DESY 06-011
KAIST-TH 2006/02

DESY 06–011
KAIST–TH 2006/02
hep-ph/0602109
February 13, 2006

Heavy Higgs Resonances for the Neutralino Relic Density in the Higgs Decoupling Limit of the CP–noninvariant Minimal Supersymmetric Standard Model

S.Y. Choi[1*] and Y.G. Kim[2†]

[1] *Deutches Elektronen–Synchrotron DESY, D–22603 Hamburg, Germany and Physics Department and RIPC, Chonbuk National University, Jeonju 561-756, Korea*

[2] *Department of Physics, KAIST, Daejon 305-701, Korea*

Abstract

The lightest neutralino is a compelling candidate to account for cold dark matter in the universe in supersymmetric theories with R–parity. In the CP–invariant theory, the neutralino relic density can be found in accord with recent WMAP data if neutralino annihilation in the early universe occurs via the s–channel A funnel. In contrast, in the CP–noninvariant theory two heavy neutral Higgs bosons can contribute to the Higgs funnel mechanism significantly due to a CP–violating *complex* mixing between two heavy states, in particular, when they are almost degenerate. With a simple analytic and numerical analysis, we demonstrate that the CP–violating Higgs mixing can modify the profile of the neutralino relic density *considerably* in the heavy Higgs funnel with the neutralino mass close to half of the heavy Higgs masses.

*E-mail: sychoi@chonbuk.ac.kr
†E-mail: ygkim@muon.kaist.ac.kr

The nature of the dark matter is one of the most important questions at the interface of particle physics and cosmology. Recently there have been big improvements in the astrophysical and cosmological data, most notably due to the Wilkinson microwave anisotropy probe (WMAP)[1] and the Sloan digital sky survey (SDSS)[2]. With the data we can infer the following 2σ range for the density of cold dark matter normalized by the critical density

$$0.094 < \Omega_{\rm CDM} h^2 < 0.129, \tag{1}$$

where $h \approx 0.7$ is the (scaled) Hubble constant in units of 100 km/sec/Mpc. Such a precise determination of $\Omega_{\rm CDM} h^2$ imposes severe constraints on any model that tries to explain it.

In supersymmetric theories with R–parity [3], the lightest supersymmetric particle (LSP), which is typically the lightest neutralino $\tilde{\chi}_1^0 \equiv \chi$, is stable and it serves as an excellent cold dark matter (CDM) candidate [4, 5]. However, typical mSUGRA models in the parameter space of minimal SUSY predict much larger values for the neutralino relic density than the values in the range (1). Some specific mechanisms leading to strongly enhanced neutralino annihilation are required to produce the observed dark matter relic density [6]. Such an enhancement might be due to the presence of light sleptons, enhancing the LSP annihilation into leptons, to an accidental degeneracy of the LSP and the lighter stau (or stop), leading to enhanced LSP–stau (or stop) co–annihilation, to the LSP with significantly mixed gaugino–higgsino components, enhancing the annihilation into gauge bosons, or to an accidental degeneracy $M_A \approx 2m_\chi$ with large $\tan\beta$, leading to enhanced annihilation through an s–channel pseudoscalar A in the CP–invariant theory.

In particular, the enhanced LSP annihilation via a A funnel in the CP–invariant case is due to two reasons; (i) the S–wave amplitude for $\chi\chi \to A$ is not suppressed near threshold while the P–wave amplitude for $\chi\chi \to H$ is suppressed near threshold and (ii) the total A decay width[‡] becomes large as the $A \to b\bar{b}$ decay mode is greatly enhanced for large $\tan\beta$.

The generic feature of the A funnel enhancement could, however, be greatly modified due to the CP–violating mixing among neutral Higgs bosons as well as due to the CP–violating Higgs couplings to neutralino pairs in the CP–noninvariant theory [7, 8, 9, 10, 11]. In this work we analyze, both analytically and numerically, the impact on the LSP relic density by the CP–violating Higgs mixing, loop–induced at the loop level in the CP–noninvariant MSSM [12]. To be specific, we consider the case when two (almost) degenerate heavy

[‡]The decays, $A \to W^+W^-$ and $A \to ZZ$, are forbidden, leading to a small A width for small $\tan\beta$ (unless the decay $A \to t\bar{t}$ is open).

neutral Higgs bosons H and A are essentially decoupled from the lightest neutral Higgs boson[§] and their masses are very close to twice the LSP mass.

