PHYSICAL REVIEW D **83,** 063527 (2011)

Inflation in string theory: A graceful exit to the real world

Michele Cicoli[1] and Anupam Mazumdar[2,3]
[1]*Deutsches Elektronen-Synchrotron DESY, Notkestrasse 85, 22603 Hamburg, Germany*
[2]*Physics Department, Lancaster University, Lancaster, LA1 4YB, United Kingdom*
[3]*Niels Bohr Institute, Copenhagen University, Blegdamsvej-17, DK-2100, Denmark*
(Received 4 November 2010; published 31 March 2011)

The most important criteria for a successful inflation are: explaining the observed temperature anisotropy in the cosmic microwave background radiation, and exiting inflation in a vacuum where it can excite the standard model quarks and leptons required for the success of big bang nucleosynthesis. In this paper, we provide the first ever closed-string model of inflation where the inflaton couplings to hidden sector, moduli sector, and visible sector fields can be computed, showing that inflation can lead to reheating the standard model degrees of freedom before the electro-weak scale.

DOI: 10.1103/PhysRevD.83.063527 PACS numbers: 98.80.Cq, 11.25.Wx, 12.60.Jv, 98.80.Bp

I. INTRODUCTION

The inflationary paradigm has played a key role in explaining the large scale structure of the Universe and the temperature anisotropy in the cosmic microwave background (CMB) radiation [1]. It is believed that a scalar field, the inflaton, drives the inflationary dynamics. During inflation the quantum fluctuations of the inflaton seed the initial perturbations for the structure formation, and after inflation its coherent oscillations lead to reheat the Universe with the observed light degrees of freedom of the standard model quarks and leptons (for a review, see [2]). These degrees of freedom need also to thermalize before big bang nucleosynthesis (BBN) [3].

In this regard, visible sector models of inflation where inflation ends in the standard model *gauge invariant* vacuum are preferred with respect to the ones where the inflaton belongs to the hidden sector and couples to all particles in the model. In the former case, there are very few models based on low scale supersymmetry (SUSY) as in the minimal supersymmetric standard model (MSSM), where there exists *only* 2 inflaton candidates which carry standard model charges [4], and their decay populates the Universe with MSSM particles and dark matter [5].

On the contrary, there are many models of inflation where the inflaton is a standard model gauge singlet [2]. Promising examples are closed-string inflationary models in type IIB flux compactifications [6,7]. The inflaton is a Kähler modulus parameterizing the size of an internal cycle, and inflation can be embedded within a uv complete theory where the η-problem can be solved due to the no-scale structure of the potential and the fact that the volume mode of the internal manifold is kept stable during inflation. Moreover, the order of magnitude of all the inflaton couplings can be computed [8,9].

Despite all these successes, it is *still* hard to explain how the inflaton energy gets transferred primarily to visible degrees of freedom, and not to hidden ones, since *a priori a gauge singlet inflaton has no preference to either the visible or the hidden sector.*

Thus, the main challenge to build any hidden sector model of inflation is to reheat the visible sector so that the thermal bath prior to BBN is filled with the standard model hadrons.

In order to fulfill this, the conditions are:

(i) The inflaton must decay primarily into the visible sector, and its coupling to the hidden sector must be weak enough to prevent an overproduction of its degrees of freedom nonthermally or nonperturbatively, as in the case of preheating (for a review see [10]).

(ii) The hidden sector must not contain very light species since their presence at the time of BBN could modify the light-element abundance. The current LEP and BBN constraints on extra light species is very tight, i.e. ≤ 4 [3].

(iii) The visible sector must be sequestered from the hidden sector so that the late decay of the hidden degrees of freedom does not spoil the success of BBN or overpopulate the dark matter abundance.

We stress that hidden sectors arise naturally in string compactifications since they come along with many internal cycles which have to be stabilized. This is generically achieved by wrapping stacks of Dp-branes around $(p-3)$-cycles in order to generate perturbative and nonperturbative effects that lead to moduli fixing. The presence of such Dp-branes is also in general needed to achieve tadpole cancellation, which ensures the absence of anomalies in the $4D$ effective field theory. Given that each stack of D-branes supports a different field theory, the presence of hidden sectors turns out to be very generic in any such constructions.

