Gravitino production in the warm inflationary scenario

A. N. Taylor
Institute for Astronomy, Royal Observatory, Blackford Hill, Edinburgh EH9 3HJ, United Kingdom

Andrew R. Liddle
Astronomy Centre, University of Sussex, Brighton, BN1 9QJ, United Kingdom

(Received 23 November 2000; published 19 June 2001)

We estimate the production of gravitinos during and after the end of a period of warm inflation, a model in which radiation is produced continuously as the field rolls down the potential producing dissipation. We find that gravitino production is efficient for models in the strong dissipation regime, with the result that standard nucleosynthesis is disrupted unless the magnitude of the inflaton potential is very small. Combining this with the constraint from the thermal production of adiabatic density perturbations, we find the dissipation rate must be extraordinarily strong, or that the potential is very flat.

DOI: 10.1103/PhysRevD.64.023513 PACS number(s): 98.80.Cq, 98.80.Es

I. INTRODUCTION

The *warm inflation* scenario [1] is an unusual variant on inflationary cosmology, in which the inflaton has significant interactions during the inflationary epoch leading to continuous production of radiation. The back-reaction of this production on the inflaton field appears as viscosity, slowing down the scalar field evolution and hence aiding slow-roll inflation [2]. In such a scenario, inflation can proceed with potentials steeper than those in standard chaotic scenarios.

Issues concerning the implementation of warm inflation within a realistic particle physics context have yet to be studied to the same depth as the standard inflationary scenario [3,4]; indeed a considerable number of obstructions to such an implementation have been described [5]. Nevertheless, given that the warm inflationary scenario is very different phenomenologically from the usual picture, it makes good sense to examine the extent to which its phenomenology is consistent with observations. The two main purposes of inflation are to provide a large, nearly homogeneous patch in the Universe, within which structure formation can take place, and to ensure that unwanted relic particles do not spoil the successes of standard hot big-bang cosmology. The first of these has seen a reasonable amount of study [6], and so we will consider an example of the latter.

In the context of modern particle physics, the most troublesome relics are the gravitino and the moduli fields [7]. We will consider the gravitino, whose existence arises as the supersymmetric partner of the graviton, and whose mass is expected to be an order of 1 TeV. It is a cosmological threat because if produced in enough abundance in the early Universe, it is sufficiently long lived to survive until after nucleosynthesis, at which point its decays to spoil the element abundances [8]. To avoid this, the ratio of gravitino to photon number densities must be below about 10^{-12}. The gravitino may be produced both by interactions within a thermal bath [11] and by various nonthermal processes [12]. In conventional inflationary scenarios, the former gives an important upper limit on the reheat temperature, while the latter may constrain many possible physical processes.

In this paper we explore the consequences of gravitino production during and after warm inflation. Warm inflation differs from conventional inflation in that radiation is constantly produced during inflation, and the radiation density decreases monotonically throughout the evolution, with inflation ending when the radiation density overtakes the inflaton energy density. There is therefore continuous gravitino production during inflation, and also no delay in post-inflationary thermal production due to an intervening (p)reheating period. Consequently, the gravitino bound is much harder to satisfy. We will show that the abundance of gravitinos produced during inflation is similar to that produced after inflation, and assess the strength of the constraints this imposes on warm inflation model building.

II. EVOLUTION OF FIELDS DURING WARM INFLATION

A. Dissipation effects during inflation

We review the dynamics of warm inflation closely following Taylor and Berera [6], where full details can be found. Warm inflation is distinguished from ordinary inflation by the presence of a viscous damping during the inflationary evolution, so that the inflaton field ϕ satisfies the equation

$$\ddot{\phi}+(3H+\Gamma)\dot{\phi}+V'=0. \tag{1}$$

Here $H\equiv\dot{a}/a$ is the Hubble parameter, a is the cosmological expansion factor, and Γ is the dissipation coefficient. $V(\phi)$ is the potential of the inflationary field. For simplicity, we assume a spatially flat universe throughout. A dot denotes differentiation with respect to time and a prime with respect to ϕ.

