CERN-TH/2002–008
DCPT/02/12
FTUAM 02/02
IFUM-76/FT
IFT-UAM/CSIC-02-03
IPPP/02/06
February 2002

Relic Neutralino Density in Scenarios with Intermediate Unification Scale

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Abstract

We analyse the relic neutralino density in supersymmetric models with an intermediate unification scale. In particular, we present concrete cosmological scenarios where the reheating temperature is as small as $\mathcal{O}(1 - 1000 \text{ MeV})$. When this temperature is associated to the decay of moduli fields producing neutralinos, we show that the relic abundance increases considerably with respect to the standard thermal production. Thus the neutralino becomes a good dark matter candidate with $0.1 \lesssim \Omega h^2 \lesssim 0.3$, even for regions of the parameter space where large neutralino-nucleon cross sections, compatible with current dark matter experiments, are present. This is obtained for intermediate scales $M_I \sim 10^{11} - 10^{14}$ GeV, and moduli masses $m_\phi \sim 100 - 1000$ GeV. On the other hand, when the above temperature is associated to the decay of an inflaton field, the relic abundance is too small.

1 Introduction

As it is well known, the lightest neutralino, $\tilde{\chi}_1^0$, is a weakly interacting massive particle (WIMP), and therefore a very interesting candidate for dark matter in the universe. In fact, many experimental efforts are being carried out in order to detect WIMPs through elastic scattering with nuclei in a detector [1]. In this sense the theoretical analysis of the neutralino–nucleus cross section $\sigma_{\tilde{\chi}_1^0-N}$ is very important. In particular, these analyses in the context of the minimal supersymmetric standard model (MSSM) are usually performed assuming the unification scale $M_{GUT} \approx 10^{16}$ GeV for the running of the universal soft supersymmetry (SUSY)–breaking terms. However, it was pointed out recently [2] that this cross section is very sensitive to the variation of the unification scale. For instance, by taking an intermediate unification scale $M_I \approx 10^{10-12}$ GeV the cross section increases substantially, being compatible for large regions of the parameter space of the MSSM with the sensitivity of current dark matter experiments $\sigma_{\tilde{\chi}_1^0-N} \approx 10^{-7}$–$10^{-6}$ GeV^{-2}, for $\tan\beta \gtrsim 3$ and $m_{\tilde{\chi}_1^0} \approx 100$ GeV. For larger values of the scale, as e.g. $M_I = 10^{14}$ GeV, a similar result is obtained for $\tan\beta \gtrsim 10$. Explicit scenarios with intermediate scales, arising in D-brane constructions from type I strings, were analysed in ref. [3]. Although compatibility with the experiments may also be obtained within the usual MSSM scenario with the scale $\approx 10^{16}$ GeV, it requires large values of $\tan\beta$ ($\tan\beta \gtrsim 20$) [4]-[6] or a specific non–universal structure of the soft terms [4, 5, 7, 8].

In all the above works the relic neutralino density was also discussed. In these scenarios with a large cross section in some regions of the parameter space, generically $\Omega_{\tilde{\chi}_1^0} h^2 \lesssim 0.01$. Of course, this might be a potential problem for the consistency of those regions given the observational bounds[1] $0.1 \lesssim \Omega_{\tilde{\chi}_1^0} h^2 \lesssim 0.3$.

This result is obtained because in the usual early–universe model thermal production of neutralinos gives rise to $\Omega_{\tilde{\chi}_1^0} h^2 \propto 1/\langle \sigma_{\tilde{\chi}_1^0}^{\rm ann} v \rangle$, where $\sigma_{\tilde{\chi}_1^0}^{\rm ann}$ is the cross section for annihilation of a pair of neutralinos, v is the relative velocity between the two neutralinos, and $\langle .. \rangle$ denotes thermal averaging. Therefore, in this scheme the relic density is inversely proportional to the annihilation cross section. Let us recall that crossing arguments, when the main annihilation channel is into quarks, ensure that the cross sections of annihilation and scattering with nucleons are similar. Thus a large scattering cross section $\sigma_{\tilde{\chi}_1^0-N}$ leads generically to a large annihilation cross section $\sigma_{\tilde{\chi}_1^0}^{\rm ann}$, and as a consequence to a small relic density.

However, it is important to remark that this result depends on assumptions about the evolution of the early universe. In principle, different cosmological scenarios might give rise to different results. To address this question is precisely the aim of this paper. We will study the relic density in the context of some non-standard cosmological scenarios. In particular, we will show that, when intermediate scales are present, results different from the usual ones summarised above may be produced. This is because a low reheating temperature, below the freeze-out temperature, can be obtained. We will see that, in the case of one of the scenarios, values of the relic density within the observational bounds

[1] It is worth noticing, however, that more conservative lower bounds, $\Omega_{\tilde{\chi}_1^0} h^2 \approx 0.01$, have also been quoted in the literature. For a brief discussion on this issue see e.g. ref. [8] and references therein.

are possible, even for regions of the parameter space with a large neutralino–nucleus cross section $\sigma_{\tilde{\chi}_1^0-N} \approx 10^{-7}$ GeV^{-2}.

