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, Wisconsin 53706

Mihoko M. Nojiri*
Theory Group, KEK, Oho 1-1, Tsukuba, Ibaraki 305, Japan
(Received 6 September 1994)

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 $b \to s\gamma$. The recent measurement of this branching ratio, therefore, indicates that searches for supersymmetric dark matter might be even more difficult than previously anticipated.

PACS number(s): 95.35.+d, 12.15.Ji, 14.80.Ly

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 LSP's 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 as

$$0.025 \leq \Omega_{\rm LSP} h^2 \leq 1 \ , \tag{1}$$

where $\Omega_{\rm LSP}$ is the relic density in units of the closure density, and h is the Hubble constant in terms of 100 km/(sec Mpc). Observations imply $0.5 \leq h \leq 1$, the lower range, perhaps, being favored. The lower bound in (1) then follows from the requirement that there be enough relic LSP's to form the dark matter halos of galaxies ($\Omega_{\rm LSP} \geq 0.1$). 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 a 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-boson-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 μ and ratio of Higgs vacuum expectation values $\tan\beta$ entering the neutralino mass matrix[1] [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 psuedoscalar Higgs boson. As 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 top squark $\tilde{t}$ (see below). In our analysis we include these corrections using the effective potential method[2] [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

*Electronic address: NOJIRIN@JPNKEKVX

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

[2]Since not only corrections growing as ln $(m_{\tilde{t}}/M_t)$ are included, it is technically easier to present our results for mixed m_P, rather than for fixed mass of one of the scalar Higgs bosons.

supersymmetry-breaking contribution m^2 along the diagonal of their respective mass matrices. Our 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 check first for consistency with experimental and theoretical constraints before we use this spectrum to predict LSP detection rates. In particular, M_2, μ, and $\tan\beta$ must be chosen such that charginos and neutralinos escape detection at the CERN e^+e^- collider 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 measured [12] the branching ratio for inclusive $b \to s\gamma$ decays to be

$$B(b \to s\gamma) = (2.32 \pm 0.51 \pm 0.29 \pm 0.32) \times 10^{-4} , \quad (2)$$

where the errors are statistical, experimental systematics, and theoretical systematics (due to the extrapolation from the observed part of the photon spectrum), respectively. Adding all errors in quadrature, this implies 95% C.L. upper and lower limits on this branching ratio of 3.4×10^{-4} and 1.2×10^{-4}, respectively. These bounds are relevant for LSP searches since within the MSSM the $B(b \to s\gamma)$ 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[3]

$$m_{H^\pm}^2 = m_P^2 + m_W^2 , \quad (3)$$

where $m_W \simeq 80$ GeV is the mass of the W bosons. In particular, a light-charged Higgs boson gives a large positive contribution to the amplitude $\mathcal{A}(b \to s\gamma)$. 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 Higgs bosons when sparticles are taken to be heavy, although well suited for DM detection, tend to give results for the $B(b \to s\gamma)$ 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 $b \to s\gamma$, can occur. As mentioned, mimicking as closely as possible the treatment of squark masses in previous analyses 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 $b \to s\gamma$ can come from loops mediated by neutral gauginos (gluinos or neutralinos). Flavor mixing 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, affects only left-handed u-type squarks, $\tilde{u}$ in the 6×6 mass matrix:

$$\mathcal{M}_{\tilde{u}}^2 = \begin{pmatrix} \mathcal{M}_{\tilde{u}L}^2 & \mathcal{M}_{\tilde{u}LR}^2 \\ (\mathcal{M}_{\tilde{u}LR}^2)^\dagger & \mathcal{M}_{\tilde{u}R}^2 \end{pmatrix} . \quad (4)$$

