Identifying the curvaton within MSSM

Rouzbeh Allahverdi\textsuperscript{1,2}, Kari Enqvist\textsuperscript{3}, Asko Jokinen\textsuperscript{4} and Anupam Mazumdar\textsuperscript{4}

\textsuperscript{1} Perimeter Institute for Theoretical Physics, Waterloo, ON, N2L 2Y5, Canada.
\textsuperscript{2} Department of Physics and Astronomy, McMaster University, Hamilton, ON, L8S 4M1, Canada.
\textsuperscript{3} Department of Physical Sciences, University of Helsinki, and Helsinki Institute of Physics, P.O. Box 64, FIN-00014 University of Helsinki, Finland
\textsuperscript{4} NORDITA, Blegdamsvej-17, Copenhagen-2100, Denmark.

E-mail: rallahverdi@perimeterinstitute.ca, kari.enqvist@helsinki.fi, ajokinen@nordita.dk, anupamm@nordita.dk

Abstract. We consider inflaton couplings to MSSM flat directions and the thermalization of the inflaton decay products, taking into account gauge symmetry breaking due to flat direction condensates. We then search for a suitable curvaton candidate among the flat directions, requiring an early thermally induced start for the flat direction oscillations to facilitate the necessary curvaton energy density dominance. We demonstrate that the supersymmetry breaking $A$-term is crucial for achieving a successful curvaton scenario. Among the many possible candidates, we identify the $u_1dd$ flat direction as a viable MSSM curvaton.
1. Introduction

Minimal Supersymmetric Standard Model (MSSM) provides nearly 300 gauge invariant flat directions, whose potentials are vanishing in a perfect supersymmetric (SUSY) limit (for a review, see [1]). However, in the early Universe SUSY is broken and the flatness of the potentials is spoiled, but the directions still remain flat as compared to the curvature scalar of the Universe. Soft SUSY breaking terms induce $m_0 \sim O(\text{TeV})$ masses to the flat direction. Flatness is also lifted by non-renormalizable superpotential terms, the form of which is dictated by the gauge properties of the flat direction. When the Hubble expansion rate becomes equal to the low energy SUSY scale, all the flat directions are trapped in their local minima, after which they start to oscillate and ultimately decay. The flat directions may play a crucial role in cosmological issues ranging from baryogenesis to dark generation matter. Recently it has also been pointed out that they are also important for understanding the full thermalization history of the Universe after inflation [2, 3].

An MSSM flat direction may also account for generating adiabatic density perturbations [4, 5, 6, 7] through the curvaton mechanism [8, 9, 10]. Since during inflation the Hubble expansion rate $H_I \gg m_0$ it does not cost anything in energy, quantum fluctuations are free to accumulate (in a coherent state) along a flat direction and form a condensate with a large VEV, $\varphi_0$. Because inflation smoothes out all the gradients, only the homogeneous condensate mode survives. However, the zero point fluctuations of the condensate impart a small, and in inflationary models a calculable, spectrum of perturbations on the condensate.

After inflation, $H \propto t^{-1}$ and the flat direction stays at a relatively large VEV due to large Hubble friction term until $H \simeq m_0$, whence the flat direction condensate starts oscillating around the origin with an initial amplitude $\sim \varphi_0$. From then on $|\varphi|$ is redshifted by the Hubble expansion $\propto H$ for matter dominated and $\sim H^{3/4}$ for radiation dominated Universe. The energy of the oscillating flat direction may eventually start to dominate over the inflaton decay products. When the flat direction decays, its isocurvature perturbations will be converted to the usual adiabatic perturbations of the decay products, which thus should ultimately contain also Standard Model (SM) degrees of freedom. However, such an evolution is not automatic but is subject to various constraints [4, 5], with the outcome that in general it is very difficult to have a MSSM flat direction curvaton that could give rise to the desired spectrum of density perturbations.

The key ingredient in these considerations is the process of thermalization of the ambient plasma. The state of the plasma before complete thermalization depends on the nature of flat direction which is developing a large VEV. For instance, MSSM Higgses developing a VEV give masses to the gauge bosons of $SU(2)_W \times U(1)_Y$, but gluons remain massless. The decay products can then reach partial thermalization through the gluons. Therefore one needs to study various possibilities case by case. It is important to know the details of the inflaton couplings to the MSSM degrees as well, since not all the couplings are renormalizable.

Thermal effects are important for the MSSM curvaton mechanism [11, 12] because large thermal corrections can evaporate or dissociate the MSSM curvaton candidate before it has a chance to dominate the energy density. These effects obviously depend on the thermalization rate. Recently it has been pointed out that the VEVs of the gauge invariant flat directions actually slow down the thermalization rate of the inflaton decay products [2, 3]. This is due to the fact that, in general, the flat direction
VEV gives masses to gauge bosons and gauginos which decrease the rates governed by $2 \leftrightarrow 2$ and $2 \rightarrow 3$ scatterings mediated via gauge bosons and gauginos. After the inflaton decay the initial plasma will be in a quasi-thermal phase for a long period during which kinetic equilibrium is established but not the full chemical equilibrium and the initial plasma is far from thermal equilibrium.

Another consequence of thermal corrections is that if the flat direction survives long enough, then thermally induced mass corrections can trigger early oscillations, which may hasten the dominance of the curvaton over the inflaton decayed products. These subtleties has never been systematically considered in the context of an MSSM curvaton. It is the purpose of the present paper to make use of the results of ref. and study these effects properly.

In previous work the evolution of the MSSM curvaton has been discussed without due consideration of the SUSY breaking $A$-term in the potential, which is induced by the presence of the non-renormalizable terms in the superpotential. Here we take the $A$-terms into account and show that they greatly modify the dynamical evolution and in particular the initial amplitude of the curvaton oscillations. As a consequence, many of the obstacles for an MSSM flat direction curvaton are now removed, and considering the various suitable directions case by case we are able to pick out the $u_d d_d$ flat direction as the most promising curvaton candidate.

The paper is organized as follows. In Section 2 we discuss how the inflaton couples to the MSSM degrees of freedom, emphasizing that there are certain combinations of slepton and squark superfields that do not have renormalizable couplings to a gauge singlet inflaton. In Section 3 we disseminate the potential flat direction curvaton candidates and choose those which give rise to thermal effects that can trigger early curvaton oscillations. Section 4 explains the importance of the $A$-term for MSSM curvaton dynamics, while in Section 5 we discuss the thermal corrections to flat directions, which singles out the $u_d d_d$ direction as the curvaton. In Section 6 we discuss the dynamics of the $u_d d_d$ curvaton and find its initial amplitude and a lower bound on the renormalizable inflaton coupling. Our conclusions are presented in Section 7.

2. Inflaton decay

In almost all $F$-term or $D$-term models of inflation the inflaton, $\Phi$, is considered to be an absolute gauge singlet. Then the pertinent question is how the inflaton couples to the matter.