The content of the paper is as follows. In Section 2 we will briefly review the usual cosmological scenario where thermal production of neutralinos is assumed. Several well-known formulas will be explicitly written since we will use them in the discussions of the next sections. Then, in Section 3, we will discuss the modifications introduced in the relic density analysis by considering non–standard cosmological scenarios in the case of intermediate scales. In particular, we will study the situation when an inflation or a modulus field produce low reheating temperatures, close to the nucleosynthesis one. Finally, the conclusions are left for Section 4.

2 Thermal production of neutralinos

Let us briefly review the standard computation of the cosmological abundance of neutralinos [9]. Neutralinos were in thermal equilibrium with the standard model particles in the early universe, and decoupled when they were non-relativistic. The process was the following. When the temperature T of the universe was larger than the mass of the neutralino, the number density of neutralinos and photons was roughly the same, $n^{eq}_{\tilde{\chi}_1^0} \propto T^3$, and the neutralino was annihilating with its own antiparticle into lighter particles and vice versa. However, shortly after the temperature dropped below the mass of the neutralino, $m_{\tilde{\chi}_1^0}$, its number density dropped exponentially, $n^{eq}_{\tilde{\chi}_1^0} \propto e^{-m_{\tilde{\chi}_1^0}/T}$, because only a small fraction of the light particles mentioned above had sufficient kinetic energy to create neutralinos. As a consequence, the neutralino annihilation rate $\Gamma_{\tilde{\chi}_1^0} = \langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle n_{\tilde{\chi}_1^0}$ dropped below the expansion rate of the universe, $\Gamma_{\tilde{\chi}_1^0} \lesssim H$, where H is the Hubble expansion rate. At this point neutralinos came away, they could not annihilate, and their density is the same since then. This can be obtained using the Boltzmann equation:

$$\frac{dn_{\tilde{\chi}_1^0}}{dt} + 3Hn_{\tilde{\chi}_1^0} = -\langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle \left[(n_{\tilde{\chi}_1^0})^2 - (n^{eq}_{\tilde{\chi}_1^0})^2 \right] \ . \tag{1}$$

One can discuss qualitatively the solution using the freeze-out condition $\Gamma_{\tilde{\chi}_1^0} = \langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle_F n_{\tilde{\chi}_1^0} = H$. Then $\Omega_{\tilde{\chi}_1^0} h^2 = (\rho_{\tilde{\chi}_1^0}/\rho_c) h^2$, where $\rho_{\tilde{\chi}_1^0}$ is the current neutralino mass density and ρ_c is the critical density, turns out to be

$$\Omega_{\tilde{\chi}_1^0} h^2 = \frac{m_{\tilde{\chi}_1^0}\ H}{(2\pi^2/45)\ g_\star(T_F)\ T_F^3\ \langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle_F}\ \frac{h^2}{\rho_c/s_0} = \frac{m_{\tilde{\chi}_1^0}\ (45/6\pi\sqrt{10})}{M_P\ g_\star^{1/2}(T_F)\ T_F\ \langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle_F}\ \frac{h^2}{\rho_c/s_0}\ , \tag{2}$$

where $g_*(T)$ is the effective number of degrees of freedom at temperature T (e.g. including all the standard-model degrees of freedom one gets $g_* = 106.75$), T_F is the freeze-out temperature, and s_0 is the current entropy density. In the second expression we have used the fact that $H = (\pi^2 g_*(T)/90)^{1/2}\, T^2 M_P^{-1}$, with $M_P = M_{Planck}/\sqrt{8\pi} \simeq 2.4 \times 10^{18}$ GeV the reduced Planck mass. Taking into account the current value $\rho_c/s_0 \simeq 3.6 \times 10^{-9}$ GeV$\times h^2$, and the typical freeze-out temperature $T_F \simeq m_{\tilde{\chi}_1^0}/20$, one can write the above expression as

$$\Omega_{\tilde{\chi}_1^0} h^2 \simeq 1.7 \times 10^{-10} \left(\frac{1\ GeV^{-2}}{\langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle} \right) \left(\frac{100}{g_*(T_F)} \right)^{1/2} \quad . \tag{3}$$

Since neutralinos freeze–out at $T_F \simeq m_{\tilde{\chi}_1^0}/20 << m_{\tilde{\chi}_1^0}$, they are non–relativistic and therefore the averaged annihilation cross section can be expanded as follows:

$$\langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle = \alpha_s + \alpha_p \langle v^2 \rangle \ , \tag{4}$$

where α_s describes the s-wave annihilation and α_p describes both s– and p– wave annihilation. Then, eq. (3) can alternatively be written as:

$$\Omega_{\tilde{\chi}_1^0} h^2 \ \ \simeq \ \ 8.8 \times 10^{-11} \frac{GeV^{-2}}{g_*^{1/2}(T_F) \left(\alpha_s/x_F + 3\alpha_p/x_F^2 \right)} \ , \tag{5}$$

where $x_F \equiv m_\chi/T_F$.

