Neutralino relic density from direct non-equilibrium production and intermediate scales

E. Torrente-Lujan

Dept. Física Teórica C-XI, Univ. Autónoma de Madrid, Cantoblanco, 28049 Madrid, Spain.

We review the calculation of the LSP relic density in alternative cosmological scenarios where contributions from direct decay production are important. We study supersymmetric models with intermediate unification scale. We find concrete scenarios where the reheating temperature is of order one GeV ($M_I \sim 10^{12}$) and below.

If the case that this reheating temperature is associated to the decay of oscillating moduli fields appearing in string theories, we show that the LSP relic density considerably increases with respect to the standard radiation-dominated case by the effect of the direct non-thermal production by the modulus field. The LSP can become a good dark matter candidate ($0.01 - 0.1 \lesssim \Omega h^2 \lesssim 0.3 - 1$) for $M_I \sim 10^{12} - 10^{14}$ GeV and $m_\phi \sim 1 - 10$ TeV.

1. Introduction

In a series of works [1,3,4], various implications for an intermediate unification scale in supersymmetric models have been considered. In particular, it was pointed out that the lightest neutralinos (LSP)–nucleon cross sections are very sensitive to the variation of this scale. For instance, for $M_I \sim 10^{12} - 10^{14}$ GeV the cross sections are enhanced and a large region of the parameter space of the minimal supersymmetric standard model (MSSM) is compatible with the sensitivity of the current and near future direct detection dark matter experiments (DAMA, CDMS, EDELWEISS). In fact, such results can also be obtained within the GUT unification scenario, i.e., $M_I \sim 10^{16}$ GeV, but it requires tuning the available parameter space: either large $\tan \beta$ ($\tan \beta \geq 25$) or a specific choice for non–universal soft SUSY breaking terms so that the value of $\mu$ term is reduced and hence a larger Higgsino components of the LSP is obtained, which is essential for enhancing the LSP–nucleon cross sections [1].

One problem is than in GUT unification scenarios with standard Big Bang (BB) cosmological expansion where the LSP is produced in a thermal environment dominated by radiation, such large cross section leads to a very small LSP relic density (outside the cosmological and astrophysical interesting region $0.1 \leq \Omega h^2 \leq 0.3$). In this scheme there is a generic anti-correlation between the LSP relic density and its quark or nucleon cross sections: one usually obtains

$$\Omega_\chi h^2 \simeq \frac{C}{\langle \sigma_{\text{ann}} \chi \rangle}$$

(1)

where $\sigma_{\text{ann}}$ is the annihilation cross section for neutralino pair, $v$ is the pair relative velocity, and $\langle \ldots \rangle$ denotes thermal averaging. Scattering ($\sigma_{\chi-q}$) and annihilation $\sigma_{\chi}^\text{ann}$ LSP cross sections can be related by a cross symmetry. Therefore, a large scattering cross section $\sigma_{\chi-q}$ leads generically to a large annihilation cross section and a small relic density as Eq. 1 suggest. Indeed, in these models, it was observed that LSP–proton cross sections of order $10^{-6}$ Pb correspond to too low relic densities ($\lesssim 0.005 - 0.01$).

We want to address here the relic density question: in the context of diverse non-standard cosmological scenarios, what is the effect of changing the unification scale from GUT scale ($\sim 10^{16}$ GeV) down to intermediate scale ($\sim 10^{12}$ GeV)
and whether we still have any anti-correlation between $\Omega_\chi h^2$ and $\sigma_{\chi}^{ann}$.

2. The standard BB scenario

The LSP relic density is given by

$$\Omega_\chi h^2 = \frac{\rho_\chi}{\rho_{\text{crit}}},$$

where $\rho_\chi$ is the LSP mass density and $\rho_{\text{crit}} = 3H_0^2/8\pi G \simeq 1.9 \times 10^{-6} h^2$ $g/cm^3$. The evolution of the number density $n_\chi = \rho_\chi/m_\chi$ can be described to a precision sufficient for our purposes by the Boltzmann equation

$$\frac{dn_\chi}{dt} + 3Hn_\chi = -\langle \sigma_{\chi}^{ann}v \rangle \left[ (n_\chi)^2 - (n^c_\chi)^2 \right],$$

where $H$ is the Hubble expansion rate. In the standard calculation, it is assumed that in the early universe ($T \gg m_\chi$) the LSP density was initially equal to its equilibrium value $n^c_\chi = n^eq$.