Gauge symmetry implies that the inflaton must be coupled to a gauge-invariant combination of the SM fields and their SUSY partners. The field content of MSSM is governed by the following superpotential:

$$W_{MSSM} = \lambda_u Q H_u u + \lambda_d Q H_d d + \lambda_e L H_d e + \mu H_u H_d,$$

where $H_u, H_d, Q, L, u, d, e$ in Eq. (1) are chiral superfields representing the two Higgs fields (and their Higgsino partners), $L, (s)$quark doublets, $R, (s)$up- and down-type $(s)$quarks, $L, (s)$lepton doublets and $H, (s)$leptons respectively. The dimensionless

‡ Within SUSY the inflaton predominantly decays perturbatively. Non-perturbative decay of the inflaton to MS(SM) or through selfcoupling via parametric resonance is kinematically blocked because the inflaton decay products can couple (in a gauge invariant fashion) to MSSM flat directions. Large VEV of the flat direction induces large masses to the inflaton decay products which blocks preheating completely.

§ One exceptional example is ref. where the gauge invariant combinations of $SO(N)$ flat directions are responsible for driving inflation.
Yukawa couplings $\lambda_u, \lambda_d, \lambda_e$ are $3 \times 3$ matrices in the flavor space, and we have omitted the gauge and flavor indices. The last term is the $\mu$ term, which is a supersymmetric version of the SM Higgs boson mass.

There exist two gauge-invariant combinations of only two superfields:

$$H_u H_d, H_u L.$$  \hspace{2cm} (2)

The combinations which include three superfields are:

$$H_u Q u, H_d Q d, H_u L e, Q L d, u d d, L L e.$$  \hspace{2cm} (3)

SUSY together with gauge symmetry requires that the inflaton superfield $\Phi$ is coupled to these combinations or to their superpositions. The terms $\Phi H_u H_d$ and $\Phi H_u L$ have dimension four, and hence are renormalizable. On the other hand, the interaction terms that couple the inflaton to the combinations in Eq. (3) have dimension five and are non-renormalizable. They can arise after integrating out the heavy degrees of freedom which are associated to the new physics at high scales (for example GUT). Such terms are also generically induced by gravity, notably in supergravity, in which case they will be Planck mass suppressed $\parallel$.

The simplest case arises when the inflaton is coupled to matter via superpotential terms of the form:

$$h \Phi H_u H_d, h \Phi H_u \Psi,$$  \hspace{2cm} (4)

where we have defined

$$\Psi = \frac{1}{h} \sum_i h_i L_i, \quad h = \left[ \sum_i (h_i)^2 \right]^{1/2}.$$  \hspace{2cm} (5)

Here $1 \leq i \leq 3$ is a flavor index and $h$ can be as large as $O(1)$.

It is important to realize that one has to go from the $\{L_1, L_2, L_3\}$ basis, in which $\lambda_e$ in Eq. (1) are diagonalized, to the $\{\Psi, L'_1, L'_2\}$, where $L'_1$ and $L'_2$ are orthogonal to $\Psi$, and hence have no renormalizable couplings to the inflaton; only $\Psi$ couples to the inflaton. (for a discussion on multi-flat directions, see [16]). This is just a unitary transformation.

Besides the SM gauge group the MSSM Lagrangian is also invariant under R-parity ($(-1)^{3B+L+2S}$). Preserving R-parity (at least) at the renormalizable level further constrains inflaton couplings to matter. Note that $H_u H_d$ is assigned $+1$ under R-parity, while $H_u \Psi$ has the opposite assignment $-1$. Therefore only one of the couplings in Eq. (4) preserves R-parity: $\Phi H_u H_d$ if the inflaton superfield carries no lepton or baryon number, which is generically the case. In contrast, the $\Phi H_u \Psi$ term could be relevant for e.g. models where the RH sneutrino plays the role of the inflaton [17].

In the context of a curvaton scenario the fact that the inflaton can have renormalizable couplings to (some of) the MSSM fields is an auspicious sign. The interactions in Eq. (4) result in an inflaton decay rate:

$$\Gamma_d = \left( \frac{h^2}{4\pi} \right) m_\phi.$$  \hspace{2cm} (6)

On the other hand, non-renormalizable interactions from Eq. (3) result in a much smaller decay rate $\sim m_\phi^3 / M_P^2$. Such a late decay would imply a longer period of

$\parallel$ It is possible that the inflaton decays mainly to another singlet (for example, the right-handed neutrino) superfield. Then, since the ultimate goal is to create matter, this singlet must be coupled to MSSM fields. Supersymmetry and gauge symmetry again require couplings to gauge-invariant combinations in Eqs. (2) and (3) along the lines above.
Identifying the curvaton within MSSM

inflaton domination during which the ratio of the curvaton energy density to the total energy density of the Universe does not grow. Moreover, a small decay rate would result in a low reheat temperature, and hence would diminish the thermal effects necessary to trigger early oscillations of the curvaton field, as will be discussed below.

3. Curvaton candidates

Having discussed the inflaton couplings to matter, let us now attempt to identify the curvaton candidates within MSSM. An important point is that generating density perturbations of the correct size from curvaton decay requires that during inflation its VEV \( \varphi_0 \sim 10^5 H_1 \). Here \( H_1 \) is the Hubble expansion rate during the inflationary epoch.

A curvaton is a light modulus, lighter than the Hubble expansion rate during inflation. The total potential during inflation is given by

\[
V_{total} = V(I) + V(\varphi) \quad (7)
\]

where \( V(I) \) is due to inflation and \( V(\varphi) \) is due to the curvaton. The curvaton carries isocurvature perturbations which sources the curvature perturbations. In order not to have any residual isocurvature perturbations left over the curvaton must decay dominantly into the SM degrees of freedom. The interesting quantity to study is the ratio of the perturbation and the background field value of the curvaton, since this is related to the curvature perturbation \([8, 9]\), which can be constrained from the amplitude of the CMB anisotropy. If the perturbations in the curvaton do not damp then the final curvature perturbation will be given by

\[
\delta = \frac{H_{\text{inf}}}{\varphi_{\text{inf}}} \sim 10^{-5}, \quad (8)
\]

where \( 10^{-5} \) arises from the COBE normalization. In this we will present a scenario where the perturbations are not damped away and the curvaton dominates while decaying.

There are three important points to note:

- **The flat direction that plays the role of the curvaton cannot include those fields which have a renormalizable coupling to the inflaton.** The reason is that such fields generically acquire a mass \( h_\varphi \gg H_{\text{inf}} \) through their couplings to the inflaton, and hence cannot develop the required VEV. Therefore Eq. \([4]\) implies that a curvaton candidate cannot include \( H_u, H_d \) and \( \Psi \). As pointed out earlier, such a late decay would be harmful to the curvaton scenario.