As it is well known, in most of the parameter space of the MSSM the neutralino is mainly pure bino, and as a consequence it will mainly annihilate into lepton pairs through t–channel exchange of right–handed sleptons. The p–wave dominant cross section is given by [10, 11]

$$\langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle \simeq 8\pi\alpha'^2 \frac{1}{m^2_{\tilde{\chi}_1^0}} \frac{1}{\left(1 + x_{\tilde{l}_R}\right)^2} \ \langle v^2 \rangle \ , \tag{6}$$

where $x_{\tilde{l}_R} \equiv m^2_{\tilde{l}_R}/m^2_{\tilde{\chi}_1^0}$ and α' is the coupling constant for the $U(1)_Y$ interaction. Taking $m_{\tilde{l}_R} \sim m_{\tilde{\chi}_1^0} \sim 100$ GeV, $\langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle$ in eq. (6) becomes of the order of 10^{-9} GeV^{-2} or smaller. Using eq. (3) an interesting relic abundance, $\Omega_{\tilde{\chi}_1^0} h^2 \gtrsim 0.1$, is obtained.

However, in the special regions mentioned in the Introduction, with non-universality and/or large $\tan\beta$, the lightest neutralino may have an important Higgsino component, producing a larger cross section. This is also the case of scenarios with an intermediate unification scale. The upper bound for the annihilation cross section, obtained when the neutralino is Higgsino–like, is given by [10]

$$\langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle \simeq \frac{\pi\alpha_2^2}{2} \frac{1}{m^2_{\tilde{\chi}_1^0}} \frac{\left(1 - x_W\right)^{3/2}}{\left(2 - x_W\right)^2} \ , \tag{7}$$

where $x_W \equiv m_W^2/m^2_{\tilde{\chi}_1^0}$ and α_2 is the coupling constant for the $SU(2)_L$ interaction. Here one is considering that the Higgsino dominantly annihilates into W-boson pairs. Since now $\langle \sigma^{ann}_{\tilde{\chi}_1^0} v \rangle$ in eq. (7) is of the order of 10^{-8} GeV^{-2}, the relic abundance given by eq. (3) turns out to be small, $\Omega_{\tilde{\chi}_1^0} h^2 \approx 0.01$, as expected.

3 Non–standard cosmological scenarios

In the standard computation reviewed in the previous section, one is tacitly assuming that the radiation–dominated era is the result of a reheat process in the early universe, where the reheating temperature T_{RH} is very large, in particular $T_{RH} >> T_F \sim 10$ GeV. The scalar field ϕ, whose decay leads to reheating, is usually assumed to be the inflaton field. One can estimate the reheating temperature as a function of the decay width Γ_ϕ as [12]

$$T_{RH} = \left(\frac{90}{\pi^2 g_*(T_{RH})}\right)^{1/4} (\Gamma_\phi M_P)^{1/2} \ . \tag{8}$$

However, the only constraint on the reheating temperature is $T_{RH} \gtrsim 1$ MeV in order not to affect the successful predictions of big–bang nucleosynthesis. This allows in principle to consider cosmological scenarios with a low reheating temperature [11], $T_{RH} < T_F$. On the other hand, the reheating process can also be associated with the decay of moduli fields, as e.g. those appearing in string theory. Thus the relic abundance could receive contributions from this source [13].

In what follows, we will show that scenarios with intermediate unification scales are explicit examples for the two non–standard cosmological possibilities mentioned above, namely, decay of the inflaton and moduli fields producing a low reheating temperature. For this analysis eq. (8) is still valid using Γ_ϕ as given by the corresponding scenario. This is also true for relation $\Omega_{\tilde{\chi}_1^0} h^2 \propto 1/\langle\sigma_{\tilde{\chi}_1^0}^{\rm ann} v\rangle$ in the case of neutralino production through modulus decay. Notice that in this case $\Omega_{\tilde{\chi}_1^0} \propto 1/T_F$, as shown in eq. (2), and therefore if this mechanism produces a temperature smaller than the typical $T_F \simeq m_{\tilde{\chi}_1^0}/20$, the value of the relic neutralino density will be increased [2]. We will see below that temperatures as required can be obtained in scenarios with intermediate scales. In ref. [13] this mechanism was applied in order to obtain reasonable values of the relic wino density in anomaly-mediated SUSY breaking scenarios, using $m_\phi \approx 100$ TeV. In our case, standard masses in supergravity scenarios, $m_\phi \approx 1$ TeV, will be used.

On the other hand, in the scenario where the low reheating temperature is obtained through an inflaton field, the result for the relic density is quite different from the usual one with large T_{RH}. In fact, in certain cases, the usual relation $\Omega_{\tilde{\chi}_1^0} h^2 \propto 1/\langle\sigma_{\tilde{\chi}_1^0}^{\rm ann} v\rangle$ is not even valid, and the relic abundance may well be proportional to the annihilation cross section [11]. We will see below, however, that the relic abundance will not increase when intermediate scales are considered.

