Heavy particle production during reheating

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We discuss production of heavy particles during reheating. We find that the very energetic inflaton decay products can contribute to the production of massive stable particles, either through collisions with the thermal plasma, or through collisions with each other. If such reactions exist, the same massive particles can also be produced directly in inflaton decay, once higher–order processes are included. We show that these new, non–thermal production mechanisms often significantly strengthen constraints on the parameters of models containing massive stable particles.

I. INTRODUCTION

After inflation [1], coherent oscillations of the inflaton dominate the energy density of the Universe. At some later time these coherent oscillations decay to the fields to which they are coupled, and their energy density is transferred to relativistic particles; this reheating stage results in a radiation–dominated Friedmann–Robertson–Walker (FRW) universe.

Until a few years ago, reheating was treated as the perturbative, one particle decay of the inflaton with decay rate Γ_d, leading to the simple estimate $T_\mathrm{R} \sim (\Gamma_\mathrm{d} M_\mathrm{Planck})^{1/2}$ for the reheat temperature [2], where $M_\mathrm{Planck} = 2.4 \times 10^{18}$ GeV represents the reduced Planck mass. It has been noticed in recent years that the initial stage of inflaton decay might occur through a complicated and non–perturbative process called parametric resonance [3]. However, it is generally believed that an epoch of (perturbative) reheating from the decay of massive particles (or coherent field oscillations, which amounts to the same thing) is an essential ingredient of any potentially realistic cosmological model [4]. In what follows we generically call the decaying particle the "inflaton". However, it should be clear that our results are equally well applicable to any other particle whose (late) decay results in entropy production.

Even before all inflatons decay, the decay products form a plasma which has the instantaneous temperature

$$T \sim \left(g_*^{-1/2} H \Gamma_\mathrm{d} M_\mathrm{Planck}^2 \right)^{1/4} , \tag{1}$$

where H is the Hubble parameter and g_* denotes the number of relativistic degrees of freedom in the plasma [5]. This temperature reaches its maximum T_max soon after the inflaton field starts to oscillate. In addition to this thermalized plasma there are inflaton decay products with energy $\simeq m_\phi/2$, which will eventually come into equilibrium with the thermal bath.

In recent years several mechanisms have been put forward for creating very heavy, even superheavy, particles in cosmologically interesting abundances [6]. Here we will focus on production of very massive particles from various processes, including a thermal bath, during perturbative reheating. Note that particle production from other sources, if present, would further strengthen the bounds which we will derive as they simply add to production from mechanisms discussed here.

II. HEAVY PARTICLE PRODUCTION

In Ref. [7] out of equilibrium production of χ from scatterings of "soft" particles in the thermal bath (with energy $E \sim T$) is studied and the final result is found to be (the superscript "ss" stands for χ production from "soft–soft" scattering)

$$\Omega_\chi^{\mathrm{ss}} h^2 \sim \left(\frac{200}{g_*}\right)^{3/2} \alpha_\chi^2 \left(\frac{2000 T_{\mathrm{R}}}{m_\chi}\right)^7$$
$$(\chi \text{ not in equilibrium}). \qquad (2)$$

In the opposite situation, with χ being initially in equilibrium with the thermal bath, today's χ relic density is given [8] by

$$\Omega_\chi^{\mathrm{ss}} h^2 \sim \left(\frac{200}{g_*}\right)^{1/2} \frac{T_{\mathrm{R}} x_{\mathrm{f}}^{4+a}}{m_\chi \alpha_\chi^2} \left(\frac{T_{\mathrm{R}}}{8 \cdot 10^5 \text{ GeV}}\right)^2$$
$$(\chi \text{ in equilibrium}), \qquad (3)$$

where the exponent $a = 0$ (1) if $\chi\chi$ annihilation proceeds from an S (P) wave initial state. The freeze–out temperature is now given by $x_{\mathrm{f}} \equiv m_\chi/T_{\mathrm{f}} \simeq \log(0.08 g_*^{-1/2} \alpha_\chi^2 x_{\mathrm{f}}^{2.5-a} M_{\mathrm{Planck}} T_{\mathrm{R}}^2/m_\chi^3)$.

