Light scalar top quarks and supersymmetric dark matter

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A stable neutralino $\tilde{\chi}_1^0$, assumed to be the lightest supersymmetric particle, is a favored particle physics candidate for cosmological dark matter. We study coannihilation of the lightest neutralino with the lighter scalar top quark $\tilde{t}_1$. We show that for natural values of the neutralino mass, $\lesssim 300$ GeV, the $\tilde{\chi}_1^0$-$\tilde{t}_1$ mass difference has to exceed ~ 10 to 30 GeV if $\tilde{\chi}_1^0$ is to contribute significantly to the dark matter. Scenarios with smaller mass splitting, where $\tilde{t}_1$ is quite difficult to detect at collider experiments, are thus cosmologically disfavored. On the other hand, for small $\tilde{t}_1$-$\tilde{\chi}_1^0$ mass splitting, we show that coannihilation allows superparticle masses well beyond the reach of the CERN LHC, $m_{\tilde{\chi}_1^0} \sim 5$ TeV, without ''overclosing'' the Universe.

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There is convincing evidence [1] that most matter in the Universe is dark (nonluminous):

$$0.2 \lesssim \Omega_{\rm DM} \lesssim 1, \tag{1}$$

where $\Omega_{\rm DM}$ is the dark matter (DM) density in units of the critical density, so that $\Omega = 1$ corresponds to a flat Universe. On the other hand, analyses of big bang nucleosynthesis [2] imply that most DM is nonbaryonic (although dark baryons probably exist as well).

One of the favorite particle physics candidates for DM is the lightest neutralino $\tilde{\chi}_1^0$ [3], assumed to be the lightest supersymmetric particle (LSP). It is stable if R parity is conserved [4]; this is also a sufficient (although not necessary) condition for avoiding very fast nucleon decay in supersymmetric theories. The LSP makes an attractive DM candidate since the primary motivation for its introduction comes from particle physics arguments [5]: supersymmetry stabilizes the huge hierarchy between the weak and grand unification scales against radiative corrections, and if it is broken at a sufficiently high scale, it allows us to understand the origin of the hierarchy in terms of radiative breaking of the standard model (SM) electroweak $SU_L(2) \times U(1)_Y$ gauge symmetry; furthermore it allows for a consistent unification of the gauge couplings.

Supersymmetric contributions to DM then come as extra bonus, and for wide regions of parameter space, the LSP relic density falls in the preferred range Eq. (1). This is true in particular if the LSP is mostly a superpartner of the $U(1)_Y$ gauge boson, i.e., B-ino-like, and if both $m_{\tilde{\chi}_1^0}$ and the masses of SU(2) singlet scalar leptons fall in the natural range below a few hundred GeV [6] (but above [7] the mass range excluded by the CERN e^+e^- collider LEP experiments).

In most of this cosmologically favored region the cross section for the production of superparticles at the CERN Large Hadron Collider (LHC), as well as at future TeV-scale lepton colliders, would be quite large [8]. However, a large cross section by itself is not sufficient to guarantee discovery of a given superparticle. One also needs a signature that allows to discriminate between superparticle production and backgrounds from standard model processes. Most search strategies rely on the assumption that a large amount of visible energy is released when the superparticle one is searching for decays; this in turn requires a large mass splitting between this superparticle and the LSP.

The mass splitting between the LSP and the next-to-lightest superparticle (NLSP) $\tilde{P}$ also affects the estimate of the LSP relic density. Our previous statement about the cosmologically favored region of parameter space assumes that $\tilde{\chi}_1^0 \tilde{\chi}_1^0$ annihilation reactions are the only processes that change the number of superparticles at temperatures around $T_F \simeq m_{\tilde{\chi}_1^0}/20$, where the neutralino $\tilde{\chi}_1^0$ decouples from the plasma of SM particles. It has been known for some time [9] that this is not true if the LSP-NLSP mass splitting is small. In this case, reactions of the type

$$\tilde{\chi}_1^0 + X \leftrightarrow \tilde{P} + Y, \tag{2}$$

where X, Y are SM particles, occur much more frequently at a temperature $T \sim T_F$ than $\tilde{\chi}_1^0 \tilde{\chi}_1^0$ annihilation reactions do. The rate of the latter kind of process is proportional to two powers of the Boltzmann factor $\exp(-m_{\tilde{\chi}_1^0}/T_F) \simeq \exp(-20)$, whereas for $m_{\tilde{\chi}_1^0} \simeq m_{\tilde{P}}$ the rate for reaction (2) is linear in this factor. These reactions will therefore maintain *relative* equilibrium between the states $\tilde{\chi}_1^0$ and $\tilde{P}$ until long after all superparticles decouple from the standard model plasma.