- **A curvaton candidate should not induce a mass \( \geq m_\varphi/2 \) for the inflaton decay products, otherwise the two-body inflaton decay will be kinematically blocked.** The decay will be delayed until the relevant flat direction has started its oscillation and its VEV has been redshifted to sufficiently small values \([14]\).

- **The flat direction which plays the role of the curvaton should not break all of the SM gauge symmetry.** The reason is that in this case all gauge integrations will get decoupled and full (i.e. both kinetic and chemical) thermal equilibrium will be established much later than the inflaton decay \([3]\). However thermal effects which are necessary to trigger early oscillations of the flat direction \([11, 12]\).

\[\footnote{If only part of the SM group is broken, gauge interactions of the unbroken part will bring fields which carry those gauge quantum numbers, and subsequently other fields, into full equilibrium.}\]
crucially depend on the presence of a thermal bath consisting of some MSSM degrees of freedom in full equilibrium when the inflaton decay is completed.

Note that a flat direction, which is a combination of the squark and slepton fields, has couplings to other MSSM fields through which it induces a large mass for them. D-terms lead to flat direction couplings of gauge strength to the gauge fields and gauginos associated with the gauge (sub)group which is spontaneously broken by the flat direction VEV, as well as to the orthogonal directions and their supersymmetric partners. F-terms result in flat direction couplings of Yukawa strength to the scalars that are not included in the monomial representing the flat direction, as well as to their fermionic partners.

To elucidate these points, let us consider the particularly simple example of the $H_u H_d$ flat direction*. In this case the flat direction and the orthogonal directions are defined as $(H_{u,1} + H_{d,2}) / \sqrt{2}$, and $(H_{u,1} - H_{d,2}) / \sqrt{2}$, $H_{u,2}$ and $H_{d,1}$, respectively, where 1, 2 are the two components of $H_u$ and $H_d$. The flat direction breaks $SU(2)_W \times U(1)_Y$ down to $U(1)_{em}$. Then the gauge fields and gauginos associated with the broken subgroup, as well as $(H_{u,1} - H_{d,2}) / \sqrt{2}$, $H_{u,2}$, $H_{d,1}$ and their fermionic partner, acquire a mass $\sim g \langle \varphi \rangle$, where $g$ is a gauge coupling. The $Q_i$, multiplets, which are coupled to both $H_u$ and $H_d$, acquire a mass $\sqrt{(\lambda_{q,i}^2 + \lambda_{\tilde{d},i}^2)} \langle \varphi \rangle$ while $u$, $d$, and $e$, multiplets obtain respectively masses $\lambda_{u,i} \langle \varphi \rangle$, $\lambda_{d,i} \langle \varphi \rangle$, and $\lambda_{e,i} \langle \varphi \rangle$.

Let us denote the flat direction superfield by $\varphi$. Then we have

$$W_{\text{MSSM}} \supset \lambda_1 H_u \varphi \Sigma_1 + \lambda_2 H_d \varphi \Sigma_2 + \lambda_3 \Psi \varphi \Sigma_3,$$  \hspace{1cm} (9)

where $\Sigma_{1,2,3}$ are MSSM superfields. In general the relationship $m_\phi \leq H_{inf}$ holds*.

For the curvaton mechanism to work, we require $\varphi_{inf} \sim 10^5 H_{inf}$. Note that the VEV of the flat direction induces VEV dependent SUSY preserving masses to the MSSM particles. Therefore, for the inflaton decay to be kinematically allowed, one needs

$$\lambda_1, \lambda_2 \leq 10^{-5},$$ \hspace{1cm} (10)

if the $\Phi H_u H_d$ coupling is allowed by R-parity, and

$$\lambda_1, \lambda_3 \leq 10^{-5},$$ \hspace{1cm} (11)

if the $\Phi H_u \Psi$ coupling is allowed.

The above conditions considerably restrict the curvaton candidates within the MSSM as only the first generations of $(s)$quarks and $(s)$leptons which have a Yukawa coupling $\lesssim 10^{-5}$.

Let us now identify the flat directions which satisfy all of the above mentioned requirements. An important point is that the acceptable flat directions should include only one $Q$ or one $u$. The reason is that $D$- and $F$-flatness of directions which involve two or more $Q$ and/or $u$ requires them to be of different flavors (for details, see [13]). This implies the presence of up-type squarks from the second and/or third generation which, according to Eq. (9), will lead to $\lambda_1 \geq 10^{-3}$. This violates the condition for two-body inflaton decay via either of the $\Phi H_u H_u$ or $\Phi H_u \Psi$ terms, given in Eqs. (10)(11). A large number of MSSM flat directions will therefore be excluded by this consideration. The only flat directions with $\lambda_1 \leq 10^{-5}$ are as follows:

* Note that, as we mentioned, this flat direction cannot obtain a large VEV. We only consider this example to demonstrate how the flat direction VEV induces mass to other MSSM fields.

* This is strictly true in models of chaotic inflation and (supersymmetric) hybrid inflation. It is possible to have $m_\phi \gg H_{inf}$ in models of new inflation, but this is a rather contrived situation.
• **udd**: This monomial represents a subspace of complex dimension 6 \([18]\). \(D\)-flatness requires that the two \(d\) are from different generations (hence at least one of them will be from the second or third generation). This implies that \(\lambda_2 \geq 10^{-3}\), see Eq. (9), for all flat directions classified by \(udd\). As a consequence the two-body inflaton decay via the \(\Phi H_u H_d\) term will be kinematically forbidden. Note however that these flat directions are not coupled to \(L_{1,2,3}\). Therefore the inflaton decay via the \(\Phi H_u \Psi\) term can proceed for the \(u_1 dd\) directions, with \(u_1\) being the RH up squark. We also note that these directions leave the \(SU(2)_W\) unbroken, so that the \(SU(2)_W\) degrees of freedom can completely thermalize.

• **QLd**: This monomial represents a subspace of complex dimension 19 \([18]\). \(F\)-flatness requires that \(Q\) and \(d\) belong to different generations. Then, since \(Q\) and \(d\) are both coupled to \(H_d\), Eq. (9) implies that \(\lambda_2 \geq 10^{-3}\). Thus two-body inflaton decay via the \(\Phi H_u H_d\) term will be kinematically forbidden. On the other hand it can proceed via the \(\Phi H_u \Psi\) term for \(Q, L'dd\) directions, where \(Q_1\) is the doublet containing LH up and down squarks. Note that here we have to rotate to the \(\{\Psi, L'_1, L'_2\}\) basis where only \(L'_1\) and \(L'_2\) can acquire a large VEV during inflation. We also note that these directions completely break the \(SU(2)_W \times U(1)_Y\), but leave a \(SU(2)\) subgroup of the \(SU(3)_C\) unbroken. Therefore the associated color degrees of freedom can completely thermalize.