The energy density of relativistic species $\rho_{\rm rad}$ follows from energy conservation as

$$\dot{\rho}_{\rm rad}+4H\rho_{\rm rad}=\Gamma\dot{\phi}^2. \tag{2}$$

These equations are completed by the Friedmann equation

$$H^2=\frac{8\pi}{3m_{\rm Pl}^2}(\rho_\phi+\rho_{\rm rad}), \tag{3}$$

where $m_{\rm Pl}$ is the Planck mass and the energy density of the inflaton field is

$$\rho_\phi = V(\phi) + \tfrac{1}{2}\dot{\phi}^2. \tag{4}$$

During an inflationary era, the potential field dominates both the kinetic energy of the inflationary field and the energy density of the radiation, so the Friedmann equation can be reduced to

$$H^2 \simeq \frac{8\pi}{3m_{\rm Pl}^2} V. \tag{5}$$

Assuming the slow-roll condition $\ddot{\phi} \ll V'$, the equation of motion for the inflaton reduces to

$$\dot{\phi} \simeq -\frac{V'}{3H(1+r)}, \tag{6}$$

where

$$r \equiv \frac{\Gamma}{3H} \tag{7}$$

is a dimensionless dissipation coefficient, whose value is, in general, time dependent. During the inflationary period the production of radiation will soon settle into a stable state, where $\dot{\rho}_{\rm rad} \ll \Gamma \dot{\phi}^2$, giving

$$\rho_{\rm rad} \equiv \frac{\pi^2}{30} g_* T^4 = \frac{3}{4} r \dot{\phi}^2, \tag{8}$$

where g_* is the number of relativistic degrees of freedom. In the standard model of particle physics $g_* = 106.75$ for $T \geqslant 300$ GeV, while in the minimal supersymmetric model (MSSM) we shall assume here this rises to $g_* = 228.75$ once the temperature is above the mass of the supersymmetric particles.

Combining Eqs. (3), (6), and (8) we find

$$\rho_{\rm rad} = \frac{1}{2}\left[\left(1 + \frac{2\epsilon r}{(1+r)^2}\right)^{1/2} - 1\right]\rho_\phi, \tag{9}$$

where

$$\epsilon \equiv \frac{m_{\rm Pl}^2}{16\pi}\left(\frac{V'}{V}\right)^2, \tag{10}$$

is the usual inflationary slow-roll parameter. Equation (9) holds for all values of r. In the limit $r \to 0$ dissipation is switched off, and the radiation field vanishes as

$$\rho_{\rm rad} = \frac{\epsilon r}{2}\rho_\phi. \tag{11}$$

In this regime, increasing the dissipation factor r increases the decay of the inflaton field into radiation, while having no effect on the evolution of the inflaton field. For fixed ϵ the fractional density of radiation is at a maximum when $r = 1$ (equivalently $\Gamma = 3H$).

B. The strong dissipation regime

Our main focus will be the strong dissipation regime, where the differences from standard inflation are most pronounced. In the regime of strong dissipation $r \gg 1$, the radiation field is given by [6]

$$\rho_{\rm rad} = \frac{\epsilon}{2r}\rho_\phi. \tag{12}$$

The major effect of increasing the dissipation factor is to heavily dampen the evolution of the inflaton field, slowing its evolution down the potential and decreasing the decay into radiation.

The conditions for slow roll and warm inflation to occur are

$$\epsilon < 2r, \tag{13}$$

and for an extended period of inflation we need

$$|\eta| \ll 3r^2, \tag{14}$$

where

$$\eta \equiv \frac{m_{\rm Pl}^2}{8\pi}\frac{V''}{V}. \tag{15}$$

Equations (13) and (14) relax the usual constraints on the inflationary potential. Supercooled inflation ends when $\epsilon \approx 1$, while warm inflation takes place until

$$\epsilon \approx 2r \tag{16}$$

when the radiation energy density starts to dominate the energy density of the inflaton field. At this point the universe makes a smooth transition from the inflation phase to a radiation-dominated, hot Friedmann model.

