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Chaotic inflation with a fractional power-law potential in strongly coupled gauge theories

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ABSTRACT

Models of chaotic inflation with a fractional power-law potential are not only viable but also testable in the foreseeable future. We show that such models can be realized in simple strongly coupled supersymmetric gauge theories. In these models, the energy scale during inflation is dynamically generated by the dimensional transmutation due to the strong gauge dynamics. Therefore, such models not only explain the origin of the fractional power in the inflationary potential but also provide a reason why the energy scale of inflation is much smaller than the Planck scale.

1. Introduction

Cosmic inflation [1] is a very successful paradigm of modern cosmology which explains the origin of the anisotropies of the Cosmic Microwave Background (CMB) as well as of the Large Scale Structure of the Universe [2,3]. At present, a realistic and complete theory of inflation is, however, still pending and hence, inflationary model building remains one of the most important tasks of particle physics and cosmology.

Among the various classes of inflation models proposed so far, the chaotic inflation scheme [4] is one of the most attractive classes since it can realize an inflationary expansion even in the presence of large quantum fluctuations at the Planck time. Moreover, the large field values typically encountered in models of chaotic inflation imply a large contribution from gravitational waves to the CMB power spectrum [5], rendering these models testable in the foreseeable future. However, according to the precise observations of the CMB anisotropies, the simplest versions of chaotic inflation, i.e. the models with a quadratic potential or a quartic potential, are now somewhat disfavored [6]. With the forthcoming data provided by the Planck satellite experiment [7], the constraints on those simplest versions will be improved upon.

In light of this situation, a more general version of chaotic inflation has been gathering attention, in which the inflaton potential comes with a fractional power. As analyzed in Refs. [8], the models of chaotic inflation with a fractional power-law potential are more favored than the simplest versions of chaotic inflation. Despite such successes, these models, however, lack a firm field theoret-

ical foundation, which apparently seems difficult to be achieved from regular field theories.[1]

In this Letter, we show that such fractional power-law chaotic inflation models can be realized in simple strongly coupled supersymmetric gauge theories. There, the energy scale during inflation is generated by the dimensional transmutation due to the strong gauge dynamics. Thus, these models not only explain the origin of the fractional power in the inflationary potential but also provide a reason why the energy scale of inflation is much smaller than the Planck scale.[2]

The organization of the Letter is as follows. First, we derive the fractional power-law potential for the inflaton in strongly coupled gauge theories. Next, we discuss distortions of the inflaton potential due to supergravity contributions. Then, we outline the phenomenology of chaotic inflation with the dynamically generated potential and summarize its observational consequences. The final section is devoted to conclusions and discussion.

2. Dynamical generation of the inflaton potential

Let us discuss how the fractional power-law potential of the inflaton is generated dynamically. First, we consider an $SP(N)$ supersymmetric gauge theory[3] with $2(N+2)$ chiral superfields in the fundamental representation, Q^I ($I = 1 \cdots 2(N + 2)$). Besides the

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[1] A recipe for embedding chaotic inflation with a fractional power-law potential into supergravity is provided in Refs. [9,10]; see also Ref. [11], in which a fractional power-law potential is obtained from a running kinetic term for the inflaton. For fractional power-law potentials derived in string theories, see Ref. [12].

[2] For examples of models in which the scale of inflation is generated dynamically, see Refs. [13].

[3] We use the convention where $SP(1)$ is $SU(2)$.

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fundamental representations, we also introduce $(N + 2)(2N + 3)$ gauge-singlet chiral superfields $\mathcal{Z}_{IJ}(= -\mathcal{Z}_{JI})$ which couple to the fundamental representations in the superpotential via

$$W = \frac{1}{2}\lambda_{IJ}\mathcal{Z}_{IJ}Q^I Q^J, \tag{1}$$

with coupling constants λ_{IJ}. It should be noted that all the quantum moduli, $Q^I Q^J$, are lifted by the couplings to the gauge singlets $\mathcal{Z}^{IJ}$.