With the lightest neutral Higgs boson decoupled, the CP–violating mixing of the two nearly–degenerate heavy Higgs bosons is described by a 2×2 *complex* mass matrix, composed of a real dissipative part and an imaginary absorptive part [10]. This mixing can be very large, generating frequent mutual transitions inducing large CP–odd mixing effects, which are quantitatively described by the complex mixing parameter X:

$$X = \frac{1}{2} \tan 2\theta = \frac{\Delta^2_{HA}}{M^2_H - M^2_A - i\,[M_H \Gamma_H - M_A \Gamma_A]}, \tag{2}$$

where the complex off–diagonal term Δ^2_{HA} of the Higgs mass matrix couples two Higgs states.

The Higgs masses and widths are then shifted in a characteristic pattern by the CP–violating mixing [14], of which the individual shifts can be obtained by separating real and imaginary parts in the relations:

$$\begin{aligned} \Big[M^2_{H_2} - i M_{H_2}\Gamma_{H_2}\Big] - \Big[M^2_H - i M_H \Gamma_H\Big] &= -\Big\{\Big[M^2_{H_3} - i M_{H_3}\Gamma_{H_3}\Big] - \Big[M^2_A - i M_A \Gamma_A\Big]\Big\} \\ &= -\Big\{\Big[M^2_A - i M_A \Gamma_A\Big] - \Big[M^2_H - i M_H \Gamma_H\Big]\Big\} \times \tfrac{1}{2}\,[\sqrt{1+4X^2} - 1] \end{aligned} \tag{3}$$

In such a non–Hermitian mixing the ket and bra mass eigenstates have to be defined separately: $|H_i\rangle = C_{i\alpha}|H_\alpha\rangle$ and $\langle \widetilde{H}_i| = C_{i\alpha}\langle H_\alpha|$ ($i = 2, 3$ and $H_\alpha = H, A$); $C_{2H} = \cos\theta, C_{2A} = \sin\theta, C_{3H} = -\sin\theta$ and $C_{3A} = \cos\theta$ in terms of the complex mixing angle θ.

As two mass eigenstates have no definite CP parity and an enlarged mass splitting, the profile of the LSP relic density can considerably be modified in the heavy Higgs funnel. For a simple analytic and numerical illustration, we consider a specific scenario within the CP–violating MSSM [MSSM–CP], while a more comprehensive analysis is separately given in a future publication. We assume the source of CP violation to be localized entirely in the complex stop trilinear coupling A_t but all the other interactions to be CP conserving.[¶]

In this situation, CP violation is transmitted through stop–loop corrections to the effective Higgs potential, generating three CP–odd complex quartic parameters. The effective parameters have been calculated in Ref. [8] to two–loop accuracy and, with $t/\tilde{t}$ contributions, the parameters are determined by the parameters; the SUSY scale M_S which is taken

[§] This situation is naturally realized in the MSSM in the decoupling limit with $M_A > 2m_Z$ [10, 13].

[¶] This assignment is compatible with the bounds from the electric dipole moment measurements [15].

to be essentially the average of two stop masses–squared, the higgsino parameter μ, the stop trilinear parameter A_t and the top Yukawa coupling $h_t = \sqrt{2}\bar{m}_t/v \sin\beta$ defined with the running $\overline{\rm MS}$ mass $\bar{m}_t$ and the Higgs vacuum expectation value $v \approx 246$ GeV. The one–loop improved Born Higgs mass matrix is derived from this effective Higgs potential and then the matrix elements are shifted to the pole–mass parameters by including dispersive contributions from Higgs self–energies.

Before evaluating the impact of the complex H/A mixing on the LSP relic density in the heavy–Higgs funnel, we describe an approximate procedure for estimating the relic density [16]. The LSP number density is evolved in time according to the Boltzmann equation. When the temperature of the Universe is higher than the LSP mass, the number density is simply given by its thermal–equilibrium density. However, once the temperature drops below the LSP mass, the number density drops exponentially. As a result, the LSP annihilation rate becomes smaller than the Hubble expansion rate at a certain point when the LSP neutralinos fall out of equilibrium and the LSP number density in a co–moving volume remains constant. The present LSP relic abundance is then approximately given by