As the universe expands the LSP density traces its equilibrium value. When the particle becomes non relativistic ($T < m_\chi$) the number density of $\chi$'s drops exponentially. The annihilation rate $\Gamma_\chi \equiv \langle \sigma_{\chi}^{ann}v \rangle n_\chi$ becomes smaller than the expansion rate, $\Gamma_\chi \lesssim H$. At this point, the LSP decouples and fall out from being in equilibrium with the radiation, their number remains constant after that and its density is only diluted by the universe expansion.

As long as the LSP is non-relativistic, we can expand the averaged cross section $\langle \sigma_{\chi}^{ann}v \rangle$ as follows

$$\langle \sigma_{\chi}^{ann}v \rangle = \alpha_s + \alpha_p \langle v^2 \rangle,$$

where $\alpha_s$ is the s-wave contribution at zero relative velocity and $\alpha_p$ contains contribution from both s- and p- waves. The relic density solution to the Boltzmann equation is accurately given by the approximate formula (c.f. Eq.1)

$$\Omega_\chi h^2 \simeq \frac{8.8 \times 10^{-11} \text{GeV}^{-2}}{\sqrt{g_*(\alpha_s/x_F + 3\alpha_p/x_F^2)}},$$

where the freezing-out temperature $x_F \equiv m_\chi/T_F$ can be estimated as

$$x_F \sim \ln \frac{c(\alpha_s + 6\alpha_p/x_F)}{\sqrt{g_*(x_F)}},$$

In these formulas $g_*$ is the effective number of relativistic degrees of freedom at $T_F$ and $c$ is constant of order $10^{-1}$.

3. Alternative scenarios

In alternative cosmological scenarios, it is often assumed the existence of one (or more) epoch before the radiation-dominated era of the universe where the energy is dominated by coherent oscillating fields, the so called reheating era. A reheating process is associated by example with the final stage of Inflation. It can be also, independently or not, associated with the oscillations of moduli or Polony fields appearing in string theory after flat directions acquire a mass from SUSY breaking.

The details and physical origin of these reheating are largely unimportant. Its effects can be characterized by a reheating temperature $T_{RH}$ which is conventionally defined (see the definition below) as the temperature at which the oscillating field energy ceases to dominate the cosmological evolution and starts the radiation dominated epoch.

The standard BB scenario presented in the previous section can be seen as a special case when $T_{RH}$ is too much higher than $T_F$ (of order $10^9$ to $10^{10}$ GeV). In this case, the reheating epoch has no relevance in the final output of the relic density.

A problem could appear if the decay of the oscillating field occurs after or during the primordial BBN nucleosynthesis. This is drastically solved imposing from the onset that the reheating temperature is larger than $\sim 1$ MeV.

Finally, in interesting intermediate cases, a relatively low value of $T_{RH}$ such that $T_F > T_{RH}$ can have qualitative and quantitative implications on the predictions of the relic abundance of the LSP as discussed in Ref. [13]. The main reason for this is that now the decoupling of the LSP particle occurs in an expanding universe dominated by matter instead than radiation and that has effects. We will see below that theories with intermediate unification scale predict low reheating temperature ($\lesssim \text{GeV}$ and below) and they can be considered as interesting scenarios.

It is useful to use the following expression which can be interpreted as a quantitative definition of
the temperature $T_{RH}$ and which connects it to the decay width of the oscillating field:

$$
\Gamma_\phi = \left( \frac{4\pi^3 g_s(T_{RH})}{45} \right)^{-1/2} \frac{T_{RH}^2}{M_P}.
$$

(6)

Given any concrete theory, the reheating temperature can be related to the their physical parameters inverting Eq.6 for $T_{RH}$ as we can see on the following.

