Flatt Direction Dynamics in a
Non-Topological Soliton-Dominated
Universe

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Abstract

In running mass inflation and hybrid inflation models it is possible that the
inflaton field will fragment into non-topological solitons, resulting in a highly
inhomogeneous post-inflation era prior to reheating. In supersymmetric models
with a conventional homogeneous post-inflation era, the dynamics of flat direc-
tion scalars are determined by $cH^2$ corrections to the mass squared terms (where
$|c| \approx 1$), coming from F-terms in the early Universe combined with Planck-scale
suppressed interactions. Here we reconsider the mass squared corrections for a
Universe dominated by inflatonic non-topological solitons. We show that in this
case the dynamics of a coherently oscillating flat direction scalar are typically the
same as for the case where there is no significant mass squared correction, even
in the vicinity of the non-topological solitons. Therefore the dynamics of flat
direction scalars in a Universe dominated by inflatonic non-topological solitons
are equivalent to the case $c = 0$ of a homogeneous Universe.

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1 Introduction

In the Minimal Supersymmetric (SUSY) Standard Model (MSSM) [1] it is known that there are a large number of possible flat directions formed from combinations of squark, slepton and Higgs fields [2]. In the early Universe it is likely that these flat directions will develop a non-zero expectation value during inflation, and will subsequently form Bose condensates corresponding to coherently oscillating scalar fields. The cosmology of the formation and decay of flat direction condensates is of fundamental importance to the cosmology of the MSSM and its extensions [3]. They may be the origin of the baryon asymmetry of the Universe via the Affleck-Dine (AD) mechanism [4]. Flat direction condensates may also account for a density of dark matter via the formation [5] and late decay [6, 7] of Q-balls, allowing the baryon and dark matter number densities to be related. In addition, a SUSY flat direction condensate could possibly account for the energy density perturbations responsible for structure formation, serving as a SUSY curvaton [8, 9, 10, 12, 13, 14, 15].

The dynamics of flat direction fields in the post-inflation era are largely determined by corrections to the mass squared terms coming from SUSY breaking F-terms associated with fields which contribute to the energy density, $\rho$. These F-terms, when combined with Planck-scale suppressed interactions, result in a mass squared term for the flat direction scalar of the order of $|F|^2/M^2$, where $M = M_{Pl}/\sqrt{8\pi}$ and where $M_{Pl}$ is the Planck mass. For a conventional post-inflation Universe with a homogeneous energy density due to coherent inflaton oscillations after inflation, these corrections are of the form $cH^2$ with $|c|$ of the order of 1 [16, 17].

However, in important examples of SUSY inflation models it has become clear that the energy density in the inflaton field may not remain homogeneous after inflation ends. It has been shown that in SUSY running mass inflation models [18] it is possible that quantum fluctuations of the inflaton field will grow, resulting in formation of non-topological solitons (NTS) [19, 20, 21]. In addition, it has been shown that in SUSY hybrid inflation models [22, 19] quantum fluctuations of the inflaton can grow to form non-linear fragments. It possible that such fragments will evolve into quasi-
stable inflatonic NTS (for the case of a real inflaton field we refer to these as 'inflaton condensate lumps' [19]). This has yet to be established numerically for fully realistic SUSY hybrid inflation models [23], although some numerical evidence for NTS formation in hybrid inflation models exists [20]. For the case of D-term inflation models [24], fragmentation is expected to occur if $\lambda \gtrsim 0.1g$ [19, 25], whilst for F-term inflation models [26] it is expected to always occur [19, 25].

In SUSY inflation models the inflaton is generally a complex scalar field. As a result, it is possible for Q-balls to form [18]. This was demonstrated numerically for the case of a running mass chaotic inflation model, where it was shown that initially neutral inflaton condensate lumps fragment into Q-ball, anti-Q-ball pairs (the most stable configuration) [18]. SUSY hybrid inflation models also have Q-ball solutions, formed from a combination of a complex inflaton field and a real symmetry breaking scalar field [27]. If neutral inflaton condensate lumps form at the end of SUSY hybrid inflation then it is likely that these will subsequently fragment into inflatonic Q-ball pairs, as in the running mass case.

Thus it is possible that in both running mass inflation models and hybrid inflation models the post-inflation era will be highly inhomogeneous, formed initially of inflaton condensate lumps and subsequently of inflatonic Q-balls, which eventually decay to reheat the Universe [19, 29].

