are always very small and practically independent of the chargino mass, leading to the observed convergence of the two curves at the higher end of  $m$ . Note that the entire curve for the heavy LSP is once again above the experimental upper bound on the branching ratio.

In fig. 1d we show the dependence on the parameter  $A \equiv A_t/m$ . Compared to the previous three figures we observe a rather mild variation of the LSP counting rate. The effect is almost entirely due to the radiative corrections to  $m_{h^0}$ , which depend on  $A_t$  in a nontrivial way [9, 10]. In particular,  $m_{h^0}$  reaches a minimum when the combination  $A_t + \mu \cot\beta$  of eq.(4c) is very small. Increasing it, at first increases  $m_{h^0}$  which then reaches a maximum at some finite value of  $A$ . Increasing the combination  $|A_t + \mu \cot\beta|$  even further, reduces  $m_{h^0}$ , as it is very visibly shown by the curve relative to the heavy LSP case. The upper bound on  $A$ , at which both sets of curves are stopped, is determined by the requirement that the lightest squark be heavier than the lightest neutralino. The observed variation of the counting rate then follows from the fact that the  $h^0$ -exchange contribution to the LSP-nucleus cross section scales like the *inverse* of  $m_{h^0}^4$ .

In contrast, the  $Br(b \rightarrow s\gamma)$  increases or decreases monotonically with increasing  $A$ . Once again only the contribution from chargino-squark loops changes when  $A$  is varied. The absolute value of this contribution reaches a minimum at small values of  $|A_t + \mu \cot\beta|$  where the lightest squark mass is maximal (recall that off-diagonal entries in the squark mass matrix tend to reduce the smallest eigenvalue and increase the largest one). We observe that the sign of this contribution is flipped when going from positive to negative  $A$ , since the sign of the left-right mixing terms changes for the set of supersymmetric parameters chosen here. The sign of this contribution depends also on the sign of  $\mu$ , and its absolute size is larger for the case with light LSP. Hence, the slope of the curve for the light LSP scenario, for which we have taken  $\mu < 0$ , is opposite in sign and larger in magnitude than the slope for the curve relative to the heavy LSP and positive  $\mu$ . Notice that this latter scenario satisfies the upper bound on the branching ratio if  $A$  is very close to its maximal allowed value, while in the case of a light LSP this bound is violated for  $A > 0.4$ .

Finally, the dependence on  $M_2$  and  $\mu$  is shown in fig. 2, where we have fixed  $A_t = 0$ ,  $\tan\beta = 2$ ,  $m_p = 100$  GeV and  $m_t = 175$  GeV and we have chosen  $m = \min(120 \text{ GeV}, 2m_{\text{LSP}})$ .

The shaded regions in the three frames are excluded by LEP searches for charginos, neutralinos and Higgs bosons [11] (region of small  $|\mu|$  or small  $M_2$ ), or by the requirement that the lightest squark (which again is mostly a  $\tilde{t}$ -squark) is heavier than the lightest neutralino (regions of large  $|\mu|$  and small or moderate  $M_2$ ). The dashed lines in the three frames indicate contours of constant  $Br(b \rightarrow s\gamma) = 5.4 \cdot 10^{-4}$ : larger values of  $Br(b \rightarrow s\gamma)$  are obtained below these lines and smaller above. We also show contours of constant LSP counting rate = 1, 0.1, and 0.01 events/(kg-day) in the upper, middle and lower frame, respectively (solid lines). The dotted lines are contours of constant  $\Omega_{\text{LSP}} h^2 = 0.025$ . In the regions of small  $|\mu|$  and large  $M_2$  the LSP is higgsino-like and its relic density is too small to be of cosmological interest [3].

We observe the well-known [16] correlation between large LSP counting rate and small LSP relic density; in particular, the region where the counting rate exceeds 1 events/(kg-day) (for fixed neutralino flux!) lies completely within the region with  $\Omega_{\text{LSP}} h^2 < 0.025$  (upper frame). It has been suggested in the literature to re-scale the counting rate in such regions in order to take into account the reduced LSP flux. We prefer to discard these regions altogether, since here the LSP can only make an almost negligible contribution to the solution of the DM problem<sup>§</sup>.