Let us consider first an inflation scenario. The inflaton ($\phi$) decay width which $\Gamma_\phi$ can be written, in some some supersymmetric hybrid inflation which has been considered [17]. as

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

where $m_\phi$ is the inflaton mass. Here $M_f$ is the mass of the particle $f$ that the inflaton decay to (i.e. a right handed neutrino or a sneutrino). $M_f$ should be less than the inflaton mass to allow for the decay $\phi \to ff$. Finally $\langle \phi \rangle$ is the vev of the inflaton field which is of the order of the unification scale.

In Ref. [18] it has been shown that a intermediate unification scale of order $M_I \sim 10^{11-12}$ GeV is favored by inflation. In this case, recalling that the inflaton mass is constrained by

$$
m_\phi < M_I^2/M_P,
$$

one obtains ($\langle \phi \rangle \sim 1.7 \times 10^{16}$ GeV [17]) an inflaton mass of the order $m_\phi \sim 10^2$ GeV. For the ‘standard GUT scenario’ with $M_I \sim 10^{16}$ one obtains however $m_\phi \sim 10^{13}$ GeV. From Eq.6 we find that $T_{RH} \sim O(1)$ GeV; a value much smaller than the values quoted in the standard GUT scenario ($T_{RH} \sim 10^{11}$). This reheating temperature is of the same order or lower than the typical freeze-out temperature $T_F \sim m_{\chi}/20 \sim O(1-5)$ GeV.

A detailed analysis of the relic density with a low reheating temperature has been considered in Ref. [13] with different possible generic scenarios. The first scenario corresponds to a case where the LSP is never in chemical equilibrium either before or after reheating. In the second the LSP’s reach chemical equilibrium but they freeze out before the completion of the reheat process $T_F > T_{RH}$.

Each of these scenarios leads to different qualitative and quantitative predictions for the relic density which can become quite different from the standard which we summarized in Eq.4.

CASE A) In case of non-equilibrium production and freeze out at the early times and supposing that the number density of the LSP $n_{\chi}$ is much smaller that $n_{\chi}^{eq}$, the Boltzmann equation [4] can be approximated and solved. In this case one gets [13]

$$
\Omega h^2 = 2.1 \times 10^4 \left( \frac{\alpha_s + \alpha_p}{4} \right),
$$

(7)

where

$$
\lambda_* = \frac{g^2}{4} \left( \frac{g_s(T_{RH})}{10} \right)^{3/2} \left( \frac{10}{g_s(T_{RH})} \right)^3 \left( \frac{10^3 T_{RH}}{m_\chi^2} \right)^2,
$$

(8)

and where $g$ is the number of degrees of freedom of the LSP and $T_*$ is the temperature at which most of the LSP production takes place, it is given by $T_* \sim 4m_\chi/15$. As we can see this $\Omega h^2$ is proportional to the annihilation cross section, instead of being inversely proportional, as in Eq.3. However, the assumption that $n_{\chi} << n_{\chi}^{eq}$ leads to a sever constraint on the annihilation cross section [13]; $\alpha_s < \alpha_p$ and $\alpha_p < \alpha_{\chi}$, where $\alpha_s \simeq 10^{-10} - 10^{-9}$ and $\alpha_p \simeq 10^{-9} - 10^{-8}$. So that for large neutralino cross sections, which we are interested in, Eq.4 can not be applied.

CASE B) With large annihilation cross section the LSP reaches equilibrium before reheating as discussed in Ref. [13] and its relic density $\Omega h^2$ is very close to the one obtained by using the standard calculation of relic density in Eq.4. In this case, $T_F$ is obtained by solving a different equation as before [13]

$$
x_F = \ln \left[ C \left( \alpha_s x_F^{5/2} + \frac{5}{4} \alpha_p x_F^{3/2} \right) \right],
$$

(9)

where $C$ is an unimportant constant now. The relic density $\Omega h^2$ is given by

$$
\Omega h^2 = \frac{g^2 g_s(T_{RH})}{g_s(T_F) m_\chi^2 (\alpha_s x_F^{5/2} + 4\alpha_p x_F^{3/2})^2}.
$$

(10)

In this case the relic density is again inversely proportional to the annihilation cross section as in Eq.4. Moreover, it has a further suppression
due to the very low reheating temperature effect. The typical value for the relic density in this case for $\sigma_{NN} \sim 10^{-6}$ GeV$^{-2}$ is $\Omega h^2 \lesssim 10^{-3}$. So it is even worse result than the BB standard computation’s one.