For a later purpose, we decompose the above fields into,

$$
\begin{aligned}
Z_{ij} &= \mathcal{Z}_{IJ=ij}, & Q^i &= Q^{I=i}, \\
T_i &= \mathcal{Z}_{IJ=i(2N+3)}, & P &= Q^{I=2N+3}, \\
\bar{T}_i &= \mathcal{Z}_{IJ=i(2N+4)}, & \bar{P} &= Q^{I=2N+4}, \\
S &= \mathcal{Z}_{IJ=(2N+3)(2N+4)}, &&
\end{aligned}
\tag{2}
$$

where $i, j = 1 \cdots 2(N + 1)$. In terms of these fields, the above superpotential is now rewritten as

$$W = \frac{1}{2}\lambda Z_{ij}Q^i Q^j + \lambda_T T_i Q^i P + \lambda_{\bar{T}}\bar{T}_i Q^i \bar{P} - \lambda_S S P \bar{P}. \tag{3}$$

We have assumed $\lambda_{ij} = \lambda$ for simplicity and we also assume that $\lambda_{T,\bar{T}}$ are larger than λ in the following. The sign convention for λ_S is just for later convenience. As we will show, the scalar potential for S generated by the strong dynamics features a fractional power and eventually plays the role of the inflaton potential.

To see how the scalar potential is generated, let us remember that the $SP(N)$ gauge theory with $2(N + 2)$ fundamental representations exhibits the so-called s-confinement [14,15] at low energies below the dynamical scale Λ. In this phase, the model is well described by the composite fields, $M \propto QQ$, $M_P \propto QP$, $M_{\bar{P}} \propto Q\bar{P}$ and $M_{P\bar{P}} \propto P\bar{P}$ [14], which may be assembled in the same way in an antisymmetric matrix $\mathcal{M}$ as the gauge singlets Z, T, $\bar{T}$, and S are assembled in the antisymmetric matrix $\mathcal{Z}$. Their effective superpotential is given by[4]

$$
\begin{aligned}
W_{\text{eff}} = {} & \frac{Pf^{(N+2)}(\mathcal{M})}{\Lambda^{(N+2)-3}} + \frac{1}{2}\lambda\Lambda Z_{ij}M^{ij} + \lambda_T \Lambda T_i M_P^i + \lambda_{\bar{T}}\Lambda\bar{T}_i M_{\bar{P}}^i \\
& - \lambda_S \Lambda S M_{P\bar{P}}.
\end{aligned}
\tag{4}
$$

The first term is the non-perturbative potential generated by the s-confinement.[5] The other four terms in Eq. (4) can be regarded as mass-mixing operators between the composite fields and the gauge singlets, inducing supersymmetric masses of $\mathcal{O}(\lambda\Lambda, \lambda_T\Lambda, \lambda_{\bar{T}}\Lambda, \lambda_S\Lambda)$, respectively. The effective superpotential shows that the model possesses a supersymmetric vacuum in which all of the M's and singlets vanish. That is, as expected, there is a vacuum with vanishing vacuum energy.

Now, let us consider an effective potential for $S \neq 0$ and $X(\propto J^{ij}Z_{ij}) \neq 0$ around the vacuum. Here, J_{ij} are the components of the symplectic form $J = \mathbb{1}_n \otimes i\sigma_2$. Notice that the mesons M_P and $M_{\bar{P}}$ are still fixed at

$$M_P = 0, \qquad M_{\bar{P}} = 0, \tag{5}$$

since we have assumed $\lambda_{T,\bar{T}} \gg \lambda$. Thus, the effective potential in Eq. (4) is reduced to

$$W_{\text{eff}} = M_{P\bar{P}}\left(\frac{Pf^{(N+1)}(M)}{\Lambda^{(N+2)-3}} - \lambda_S \Lambda S\right) + \frac{1}{2}\lambda\Lambda Z_{ij}M^{ij}. \tag{}$$