To illustrate the evolution of warm inflation, we will consider polynomial potentials of the form

$$V(\phi) = \lambda m^4 \left(\frac{\phi}{m}\right)^\alpha, \tag{17}$$

where we allow α to be a positive real number. It is important to note that the only fully complete model of warm inflation, where a period of warm inflation comes to a natural end with the field finishing in the minimum of its potential, occurs when $\alpha = 2$. For other positive even integers, the potential has a suitable minimum at the origin. However, for $\alpha > 4$, the radiation density falls compared to the inflaton energy and so warm inflation does not take hold, as can be seen from Eq. (12) since $\epsilon \propto \phi^{-2}$ and $r \propto \phi^{-\alpha/2}$ [6]. The case $\alpha = 4$ is special in that the densities of the two components remain in fixed proportion; inflation proceeds forever with the dissipation easing the inflaton asymptotically into the minimum. Modification to the potential would be required to

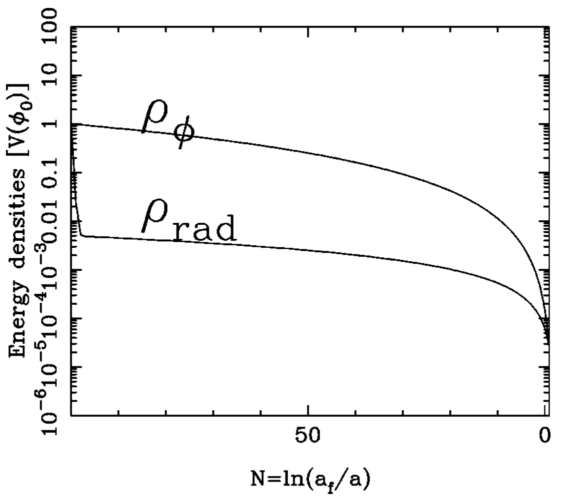

FIG. 1. The evolution of energy density of the inflaton and radiation fields as a function of number of e-folds from the end of warm inflation, for the inflationary potential $V=m^2\phi^2/2$ (see the text for details). The energy density of the inflationary field (upper curves) and the radiation field (lower curves) are taken as initially equal, and the vertical scale is arbitrary.

end inflation. We will also consider odd and nonintegral values of α (taking the field to have positive values). Such potentials have no minimum, and indeed may be ill defined for negative ϕ; we include them merely to illustrate the effects of modifications to the slope of the inflaton potential. Our range of investigation will cover $1\leqslant\alpha\leqslant4$.

For the polynomial potential, the slow-roll parameter $\epsilon=\alpha^2 m_{\rm Pl}^2/16\pi\phi^2$, and the number of e-folds of expansion to the end of inflation $N(\phi)$ is [6]

$$N(\phi)\equiv\ln\frac{a_{\rm end}}{a}\approx\frac{\alpha}{4-\alpha}\,\frac{\rho_\phi}{\rho_{\rm rad}},\tag{18}$$

where ρ_ϕ and $\rho_{\rm rad}$ are the energy densities near the start of warm inflation, once stable radiation production has been established. If warm inflation starts with an initial stable ratio $\rho_\phi/\rho_{\rm rad}$, it will take $N\sim\rho_\phi/\rho_{\rm rad}$ e-folds before the radiation and the vacuum energy are equal.