The total number of superparticles can then not only be changed by $\tilde{\chi}_1^0 \tilde{\chi}_1^0$ annihilation, but also by the ''coannihilation'' processes

$$\tilde{\chi}_1^0 + \tilde{P} \leftrightarrow X + Y \quad \text{and} \quad \tilde{P} + \tilde{P}^{(*)} \leftrightarrow X + Y. \tag{3}$$

Eventually all particles $\tilde{P}$ and $\tilde{P}^*$ will decay into $\tilde{\chi}_1^0$ (plus SM particles). In order to compute today's LSP relic density, we therefore only have to solve the Boltzmann equation for the sum $n_{\rm SUSY}$ of densities n_i of all relevant species of superparticles. In this sum contributions from reactions (2) cancel, since they do not change the total number of superparticles. One thus has [9]

$$\frac{dn_{\rm SUSY}}{dt} = -3Hn_{\rm SUSY} - \sum_{i,j} \langle \sigma_{ij} v\rangle (n_i n_j - n_i^{\rm eq} n_j^{\rm eq}) = -3Hn_{\rm SUSY} - \langle \sigma_{\rm eff} v\rangle (n_{\rm SUSY}^2 - n_{\rm SUSY}^{\rm eq^2}). \quad (4)$$

Here, H is the Hubble parameter, $\langle \cdots \rangle$ denotes thermal averaging, v is the relative velocity between the two annihilating superparticles in their center-of-mass frame, and the superscript ''eq'' indicates the equilibrium density. In the second step we made use of the fact that, as argued above, all relevant heavier superparticles maintain relative equilibrium to the neutralino LSP until long after the temperature T_F. This allowed us to sum all superparticle annihilation processes into an ''effective'' cross section; schematically [9]

$$\sigma_{\rm eff} \propto g_{\tilde\chi\tilde\chi}\sigma(\tilde\chi_1^0\tilde\chi_1^0) + g_{\tilde\chi\tilde P} B_{\tilde P}\sigma(\tilde\chi_1^0\tilde P) + g_{\tilde P\tilde P}(B_{\tilde P})^2\sigma(\tilde P\tilde P^{(*)}), \quad (5)$$

where the g_{ij} are multiplicity factors, and

$$B_{\tilde P} = (m_{\tilde P}/m_{\tilde\chi_1^0})^{3/2} e^{-(m_{\tilde P} - m_{\tilde\chi_1^0})/T} \quad (6)$$

is the temperature-dependent relative Boltzmann factor between the $\tilde P$ and $\tilde\chi_1^0$ densities. The final LSP relic density $\Omega_{\tilde\chi} h^2$, where $h = 0.65 \pm 0.15$ is the scaled Hubble constant, is then essentially inversely proportional to $\langle \sigma_{\rm eff} v\rangle$ at $T_F \simeq m_{\tilde\chi_1^0}/20$. Coannihilation can therefore reduce the LSP relic density by a large factor, if $\delta m \equiv m_{\tilde P} - m_{\tilde\chi_1^0} \ll m_{\tilde\chi_1^0}$ and $\sigma(\tilde\chi_1^0\tilde P) + \sigma(\tilde P\tilde P^{(*)}) \gg \sigma(\tilde\chi_1^0\tilde\chi_1^0)$. This is true in particular if $\tilde\chi_1^0$ is a light, $m_{\tilde\chi_1^0} < M_W$, Higgsino [6,10] or SU(2) gaugino [11]. More recently it has been pointed out [12,13] that coannihilation with light sleptons can reduce the relic density of a B-ino-like LSP by about one order of magnitude.

In this paper we study coannihilation of neutralinos with the lighter scalar top (top squark) eigenstate $\tilde t_1$. Compared to the other squarks, $m_{\tilde t_1}$ is reduced [5] by contributions of the large top quark Yukawa coupling to the relevant renormalization group equations, as well as by mixing between SU(2) doublet and singlet top squarks. While we do not know of any model that predicts $m_{\tilde t_1} \simeq m_{\tilde\chi_1^0}$, a close mass degeneracy is possible in many models, e.g., in the popular minimal supergravity (MSUGRA) model [5]. Moreover, scenarios with small $\tilde t_1$-$\tilde\chi_1^0$ mass splitting are of great concern for experimenters, since $\tilde t_1$ decays then release little visible energy, making $\tilde t_1$ production very difficult to detect at both e^+e^- [14] and hadron [15] colliders.