The 3×3 left-left, right-right, and left-right mixing submatrices $\mathcal{M}_{\tilde{u}L}^2$, $\mathcal{M}_{\tilde{u}LR}^2$, and $\mathcal{M}_{\tilde{u}R}^2$ are given by

$$\left(\mathcal{M}_{\tilde{u}L}^2\right)_{ij} = (m^2 + 0.35 m_Z^2 \cos2\beta)\delta_{ij} + m_t^2 V_{i3}^* V_{3j} , \quad (5a)$$

$$\left(\mathcal{M}_{\tilde{u}R}^2\right)_{ij} = (m^2 + 0.15 m_Z^2 \cos2\beta)\delta_{ij} + m_t^2 \delta_{i3}\delta_{j3} , \quad (5b)$$

$$\left(\mathcal{M}_{\tilde{u}LR}^2\right)_{ij} = -(A_t + \mu\cot\beta) m_t V_{3i}^* \delta_{j3} . \quad (5c)$$

The symbols V_{ij} indicate here elements of the Cabibbo-Kobayushi-Maskawa (CKM) mixing matrix, μ is the mass parameter entering the neutralino mass matrix, and A_t is a soft supersymmetry-breaking parameter of order m. When writing Eqs. (5) we have neglected all Yukawa couplings except for the top quark. Similarly, the left-right mixing in (5c) 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 $B(b \to s\gamma)$. For definiteness we focus on a detector consisting of isotopically pure ^{76}Ge since such a device is now under construction. The impact of the measurement of $B(b \to s\gamma)$ on the prospects for direct relic LSP detection is very similar for detectors of different materials as long as the total LSP-nucleus cross section is dominated by spin-independent interactions, which is true in almost all cases. 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 Fig. 1 the LSP counting rate in such a detector (solid lines) as well as the branching ratio for $b \to s\gamma$ (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 $m_{\rm LSP} \simeq 200$ GeV), and the case of a much lighter one (M_2=100 GeV, $\mu = -100$ GeV giving $m_{\rm LSP} \simeq 50$ GeV). We fix the remaining supersymmetric parameters to be, in general, $A_t = 0$, $\tan\beta = 2$, $m_P = 150$ 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. 1(a)], $\tan\beta$ [1(b)], m [1(c)], or A_t [1(d)].

We have chosen $|\mu| > M_1$ so that the LSP is predominantly a gaugino; this is necessary [3] to satisfy the lower bound on $\Omega_{\rm LSP}h^2$ 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.

[3]Equation (3) holds at the tree level. Radiative corrections to this relation are very small [10] unless one somewhat artificially allows $\tan\beta < 1$.

The LSP detection rate can therefore be quite large *if* Higgs bosons are light.

This is illustrated in Fig. 1(a), where the dependence on m_P is shown. If M_2, m, and μ are large (upper pair of curves), m_P can be chosen as small as 50 GeV without violating direct Higgs boson 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 $B(b \to s\gamma)$ be $\leq 3.4 \times 10^{-4}$ implies $m_P \geq 500$ GeV and a counting rate of less than 0.01 events/(kg day), a value too small to be detectable in the near future. Notice, however, that no error in the theoretical prediction for $B(b \to s\gamma)$ 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 $b \to s\gamma$ 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 $m_P > 350$ GeV, values of $B(b \to s\gamma)$ *below* the standard model (SM) prediction of 2.9×10^{-4} [18]. Moreover, the lower bound on m_P is relaxed to 250 GeV in this scenario. It should be noted 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 μ (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. As a result the lower bound on m_P corresponds to a maximal LSP counting rate as small as 0.007 events/(kg day) for this case.

In Fig. 1(b) we show the tanβ dependence for our two choices of the LSP mass. Here we have chosen $\mu = -300$ GeV in the light LSP case, in order to ensure $\Omega_{\rm LSP}h^2 \geq 0.025$ for a sizable range of tanβ. 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 tanβ is given by the requirement that the Higgs boson escapes detection at LEP [11]. We see that the tanβ dependence of the LSP counting rate is quite similar in both cases. At first, the rate decreases with increasing tanβ since the mass M_{h^0} of the light neutral Higgs boson h^0 increases. This mass, however, remains essentially constant for tan$\beta \geq 5$. At the same time, the coupling of

