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 χ 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 χ particle with renormalizable (gauge or Yukawa) interactions. This contribution to the χ density often exceeds the contribution from thermal χ production, leading to significantly stronger constraints on model parameters than those resulting from thermal χ 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 universe, as in the hot big bang model.

Initially, reheating was treated as the perturbative, one particle decay of the inflaton with decay rate Γ_d, resulting in $T_R \sim (\Gamma_d M_P)^{1/2}$ for the reheat temperature [1,2], 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 nonperturbative resonance processes [4]. They typically lead to a highly nonthermal distribution of particles, including inflatons with large momentum [5]. However, after sufficient redshifting the energy density of the Universe would again be dominated by nonrelativistic, 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] $T \sim (g_*^{-1/2} H \Gamma_d M_P^2)^{1/4}$, where H is the Hubble parameter and g_* denotes the number of relativistic degrees of freedom in the plasma. This temperature reaches its maximum T_{max} soon after the inflaton field ϕ starts to oscillate, which happens for a Hubble parameter $H_I \leq m_\phi$, with m_ϕ 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 $T_{max} < m_\phi/2$. (The resulting upper bound on Γ_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_{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 (nonrelativistic) inflatons that have not decayed yet. The Universe remains in this phase as long as $H > \Gamma_d$. During that epoch particles χ with mass $T_{max} > m_\chi > T_R$ can be produced copiously from the thermal plasma [8–11]. Here we point out that χ particles can also be produced directly in inflaton decays. We will show that the χ abundance from inflaton decay often exceeds that from thermal production, even if the branching ratio for $\phi \to \chi$ decays is very small.

We begin our argument by pointing out that T_{max} is frequently well below m_ϕ. This is important, since thermal production is obviously efficient only if $m_\chi \lesssim T_{max}$, while inflaton decay can produce pairs of χ 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 $2 \to N$ 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 α, the most important reactions are those with $N = 3$. These reactions have recently been studied in Ref. [12] 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

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rate for these reactions can be large due to the t-channel pole of the scattering matrix element, regulated by a cutoff on the exchanged momentum, naturally taken to be the inverse of the average separation between two particles in the plasma [12]. It turns out that the largest possible $T_{\max}$ is given by [13]

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

Even if m_ϕ is near its upper bound of $\sim 10^{13}$ GeV [14], for a chaotic inflation model, and T_R is around 10^9 GeV (saturating the gravitino bound), $T_{\max}$ 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 $T_{\max} < m_\phi/2$. Together with Eq. (1), taking $\alpha \lesssim 0.1$, this gives $T_{\max} \leq 10^{11}(10^5)$ GeV for $T_R = 10^9(1)$ GeV. This implies, in particular, that there will be no "wimpzilla" production [8] from *thermalized* inflaton decay products, since in this case $m_\chi > T_{\max}$.

On the other hand, for $m_\chi \lesssim 20 T_R$ the standard calculation [2] of the density of stable relics applies. Scenarios with $T_{\max} \gtrsim m_\chi \gtrsim 20 T_R$ have been investigated only relatively recently in Refs. [8,9,10,11], which studied χ production from the thermal plasma with $T > T_R$. If the χ density was always well below the equilibrium density, one finds

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

Here Ω_χ is the χ mass density in units of the critical density and h is the Hubble constant in units of 100 km/(s · Mpc). We have taken the cross section for χ pair production or annihilation to be $\sigma \simeq \alpha_\chi^2/m_\chi^2$. Note that Ω_χ is suppressed only by $(T_R/m_\chi)^7$ rather than by $\exp(-m_\chi/T_R)$. A stable particle with mass $m_\chi \sim 2000 T_R \alpha_\chi^{2/7}$ 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 $T_{\max} \gtrsim 1000 T_R$ is possible only if $T_R < 2 \times 10^{-12} M_P$. Equation (2) is no longer applicable [9] if the coupling α_χ is so large that χ 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 χ particles in inflaton decay. (Other mechanisms for nonthermal production of superheavy particles have been discussed in [15].) Most inflatons decay at $T \simeq T_R$; moreover, the density of χ particles produced in earlier inflaton decays will be greatly diluted. Since inflaton decay conserves energy, the density of inflatons can be estimated as $n_\phi \simeq 0.3 g_* T_R^4/m_\phi$. Let us denote the average number of χ particles which are produced in each ϕ decay by $B(\phi \to \chi)$. We translate the χ density at $T = T_R$ into the present χ relic density using the relation [2]

$$\Omega_\chi h^2 = 6.5 \times 10^{-7} \frac{200}{g_*} \frac{m_\chi n_\chi(T_R)}{T_R^3 T_{\text{now}}}. \tag{3}$$