• **LLe**: This monomial represents a subspace of complex dimension three \([18]\). \(D\)-flatness requires that the two \(L\)s are from different generations, while \(F\)-flatness requires that \(e\) belongs to the third generations (therefore all the three lepton generations will be involved). Eq. (9) then implies that \(\lambda_2 \simeq 10^{-2}\), which kinematically blocks two-body inflaton decay via the \(\Phi H_u H_d\) term. The decay can nevertheless proceed via the \(\Phi H_u \Psi\) term. However, since \(\Psi\) cannot develop a large VEV, we should actually rotate to the \(\{\Psi, L'_1, L'_2\}\) basis instead and consider the \(L'_1 L'_2 e\) monomial. Such a change of basis has no impact on \(D\)-flatness but will affect the \(F\)-flatness. The reason is that in general \(\lambda_e\) is not diagonal in the \(\{\Psi, L'_1, L'_2\}\) basis, and hence \(L'_1\) and \(L'_2\) will be coupled to all flavors of \(e\). This can be circumvented if the inflaton dominantly couples to one of the \(L_i\) and \(\Psi\) is mainly \(L_i\). More specifically, a feasible curvaton candidate will be obtained along the \(L_2 L_3 e_1\) direction if \(\Psi \approx L_i\). For this flat direction we have \(\lambda_1 = 0\) and \(\lambda_3 \sim 10^{-5}\) (see Eq. (9)). This implies that two-body inflaton decay will proceed via the \(\Phi H_u L_1\) term without trouble. We note that this flat direction completely breaks the electroweak symmetry \(SU(2)_W \times U(1)_Y\), while not affecting \(SU(3)_C\). Therefore color degrees of freedom can fully thermalize. In passing we note that this monomial could be useful to generate primordial magnetic field, see \([19]\).

• **LLddd**: This monomial represents a subspace of complex dimension three \([18]\). \(D\)-flatness requires that the two \(L\)s and the three \(d\)s are all from different generations. This implies that that \(\lambda_2 \simeq 10^{-2}\) and \(\lambda_3 = \lambda_3 = 0\) (see Eq. (9)). Therefore two-body inflaton decay via the \(\Phi H_u H_d\) term will be kinematically forbidden, but the decay can proceed via the \(\Phi H_u \Psi\) term. Again note that we should rotate to the \(\{\Psi, L'_1, L'_2\}\) basis, and hence the \(L'_1 L'_2 ddd\) direction will be the relevant direction. However this direction breaks all of the SM gauge group. This results in late thermalization of the Universe \([3]\) and the absence of thermal effects which, as a consequence, does not yield early oscillations of
Identifying the curvaton within MSSM

the flat direction. This is an undesirable feature for our scenario in which, the curvaton oscillations are triggered by thermal effects, as we will discuss in the next Section.

To summarize, after taking all considerations into account, we find that the most promising candidates for the curvaton within MSSM are the $u_1dd$ and $Q_1L'd$ (and possibly $L_2L_3e_1$) flat directions.

4. Importance of the $A$-term

Let us now discuss the role of the $A$-term in constructing a curvaton model. In models with gravity and anomaly mediation, SUSY breaking results in the usual soft term, $m_0^2||\varphi||^2$, in the scalar potential with $m_0 \approx 100 \text{ GeV} - 1 \text{ TeV}$. There is also a new contribution arising from integrating out heavy modes beyond the scale $M$, which usually induces non-renormalizable superpotential terms of the form

$$W \sim \lambda_n \frac{\tilde{\varphi}^n}{nM^{n-3}},$$

where $\tilde{\varphi}$ denotes the superfield which comprises the flat direction $\varphi$. In general $M$ could be the string scale, below which we can trust the effective field theory, or $M = M_P$ (in the case of supergravity).

In addition, there are also inflaton-induced supergravity corrections to the flat direction potential. All these terms provide a general contribution to the flat direction potential which are of the form

$$V(\varphi) \sim H^2M_P^2f\left(\frac{\varphi}{M_P}\right), \quad V(\varphi) \sim HM_P^3f\left(\frac{\varphi^n}{M_P^{n-3}}\right),$$

where $f$ is some function. The first contribution usually gives rise to a Hubble induced correction to the mass of the flat direction, $c_H H^2|\varphi|^2$, with an unknown coefficient, $c_H$, which depends on the nature of the Kähler potential. The second contribution is the Hubble-induced $A$-term.

Note that $c_H$ can have either sign. If $c_H \sim 1$, the flat direction mass is $> H$. It therefore settles at the origin during inflation and remains there. Since $|\varphi| = 0$ at all times, the flat direction will have no interesting consequences in this case. The positive Hubble induced mass to the flat direction has a common origin to the Hubble induced mass correction to the inflaton in supergravity models. This is the well known $\eta$-problem, which arises because of the canonical form of the inflaton part of the Kähler potential. A large $\eta$ generically spoils slow roll inflation. In order to have a successful slow roll inflation, one needs $\eta \ll 1$.

The Hubble induced terms can be eliminated completely from inflation and MSSM flat direction sectors completely if there is a shift symmetry in the Kähler potential, or, the Heisenberg symmetry. Note that string theory generically gives rise to No-scale type Kähler potential which preserves Heisenberg symmetry at tree level. In

A nice realization of chaotic inflation within supergravity is obtained by implementing a shift symmetry. If the inflaton Kähler potential has the form $K = (\phi + \phi^*)^2 / M_P^2$, instead of the minimal one $K = \phi^*\phi / M_P^2$, the scalar potential along the imaginary part of $\phi$ remains flat even for Transplanckian field values. Therefore it can play the role of the inflaton in a chaotic model. Note that a shift symmetry also ensures that the (positive) Hubble induced corrections to the mass of flat directions vanishes at the tree-level and the terms in Eq. (13) disappear. It is also possible to realize chaotic inflation for sub Planckian field values in supergravity. For example, see the multi-axions driven assisted inflation.
either cases the inflaton and the MSSM sectors are free from Hubble-induced mass and $A$-terms. In the following we will consider examples where there are no Hubble-induced terms in the potential.

The relevant part of the scalar potential is then given by

\[
V = m_0^2|\phi|^2 + \lambda_n^2 \frac{|\phi|^{2(n-1)}}{M_P^{2(n-3)}} + \left( A\lambda_n \frac{\phi^n}{M_P^{n-3}} + \text{h.c.} \right),
\]

(14)

where $\lambda_n \sim \mathcal{O}(1)$ and $n \geq 4$. Note that the low energy $A$-term is a dimensionful quantity, i.e., $A \sim m_0 \sim \mathcal{O}(100 \text{ GeV} - 1 \text{ TeV})$, and depends on a phase. As shown in [18], all of the MSSM flat directions are lifted by higher-order terms with $n \leq 9$. If a flat direction is lifted at the superpotential level $n$, the VEV that it acquires during inflation will also depend on the presence of the $A$-term.

