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Inflationary Fluctuations, Entropy Generation and Baryogenesis

in a Cold Universe

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Abstract

We study the implications of a generic inflationary model for scenarios of baryogenesis based on the decays of coherent oscillations of squark and slepton fields. We consider the effects of De Sitter fluctuations on the magnitudes of the coherent oscillations of squarks and sleptons. We see that the largest contribution to the entropy density is due to inflaton decays which together with the value of the oscillation amplitude determined by the De Sitter fluctuations leads to a baryon to entropy ratio $O(10^{-10})$. The isothermal density fluctuations produced by the coherent oscillations are found to be negligible compared with the adiabatic fluctuations produced during inflation.

A fundamental test of any unified theory is its ability to produce a Universal baryon asymmetry¹⁾. In conventional grand unified theories, baryogenesis is most simply achieved by the out-of-equilibrium decays²⁾ of superheavy gauge or Higgs bosons at temperatures $T \sim M_G \sim 10^{15}$ GeV. Cosmological inflation already complicates this simple picture if the temperature of the Universe after inflation is $T_R \ll M_G$ as it is in most inflationary models.³⁾ In superstring models⁴⁾, the problem is further compounded. The gauge and Higgs bosons with masses much less than M_p that are needed for the out-of-equilibrium decays are not present. In addition, although other fields are present that could in principle have baryon non-conserving interactions, proton stability requires that these couplings be absent⁴⁾. Hence the only interactions violating baryon number are mediated by particles with masses $O(M_p)$. In this case,⁵⁾ one is forced to consider the possibility of generating a baryon asymmetry at low temperatures ($T \sim 10^3 \pm 1$ GeV) which is the only significant dynamical scale below M_p .⁶⁾

In supersymmetric models, a mechanism for low temperature baryogenesis has been proposed by Affleck and Dine (AD)⁷⁾. In this scenario, squark and slepton fields obtain large vacuum expectation values along flat directions of the scalar potential. The decay of these expectation values at $T \leq 10^4$ GeV when supersymmetry breaking effects become important will produce a large baryon asymmetry if C and CP are violated. Indeed, if such flat directions are present then large vacuum expectation values are naturally obtained⁸⁾ during inflation as a result of De Sitter fluctuations. In refs. 5 and 9, it was shown that such flat directions are in fact generically present in superstring models, sources of C and CP violation are

also readily available and that subsequent phase transitions do not significantly lower the baryon to entropy ratio, so that a maximal value for $n_B/s \sim 0(1)$ was obtainable. In this paper, we look more closely at the effects of inflation on the AD mechanism. In particular, we examine the magnitudes of the squark and slepton expectation values and the generation of entropy at the end of inflation. Although our results are applicable to superstring models, they also apply more generally to any supersymmetric model*.

In order not to tie ourselves down to any particular model of inflation we consider a generic inflaton scalar potential $V(\psi)$ of the form

$$V(\psi) = \mu^4 P(\psi) \tag{1}$$

where ψ is the inflaton and $P(\psi)$ is some polynomial in ψ having the right properties so as to inflate. We expect that P lacks a quadratic term and that other terms have coefficients $0(1)$. If the leading term in $P(\psi)$ is ψ^3 or ψ^4 the adiabatic density fluctuations¹² produced during inflation $\delta\rho_\psi/\rho_\psi \approx 10^{-4}$ as needed for galaxy formation if $\mu^2 \approx 10^{-7} M_P^2$. Also the Hubble

 *Note that although a generic scenario for inflation in supersymmetric models is well known¹⁰⁾, its extension to superstring inspired models is not¹¹⁾.

parameter during inflation is $H^2 \approx V(0)/M_p^2 = \mu^4/M_p^2$ so that $H \approx \mu^2/M_p \approx 10^{-7} M_p$ and the mass of the inflaton at the global minimum of $V(\psi)$ is $m_\psi^2 = d^2V/d\psi^2|_{\langle\psi\rangle} \approx 0(1)M_p^2 = \mu^4/M_p^2$ so that $m_\psi \approx \mu^2/M_p \approx 10^{-7} M_p$.