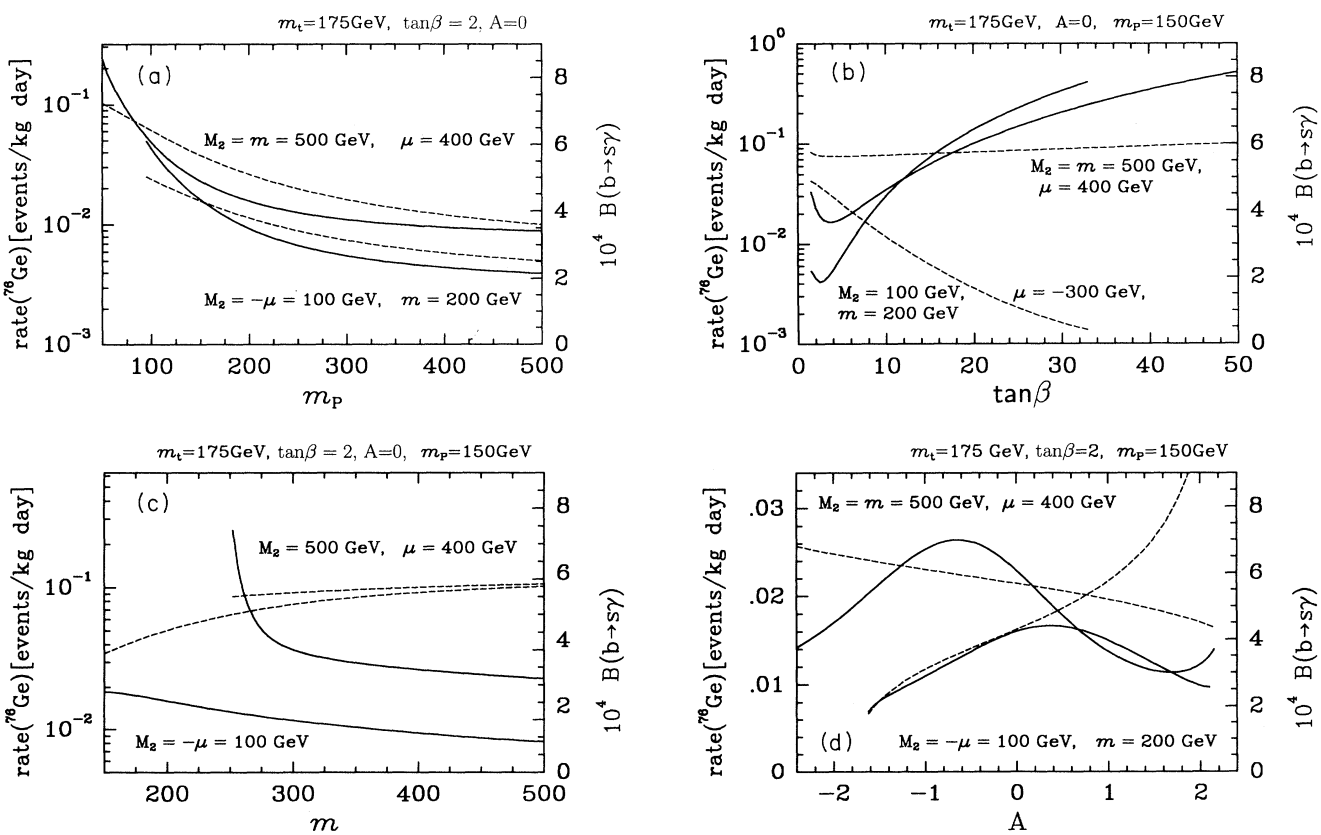

FIG. 1. Dependence of the LSP counting rate in a ^{76}Ge detector (solid lines) and $B(b \to s\gamma)$ (dashed lines) on (a) the mass m_P of the pseudoscalar Higgs boson, (b) the ratio of vacuum expectation values tanβ, (c) the scalar mass parameter m, and (d) the soft breaking parameter $A \equiv A_t/m$. Results are presented for $m_t = 175$ GeV and for the case of a heavy ($M_2 = 500$ GeV) and a light ($M_2 = 100$ GeV) LSP. The remaining MSSM parameters are specified in the text.