However the angular direction of the potential being flat, during inflation it obtains random fluctuations. There will be equally populated domains of Hubble patch size where the phase of the $A$-term is positive and negative. In either case during inflation, the flat direction VEV is given by:

\[
\phi_{\text{inf}} \sim \left( m_0 M_P^{n-3} \right)^{1/n-2}.
\]

(15)

However there is a distinction between a positive and a negative phase of the $A$-terms. The difference in dynamics arises after the end of inflation. In the case of positive $A$-term the flat direction starts rolling immediately, but in the case of a negative $A$-term, the flat direction remains in a false vacuum for which the VEV is given by Eq. (15).

The mass of the flat direction around this false minimum is very small compared to the Hubble expansion rate during inflation, i.e. $(3n^2 - 9n + 8)m_0^2 \ll H_{\text{inf}}^2$ where $n > 3$. During inflation the flat direction obtains quantum fluctuations whose amplitude is given by Eq. (5).

The flat direction can exit such a metastable minimum only if thermal corrections are taken into account. From now onwards we will assume that the flat direction is locked in a false vacuum with a VEV $\phi_{\text{inf}}$ during and right after inflation. In the next Section we will discuss the various thermal corrections a flat direction obtains.

5. Thermal corrections to the flat direction

If part of the SM gauge group remains unbroken by the flat direction VEV, the associated gauge fields and gauginos will not receive an induced mass. Together with the light particles and sparticles, they reach full thermal equilibrium through gauge interactions (for related studies, see refs. [27, 28]). The most important processes are $2 \to 2$ and $2 \to 3$ scatterings with gauge boson exchange in the $t$-channel [27]. For massless gauge bosons (as happens for the unbroken subgroup) these scatterings are extremely efficient and lead to an almost instant thermalization of particles upon their production in inflaton decay [28] ††.

The presence of a thermal bath with a temperature $T$ affects the flat direction dynamics. This happens through the back-reaction of fields which are coupled to the flat direction [11, 12]. The flat direction VEV naturally induces a mass $y\phi_{\text{inf}}$ to the field which are coupled to it, where $y$ is a gauge or Yukawa coupling. Depending on whether $y\phi_{\text{inf}} \leq T$ or $y\phi_{\text{inf}} > T$, different situations arise.

††However particles carrying quantum numbers associated with the broken gauge subgroup reach full equilibrium much later. The reason is that corresponding gauge fields (and gauginos) acquire a mass from the flat direction VEV which suppresses the thermalization rate of these particles [2, 3].
• \( y\varphi_{\text{inf}} \leq T \): Fields which have a mass smaller than temperature are kinematically accessible to the thermal bath. They will reach full equilibrium and result in a thermal correction \( V_{\text{th}} \) to the flat direction potential

\[
V_{\text{th}} \sim \mp y^2 T^2 |\varphi|^2.
\]  

(16)

The flat direction then starts oscillating, provided that \( yT > H \) \[11\]. For \( H \leq \Gamma_d \) the Universe is in a radiation-dominated phase, and hence \( T \propto H^{3/2} \). For \( H > \Gamma_d \) the inflaton has not completely decayed and the plasma from partial inflaton decay has a temperature \( T \sim (\Gamma_d M_{\text{Pl}}^2)^{1/4} \) \[29\], which implies that \( T \propto H^{1/4} \). Therefore in both cases \( yT > H \) will remain valid once oscillations start.

• \( y\varphi_{\text{inf}} > T \): Fields which have a mass larger than temperature will not be in equilibrium with the thermal bath. For this reason they are also decoupled from the running of gauge couplings (at finite temperature). This shows up as a correction to the free energy of gauge fields, which is equivalent to a logarithmic correction to the flat direction potential \[12\]

\[
V_{\text{th}} \sim \pm \alpha T^4 \ln \left( |\varphi|^2 \right),
\]  

(17)

where \( \alpha \) is a gauge fine structure constant. Decoupling of gauge fields (and gauginos) results in a positive correction, while decoupling of matter fields (and their superpartners) results in a negative sign. The overall sign then depends on the relative contribution of decoupled fields \( \dagger \). Obviously only corrections with a positive sign can lead to flat direction oscillations around the origin (as we require). Oscillations begin when the second derivative of the potential exceeds the Hubble rate-squared which, from Eq. (17), reads \( (\alpha T^2/\varphi_{\text{inf}}) > H \).

Note that thermal effects of the first type require that fields which have Yukawa couplings to the flat direction are in full equilibrium, while those of the second type require the gauge fields (and gauginos) be in full equilibrium. If all of the SM group is broken by flat direction(s), the reheated plasma will reach full equilibrium much later after the inflaton decay \[3\]. Instead it enters a long phase of quasi-adiabatic evolution during which the plasma remains dilute. This implies that fields which are coupled to the flat direction, and hence have a large mass, decay quickly and are not produced again as the inverse decays are inefficient. As a consequence, thermal effects, which crucially depend on the presence of these degrees of freedom, will be weakened. This is the reason why we want a subgroup of the SM gauge symmetry to remain unbroken.

Let us now examine the thermal effects for the curvaton candidates identified in Section III. Note that the instantaneous temperature of the thermal bath from (partial) inflaton decay is \( T \leq m_\varphi \) for a perturbative inflaton decay. Therefore the kinematical condition for inflaton decay via the \( \Phi H_u H_d \) and \( \Phi H_u \Psi \) terms, given in Eqs. (10, 11) respectively, implies that only fields with a Yukawa coupling \( \lambda \leq 10^{-5} \) to the flat direction can be in equilibrium with the thermal bath. We only have two possibilities:

• \( u_1 d d \): \( SU(2)_W \) remains unbroken in this case. This implies that the corresponding gauge fields and gauginos, \( H_u \) and \( H_d \) (plus the Higgsinos) and the LH (s)leptons reach full thermal equilibrium. The back-reaction of \( H_u \) results in a

\( \dagger \) For example, consider the \( H_u H_d \) flat direction. This direction induces large masses for the top (s)quarks which decouples them from the thermal bath, while not affecting gluons and gluinos. Therefore this leads to a positive contribution from the free energy of the gluons.
Identifying the curvaton within MSSM

thermal correction $\lambda^2 T^2 |\varphi|^2$, see Eq. (9), with $\lambda_1 \sim 10^{-5}$. The free energy of the $SU(2)_W$ gauge fields result in a thermal correction $+\alpha_W T^4 \ln \left( |\varphi|^2 \right)$. Note that the sign is positive since the flat direction induces a mass which is $> T$ (through the $d$) for the LH (s)quarks but not the $SU(2)_W$ gauge fields and gauginos.