The first term leads to the so-called deformed moduli constraint on the moduli space of the $SP(N)$ gauge theory with $2(N + 1)$ fundamental representations [14]. Therefore, for a given non-vanishing S, the above model is nothing but the dynamical supersymmetry breaking model of Ref. [16]. By solving the quantum deformed constraint for $S \neq 0$, we obtain

$$M_{ij} = \Lambda\left(\frac{\lambda_S S}{\Lambda}\right)^{\frac{1}{N+1}} \times J_{ij}, \tag{6}$$

which leads to

$$W_{\text{eff}} = \lambda(N + 1)^{1/2}\Lambda^2\left(\frac{\lambda_S S}{\Lambda}\right)^{\frac{1}{N+1}} X, \tag{7}$$

where X is defined by $X = Z_{ij}J^{ij}/(2(N + 1)^{1/2})$. Hence, for $S \neq 0$, supersymmetry is broken by the F-component of X, which leads to a scalar potential for S,

$$V \simeq \lambda^2(N + 1)\Lambda^4\left(\frac{\lambda_S|S|}{\Lambda}\right)^{\frac{2}{N+1}}. \tag{8}$$

As promised, we find that the scalar component of the singlet S obtains a fractional power-law potential,

$$V \propto |S|^p. \tag{9}$$

Its power is solely determined by the size of the $SP(N)$ gauge group,

$$p = \frac{2}{N+1}. \tag{10}$$

In the above analysis, we have tacitly assumed that the field value of S is around or below the dynamical scale, i.e. $\lambda_S S \lesssim \Lambda$. The above scalar potential is, however, also obtained for $\lambda_S S \gg \Lambda$. For that purpose, let us remember that P and $\bar{P}$ are heavier than the dynamical scale for $\lambda_S S \gg \Lambda$ and decouple perturbatively. Thus, the effective theory below $\lambda_S S$ consists of the $SP(N)$ gauge theory with $2(N + 1)$ fundamental representations whose effective dynamical scale is given by [17]

$$\Lambda_{\text{eff}} = \Lambda \times \left(\frac{\lambda_S|S|}{\Lambda}\right)^{\frac{1}{2(N+1)}}. \tag{11}$$

Then, since the effective theory below $\lambda_S S$ is again the dynamical supersymmetry breaking model, we again reach the effective superpotential in Eq. (7). Therefore, again, supersymmetry is broken by the F-component of X at $S \neq 0$, which again leads to[6]

$$V \simeq \lambda^2(N + 1)\Lambda^4\left(\frac{\lambda_S|S|}{\Lambda}\right)^{\frac{2}{N+1}}. \tag{12}$$

In the following sections, we assume that S plays the role of the inflaton. After inflation, S reaches its origin, which leads to the restoration of supersymmetry.

Before discussing the details of our scenario of chaotic inflation, let us comment on a seemingly contradictory point concerning the inflaton potential. According to Eq. (4) all fields have masses.

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[4] It should be noted that we have neglected $O(1)$ differences between the λ's in Eq. (3) and the ones in Eq. (4) due to non-perturbative effects. We also assume that the composite fields in $\mathcal{M}$ are close to the canonically normalized ones.

[5] In this Letter, we define the Pfaffian of a $2n \times 2n$ antisymmetric matrix, $Pf^{(n)}$, so that the symplectic form J, where $J = \mathbb{1}_n \otimes i\sigma_2$ with $\mathbb{1}_n$ being the $n \times n$ unit matrix and σ_2 the second Pauli matrix, satisfies $Pf^{(n)}(J) = 1$.

[6] The agreement between the powers in the potential in the two different field regimes is not a coincidence but can be understood by remembering that the effective superpotential consistent with (anomalous) R-symmetries, holomorphicity and dimensional analysis should have the form in Eq. (7), regardless of the sizes of $X \neq 0$ and $S \neq 0$.