In this paper we wish to discuss how SUSY flat direction scalar field dynamics, in particular the dynamics of a coherently oscillating scalar field, would be modified during an inflatonic NTS-dominated post-inflation era. We first give a qualitative discussion of how the scalar field dynamics might be expected to change. The energy density of the Universe after inflation is entirely in the form of inflaton condensate lumps or inflatonic Q-balls, with little or no energy outside the NTS where the inflaton field is strongly exponentially suppressed. Thus we might expect that there will effectively be no scalar mass squared corrections outside the NTS, since $\rho$ and $|F|^2$ are very close to zero outside the NTS. However, in the vicinity of the NTS a much larger than average inflaton energy density will occur. This suggests that the now space-dependent inflaton energy density and associated F-term will induce a very
large mass squared correction ($\gg H^2$) for the flat direction field in the vicinity of the inflatonic NTS. If so, this would be expected to have a strong effect on the dynamics of the coherently oscillating flat direction fields. However, we will show that this is not the case; even at the centre of the inflatonic NTS the dynamics of a coherently oscillating scalar field will be shown to be essentially the same as in the absence of any mass squared correction. This is because the gradient energy term, due to the spatial distortion of the flat direction field in the vicinity of the inflatonic NTS, cancels the effect of the effective mass squared term in the vicinity of the NTS such that the amplitude of the coherent oscillations is essentially unaltered. As a result, the dynamics of a coherently oscillating flat direction scalar in a NTS-dominated Universe are equivalent to the case of a homogeneous Universe with $c = 0$.

The paper is organised as follows. In Section 2 we discuss the solution of the flat direction scalar field equation for a coherently oscillating flat direction scalar in the presence of an inflatonic NTS and show that the oscillation amplitude is typically unaffected by the presence of the NTS. We also give a physical argument to support this conclusion. In Section 3 we consider the consequences for Affleck-Dine baryogenesis and SUSY curvaton dynamics. In Section 4 we present our conclusions.

## 2 Gaussian Non-Topological Solitons and Flat Direction Condensates

In order to understand the effect of inflatonic NTS on the dynamics of a coherently oscillating flat direction condensate, we will consider the case of a NTS with a Gaussian profile. We expect that the Gaussian profile will be a reasonable approximation to the thick-walled NTS which are expected to form during inflaton condensate fragmentation. A particular example is the inflatonic Q-ball which forms at the end of running mass chaotic inflation [18], which has the same form as the Q-balls which form from squark and slepton flat direction condensates during Affleck-Dine baryogenesis [7]. SUSY hybrid inflation Q-balls also have a near Gaussian profile for the inflaton field [27].
We will consider the potential of the complex inflaton $S$ to consist in general of a mass squared term plus an attractive self-interaction term which allows the formation of NTS,

$$V(S) = m_S^2 |S|^2 + V_{int}(S).$$  \(1\)

For the case of the running mass chaotic inflation model the interaction potential has the form [18],

$$V_{int}(S) = K m_S^2 |S|^2 \ln \left( \frac{|S|^2}{\Lambda^2} \right),$$  \(2\)

where $K < 0$ and $\Lambda$ is a renormalization scale. (This has the same form as the potential of MSSM flat directions involving squarks [6, 7].) $m_S^2$ is assumed to originate from soft SUSY breaking [18]. Eq. (2) leads to Gaussian Q-ball solutions of the $S$ field equation of the form [7]

$$S = \frac{s(r)}{\sqrt{2}} e^{i\omega t},$$  \(3\)

where

$$s(r) \approx s_0 e^{-r^2/R^2},$$  \(4\)

$$R \approx \frac{\sqrt{2}}{|K|^{1/2} m_S},$$  \(5\)

$$\omega^2 = \omega_0^2 + m_S^2 (1 + K),$$  \(6\)

and

$$\omega_0^2 \approx 3 |K|m_S^2.$$  \(7\)

The values of $\omega$ and $R$ above are for the specific case of the interaction potential of Eq. (2). However, the form of the Gaussian Q-ball solution given by Eq. (3) and Eq. (4) is quite general. Moreover, we expect the magnitudes of $\omega$ and $R$ to be determined in most cases by the same dynamical mass scale in the field equations, such that $\omega R$ will be typically of the order of 1. Thus our results should apply to thick-walled inflatonic Q-balls in general, with the interaction potential of Eq. (2) providing a specific example.