The χ density from ϕ decay is therefore [10]

$$\Omega_\chi^{\text{decay}} h^2 \simeq 2 \times 10^8 B(\phi \to \chi) \frac{m_\chi}{m_\phi} \frac{T_R}{1\ \text{GeV}}. \tag{4}$$

Equation (4) holds if the χ annihilation rate is smaller than the Hubble expansion rate at $T \simeq T_R$, which requires

$$\frac{m_\phi}{M_P} > 5 B(\phi \to \chi) \alpha_\chi^2 \left(\frac{T_R}{m_\chi}\right)^2 \left(\frac{g_*}{200}\right)^{1/2}. \tag{5}$$

This condition will be satisfied in chaotic inflation models with $m_\phi \sim 10^{-5} M_P$, if m_χ is large enough to avoid overclosure from thermal χ production alone. It might be violated in models with light inflaton. In that case the true χ density at T_R can be estimated by equating the annihilation rate with the expansion rate:

$$\Omega_\chi^{\max} \simeq \frac{5 \times 10^7}{\alpha_\chi^2} \frac{m_\chi^3}{(1\ \text{GeV}) M_P T_R} \left(\frac{200}{g_*}\right)^{1/2}. \tag{6}$$

This maximal density violates the overclosure constraint $\Omega_\chi < 1$ badly for the kind of weakly interacting ($\alpha_\chi \lesssim 0.1$), massive ($m_\chi \gg T_R$ and $m_\chi \gtrsim 1$ TeV) particles we are interested in. [Equation (6) describes the maximal χ density if χ decouples at $T \sim T_R$. It is not applicable to weakly interacting massive particles decoupling at $T < T_R$.] For the remainder of this Letter we will therefore estimate the χ density from inflaton decay using Eq. (4).

Our remaining task is to estimate $B(\phi \to \chi)$. This quantity is obviously model dependent, so we have to investigate several scenarios. The first, important special case is where χ is the lightest supersymmetric particle (LSP). If m_ϕ is large compared to typical visible-sector superparticle masses, ϕ will decay into particles and superparticles with approximately equal probability. (This statement is true so long as the superpotential is quadratic or higher in the inflaton superfield [13].) Moreover, all superparticles will decay into one χ particle and some standard particle(s) at a time scale which is shorter than the superparticle annihilation time scale [16], as long as $m_\chi > T_R$, even if $\alpha_\chi \simeq 0.1$. 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 nonrenormalizable 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 (4) 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 if χ has electroweak gauge interactions, where f stands for some gauge nonsinglet with tree-level coupling to ϕ. A diagram contributing to this decay is shown in Fig. 1. Note that the part of the diagram

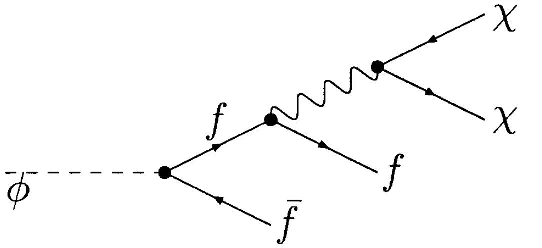

FIG. 1. Sample diagram for χ production in four-body inflaton decay.

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:

$$B(\phi \to \chi)_4 \sim \frac{C_4 \alpha_\chi^2}{96\pi^3}\left(1 - \frac{4m_\chi^2}{m_\phi^2}\right)^2\left(1 - \frac{2m_\chi}{m_\phi}\right)^{5/2}, \quad (7)$$

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 χ particles from ϕ decays. For example, if χ has only strong interactions but ϕ couples only to SU(3) singlets, $\chi\bar{\chi}$ pairs can be produced only 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

$$B(\phi \to \chi)_6 \sim \frac{C_6 \alpha_\chi^2 \alpha_W^2}{1.1 \times 10^7}\left(1 - \frac{4m_\chi^2}{m_\phi^2}\right)^4\left(1 - \frac{2m_\chi}{m_\phi}\right)^{9/2}. \quad (8)$$

Another example where $\chi\bar{\chi}$ pairs can be produced only in ϕ decays into six-body final states occurs if the inflaton couples only to fields that are singlets under the standard model gauge group, e.g., right-handed (s)neutrinos ν_R [17]. [Since ν_R decays very quickly, $B(\phi \to \nu_R) \sim 1$ does not cause any problem.] Since ν_R has only Yukawa interactions, the factor α_W^2 in Eq. (8) would have to be replaced by the combination of Yukawa couplings $\lambda_{\nu_R}^2\lambda_t^2/(16\pi^2)$. If $2m_\chi < m_{\nu_R}$, $\chi\bar{\chi}$ pairs can already be produced in four-body final states from ν_R decay. The effective $\phi \to \chi$ branching ratio would then again be given by Eq. (7), with m_ϕ replaced by m_{ν_R} in the kinematical factors.