3.1 Inflation scenario

Let us consider for example the SUSY hybrid inflation scenario studied in ref. [15]. There, the inflaton decay width can be computed with the result

$$\Gamma_\phi = \frac{1}{8\pi}\left(\frac{m_f}{\langle\phi\rangle}\right)^2 m_\phi \ , \tag{9}$$

[2]For another alternative cosmological scenario with the potential of increasing the relic density see ref. [14], where the decay of cosmic strings producing neutralinos is considered.

where $\langle\phi\rangle$ is the vacuum expectation value of the inflaton field, which is of the order of the unification scale, m_ϕ is the inflaton mass, and m_f is the mass of the particle f that the inflaton decay to (in this case a right handed neutrino or sneutrino). Obviously, m_f should be smaller than the inflaton mass to allow for the decay $\phi \to ff$.

Now, using eqs. (8) and (9) one obtains the following reheating temperature:

$$T_{RH} = 1.7 \times 10^9\ GeV \left(\frac{100\ GeV}{\langle\phi\rangle}\right) \left(\frac{m_f}{100\ GeV}\right) \left(\frac{m_\phi}{100\ GeV}\right)^{1/2} \left(\frac{100}{g_*(T_{RH})}\right)^{1/4} . \quad (10)$$

In ref. [16] it has been shown that an intermediate unification scale of the order of $M_I \sim 10^{11}$ GeV is favoured by inflation. Then, recalling that the inflaton mass is constrained by $m_\phi \lesssim M_I^2/M_{Planck}$, we obtain $m_\phi \sim 10^2$ GeV. From Eq.(10) we find that $T_{RH} \sim 1$ GeV, since now $\langle\phi\rangle \sim 10^{11}$ GeV. This reheating temperature is lower than the typical freeze-out temperature $T_F \simeq m_\chi/20$. Notice that in the standard GUT scenario discussed in ref. [15], one has $m_\phi \sim 10^{11}$ GeV, $\langle\phi\rangle \sim 10^{16}$ GeV, and therefore $T_{RH} \sim 10^9$ GeV.

As mentioned above, a detailed analysis of the relic density with a low reheating temperature has been carried out in ref. [11] by Giudice, Kolb and Riotto. They study two possible non-relativistic cases, depending on whether or not the dark-matter particles are in chemical equilibrium. In the first one they are never in equilibrium, either before or after reheating. In the second one the dark-matter particles reach chemical equilibrium, but then freeze out before the completion of the reheat process. These scenarios not only lead to different qualitative and quantitative predictions for the relic density, but also these predictions are quite different from the standard ones summarised in eq. (5).

In the case of non–equilibrium production, the number density of neutralinos $n_{\tilde{\chi}_1^0}$ is much smaller than $n^{eq}_{\tilde{\chi}_1^0}$, thus the relevant Boltzmann equations can be approximated and solved. One gets [11]

$$\Omega_{\tilde{\chi}_1^0} h^2 = 2.1 \times 10^4 \left(\frac{g}{2}\right)^2 \left(\frac{g_*(T_{RH})}{10}\right)^{3/2} \left(\frac{10}{g_*(T_*)}\right)^3 \frac{(10^3 T_{RH})^7}{m^5_{\tilde{\chi}_1^0}} \left(\alpha_s + \frac{\alpha_p}{4}\right) , \quad (11)$$

where g is the number of degrees of freedom of the neutralino and T_* is the temperature at which most of the neutralino production takes place, it is given by $T_* \sim 4m_{\tilde{\chi}_1^0}/15$. As we can see $\Omega_{\tilde{\chi}_1^0} h^2$ is proportional to the annihilation cross section, instead of being inversely proportional as in eq. (5). This raises the hope that the relic abundance could be increased in scenarios with intermediate scales where generically it is low. Unfortunately, the assumption that $n_{\tilde{\chi}_1^0} << n^{eq}_{\tilde{\chi}_1^0}$ leads to a severe constraint on the annihilation cross section [11]. Namely $\alpha_s < \bar{\alpha_s}$ and $\alpha_p < \bar{\alpha_p}$, where $\bar{\alpha_s}$ and $\bar{\alpha_p}$ are of the order of 10^{-15} GeV^{-2} and 10^{-14} GeV^{-2}, respectively, for $T_{RH} \sim 1$ GeV. Since we are interested in large cross sections, of the order of 10^{-7} GeV^{-2}, eq. (11) cannot be applied.

With a large annihilation cross section ($\alpha_s > \bar{\alpha_s}$ or $\alpha_p > \bar{\alpha_p}$), the neutralino reaches equilibrium before reheating as discussed in ref. [11], and its relic density is given by

$$\Omega_{\tilde{\chi}_1^0} h^2 = 2.3 \times 10^{-11} \frac{g_*^{1/2}(T_{RH})}{g_*(T_F)} \frac{T^3_{RH} GeV^{-2}}{m^3_{\tilde{\chi}_1^0} \left(\alpha_s x_F^{-4} + 4\alpha_p x_F^{-5}/5\right)} , \quad (12)$$

Now the relic density is again inversely proportional to the annihilation cross section as in eq. (5). Moreover, it has a further suppression because of the low reheating temperature $T_{RH} \sim 1$ GeV, and as a consequence we expect a result even worse than the one obtained in the standard computation discussed below eq. (5). Indeed, for a neutralino–nucleus cross section of the order of 10^{-7} GeV^{-2} we obtain $\Omega_{\tilde{\chi}_1^0} h^2 \approx 10^{-5}$.

