Production of massive stable particles in inflaton decay

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We point out that inflaton decays can be a copious source of stable or long-lived particles $\chi$ with mass exceeding the reheat temperature $T_R$ but less than half the inflaton mass. Once higher order processes are included, this statement is true for any $\chi$ particle with renormalizable (gauge or Yukawa) interactions. This contribution to the $\chi$ density often exceeds the contribution from thermal $\chi$ production, leading to significantly stronger constraints on model parameters than those resulting from thermal $\chi$ production alone, particularly in models containing stable charged particles.

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According to inflationary models [1], which were first considered to address the flatness, isotropy, and monopole problems of the hot Big Bang model, the Universe has evolved through several stages. During inflation, the energy density of the Universe is dominated by the potential energy of the inflaton and the Universe experiences a period of superluminal expansion. Immediately after inflation, coherent oscillations of the inflaton dominate the energy density of the Universe. These oscillations eventually decay, and their energy density is transferred to relativistic particles; this reheating stage results in a radiation–dominated Friedmann–Robertson–Walker (FRW) Universe, as in the hot Big Bang model.

Initially reheating was treated as the perturbative, one particle decay of the inflaton with decay rate $\Gamma_d$, resulting in $T_R \sim (\Gamma_d M_p)^{1/2}$ for the reheat temperature [2], where $M_p = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. $T_R$ should be low enough so that the original monopole problem is avoided. Moreover, in many supersymmetric models $T_R \leq 10^7 - 10^9$ GeV, in order to avoid gravitino overproduction which would destroy the success of nucleosynthesis [3]. Later it has been noticed that the initial stages of inflaton decay might involve non–perturbative resonance processes [4]. They typically lead to a highly non–thermal distribution of particles, including inflatons with large momentum [5]. However, after sufficient red–shifting the energy density of the Universe would again be dominated by non–relativistic, massive particles. It is therefore 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 [6]. In what follows we generically call the decaying particle the “inflaton”, since we are (almost) sure that inflatons indeed exist. Note also that in a large class of well–motivated models, where the inflaton resides in a “hidden sector” of a supergravity theory [7], its couplings are suppressed by inverse powers of $M_p$, and hence are so weak that inflaton decays are purely perturbative. However, it should be clear that our results hold equally well for any other (late) decaying particle.

Even before all inflatons decay, their decay products form a plasma which, upon a very quick thermalization, has the instantaneous temperature [2] $T \sim \left( g_\ast^{-1/2} H I_d M_p^2 \right)^{1/4}$, where $H$ is the Hubble parameter and $g_\ast$ denotes the number of relativistic degrees of freedom in the plasma. This temperature reaches its maximum $T_{\text{max}}$ soon after the inflaton field $\phi$ starts to oscillate, which happens for a Hubble parameter $H_I \leq m_\phi$, with $m_\phi$ being the frequency of inflaton oscillations about the global minimum of the potential. We will assume that all inflaton decays can be described by perturbation theory in a trivial vacuum, which implies $T_{\text{max}} < m_\phi/2$. The resulting upper bound on $\Gamma_d$ also implies that a vacuum expectation value of the inflaton field does not induce large masses to the particles to which it couples. However, $T_{\text{max}}$ can be much larger than $T_R$. As long as $T > T_R$ the energy density of the Universe is still dominated by the (non–relativistic) inflatons that haven’t decayed yet. The Universe remains in this phase as long as $H > H_I$. During that epoch particles $\chi$ with mass $T_{\text{max}} > m_\chi > T_R$ can be produced copiously from the thermal plasma [8]. Here we point out that $\chi$ particles can also be produced directly in inflaton decays. We will show that the $\chi$ abundance from inflaton decay often exceeds that from thermal production, even if the branching ratio for $\phi \rightarrow \chi$ decays is very small.

We begin our argument by pointing out that $T_{\text{max}}$ is frequently well below $m_\phi$. This is important, since thermal production is obviously only efficient if $m_\chi \lesssim T_{\text{max}}$, while inflaton decay can produce pairs of $\chi$ particles as long as $m_\chi < m_\phi/2$. For perturbative inflaton decay thermalization increases the number density and reduces the mean energy of the decay products. Complete thermalization (i.e. both chemical and kinetic) therefore requires $2 \rightarrow N$ reactions, which change the number of particles, to be in equilibrium. Since the rate for higher order processes is suppressed by powers of the relevant coupling constant $\alpha$, the most important reactions are those with $N = 3$. These reactions have recently been studied in Ref. [2] where the scattering of two matter.
fermions with energy \( \simeq m_\phi / 2 \) (from inflaton decay) to two fermions, plus one gauge boson with typical energy \( E \ll m_\phi \) is considered. The rate for these reactions can be large due to the \( t \)-channel pole of the scattering matrix element, regulated by a cut–off on the exchanged momentum, naturally taken to be the inverse of the average separation between two particles in the plasma \[13\]. It turns out that the largest possible \( T_{\text{max}} \) is given by \[13\]:

\[
T_{\text{max}} \sim T_R \left( \frac{\alpha^3 (g_*)^{1/3}}{3} \frac{M_P}{m_\phi^{1/3} T_R^{2/3}} \right)^{3/8}.
\] (1)