The initial spacing between the NTS when they first form will typically be of the order of the radius of the NTS [19]. Thus after a period of Universe expansion the
spacing will be large compared with the NTS radius. From the Gaussian form for the inflatonic NTS amplitude, Eq. (4), the inflaton field will be exponentially suppressed outside an NTS, such that its value is effectively zero. So long as the spacing between the NTS is sufficiently large compared with their radius ($\delta x/R \approx 10$ will result in a suppression factor $e^{-100}$, where $\delta x$ is the spacing between the NTS), including more than one NTS will not alter the nearly zero value of the inflaton field and so mass squared correction outside the NTS. The NTS effectively do not see each other because of the exponential suppression of the inflaton field outside the NTS. As a result, we can study flat direction scalar dynamics by considering the flat direction scalar field equation in the background of a single NTS.

An important issue is the equation of motion of the flat direction scalar in the inhomogeneous background of a NTS-dominated Universe. In this case the usual Friedmann-Robertson-Walker metric is not strictly valid and a non-trivial calculation of the inhomogeneous metric due to the ensemble of NTS is necessary to obtain the correct scalar field equation including the effect of the expansion of the Universe. However, we will be concerned with the effect of the flat direction scalar mass term induced in the vicinity of the inflatonic NTS on the dynamics of a coherently oscillating flat direction field. Since this mass term is much larger than the expansion rate, the effect of gravitational corrections to the scalar field equation of motion due to the expansion of the Universe will be negligible (i.e. the time scale over which the flat direction field changes due to the effect of the NTS will be small compared with $H^{-1}$). In addition, it has been shown that gravitational effects typically play no role in the NTS solution itself [28]. Therefore flat space may be considered when solving the field equation for the flat direction scalar.

The potential of a SUSY flat direction scalar $\Phi$ is expected to purely consist of a conventional gravity-mediated SUSY breaking mass squared term,

$$V(\Phi) = m_{\phi}^2 |\Phi|^2,$$

where $m_{\phi} \approx 100 \text{ GeV} - 1 \text{ TeV}$. (We assume the coherent oscillations have an amplitude small enough that possible non-renormalizable corrections to the flat direction superpotential may be neglected.) We will consider a non-renormalizable Planck-scale
suppressed interaction term between the inflaton field and flat direction field of the form which is generally expected to arise in the low energy effective theory from supergravity [2, 16, 17] and which is not excluded by any symmetry,

\[ \mathcal{L}_{\text{int}} = -\frac{\lambda}{M^2} \int d^4\theta S^\dagger S \Phi^\dagger \Phi \equiv -\frac{\lambda}{M^2} |F_S|^2 \Phi^\dagger \Phi, \]  

where \( M = M_{\text{Pl}}/\sqrt{8\pi} \) and \( \lambda \approx 1 \). (This is Equation 8 of [2]) \( F_S \) is the F-term of the inflaton scalar,

\[ |F_S|^2 = \left( \partial_\mu S^\dagger \partial^\mu S + \left| \frac{\partial W}{\partial S} \right|^2 \right) |\Phi|^2, \]  

where \( W \) is the superpotential. This term is of the generic form which leads to the \( cH^2 \) correction to scalar mass squared terms in the case of a conventional homogeneous post-inflation Universe. However, its effect must be reconsidered in the case of an inhomogeneous NTS-dominated Universe.

In the case of a conventional homogeneous post-inflation Universe dominated by a coherently oscillating inflaton condensate we have \[ |F_S|^2 = \dot{\hat{S}}^2/2 + V(S) \equiv \rho_S, \]  
where \( V(S) = |\partial W/\partial S|^2 \) is the SUSY inflaton potential. Therefore \[ |F_S|^2/M^2 = \rho_S/M^2 \equiv 3H^2/M^2, \]  
resulting in a mass-squared correction of the order of \( H^2 \). \( \lambda > 0 \) \((< 0)\) would then correspond to having \( c < 0 \) \((c > 0)\) in the \( cH^2 \) term. Other possible Planck-scale suppressed interactions with the SUSY breaking inflaton F-term all lead to similar corrections to the mass-squared terms [2]. In the following we will consider Eq. (9) as a typical example.