3.2 Modulus–decay scenario

It has been assumed in the above computation that the relic abundance does not receive any contribution from other sources. However, as shown by Moroi and Randall [13], the production of neutralinos through moduli decay can modify those results.

Let us recall first that moduli fields are present for example in string theory [3]. Since moduli acquire masses through SUSY breaking effects, these masses, m_ϕ, are expected to be of the order of the gravitino mass, i.e. $\mathcal{O}(100 - 1000$ GeV). On the other hand, their couplings with the MSSM matter are suppressed by a high energy scale. Thus one can parameterise the moduli decay width as

$$\Gamma_\phi = \frac{1}{2\pi} \frac{m_\phi^3}{M_I^2} \ , \tag{13}$$

where we denote with M_I the effective suppression scale. Since we are interested in the analysis of scenarios with intermediate unification scales, we will consider the case of $M_I \approx 10^{11-14}$ GeV. An explicit example where this situation arises is the case of type I string constructions. There, twisted moduli are present with interactions suppressed by the string scale. As discussed in ref. [18] intermediate values for this scale can be obtained, and they have very interesting phenomenological implications [18, 19, 3, 16].

Using eqs. (8) and (13) one obtains the following reheating temperature:

$$T_{RH} = 1\ MeV \times \left(\frac{6 \times 10^{14}\ GeV}{M_I} \right) \left(\frac{m_\phi}{100\ GeV} \right)^{3/2} \left(\frac{10.75}{g_*(T_{RH})} \right)^{1/4} \ , \tag{14}$$

where note that $g_* = 10.75$ for $T \sim \mathcal{O}(1 - 10)$ MeV. This reheating temperature is shown as a function of the modulus mass in Fig. 1 for different values of M_I. The request that the modulus mass is larger than $\sim 100 - 500$ GeV in order to allow for kinematical decays into neutralinos of suitable mass $m_{\tilde{\chi}_1^0} \sim 50 - 200$ GeV, limits in practice the reheating temperature to be above ~ 3 GeV for the lowest scale on consideration $M_I = 10^{11}$ GeV. This has important consequences for the relic density computation. As discussed in the introduction of this section, we need a temperature smaller than the typical $T_F \simeq m_{\tilde{\chi}_1^0}/20 \sim 3 - 10$ GeV, in order to increase the relic neutralino density. This implies that the scale $M_I \sim 10^{11}$ GeV is in the border of validity. On the other hand, for larger values we can obtain very easily interesting reheating temperatures. For example, for $M_I = 10^{12}$ GeV we have $T_{RH} \gtrsim 0.3$ GeV. For the highest scale with an interesting phenomenological value of the neutralino–nucleus cross section, in the case of universality

[3]See other examples of moduli fields, for instance in GUTs, in ref. [17].

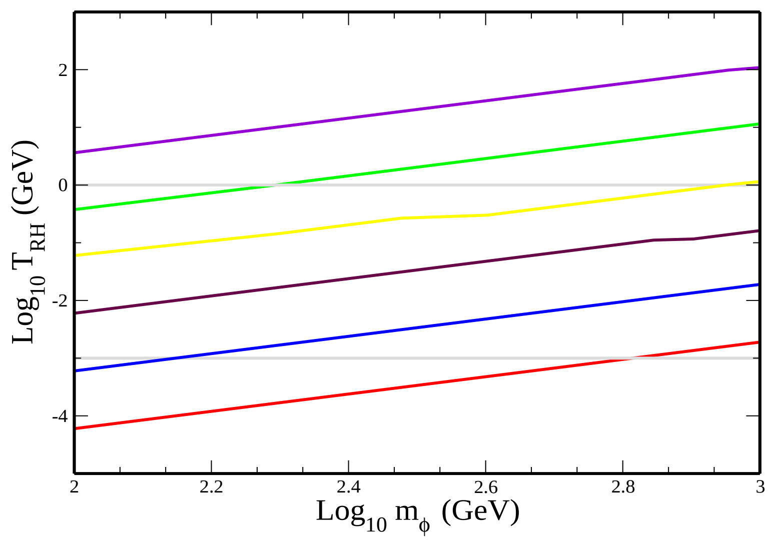

Figure 1: The reheating temperature T_{RH} as a function of the modulus mass m_ϕ. The six curves correspond, from top to bottom, to $M_I = 10^{11}, 10^{12}, 10^{13}, 10^{14}, 10^{15}, 10^{16}$ GeV, respectively. The region bounded by the horizontal lines corresponds to a reheating temperature larger than 1 MeV, a value close to the the nucleosynthesis limit, and smaller than 1 GeV, a value close to the freeze-out limit, as explained in the text.

and moderate $\tan\beta$ [2], $M_I = 10^{14}$ GeV, the lowest value of the reheating temperature corresponds to $T_{RH} \sim 6$ MeV. Larger values of the scale, $M_I \gtrsim 10^{15}$ GeV, producing also a large cross section, are possible in D-brane scenarios since non-universality in soft terms is generically present [3]. In this case the constraint $T_{RH} \gtrsim$ 1 MeV from nucleosynthesis can be translated into a constraint on the modulus mass $m_\phi \gtrsim 140$ GeV.