4. Our Model

In superstring theory, which is the first motivation for the idea of considering an intermediate scale $\mathcal{M}$, there are many moduli fields which acquire masses from SUSY breaking effects. Their masses are expected to be of order the gravitino mass but their couplings with the MSSM matter content are suppressed by a high energy scale $\mathcal{M}$. The moduli ($\phi$) decay width will now be parameterized now as

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

where $M_I$ is the unification scale ($\sim 10^{12}$ GeV) which acts as an effective suppression scale.

The reheating temperature is obtained using eqs. (8) and (11)

$$\frac{T_{RH}}{5 \text{ MeV}} = \frac{10^{14} \text{ GeV}}{M_I} \left( \frac{m_\phi}{100 \text{ GeV}} \right)^{3/2} \left( \frac{10.75}{g_*(T_{RH})} \right)^{1/4}$$

where, for example, $g_* = 10.75$ 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 $m_\chi = 50 - 200$ GeV, limits in practice the reheating temperature to be above $\sim 3$ 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 $T_F \approx m_\chi / 20 \sim 3 - 10$ 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 and moderate $\tan \beta$, $M_I = 10^{14}$ 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. 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.

![Figure 1](image)

Figure 1. The reheating temperature $T_{RH}$ as a function of the modulus mass $m_\phi$ for different values of the effective scale $M_I = 10^{11}, 10^{14}, 10^{16}$ GeV. The horizontal at $O(1)$ MeV is the nucleosynthesis limit.

When considering the direct decay of the moduli fields into LSP (and therefore, an additional quantity of non thermal LSP production) the evolution of the cosmological LSP abundance becomes more complicated. We have to consider now the coupled Boltzmann equations for the LSP, the moduli field and the radiation. The detailed equations used by us appear explicitly in Ref. (11). These equations depend on $m_\chi$, the LSP mass, $N_{LSP}$, the averaged number of LSP particles produced in the decay of the modulus field.

Let us discuss qualitatively the solution following the arguments used in ref. (11). For a $T_{RH}$ higher than $T_F$ the relic density will roughly reproduce the usual result given by eq. (11). How-
ever, 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 \( (\sigma v)_{\chi_1^n} \) drops below the expansion rate of the universe, \( H \), the neutralino freezes out. Then the relic density can be estimated as

\[
\Omega_{\chi_1^0} h^2 = \frac{3 m_{\chi_1^0} \Gamma_{\phi}^2}{2 (2\pi^2/45) g_* T_{RH}^4 (\sigma v)_{\chi_1^0} \rho_c / s_0} \quad (13)
\]

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

\[
\Omega_{\chi_1^0} h^2 = \frac{3 \tilde{N}_{\chi_i^0} m_{\chi_i} \Gamma_{\phi}^2 M_{\phi}^2}{(2\pi^2/45) g_* T_{RH}^3 m_{\phi} \rho_c / s_0} \quad (14)
\]

The actual relic density is estimated as the minimum of (13) and (14).

Now we can apply the above equations to our case with intermediate scales. Using eqs. (11) and (12), we can write expressions (13) and (14) as

\[
\Omega_{\chi_1^0} h^2 \propto \frac{M_I m_{\chi_i^0}^{3/2} \Gamma_{\phi}^{1/4}}{g_*} (\sigma v)_{\chi_1^0}^{1/8} \quad (15)
\]

\[
\Omega_{\chi_1^0} h^2 \propto \frac{M_{\chi_i^0}^{1/2} m_{\phi}}{g_*} \frac{1}{M_I} \quad (16)
\]

From these equations we can see that even with a large annihilation cross section, \( (\sigma v)_{\chi_1^0} \) \( \sim \) \( 10^{-8} \) GeV\(^{-2} \), we are able to obtain the cosmologically interesting value \( \Omega_{\chi_1^0} h^2 \) \( \sim \) \( 1 \). For example for \( M_I \) \( \sim \) \( 10^{13} \) GeV we obtain it when \( \tilde{N}_{\chi_i^0} \) \( \sim \) \( 1 \), using eq. (15).