On the other hand, Eq. (8) shows that the potential for S is not quadratic in S around the origin, but rather features a singular first derivative. This difference is due to the fact that we have ignored a non-minimal contribution to the Kähler potential in the derivation of Eq. (8). When integrating out the meson fields, we have used the constraint (6), which also generates the following non-minimal Kähler potential for S,

$$\delta K = \frac{1}{2} M_{ij}^{\dagger} M_{ij} = (N+1)\Lambda^2 \left(\frac{\lambda_S |S|}{\Lambda} \right)^p. \tag{13}$$

For $\lambda_S S \ll \Lambda$, this contribution dominates over the canonical Kähler potential. The potential for the reparametrized field $S' = \sqrt{N+1}\Lambda(\lambda_S S/\Lambda)^{p/2}$, which has a canonical Kähler potential for $\lambda_S S \ll \Lambda$, is then given by

$$V \simeq \lambda^2 \Lambda^2 |S'|^2, \tag{14}$$

which is quadratic in S', thus implying that the canonically normalized field S' is massive. On the other hand, at larger values of S, the canonical Kähler potential of S dominates over the non-canonical one in Eq. (13) and the potential in Eq. (12) is valid after all.

3. Supergravity contributions

In our model, we apply the above obtained fractional potential to the chaotic inflation scenario, where the field value of the inflaton exceeds the Planck scale M_{Pl}. For such a large field value, we need to carefully examine the supergravity contributions to the scalar potential, which could change the potential drastically from the fractional power-law potential. In fact, if we assume, for example, a minimal Kähler potential $S^{\dagger}S$ for the inflaton S, the inflaton potential is modified in supergravity as follows,

$$V \simeq e^{|S|^2/M_{\mathrm{Pl}}^2} \times \lambda^2 (N+1)\Lambda^4 \left(\frac{\lambda_S |S|}{\Lambda} \right)^{\frac{2}{N+1}}, \tag{15}$$

which is too steep for chaotic inflation for $S \gg M_{\mathrm{Pl}}$.[7]

The most attractive way to avoid such supergravity contributions is to introduce a shift symmetry in the direction of S [18,19],

$$S \rightarrow S + ic, \quad c \in \mathbb{R} \tag{16}$$

(see also Ref. [9] for recent developments). With this shift symmetry, the Kähler potential is a function of $S + S^{\dagger}$,

$$K = \frac{1}{2} |S + S^{\dagger}|^2 + \cdots, \tag{17}$$

and hence, it does not depend on the imaginary component of S, $\Im(S)$. Therefore, the imaginary component of S has a fractional power-law potential even for $\Im(S) \gg M_{\mathrm{Pl}}$, while the real part of S obtains a Hubble mass term around $\Re(S) = 0$. In the following, $\sqrt{2}\Im(S)$ plays the role of the inflaton in the chaotic inflation scenario with the dynamically generated fractional power-law potential in Eq. (12).[8]

One caveat to the above argument is that the shift symmetry is explicitly broken by the interaction

$$W \supset -\lambda_S S P \bar{P} \tag{18}$$

in Eq. (3). With this explicit breaking, the Kähler potential in Eq. (17) obtains a radiative correction which breaks the symmetry,

$$\delta K \sim \frac{2N\lambda_S^2}{16\pi^2} |S|^2 \log\left(\frac{\mu^2}{M_{\mathrm{Pl}}^2} \right), \tag{19}$$

where μ is a renormalization scale.[9] This breaking term causes a steep exponential potential of $\Im(S)$ for $\Im(S) \gg M_{\mathrm{Pl}}$ unless λ_S is small enough. Thus, to avoid a too large breaking of the shift symmetry, we assume in the following analysis that λ_S is rather suppressed.[10]

Furthermore, we note that in general the mere introduction of a shift symmetry in the direction of the inflaton field S does not suffice to protect chaotic inflation from receiving disastrously large supergravity corrections. In addition we have to require that the superpotential be of the form $W = Xf(S)$ [18,9], where f is an arbitrary holomorphic function of S and X is a gauge singlet that can be identified as the goldstino superfield responsible for the spontaneous breaking of supersymmetry during inflation [10]. Evidently, the effective superpotential in Eq. (7), which we equally obtained in the small-S as well as in the large-S regime, is just of the required form, with $X \propto J^{ij}Z_{ij}$ playing the role of the goldstino field. That is why, after supplementing our model with a shift symmetry in the direction of the inflaton field S, all necessary conditions for the successful implementation of chaotic inflation into supergravity are satisfied.