Finally, in supergravity models there in general exists a coupling between ϕ and either χ itself or, for fermionic χ, its scalar superpartner, of the form $a(m_\phi m_\chi/M_P)\phi_{\chi\chi}$ + H.c. in the scalar potential [18]. A reasonable estimate for the coupling strength is [18] $a \sim \langle\phi\rangle/M_P$, unless an R symmetry suppresses a. Assuming that most inflatons decay into other channels, so that $\Gamma_d \sim \sqrt{g_*}T_R^2/M_P$ remains valid, this gives

$$B(\phi \to \chi) \sim \frac{a^2 m_\chi^2 m_\phi}{16\pi\sqrt{g_*}M_P T_R^2}\left(1 - \frac{4m_\chi^2}{m_\phi^2}\right)^{1/2}. \quad (9)$$

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 $m_\chi < T_{max}$, χ production from the thermal plasma (2) will be subdominant if

$$\frac{B(\phi \to \chi)}{\alpha_\chi^2} > \left(\frac{100T_R}{m_\chi}\right)^6 \frac{m_\phi}{m_\chi}\frac{1\ \text{TeV}}{m_\chi}. \quad (10)$$

The first factor on the right-hand side of (10) must be $\lesssim 10^{-6}$ in order to avoid overproduction of χ from thermal sources alone. Even if $\phi \to \chi$ decays occur only in higher orders of perturbation theory, the left-hand side (lhs) of (10) will be of order 10^{-4} (10^{-10}) for four-(six-)body final states; see Eqs. (7) and (8); if $\phi \to \chi\bar{\chi}$ decays at tree level, the lhs of (10) will usually be bigger than unity. We thus see that even for $m_\phi \sim 10^{13}$ GeV, as in chaotic inflation models, and for $m_\chi \simeq 10^3 T_R$, χ 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 5 or so of 200 GeV. Recall that in this case $B(\phi \to \chi) = 1$. Taking $\alpha_\chi \sim 0.01$, we see that χ production from decay will dominate over production from the thermal plasma if $m_\phi < 6 \times 10^7$ GeV for $T_R = 1$ GeV; this statement will be true for all $m_\phi \lesssim 10^{13}$ GeV if $T_R \lesssim 100$ MeV.

Let us now assume that Eq. (4) indeed gives the dominant contribution to χ 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 $m_\chi \simeq 200$ GeV, Eq. (4) with $B(\phi \to \chi) = 1$ implies $m_\phi/T_R > 4 \times 10^{10}$. Such a large ratio m_ϕ/T_R in turn requires $\Gamma_d < 10^{-21}m_\phi^2/M_P$, which indicates that ϕ would have to decay through higher dimensional operators. Of course, this constraint is no longer valid if χ reaches equilibrium with the plasma at temperatures $\lesssim T_R$.

Another dark matter candidate is a very massive particle, with $m_\chi \sim 10^{12}$ GeV; decays of this particle could give rise to the observed very energetic cosmic rays [19] 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 $m_\phi = 10m_\chi = 10^{13}$ GeV, we see that we need a branching ratio as small as $5 \times 10^{-8}\ \text{GeV}/T_R$, which implies quite a severe upper bound on T_R even if χ pairs can be produced only in six-body decays of the inflaton. Even taking $T_R = 1$ MeV, the lowest value compatible with successful nucleosynthesis, this requires $B(\phi \to \chi) < 10^{-4}$. Finally, if χ is produced only through M_P suppressed interactions, Eq. (9) implies $a^2 < (3.5 \times 10^{-6}\ \mathrm{GeV}) M_P T_R / m_\chi^3$, 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 Ω_χ. 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 $\lesssim m_\chi \lesssim$ 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, $m_\phi / T_R > 4 \times 10^{30} B(\phi \to \chi)$ 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. (8). We saw above that arranging χ 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 10 orders of magnitude. Equation (4) shows that the situation for larger m_χ 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 lead to an acceptable cosmology only 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 χ is lighter than this additional ϕ' particle [16], all our expressions go through with the obvious replacement $\phi \to \phi'$ everywhere. More generally our result holds if ϕ decays result in a radiation-dominated era with $T_R > m_{\phi'}$. If ϕ' is sufficiently long-lived, the Universe will eventually enter a second matter-dominated epoch. ϕ' decays then give rise to a second epoch of reheating, leading to a radiation-dominated Universe with final reheating temperature T_{R_f}, and increasing the entropy by a factor $m_{\phi'}/T_{R_f}$. This could be incorporated into Eq. (4) by replacing $T_R \to T_R T_{R_f} / m_{\phi'} > T_{R_f}$. 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_f}. The only way out would be to allow ϕ' to be essentially the only decay product of ϕ, where ϕ' itself does not have renormalizable interactions with standard particles or their superpartners (so that higher order ϕ decays are negligible) and $2m_\chi > m_{\phi'}$. However, there is presently no motivation for considering such baroque models.

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