the heavier neutral Higgs boson H^0 to down-type quarks increases roughly as tanβ. 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 tanβ 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 $B(b \to s\gamma)$. The contribution from $H^\pm$ loops becomes independent of tanβ 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 tanβ. Notice that the entire curve lies well above the experimental upper bound of 3.4×10^{-4}. In contrast, in the case with a light LSP and comparatively light charginos and squarks, sparticle loops do contribute significantly. The observed strong tanβ dependence of this contribution is due to the fact that the chargino-b-$\tilde{t}$ interaction contains [8,13] a term proportional to the bottom Yukawa coupling, which grows as tanβ for large tanβ. For the given values of M_2 and m but positive μ the contribution from this term would be positive, leading to predictions for $B(b \to s\gamma)$ rapidly growing and quickly exceeding the experimental upper bound. For the given case of negative μ this contribution interferes destructively with the W and $H^\pm$ loops; for tan$\beta > 24$ one gets into conflict with the experimental lower bound which implies a counting rate of less than 0.2 events/(kg day) in the case with light LSP.

The dependence on the parameter m is shown in Fig. 1(c). 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 $\mathcal{O}(m_{\tilde{q}}^{-4})$ 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 $B(b \to s\gamma)$ 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, while in the case of light LSP only the region close to the lower bound on m imposed by Higgs searches is marginally compatible with the upper bound on the branching ratio.

In Fig. 1(d) 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. (5c) 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 $B(b \to s\gamma)$ increases or decreases

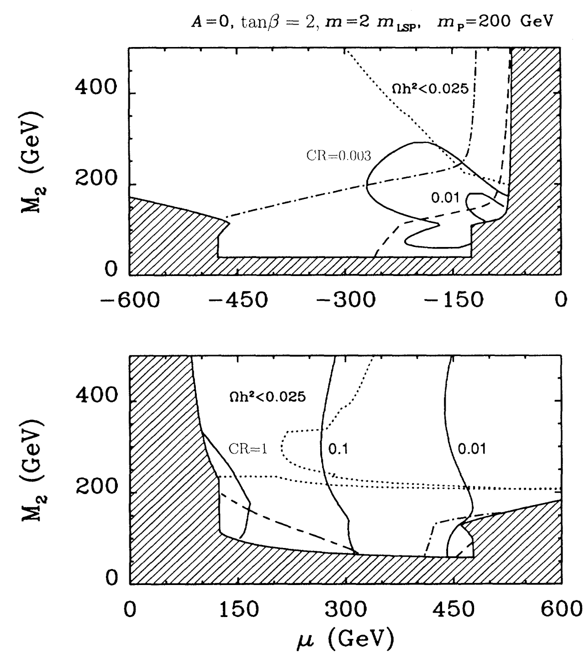

FIG. 2. Contours of LSP counting rate in a ^{76}Ge detector (solid lines) equal to 10.10.01 events/(kg day) (for positive μ), 0.01,003 events/(kg day) (for negative μ) and of constant $B(b \to s\gamma) = 3.4 \times 10^{-4}$ (dashed and dot-dashed lines) in the (M_2, μ) plane. The values of the other parameters are $M_t = 175$ GeV, $A_t = 0$, tan$\beta = 2$, $m_P = 200$ GeV, and $m = \min(2m_{\rm LSP}, 120$ GeV). In each frame the shaded regions are excluded by sparticle and Higgs boson searches and by the requirement that the LSP be neutral; the regions enclosed by the dotted lines have a very small LSP relic density, $\Omega_{\rm LSP}h^2 < 0.025$. The short-dashed contours of the $B(b \to s\gamma)$ have been computed using $Q_0 = 5$ GeV; the dot-dashed ones allow for some theoretical uncertainty, as described in the text, and the long-short-dashed contours correspond to $B(b \to s\gamma) = 4.2 \times 10^{-4}$.

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 μ, 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 μ. Notice that this latter scenario again violates the upper bound on the branching ratio over the entire parameter range shown here, while in the case of a light LSP this bound is violated for $A > -0.7$.