In order to address all these issues related to the hidden sector, it is crucial to know the inflaton couplings to all hidden and visible degrees of freedom. In this paper we shall present a type IIB closed-string inflation model where all the inflaton couplings can be derived within a uv complete theory [8]. The knowledge of these couplings allows us to study the reheating of the visible degrees of freedom, which can be achieved with a temperature above the BBN temperature (for other studies on (p)reheating in string cosmology, see [11,12]).

The beauty of this *top-down* setup is also that the flatness of the inflaton potential can be checked, and all the main phenomenological scales can be generated with the following achievements:

(i) Correct amount of CMB density perturbations;
(ii) Right scale for grand unification theories (GUT);
(iii) TeV-scale SUSY;
(iv) No cosmological moduli problem (CMP).

II. TYPE IIB LARGE VOLUME SCENARIOS

String compactifications typically give rise to a large number of moduli which can be ideal candidates to drive inflation. It is crucial to lift these flat directions in order to determine the features of the low energy effective field theory (like masses and coupling constants) and to avoid the presence of unobserved long range fifth forces.

A. Moduli stabilization

Moduli stabilization is best understood in the context of type IIB where background fluxes fix the dilaton and the complex structure moduli at tree level. On the other hand, the stabilization of the Kähler moduli requires a consideration of perturbative and nonperturbative effects [13–15]. Expressing the Kähler moduli as $T_i = \tau_i + i b_i$, $i = 1, \ldots, h_{1,1}$, with τ_i the volume of an internal 4-cycle Σ_i and $b_i = \int_{\Sigma_i} C_4$, we shall focus on compactification manifolds whose volume reads (with $\alpha > 0$, $\gamma_i > 0\ \forall i$):

$$\hat{\mathcal{V}} = \alpha(\tau_1^{3/2} - \gamma_2 \tau_2^{3/2} - \gamma_3 \tau_3^{3/2} - \gamma_4 \tau_4^{3/2}). \qquad (1)$$

Assuming that the tadpole-cancellation condition can be satisfied by an appropriate choice of background fluxes, we wrap a hidden sector $D7$-stack that undergoes gaugino condensation both around τ_2 and τ_3. This brane setup induces a superpotential of the form (the dilaton and the complex structure moduli are flux-stabilized at tree level and so they can be integrated out):

$$W = W_0 + A_2 e^{-a_2 T_2} + A_3 e^{-a_3 T_3}. \qquad (2)$$

We focus on the case when the cycle τ_4 supporting the visible sector shrinks down at the singularity. The visible sector is built via $D3$-branes at the quiver locus and the gauge coupling is given by the dilaton s (with $\langle s \rangle = g_s^{-1}$) while τ_4 enters as a flux-dependent correction: $4\pi g^{-2} = s + h(F)\tau_4$. The Kähler potential with the leading order α' correction can be expanded around $\tau_4 = 0$ [16]:

$$K = -2\ln\left(\mathcal{V} + \frac{\xi s^{3/2}}{2}\right) + \lambda \frac{\tau_4^2}{\mathcal{V}} - \ln(2s), \qquad (3)$$

with $\mathcal{V} = \alpha(\tau_1^{3/2} - \gamma_2 \tau_2^{3/2} - \gamma_3 \tau_3^{3/2})$. In the absence of standard sodel singlets, which can get a nonvanishing vacuum expectation value (VEV), an anomalous $U(1)$ on the visible sector cycle generates a D-term potential with a Fayet-Iliopoulos term:

$$V_D = \left(\frac{2\pi}{s + h(F)\tau_4}\right)\xi_{FI}^2 \quad \text{with} \quad \xi_{FI} = \frac{Q_{\tau_4}\tau_4}{\mathcal{V}}, \qquad (4)$$

while the leading order F-term potential reads (after minimizing the axion directions):

$$V_F = \sum_{i=2}^{3} \frac{8 a_i^2 A_i^2}{3\alpha\gamma_i} \frac{\sqrt{\tau_i} e^{-2a_i\tau_i}}{\mathcal{V}} - 4W_0 \sum_{i=2}^{3} a_i A_i \frac{\tau_i e^{-a_i\tau_i}}{\mathcal{V}^2} + \frac{3\xi W_0^2}{4 g_s^{3/2} \mathcal{V}^3}. \qquad (5)$$