In order to study the effect of the inflatonic Q-ball (which we refer to as the S-ball in the following) on the dynamics of a coherently oscillating flat direction scalar \( \Phi \), we will introduce a single Gaussian S-ball solution into the \( \Phi \) scalar field equation. As discussed above, this is justified if the separation of the S-balls is sufficiently large compared with their radius \( R \) \((\delta x > 10R)\), in which case there will be an extreme exponential suppression of the contribution of the other S-balls to the inflaton field in the vicinity of a given S-ball. In the space between the S-balls the extreme suppression will result in effectively no inflaton field or energy density.

We will consider the case where the interaction \( \mathcal{L}_{\text{int}} \) comes purely from the derivative terms in Eq. (9). This is exactly true for the case where the \( S \) scalar potential
is assumed to come purely from soft SUSY breaking terms and radiative corrections, such that there is no superpotential for $S$. More generally, the effect of a superpotential will be to introduce a potential term into the inflaton F-term which will have at most the same magnitude as the derivative terms in the S-ball solution. Thus we expect similar results in the more general but model dependent case where there is a superpotential for the inflaton.

Since we are interested in the effect of the S-ball on a coherently oscillating flat direction condensate, we will consider a solution for $\Phi$ in the presence of an S-ball of the form,

$$\Phi(r,t) = \frac{\phi(r)}{\sqrt{2}} \sin(m\phi t),$$

where $\phi(r)$ is a space-dependent amplitude for the coherent oscillation in the presence of the S-ball and $r$ is the radius from the S-ball centre. The $\Phi$ scalar field equation is then

$$\ddot{\Phi} - \nabla^2 \Phi = - \left( m^2 \phi \Phi - \frac{\lambda}{M^2} f(S) \Phi \right),$$

where

$$f(S) = |\dot{S}|^2 + |\nabla S|^2,$$

and where as discussed above we consider the flat space scalar field equation. With $S$ given by the S-ball solution, Eq. (3), $f(S)$ becomes

$$f(S) = \left( \omega^2 + \frac{4r^2}{R^4} \right) \frac{s^2}{2} e^{-2r^2/R^2}.$$

Substituting the coherently oscillating $\Phi$ solution, Eq. (11), into Eq. (12) then gives the equation for the flat direction scalar oscillation amplitude in the background of an S-ball,

$$\phi''(r) + \frac{2\phi'(r)}{r} = - \frac{\lambda}{M^2} \left( \omega^2 + \frac{4r^2}{R^4} \right) \frac{s^2}{2} e^{-2r^2/R^2} \phi(r),$$

where $'$ denotes differentiation with respect to $r$.

We will look for a solution which is valid at $r/R \ll 1$, since the greatest effect of the inflaton energy density on the amplitude of coherent oscillations will be found at the centre of the S-ball. Suppose we consider $\phi(r) = \phi_0 s_o(r)$, where $s_o$ is the value of
the inflaton field at the centre of the S-ball. Far from the S-ball \((r/R >> 1)\) we expect that the function \(g_{s_o}(r) \to 1\) if \(\phi_o\) is the amplitude of the coherent oscillations in the absence of the S-ball. This is reasonable both physically and by inspection of the right hand side of Eq. (15), which rapidly tends to zero as \(r/R\) becomes large compared with 1, such that \(\phi = \phi_o\) (constant) becomes a solution at large \(r\). In addition, as \(s_o \to 0\) the function \(g_{s_o}(r) \to 1 \forall r\), since there is no S-ball in this limit. Therefore if we obtain a solution \(\phi(r)\) valid at small \(r\) then in order to relate it to the coherent oscillation amplitude far outside the S-ball we need only identify \(g_{s_o}(r)\) by taking \(s_o \to 0\) and setting \(\phi(r) = \phi_o\) in this limit.