We also note that our scenario is free of (pseudo-) moduli. One candidate for a modulus is X, which has a flat potential at $\lambda X \gg \Lambda$ and $S \gg M_{\mathrm{Pl}}$ in global supersymmetry. In supergravity however, the canonical Kähler potential for X yields a Hubble induced mass term, stabilizing X at its origin during inflation. Note that this argument fails if the Kähler potential contains a term linear in X. But, assigning a non-zero R charge to X, we are fortunately able to prevent such a term, cf. footnote 8.

4. Chaotic inflation with a fractional power-law potential

As we have shown, simple strongly coupled gauge dynamics are able to generate an inflationary potential featuring a fractional power. We have also discussed the shift symmetry of the model which suppresses the distortions of the inflaton potential due to the supergravity contributions. Let us now outline the phenomenology of the model, summarize its predictions for the inflationary observables encoded in the CMB power spectrum and, in relation to that, discuss its testability.

Inflation starts out at an arbitrary initial value of the inflaton field above the Planck mass $S \gg M_{\mathrm{Pl}}$. At its early stages, i.e. as long as $\lambda_S S \gg \Lambda$, the $SP(N)$ gauge interactions are in the perturbative regime and inflation is characterized by the slow-roll motion of the inflaton in the effective potential in Eq. (12). Similarly, we know that at small field values, i.e. when $\lambda_S S \ll \Lambda$, the system is in the s-confinement phase, in which S and the composite mesons have masses of $\mathcal{O}(\lambda_{IJ}\Lambda)$. In this case the inflaton potential is given by Eq. (8). In the intermediate regime, where $\lambda_S S \simeq \Lambda$, we however lack the ability to precisely calculate the inflaton potential, which is why we do not exactly know how the transition

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[7] Even for $S < M_{\mathrm{Pl}}$, there is an eta problem. This problem is also avoided by the solution mentioned below.

[8] One may assume R-symmetry with the following charge assignments: $Q(0)$, $P(1)$, $\bar{P}(1)$, $Z(2)$, $T(1)$, $\bar{T}(1)$, and $S(0)$. With these charge assignments, R-symmetry allows the shift symmetry only for the singlet S, which explains why only the imaginary part of S can exceed the Planck scale.

[9] Here, we have assumed that the Kähler potential in Eq. (17) with the shift symmetry is defined around the Planck scale.

[10] A second reason why we have to assume λ_S to be small is the fact that the masses of the P and $\bar{P}$ quarks, $m_P = m_{\bar{P}} \simeq \lambda_S S$, must be smaller than the Planck scale at all times, even though S might take huge values, $S \sim 10 \cdots 100\, M_{\mathrm{Pl}}$, during inflation. That is, too large λ_S would entail $m_P, m_{\bar{P}} \gtrsim M_{\mathrm{Pl}}$ at some point during inflation, causing our approach based on ordinary quantum field theory to break down.

from the large-S to the small-S regime takes place. For instance, it might be that towards the end of inflation S becomes trapped in a metastable vacuum at a field value around Λ/λ_S such that it actually never reaches the small-S regime. Assuming that the effective inflaton potential exhibits no such peculiar features around Λ/λ_S, we are led to the conclusion that inflation continues without any hindrance until the slow-roll conditions become violated at small values of the inflaton field.

The end of inflation marks the onset of preheating, which proceeds in a rather unconventional way in our scenario due to the negative curvature of the inflaton potential. In fact, so far only small-field and hybrid models of inflation have been studied in connection with a negatively curved scalar potential, where it was found that preheating occurs via *tachyonic oscillations* [20] of the inflaton field or *tachyonic preheating* [21], respectively. As for large-field, i.e. chaotic inflation only the case of a positively curved inflaton potential, for which preheating occurs via *parametric resonance* [22], has been considered up to now. We presume that in our large-field model featuring a negatively curved inflaton potential preheating ends up being a combination of both, tachyonic inflaton oscillations as well as parametric resonance. The verification of this conjecture certainly requires a more comprehensive and ultimately numerical study.