The D-term potential scales as $V_D \sim \mathcal{O}(\mathcal{V}^{-2})$ while in the regime $\mathcal{V} \sim e^{a_i\tau_i}$, $i = 2, 3$, $V_F \sim \mathcal{O}(\mathcal{V}^{-3})$. Hence, at leading order $\xi_{FI} = 0$ leads to $\tau_4 \to 0$, fixing this cycle at the quiver locus [16]. On the other hand, V_F completely stabilizes τ_2, τ_3 and the volume $\mathcal{V} \simeq \alpha \tau_1^{3/2}$ at:

$$a_i \langle \tau_i \rangle = \frac{1}{g_s}\left(\frac{\xi}{2\alpha J}\right)^{2/3} \quad \text{with} \quad J = \sum_{i=2}^{3} \gamma_i / a_i^{3/2}, \qquad (6)$$

and $\langle \mathcal{V} \rangle = (\frac{3\alpha\gamma_i}{4 a_i A_i}) W_0 \sqrt{\langle \tau_i \rangle} e^{a_i \langle \tau_i \rangle}$, $\forall i = 2, 3$.

Moduli stabilization is performed without fine-tuning the internal fluxes ($W_0 \sim \mathcal{O}(1)$) and the volume is fixed exponentially large in string units. As a consequence, one has a very reliable effective field theory, as well as a tool for the generation of phenomenologically desirable hierarchies.

B. Particle physics phenomenology

The particle phenomenology is governed by the background fluxes, which break SUSY by the F-terms of the Kähler moduli and the dilaton, which then mediate this breaking to the visible sector. However, the F-term of τ_4 vanishes since it is proportional to $\xi_{FI} = 0$: $F^4 \sim e^{K/2} K^{4\bar{4}} W_0 \xi_{FI} = 0$. Thus, there is no local SUSY-breaking and the visible sector is sequestered, implying that the soft terms can be suppressed with respect to $m_{3/2}$ by an inverse power of the volume. The main scales in the model are:

(i) GUT-scale: $M_{\rm GUT} \sim M_P / \mathcal{V}^{1/3}$,
(ii) String-scale: $M_s \sim M_P / \mathcal{V}^{1/2}$,
(iii) Kaluza-Klein scale: $M_{KK} \sim M_P / \mathcal{V}^{2/3}$,
(iv) Gravitino mass: $m_{3/2} \sim M_P / \mathcal{V}$,
(v) blowup modes: $m_{\tau_i} \sim m_{3/2}$, $i = 2, 3$,
(vi) Volume mode: $m_{\mathcal{V}} \sim M_P / \mathcal{V}^{3/2}$,
(vii) Soft-terms: $M_{\rm soft} \sim m_{3/2}^2 / M_P \sim M_P / \mathcal{V}^2$.

Setting the volume $\mathcal{V} \simeq 10^{6-7}$ in string units, corresponding to $M_s \simeq 10^{15}$ GeV, one can realize GUT theories, TeV-scale SUSY, and avoid any CMP [16].

C. Inflationary cosmology

A very promising inflationary model can be embedded in this type IIB scenario with the inflaton, which is the size of the blowup τ_2 [6]. Displacing τ_2 far from its minimum, due to the exponential suppression, this field experiences a very flat direction which is suitable for inflation. The other blowup τ_3, which sits at its minimum while τ_2 is slow-rolling, has been added to keep the volume stable during inflation. In terms of the canonically normalized inflaton ϕ, the potential looks like [6]:

$$V \simeq V_0 - \beta\left(\frac{\phi}{\mathcal{V}}\right)^{4/3} e^{-a\mathcal{V}^{2/3}\phi^{4/3}}. \tag{7}$$

This is a model of small field inflation, and so no detectable gravity waves are produced: $r \equiv T/S \ll 1$. The spectral index is in good agreement with the observations: $0.960 < n_s < 0.967$, and the requirement of generating enough density perturbations fixes $\mathcal{V} \simeq 10^{6-7}$, which is the same value preferred by particle physics.

Potential problems come from g_s corrections to K [17] which dominate the potential, spoiling its flatness once τ_2 is displaced far from its minimum. The only way out is to fine-tune these g_s corrections on a small scale [8].