When considering the decay of the modulus field producing neutralinos, the evolution of the cosmological abundance of the latter becomes more complicated than in the usual thermal-production case reviewed in Section 2. Now one has to solve the coupled Boltzmann equations for the neutralino, the moduli field and the radiation [13, 20]:

$$\frac{dn_{\tilde{\chi}_1^0}}{dt} + 3Hn_{\tilde{\chi}_1^0} = \bar{N}_{\tilde{\chi}_1^0}\Gamma_\phi n_\phi - \langle\sigma^{ann}_{\tilde{\chi}_1^0}v\rangle \left[(n_{\tilde{\chi}_1^0})^2 - (n^{eq}_{\tilde{\chi}_1^0})^2\right] , \quad (15)$$

$$\frac{dn_\phi}{dt} + 3Hn_\phi = -\Gamma_\phi n_\phi , \quad (16)$$

$$\frac{d\rho_{rad}}{dt} + 4H\rho_{rad} = (m_\phi - \bar{N}_{\tilde{\chi}_1^0} m_{\tilde{\chi}_1^0})\Gamma_\phi n_\phi + 2m_{\tilde{\chi}_1^0}\langle\sigma^{ann}_{\tilde{\chi}_1^0}v\rangle \left[(n_{\tilde{\chi}_1^0})^2 - (n^{eq}_{\tilde{\chi}_1^0})^2\right] , \quad (17)$$

where $\bar{N}_{\tilde{\chi}_1^0}$ is the averaged number of neutralinos produced in the decay of one modulus field.

Let us discuss qualitatively the solution following the arguments used in ref. [13]. For a T_{RH} higher than T_F the relic density will roughly reproduce the usual result given by eq. (3). However, for the interesting case for us when T_{RH} is lower than T_F, neutralinos produced from modulus decay are never in chemical equilibrium, unlike the thermal production case reviewed in Section 2. As a consequence, its number density always decreases through pair annihilation. When the annihilation rate $\langle\sigma^{ann}_{\tilde{\chi}^0_1}v\rangle n_{\tilde{\chi}^0_1}$ drops below the expansion rate of the universe, H, the neutralino freezes out. Then the relic density can be estimated as [13]

$$\Omega_{\tilde{\chi}^0_1}h^2 = \frac{3\ m_{\tilde{\chi}^0_1}\ \Gamma_\phi}{2\ (2\pi^2/45)\ g_\star\ T^3_{RH}\ \langle\sigma^{ann}_{\tilde{\chi}^0_1}v\rangle}\ \frac{h^2}{\rho_c/s_0}\ . \tag{18}$$

This result is valid when there is a large number of neutralinos produced by the modulus decay. When the number is insufficient, they do not annihilate and therefore all the neutralinos survive. The result in this case is given by

$$\Omega_{\tilde{\chi}^0_1}h^2 = \frac{3\ \bar{N}_{\tilde{\chi}^0_1}\ m_\chi\ \Gamma^2_\phi\ M^2_P}{(2\pi^2/45)\ g_\star\ T^3_{RH}\ m_\phi}\ \frac{h^2}{\rho_c/s_0}\ . \tag{19}$$

The actual relic density is estimated [13] as the minimum of (18) and (19).

Now we can apply the above equations to our case with intermediate scales. Using eqs. (13) and (14), we can write expressions (18) and (19) as

$$\Omega_{\tilde{\chi}^0_1}h^2 = \left(\frac{M_I}{1.5\times 10^{20}\ GeV}\right)\ \left(\frac{1\ GeV^{-2}}{\langle\sigma^{ann}_{\tilde{\chi}^0_1}v\rangle}\right)\left(\frac{100\ GeV}{m_\phi}\right)^{3/2}\ \left(\frac{10.75}{g_*}\right)^{1/4}\ \left(\frac{m_{\tilde{\chi}^0_1}}{100\ GeV}\right)\ , \tag{20}$$

$$\Omega_{\tilde{\chi}^0_1}h^2 = \bar{N}_{\tilde{\chi}^0_1}\left(\frac{1.2\times 10^{20}\ GeV}{M_I}\right)\ \left(\frac{m_\phi}{100\ GeV}\right)^{1/2}\ \left(\frac{10.75}{g_*}\right)^{1/4}\ \left(\frac{m_{\tilde{\chi}^0_1}}{100\ GeV}\right)\ . \tag{21}$$