$$\Omega h^2 \simeq \frac{1.07 \times 10^9 \ {\rm GeV}^{-1}}{{\rm J} g_*^{1/2} \, {\rm M_{PL}}}, \tag{4}$$

where $g_* = 81$ is the number of relativistic degrees of freedom and ${\rm M_{PL}} = 1.22 \times 10^{19}$ GeV is the Planck mass. And the integral J is given by

$$J(x_f) = \int_{x_f}^{\infty} \frac{\langle \sigma v \rangle}{x^2} dx, \tag{5}$$

where $\langle \sigma v \rangle$ is the thermally averaged LSP annihilation cross section times the relative velocity v of two annihilating LSPs, and $x_f = m_\chi/T_f$ with the freeze–out temperature $T_f \simeq m_\chi/25$ for typical weak–scale numbers. We take $x_f = 25$ in the following numerical demonstration.

When the heavy Higgs boson masses are large and close to twice the LSP mass, the LSP annihilation is dominated by heavy Higgs–boson exchanges. The LSP annihilation rate can then be estimated with reasonable approximation by including only the s–channel heavy Higgs boson exchanges. In the decoupling limit the $H\chi\chi$ and $A\chi\chi$ couplings read

$$\begin{aligned}
\langle \chi_L | H | \chi_R \rangle &= \langle \chi_R | H | \chi_L \rangle^* \simeq -\frac{g}{2} \, (N_{12} - N_{11} \tan\theta_W)(\sin\beta N_{13} + \cos\beta N_{14}), \\
\langle \chi_L | A | \chi_R \rangle &= \langle \chi_R | A | \chi_L \rangle^* = -\frac{g}{2} \, i(N_{12} - N_{11} \tan\theta_W)(\sin\beta N_{13} - \cos\beta N_{14}),
\end{aligned} \tag{6}$$

in terms of $\tan\beta$ and the neutralino mixing matrix $N_{i\alpha}$ (i, α = 1-4) diagonalizing the neutralino mass matrix $\mathcal{M}_N$ as $N^* \mathcal{M}_N N^\dagger = \mathcal{M}_{diag}$ [17]. The LSP annihilation rate multiplied by the relative velocity v of two LSPs can be expressed as

$$\sigma v = \frac{1}{2} \sum_{a,b=H,A} \mathcal{P}_a \mathcal{P}_b^* \, \frac{\Gamma_{ab}(\sqrt{s})}{\sqrt{s}}, \tag{7}$$

where the relative velocity v is taken to be $2\beta = 2\sqrt{1-4m_\chi^2/s}$, and the production amplitudes $\mathcal{P}_{a,b}$ and the transition decay widths Γ_{ab} are defined as

$$\begin{aligned} \mathcal{P}_a &= \sum_{i=2,3} \sum_{b=H,A} C_{ia} \, \Pi_i \, C_{ib} \, P(\chi\chi \to a), \\ \Gamma_{ab} &= \frac{1}{2\sqrt{s}} \sum_F \oint d\Phi_F \, D(a \to F) D^*(b \to F), \end{aligned} \tag{8}$$

with the Higgs propagators $\Pi_i = 1/(s - M_{H_i}^2 + i M_i \Gamma_{H_i})$. Here, $P(\chi\chi \to H, A)$ are the $\chi\chi \to H, A$ production amplitudes, determined by the couplings (6), and $D(H, A \to F)$ the $H, A \to F$ decay amplitudes, for any kinematically and dynamically allowed decay mode F. Evaluating $J(x_f)$ in Eq. (5) with the event rate (7) and inserting its value into Eq. (4) yields the present neutralino relic density.

Although it is possible to calculate the masses and (transition) decay widths of the heavy Higgs bosons fully, we estimate them in the present work with a few approximations, which are reliable in the Higgs decoupling limit. In general, the light Higgs boson, the fermions and electroweak gauge bosons, and in supersymmetric theories, gauginos, higgsinos and scalar states may contribute to the loops in the complex mass matrix. In the decoupling limit, the couplings of the heavy Higgs bosons to gauge bosons and their superpartners are suppressed. Assuming all the other supersymmetric particles to be suppressed either by couplings or by phase space, we consider only loops by the LSP neutralino, the light Higgs boson and the top/bottom quark for the absorptive parts as characteristic examples; loops from other (s)particles could be treated in the same way of course.