Finally, the dependence on M_2 and μ is shown in Fig. 2, where we have fixed $A_t = 0$, $\tan\beta = 2$, $m_P = 200$ GeV, and $m_t = 175$ GeV, and we have chosen m =min(120 GeV, $2m_{\rm LSP}$). The shaded regions in both 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 short-dashed lines indicate contours of constant $B(b \to s\gamma) = 3.4\times10^{-4}$: larger values of $B(b \to s\gamma)$ are obtained above these lines and smaller below. We also show contours of constant LSP counting rate = 1, 0.1, and 0.01 events/(kg day) for $\mu > 0$, and 0.01 and 0.003 events/(kg day) for $\mu < 0$ (solid lines). The dotted lines are contours of constant[4] $\Omega_{\rm LSP}h^2 = 0.025$. 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 about 50% of the region where the counting rate exceeds 0.1 events/(kg day) (for fixed neutralino flux) lies within the region with $\Omega_{\rm LSP}h^2 < 0.025$ (lower frame). It has been suggested in the literature to rescale 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.[5]

Prospects for direct LSP detection become even less promising once we require the $B(b \to s\gamma)$ to be below its experimental upper bound. Only the little region at small M_2 and $\mu \simeq 450$ GeV survives for positive μ, while for $\mu < 0$ the somewhat larger region to the right and below the short-dashed curve remains acceptable. The whole region where the counting rate exceeds 0.1 events/(kg day) is now excluded. For the allowed region with positive μ the counting rate is even below 0.01 events/(kg day). The implementation of the experimental bound on $B(b \to s\gamma)$ implies that, for the values of supersymmetric parameters chosen here, the maximal LSP counting rate in ^{76}Ge is about 0.007 and 0.02 events/(kg day) for positive and negative μ, respectively. Notice that a wider portion of parameter space survives for $\mu < 0$ where the LSP counting rate is usually smaller. In this region, in fact, the expected $B(b \to s\gamma)$ gets destructively interfering contributions from chargino-squark loops (at least in the region of small or moderate M_2).

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 GeV/cm^3 for the LSP mass density and 320 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 MeV [22]. Varying this value within the range favored by model calculations can change the prediction for the LSP counting rate by as much as a factor of 2.

The uncertainty in our prediction for $B(b \to s\gamma)$ 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 [23] 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 "1σ" uncertainty of about $\pm25\%$, 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 an already known part of higher-order QCD corrections [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 "1σ downward fluctuation" due to this theoretical uncertainty.[6] We should mention here that the *relative*

[4]To avoid figures too cluttered, we have omitted very narrow regions with $\Omega_{\rm LSP}h^2 < 0.025$ where s-channel exchange diagrams become resonant.

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

[6]Very recently another theoretical estimate of $B(b \to s\gamma)$ has appeared [24], where also parts of the next-to-leading order contributions to the relevant matrix of anomalous dimensions have been included. This introduces a strong renormalization scheme dependence. We prefer to follow here the approach of Ref. [18] where these terms are not included. The result of Ref. [24] falls within our theoretical error band.

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 $b \to s\gamma$ 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 of 3.4×10^{-4} as this conservative (low) estimate of the $B(b \to s\gamma)$ are shown by the long-short dashed lines in Fig. 2. We see that, for the given set of parameters, most of the (M_2, μ) half-plane with $\mu > 0$ is still excluded even if this lower theoretical estimate of the branching ratio is indeed correct, while for $\mu < 0$ the allowed region grows substantially. If this lower theoretical estimate is used, the maximal counting rate for $\mu > 0$ increases to about 0.02 events/(kg day); for $\mu < 0$ the upper bound on the counting rate is mostly due to direct experimental supersymmetry (SUSY) searches, but it is noteworthy that now the entire region of this half-plane where the counting rate exceeds 0.01 events/(kg day) is allowed.

In order to give the reader a feeling of how far above the experimental upper bound most of the half-plane with positive μ lies, we also show a contour where the lower theoretical estimate yields $B(b \to s\gamma) = 4.2 \times 10^{-4}$ (long-short-dashed curve), which is the 95% C.L. upper limit if all errors in Eq. (2) are added linearly. Only in the region to the left and below this line, as well as in the region where both M_2 and μ are very large so that sparticles decouple, does the lower estimate lie about this value. This means that most of the region with $\mu > 0$ is allowed only if we simultaneously assume a large downward fluctuation of the measured $B(b \to s\gamma)$ (including systematic errors) *and* a theoretical estimate for the branching ratio at the lower edge of the expected range. The contour where our *central* estimate for $B(b \to s\gamma) = 4.2 \times 10^{-4}$ almost coincides with the contour where the lower estimate yields 3.4×10^{-4}, which would again exclude most of the positive half-plane for the present choice of m_P, $\tan\beta$, m, and A.