From these equations we can see that even with a large annihilation cross section, $\langle\sigma^{ann}_{\tilde{\chi}^0_1}v\rangle \sim 10^{-8}$ GeV^{-2}, we are able to obtain the cosmologically interesting value $\Omega_{\tilde{\chi}^0_1}h^2 \sim 1$. For example for $M_I \sim 10^{13}$ GeV we obtain it when $\bar{N}_{\tilde{\chi}^0_1} \sim 1$, using eq. (20). In Fig. 2 we show in more detail these results solving numerically the Boltzmann eqs. (15)–(17), for the large annihilation cross section introduced in eq. (7). There, the contours of constant relic neutralino density $\Omega_{\tilde{\chi}^0_1}h^2$ as a function of m_ϕ and $\bar{N}_{\tilde{\chi}^0_1}$ are shown, for fixed values of M_I and $m_{\tilde{\chi}^0_1}$. In particular, we consider the cases $M_I = 10^{12}, 10^{13}, 10^{14}$ GeV, with $m_{\tilde{\chi}^0_1} = 100$ GeV. The corresponding reheating temperatures can be obtained from Fig. 1. Note that whereas many values of $\bar{N}_{\tilde{\chi}^0_1}$ correspond to a satisfactory relic density for $M_I = 10^{12} - 10^{13}$ GeV, for the case $M_I = 10^{14}$ GeV only a small range works.

Let us finally remark that the numerical value of $\bar{N}_{\tilde{\chi}^0_1}$ is in general model dependent. This was discussed in the context of supergravity in ref. [13]. In this particular case both values $\bar{N}_{\tilde{\chi}^0_1} \sim 1$ and $\bar{N}_{\tilde{\chi}^0_1} \sim 10^{-3} - 10^{-4}$ are plausible, depending on the characteristics of the supergravity theory under consideration.

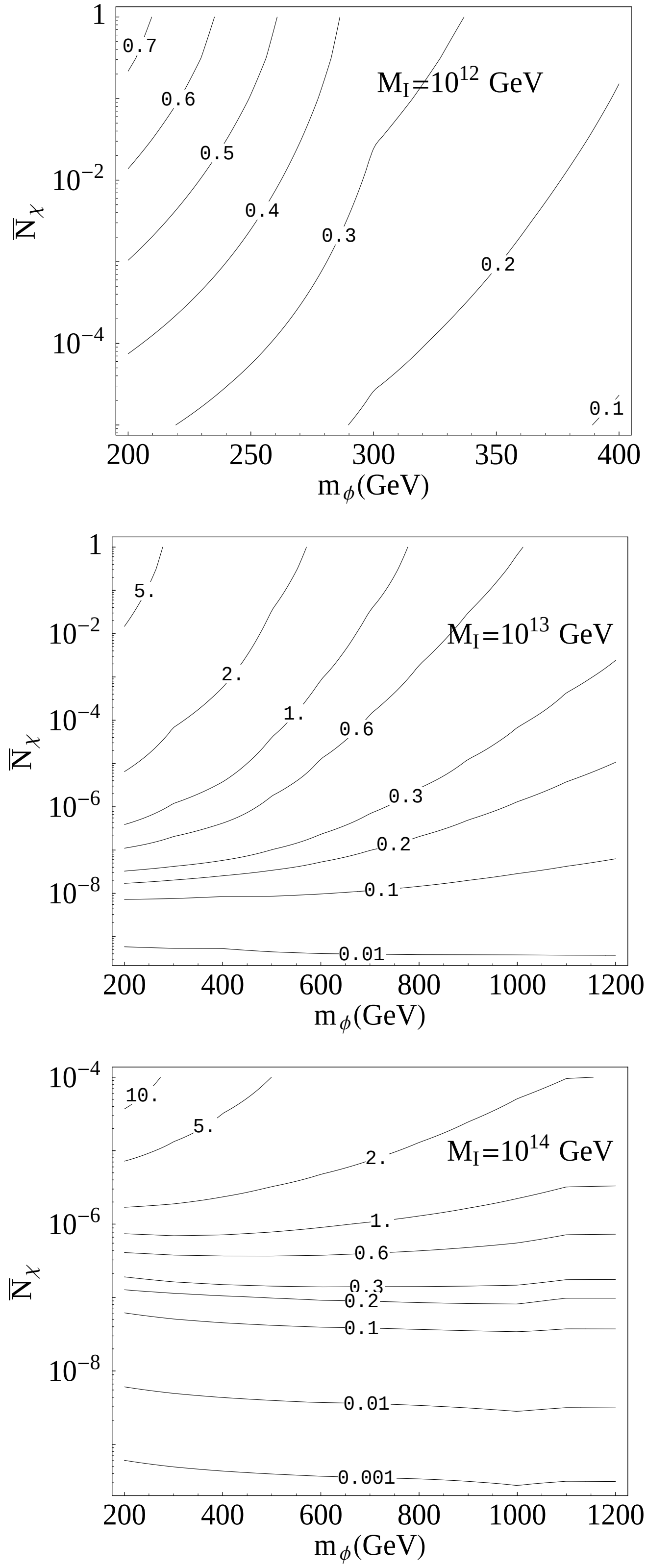

Figure 2: relic neutralino density $\Omega_{\tilde{\chi}_1^0} h^2$ contours as a function of m_ϕ and $\bar{N}_{\tilde{\chi}_1^0}$, for $m_{\tilde{\chi}_1^0} =$ 100 GeV and several possible values of the intermediate scale $M_I = 10^{12}, 10^{13}, 10^{14}$ GeV.