- $Q_1 L'd$: An $SU(2)$ subgroup of $SU(3)_C$ is unbroken in this case. Therefore only the corresponding gauge fields and gauginos plus some of the (s)quark fields reach full thermal equilibrium. Then the back-reaction of $u_1$ and $d_1$ results in a thermal correction $(\lambda_1^2 + \lambda_2^2) T^2 |\varphi|^2$ according to Eq. (9), where $\lambda_1 \sim \lambda_2 \sim 10^{-5}$. Note that logarithmic thermal corrections will not be useful in this case as decoupling of a number of gluons (and gluinos) from the running of strong gauge coupling results in a negative contribution to the free energy of the unbroken part of $SU(3)_C$.

Therefore, within MSSM and from the point of view of thermal effects, the $u_1 dd$ flat direction is the most suitable curvaton candidate.

6. $u_1 dd$ as the MSSM curvaton

The flat directions represented by the $udd$ monomial are partly lifted by $n = 4$ superpotential terms, and eventually lifted at the $n = 6$ level. The former are lifted by the $udde$ superpotential term which implies a zero $A$-term because $udd$ achieves a large VEV which renders a large mass for $e$ and consequently $\langle e \rangle = 0$. Therefore the scalar potential along these directions will have no metastable minimum. The latter, on the other hand, are lifted by the $(udd)^2$ superpotential term which also induces a non-zero $A$-term (as required). As we shall see, for these directions sufficient thermal corrections are induced to lift the flat direction from its false minimum.

The only channel for two-body inflaton decay is through the $\Phi H_u \Psi$ term. The existence of two generations of $d$ in the flat direction kinematically blocks two-body decay via the other renormalizable term $\Phi H_u H_d$ (see the discussion in Section 3). Note that the $SU(2)_W$ gauge fields and gauginos are massless as this subgroup remains unbroken by the $u_1 dd$ direction. Therefore degrees of freedom which carry $SU(2)_W$ gauge quantum numbers, and their induced mass by the curvaton is $\lesssim T$, fully thermalize instantly (compared with the Hubble expansion rate) via $2 \rightarrow 2$ and $2 \rightarrow 3$ scatterings with $SU(2)_W$ gauge bosons in the t-channel. On the other hand, the $SU(3)_C \times U(1)_Y$ gauge fields and gauginos, plus the right-handed (s)quarks and (s)leptons, reach full equilibrium much later as their gauge interactions are suppressed by the flat direction induced mass. In the following $T$ refers to the temperature of $SU(2)_W$ degrees of freedom which are in full thermal equilibrium.

From the discussion in the previous Section it follows that thermal effects result in a correction to the effective thermal flat direction potential given by

$$V_{\text{th}} \sim \lambda_1^2 T^2 |\varphi|^2 + \alpha_W T^4 \ln \left( |\varphi|^2 \right),$$

where $\lambda_1 \sim 10^{-5}$ and $\alpha_W \sim 10^{-2}$.

† As pointed earlier, the $L_L \bar{L}_L \bar{e}_1$ flat direction can obtain a large VEV in the specific case where $\Psi = L_1$. The $SU(3)_C$ part of the SM gauge symmetry remains unbroken for this flat direction, and hence gluons, gluinos and (s)quarks will reach full equilibrium. We note that neither of $L_1$ and $e$ are coupled to the color degrees of freedom. This implies that there will be no $T^2 \varphi^2$ or logarithmic correction to the flat direction potential in this case. This excludes the $L_L \bar{L}_L \bar{e}_1$ flat direction from being a successful curvaton candidate.
Identifying the curvaton within MSSM

Note that the first and foremost condition for lifting the flat direction from its metastable minimum is that $V_{\text{th}} > m_0 \phi_{\text{inf}}^2$. On the other hand, a perturbative inflaton decay yields a radiation-dominated Universe whose temperature is $T \leq T_R \leq m_\phi$. The reheat temperature $T_R \sim \sqrt{\Gamma_d M_P}$ is the largest temperature of the Universe in a quasi-thermal radiation-dominated phase, where $\Gamma_d$ is the inflaton decay rate given in Eq. (6).

If $T_R \sim m_\phi$, the first term on the right-hand side of Eq. (18) dominates over the flat direction energy density given by Eq. (14) for $n = 6$ right after the inflaton has completely decayed. Note that $\phi_{\text{inf}} \sim 10^9 m_\phi$ from Eq. (8). Note however that very soon we will have $T < \lambda_1 \phi_{\text{inf}}$ since the Hubble expansion rate is gradually decreasing and the thermal contribution, i.e. $\lambda^2 T^2 |\phi|^2$, becomes ineffective, since the fields with coupling $\lambda_1$ to the $u_1d_d$ direction have a mass $> T$ and hence drop out of quasi-thermal equilibrium. Eventually the logarithmic thermal correction would take over very quickly, even if $T \approx m_\phi$. For this reason we focus on the second term of Eq. (18) in the discussion below.

First of all we need $V_{\text{th}} > m_0 \phi_{\text{inf}}^2$, so that the thermal effects will overcome the potential barrier. This leads to the requirement

$$\alpha_W T^4 > m_0^2 \phi_{\text{inf}}^2.$$  \hspace{1cm} (19)

Second, the thermal mass should trigger flat direction oscillations. If oscillations do not begin, the flat direction simply sits at a field value $\phi \sim \phi_{\text{inf}}$ while $V_{\text{th}}$ is redshifted $\propto H^2$ in a radiation-dominated Universe. The flat direction will be trapped again in the metastable minimum when $V_{\text{th}} < m_0^2 \phi_{\text{inf}}^2$.

In order for the flat direction oscillations to start we must have $d^2 V_{\text{th}}/d|\varphi|^2 > H^2$. This leads to the condition

$$\alpha_{\text{W}}^{1/2} \frac{T^2}{\varphi_{\text{inf}}} > H(T),$$  \hspace{1cm} (20)

which always holds in a radiation-dominated phase where $H \simeq T^2/M_P$ (note that $\phi_{\text{inf}} \ll M_P$). This implies that the $u_1d_d$ direction starts oscillating once the condition given in Eq. (19) is satisfied. This happens, when temperature of the Universe is given by

$$T_{\text{osc}} \sim \left( \frac{\phi_{\text{inf}}}{10^{14} \text{ GeV}} \right)^{1/2} \times 10^9 \text{ GeV}.$$  \hspace{1cm} (21)

After taking into account of the constraints $T_{\text{osc}} \leq m_\phi$, which holds for a perturbative inflaton decay, and $m_\phi \sim 10^{-5} \phi_{\text{inf}}$, which arises from the amplitude of the CMB anisotropy, we obtain

$$\phi_{\text{inf}} \geq 10^{14} \text{ GeV}.$$  \hspace{1cm} (22)

Interestingly enough for directions which are lifted at the $n=6$ superpotential level we have $\phi_{\text{inf}} \sim 3 \times 10^{14} \text{ GeV}$, see Eq. (17). Then the inflaton mass is determined from Eq. (8) to be $m_\phi \sim 3 \times 10^9 \text{ GeV}$. Also the requirement that $T_R \geq T_{\text{osc}}$ (a radiation-dominated Universe is formed only after the inflaton has completely decayed) restricts the inflaton decay rate according to Eqs. (6), (24). A lower bound on the renormalizable inflaton coupling is $h \geq 10^{-4}$ from Eqs. (4), (6), which is perfect for the validity of a perturbative inflaton decay [3]. Note that the flat direction VEV prevents non-perturbative inflaton decay as described in Ref. [14].