In contrast to the cases mentioned earlier, for $\tilde P = \tilde t_1$ it is not entirely obvious that reactions of the type (2) will indeed be much faster than $\tilde\chi_1^0\tilde\chi_1^0$ annihilation processes. In the absence of flavor mixing, one would have to chose $X = W, Y = b$ or vice versa. However, for a temperature $T < M_W$, the W density is itself quite small, so reaction (2) would be much faster than $\tilde\chi_1^0\tilde\chi_1^0$annihilation only for $m_{\tilde\chi_1^0}$ significantly above M_W. On the other hand, most supersymmetric models predict some amount of flavor mixing in the squark sector, even if it is absent at some high-energy scale. As a result, for small δm the dominant $\tilde t_1$ decay mode is usually its flavor changing two-body decay into $\tilde\chi_1^0 + c$ [16]. For $\tilde t_1$ masses of current experimental interest the dominant contribution to Eq. (2) therefore comes from $X = c, Y =$ nothing, i.e., (inverse) $\tilde t_1$ decay. If the effective $c\tilde t_1\tilde\chi_1^0$ coupling is suppressed by a small mixing angle ϵ, the condition that Eq. (2) is much faster than $\tilde\chi_1^0\tilde\chi_1^0$ annihilation reads

$$\epsilon^2 e^{-\delta m/T_F} \gg \alpha e^{-m_{\tilde\chi_1^0}/T_F}, \quad (7)$$

where the extra factor of $\alpha \sim 0.01$ occurs since we are comparing $2 \leftrightarrow 1$ reactions with $2 \leftrightarrow 2$ processes. For $\delta m \lesssim T_F \sim m_{\tilde\chi_1^0}/20$ we then only need $\epsilon > e^{-10} \simeq 5 \times 10^{-5}$. In what follows we will assume that this is true, or that $\tilde\chi_1^0$ is sufficiently heavy that $\tilde\chi_1^0 + W^+ \leftrightarrow \tilde t_1 + \overline{b}$ is fast.

Another property of the top squark is that it has strong interactions. A leading order calculation of $\sigma(\tilde\chi_1^0\tilde t_1)$ and $\sigma(\tilde t_1\tilde t_1^{(*)})$ will therefore not be very reliable. Unfortunately a full higher order calculation is highly nontrivial, since one would need to include finite temperature effects (e.g., in order to cancel Coulomb singularities in the nonrelativistic limit). We expect these unknown higher order QCD corrections to be more important than the contributions of higher partial waves. In the calculation of the cross sections $\sigma(\tilde\chi_1^0\tilde t_1)$ and $\sigma(\tilde t_1\tilde t_1^{(*)})$ we therefore only include the leading, S-wave contribution; however, the P-wave contributions to $\tilde\chi_1^0\tilde\chi_1^0$ annihilation process [6] are included. Our coannihilation cross sections will thus only be accurate to a factor of 2 or so. Because of the exponential dependence of $\sigma_{\rm eff}$ on δm, see Eqs. (5),(6), the bounds on the $\tilde t_1$-$\tilde\chi_1^0$ mass splitting that will be inferred from upper or lower bounds on $\Omega_{\tilde\chi} h^2$ should nevertheless be fairly accurate.

The existence of unknown, but probably large, higher order corrections also means that we can ignore all $\tilde t_1$ annihilation reactions that involve more than the minimal required number of electroweak gauge couplings. However, we treat the top and bottom quarks Yukawa couplings on the same footing as the strong coupling (the latter Yukawa coupling will be large only for $\tan\beta \sim m_t/m_b$). Altogether we therefore computed the cross sections for the following processes:

$$\begin{aligned} &\tilde\chi_1^0\tilde t_1 \to tg, tH_i^0, bH^+, \\ &\tilde t_1\tilde t_1 \to tt, \\ &\tilde t_1\tilde t_1^* \to gg, H_i^0H_j^0, H^+H^-, b\overline{b}, t\overline{t}, \end{aligned} \quad (8)$$

where $H_i^0 \equiv h, H, A$ is one of the three neutral Higgs bosons of the minimal supersymmetric standard model (MSSM) [17]. The cross sections for $\tilde\chi_1^0\tilde t_1^*$ and $\tilde t_1^*\tilde t_1^*$ annihilation are identical to those in the first and second lines of Eq. (8),