In order to demonstrate the effect of the CP–violating H/A mixing on the neutralino relic density in the MSSM–CP numerically, we adopt a typical set of parameters[‖],

$$M_S = 0.5 \text{ TeV}, \quad |A_t| = 1.0 \text{ TeV}, \quad \mu = 2.0 \text{ TeV}; \quad \tan\beta = 5, \tag{9}$$

while varying the pseudoscalar mass M_A, the SU(2) gaugino mass M_2, and the phase Φ_A of the trilinear term A_t, and taking $M_1 \simeq 0.5 M_2$. [By reparameterization of the fields, M_2

[‖]Analyses of electric dipole moments show that the phase of μ is quite small, unless sfermions are very heavy [15]; therefore its phase is set zero in our numerical demonstration.

is set real and positive.] For such a large μ compared to M_2, the LSP is almost bino–like and its mass is close to M_1.

The Φ_A dependence of the H/A mixing parameter X and the heavy Higgs masses and are displayed in Figs. 1(a) and (b), respectively, for $M_{2,A} = 0.5$ TeV.** The two–state system in the MSSM–CP shows a very sharp resonance CP–violating mixing, purely imaginary near $\Phi_A = 0.09\pi$ and $\Phi_A = 0.67\pi$. We note that the mass shift is indeed enhanced by more than an order of magnitude if the CP–violating phase rises to non–zero values, reaching a maximal value of the mass difference ~ 24 GeV. As a result, the two mass eigenstates become clearly distinguishable, incorporating significant admixtures of CP–even and CP–odd components mutually in the wave functions.

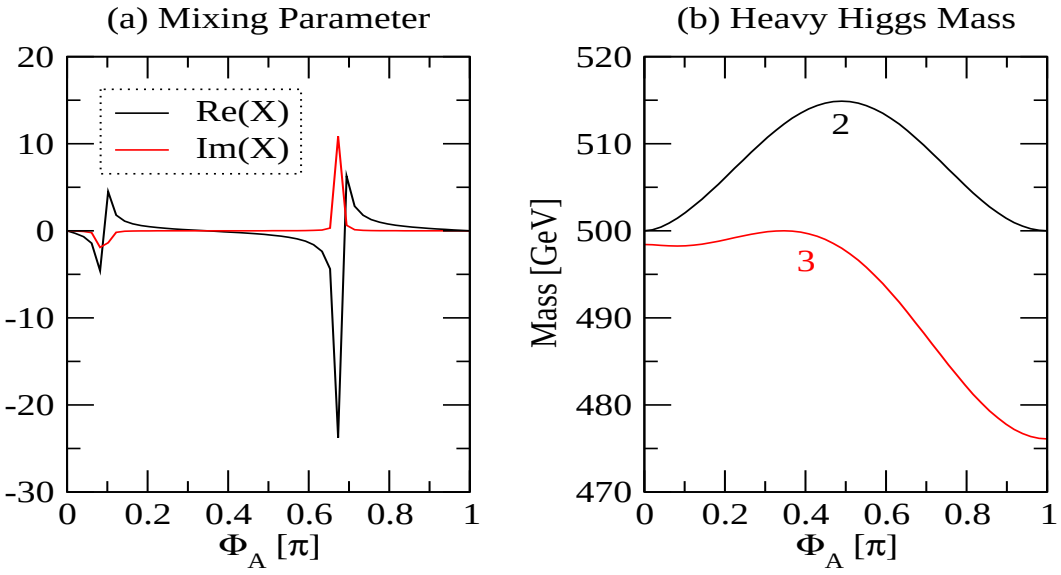

Figure 1: *The Φ_A dependence of (a) the real (black) and imaginary (red) parts of the mixing parameter X and (b) the heavy Higgs boson masses, M_{H_2} (black) and M_{H_3} (red). M_2 and M_A are set to 500 GeV. Note that $\Re\mathrm{e}/\Im\mathrm{m}X(2\pi - \Phi_A) = +\Re\mathrm{e}/-\Im\mathrm{m}X(\Phi_A)$ and the masses and widths are symmetric about $\Phi_A = \pi$.*

The left panel of Fig. 2 shows the allowed space of the phase Φ_A and the normalized mass difference $(M_A - 2m_\chi)/2m_\chi$ for the range (1). Here we have set M_2 to 0.5 TeV and