To find a solution valid at \(r/R < 1\), we first change variable to \(y = \log(\phi)\). Then Eq. (15) becomes

\[
y'' + y' \frac{2y'}{r} = -\frac{\lambda}{M^2} \left( \omega^2 + \frac{4r^2}{R^4} \right) \frac{s_o^2}{2} e^{-2r^2/R^2} .
\]

We then look for a solution of the form

\[
y = A e^{-2r^2/R^2} + C .
\]

Substituting this into Eq. (15), the left hand side becomes

\[
y'' + y' \frac{2y'}{r} \equiv \left( -\frac{12}{R^2} + \frac{16r^2}{R^4} \right) A e^{-2r^2/R^2} + \frac{16r^2}{R^4} A^2 e^{-4r^2/R^2} .
\]

Thus as \(r/R \to 0\),

\[
y'' + y' \frac{2y'}{r} \to -\frac{12}{R^2} A e^{-2r^2/R^2} .
\]

This has the same form as the right hand side of Eq. (16) in the limit \(r/R \ll \omega R/2\). Thus a solution valid for small \(r/R\) is

\[
\phi(r) = e^y = e^C \exp \left( A e^{-2r^2/R^2} \right) ,
\]

where

\[
A = \frac{\lambda \omega^2 R^2 s_o^2}{24M^2} .
\]

We see that as \(s_o \to 0\), \(\phi(r) \to e^C\). Therefore we set \(\phi_o = e^C\), such that

\[
\phi(r) = \phi_o \exp \left( \frac{\lambda \omega^2 R^2 s_o^2 e^{-2r^2/R^2}}{24M^2} \right) .
\]
Thus at the centre of the S-ball the amplitude of the coherently oscillating flat direction scalar as a function of $s_o$ is given by

$$\phi(r = 0) = \phi_o \exp(A) \equiv \phi_o \exp \left( \frac{\lambda \omega^2 R^2 s_o^2}{24M^2} \right).$$

(23)

From Eq. (23) we see that a significant change of the $\Phi$ amplitude at $r = 0$ relative to its value in the absence of the S-ball, $\phi_o$, is possible only if $A \gtrsim 1$. This requires that

$$\frac{s_o}{M} \sim \sqrt{\frac{24}{|\lambda| \omega R}}.$$  

(24)

Since $\omega$ and $R$ for a Gaussian S-ball are typically determined by the same dynamical mass scale (the $S$ mass), $\omega R$ is expected to be not very much larger than 1. Therefore unless $s_o$ is close to the Planck scale, the S-ball will typically have little effect on the amplitude of flat direction coherent oscillations. Thus in the case of a NTS-dominated Universe the flat direction coherent oscillations are essentially unaltered from the case where there is no correction to the flat direction mass squared term i.e. the oscillations are equivalent to the case $c = 0$ of a homogeneous post-inflation Universe.

At first sight this result is surprising. It would seem that the large inflaton energy density and F-term would induce a large mass squared term in the vicinity of the S-ball, much larger than the order $H^2$ correction expected in the case of a homogeneous post-inflation Universe. The reason that the S-ball has little effect on the coherent oscillations of the flat direction scalar is that as the flat direction oscillation amplitude distorts under the influence of the induced mass squared term in the vicinity of the S-ball, the gradient energy of the now space-dependent amplitude counteracts the effect of the mass squared term. It happens that the energy density of the flat direction scalar is minimized when the distortion in the amplitude is negligibly small. To see this physically, consider the case where $\lambda > 0$ in Eq. (9), corresponding to a negative mass squared term for the flat direction scalar in the vicinity of the S-ball. The contribution of the induced mass squared term $m_{eff}^2 \approx -\rho_S/M^2$ in the vicinity of the S-ball (energy density $\approx \rho_S$) to the energy density of the coherently oscillating flat direction scalar is then,

$$\rho_{m_{eff}} \approx m_{eff}^2 \phi_o + \Delta \phi_o)^2 \approx -\frac{\rho_S}{M^2}(\phi_o + \Delta \phi_o)^2 \approx -\frac{s_o^2}{M^2 R^2}(\phi_o + \Delta \phi_o)^2,$$

(25)
where the energy density inside the S-ball is expected to be of the order of \(s_o^2/R^2\). Here \(\phi_o\) is the amplitude in the absence of the S-ball and \(\Delta \phi_o\) is the change in the amplitude in the vicinity of the S-ball. The negative mass squared term will cause the amplitude of the coherent oscillations to increase in the vicinity of the S-ball, resulting in a gradient energy term \(\rho_{\text{grad}} \approx (\Delta \phi_o/R)^2\). Thus the energy density as a function of \(\Delta \phi_o\) in the vicinity of the S-ball is given by

\[
\rho(\Delta \phi_o) \approx \rho_{\text{m,eff}} + \rho_{\text{grad}} = -\frac{s_o^2(\phi_o + \Delta \phi_o)^2}{M^2 R^2} + \frac{\Delta \phi_o^2}{R^2}.
\]

(26)

This is minimized at

\[
\frac{\Delta \phi_o}{\phi_o} \approx \frac{s_o^2}{M^2}.
\]

(27)

Therefore if \(s_o\) is small compared with the Planck scale then the shift of flat direction oscillation amplitude is negligible, in agreement with Eq. (24).