Figure 1 shows the evolution of the energy densities of the inflaton and radiation fields during an inflationary era, when $\alpha=2$. We numerically solved Eqs. (1), (2), and (3) as in Ref. [6]. The number of e-folds is $N=100$, and we began the model with $\rho_{\rm rad}=\rho_\phi$, although the evolution is insensitive to the initial conditions and soon settles into its stable configuration. The choice of parameters was made so that the amplitude of thermally produced adiabatic perturbations generated during warm inflation agrees with the amplitude of temperature fluctuations in the cosmic microwave background measured by the Cosmic Background Explorer, $\delta_{\rm H}=2\times10^{-5}$ (see Sec. IV B and Ref. [6] for details), with $\Gamma=10^2 m$, $m=10^{-8}m_{\rm Pl}$, and $V^{1/4}\sim10^{-4}m_{\rm Pl}$.

III. THERMAL PRODUCTION OF GRAVITINOS

Gravitinos are too weakly interacting to be able to reach thermal equilibrium with a radiation bath unless the temperature is around the Planck temperature. However, although interactions are negligible once they form, they can be created by two-body processes such as

$$\gamma+\gamma\rightarrow\tilde{g}+\bar{\tilde{g}}.\tag{19}$$

The single-particle decay rate for gravitinos gives a lifetime of order $m_{3/2}^3/m_{\rm Pl}^2$.

A. Thermal production during warm inflation

The number density $n_{3/2}$, of gravitinos produced from the thermal bath, is governed by the equation [11]

$$\dot{n}_{3/2}+3Hn_{3/2}=\langle\sigma_{3/2}|v|\rangle n_{\rm rad}^2,\tag{20}$$

where

$$n_{\rm rad}=\frac{\zeta(3)}{\pi^2}g_*T^3\approx0.28g_*^{1/4}\rho_{\rm rad}^{3/4}\tag{21}$$

is the number density of particles in the thermal bath, and $\zeta(3)=1.202$. Provided the typical particle energies well exceed the gravitino mass, which is always an excellent approximation for us, the effective total cross section in a thermal bath for gravitino production, including all particle channels, has been computed as $\sigma_{3/2}^{\rm tot}\approx250/m_{\rm Pl}^2$ [9–11,13]. For convenience we have defined a mean creation rate per particle species as $\sigma_{3/2}\approx250/g_*^2m_{\rm Pl}^2$, the square arising as the production is a two-body process.[1] We will be considering temperatures hot enough that all supersymmetric species participate in the thermal bath. The factor $v\approx1$ is the velocity of the produced gravitinos. The gravitino production is small enough that backreaction on the radiation density can be neglected, as can reactions destroying gravitinos by interactions or decays.

After a short period of inflation, the gravitino production rate becomes stable, $\dot{n}_{3/2}\ll\langle\sigma_{3/2}|v|\rangle n_{\rm rad}^2$, and the gravitino number density is given by

$$n_{3/2}=r_g n_{\rm rad},\tag{22}$$

where

$$r_g=\frac{\langle\sigma_{3/2}|v|\rangle}{3H}n_{\rm rad}\tag{23}$$

is the dimensionless production rate in units of the Hubble expansion. We define the yield of gravitinos as

$$Y_{3/2}\equiv\frac{n_{3/2}}{s},\tag{24}$$

where the entropy s is equal to $3.6n_{\rm rad}$ in the high-energy regime. Hence, during warm inflation, the yield is simply

[1]Our notation differs from that of Ref. [11], who define $n_{\rm rad}$ to be the number density of a single degree of freedom and use the total cross section.

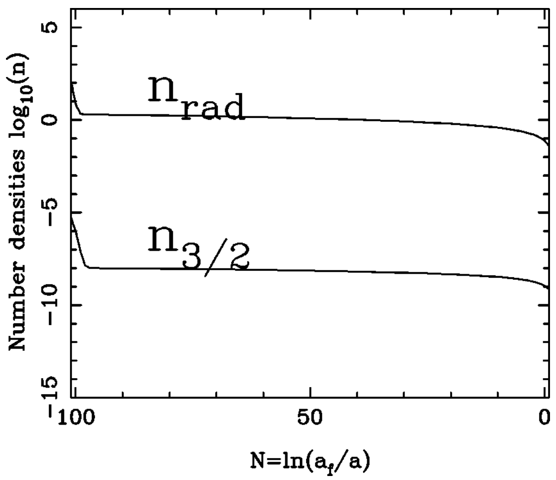

FIG. 2. The evolution of number density of the radiation and gravitinos, as a function of number of e-folds from the end of warm inflation, for the inflationary potential $V=m^2\phi^2/2$. The number densities are expressed in units of the radiation number density at the 50th e-fold.