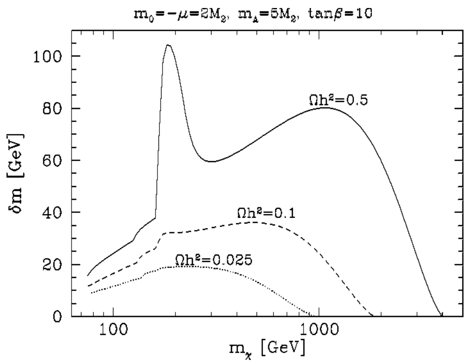

FIG. 1. Contours of constant $\Omega_{\tilde\chi}h^2=0.5$ (solid), 0.1 (dashed), and 0.025 (dotted) in the $(m_{\tilde\chi_1^0},\delta m)$ plane, where $\delta m=m_{\tilde t_1}-m_{\tilde\chi_1^0}$. We took μ, m_0, and m_A to be fixed multiples of $M_2\simeq 2m_{\tilde\chi_1^0}$, as indicated, whereas $\tan\beta=10$ has been kept fixed. The parameter A_0 varies between about $2.5m_0$ and $3.2m_0$, with larger A_0 values corresponding to smaller values of δm.

respectively. We have performed two independent calculations of these cross sections. One calculation was based on trace techniques and the usual polarization sum for external gluons; here the nonrelativistic limit (to extract the S-wave contribution [3]) was only taken at the end. The second method uses helicity amplitudes [6]; in this case the nonrelativistic limit can already be taken at the beginning of the calculation. (Note that the cross sections for $\tilde t_1\tilde t_1^*\to H_i^0 g$ vanish in this limit.) Explicit expressions for these cross sections will be published elsewhere. (Note that Refs. [12,13] do not keep the mass of the relevant SM fermion, in their case the τ lepton, whereas we have to keep a finite value for the top quark mass, $m_t\neq 0$. Reference [12] also did not include $\tilde f_L$-$\tilde f_R$ mixing, which in our case is crucial for obtaining a light $\tilde t_1$. In the relevant limit we agree with Refs. [12,13].) Given the cross sections for reactions (8), our calculation of the relic density closely follows the one of Ref. [9].

We use a variant of the minimal supergravity model [5] for our numerical analysis. In particular, we assume a common gaugino mass, a common sfermion mass m_0, and a common trilinear soft breaking parameter A_0 at the grand unification scale $M_X=2\times 10^{16}$ GeV. However, we allow the soft breaking masses of the two Higgs doublets to differ from m_0. In practice, this means that we keep the Higgsino mass parameter μ and the mass m_A of the CP-odd Higgs boson as free parameters at the weak scale. The final free parameter is the ratio $\tan\beta$ of vacuum expectation values of the two Higgs fields.

For illustration, we take $\mu=-2M_2$, where $M_2\simeq 2m_{\tilde\chi_1^0}$ is the SU(2) gaugino mass. This implies that the LSP is B-ino-like, which is the most natural choice for this type of model [18]. It is also conservative, since a Higgsino-like LSP will have larger couplings to top (quarks and squarks), and hence even larger co–annihilation cross sections Eq. (8). We also chose a large sfermion mass $m_0=2M_2$. In the absence of coannihilation this choice is usually incompatible [6] with the upper bound on the LSP relic density, which we conservatively take as $\Omega_{\tilde\chi}h^2\leqslant 0.5$.

In Fig. 1 we show contours of constant $\Omega_{\tilde\chi}h^2$ in the $(m_{\tilde\chi_1^0},\delta m)$ plane for a scenario with moderate $\tan\beta$ and a very heavy Higgs boson spectrum $m_A=5M_2$. This latter choice implies that the only Higgs boson relevant for the calculation of the LSP relic density is the light CP-even scalar h, with mass $m_h\leqslant 130$ GeV. This is a conservative scenario in the sense that it minimizes the number of final states contributing in Eqs. (8), and also leads to a small $\tilde\chi_1^0\tilde\chi_1^0$ annihilation cross section. We see that scenarios with very large δm values are indeed excluded by the upper bound on $\Omega_{\tilde\chi}h^2$. The peak in the contour $\Omega_{\tilde\chi}h^2=0.5$ at $m_{\tilde\chi_1^0}\simeq m_t$ is due to $\tilde\chi_1^0\tilde\chi_1^0\to t\bar t$, which has a sizable S-wave cross section if $\tilde t_1$ is not too heavy and $m_{\tilde\chi_1^0}$ is not much above m_t. The much smaller bumps at $m_{\tilde\chi_1^0}\simeq 130$ GeV are due to hh final states becoming accessible.