**With one common phase Φ_A, the complex mixing parameter X obeys the relation $X(2\pi-\Phi_A) = X(\Phi_A)$ so that all CP–even quantities are symmetric when switching from Φ_A to $2\pi-\Phi_A$. Therefore we can restrict the discussion to the range $0 \leq \Phi_A \leq \pi$.

have scanned the parameter space where $450\,\mathrm{GeV} \leq M_A \leq 550\,\mathrm{GeV}$ and $0 \leq \Phi_A \leq \pi$. The allowed region for $\pi \leq \Phi_A \leq 2\pi$ is simply obtained by reflecting the allowed region for $0 \leq \Phi_A \leq \pi$ with respect to $\Phi_A = \pi$. The green strip is for the range (1) and the blue region for $\Omega h^2 < 0.095$. In the other remaining region, we have $\Omega h^2 > 0.129$. One can clearly see that (i) the neutralino relic density is indeed greatly suppressed for $M_A \sim 2m_\chi$ due to the Higgs resonances and the detailed prediction for the relic density depends strongly on the value of the phase Φ_A as well as the mass difference between M_A and $2m_\chi$.

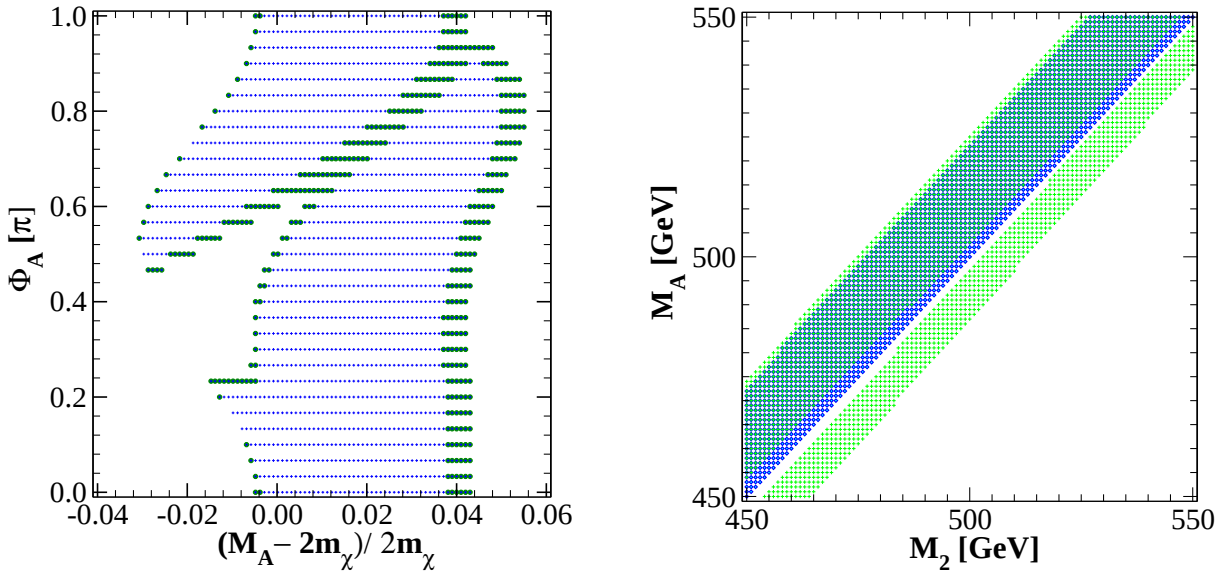

Figure 2: *Left panel: The allowed phase space of the CP phase Φ_A and the normalized mass difference $(M_A - 2m_\chi)/2m_\chi$ for the range (1). The green area is for the range (1), but the blue area is for the enlarged range with the lower bound ignored. Right panel: The allowed region of the (M_2, M_A) plane for the bound $\Omega h^2 < 0.129$ in the CP–invariant case with $\Phi_A = 0$ (a blue strip) and CP–noninvariant case with $\Phi_A = 0.55\pi$ (two green strips). The values of the other relevant parameters are given in the text.*