However, as mentioned earlier, "hard" particles with energy $E \simeq m_\phi/2 \gg T$ are continuously created by inflaton decay for $H \geq \Gamma_{\mathrm{d}}$. These particles eventually thermalize with the bath, but this takes a finite amount of time. The presence of hard inflaton decay products can therefore affect heavy particle production in two ways. Firstly, χs can be produced from $2 \to 2$ scatterings of a hard particle off either soft particles in the thermal bath (if kinematically allowed), or off other hard particles [9]. Moreover, χs might be directly produced from inflaton decay [10].

A. Particle production from hard–soft scatterings

In order to estimate the rate of heavy particle production from the "hard" inflaton decay products, we also have to know the time needed to reduce their energy from a value $\sim m_\phi/2$ to a value near T. As shown in Ref. [11], $2 \to 2$ scattering reactions are *not* very efficient in this respect. The reaction rate is large, but the average energy loss per scattering is only $\mathcal{O}(T^2/m_\phi)$, giving a slow–down time of order $[\alpha^2 T^2/m_\phi]^{-1}$ (up to logarithmic factors). On the other hand, inelastic $2 \to 3$ reactions allow large energy losses (in nearly collinear particles) even if all virtual particles only have virtuality of

order T. The slow–down rate is thus given by:

$$\Gamma_{\text{slow}} \simeq 3\alpha^3 T \left(\frac{g_*}{200}\right)^{1/3}. \tag{4}$$

Next let us estimate the rate for χ pair production from hard–soft scatterings. This process is kinematically allowed so long as $ET \geq 4m_\chi^2$, where E is the energy of the hard particle so that the square of the center–of–mass energy is typically a few times ET. The hard particle initially has energy $E \simeq m_\phi/2$ and average number density $\bar{n}_h \sim g_* T^4/(3m_\phi)$, just after its production from inflaton decay. On the other hand, the rate for χ production from hard–soft scatterings is approximately given by

$$\Gamma_\chi^{\text{hs}} \sim \left(\frac{\alpha_\chi^2}{Tm_\phi} + \frac{\alpha\alpha_\chi^2}{m_\chi^2}\right) 0.2T^3. \tag{5}$$

The two contributions in (5) describe $2 \to 2$ reactions with squared center–of–mass energy $\sim m_\phi T$ and "radiative return" $2 \to 3$ reactions, respectively; in the latter case the hard particle emits a collinear particle prior to the collision, thereby reducing the effective cms energy of the collision to a value near m_χ. In order to make a safe (under)estimate we choose the temperature $T_0 = 2T_{\text{thr}}$ for presenting our results; note that the χ pair production cross section at threshold, $s = 4m_\chi^2$, is suppressed kinematically.

Our final results for the contribution of hard–soft collisions to the χ relic density are [9]: for $T_0 < T_{\text{R}}$:

$$\Omega_\chi^{\text{hs}} h^2 \sim \left(\frac{200}{g_*}\right)^{1/3} \frac{\alpha_\chi^2}{\alpha^2} \left(\frac{T_{\text{R}}}{10^4 \text{ GeV}}\right)^2 \frac{10^{13} \text{ GeV}}{m_\phi}$$
$$\cdot \frac{100T_{\text{R}}}{m_\chi} \left(1 + \frac{m_\chi^2}{\alpha T_{\text{R}} m_\phi}\right), \quad (T_0 < T_{\text{R}}) \tag{6}$$

in the opposite situation, we have

$$\Omega_\chi^{\text{hs}} h^2 \sim \left(\frac{200}{g_*}\right)^{1/3} \frac{\alpha_\chi^2}{\alpha^3} \frac{m_\phi}{10^{13} \text{ GeV}} \left(\frac{3000T_R}{m_\chi}\right)^5,$$
$$(T_0 > T_{\text{R}}) \tag{7}$$

B. Particle production from hard–hard scatterings

If $T_0 > T_{\text{R}}$, we should also consider χ production from scattering of two hard particles. Collisions of these particles with each other can produce χ pairs if $m_\chi < m_\phi/2$. Note that this constraint is independent of the temperature. On the other hand, it's also possible that $T_0 > T_{\text{max}}$, in which case hard–soft scattering (and soft–soft scattering) does not produce any χ particles.