Prospects for direct LSP detection become even less promising once we require the  $Br(b \rightarrow s\gamma)$  to be below its experimental upper bound. Only the little region at small  $M_2$  and  $\mu \simeq 450$  GeV survives for positive  $\mu$ , while for  $\mu < 0$  the somewhat larger region to the right and below the long-dashed curve remains acceptable. In particular, the whole region where the counting rate exceeds 0.1 events/(kg-day) is now excluded (middle frame). For most of the allowed region with negative  $\mu$  the counting rate is even below 0.01 events/(kg-day) (lower frame). The implementation of the experimental bound on  $Br(b \rightarrow s\gamma)$ , together with the requirement of a cosmologically interesting relic density, implies that, for the values of supersymmetric parameters chosen here, the maximal LSP counting rate in  $^{76}\text{Ge}$  is about 0.04 and 0.1 events/(kg-day) for positive and negative  $\mu$ , respectively. Notice that the maximum occurs for negative  $\mu$  where the LSP counting rate is usually smaller. In this region, in fact, the expected  $Br(b \rightarrow s\gamma)$  gets destructively interfering contributions from chargino-squark loops (at least in the region of small or moderate  $M_2$ ) thereby leaving open wider portions of parameter space.

At this point we should warn the reader that our predictions for both the LSP counting rate and the branching ratio of radiative  $b$  decays are fraught with substantial theoretical uncertainties. The LSP counting rate obviously depends on the local density and velocity distribution of relic neutralinos. In our calculation we have used the standard values [21] of  $0.3 \text{ GeV/cm}^3$  for the LSP mass density and  $320 \text{ km/sec}$  for their velocity dispersion. The calculation of the LSP-nucleus scattering cross section also suffers from uncertainties, the most important one being the value of the nucleonic matrix element  $\langle p | m_s \bar{s} s | p \rangle$ , which we have taken to be  $130 \text{ MeV}$  [22]. Varying this value within the range favoured by model calculations can change the prediction for the LSP counting rate by as much as a factor of 2.

<sup>§</sup>Recall that there is now rather solid evidence [20] that  $\Omega > 0.1$  on bigger than galactic length scales.

The uncertainty in our prediction for  $Br(b \rightarrow s\gamma)$  stems primarily from the fact that the present calculation is in some sense still at the leading order in perturbative QCD. As a result there is a rather strong dependence on the value of the renormalization scale  $Q_0$  that is used in the calculation. The resulting uncertainty has been emphasized by Ali and Greub, and has more recently been elaborated by Buras et al. [18], who have also included uncertainties due to the experimental errors on parameters entering the prediction of this branching ratio in their analysis. Within the SM they find an overall theoretical “1  $\sigma$ ” uncertainty of about  $\pm 25\%$ , which includes the uncertainty that results from varying  $Q_0$  from 2.5 to 10 GeV (i.e. approximately from  $m_b/2$  to  $2m_b$ ). We have (linearly) added an additional 8% uncertainty, which is the size of that part of higher order QCD corrections already known [18]. Although strictly speaking a statistical meaning cannot be assigned to the theoretical uncertainty, nevertheless, for the time being, one can obtain a conservative estimate of the branching ratio by allowing a “1  $\sigma$  downward fluctuation” due to this theoretical uncertainty. We should mention here that the *relative* theoretical uncertainty is often smaller in the MSSM than in the standard model. The reason is that a purely QCD-induced additive contribution to the  $b \rightarrow s\gamma$  matrix element, which contributes greatly to the QCD uncertainty, becomes less important when additional terms are added with the same sign as the  $W$ -loop contribution present in the SM.