§ Obviously the $T^2 |\phi|^2$ correction will not arise at all if $T_R \ll m_\phi$, since in this case we will have $\lambda_1 \phi_{\text{inf}} \gg T$. 


Last but an important reminder to the readers is that the evaporation of the flat direction is not an important issue in our case. The interaction term that leads to evaporation is of the form $\lambda_1^2 \varphi^2 \chi^2$, with $\varphi$ being the flat direction and, $\chi$ collectively denotes the fields in thermal equilibrium. The rate for evaporation is then given by:

$$\Gamma \simeq \lambda_1^4 \frac{n_\chi}{E_\chi E_\varphi},$$

(23)

where $n_\chi \sim T^3$, $E_\chi \sim T$ and $E_\varphi \sim T^2/\langle \varphi \rangle$. Note that the latter being the mass of flat direction due to logarithmic thermal corrections, where $\langle \varphi \rangle$ is the the amplitude of the flat direction oscillations. We then find:

$$\Gamma \sim \lambda_1^4 \langle \varphi \rangle \ll H(T).$$

(24)

since $\lambda_1 \sim 10^{-5}$ for $u_1 dd$ flat direction. It takes a large number of oscillations for the flat directions to evaporate. Further note that during radiation domination, the flat direction VEV scales like: $\langle \varphi \rangle \propto H^{3/4}$, and $\langle \varphi \rangle \propto H$ during matter domination and/or when the flat flat direction oscillates. Therefore, thermal evaporation is not an important threat before the curvaton oscillations dominate the Universe ||.

Hence we conclude that the $u_1 dd$ flat direction is a realistic curvaton candidate in the MSSM. At the time of its decay, it dominates the energy density of the universe so that its isocurvature perturbations are converted into the observed curvature perturbations in CMB. During the last stages of the $u_1 dd$ oscillations, the flat direction generically fragments into $Q$-balls [30]. The $Q$-balls eventually decay into LSPs through surface evaporation, which can lead to the observed dark matter. Since the $Q$-balls behave like non-relativistic matter, they again start dominating the radiation energy density until they completely evaporates. However, this does not modify the super-Hubble perturbations. The fragmentation of the flat direction into $Q$-balls is strictly a sub-Hubble process.

7. Conclusion

We have considered thermal corrections in the presence of a MSSM flat direction condensate, which gives rise to gauge symmetry breaking and a slowing down of thermalization rates. We have paid particular attention to the coupling of the inflaton to ordinary SM matter and identified among the nearly 300 MSSM flat directions the $u_1 dd$ direction lifted by $n = 6$ non-renormalizable terms as a successful MSSM curvaton candidate that survives the constraints of energy density dominance and the condensate non-dissociation by the ambient plasma. To our knowledge this is the first paper where thermal corrections to the MSSM flat direction curvature are accounted for properly.

The $u_1 dd$ direction provides masses for the gauge bosons of $SU(3)_C \times U(1)_Y$ while leaving $SU(2)_W$ unbroken. It is thus only the $SU(2)_W$ degrees of freedom that contribute to the thermal correction of the flat direction potential, as given in Eq. (18). The dominance of $u_1 dd$ over thermal plasma comes about because of two important factors. First, the condensate has an initial amplitude which is large, i.e. $\varphi_{inf} \sim 3 \times 10^{14}$ GeV. Second, there is an $A$-term which creates a false vacuum during and after inflation which traps the flat direction VEV for a sufficiently long time.

The CMB fluctuations imparted by $u_1 dd$ not only have cosmological implications but also astrophysical ones. The flat direction is also a well motivated candidate ||. Along the same lines, it is also true that in the case of a quadratic thermal correction, the evaporation rate is again subdominant compared to the Hubble expansion rate.
for generating the cold dark matter through $Q$-ball evaporation. Future collider-based and astrophysical experiments will hopefully pin down the nature of dark matter and the physics beyond the SM, but for the time being we are assured that minimal supersymmetric Standard Model can provide a cosmologically viable curvaton candidate, $udd$ responsible for the CMB fluctuations and reheating, and possibly also accounting for the dark matter.

As in all curvaton scenarios, the spectral index $n_s$ for the $udd$ curvaton is very close to 1 whereas WMAP 3-year data indicates $n_s = 0.951^{+0.015}_{-0.019}$. However, in case of the $udd$ curvaton the spectral index of the power spectrum depends also on the yet unspecified inflaton sector with possible tensor perturbations.

Acknowledgments

The work of R.A. is supported by the Natural Sciences and Engineering Research Council of Canada (NSERC). K.E. is supported in part by the Academy of Finland grant no. 205800. A.M. would like to thank CERN, University of Padova, Perimeter Institute and McGill University for their kind hospitality where parts of this project were carried out.