Previous expressions are useful but approximate. We have solved numerically the Boltzmann equations in several cases. In Fig. 2 we show in detail the results of the numerical computation. We have used a particular expression for the annihilation cross section introduced which is explained below. In this figure, the contours of constant relic neutralino density \( \Omega_{\chi_1^0} h^2 \) as a function of \( m_{\phi} \) and \( \tilde{N}_{\chi_i^0} \) are shown, for fixed values of \( M_I \) and \( m_{\chi_i^0} \). In particular, we consider the cases \( M_I = 10^{12}, 10^{13}, 10^{14} \) GeV, with \( m_{\chi_i^0} = 100 \) GeV. The corresponding reheating temperatures can be obtained from Fig. 1. Note that whereas many values of \( \tilde{N}_{\chi_i^0} \) 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 finish this section with a few words about the neutralino annihilation cross section. 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

\[
(\sigma v)_{\chi_1^0} \simeq 8\pi a^2 \frac{1}{m_{\chi_i^0}^2} \frac{1}{1 + x_{\chi_i}^0} \langle u^2 \rangle \quad (17)
\]
where \( m_{1/2} = m_{\tilde{\chi}_1}^2 / m_{\tilde{\chi}_0} \) and \( \alpha' \) is the coupling constant for the \( U(1)_Y \) interaction. Taking \( m_{1/2} \sim m_{\tilde{\chi}_0} \sim 100 \) GeV, \( \langle \sigma_{\text{ann}} v \rangle \) in eq. (7) becomes of the order of \( 10^{-9} \) GeV\(^{-2}\) or smaller. Using eq. (3) an interesting relic abundance, \( \Omega_{\chi_1} h^2 \gtrsim 0.1 \), is obtained.

However, in Refs. [2,4], we obtained that for \( M_1 \sim 10^{12} - 10^{14} \) GeV the neutralino LSP acquires a sizeable Higgsino component, as we commented before this is essential for enhancing the LSP-nucleon cross section. For a Higgsino-like LSP the dominant annihilation channel is the decay into a W boson pair. An upper bound on the cross section can be obtained [13]:

\[
\langle \sigma_{\text{ann}} v \rangle \leq \frac{\pi \alpha_2^2}{2 m_{\tilde{\chi}}^2} \frac{1}{(1 - x_W)^{3/2}} \frac{1}{(2 - x_W)^2},
\]

where \( \alpha_2 \) is the \( SU(2) \) constant and \( x_W = m_{\tilde{\chi}}^2 / m_{\tilde{\chi}}^2 \). This is last expression for the cross section the one used in the numerical example presented in Fig. 1. For a LSP with \( m_{\tilde{\chi}} \sim 100 \) GeV, the formula gives a maximum \( \sim 2 \times 10^{-8} \) GeV\(^{-2}\). In the case that the LSP have sizeable components of both gaugino and Higgsino one expects smaller cross sections.

5. Conclusions and Comments

In conclusion we have analyzed the LSP relic density in supersymmetric models with intermediate unification scale. We find that the reheating temperature is of order 1 GeV (\( M_1 \sim 10^{12} \)) and below down to the nucleosynthesis threshold. We have shown that in generic low \( T_{\text{RH}} \) scenarios without non-thermal direct production is expected that the LSP relic density is too low. As an alternative, we have presented a concrete scenario where the reheating temperature is associated to the decay of oscillating moduli fields appearing in string theories, we show that the LSP relic density considerably increases with respect to the standard radiation-dominated case by the effect of the direct non-thermal production by the modulus field. The LSP can become a good dark matter candidate (0.01 – 0.1 \( \gtrsim \Omega h^2 \lesssim 0.3 – 1 \)) for \( M_1 \sim 10^{12} - 10^{14} \) GeV and \( m_{\phi} \sim 1 - 10 \) TeV.