On the other hand, for very small values of δm and $m_{\tilde\chi_1^0}$ in the range indicated by naturalness arguments ($\lesssim 0.3$ TeV, corresponding to a gluino mass $m_{\tilde g}\lesssim 2$ TeV), we find that the LSP cannot contribute significantly to the solution of the dark matter puzzle, since its relic density is too small. In particular, one needs [3] $\Omega_{\tilde\chi}h^2>0.025$ for $\tilde\chi_1^0$ to form galactic haloes. We see that even for the present very conservative choice of parameters one needs a $\tilde t_1$-$\tilde\chi_1^0$ mass splitting of at least 9 to 19 GeV (6 to 10 %) to satisfy this lower bound on $\Omega_{\tilde\chi}h^2$. This mass splitting is large enough for standard $\tilde t_1$ search methods at e^+e^- colliders [14,19] to have reasonably high efficiency. If we require that $\Omega_{\tilde\chi}h^2$ lies in the currently favored ''best fit'' range between about 0.1 and 0.2, δm has to be between 11 and 33 GeV. Unfortunately this is still not high enough for current $\tilde t_1$ search strategies at the Tevatron [15] to be sensitive.

So far we have focused on LSP masses in the range favored by naturalness arguments. It is sometimes claimed [8] that the upper bound on $\Omega_{\tilde\chi}h^2$ implies that LHC experiments must find superparticles if the MSSM is correct and $\tilde\chi_1^0$ is B-ino-like. Unfortunately this is not true; for $\delta m\to 0$ an LSP mass up to 4 TeV, corresponding to a gluino mass in excess of 20 TeV, cannot be excluded from this cosmological argument. (As noted above, our estimates for $\tilde t_1$annihilation cross sections are not very reliable. However, even if we overestimated them by a factor of 2, the bound on $m_{\tilde\chi_1^0}$ would only be reduced by a factor of $\sqrt{2}$, and would thus still allow sparticle masses far above the range to be covered by the LHC.)

In Fig. 2 we show analogous results for a light spectrum of Higgs bosons and large $\tan\beta$, where the bottom Yukawa coupling is sizable; the choice $m_A=0.35M_2\simeq 0.7m_{\tilde\chi_1^0}$ ensures that all Higgs pair final states will be accessible for $m_{\tilde\chi_1^0}>100$ GeV. However, we keep the previous (large) values for $|\mu|$ and m_0. Nevertheless we see that for natural values of $m_{\tilde\chi_1^0}$, requiring $\Omega_{\tilde\chi}h^2>0.025$ now implies $\delta m>20$ GeV. Moreover, the LSP makes a good DM candidate, i.e., $\Omega_{\tilde\chi}h^2\sim 0.1$, only for $\delta m\gtrsim 40$ GeV; this is sufficiently large to permit $\tilde t_1$ searches at the Tevatron [14,20]. Finally, for $\delta m\to 0$, cosmology now allows an LSP mass up to 6 TeV, corresponding to a gluino mass of about 30 TeV. Ob-

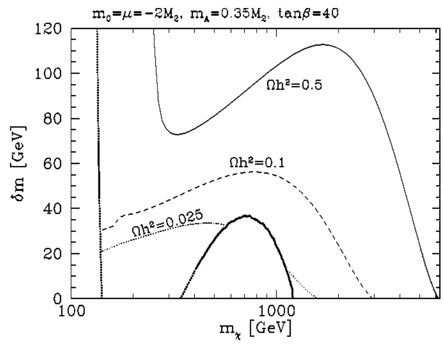

FIG. 2. As in Fig. 1, except that we took a large value of tan β and a light Higgs boson spectrum. The regions below and to the left of the heavy dotted lines are excluded by Higgs boson searches at LEP.

viously the upper bound on δm that follows from $\Omega_{\tilde{\chi}} h^2 > 0.025$ for natural values of $m_{\tilde{\chi}_1^0}$, as well as the absolute upper bound on $m_{\tilde{\chi}_1^0}$ that follows from $\Omega_{\tilde{\chi}} h^2 \leqslant 0.5$, are even higher if we chose smaller values for m_0 and/or $|\mu|$.