For the case of the S-ball associated with running mass inflation the condition Eq. (24) becomes

\[
\frac{s_o}{M} \gtrsim \sqrt{\frac{12}{|\lambda|}} \frac{|K|^{1/2}}{(1 + 2|K|)^{1/2}}.
\]

(28)

The value of \(s_o\) when the S-balls form is \(s_o \approx 10^{-2}|K|M\), with \(|K| \approx 0.01 - 0.1\) [18]. Thus in this model the condition Eq. (24) is not satisfied and so there will be no significant effect on the dynamics of a coherently oscillating flat direction scalar i.e. the dynamics are equivalent to the case \(c = 0\) of a conventional homogeneous post-inflation era. Similarly, if we consider the case of D-term hybrid inflation and assume that the value of \(s_o\) is characterised by the value of the inflaton field at the end of inflation, \(s_o \approx 10^{16}\) GeV [19, 24], then the flat direction scalar dynamics will again effectively correspond to the case \(c = 0\). Therefore we expect that the flat direction dynamics during an NTS-dominated post-inflation era will typically be equivalent to the case \(c = 0\) in a conventional homogeneous post-inflation cosmology.

The end of the NTS-dominated era is expected to occur via the decay of the NTS. The value of \(H\) at reheating is then given by \(H(T_R) \approx k_{T_R} T_R^2/M_{\text{Pl}}\), where \(T_R\) is the reheating temperature and \(k_{T_R} \approx 20\). After reheating the energy density will be homogeneous and conventional \(cH^2\) corrections with \(|c| \approx 1\) will apply. Assuming
that $T_R \lesssim 10^{8-9}$ GeV in order to evade the thermal gravitino problem [30] implies that
reheating will occur at $H(T_R) \lesssim 1$ GeV. Therefore when the $cH^2$ terms switch on they
will already be small compared with the gravity-mediated soft SUSY breaking mass
squared terms in the flat direction potential (of the order of $m_{\tilde{W}}^2$) and so will play no
role in the dynamics of the flat direction scalars.

In the above we have assumed that the flat direction condensate does not modify
the Gaussian NTS solution. The interaction term Eq. (9) also contributes terms to
the $S$ field equation,

$$\mathcal{L}_{\text{int}} = \ldots + \frac{\lambda}{M^2} \left( \partial_\mu \Phi^\dagger \partial^\mu \Phi \right) |S|^2 \rightarrow \frac{\lambda}{2M^2} m_\phi^2 \phi_o^2 \cos^2 (m_\phi t) |S|^2 .$$

(29)

The condition for the flat direction condensate to have no effect on the S-ball solution
is that this term is small compared with $m_S^2 |S|^2$. This requires that,

$$\frac{\phi_o}{M} < \frac{m_S}{\sqrt{\lambda m_\phi}} .$$

(30)

Since $m_S \gtrsim m_\phi$ is expected (where $m_\phi \approx 100$ GeV is the SUSY breaking mass term),
this will typically be satisfied for all $\phi_o \lesssim M$.