Contours where we take the value  $5.4 \cdot 10^{-4}$  to be as this conservative (low) estimate of the  $Br(b \rightarrow s\gamma)$  are shown by the dash-dotted lines in fig. 2. We see that, for the given set of parameters, most of the  $(M_2, \mu)$  plane is allowed if this lower theoretical estimate of the branching ratio is indeed correct. The regions where both  $M_2$  and  $|\mu|$  are large, however, are still excluded for the assumed value of  $m_P$ . Moreover, also for less extreme choices of  $M_2$  and  $\mu$ , there still remain significant constraints on parameter space even if this lower theoretical estimate is used. For example, for the heavy LSP case shown in fig. 1a a bound  $m_P \geq 95$  GeV would result, leading to a reduction of the maximal counting rate by a factor of four. Finally, it should be clear from the above discussion that our quantitative results depend on the ansatz for the squark mass matrices. It is conceivable that *ad hoc* modifications of this simple ansatz would allow to partially circumvent the constraints imposed by the experimental bounds on  $Br(b \rightarrow s\gamma)$ . It remains still true, however, that in general one is likely to get into conflict with these bounds by simultaneously choosing a heavy sparticle spectrum and light Higgs bosons.

In conclusion, we have pointed out that the experimental upper bound [12] on the branching ratio for inclusive  $b \rightarrow s\gamma$  decays imposes somewhat strong constraints on the region of parameter space where sizable counting rates for relic neutralinos are expected. In some cases, the lower bound on the branching ratio is also relevant. In spite of considerable theoretical uncertainties in estimates of both the LSP counting rate and the  $Br(b \rightarrow s\gamma)$ , it seems fair to say that prospects for the next round of LSP detection experiments look much bleaker once the CLEO constraint is incorporated in the analysis. The main reason for this is that both the counting rate and the branching ratio increase with decreasing mass of the Higgs bosons in the theory; the experimental upper limit on the latter hence reduces the maximal possible value of the former.

In this paper we focussed on direct LSP detection experiments. The expected signal rate in experiments looking for LSP annihilation in the center of the Earth or Sun [23]

is also proportional to LSP-nucleus scattering cross sections. The upper limit on the  $Br(b \rightarrow s\gamma)$  therefore will tend to reduce the maximal signal that can be expected in such experiments as well.

Note that we have assumed in our analysis that sparticle masses,  $\mu$ , and Higgs masses can all be varied independently from each other; the same assumption has been made in almost all previous studies of LSP detection. Given that the main motivation for the introduction of weak scale supersymmetry is to help understanding electroweak symmetry breaking, such an assumption appears quite unnatural. One would rather expect these masses to be of roughly the same size. This statement can be quantified in so-called minimal supergravity theories [24], where the mechanism of radiative gauge symmetry breaking ensures that sparticle masses,  $m_P$  and  $\mu$  are strongly correlated. In such models the upper bound on  $Br(b \rightarrow s\gamma)$  is more easily satisfied [13, 25]. The price one has to pay, though, is that expected LSP counting rates are usually very low [7], of order  $10^{-3}$  events/(kg-day) or even less in a  $^{76}\text{Ge}$  detector. The main result of this paper is that now a purely *experimental* constraint seems to force us closer to regions of parameter space favoured by these supergravity models, which are at the same time theoretically very appealing but difficult to probe by LSP detection experiments.