If we let ϕ denote generically our squark and slepton fields, then we can consider directions (in ϕ space) such that $V(\phi)$ is flat or at most $V(\phi) = \tilde{m}^2 \phi^2$ due to supersymmetry breaking effects. (That these directions exist has been explicitly demonstrated in supersymmetric⁷⁾ and superstring models⁵⁾.) It has been argued that during inflation quantum gravitational effects¹³⁾ can produce a non-zero value $\langle\phi^2\rangle$ for all scalar fields. In the limit that $m_\phi^2 \ll H^2$ these fluctuations can be as large as¹³⁾

$$\langle\phi^2\rangle \approx H^4/m_\phi^2 \quad (2)$$

which in our case for $m_\phi \sim \tilde{m} \sim 10^2$ GeV is $\langle\phi^2\rangle \sim 10^4 M_p^2$. It has also been argued that fluctuations with wavelength much greater than the horizon cannot be distinguished from a homogeneous classical field¹⁴⁾ with amplitude $\langle\phi^2\rangle^{1/2} = \phi_0$. In previous analyses^{7,8,5,9)} of the AD mechanism, it has been assumed that $\phi_0 \approx M_p$ as a working hypothesis.

There are, however, two reasons why ϕ_0 might be much smaller than M_p .

1) In reality inflation lasts only a finite time and may not last long enough to give the maximum value of ϕ_0 although still long enough to solve the usual cosmological problems.

2) Higher order (non-renormalizable) effects may prevent $\langle\phi^2\rangle$ from reaching its maximum value.

On the first point, the time dependence of $\langle \phi^2(t) \rangle$ in a De Sitter background has been calculated¹³⁾ and for time-scales $t \ll H/m_\phi^2$, the growth of $\langle \phi^2 \rangle$ is linear $\langle \phi^2(t) \rangle \approx H^3 t$, where $t = 0$ corresponds to the time when the temperature corrected bare mass of ϕ becomes less than $2H^2$. It is obvious that this moment coincides approximately with the onset of the inflationary epoch. At later times $\langle \phi^2 \rangle$ approaches exponentially fast to its asymptotic value. The time scale for the duration of inflation is given by $\tau \approx H/m_\psi^2(\psi \approx H)$. If $V(\psi)$ contains a cubic (quartic) term, $m_\psi^2(\psi \approx H) \approx \mu^4 H/M_P^3 \approx \mu^6/M_P^4$ ($m_\psi^2 \sim \mu^4 \langle \psi^2 \rangle / M_P^4 \approx \mu^6/M_P^4$) and thus $\tau \approx M_P^3/\mu^4 \ll H/m_\phi^2$ so that we are in the linear regime of the growth of $\langle \phi^2 \rangle$. Thus we would find

$$\phi_0^2 = (H^3 \tau) = \mu^2 \sim 10^{-7} M_P^2 \quad (3)$$

A second limiting factor to ϕ_0 is the effect of non-renormalizable interactions involving ϕ . The flat directions we are considering have only been shown to be flat through $O(\phi^4)$ in the scalar potential. This does not preclude the presence of higher order terms

$$V \approx \phi^{2n+2}/M_P^{2n-2} \quad (4)$$

with $n \geq 2$. If such a term is present, we can as a first approximation take $m_\phi^2 \sim \tilde{m}^2 + \langle \phi^{2n} \rangle / M_P^{2n-2} \sim \tilde{m}^2 + \langle \phi^2 \rangle^n / M_P^{2n-2} \sim \langle \phi^2 \rangle^n / M_P^{2n-2}$. In this case we can use eq. (2) to solve for $\langle \phi^2 \rangle$

$$\langle \phi^2 \rangle \sim H^{4/(n+1)} M_P^{\frac{2n-2}{n+1}} \quad (5)$$

For $H \sim 10^{-7} M_P$, we have $\langle \phi^2 \rangle \sim 0(10^{-9})$ if $n = 2$; $\langle \phi^2 \rangle \sim 0(10^{-7})$ if $n = 3$ etc. for the asymptotic values of the ϕ^2 fluctuations. These asymptotic values are reached in a time $t \sim 3HM_P^{2n-2}/\langle \phi^2 \rangle^n \sim H^{(1-3n)/n+1} M_P^{(2n-2)/(n+1)}$ which must be compared with the duration of inflation.