Even if \( m_\phi \) is near its upper bound of \( \sim 10^{13} \text{ GeV} \[4\], for a chaotic inflation model, and \( T_R \) is around \( 10^9 \text{ GeV} \) (saturating the gravitino bound) \( T_{\text{max}} \) will exceed \( T_R \) if the coupling \( \alpha^3 \gtrsim 10^{-8} \). This is easily accommodated for particles with gauge interactions. On the other hand, recall that \( T_{\text{max}} < m_\phi / 2 \). Together with eq.(1), taking \( \alpha \lesssim 0.1 \), this gives \( T_{\text{max}} \lesssim 10^{11} (10^5) \text{ GeV} \) for \( T_R = 10^9 (1) \text{ GeV} \). This implies in particular that there will be no “wimpzilla” production \[3\] from thermalized inflaton decay products, since in this case \( m_\chi > T_{\text{max}} \).

On the other hand, for \( m_\chi \lesssim 20 T_R \) the standard calculation \[3\] of the density of stable relics applies. Scenarios with \( T_{\text{max}} \gtrsim m_\chi \gtrsim 20 T_R \) have only been investigated relatively recently in refs. \[3\] \[13\], which studied \( \phi \) production from the thermal plasma with \( T > T_R \). If the \( \chi \) density was always well below the equilibrium density, one finds

\[
\Omega_\chi^{\text{therm}} h^2 \sim \left( \frac{200}{g_*} \right)^{3/2} \frac{\alpha_\chi^2}{2000 T_R} \left( \frac{\chi}{m_\chi} \right)^7.
\] (2)

Here \( \Omega_\chi \) is the \( \chi \) mass density in units of the critical density and \( h \) is the Hubble constant in units of \( 100 \text{ km/s/Mpc} \). We have taken the cross section for \( \chi \) pair production or annihilation to be \( \sigma \simeq \alpha_\chi^2 / m_\chi^2 \). Note that \( \Omega_\chi \) is only suppressed by \( (T_R / m_\chi)^2 \) rather than by \( \exp (-m_\chi / T_R) \). A stable particle with mass \( m_\chi \sim 2000 T_R \cdot \alpha_\chi^2 / m_\chi \) might thus act as the Dark Matter in the Universe (i.e. \( \Omega_\chi \simeq 0.3 \)). However, eq.(1) with \( \alpha = 0.05 \) implies that \( T_{\text{max}} \gtrsim 1000 T_R \) is only possible if \( T_R < 2 \cdot 10^{-12} M_P \). Eq.(2) is no longer applicable \[3\] if the coupling \( \alpha_\chi \) is so large that \( \chi \) reached chemical equilibrium; however, it can then still be used as an upper bound on \( \Omega_\chi^{\text{therm}} \).

We now discuss the direct production of \( \chi \) particles in inflaton decay. (Other mechanisms for nonthermal production of superheavy particles have been discussed in \[13\].) Most inflatons decay at \( T \simeq T_R \); moreover, the density of \( \chi \) particles produced in earlier inflaton decays will be greatly diluted. Since inflaton decay conserves energy, the density of \( \chi \) particles produced in earlier inflaton decays will be greatly diluted. Since inflaton decay conserves energy, the density of \( \chi \) particles produced in earlier inflaton decays will be greatly diluted. Since inflaton decay conserves energy, the density of \( \chi \) particles produced in earlier inflaton decays will be greatly diluted. Since inflaton decay conserves energy, the density of \( \chi \) particles produced in earlier inflaton decays will be greatly diluted. 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with tree–level coupling to \( \phi \). 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:

\[
B(\phi \to \chi)_4 \sim \frac{C_4 \alpha^2_{\chi}}{96\pi^3} \left( 1 - 4\frac{m^2_{\chi}}{m^2_{\phi}} \right)^2 \left( 1 - 2\frac{m_{\chi}}{m_{\phi}} \right)^\frac{3}{2},
\]