If the lower theoretical estimate for $B(b \to s\gamma)$ turns out to be correct the interpretation of Fig. 1 would also change somewhat. For example, for the heavy LSP case shown in Fig. 1(a) a bound $m_P \geq 300$ GeV would result, leading to a reduction of the maximal counting rate by a factor of 20. The heavy LSP case in Figs. 1(b)–(d) would still be excluded over the whole range of parameters shown, however. 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 $B(b \to s\gamma)$. 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 $b \to s\gamma$ 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 the estimates of both the LSP counting rate and the $B(b \to s\gamma)$, 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 focused on direct LSP detection experiments. The expected signal rate in experiments looking for LSP annihilation in the center of the Earth or Sun [25] is also proportional to LSP-nucleus scattering cross sections. The upper limit on the $B(b \to s\gamma)$ 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, μ, and Higgs boson 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 roughly the same size. This statement can be quantified in so-called minimal supergravity theories [26], where the mechanism of radiative gauge symmetry breaking ensures that sparticle masses m_P and μ are strongly correlated. In such models the upper bound on $B(b \to s\gamma)$ is more easily satisfied [13,27]. The price one has to pay, though, is that the expected LSP counting rates are usually very low [7], of order 10^{-3} events/(kg day) or even less in a ^{76}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 favored by these supergravity models, which are at the same time theoretically very appealing but difficult to probe by LSP detection experiments.

M.D. acknowledges the hospitality of theory group at DESY, where part of this work was carried out. His work was supported in part by the U.S. Department of Energy under Contract No. DE-AC02-76ER00881, by the Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation, as well as by a grant from the Deutsche Forschungsgemeinschaft under the Heisenberg program. F.M.B.'s work was supported by Grant No. 05 5HH 91P(8) from the Bundesministerium für Forschung and Technologie, Bonn, Germany. M.M.N. was supported by Grant no. 06740326 from the Japanese Ministry of Education, Science and Culture.