After inflation the inflaton decays through its coupling to the Higgs fields H_u and H_d of the supersymmetric standard model in the Kähler potential, $K \supset (S + S^\dagger)H_u H_d$, at a rate

$$\Gamma_S \sim \frac{M_S^3}{M_{\rm Pl}^2}, \tag{20}$$

with $M_S \simeq \lambda_S \Lambda$ denoting the inflaton mass.[11] Interactions with standard model gauge multiplets through gauge kinetic functions [19] contribute to the S decay rate at the same order. Hence, the inflaton eventually reaches the supersymmetric vacuum, in which $S = \mathcal{M} = 0$. The rate of the perturbative inflaton decays directly determines the reheating temperature, $T_{\rm R} \sim \sqrt{\Gamma_S M_{\rm Pl}}$, or, using Eq. (20),

$$T_{\rm R} \sim 10^7 \text{ GeV} \left(\frac{\lambda_S}{10^{-4}}\right)^{3/2} \left(\frac{\Lambda}{10^{15} \text{ GeV}}\right)^{3/2}. \tag{21}$$

It is interesting to note that non-perturbative effects, i.e. the formation and evaporation of Q-balls near the end of inflation [23], could speed up the decay of the inflaton field, thus leading to a reheating temperature much higher than in the mere perturbative picture. In principle, the formation of Q-balls is feasible in our model since our effective inflaton potential is shallower than a quadratic one. Nonetheless, we suppose that no Q-balls emerge towards the end of inflation because, owing to its Hubble induced mass term, the real part of S, $\Re(S)$, is stabilized at zero. This presumably renders it impossible to induce inspiraling orbits in S field space continuously connected to the inflationary trajectory, which would be a necessary prerequisite for Q-balls to occur [24]. A further study of this issue is beyond the scope of this Letter and shall be carried out elsewhere.

Finally, let us summarize the implications of our fractional power-law inflaton potential for the CMB observables and discuss the testability of our model. Given a potential $V(S) \propto \Lambda^4(|S|/\Lambda)^p$, with $p = 2/(N + 1)$, one finds for the power spectrum $\mathcal{P}_\zeta$ of the curvature perturbations ζ [3]

$$\mathcal{P}_\zeta = \frac{1}{12\pi^2 p^2} \left(\frac{\Lambda}{M_{\rm Pl}}\right)^{4-p} (2p N_e)^{1+p/2}, \tag{22}$$

where N_e is the number of e-foldings. The observational result $\mathcal{P}_\zeta = 2.42 \times 10^{-9}$ [6] then requires the dynamical scale Λ to be shortly below the GUT scale,

$$\Lambda \simeq 10^{15} \text{ GeV}. \tag{23}$$

The spectral index n_s and the tensor-to-scalar ratio r of the curvature perturbations are respectively given by

$$n_s = 1 - \frac{p+2}{2N_e}, \qquad r = \frac{4p}{N_e}. \tag{24}$$

For $N_e = 50$ and $N \geqslant 1$, such that $0 < p \leqslant 1$, we obtain $n_s = 0.97 \cdots 0.98$ and $r = 0.16/(N + 1)$, which is consistent with the recent CMB observations [6]. The Planck experiment is expected to detect the presence of tensor modes if $r > 0.05$ [7], that is, if the inflaton potential is generated by the strong dynamics of an $SP(1) \cong SU(2)$ gauge theory. Future experiment such as CMBPol [25] and LiteBIRD [26], which are expected to reach sensitivities to r of $\mathcal{O}(10^{-3})$, will detect tensor modes unless the underlying gauge group is very large, $N > \mathcal{O}(100)$.