There are two reasons why inflation cannot last arbitrarily long. The first reason is that potentials with no barriers can also lead to an Euclidean bounce action,¹⁵⁾ so that the metastable inflaton state can decay directly by tunneling instead of rolling. The tunneling time t_{tunn} is proportional to M^{-1} times a large exponential factor, whose value depends on the exact details of the model, where M is the length of the flat part of the potential. Hence in primordial inflation $t_{\text{tunn}} \gg 0(1)$ and tunneling is unlikely to be of interest. More important limits can be obtained by considering the quantum fluctuations of the inflaton field. As we have seen in this case $\tau \leq M_P H^{-2}$.

We thus find that the time scales of the full sfermion field fluctuations, t , and the duration of inflation τ compare as

$$t/\tau \sim H^{\frac{3-n}{n+1}} M_p^{\frac{n-3}{n+1}} \quad (6)$$

Hence for $n > 3$, (remember that eq. (5) is only valid for $n > 1$, i.e. the flat potential cases) the region of linear growth is valid during the whole of inflation, and ϕ_o^2 is given by eq. (3). If $n = 2, 3$ inflation can (although it need not) last long enough for $\langle \phi^2 \rangle$ to reach its asymptotic value. We will assume for the remainder of the discussion that $\phi_o^2 \approx 10^{-7} M_p^2$.

At the end of inflation, we have (at least) two (nearly) homogeneous scalar fields ψ and ϕ oscillating about their global minima. For the inflaton, we expect that $\langle \psi \rangle \sim O(M_p)$ so that the energy density stored in the ψ oscillations is

$$\rho_\psi = m_\psi^2 \psi^2 \quad (7)$$

with an initial value $\psi_o \sim O(M_p)$. For our sfermion fields ϕ we have

$$\rho_\phi = \tilde{m}^2 \phi^2 \quad (8)$$

with the initial amplitude given by ϕ_o . Scalar field oscillations of ψ (ϕ) will begin when the expansion rate of the Universe drops below the scalar mass m_ψ (m_ϕ). Because $m_\psi > m_\phi = \tilde{m}$, inflaton oscillations begin first, when

$H \approx m_\psi \psi / M_p \sim m_\psi$ or when $\psi \sim O(M_p)$. During ψ oscillations, the Universe expands as if it were matter dominated so that for $R > R_\psi$ (R is the Robertson-Walker scale factor and R_ψ is its value when $\psi \sim O(M_p)$)

$$\rho_\psi = m_\psi^2 M_p^2 (R_\psi / R)^3 \quad (9)$$

During this period the Hubble parameter is given by

$$H \approx m_\psi (R_\psi / R)^{3/2} \quad (10)$$

Sfermion oscillations begin when $H \sim \tilde{m}$ or at

$$R = R_\phi \approx (m_\psi / \tilde{m})^{2/3} R_\psi \quad (11)$$

and the energy density in ϕ oscillations at $R > R_\psi$ is given by

$$\rho_\phi \approx \tilde{m}^2 \phi_o^2 (R_\phi^2 / R)^3 \approx m_\psi^2 \phi_o^2 (R_\psi / R)^3 \quad (12)$$

The above expression assumed that ϕ oscillations begin before ψ decays. We can check this by calculating the scale factor $R_{d\psi}$ when ψ 's decay. The inflaton decay rate is given by $\Gamma_\psi = m_\psi^3 / M_p^2$, and setting this equal to the Hubble parameter in eq (10) we find

$$R_{d\psi} \approx (M_p / m_\psi)^{4/3} R_\psi \quad (13)$$

Comparing $R_{d\psi}$ to R_ϕ , we see that $R_\phi/R_{d\psi} \approx m_\psi^2/M_p^{4/3} \tilde{m}^{-2/3} < 1$ for $m_\psi \sim 10^{-7} M_p$ and $\tilde{m} \sim 10^{-16} M_p$. At $R \geq R_\phi$ both ϕ and ψ are oscillating and if we compare eqs. (9) and (12) we see that $\rho_\psi/\rho_\phi \approx M_p^2/\phi_o^2 \gg 1$, so the Universe is still dominated by inflaton oscillations.