$$Y_{3/2}=\frac{r_g}{3.6}. \tag{25}$$

As long as the stable production hypothesis is valid, the ultimate yield will only depend on the situation at the end of warm inflation; this means that we do not need to take into account any running of the gravitino interaction cross-section in obtaining our constraints.

Figure 2 shows the evolution of the number densities of the radiation and gravitino populations during warm inflation. The number densities are normalized to the number density of radiation at the 50th e-fold, and are typically $n_{\rm rad}\sim 10^{-12}m_{\rm Pl}^3$, and $n_{3/2}\sim 10^{-20}m_{\rm Pl}^3$. The choice of model parameters is the same as for Fig. 1.

The gravitino yield is given by Eqs. (21), (23), (25), (12), and (5):

$$Y_{3/2}=2\,g_*^{-7/4}\left(\frac{\epsilon}{2r}\right)^{3/4}\frac{V^{1/4}}{m_{\rm Pl}}. \tag{26}$$

For a fixed potential, the yield during warm inflation increases with increasing potential magnitude and slope, as there is greater dissipation into radiation, while the yield decreases with increasing dissipation factor, as strong dissipation will dampen the decay process.

For polynomial potentials of the form of Eq. (17), the yield is given by

$$Y_{3/2}=0.4\,g_*^{-7/4}\lambda^{5/8}\alpha^{3/2}\frac{m}{m_{\rm Pl}}\left(\frac{m_{\rm Pl}}{\Gamma}\right)^{3/4}\left(\frac{\phi}{m}\right)^{(5\alpha-12)/8}. \tag{27}$$

For this type of potential we see that $\alpha=12/5$ is a critical slope for gravitino production, leading to a constant yield as a function of time. Expressing Eq. (27) in terms of the number of e-folds until the end of warm inflation

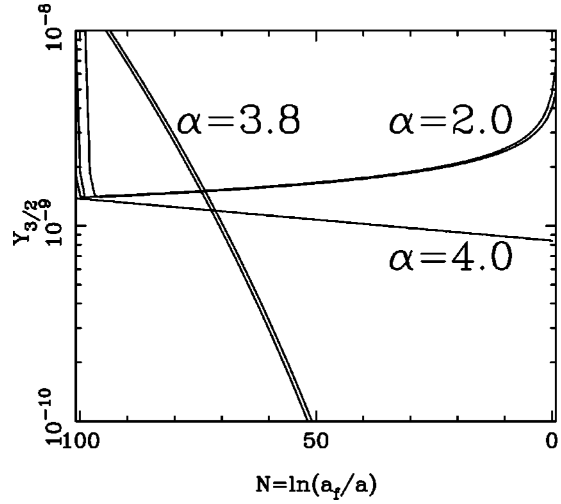

FIG. 3. The evolution of gravitino yield for $\alpha=2$, 3.8, and 4. The thick lines are calculated numerically, while the lighter lines are from the analytical expression [Eq. (27)]. The warm inflationary parameters are chosen to produce the observed amplitude of adiabatic density perturbations (see the text).

$$Y_{3/2}\sim\left[1+\frac{(4-\alpha)}{2\alpha}N\right]^{(5\alpha-12)/4(4-\alpha)}. \tag{28}$$

For $\alpha<12/5$, the yield increases as a function of time, or e-folds, with a maximum at the end of warm inflation. For $\alpha>12/5$, the yield is a decreasing function of time.