5. Conclusions and discussion

In this Letter, we have shown that fractional power-law chaotic inflation models can be realized in simple supersymmetric gauge theories. In this class of models, the energy scale during inflation is dynamically generated by the dimensional transmutation due to the strong gauge dynamics. Therefore, the models not only explain the origin of the fractional power in the inflationary potential but also provide a reason why the energy scale of inflation is much smaller than the Planck scale. We also discussed how well the model fits together with the current data on the inflationary observables.

Several comments are in order. In our analysis, we have confined ourselves to models with $2(N + 2)$ fundamental representations. One of the reasons for this assumption is that for models with more fundamental representations the system is in the so-called conformal window [27], and hence, exhibits conformal symmetry after inflation. In such cases, the inflaton becomes an unparticle [28] after inflation, which may change the evolution of the universe drastically compared to conventional inflaton scenarios. Although such a possibility is intriguing, we do not pursue it further since it is beyond the scope of this Letter.

In our scenario, we have considered models in which supersymmetry breaking in the inflaton sector vanishes after inflation. It is, however, possible that some portion of supersymmetry breaking in the inflaton sector remains non-vanishing even after inflation, providing the dominant source of supersymmetry breaking in the true vacuum. In this case, we can consolidate two well motivated new physics, supersymmetry breaking and inflation into one model. We will discuss this possibility elsewhere.

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[11] The same coupling leads to the non-adiabatic production of radiation during preheating. As it is strongly suppressed, we assume that during preheating most of the initial vacuum energy is transferred into non-relativistic inflaton particles and only a small fraction into radiation. This implies in particular that the standard definition of the reheating temperature is applicable.

References

[1] A.H. Guth, Phys. Rev. D 23 (1981) 347;
A.D. Linde, Phys. Lett. B 108 (1982) 389;
A. Albrecht, P.J. Steinhardt, Phys. Rev. Lett. 48 (1982) 1220.

[2] V.F. Mukhanov, G.V. Chibisov, JETP Lett. 33 (1981) 532, Pisma Zh. Eksp. Teor. Fiz. 33 (1981) 549;
A.D. Linde, Contemp. Concepts Phys. 5 (1990) 1, hep-th/0503203.

[3] D.H. Lyth, A. Riotto, Phys. Rept. 314 (1999) 1, hep-ph/9807278.

[4] A.D. Linde, Phys. Lett. B 129 (1983) 177;
A.D. Linde, JETP Lett. 38 (1983) 176, Pisma Zh. Eksp. Teor. Fiz. 38 (1983) 149.

[5] A.A. Starobinsky, Sov. Astron. Lett. 11 (1985) 133;
D.H. Lyth, Phys. Rev. Lett. 78 (1997) 1861, hep-ph/9606387.

[6] E. Komatsu, et al., WMAP Collaboration, Astrophys. J. Suppl. 192 (2011) 18, arXiv:1001.4538 [astro-ph.CO].

[7] Planck Collaboration, astro-ph/0604069.

[8] L. Alabidi, I. Huston, JCAP 1008 (2010) 037, arXiv:1004.4794 [astro-ph.CO];
J. Martin, C. Ringeval, R. Trotta, Phys. Rev. D 83 (2011) 063524, arXiv:1009.4157 [astro-ph.CO].

[9] R. Kallosh, A. Linde, JCAP 1011 (2010) 011, arXiv:1008.3375 [hep-th].

[10] R. Kallosh, A. Linde, T. Rube, Phys. Rev. D 83 (2011) 043507, arXiv:1011.5945 [hep-th].

[11] F. Takahashi, Phys. Lett. B 693 (2010) 140, arXiv:1006.2801 [hep-ph].

[12] E. Silverstein, A. Westphal, Phys. Rev. D 78 (2008) 106003, arXiv:0803.3085 [hep-th].

[13] S. Dimopoulos, G.R. Dvali, R. Rattazzi, Phys. Lett. B 410 (1997) 119, hep-ph/9705348;
K.I. Izawa, M. Kawasaki, T. Yanagida, Phys. Lett. B 411 (1997) 249, hep-ph/9707201;
K.I. Izawa, Prog. Theor. Phys. 99 (1998) 157, hep-ph/9708315.