The rate of χ production from hard–hard scattering is quadratic in the density of hard particles. Therefore we can not use our earlier approximation of the density $\bar{n}_h$ of hard particles produced in one Hubble time in the presence of a thermalized plasma, since the actual density $n_h(t)$ at any given time will be much smaller than this. In a plasma with

temperature T a hard particle will only survive for a time $\sim 1/\Gamma_{\text{slow}}$, see eq.(4). For $T_{\text{max}} > T > T_{\text{R}}$, the production of hard particles from inflaton decays and their slow–down will be in equilibrium, i.e. the instantaneous density $n_h(t) = 2\Gamma_d n_\phi(t)/\Gamma_{\text{slow}}$, where n_ϕ is the density of inflatons. By taking into account χ production prior to the bulid up of a thermalized plasma, our final estimate will be [9]

$$\Omega_\chi^{\text{hh}} h^2 \sim 6 \cdot 10^{27} \cdot \left(\frac{g_*}{200}\right)^{1/2} \sigma_{\chi\chi} \frac{m_\chi T_{\text{R}}^7}{m_\phi^2 T_{\text{max}}^4}. \tag{8}$$

In Fig. (2) we present three numerical examples to compare the significance of hard-soft and hard-hard scatterings with that of soft-soft scatterings. A detailed discussion on this figure and features observed in it can be found in Ref. [9].

C. Particle production from inflaton decay

We now discuss the direct production of χ particles in inflaton decay whose importance has recently been noticed [10]. Let us denote the average number of χ particles produced in each ϕ decay by $B(\phi \to \chi)$. The χ density from ϕ decay can then be estimated as:

$$\Omega_\chi^{\text{decay}} h^2 \simeq 2 \cdot 10^8 B(\phi \to \chi) \frac{m_\chi}{m_\phi} \frac{T_{\text{R}}}{1 \text{ GeV}}. \tag{9}$$

Eq. (9) holds if the χ annihilation rate is smaller than the Hubble expansion rate at $T \simeq T_{\text{R}}$.

We now discuss estimates of $B(\phi \to \chi)$. This quantity is obviously model dependent, so we have to investigate several scenarios. The first, important special case is where χ is the LSP. If m_ϕ is large compared to typical visible–sector superparticle masses, ϕ will decay into particles and superparticles with approximately equal probability [10]. Moreover, all superparticles will quickly decay into the LSP and some standard particle(s). As a result, if χ is the LSP, then $B(\phi \to \chi) \simeq 1$, independently of the nature of the LSP.

Another possibility is that the inflaton couples to all particles with more or less equal strength, e.g. through non–renormalizable interactions. In that case one expects $B(\phi \to \chi) \sim 1/g_* \sim 1/200$. However, even if ϕ has no direct couplings to χ, the rate (9) can be large. The key observation is that χ can be produced in ϕ decays that occur in higher order in perturbation theory whenever χ can be produced from annihilation of particles in the thermal plasma. In most realistic cases, $\phi \to f\bar{f}\chi\bar{\chi}$ decays will be possible
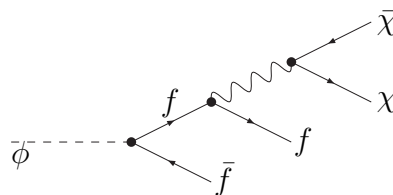

FIG. 1: Sample diagram for χ production in four-body inflaton decay.