3 Consequences for Affleck-Dine Baryogenesis and
SUSY Curvatons

One common application of SUSY flat direction dynamics is to Affleck-Dine (AD)
baryogenesis [4]. In the conventional homogeneous case we expect $|c| \approx 1$ after the
end of inflation. If $c$ is positive ($c \approx 1$) then the effect of the positive order $H^2$ term
will be to drive damped oscillations of the flat direction scalar, such that the amplitude
of the AD scalar oscillations is strongly suppressed by the time the gravity-mediated
SUSY breaking mass squared term comes to dominate the dynamics and the baryon
asymmetry in the condensate is fixed. As a result, AD baryogenesis is effectively
ruled out for $c \approx 1$. Therefore the possibility of AD baryogenesis depends upon the
(typically unknown) sign of $c$. For $c \approx -1$, the flat direction field will roll away from
zero until the effect of non-renormalizable terms in the scalar potential stabilises the
field at the minimum of its potential. In this case the initial amplitude of the AD scalar oscillations is fixed by the dimension of the non-renormalizable superpotential terms which lift the flat direction potential and introduce the B and CP violation necessary for baryogenesis [2, 4, 6]. As a result, the baryon asymmetry is fixed by the reheating temperature $T_R$ and the dimension $d$ of the non-renormalizable superpotential term. For $d = 4$ the reheating temperature must be of order $10^8$ GeV in order to generate the observed asymmetry, whilst for $d = 6$ the reheating temperature must be around 1 GeV. (Even values of $d$ are required if R-parity is conserved [2, 7], as suggested by the absence of dangerous renormalizable B- and L-violating contributions to the MSSM superpotential [1].) Thus in the case of a homogeneous post-inflation cosmology with $|c| \approx 1$, the Affleck-Dine mechanism can only function if $c$ is negative and the reheating temperature must then either be high (close to the gravitino upper limit) or low (1 GeV or less) if R-parity is conserved.

In contrast, in the NTS-dominated case we have effectively $c \approx 0$. Therefore, the amplitude of the AD scalar remains fixed until the gravity-mediated SUSY breaking mass squared term comes to dominate the dynamics, as in the original Affleck-Dine baryogenesis scenario [4]. Thus in the NTS-dominated Universe the AD mechanism can, at least in principle, always generate the baryon asymmetry i.e. the AD scalar is never damped to zero. In addition, the asymmetry is not purely determined by $d$ and $T_R$ as in the $|c| \approx 1$ case, but also by the amplitude of the AD scalar at the end of inflation. This should allow a wider range of MSSM flat directions and reheating temperatures to be compatible with the observed asymmetry.

More recently, it has been suggested that a coherently oscillating scalar field (a 'curvaton') could be the source of cosmological density perturbations [8, 9, 10, 11, 12, 13, 14, 15]. A SUSY flat direction scalar could serve as a curvaton if its coherent oscillations can dominate the Universe before they decay, which requires a sufficiently large initial amplitude. Similar considerations then apply as in the case of the AD mechanism. In a conventional homogeneous post-inflation era with $|c| \approx 1$, a positive value of $c$ will cause the curvaton to rapidly damp away after inflation. In this case the flat direction scalar energy density will be too small to dominate the Universe before
it decays. Thus $c \approx -1$ is necessary. However, this case also has a possible problem. If the curvaton rolls to the minimum of its potential, as determined by the negative mass squared term and non-renormalizable superpotential terms, then oscillations of the curvaton about this minimum will damp the quantum fluctuations which lead to the energy density perturbations [15]. This will then require a large expansion rate during inflation to produce a sufficiently large quantum fluctuation, which may not be consistent with small enough adiabatic energy density perturbations from inflaton quantum fluctuations [15]. Both of these problems may be avoided in the case of a NTS-dominated post-inflation era, since effectively $c = 0$ in this case.

These examples make it clear that an inhomogeneous post-inflation era would have significant consequences for SUSY cosmology.

4 Conclusions

We have considered the effect of a non-topological soliton-dominated post-inflation era on the dynamics of a coherently oscillating SUSY flat direction scalar. We have shown that the flat direction dynamics during a NTS-dominated era are typically not affected by the Planck-suppressed interactions which would generate a $cH^2$ mass squared term in the case of a conventional post-inflation cosmology. Thus during the NTS-dominated era we expect that the flat direction dynamics will effectively be equivalent to the case $c = 0$ of a homogeneous post-inflation era.

We have considered some consequences of a NTS-dominated era for SUSY cosmology. Affleck-Dine baryogenesis will occur as in the original scenario with no order $H^2$ corrections. As a result, AD baryogenesis will become a more general possibility, since there will be no positive $H^2$ correction to drive the AD scalar amplitude to zero before the asymmetry forms. In addition, the initial expectation value of the AD field is not fixed by the dimension of the terms lifting the flat direction, allowing a wider range of flat direction and reheating temperature to generate the observed baryon asymmetry. The dynamics of a curvaton in a NTS-dominated Universe will be similarly modified, allowing a SUSY curvaton to have a large initial amplitude without quantum fluctua-
tions of the curvaton being damped by oscillations about the minimum of the curvaton potential.

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