The right panel of Fig. 2 shows the allowed regions of the (M_2, M_A) plane for $\Omega h^2 < 0.129$ in the CP–invariant case (one blue strip) with $\Phi_A = 0$ and in the CP–noninvariant case with $\Phi_A \simeq 0.55\pi$ (two green strips). Clearly, in order to satisfy the relic density constraint, the LSP mass, which is approximately $0.5M_2$, should be close to half of the Higgs masses. In the CP–invariant case only the CP–odd Higgs boson A is active for the Higgs funnel mechanism and so only one allowed strip with its width of about 20 GeV is developed. In contrast, in the CP–noninvariant case with $\Phi_A = 0.55\pi$, both of the heavy Higgs bosons

become active for the funnel mechanism, leading to two strips; one strip is almost identical to the strip in the CP–invariant case, but the other is newly developed as the H_3 state, which is purely CP–even in the CP–invariant case, has a significant CP–odd component due to the CP–violating Higgs mixing. The combined width of two strips is widened due to the enlarged mass splitting between two mass eigenstates in the CP–noninvariant case.[††]

To summarize. We have examined the effect of the CP–violating H/A mixing on the LSP annihilation cross section in the Higgs decoupling limit. By a simple analysis with a specific parameter set (9) we have demonstrated that the CP–violating mixing can modify the profile of the LSP relic density *considerably* in the heavy Higgs funnel with the LSP mass close to half of the Higgs masses. Therefore, in order to elucidate the Higgs funnel mechanism through high–energy experiments on the supersymmetric particles, it is necessary to determine with good accuracy the complex mixing angle between two Higgs states in addition to the LSP and heavy Higgs boson masses and couplings [18].

Acknowledgments

The authors are grateful to P.M. Zerwas for his valuable comments and suggestions. The work of SYC was supported partially by KOSEF through CHEP at Kyungpook National University and by the Korea Research Foundation Grant by the Korean Government (MOEHRD) (KRF–2006–013–C00097) and the work of YGK was supported partially by the KRF Grant funded by the Korean Government (KRF–2005–201–C00006) and by the KOSEF Grant (KOSEF R01–2005–000–10404–0).

References

[1] WMAP Collaboration, D.N. Spergel *et al.*, Astrophys. J. Suppl. Ser. **148** (2003) 175.

[2] SDSS Collaboration, M. Tegmark *et al.*, Phys. Rev. D **69** (2004) 103501.

[3] For a recent review, we refer to D.J.H. Chung, L.L. Everett, G.L. Kane, S.F. King, J.D. Lykken and L.T. Wang, Phys. Rept. **407** (2005) 1.

[4] H. Goldberg, Phys. Rev. Lett. **50** (1983) 1419; J. Ellis, J. Hagelin, D. Nanopoulos and M. Srednicki, Phys. Lett. B **127** (1983) 233; J. Ellis, J. Hagelin, D. Nanopoulos, K. Olive and M. Srednicki, Nucl. Phys. B **238** (1984) 453.

[††]Two Higgs–boson masses are approximately $M_{H_2} \approx M_A + 15\,\mathrm{GeV}$ and $M_{H_3} \approx M_A - 5\,\mathrm{GeV}$, respectively.

[5] For reviews on supersymmetric dark matter, see, for example, G. Jungman, M. Kamionkowski and K. Griest, Phys. Rept. **267** (1996) 195; G. Bertane, D. Hooper and J. Silk, Phys. Rept. **405** (2005) 279.

[6] H. Baer, C. Balazs, A. Belyaev, J.M. Mizukoshi, X. Tata and Y. Wang, JHEP **0207** (2002) 050; H. Baer and C. Balazs, JCAP **0305** (2003) 006; U. Chattopadhyay , A. Corsetti and P. Nath, Phys. Rev. D **68** (2003) 035005; J. Ellis, K.A. Olive, Y. Santoso and V.C. Spanos, Phys. Lett. B **565** (2003) 176; M. Battaglia, A. De Roeck, J.R. Ellis, F. Gianotti, K.A. Olive and L. Pape, Eur. Phys. J. C **33** (2004) 273; R. Arnowitt and B. Dutta and B. Hu, hep–ph/0310103; J.R. Ellis, K.A. Olive, Y. Santoso and V.C. Spanos, Phys. Rev. D **69** (2004) 095004; M.E. Gomez, T. Ibrahim, P. Nath and S. Skadhauge, Phys. Rev. D **70** (2004) 035014; J.R. Ellis, S. Heinemeyer, K.A. Olive and G. Weiglein, JHEP **0502** (2005) 013; A. Djouadi, M. Drees and J.L. Kneur, JHEP **0108** (2001) 055; C. Boehm, A. Djouadi and M. Drees, Phys. Rev. D **62** (2000) 035012; R. Arnowitt, B. Dutta and Y. Santoso, Nucl. Phys. B **606** (2001) 59; J.R. Ellis, K.A. Olive and Y. Santoso, Astropart. Phys. **18** (2003) 395; A. Djouadi, M. Drees and J.L. Kneur, Phys. Lett. B **624** (2005) 60.