if χ has gauge interactions, where f stands for some gauge non–singlet with tree–level coupling to ϕ. A diagram contributing to this decay is shown in Fig. (1). Note that the part of the diagram describing $\chi\bar{\chi}$ production is identical to the diagram describing $\chi\bar{\chi} \leftrightarrow f\bar{f}$ transitions. This leads to the following estimate:

$$B(\phi \to \chi)_4 \sim \frac{C_4 \alpha_\chi^2}{96\pi^3} \left(1 - \frac{4m_\chi^2}{m_\phi^2}\right)^2 \left(1 - \frac{2m_\chi}{m_\phi}\right)^{\frac{5}{2}}, \tag{10}$$

where C_4 is a multiplicity (color) factor. The phase space factors have been written in a fashion that reproduces the correct behavior for $m_\chi \to m_\phi/2$ as well as for $m_\chi \to 0$. This estimate provides a lower bound on $B(\phi \to \chi)$ under the conditions assumed for our calculation of Ω_χ^{hs} and Ω_χ^{hh}; whenever a primary inflaton decay product can interact with a particle in the thermal plasma, or with another primary decay product, to produce a $\chi\bar{\chi}$ pair, $\phi \to \chi$ four–body decays must exist.

Occasionally one has to go to even higher order in perturbation theory to produce χ particles from ϕ decays. For example, if χ has only strong interactions but ϕ only couples to $SU(3)$ singlets, $\chi\bar{\chi}$ pairs can only be produced in six body final states, $\phi \to f\bar{f}q\bar{q}\chi\bar{\chi}$. A representative diagram can be obtained from the one shown in Fig. (1) by replacing the χ lines by quark lines, attaching an additional virtual gluon to one of the quarks which finally splits into $\chi\bar{\chi}$. The branching ratio for such six body decays can be estimated as

$$B(\phi \to \chi)_6 \sim \frac{C_6 \alpha_\chi^2 \alpha^2}{1.1 \cdot 10^7} \left(1 - \frac{4m_\chi^2}{m_\phi^2}\right)^4 \left(1 - \frac{2m_\chi}{m_\phi}\right)^{\frac{9}{2}}. \tag{11}$$

Finally, in supergravity models with explicit (supersymmetric) χ mass term there in general exists a coupling between ϕ and either χ itself or, for fermionic χ, to its scalar superpartner, resulting in the estimate [12]

$$B(\phi \to \chi) \sim \frac{v^2 m_\chi^2 m_\phi}{16\pi \sqrt{g_*} M_{\text{Planck}}^3 T_{\text{R}}^2} \left(1 - \frac{4m_\chi^2}{m_\phi^2}\right)^{\frac{1}{2}}, \tag{12}$$

where v denotes the vacuum expectation value of the inflaton at the true minimum of its potential.

III. DISCUSSION

The production of χ particles from inflaton decay will be important for large m_χ and large ratio m_χ/T_{R}, but tends to become less relevant for large ratio m_ϕ/m_χ. Even if $m_\chi < T_{\text{max}}$, χ production from the thermal plasma (2) will be subdominant if

$$\frac{B(\phi \to \chi)}{\alpha_\chi^2} > \left(\frac{100 T_{\text{R}}}{m_\chi}\right)^6 \frac{m_\phi}{m_\chi} \frac{1 \text{ TeV}}{m_\chi}. \tag{13}$$

Note that the first factor on the r.h.s. of (13) must be $\lesssim 10^{-6}$ in order to avoid over–production of χ from thermal sources alone.

In [10] we showed that the decay contribution (9) by itself leads to very stringent constraints on models with massive stable χ particles. In particular, charged stable particles with mass below ~ 100 TeV seem to be excluded, unless $m_\chi > m_\phi/2$. In case of a (neutral) LSP with mass around 200 GeV, the overclosure constraint implies $m_\phi/T_R > 4 \cdot 10^{10}$, i.e. a very low reheat temperature, unless χ was in thermal equilibrium below T_R; recall that $B(\phi \to \chi) = 1$ in this case. Finally, if $m_\phi \sim 10^{13}$ GeV a "wimpzilla" with mass $m_\chi \sim 10^{12}$ GeV will be a good Dark Matter candidate only if it has a very low branching ratio, $B(\phi \to \chi) \sim 5 \cdot 10^{-8}$ GeV$/T_R$, i.e. if its couplings to ordinary matter are very small.