Figure 3 shows the evolution of the yield $Y_{3/2}$ during the warm inflation phase, for polynomial potentials with $\alpha=2$, 3.8, and 4. The solid lines are calculated numerically, while the lighter lines are from the analytic expression [Eq. (27)]. Since the gravitino number density quickly settles into its stable production state, the analytical expression accurately describes the evolution.

As the slope is increased to $\alpha\to 4$, the timescale for inflation is asymptotically stretched out, as discussed in Sec. II B. The dependence on e-folds asymptotically becomes $Y_{3/2}\sim e^{N/4}=(a/a_{\rm end})^{-1/4}$. This weak dependence can be seen in Fig. 3.

At the end of warm inflation, $\epsilon=2r$ and the yield is given by

$$Y_{3/2}^{\rm end}=2\,g_*^{-7/4}\frac{V^{1/4}}{m_{\rm Pl}}. \tag{29}$$

Hence at the end of warm inflation, the yield only depends on the magnitude of the potential. The end yield is directly related to the final temperature $T_{\rm end}$ by

$$Y_{3/2}(T_{\rm end})\simeq 1.5\,g_*^{-3/2}\left(\frac{T_{\rm end}}{m_{\rm Pl}}\right). \tag{30}$$

B. Evolution of gravitinos after warm inflation

Gravitinos have a sufficiently long decay time that they survive beyond nucleosynthesis. Their decay destroys ^{4}He and D nuclei by photodissociation, and if the gravitino abundance is high enough, this will disrupt the successful predictions of nucleosynthesis. The important quantity is the ratio of the gravitino to entropy densities.

The entropy is a particularly useful quantity to follow, as the comoving entropy is conserved not only during normal expansion but also during epochs where species fall out of thermal equilibrium and annihilate, changing the number of particle species in the thermal bath, so $s \propto 1/a^3$ always.[2] The evolution of gravitinos, Eq. (20), can be rewritten

$$\dot{n}_{3/2} + 3Hn_{3/2} = s\dot{Y}_{3/2} = \langle \sigma_{3/2}|v| \rangle n_{\rm rad}^2. \tag{31}$$

The yield produced, once radiation domination begins, is readily calculated using the normal radiation-dominated solution. It is dominated by early time production, as is well known, and the total yield from the beginning of radiation domination is

$$Y_{3/2} - Y_{3/2}^{\rm end} \simeq \frac{3g_*^{-11/6}}{m_{\rm Pl}}[n_{\rm rad}^{1/3}(T_{\rm end}) - n_{\rm rad}^{1/3}(T)]. \tag{32}$$

Once $T \ll T_{\rm end}$, the right-hand side of Eq. (32) equals the yield at the end of warm inflation, given by Eq. (29), since at the end of warm inflation, $\rho_\phi \simeq \rho_{\rm rad}$. We conclude therefore that the production of gravitinos during warm inflation can only lead to at most a factor of 2 enhancement over the total production at the end of the warm inflation era.

We note at this point that gravitinos might be produced by a variety of nonthermal mechanisms [12], which would add to the abundance. However, we will see that the thermal production already gives such strong constraints that there is no motivation to consider these effects.

IV. CONSTRAINTS ON WARM INFLATION

Having calculated the yield of gravitinos during the warm inflationary era, and shown that the yield does not significantly change afterwards, we now use this to constrain the warm inflationary parameters. In standard inflation, there are essentially two free parameters, the amplitude of the inflaton potential V and its slope ϵ. In general, these can be constrained by the gravitino yield and the amplitude of adiabatic perturbations, the latter being constrained by the observed fluctuations in the microwave background. In Sec. IV A we shall show that the main constraint from the gravitino production is on the magnitude of the inflaton potential, while the adiabatic density perturbations constrain the dissipation factor Γ. We begin with the constraint from the gravitino production.