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 \( m_{\chi} \to 0 \). Occasionally one has to go to even higher order in perturbation theory to produce \( \chi \) particles from \( \phi \) decays. For example, if \( \chi \) has only strong interactions but \( \phi \) only couples to \( SU(3) \) singlets, \( \chi \bar{\chi} \) pairs can only be produced in six body final states, \( \phi \to f f q \bar{q} \chi \bar{\chi} \). A representative diagram can be obtained from the one shown in Fig. 1 by replacing the \( \chi \) 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

\[
B(\phi \to \chi)_6 \sim \frac{C_6 \alpha^2_{\chi} \alpha_W^2}{1.1 \cdot 10^7} \left( 1 - 4\frac{m^2_{\chi}}{m^2_{\phi}} \right)^4 \left( 1 - 2\frac{m_{\chi}}{m_{\phi}} \right)^2.
\]

Another example where \( \chi \bar{\chi} \) pairs can only be produced in \( \phi \) decays into six body final states occurs if the inflaton only couples to fields that are singlets under the standard model gauge group, e.g. right–handed (s)neutrinos \( \nu_R \) \[17\]. [Since \( \nu_R \) decays very quickly, \( B(\phi \to \nu_R) \sim 1 \) does not cause any problem.] Since \( \nu_R \) only has Yukawa interactions, the factor \( \alpha_W^2 \) in eq. \[8\] would have to be replaced by the combination of Yukawa couplings \( \lambda_{\nu_R}^2 \lambda_{\phi}^2/(16\pi^2) \).

If \( 2m_{\chi} < m_{\nu_R} \), \( \chi \bar{\chi} \) pairs can already be produced in four body final states from \( \nu_R \) decay. The effective \( \phi \to \chi \) branching ratio would then again be given by eq. \[6\], with \( m_{\nu_R} \) replaced by \( m_{\nu_R} \) in the kinematical factors.

Finally, in supergravity models there in general exists a coupling between \( \phi \) and either \( \chi \) itself or, for fermionic \( \chi \), to its scalar superpartner, of the form \( a (m_{\phi} m_{\chi}/M_P) \phi \chi \chi + \text{h.c.} \) in the scalar potential \[13\]. A reasonable estimate for the coupling strength is \[13\] \( a \sim \langle \phi^2 \rangle/M_P \), unless an \( R \)–symmetry suppresses \( a \).

that most inflatons decay into other channels, so that \( \Gamma_d \sim \sqrt{\pi} T^2_{\phi}/M_P \) remains valid, this gives

\[
B(\phi \to \chi) \sim \frac{a^2 m_{\phi}^2}{16\pi \sqrt{g_s} M_P T_{\phi}^4} \left( 1 - 4\frac{m_{\chi}^2}{m_{\phi}^2} \right)^\frac{1}{2}.
\]

The production of \( \chi \) particles from inflaton decay will be important for large \( m_{\chi} \) and large ratio \( m_{\chi}/T_{\phi} \), but tends to become less relevant for large ratio \( m_{\chi}/m_{\phi} \). Even if \( m_{\chi} < T_{\max} \), \( \chi \) production from the thermal plasma \[9\] will be subdominant if

\[
B(\phi \to \chi) > \left( \frac{100 T_{\phi}}{m_{\chi}} \right)^6 \frac{m_{\phi}^6}{m_{\chi}^4} 1 \text{ TeV}.
\]

The first factor on the r.h.s. of \[10\] must be \( \lesssim 10^{-6} \) in order to avoid over–production of \( \chi \) from thermal sources alone. Even if \( \phi \to \chi \) decays only occur in higher orders of perturbation theory, the l.h.s. of \[10\] will be of order \( 10^{-4} \) (\( 10^{-10} \)) for four (six) body final states, see eqs. \[6\], \[8\]; if \( \phi \to \chi \) decays at tree–level, the l.h.s. of \[10\] will usually be bigger than unity. We thus see that even for \( m_{\phi} \sim 10^{13} \text{ GeV} \), as in chaotic inflation models, and for \( m_{\chi} \sim 10^3 T_{\phi} \), \( \chi \) production from decay will dominate if \( m_{\chi} \gtrsim 10^7 (10^{10}) \) GeV for four (six) body final states. As a second example, consider LSP production in models with very low reheat temperature. The LSP mass should lie within a factor of five or so of 200 GeV. Recall that in this case \( B(\phi \to \chi) = 1 \). Taking \( \alpha_W \sim 0.01 \), we see that \( \chi \) production from decay will dominate over production from the thermal plasma if \( m_{\phi} < 6 \cdot 10^7 \text{ GeV} \) for \( T_{\phi} = 1 \) GeV; this statement will be true for all \( m_{\phi} \lesssim 10^{13} \text{ GeV} \) if \( T_{\phi} \lesssim 100 \text{ MeV} \).