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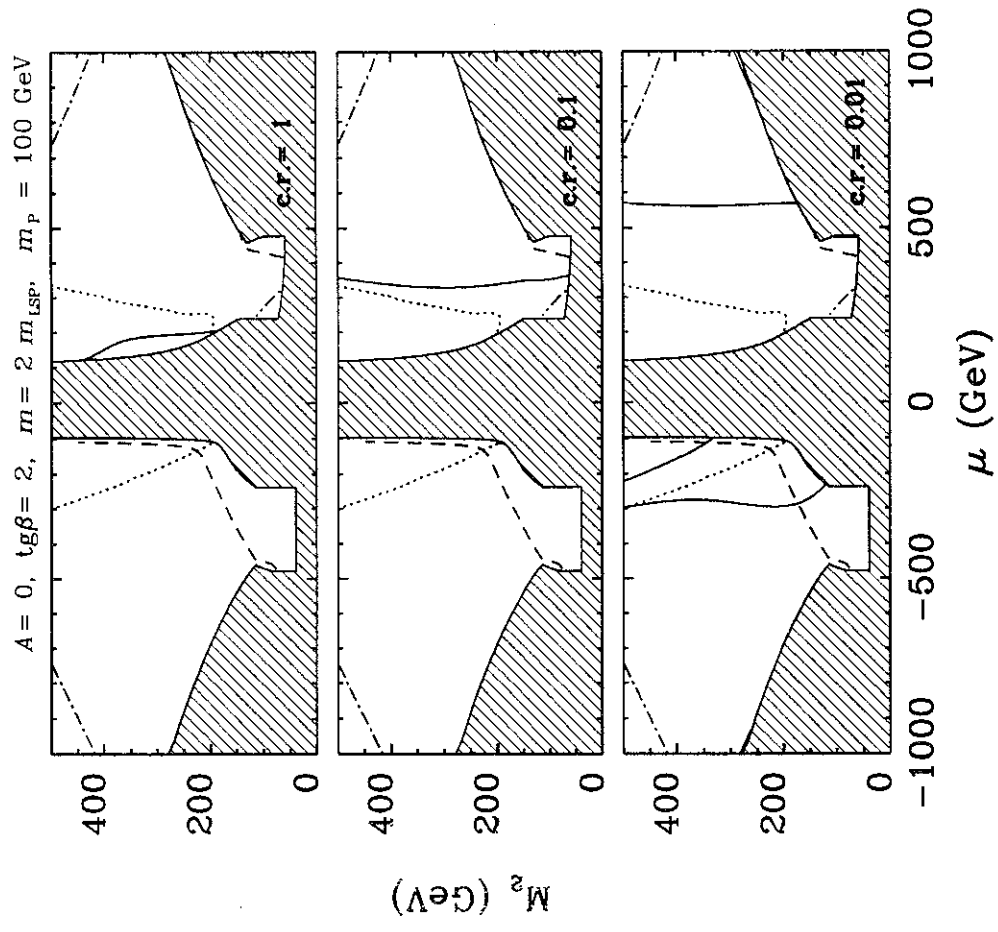


Fig 2

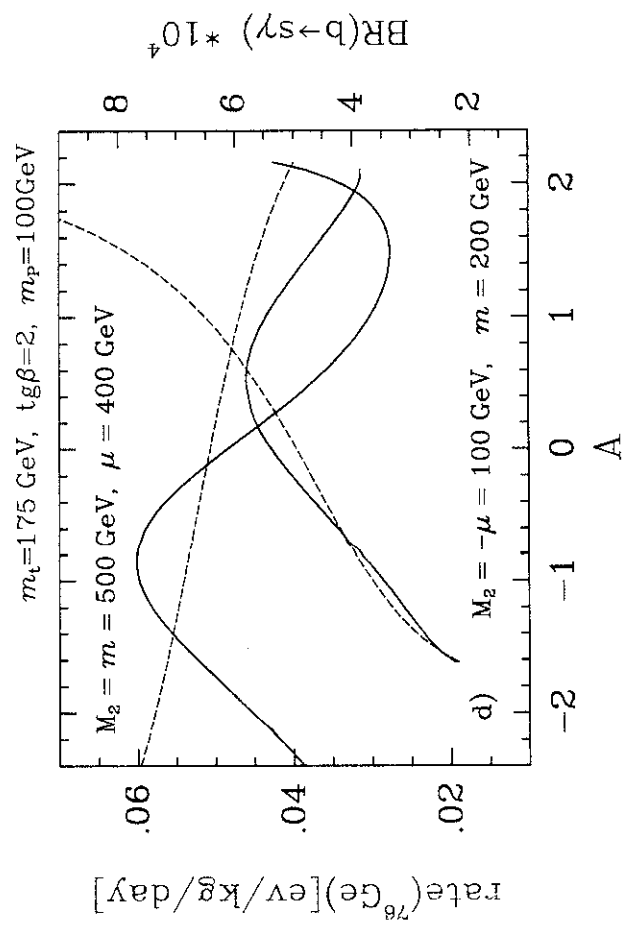


Fig 1d)