The baryon number density stored in the sfermion oscillations is⁷⁾

$$n_B = \epsilon(\phi_o^2/M_G^2) \rho_\phi / \tilde{m} \quad (14)$$

where $\epsilon(\phi_o^2/M_G^2)$ is the net baryon number produced in a ϕ decay with the C and CP violation given by ϵ . We expect the grand unification scale M_G to be within one order of magnitude of M_p in superstring models¹⁶⁾, whereas it could be lower in a general supersymmetric GUT. If ϕ 's decay first, the baryon to entropy ratio will be determined by eq. (14) and the entropy produced by inflaton decays. If ψ 's decay first, then we must check whether or not the radiation produced in ψ decays thermalize⁹⁾ before ϕ decay. If thermalization occurs before ϕ decay with the energy density in radiation $\rho_{r\psi} > \rho_\phi$, the AD mechanism will not be possible. On the other hand if thermalization takes place after ϕ decay and $\rho_{r\psi} > \rho_\phi$ the final baryon to entropy ratio will be again given by (14) and the entropy due to ψ decay.

To determine which field decays first we need the decay rate Γ_ψ and the decay rate for ϕ 's, $\Gamma_\phi \approx \tilde{m}^3/\phi^2$. From eqs. (8) and (12), $\phi^2 = \phi_o^2 (m_\psi^2/\tilde{m}^2) (R_\psi/R)^3$. Comparing Γ_ψ with Γ_ϕ , $\Gamma_\psi/\Gamma_\phi = \phi_o^2 m_\psi^5 (R_\psi/R)^3 / M_p^2 \tilde{m}^5 > 1$ implies that $R/R_\psi < 0(10^{12})$. Thus ψ 's will decay first if $R_{d\psi}/R_\psi < 0(10^{12})$. As one can see from eq. (13) $R_{d\psi}/R_\psi \approx 0(10^9)$. After ψ decays, the Universe becomes radiation-dominated with

$$\rho_{r\psi} = m_\psi^{2/3} M_P^{10/3} (R_\psi/R)^4 \quad (15)$$

for $R > R_{d\psi}$. We must now make sure that ϕ 's decay before the radiation in $\rho_{r\psi}$ thermalizes.

In the radiation-dominated period, the Hubble parameter is given by

$$H = m_\psi^{1/3} M_P^{2/3} (R_\psi/R)^2 \quad (16)$$

To find $R_{d\phi}$, the scale factor at ϕ decay, we must compare H and Γ_ϕ ,

$$R_{d\phi} = [m_\psi^{7/15} \phi_o^{2/5} M_P^{2/15} / \tilde{m}] R_\psi \quad (17)$$

using $\rho_{r\psi} > \rho_\phi$ at $R_{d\phi}$ which is easily verified. Following the argument in ref. 9, the scale factor at thermalization R_T is determined by comparing the thermalization rate to the Hubble parameter H in eq. (16). The thermalization rate is

$$\Gamma_T \sim n_\psi \sigma m_\psi M_P^2 (R_\psi/R)^3 (\alpha^2/m_\psi^2) (R/R_{d\psi})^2 \sim \alpha^2 (m_\psi^{5/3}/M_P^{2/3}) (R_\psi/R) \quad (18)$$

where the number density of relativistic particles initially is n_ψ , σ is their elastic scattering cross-section and α the gauge fine structure constant. Thus thermalization occurs at

$$R_T \sim [(M_P/m_\psi)^{4/3}/\alpha^2] R_\psi \sim R_{d\psi}/\alpha^2 \quad (19)$$

We can now check that $R_{d\phi} < R_T$, since $R_{d\phi}/R_T \sim m_\psi^{9/5} \phi_o^{2/5} \alpha^2 / \tilde{m} M_p^{6/5} \sim 10^2 \alpha^2 \sim 10^{-1}$ for $\alpha^2 \sim 10^{-3}$ so that indeed ϕ 's decay before thermalization.

At R_T , the final baryon to entropy ration can be computed using eq. (14) for n_B and eq. (15) to find the entropy density. At R_T the entropy density is

$$s = [\rho_{r\psi}(R_T)]^{3/4} \approx m_\psi^{9/2} \alpha^6 / M_p^{3/2} \quad (20)$$

and the baryon to entropy ratio is

$$\frac{n_B}{s} = \frac{\epsilon \phi_o^4 m_\psi^{3/2}}{m_G^2 M_p^{5/2} \tilde{m}} \quad (21)$$

Taking $\epsilon \sim 10^{-3}$, $\phi_o^2 \sim 10^{-7} M_p^2$, $M_G \sim 10^{-1} M_p$, $m_\psi \sim 10^{-7} M_p$, $\tilde{m} \sim 10^{-16} M_p$, we find amazingly that $n_B/s \sim 3 \times 10^{-10}$! This should be compared with the ratio preferred by big bang nucleosynthesis¹⁷⁾, $n_B/s \approx 7 \times 10^{-11}$ today. Clearly such a solution ends the conflict between the baryon asymmetry and reheating after inflation.