A. The constraint from nucleosynthesis

Avoiding overproduction of D+^{3}He constrains the ratio of gravitinos to photons at the end of warm inflation. The details of the constraint depend on the gravitino mass [11], but for our purposes we can safely adopt a conservative limit

$$Y_{3/2} \lesssim 10^{-12}. \tag{33}$$

Combining this with Eq. (29) for the yield at the end of warm inflation, and Eq. (32), for the subsequent production during radiation domination, we find the following constraint on the magnitude of the inflation potential:

$$V^{1/4} \lesssim 2 \times 10^{-13} g_*^{7/4} m_{\rm Pl}. \tag{34}$$

This can be expressed in terms of the temperature at the end of the warm inflationary phase, via Eq. (30), giving

$$T_{\rm end} \lesssim 8 \times 10^6 g_*^{3/2}\ \ {\rm GeV}. \tag{35}$$

With $g_* = 228.75$ for the MSSM, this gives a constraint in good agreement with the standard result [11], though slightly weaker due to our adoption of a conservative constraint on $Y_{3/2}$. This is as expected since warm inflation has not greatly enhanced the gravitino yield. The main difference in warm inflation is how this constraint interacts with other constraints on the scenario.

B. The constraint from the amplitude of adiabatic density perturbations

In addition to the yield, we can add the independent constraint on warm inflation parameters from the amplitude of perturbations, δ_H^2, produced from thermal fluctuations during the warm inflationary era [6],

$$\delta_H^2 = 0.57 g_*^{-1/4}\left(\frac{r}{\epsilon}\right)^{3/4}\left(\frac{r^2 V}{m_{\rm Pl}^4}\right)^{3/4}. \tag{36}$$

The observational constraint from adiabatic density perturbations does not actually constrain the amplitude of the potential field alone [6], since the amplitude is observed at a fixed e-fold from the end of warm inflation, but rather with the combination $r^2 V \equiv \Gamma^2 m_{\rm Pl}^2/24\pi$. As we usually assume the amplitude of perturbations is measured at the 50th e-fold from the end of warm inflation, some scaling has to be made to match the constraint from the yield at the end of warm inflation. We may write the number of e-folds from the end of warm inflation as

$$N \simeq A\left(\frac{r}{\epsilon} - 1\right), \tag{37}$$

where A depends on the shape of the potential. For polynomial potentials

$$A = \frac{2\alpha}{(4-\alpha)}. \tag{38}$$

The amplitude of perturbations then constrains the dissipation factor

[2]An exception to this would be if there were particles with late out-of-equilibrium decays.

$$\Gamma = 10^{-6}\left(\frac{\delta_H^2}{4\times 10^{-10}}\right)^{2/3} g_*^{1/6}\left(1+\frac{N}{50A}\right)^{-1/2} m_{\rm Pl}. \quad (39)$$

As $r=\Gamma/H$ and $H\sim V^{1/2}$, we can combine Eq. (39) with the constraint on the gravitino yield, giving a constraint on the dimensionless dissipation rate

$$r \geqslant 2\times 10^{18}\left(\frac{\delta_{\rm H}^2}{4\times 10^{-10}}\right)^{2/3} g_*^{-7/3}\left(1+\frac{N}{50A}\right)^{-1/2}. \quad (40)$$

This equation is our main result. It shows that if warm inflation is to simultaneously produce density perturbations of a satisfactory magnitude and avoid overgenerating gravitinos, then the dimensionless dissipation must be extremely high. The main reason for this is that high dissipation increases the magnitude of density perturbations, allowing the inflationary energy scale to be normalized down. With sufficient dissipation, the energy scale becomes low enough that the gravitino yield is sufficiently suppressed.

One caveat to Eq. (40) is the factor A, relating the ratio r/ϵ to the number of e-folds before the end of warm inflation, when the density perturbations where formed. For polynomial potentials this is typically of order unity, but for potentials with very flat slopes, $\alpha\rightarrow 0$, the dissipation rate can be arbitrarily small, as the radiation production is suppressed.