In principle, the nonperturbative potential for τ_2 could also be generated by a $D3$-brane instanton. In this way, hidden sectors and g_s corrections would be absent. However, the requirement of having $\mathcal{V} \simeq 10^{6-7}$ prevents this setup since both τ_2 and τ_3 would be fixed smaller than the string-scale where the effective field theory cannot be trusted anymore [8]. Thus, we realize that hidden sectors are always present in these models.

D. Hidden sector configurations

The hidden sectors on τ_i, $i = 2, 3$ consist in a supersymmetric field theory that undergoes gaugino condensation. Broadly speaking, we can entertain 3 scenarios for the possible particle content and mass spectrum:

(i) The hidden sector is a pure $N = 1$ supersymmetric Yang-Mills (SYM) theory which, due to strong dynamics, confines in the ir at the scale Λ:

$$\Lambda_i = M_s e^{-(4\pi g^{-2})a_i/3}, \qquad i = 2, 3. \tag{8}$$

Given that $4\pi g^{-2} = \tau_i$, and at the minimum $e^{-a_i\tau_i} \sim \mathcal{V}^{-1}$, the order of magnitude of Λ_i can be estimated as $\Lambda_i \simeq M_P \mathcal{V}^{-5/6}$. The theory develops a mass gap and all particles acquire a mass of the order Λ_i and are heavier than the inflaton after inflation since $m_{\tau_2} \simeq M_P/\mathcal{V} < \Lambda_i$. Thus, the inflaton decay to hidden degrees of freedom is *kinematically forbidden.*

(ii) The hidden sector is a pure SYM theory plus a massless $U(1)$. The mass-spectrum below Λ consists of massless hidden photons and photini with an $\mathcal{O}(M_{\rm soft})$ mass due to SUSY-breaking effects.

(iii) The hidden sector is an $N = 1$ $SU(N_c)$ theory with $N_f < (N_c - 1)$ flavors. The condensates of gauge bosons and gauginos get a mass of the order Λ while all the matter condensates get an $\mathcal{O}(M_{\rm soft})$ mass due to SUSY-breaking effects except pionlike mesons, which remain massless in the presence of spontaneous chiral symmetry breaking. If chiral symmetry is explicitly broken by a low energy Higgs-like mechanism, all matter fields get a $\delta m \ll M_{\rm soft}$ correction to their masses. For an additional massless $U(1)$, there are also massless hidden photons and photini with an $\mathcal{O}(M_{\rm soft})$ mass.

III. REHEATING

We shall now focus on the study of reheating of the MSSM degrees of freedom after the end of inflation. In order to check what fraction of the inflaton energy is transferred to hidden and visible degrees of freedom, we have to derive the moduli mass spectrum and their couplings to all particles in the model.

A. Canonical normalization and mass spectrum

The first step is to canonically normalize the moduli around the minimum of their VEVs: $\tau_i = \langle\tau_i\rangle + \delta\tau_i$, $\forall i$. The fluctuations $\delta\tau_i$ are written in terms of the canonically normalized fields $\delta\phi_i$ as $\delta\tau_i = \frac{1}{\sqrt{2}}C_{ij}\delta\phi_j$, where C_{ij} are the eigenvectors of the matrix $(M^2)_{ij} \equiv \frac{1}{2}(K^{-1})_{ik}V_{kj}$ whose eigenvalues m_i^2 are the moduli mass-squareds.

The form of K, Eq. (3), and $\langle\tau_4\rangle = 0$ imply that at leading order τ_2 does not mix with τ_4. However, τ_2 mixes with s due to α' corrections to K [8]:

$$\begin{aligned}
&\delta\tau_1 \sim \mathcal{O}(\mathcal{V}^{2/3})\delta\phi_1 + \sum_i \mathcal{O}(\mathcal{V}^{1/6})\delta\phi_i,\ i = 2, 3, s,\\
&\delta\tau_i \sim \mathcal{O}(\mathcal{V}^{1/2})\delta\phi_i + \mathcal{O}(1)\delta\phi_1\\
&\qquad + \sum_j \mathcal{O}(\mathcal{V}^{-1/2})\delta\phi_j,\ i = 2, 3,\\
&\delta\tau_4 \sim \mathcal{O}(\mathcal{V}^{1/2})\delta\phi_4,\ \delta s \sim \mathcal{O}(1)\delta\phi_s + \mathcal{O}(\mathcal{V}^{-1/2})\delta\phi_1\\
&\qquad + \sum_i \mathcal{O}(\mathcal{V}^{-1})\delta\phi_i,\ i = 2, 3,
\end{aligned} \tag{9}$$

with $j = 2, 3, s$, $j \neq i$. The masses turn out to be

$$m_1^2 \simeq \frac{M_P^2}{\mathcal{V}^3}, \qquad m_i^2 \simeq \frac{M_P^2}{\mathcal{V}^2}, \qquad \forall i = 2, 3, s, \quad \text{and} \quad m_4^2 \simeq \frac{M_P^2}{\mathcal{V}}. \tag{13}$$