Once the radiation from inflaton decays thermalizes, the temperature of the Universe is $T \sim s^{1/3} \sim 0(10^4) \text{ GeV}$. If the above solution of n_B/s is to account for the present baryon asymmetry then very little entropy should have been produced between $T \sim 10^4 \text{ GeV}$ and today. In ref. 9, we discussed in detail two possible sources of entropy. One was the electroweak transition at $T \lesssim 10^3 \text{ GeV}$ in which we considered specifically the superstring-inspired model of ref. 16. There we found at most a factor of

$0(10)$ dilution was possible and that in some cases the transition is in fact second order. In standard supersymmetric models, the transition is again model-dependent and it must be checked that entropy production is not significant. A similar bound on entropy production applied to the old out-of-equilibrium decay scenarios.

A second source of entropy production was the presence of additional scalar fields in superstring models which could also oscillate, decay and produce entropy. In ref. 9 we considered the effects of gauge singlet fields X and Y (sometimes referred to as S and T) but did not specifically assume the presence of an inflaton. It is possible to show (though somewhat tedious) that these fields never dominate the Universe unless $m_{X,Y} < 10^{-12} M_p$ whereas their masses are expected to be $0(10^{-7} - 10^{-8}) M_p$. Here we will only outline the argument against entropy production by (say) Y decays. Our results are summarized in the Figure.

Depending on the mass of the Y field, we can distinguish several different cases. First, if $m_Y > m_\psi$, since the Y decay rate $\Gamma_Y \sim M_Y^3/M_p^2 > \Gamma_\psi$, the Y field would decay first. But since at the end of inflation $\rho_\psi > \rho_Y$ (by definition it is the energy density of the inflaton which dominates the Universe), and until Y decays ρ_ψ/ρ_Y is fixed, after Y decays $\rho_Y \sim R^{-4}$ while $\rho_\psi \sim R^{-3}$ and ρ_ψ/ρ_Y in fact grows. This situation is displayed by the dot-dashed curve in the Figure. Clearly Y can be neglected in this case.

If $m_Y < m_\psi$, ψ 's decay before Y 's and ρ_Y potentially has a chance of dominating. Recall that Y oscillations begin when $H \sim m_Y$ or when the scale factor $R = R_Y$, $R_Y \approx (m_\psi/m_Y)^{2/3} R_\psi$ and $R_Y/R_{d\psi} \approx m_\psi^2/(M_p^{4/3} m_Y^{2/3})$. That is, Y 's will begin to oscillate before ψ decay so long as $m_Y > \Gamma_\psi \sim 1 \text{ MeV}$. At

$R \geq R_Y$, $\rho_Y = m_\psi^2 Y_o^2 (R_\psi/R)^3$ as in eq. (12), where Y_o is the initial amplitude of Y oscillations and we see that $\rho_Y/\rho_\psi = Y_o^2/M_p^2 \ll 1$, i.e. ψ still dominates at R_Y . Now since ψ decays first ($m_\psi > m_Y$), at Y decay, depending on m_Y , ρ_Y may eventually dominate the energy density. To determine the critical value of m_Y , we can use eq. (16) for H during the radiation-dominated period to find R_{dY} , the scale factor at Y decay, $R_{dY} = (m_\psi^{1/6} M_p^{4/3}/m_Y^{3/2})R_\psi$. Then using eq. (15) we can compare $\rho_{r\psi}$ and ρ_Y , $\rho_{r\psi}/\rho_Y = m_Y^{3/2} M_p^2/m_\psi^{3/2} Y_o^2$ at R_{dY} . For $\rho_{r\psi}/\rho_Y = 1$, we see that $m_Y^{3/2} = m_\psi^{3/2} (Y_o^2/M_p^2)$, for $Y_o^2 \sim 10^{-7} M_p^2$ (true if $m_Y^2 \ll H^2$, which would be the case since $H \sim m_\psi$) and ρ_Y only dominates if $m_Y \leq 0(10^{-12})M_p$. This case is shown in the Figure by the dashed and dotted curves. We can conclude therefore that for $m_Y \geq 10^{-12}M_p$ as we expect (actually it is only the lightest scalar which is important⁹⁾), we can neglect the effects of Y oscillations on the baryon to entropy ratio. In a similar way⁹⁾ the Polonyi and gravitino problems are resolved.