[1] E.W. Kolb and M.S. Turner, *The Early Universe* (Addison-Wesley, Redwood City, CA, 1990).
[2] H.E. Haber and G.L. Kane, Phys. Rep. **117**, 75 (1985).
[3] Some recent examples are: K. Griest, M. Kamionkowski, and M.S. Turner, Phys. Rev. D **41**, 3565 (1990); K. Olive and M. Srednicki, Nucl. Phys. **B355**, 208 (1991); J. Ellis and L. Roszkowski, Phys. Lett. B **283**, 252 (1992); J.L. Lopez, D.V. Nanopoulos, and K. Yuan, Nucl. Phys. **B370**, 445 (1992); P. Nath and R. Arnowitt, Phys. Rev. Lett. **70**, 3696 (1993); M. Drees and M.M. Nojiri, Phys. Rev. D **47**, 376 (1993).
[4] See, e.g., talks by B. Cabrera and P.F. Smith, in A Critique of Dark Matter, Proceedings, Los Angeles, California, 1994, edited by D. Cline (unpublished).
[5] M.W. Goodman and E. Witten, Phys. Rev. D **31**, 3059 (1985).
[6] K. Griest, Phys. Rev. D **38**, 2357 (1988); M. Srednicki and R. Watkins, Phys. Lett. B **225**, 140 (1989); G.F. Giudice and E. Roulet, Nucl. Phys. **B316**, 429 (1989).
[7] M. Drees and M.M. Nojiri, Phys. Rev. D **48**, 3483 (1993).
[8] J.F. Gunion and H.E. Haber, Nucl. Phys. **B272**, 1 (1986); Erratum: University of Santa Cruz Report No. UCD-92-31, 1992 (unpublished).
[9] Y. Okada, M. Yamaguchi, and T. Yanagida, Prog. Theor. Phys. **85**, 1 (1991); Phys. Lett. B **262**, 54 (1991); H.E. Haber and R. Hempfling, Phys. Rev. Lett. **66**, 1815 (1991).
[10] J. Ellis, G. Ridolfi, and F. Zwirner, Phys. Lett. B **257**, 83 (1991); **262**, 477 (1991); A. Brignole, J. Ellis, G. Ridolfi, and F. Zwirner, *ibid.* **271**, 123 (1991); M. Drees and M.M. Nojiri, Phys. Rev. D **45**, 2482 (1992).
[11] For example, ALEPH Collaboration, D. Decamp *et al.*, Phys. Rep. **216**, 253 (1992); L3 Collaboration, O. Adriani *et al.*, *ibid.* **236**, 1 (1993).
[12] CLEO Collaboration, R. Ammar *et al.*, CLEO CONF No. 94-1 (unpublished).
[13] S. Bertolini, F. Borzumati, A. Masiero, and G. Ridolfi, Nucl. Phys. **B353**, 591 (1991).
[14] See, for example, S. Bertolini, F. Borzumati, and A. Masiero, in *B Decays*, 1st ed., edited by S. Stone (World Scientific, Singapore, 1992).
[15] J.L. Hewett, Phys. Rev. Lett. **70**, 1045 (1993); V. Barger, M.S. Berger, and R.J. Phillips, *ibid.* **70**, 1368 (1993).
[16] G. Gelmini, P. Gondolo, and E. Roulet, Nucl. Phys. **B351**, 623 (1991); J. Engel, S. Pittel, and P. Vogel, Int. J. Mod. Phys. E **1**, 1 (1992); A. Bottino *et al.*, Mod. Phys. Lett. A **7**, 733 (1992); V.A. Bednyakov, H.V. Klapdor-Kleingrothaus, and S.G. Kovalenko, JINR Report No. E2-93-448 (unpublished).
[17] D.O. Caldwell *et al.*, Phys. Rev. Lett. **61**, 510 (1988); **65**, 305 (1990); C. Bacci *et al.*, Phys. Lett. B **293**, 460 (1992); **295**, 330 (1992).
[18] A.J. Buras, M. Misiak, M. Münz, and S. Pokorski, Nucl. Phys. **B424**, 374 (1994).
[19] T.K. Hemmick *et al.*, Phys. Rev. D **41**, 2074 (1990).
[20] M.S. Turner, in Direct Detection of Dark Matter, Proceedings, Berkeley, CA, 1994, edited by B. Sadoulet (unpublished).
[21] M.S. Turner, Phys. Rev. D **33**, 889 (1986); R.A. Flores, Phys. Lett. B **215**, 73 (1988); For a post massive compact halo object (MACHO) study, see E. Gates and M.S. Turner, Phys. Rev. Lett. **72**, 2520 (1994).
[22] J. Gasser, H. Leutwyler, and M.E. Sainio, Phys. Lett. B **253**, 252 (1991).
[23] A. Ali and C. Greub, Z. Phys. C **60**, 19433 (1993).
[24] M. Ciuchini, E. Franco, G. Martinelli, L. Reina, and L. Silvestrini, Phys. Lett. B **334**, 137 (1994).
[25] M. Kamionkowski, Phys. Rev. D **44**, 3021 (1991); Kamiokande Collab., M. Mori *et al.*, Phys. Lett. B **270**, 89 (1991); **289**, 463 (1992); F. Halzen, in [4].
[26] H.P. Nilles, Phys. Rep. **110**, 1 (1984); L.E. Ibáñez and G.G. Ross, in *Perspectives in Higgs Physics*, edited by G.L. Kane (World Scientific, Singapore, 1993), p. 229.
[27] N. Oshimo, Nucl. Phys. **B404**, 20 (1993); J.L. Lopez, D. Nanopoulos, and G.T. Park, Phys. Rev. D **48**, 974 (1993); F.M. Borzumati, Z. Phys. C **63**, 19291 (1994); M.A. Diaz, Phys. Lett. **22B**, 207 (1994); S. Bertolini and F. Vissani, Scuola Internazionale Superiore Studi Avanzati Report No. SISSA 40/94/EP (unpublished).