[14] N. Seiberg, Phys. Rev. D 49 (1994) 6857, hep-th/9402044.

[15] K.A. Intriligator, P. Pouliot, Phys. Lett. B 353 (1995) 471, hep-th/9505006;
C. Csaki, M. Schmaltz, W. Skiba, Phys. Rev. Lett. 78 (1997) 799, hep-th/9610139.

[16] K.-I. Izawa, T. Yanagida, Prog. Theor. Phys. 95 (1996) 829, hep-th/9602180;
K.A. Intriligator, S.D. Thomas, Nucl. Phys. B 473 (1996) 121, hep-th/9603158.

[17] M.A. Shifman, A.I. Vainshtein, Nucl. Phys. B 277 (1986) 456, Sov. Phys. JETP 64 (1986) 428, Zh. Eksp. Teor. Fiz. 91 (1986) 723;

G. Veneziano, S. Yankielowicz, Phys. Lett. B 113 (1982) 231;
T.R. Taylor, G. Veneziano, S. Yankielowicz, Nucl. Phys. B 218 (1983) 493;
V.A. Novikov, M.A. Shifman, A.I. Vainshtein, V.I. Zakharov, Nucl. Phys. B 229 (1983) 381;
N. Seiberg, Phys. Lett. B 318 (1993) 469, hep-ph/9309335.

[18] M. Kawasaki, M. Yamaguchi, T. Yanagida, Phys. Rev. Lett. 85 (2000) 3572, hep-ph/0004243;
M. Kawasaki, M. Yamaguchi, T. Yanagida, Phys. Rev. D 63 (2001) 103514, hep-ph/0011104.

[19] R. Kallosh, A. Linde, K.A. Olive, T. Rube, Phys. Rev. D 84 (2011) 083519, arXiv:1106.6025 [hep-th];
A. Linde, S. Mooij, E. Pajer, arXiv:1212.1693 [hep-th].

[20] P. Brax, J.-F. Dufaux, S. Mariadassou, Phys. Rev. D 83 (2011) 103510, arXiv:1012.4656 [hep-th].

[21] G.N. Felder, J. Garcia-Bellido, P.B. Greene, L. Kofman, A.D. Linde, I. Tkachev, Phys. Rev. Lett. 87 (2001) 011601, hep-ph/0012142;
G.N. Felder, L. Kofman, A.D. Linde, Phys. Rev. D 64 (2001) 123517, hep-th/0106179;
M. Desroche, G.N. Felder, J.M. Kratochvil, A.D. Linde, Phys. Rev. D 71 (2005) 103516, hep-th/0501080.

[22] L. Kofman, A.D. Linde, A.A. Starobinsky, Phys. Rev. Lett. 73 (1994) 3195, hep-th/9405187;
L. Kofman, A.D. Linde, A.A. Starobinsky, Phys. Rev. D 56 (1997) 3258, hep-ph/9704452.

[23] S.R. Coleman, Nucl. Phys. B 262 (1985) 263;
S.R. Coleman, Nucl. Phys. B 269 (1986) 744 (Erratum);
A.G. Cohen, S.R. Coleman, H. Georgi, A. Manohar, Nucl. Phys. B 272 (1986) 301;
K. Enqvist, S. Kasuya, A. Mazumdar, Phys. Rev. D 66 (2002) 043505, hep-ph/0206272.

[24] S. Kasuya, M. Kawasaki, Phys. Rev. D 61 (2000) 041301, hep-ph/9909509;
A. Jokinen, hep-ph/0204086.

[25] D. Baumann, et al., CMBPol Study Team Collaboration, AIP Conf. Proc. 1141 (2009) 10, arXiv:0811.3919 [astro-ph].

[26] LiteBIRD project, http://cmb.kek.jp/litebird/index.html.

[27] T. Banks, A. Zaks, Nucl. Phys. B 196 (1982) 189.

[28] H. Georgi, Phys. Rev. Lett. 98 (2007) 221601, hep-ph/0703260.