B. Inflaton couplings

The inflaton coupling to visible and hidden degrees of freedom can be derived from the moduli dependence of the kinetic and mass terms of open string modes. The moduli are expanded around their VEVs and then expressed in terms of the canonically normalized fields. This procedure led to the derivation of the moduli couplings to all particles in the model [8,9], finding that the strongest moduli decay rates are to hidden gauge bosons on τ_2 and τ_3, and to visible gauge bosons at the τ_4-singularity (see Table I).

C. Moduli dynamics after inflation and reheating

At the end of inflation, due to the steepness of the potential, the inflaton τ_2, which acts like a homogeneous condensate, stops oscillating coherently around its minimum just after 2–3 oscillations due to a very violent nonperturbative production of $\delta\tau_2$ quanta [12]. The production of other degrees of freedom at preheating is instead less efficient.

According to the second of Eqs. (9), our Universe is mostly filled with $\delta\phi_2$ plus some $\delta\phi_1$ and fewer $\delta\phi_3$ and $\delta\phi_s$-particles. Thus, the energy density is dominated by $\delta\phi_2$, whose perturbative decay leads to reheating. Denoting the visible gauge bosons as g , and X_2 and X_3 as the hidden ones, the coupling of $\delta\phi_2$ to X_2X_2 is stronger than the one to X_3X_3 which, in turn, is stronger than the one to gg. This is due to the geometric separation between τ_2 and τ_3, and the sequestering of the visible sector at the τ_4-singularity.

Hence, the first decays are $\delta\phi_i \to X_iX_i$, $i = 2, 3$ with decay rate $\Gamma \sim M_P/\mathcal{V}^2$. Thus, the inflaton dumps all its energy to hidden, instead of visible, degrees of freedom without reheating the visible sector. We stress that there is no direct coupling between hidden and visible degrees of freedom since they correspond to two open string sectors localized in different regions of the Calabi-Yau, and so the reheating of the visible sector cannot occur via the decay of hidden to visible degrees of freedom. Hence, the only way out is to forbid the decay of $\delta\phi_2$ to any hidden particle. This forces us to consider both τ_2 and τ_3 as a pure $N = 1$ SYM theory that develops a mass gap, so that the decay of $\delta\phi_2$ to X_iX_i, $i = 2, 3$ is kinematically forbidden.

Then, the first decay is $\delta\phi_s \to gg$ with $\Gamma \sim M_P/\mathcal{V}^3$ but without leading to reheating since the energy density is dominated by $\delta\phi_2$. Reheating occurs only later on when $\delta\phi_2$ decays to visible gauge bosons with total decay rate $\Gamma^{\rm TOT}_{\delta\phi_2\to gg} \simeq (\ln\mathcal{V})^3 M_P \mathcal{V}^{-5}$ [8]. At the same time, $\delta\phi_3 \to gg$ without giving rise to reheating since $\delta\phi_3$ is not dominating the energy density.

TABLE I. Moduli couplings to all gauge bosons in the model.

	$\delta\phi_1$	$\delta\phi_2$	$\delta\phi_3$	$\delta\phi_4$	$\delta\phi_s$
$(F^{(2)}_{\mu\nu}F^{\mu\nu}_{(2)})$	$\frac{1}{M_P}$	$\frac{\mathcal{V}^{1/2}}{M_P}$	$\frac{1}{\mathcal{V}^{1/2}M_P}$	$\cdots$	$\frac{1}{\mathcal{V}^{1/2}M_P}$
$(F^{(3)}_{\mu\nu}F^{\mu\nu}_{(3)})$	$\frac{1}{M_P}$	$\frac{1}{\mathcal{V}^{1/2}M_P}$	$\frac{\mathcal{V}^{1/2}}{M_P}$	$\cdots$	$\frac{1}{\mathcal{V}^{1/2}M_P}$
$(F^{(4)}_{\mu\nu}F^{\mu\nu}_{(4)})$	$\frac{1}{\mathcal{V}^{1/2}M_P}$	$\frac{1}{\mathcal{V}M_P}$	$\frac{1}{\mathcal{V}M_P}$	$\frac{\mathcal{V}^{1/2}}{M_P}$	$\frac{1}{M_P}$