In conclusion, we have shown that scenarios with very small $\tilde{t}_1$-$\tilde{\chi}_1^0$ mass splitting would permit an LSP mass of several TeV without ''overclosing'' the Universe. This shows once again [6] that the upper bound on the LSP relic density does not guarantee that LHC experiments will detect superparticles, even if the MSSM is correct; of course, (third generation) superparticles with masses out of the reach of the LHC can hardly be argued to be ''natural.'' On the other hand, for $\tilde{\chi}_1^0$ and $\tilde{t}_1$ masses of present experimental interest, and indeed for the entire natural range of these masses, $\tilde{\chi}_1^0$ cannot contribute significantly to the dark matter in the Universe unless the $\tilde{t}_1$-$\tilde{\chi}_1^0$ mass difference is large enough for conventional $\tilde{t}_1$ search strategies at e^+e^- colliders to be effective. This does not imply that collider searches for $\tilde{t}_1$ nearly degenerate with $\tilde{\chi}_1^0$ [21] should not be continued; a positive signal would definitely exclude $\tilde{\chi}_1^0$ as DM candidate, which is not easy to accomplish with cosmological dark matter searches. However, since dark matter is known to exist, for natural values of $m_{\tilde{\chi}_1^0}$ a very small $\tilde{t}_1$-$\tilde{\chi}_1^0$ mass splitting would require physics beyond the MSSM.

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[1] For a recent review, see N. Bahcall, J.P. Ostriker, and S. Perlmutter, Science **284**, 1481 (1999).
[2] S. Sarkar, Rep. Prog. Phys. **59**, 1493 (1996).
[3] G. Jungman, M. Kamionkowski, and K. Griest, Phys. Rep. **267**, 195 (1996).
[4] P. Fayet, Phys. Lett. **69B**, 489 (1977).
[5] H.P. Nilles, Phys. Rep. **110**, 1 (1984).
[6] M. Drees and M.M. Nojiri, Phys. Rev. D **47**, 376 (1993).
[7] L. Roszkowski, Phys. Lett. B **278**, 147 (1992).
[8] J.D. Wells, Phys. Lett. B **443**, 196 (1998).
[9] K. Griest and D. Seckel, Phys. Rev. D **43**, 3191 (1991).
[10] S. Mizuta and M. Yamaguchi, Phys. Lett. B **298**, 120 (1993); J. Edsjö and P. Gondolo, Phys. Rev. D **56**, 1879 (1997).
[11] S. Mizuta, D. Ng, and M. Yamaguchi, Phys. Lett. B **300**, 96 (1993).
[12] J. Ellis, T. Falk, K.A. Olive, and M. Srednicki, Astropart. Phys. **13**, 181 (2000).
[13] M. Gomes, G. Lazarides, and C. Pallis, Phys. Rev. D **61**, 123512 (2000).
[14] ALEPH Collaboration, R. Barate *et al.*, Phys. Lett. B **434**, 189 (1998); DELPHI Collaboration, P. Abreu *et al.*, *ibid.* **444**, 491 (1998); L3 Collaboration, M. Acciarri *et al.*, *ibid.* **445**, 428 (1999); OPAL Collaboration, G. Abbiendi *et al.*, *ibid.* **456**, 95 (1999).
[15] DØ Collaboration, S. Abachi *et al.*, Phys. Rev. D **57**, 589 (1998); CDF Collaboration, T. Affolder *et al.*, Phys. Rev. Lett. **84**, 5704 (2000).
[16] K.-I. Hikasa and M. Kobayashi, Phys. Rev. D **36**, 724 (1987); C. Boehm, A. Djouadi, and Y. Mambrini, *ibid.* **61**, 095006 (2000).
[17] J. Gunion and H. Haber, Nucl. Phys. **B272**, 1 (1986).
[18] T. Falk, Phys. Lett. B **456**, 171 (1999).
[19] M. Berggren, R. Keranen, H. Nowak, and A. Sopczak, hep-ph/9911345.
[20] H. Baer, J. Sender, and X. Tata, Phys. Rev. D **50**, 4517 (1994).
[21] M. Drees and O. Éboli, Eur. J. Phys. **10**, 337 (1999).