References

[1] K. Enqvist and A. Mazumdar, Phys. Rept. 380, 99 (2003) [arXiv:hep-ph/0209214]. M. Dine and A. Kusenko, Rev. Mod. Phys. 76, 1 (2004) [arXiv:hep-ph/0303065].
[2] R. Allahverdi and A. Mazumdar, arXiv:hep-ph/0505059.
[3] R. Allahverdi and A. Mazumdar, arXiv:hep-ph/0512227.
[4] K. Enqvist, S. Kasuya and A. Mazumdar, Phys. Rev. Lett. 90, 091302 (2003) [arXiv:hep-ph/0211147].
[5] K. Enqvist, A. Jokinen, S. Kasuya and A. Mazumdar, Phys. Rev. D 68, 103507 (2003) [arXiv:hep-ph/0303165].
[6] K. Enqvist, S. Kasuya and A. Mazumdar, Phys. Rev. Lett. 93, 061301 (2004) [arXiv:hep-ph/0311221]. K. Enqvist, A. Mazumdar and A. Perez-Lorenzana, Phys. Rev. D 70, 103508 (2004) [arXiv:hep-th/0403041]. K. Enqvist, A. Mazumdar and M. Postma, Phys. Rev. D 67, 123503 (2003) [arXiv:astro-ph/0301127].
[7] M. Postma, Phys. Rev. D 67, 063518 (2003) [arXiv:hep-ph/0212005]. S. Kasuya, M. Kawasaki and F. Takahashi, Phys. Lett. B 578, 259 (2004) [arXiv:hep-ph/0305134]. A. Mazumdar and M. Postma, Phys. Lett. B 573, 5 (2003) [Erratum-ibid. B 585, 295 (2004)] [arXiv:astro-ph/0306509]. M. Tegami and T. Moroi, Phys. Rev. D 70, 083515 (2004) [arXiv:hep-ph/0404253]. A. Mazumdar and A. Perez-Lorenzana, Phys. Rev. Lett. 92, 251301 (2004) [arXiv:hep-ph/0306020]. A. Mazumdar and A. Perez-Lorenzana, Phys. Rev. D 70, 083526 (2004) [arXiv:hep-ph/0406154].
[8] K. Enqvist and M. S. Sloth, Nucl. Phys. B 626, 395 (2002) [arXiv:hep-ph/0109214]. A. D. Linde and V. F. Mukhanov, Phys. Rev. D 56, 535 (1997) [arXiv:astro-ph/9610219].
[9] D. H. Lyth and D. Wands, Phys. Lett. B 524, 5 (2002) [arXiv:hep-ph/0110002]. D. H. Lyth, C. Ungarelli and D. Wands, Phys. Rev. D 67, 023503 (2003) [arXiv:astro-ph/0208055].
[10] T. Moroi and T. Takahashi, Phys. Lett. B 522, 215 (2001) [Erratum-ibid. B 539, 303 (2002)] [arXiv:hep-ph/0110096].
[11] R. Allahverdi, B. A. Campbell and J. R. Ellis, Nucl. Phys. B 579, 355 (2000) [arXiv:hep-ph/0001122].
[12] A. Anisimov and M. Dine, Nucl. Phys. B 619, 729 (2001) [arXiv:hep-ph/0008058]. A. Anisimov, Phys. Atom. Nucl. 67, 640 (2004) [arXiv:hep-ph/0111233].
Identifying the curvaton within MSSM

[13] K. Enqvist, S. Kasuya and A. Mazumdar, Phys. Rev. Lett. 89, 091301 (2002) [arXiv:hep-ph/0204270]; K. Enqvist, S. Kasuya and A. Mazumdar, Phys. Rev. D 66, 043505 (2002) [arXiv:hep-ph/0206272]; R. Allahverdi, R. Brandenberger and A. Mazumdar, Phys. Rev. D 70, 083535 (2004) [arXiv:hep-ph/0407230].

[14] R. Allahverdi and A. Mazumdar, “Towards a successful reheating within supersymmetry”, arXiv:hep-ph/0603244.

[15] A. Jokinen and A. Mazumdar, Phys. Lett. B 597, 222 (2004) [arXiv:hep-th/0406074].

[16] K. Enqvist, A. Jokinen and A. Mazumdar, JCAP 0401, 008 (2004) [arXiv:hep-ph/0311336].

[17] H. Murayama, H. Suzuki, T. Yanagida and J. Yokoyama, Phys. Rev. Lett. 70, 1912 (1993); H. Murayama, H. Suzuki, T. Yanagida and J. Yokoyama, Phys. Rev. D 50, 2356 (1994) [arXiv:hep-ph/9311326].

[18] T. Gherghetta, C. Kolda and S. P. Martin, Nucl. Phys. B 468, 37 (1996) [arXiv:hep-ph/9510370].

[19] K. Enqvist, A. Jokinen and A. Mazumdar, JCAP 0411, 001 (2004) [arXiv:hep-ph/0404269].

[20] M. Dine, L. Randall and S. Thomas, Phys. Rev. D 70, 083535 (2004) [arXiv:hep-ph/0407230].

[21] M. Dine, W. Fischler, and D. Nemeschansky, Phys. Lett. B 136, 169 (1984); G. D. Coughlan, R. Holman, P. Ramond, and G. G. Ross, Phys. Lett. B 140, 44 (1984); A. S. Goncharov, A. D. Linde, and M. I. Vysotsky, Phys. Lett. B 147, 279 (1984); O. Bertolami, and G. G. Ross, Phys. Lett. B 183, 163 (1987); E. J. Copeland, A. R. Liddle, D. H. Lyth, E. D. Stewart, and D. Wands, Phys. Rev. D 49, 6410 (1994) [arXiv:astro-ph/9401011].

[22] M. Kawasaki, M. Yamaguchi and T. Yanagida, Phys. Rev. Lett. 85, 3572 (2000) [arXiv:hep-ph/0004243].

[23] S. Dimopoulos, S. Kachru, J. McGreevy and J. G. Wacker, arXiv:hep-th/0507205.

[24] A. R. Liddle, A. Mazumdar and F. E. Schunck, Phys. Rev. D 58, 061301 (1998) [arXiv:astro-ph/9804177]; E. J. Copeland, A. Mazumdar and N. J. Nunes, Phys. Rev. D 60, 083506 (1999) [arXiv:astro-ph/9904309]; A. Mazumdar, S. Panda and A. Perez-Lorenzana, Nucl. Phys. B 614, 101 (2001) [arXiv:hep-ph/0107058].

[25] M. K. Gaillard, H. Murayama and K. A. Olive, Phys. Lett. B 355, 71 (1995) [arXiv:hep-ph/9504307].

[26] A. B. Lahanas and D. V. Nanopoulos, Phys. Rept. 145, 1 (1987).

[27] J. McDonald, Phys. Rev. D 61, 083513 (2000) [arXiv:hep-ph/9909467]; R. Allahverdi, Phys. Rev. D 62, 063509 (2000) [arXiv:hep-ph/0004035].

[28] S. Davidson and S. Sarkar, JHEP 0011, 012 (2000) [arXiv:hep-ph/0009078]; R. Allahverdi and M. Drees, Phys. Rev. D 66, 063513 (2002) [arXiv:hep-ph/0205234]; P. Jaikumar and A. Mazumdar, Nucl. Phys. B 683, 264 (2004) [arXiv:hep-ph/0212269].

[29] E. W. Kolb and M. S. Turner, *The Early Universe*, Addison-Wesley, New York 1990.

[30] A. Kusenko, Phys. Lett. B 405, 108 (1997) [arXiv:hep-ph/9704273]; A. Kusenko, Phys. Lett. B 404, 285 (1997) [arXiv:hep-th/9704073]; A. Kusenko and M. E. Shaposhnikov, Phys. Lett. B 418, 46 (1998) [arXiv:hep-ph/9709492]; K. Enqvist and J. McDonald, Phys. Lett. B 425, 309 (1998) [arXiv:hep-ph/9711514].

[31] D. N. Spergel, et al., astro-ph/0603449.