4 Conclusions

Current dark matter experiments are sensitive to a large neutralino-nucleus cross section, $\sigma_{\tilde{\chi}_1^0-N} \approx 10^{-7}$–$10^{-6}$ GeV^{-2}. There are regions in the parameter space of SUSY scenarios with an intermediate unification scale, where these cross sections can be obtained. However, in these regions, the standard computation of the relic abundance of neutralinos through thermal production in the early universe, would imply a small relic abundance, $\Omega_{\tilde{\chi}_1^0} h^2 \lesssim 0.01$.

We have analysed some alternatives to solve this potential problem. Let us recall that in the standard computation one is tacitly assuming that the reheating temperature is much larger than the freeze-out temperature, $T_F \sim 10$ GeV, and originated in an inflationary process. We have shown however that, when intermediate scales are present, reheating temperatures as small as $\mathcal{O}(1-1000\,\mathrm{MeV})$ are possible. Unfortunately, although the result for the relic abundance is modified, still $\Omega_{\tilde{\chi}_1^0} h^2$ is too small. On the other hand, when the above reheating temperatures are associated to the decay of moduli fields producing out of equilibrium neutralinos, the relic abundance increases considerably with respect to the standard production. Thus the neutralino becomes a good dark matter candidate with $0.1 \lesssim \Omega h^2 \lesssim 0.3$, for intermediate scales $M_I \sim 10^{11} - 10^{14}$ GeV, and moduli mass $m_\phi \sim 100 - 1000$ GeV.

Acknowledgements
We would like to thank G. Lazarides for useful discussions. The work of S. Khalil was supported by the PPARC. The work of C. Muñoz was supported in part by the Spanish Ministerio de Ciencia y Tecnología under contract FPA2000-0980, and the European Union under contract HPRN-CT-2000-00148. The work of E. Torrente-Lujan was supported in part by the Spanish Ministerio de Ciencia y Tecnología under contract FPA2000-0980, and by a MURST research grant.

References

[1] For a recent review see, S. Khalil and C. Muñoz, hep-ph/0110122, to appear in Contemporary Physics, and references therein.

[2] E. Gabrielli, S. Khalil, C. Muñoz and E. Torrente-Lujan, Phys. Rev. D63 (2001) 025008.

[3] D. Cerdeño, E. Gabrielli, S. Khalil, C. Muñoz and E. Torrente-Lujan, Nucl. Phys. B603 (2001) 231.

[4] A. Bottino, F. Donato, N. Fornengo and S. Scopel, Phys. Rev. D59 (1999) 095004.

[5] E. Accomando, R. Arnowitt, B. Dutta and Y. Santoso, Nucl. Phys. B585 (2000) 124.

[6] M.E. Gómez and J.D. Vergados, Phys. Lett. B512 (2001) 252.

[7] A. Corsetti and P. Nath, hep-ph/0003186.

[8] D.G. Cerdeño, S. Khalil and C. Muñoz, hep-ph/0105180, to appear in the Proceedings of CICHEP Conference (Cairo), Rinton Press.

[9] For a review see, G. Jungman, M. Kamionkowski and K. Griest, Phys. Rept. 267 (1996) 195, and references therein.

[10] K.A. Olive and M. Srednicki, Phys. Lett. B230 (1989) 78;
T. Moroi, M. Yamaguchi and T. Yanagida, Phys. Lett. B342 (1995) 105.

[11] G. Giudice, E.W. Kolb and A. Riotto, Phys. Rev. D64 (2001) 023508.

[12] E.W. Kolb and M.S. Turner, *The Early Universe*, (Addison–Wesley, Menlo Park, Ca., 1990).

[13] T. Moroi and L. Randall, Nucl. Phys. B570 (2000) 455.

[14] R. Jeannerot, X. Zhang and R. Brandenberger, JHEP 12 (1999) 003.

[15] R. Jeannerot, S. Khalil, G. Lazarides and Q. Shafi, JHEP 10 (2000) 012;
R. Jeannerot, S. Khalil, and G. Lazarides, Phys. Lett. B506 (2001) 344.

[16] N. Kaloper and A. Linde, Phys. Rev. D59 (1999) 101303.

[17] D.H. Lyth and E.D. Stewart, Phys. Rev. D53 (1996) 1784.

[18] K. Benakli, Phys. Rev. D60 (1999) 104002;
C. Burgess, L.E. Ibañez and F. Quevedo, Phys. Lett. B447 (1999) 257.

[19] S.A. Abel, B.C. Allanach, F. Quevedo, L.E. Ibáñez and M. Klein, JHEP 12 (2000) 026.

[20] D.J. Chung, E.W. Kolb and A. Riotto, Phys. Rev. D60 (1999) 063504.