The maximal reheating temperature for the visible sector in the approximation of sudden thermalization can be worked out equating $\Gamma^{\rm TOT}_{\delta\phi_2\to gg}$ to $H \simeq (T^{\rm max}_{\rm RH})^2/M_P$ [8]:

$$T^{\rm max}_{\rm RH} \simeq \sqrt{\Gamma^{\rm TOT}_{\delta\phi_2\to gg} M_P} \simeq (\ln\mathcal{V})^{3/2}\frac{M_P}{\mathcal{V}^{5/2}}. \qquad (10)$$

For $\mathcal{V} \simeq 10^{6-7}$, we obtain $T^{\rm max}_{\rm RH} \simeq 10^{2-4}$ GeV which is higher than $T_{\rm BBN} \simeq 1$ MeV, so it does not create any problem if the matter-antimatter asymmetry could be realized in a nonthermal/thermal way. Later on, $\delta\phi_1$ decays to visible degrees of freedom out of thermal equilibrium without suffering from the CMP since it decays before BBN: $T_{\delta\phi_1\to gg} \simeq M_P \mathcal{V}^{-11/4} \sim 10^2$ GeV.

IV. DISCUSSION: INTERPLAY BETWEEN GLOBAL AND LOCAL ISSUES

In this paper, we did not address any of the issues related to the brane construction of the visible sector given that the actual details of the embedding of the standard model (or any generalization thereof like the MSSM or GUT theories) into string theory are completely irrelevant for the study of reheating in the context of closed-string inflationary models within the framework of type IIB compactifications.

In fact, type IIB Calabi-Yau flux compactifications are characterized by the fact that physics decouples into local (or brane) and global (or bulk) issues that can be consistently studied separately. Some examples of brane issues that depend only on the local brane construction are finding the right chiral spectrum, gauge group and Yukawa couplings, while issues like moduli stabilization, supersymmetry-breaking, inflation, and reheating are purely global.

Regarding the study of the transfer of the inflaton energy density to the standard model degrees of freedom, we emphasize that the only thing which has to be considered is the inflaton dynamics after inflation with the nonperturbative (preheating) and perturbative (reheating) particle production. In particular, the study of reheating via the inflaton decay involves only the knowledge of the inflaton mass and couplings to all the degrees of freedom in the theory, and the computation of the overall scaling of these couplings does not depend at all on any local detail of the standard model brane construction.

In fact, it is the overall volume of the Calabi-Yau compactification which sets the order of magnitude of all these couplings whose computation requires only the knowledge of the moduli dependence of the gauge kinetic function, the Kähler potential, and the superpotential, together with the form of the soft supersymmetry-breaking terms. All these are global or bulk issues.

V. CONCLUSIONS

In this paper, we presented a model of closed-string slow-roll inflation embedded in type IIB string compactifications where we have full control over the inflaton dynamics after inflation and the reheating of the visible sector degrees of freedom, regardless of the local standard model construction.

We tried to bring string inflation closer to ''the real world,'' describing how to excite the visible sector degrees of freedom in a Calabi-Yau compactification with a robust moduli stabilization mechanism that allows us to check the solution of the η-problem for inflation and compute the order of magnitude of the moduli mass-spectrum and coupling to all particles of the theory.

This paper is mostly based on [8] but it singles out for the first time the best working model of closed-string inflation which shows many interesting phenomenological features like the generation of the correct amount of density perturbations, a viable reheating of the visible sector degrees of freedom which can take place substantially before big-bang nucleosynthesis, the absence of any cosmological moduli problem, the presence of TeV-scale supersymmetry, and the right scale for grand unification theories. We stress that all these achievements can be reached without the need for any fine-tuning.

In this sense, this paper represents a step forward with respect to [8], which was just a general analysis of reheating with the discussion of several models, and the pointing out of some problems that one has to solve to have a viable reheating model.