Many of the results presented here are only semi–quantitative. Unfortunately in most cases significant improvements can only be made at great effort. For example, a proper treatment of the slow–down of primary inflaton decay products would require a careful treatment of the full momentum dependence of the particle distribution functions. On the other hand, our estimates of $B(\phi \to \chi)$ should be quite reliable if $m_\chi > T_R$ (which is required for χ not to have been in thermal equilibrium at T_R); even for many–body decays, details of the matrix elements should change our estimates only be $\mathcal{O}(1)$ factors. Fortunately this is often also the most important of the new mechanisms for the production of massive particles at the end of inflation.

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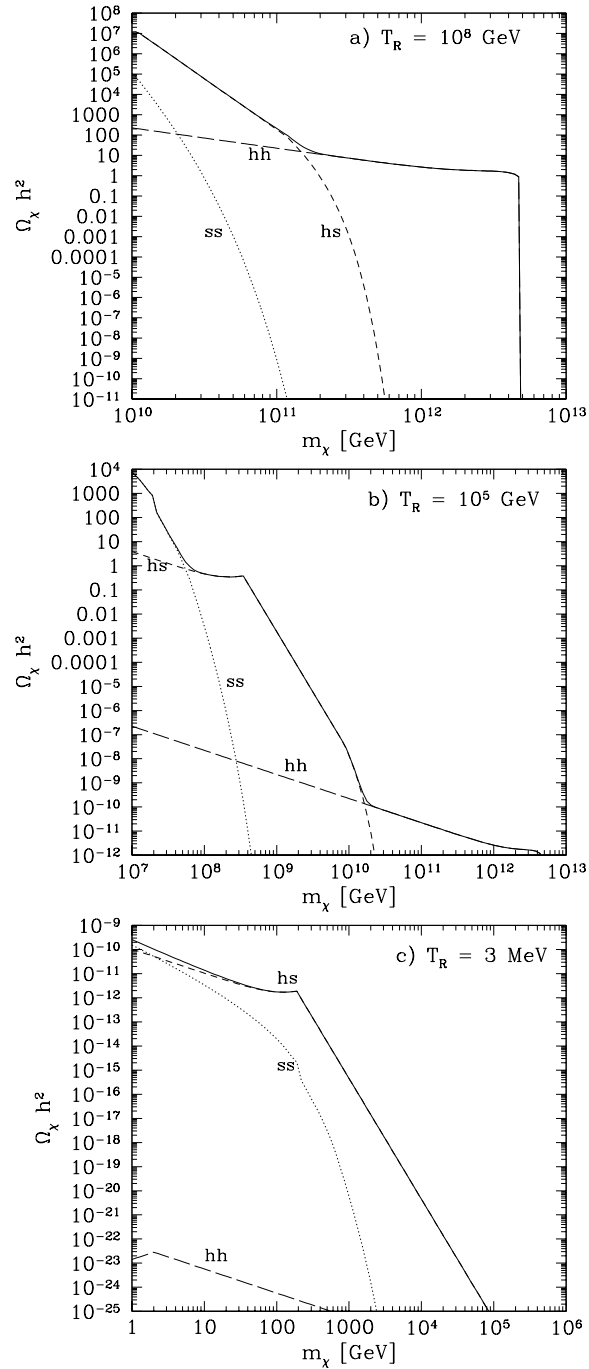[12] R. Allahverdi, K. Enqvist, and A. Mazumdar, Phys. Rev. **D 65**, 103519 (2002).

FIG. 2: The relic density χ particles would currently have if they are absolutely stable is shown as a function of m_χ for couplings $\alpha = 0.05$, $\alpha_\chi = 0.01$, and a) $(T_{\rm R}, m_\phi) = (10^8 \text{ GeV}, 10^{13} \text{ GeV})$, b)$(10^5 \text{ GeV}, 10^{13} \text{ GeV})$, c)$(3 \text{ MeV}, 10^8 \text{ GeV})$. The soft–soft, hard–soft and hard–hard contributions are shown by the dotted, short dashed and long dashed curves, respectively, while the solid curves show the sum of all three contributions.