[7] A. Pilaftsis, Phys. Rev. D **58** (1998) 096010; Phys. Lett. B 435 (1998) 88; ; D.A. Demir, Phys. Rev. D **60** (1999) 055006; S.Y. Choi, M. Drees and J.S. Lee, Phys. Lett. B **481** (2000) 57; M. Carena, J. Ellis, A. Pilaftsis and C.E.M. Wagner, Nucl. Phys. B **586** (2000) 92; Phys. Lett. B **495** (2000) 155; T. Ibrahim and P. Nath, Phys. Rev. D **63** (2001) 035009; S.W. Ham, S.K. Oh, E.J. Yoo, C.M. Kim and D. Son, Phys. Rev. D **68** (2003) 055003.

[8] A. Pilaftsis and C.E.M. Wagner, Nucl. Phys. B **553** (1999) 3.

[9] A. Pilaftsis, Nucl. Phys. B **504** (1997) 61; J.R. Ellis, J.S. Lee and A. Pilaftsis, Phys. Rev. D **70** (2004) 075010; J.S. Lee, hep-ph/0409020; J.R. Ellis, J.S. Lee and A. Pilaftsis, Nucl. Phys. B **718** (2005) 247; J.R. Ellis, J.S. Lee and A. Pilaftsis, Phys. Rev. D **71** (2005) 075007.

[10] S.Y. Choi, J. Kalinowski, Y. Liao and P.M. Zerwas, Eur. Phys. J. C **40** (2005) 555.

[11] J.F. Gunion, H.E. Haber and J. Kalinowski, in preparation; For a different point of view see also I.F. Ginzburg, M. Krawczyk and P. Osland, hep–ph/0211371; M.N. Dubinin and A.V. Semenov, Eur. Phys. J. C **28** (2003) 233.

[12] P. Gondolo and K. Freese, JHEP **0207** (2002) 052; S.Y. Choi, S.-C. Park, J.H. Jang and H.S. Song, Phys. Rev. D **64** (2001) 015006; T. Nihei and M. Sagawa, Phys. Rev. D **70** (2004) 055011; T. Nihei, hep–ph/0508285; C. Balazs, M. Carena, A. Menon, D.E. Morrissey and C.E.M. Wagner, Phys. Rev. D **71** (2005) 075002.

[13] J.F. Gunion and H.E. Haber, Phys. Rev. D **67** (2003) 075019.

[14] S. Güsken, J.H. Kühn and P.M. Zerwas, Nucl. Phys. **B** 262 (1985) 393.

[15] Y. Kizukuri and N. Oshimo, Phys. Lett. **B249** (1990) 449; I. Ibrahim and P. Nath, Phys. Rev. D **58** (1998) 111301; **60** (1998) 099902(E); **61** (2000) 093004; M. Brhlik, G.J. Good and G.L. Kane, Phys. Rev. D **59** (1999) 115004; R. Arnowitt, B. Dutta and Y. Santos, Phys. Rev. D **64** (2000) 113010; V. Barger, T. Falk, T. Han, J. Jiang, T. Li and T. Plehn, Phys. Rev. D **64** (2001) 056007; S.Y. Choi, M. Drees and B. Gaissmaier, Phys. Rev. D **70** (2004) 014010.

[16] K. Griest and D. Seckel, Phys. Rev. D **43** (1991) 3191.

[17] For details on the neutralino mixing in the MSSM, see, for instance, S.Y. Choi, J. Kalinowski, G. Moortgat–Pick and P.M. Zerwas, Eur. Phys. J. C **22** (2001) 563; **23** (2002) 769.

[18] K. Desch, T. Klimkovich, T. Kuhl and A. Raspereza, hep–ph/0406229; M. Battaglia, hep-ph/0410123; M. Battaglia and M.E. Peskin, hep–ph/0509135.