Let us finally remark that the numerical value of \( \Omega_{\chi_1} \) is in general model dependent. This was discussed in the context of supergravity in ref. [13]. In this particular case both values \( \Omega_{\chi_1} \sim 1 \) and \( \Omega_{\chi_1} \sim 10^{-3} - 10^{-4} \) are plausible, depending on the characteristics of the supergravity theory under consideration.

REFERENCES

1. S. Khalil, C. Munoz and E. Torrente-Lujan, New J. Phys. 4 (2002) 27 [arXiv:hep-ph/0202135].
2. For a recent review see, S. Khalil and C. Munoz, hep-ph/0110124, to appear in Contemporary Physics, and references therein.
3. E. Gabrielli, S. Khalil, C. Munoz and E. Torrente-Lujan, Phys. Rev. D63 (2001) 025008.
4. D. Cerdeño, E. Gabrielli, S. Khalil, C. Munoz and E. Torrente-Lujan, Nucl. Phys. B603 (2001) 231.
5. A. Bottino, F. Donato, N. Fornengo and S. Scopel, Phys. Rev. D59 (1999) 095004.
6. E. Accomando, R. Arnovitt, B. Dutta and Y. Santoso, Nucl. Phys. B585 (2000) 124.
7. M.E. Gómez and J.D. Vergados, Phys. Lett. B512 (2001) 252.
8. For recent results in neutrino physics see, P. Aliani, et al., arXiv:hep-ph/0207348, P. Aliani, et al., arXiv:hep-ph/0206308, P. Aliani, et al., arXiv:hep-ph/0205057, P. Aliani, V. Antonelli, M. Picariello and E. Torrente-Lujan, Nucl. Phys. Proc. Suppl. 110, 361 (2002) arXiv:hep-ph/0112101, P. Aliani, et al., Nucl. Phys. B 634, 393 (2002) arXiv:hep-ph/0111418, S. Khalil and E. Torrente-Lujan, J. Egyptian Math. Soc. 9, 91 (2001) arXiv:hep-ph/0012203, E. Torrente-Lujan, Phys. Rev. D 53, 4030 (1996) arXiv:hep-ph/9505200, E. Torrente-Lujan, arXiv:hep-ph/9902330, E. Torrente-Lujan, Phys. Lett. B 389, 557 (1996) arXiv:hep-ph/9604218.
9. A. Corsetti and P. Nath, hep-ph/0003186.
10. D.G. Cerdeño, S. Khalil and C. Munoz, hep-ph/0105180, to appear in the Proceedings of CIHEP Conference (cairo), Rinton Press.
11. For a review see, G. Jungman, M. Kamionkowski and K. Griest, Phys. Rept. 267 (1996) 195, and references therein.
12. K.A. Olive and M. Srednicki, Phys. Lett. B230 (1989) 78; T. Moroi, M. Yamaguchi and T. Yanagida, Phys. Lett. B342 (1995) 105.
13. G. Giudice, E.W. Kolb and A. Riotto, Phys. Rev. D64 (2001) 023508.
14. E.W. Kolb and M.S. Turner, *The Early Universe*, (Addison–Wesley, Menlo Park, Ca., 1990).
15. T. Moroi and L. Randall, Nucl. Phys. B570 (2000) 455.
16. R. Jeannerot, X. Zhang and R. Brandenberger, JHEP 12 (1999) 003.
17. 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.
18. N. Kaloper and A. Linde, Phys. Rev. D59 (1999) 101303.
19. D.H. Lyth and E.D. Stewart, Phys. Rev. D53 (1996) 1784.
20. K. Benakli, Phys. Rev. D60 (1999) 104002;
   C. Burgess, L.E. Ibañez and F. Quevedo, Phys. Lett. B447 (1999) 257.
21. S.A. Abel, B.C. Allanach, F. Quevedo, L.E. Ibañez and M. Klein, JHEP 12 (2000) 026.
22. D.J. Chung, E.W. Kolb and A. Riotto, Phys. Rev. D60 (1999) 063504.