V. CONCLUSIONS

A crucial role of inflation is to ensure that unwanted relic particles do not survive with abundances capable of spoiling the standard hot big-bang model. We have studied the production of gravitinos during and after a warm inflationary era, and combined the constraint this gives with the requirement that the density perturbations of the correct magnitude are generated.

We have found that although there is continuous gravitino production during warm inflation from interactions in the thermal bath, this does not in itself lead to a very significant extra yield of gravitinos over and above that produced at the end of warm inflation. Nevertheless, avoiding overproduction of gravitinos is much more challenging than in conventional inflationary scenarios, because the radiation density is monotonically decreasing throughout the evolution. Satisfying the gravitino bound requires that the potential energy at the end of warm inflation be very small, $V^{1/4}\lesssim 10^{-9} m_{\rm Pl}$. For density perturbations to have the correct magnitude, requires a dissipation factor $\Gamma\approx 10^{-6} m_{\rm Pl}$, and hence a dimensionless dissipation $r\gtrsim 10^{12}$, unless the slope of the potential is extremely flat. Either way, it is clear that evading overproduction of gravitinos strongly constrains the warm inflation scenario, requiring dimensionless numbers many orders of magnitude away from unity.

ACKNOWLEDGMENTS

A.N.T. thanks the PPARC Advanced Fellowship and the University of Sussex Astronomy Centre, where this work began, for its hospitality, and Arjun Berera for useful discussion. A.R.L. acknowledges CITA for hospitality while some of this work was carried out, and Beatriz de Carlos for discussions on gravitino properties.

[1] A. Berera and L. Z. Fang, Phys. Rev. Lett. **74**, 1912 (1995); A. Berera, *ibid.* **75**, 3218 (1995); Phys. Rev. D **54**, 2519 (1996); A. Berera, Nucl. Phys. **B585**, 666 (2000).
[2] J. Yokoyama and K. Maeda, Phys. Lett. B **207**, 31 (1988); A. R. Liddle, *ibid.* **220**, 502 (1989).
[3] D. H. Lyth and A. Riotto, Phys. Rep. **314**, 1 (1998).
[4] A. R. Liddle and D. H. Lyth, *Cosmological Inflation and Large-Scale Structure* (Cambridge University, Cambridge, 2000).
[5] J. Yokoyama and A. Linde, Phys. Rev. D **60**, 083509 (1999).
[6] A. N. Taylor and A. Berera, Phys. Rev. D **62**, 083517 (2000).
[7] S. Sarkar, Rep. Prog. Phys. **59**, 1493 (1996).
[8] S. Weinberg, Phys. Rev. Lett. **48**, 1303 (1982).
[9] J. Ellis, J. E. Kim, and D. Nanopoulos, Phys. Lett. **145B**, 181 (1984).
[10] J. Ellis, D. Nanopoulos, K. Olive, and S.-J. Rey, Astropart. Phys. **4**, 371 (1996).
[11] M. Kawasaki and T. Moroi, Prog. Theor. Phys. **93**, 879 (1995).
[12] A. L. Maroto and A. Mazumdar, Phys. Rev. Lett. **84**, 1655 (2000); M. Lemoine, Phys. Rev. D **60**, 103522 (1999); G. F. Giudice, I. Tkachev, and A. Riotto, J. High Energy Phys. **08**, 009 (1999); R. Kallosh, L. Kofman, A. Linde, and A. Van Proeyen, Phys. Rev. D **61**, 103503 (2000); G. F. Giudice, A. Roitto, and I. Tkachev, J. High Energy Phys. **11**, 036 (1999); D. H. Lyth, Phys. Lett. B **469**, 69 (1999); **476**, 356 (2000); R. Kallosh, L. Kofman, A. Linde, and A. Van Proeyen, Phys. Rev. D **61**, 103503 (2000).
[13] H. Fujisaki, K. Kumekawa, M. Yamaguchi, and M. Yoshimura, Phys. Rev. D **54**, 2494 (1996).