Let us now assume that eq. \[4\] indeed gives the dominant contribution to \( \chi \) production in the early Universe, and investigate the resulting constraints on model parameters. As well known, any stable particle must satisfy \( \Omega_{\chi} h^2 < 1 \), since otherwise it would “overclose” the Universe. For example, in case of a neutral LSP with \( m_{\chi} \sim 200 \text{ GeV} \), eq. \[4\] with \( B(\phi \to \chi) = 1 \) implies \( m_{\phi}/T_{\phi} > 4 \cdot 10^{10} \). Such a large ratio \( m_{\phi}/T_{\phi} \) in turn requires \( \Gamma_d < 10^{-21} m_{\phi}^2/M_P \), which indicates that \( \phi \) would have to decay through higher dimensional operators. Of course, this constraint is no longer valid if \( \chi \) reaches equilibrium with the plasma at temperatures \( \lesssim T_{\phi} \).

Another Dark Matter candidate is a very massive particle, with \( m_{\chi} \sim 10^{12} \text{ GeV} \); decays of this particle could give rise to the observed very energetic cosmic rays \[13\] if their lifetime is \( \gtrsim 10^8 \) times the age of the Universe. We noted above that such massive particles cannot be produced thermally in any realistic model of inflation. On the other hand, eq. \[4\] shows that inflaton decays might very easily produce too many of such particles. Taking \( m_{\phi} = 10 \text{ GeV} = 10^{13} \text{ GeV} \), we see that we need a branching ratio as small as \( 5 \cdot 10^{-8} \text{ GeV}/T_{\phi} \), which implies quite a severe upper bound on \( T_{\phi} \) even if \( \chi \) pairs can only be produced in six body decays of the inflaton. Even taking

\[\text{Fig. 1: Sample diagram for } \chi \text{ production in four-body \phi\chi decay.}\]
$T_R = 1$ MeV, the lowest value compatible with successful nucleosynthesis, this requires $B(\phi \to \chi) < 10^{-4}$. Finally, if $\chi$ is produced only through $M_P$ suppressed interactions, eq. (4) implies $a^2 < 3.5 \cdot 10^{-6}$ GeV $\cdot M_P T_R / m_\chi^2$, which again gives a very tight constraint if $m_\chi \sim 10^{12}$ GeV.

In some cases other considerations give an even stronger constraint on $\Omega_\chi$. For example, the abundance of charged stable particles is severely constrained from searches for exotic isotopes in sea water [11]. e.g. $\Omega_\chi h^2 \leq 10^{-20}$ for $100$ GeV $\leq m_\chi \leq 10$ TeV; for heavier particles this bound becomes weaker. This bound imposes very severe constraints on supersymmetric models with stable charged LSP. Fixing again $m_\chi = 200$ GeV from considerations of naturalness, $m_\phi / T_R > 4 \cdot 10^{30} B(\phi \to \chi)$ is required. This is clearly incompatible with the limits $T_R \gtrsim 1$ MeV, $m_\phi \lesssim 10^{13}$ GeV, even if $\phi \to \chi$ decays require six body final states, see eq. (5). We saw above that arranging $\chi$ to have been in equilibrium at $T_R$ does not help. Finally, the relic density of charged LSPs that were in thermal equilibrium at $T < T_R$ is too large by more than ten orders of magnitude. Eq. (1) shows that the situation for larger $m_\chi$ would be even worse. We thus conclude that in models where at least a significant fraction of the present entropy of the Universe originates from inflaton decay, a stable charged LSP can only lead to an acceptable cosmology if it is too massive to be produced in inflaton decays.

Our calculation is also applicable to entropy–producing particle decays that might occur at very late times. If $\chi$ is lighter than this additional $\phi'$ particle [10], all our expressions go through with the obvious replacement $\phi \to \phi'$ everywhere. More generally our result holds if $\phi$ decays result in a radiation dominated era with $T_R > m_{\phi'}$. If $\phi'$ is sufficiently long–lived, the Universe will eventually enter a second matter–dominated epoch. $\phi'$ decays then give rise to a second epoch of reheating, leading to a radiation–dominated Universe with final reheating temperature $T_R$, and increasing the entropy by a factor $m_{\phi'} / T_R$. This could be incorporated into eq. (4) by replacing $T_R \to T_R T_R / m_{\phi'} > T_R$. Our result regarding a stable charged LSP would remain valid in such a scenario even if $m_\chi > m_{\phi'}$, since the lower bound of $\sim 1$ MeV which we used now applies to $T_R$. The only way out would be to allow $\phi'$ to be essentially the only decay product of $\phi$, where $\phi'$ itself does not have renormalizable interactions with standard particles and their superpartners (so that higher order $\phi$ decays are negligible) and $2m_\chi > m_{\phi'}$. However, there is presently no motivation for considering such baroque models.

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