ACKNOWLEDGMENTS

The research of A. M. is partly supported by the UNIVERSENET (MRTN-CT-2006-035863).

[1] E. Komatsu *et al.*, Astrophys. J. Suppl. Ser., **192**, 18 (2011).
[2] A. Mazumdar and J. Rocher, Phys. Rep. **497**, 85 (2011).
[3] B. Fields and S. Sarkar, J. Phys. G **33**, 1 (2006); R. H. Cyburt, Phys. Rev. D **70**, 023505 (2004).
[4] R. Allahverdi, K. Enqvist, J. Garcia-Bellido, and A. Mazumdar, Phys. Rev. Lett. **97**, 191304 (2006); R. Allahverdi, A. Kusenko, and A. Mazumdar, J. Cosmol. Astropart. Phys. 07 (2007) 018; R. Allahverdi, K. Enqvist, J. Garcia-Bellido, A. Jokinen, and A. Mazumdar, J. Cosmol. Astropart. Phys. 06 (2007) 019.
[5] R. Allahverdi, B. Dutta, and A. Mazumdar, Phys. Rev. D **75**, 075018 (2007); R. Allahverdi, B. Dutta, and A. Mazumdar, Phys. Rev. Lett. **99**, 261301 (2007).
[6] J. P. Conlon and F. Quevedo, J. High Energy Phys. 1, (2006)146.
[7] M. Cicoli, C. P. Burgess, and F. Quevedo, J. Cosmol. Astropart. Phys. 03 (2009) 013; M. Cicoli, Fortschr. Phys. **58**, 115 (2009); C. Burgess, M. Cicoli, M. Gomez-Reino, F. Quevedo, G. Tasinato, and I. Zavala, J. High Energy Phys. 08 (2010) 045.
[8] M. Cicoli and A. Mazumdar, J. Cosmol. Astropart. Phys. 09 (2010) 025.
[9] J. P. Conlon and F. Quevedo, J. Cosmol. Astropart. Phys. 08 (2007) 019; L. Anguelova, V. Calo, and M. Cicoli, J. Cosmol. Astropart. Phys. 10 (2009) 025.
[10] R. Allahverdi, R. Brandenberger, F. Y. Cyr-Racine, and A. Mazumdar, Annu. Rev. Nucl. Part. Sci. **60**, 27 (2010).
[11] N. Barnaby, C. P. Burgess, and J. M. Cline, J. Cosmol. Astropart. Phys. 04 (2005) 007; A. Mazumdar and H. Stoica, Phys. Rev. Lett. **102**, 091601 (2009); R. H. Brandenberger, K. Dasgupta, and A. C. Davis, Phys. Rev. D **78**, 083502 (2008); L. Kofman and P. Yi, Phys. Rev. D **72**, 106001 (2005); A. R. Frey, A. Mazumdar, and R. C. Myers, Phys. Rev. D **73**, 026003 (2006); X. Chen and S. H. Tye, J. Cosmol. Astropart. Phys. 06 (2006) 011; D. R. Green, Phys. Rev. D **76**, 103504 (2007).
[12] N. Barnaby, J. R. Bond, Z. Huang, and L. Kofman, J. Cosmol. Astropart. Phys. 12 (2009) 021.
[13] S. Kachru, R. Kallosh, A. Linde, and S. P. Trivedi, Phys. Rev. D **68**, 046005 (2003).
[14] V. Balasubramanian, P. Berglund, J. P. Conlon, and F. Quevedo, J. High Energy Phys. 03 (2005) 007.
[15] M. Cicoli, J. P. Conlon, and F. Quevedo, J. High Energy Phys. 10 (2008) 105.
[16] J. P. Conlon, A. Maharana, and F. Quevedo, J. High Energy Phys. 05 (2009) 109; R. Blumenhagen, J. P. Conlon, S. Krippendorf, S. Moster, and F. Quevedo, J. High Energy Phys. 09 (2009) 007.
[17] M. Berg, M. Haack, and B. Kors, J. High Energy Phys. 11 (2005) 030; M. Berg, M. Haack, and E. Pajer, J. High Energy Phys. 09 (2007) 031; M. Cicoli, J. P. Conlon, and F. Quevedo, J. High Energy Phys. 01 (2008) 052.