Finally, we would like to discuss the possibility that density fluctuations can be produced by the sfermion oscillations. Sfermion density fluctuations are produced in a similar way as fluctuations due to axion oscillations¹⁸⁾. In the axion case, density fluctuations $\delta\rho_a/\rho_a \sim H/a_o$ where H is the Hubble parameter during inflation and a_o is the initial amplitude for axion oscillations, for $H \ll a_o$. Because the damping timescale (axion decay rate) is longer than the age of the Universe, ρ_a eventually comes to dominate the total mass density¹⁹⁾ and hence one would have an axion-dominated Universe with fluctuations $\delta\rho_a/\rho \approx \delta\rho_a/\rho_a \sim H/a_o$. In this case, the bulk of the radiation was produced during reheating after inflation with perhaps adiabatic fluctuations $\delta\rho_\psi/\rho_\psi$. Because $\delta\rho_a/\rho$ is

decoupled from the perturbations in the radiation $\delta\rho_\psi/\rho_\psi$, axion perturbations are isothermal.

In the situation we are interested in, density fluctuations $\delta\rho_\phi/\rho_\phi$ are also decoupled from $\delta\rho_\psi/\rho_\psi$ and since the bulk of the radiation again comes from ρ_ψ , the sfermion fluctuations are isothermal. The major difference between the sfermion case and the axion case is that ϕ 's never come to dominate the energy density, i.e. ϕ 's decay while $\rho_\phi < \rho_{r\psi}$. Thus $\delta\rho_\phi/\rho_{\text{tot}} \sim (\delta\rho_\phi/\rho_\phi) (\rho_\phi/\rho_{\text{tot}}) \sim (\delta\rho_\phi/\rho_\phi) (\rho_\phi/\rho_{r\psi} (R_{d\phi}))$. Taking $\delta\rho_\phi/\rho_\phi \sim H/\phi_0 \sim 0(10^{-4})$ and $\rho_\phi/\rho_{r\psi} \sim \frac{m_\psi^{9/5} \phi_0^{12/5}}{m_p^{16/5}} \sim 10^{-5}$ at $R_{d\phi}$ we find $\delta\rho_\phi/\rho_{\text{tot}} \sim 0(10^{-9})$, which is far too small to be of any significance.

In conclusion, we have calculated the baryon to entropy ratio produced by sfermion oscillations in a generic inflationary model to be $n_B/s \sim 3 \times 10^{-10} (\epsilon/10^{-3})$ where ϵ measures the amount of C and CP violation.

This includes our calculation of the initial amplitude, ϕ_0 , of these oscillations based on De Sitter fluctuations during inflation. We have also included the dilution due to entropy production by inflaton decays for a generic inflationary models. These results are valid for both standard supersymmetric models and superstring-inspired models where we have shown entropy production by additional scalar fields to be negligible. Finally, we computed the magnitude of isothermal density fluctuations produced by the sfermion oscillations and found $\delta\rho_\phi/\rho_{\text{tot}} \sim 0(10^{-9})$. At this stage, we feel that the Affleck-Dine mechanism for the production of a Universal baryon asymmetry is the most promising scenario in the context of the models studied.

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Figure Caption

The evolution of the energy density of the inflaton ψ and a scalar field Y as a function of the scale factor. The solid curve represents ρ_ψ and $\rho_{r\psi}$ (the radiation energy density produced by ψ decay at $R_{d\psi}$) the kinks in the curve occur at R_ψ when ψ 's begin to oscillate ($\rho_\psi \sim R^{-3}$) and at $R_{d\psi}$ when ψ 's decay ($\rho_{r\psi} \sim R^{-4}$). If $m_Y > m_\psi$ (dot-dashed) Y 's decay before ψ 's and ρ_Y never dominates the total energy density. If $\Gamma_\psi < m_Y < m_\psi$ (dashed), Y 's decay after ψ 's but ρ_{rY} remains less than $\rho_{r\psi}$ unless $m_Y < 0(10^{-12})M_p$ (dotted). Finally if $m_Y < \Gamma_\psi \sim 0(1) \text{ MeV}$ (double dot-dashed) then Y 's only begin to oscillate after ψ decay and eventually dominate the total energy density. We expect $m_Y \sim 0(10^{-7})M_p$.