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Linear inflation from running kinetic term in supergravity

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ABSTRACT

We propose a new class of inflation models in which the coefficient of the inflaton kinetic term rapidly changes with energy scale. This naturally occurs especially if the inflaton moves over a long distance during inflation as in the case of large-scale inflation. The peculiar behavior of the kinetic term opens up a new way to construct an inflation model. As a concrete example we construct a linear inflation model in supergravity. It is straightforward to build a chaotic inflation model with a fractional power along the same line. Interestingly, the potential takes a different form after inflation because of the running kinetic term.

The inflation has been strongly motivated by the observation [1], while it is a non-trivial task to construct a successful inflation model. A successful inflaton model must explain several features of the density perturbation, but properties of the inflaton are not well understood. It is often assumed that, in the slow-roll inflation paradigm, the inflaton is a weakly coupled field, and therefore the kinetic term is simply set to be the canonical form during inflation. This seems justified because the typical energy scale of inflation is given by the Hubble parameter, which remains almost constant during inflation. However, there is another important energy scale, namely, the inflaton field. Even in the slow-roll inflation, the motion of the inflaton is not negligible and it may travel a long distance during the whole period of inflation. In particular, in the case of large-scale inflation such as a chaotic inflation model [2], the inflaton typically moves over a Planck scale or even larger within the last 60 e-foldings [3]. Then, it seems quite generic that the precise form of the kinetic term changes during the course of inflation. In some cases, the change could be so rapid, that it significantly affects the inflaton dynamics. In this Letter, we construct a model in which the coefficient of the kinetic term grows rapidly with the inflaton field value, but in a controlled way. By doing so, we construct a linear term inflation model in the supergravity framework (see Ref. [4] for the quadratic model). The realization of the linear term inflation model in the string theory was given in Ref. [5]. We also show that a chaotic inflation model with a fractional power can be straightforwardly constructed along the same line.

Before going to a realistic inflation model, let us give our basic idea. Suppose that the inflaton field ϕ has the following kinetic term,

$$\mathcal{L}_K = \frac{1}{2} f(\phi) \partial^\mu \phi \partial_\mu \phi, \tag{1}$$

and that the inflaton field is canonically normalized at the potential minimum:

$$f(\phi_{\min}) = 1. \tag{2}$$

However, this does not necessarily mean that $f(\phi)$ remains close to 1 during inflation, especially if the inflaton moves over some high scale, e.g., the GUT or Planck scale. Suppose that the behavior of $f(\phi)$ can be approximated by $f(\phi) \approx n^2 \phi^{2n-2}$ with an integer n over a certain range of ϕ. Then, when expressed in terms of the canonically normalized inflaton field, $\chi \equiv \phi^n$, the scalar potential $V(\phi)$ is modified to be

$$V(\phi) \to V\left(\chi^{1/n}\right). \tag{3}$$

For instance, if $n = 2$, the quadratic potential, $V(\phi) \propto \phi^2$, becomes a linear term $V(\chi) \propto \chi$. Therefore, such a strong dependence of the kinetic term on the inflaton field changes the inflation dynamics significantly. In particular, the large coefficient of the kinetic term is advantageous for inflation to occur, since the effective potential becomes flatter.

Now let us construct a linear term inflation model in supergravity. In this inflation model, the inflaton field has a scalar potential linearly proportional to the inflaton, $V(\phi) \propto \phi$. For the inflation to last for the 60 e-foldings, the inflaton field must take a value greater than the Planck scale, which is difficult to implement in supergravity because of the exponential prefactor in the scalar potential. Therefore we need to introduce some sort of shift symmetry, which suppresses the exponential growth of the potential.

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We introduce a chiral superfield, ϕ, and require that the Kähler potential for ϕ is invariant under the following transformation;

$$\phi^2 \to \phi^2 + \alpha \quad \text{for } \alpha \in \mathbf{R} \text{ and } \phi \neq 0, \tag{4}$$

which means that a composite field $\chi \sim \phi^2$ transforms under a Nambu–Goldstone like shift symmetry. This is equivalent to imposing a hyperbolic rotation symmetry (or equivalently $SO(1,1)$) on (ϕ_R, ϕ_I), where ϕ_R and ϕ_I are the real and imaginary components, $\phi = (\phi_R + i\phi_I)/\sqrt{2}$.

The Kähler potential must be a function of $(\phi^2 - \phi^{\dagger 2})$:

$$K = ic\left(\phi^2 - \phi^{\dagger 2}\right) - \frac{1}{4}\left(\phi^2 - \phi^{\dagger 2}\right)^2 + \cdots, \tag{5}$$

where c is a real parameter of $O(1)$ and the Planck unit is adopted. Note that the $|\phi|^2$ term, which usually generates the kinetic term for ϕ, is forbidden by the symmetry. Instead, the kinetic term arises from the second term, and the coefficient of the kinetic term will be proportional to $|\phi|^2$. Note that the lowest component of the Kähler potential vanishes for either $\phi_R = 0$ or $\phi_I = 0$. This feature is essential for constructing a chaotic inflation model in supergravity.

Let us add a symmetry breaking term, $\Delta K = \kappa|\phi|^2$, to cure the singular behavior of the Kähler metric at the origin. Here $\kappa \ll 1$ is a real numerical coefficient, and the smallness is natural in the 't Hooft's sense [6]. There could be other symmetry breaking terms, but, throughout this Letter we assume that those symmetry breaking terms are *soft* in a sense that the shift symmetry remains a good symmetry at large enough ϕ. The kinetic term of the scalar field then becomes

$$\mathcal{L}_K = \left(\kappa + 2|\phi|^2 + \cdots\right)\partial^\mu \phi^\dagger \partial_\mu \phi, \tag{6}$$

where the higher-order terms expressed by the dots contain terms proportional to $(\phi^2 - \phi^{\dagger 2})$. Let us drop the higher-order terms for the moment. As demonstrated later, the higher-order terms do not change the form of the kinetic term. For a large field value $|\phi| \gg \sqrt{\kappa}$, the coefficient of the kinetic term grows with the field value, which makes the potential flatter. The canonically normalized field is $\chi = \phi^2/\sqrt{2}$, as expected. In a sense, χ is a more suitable dynamical variable to describe the system satisfying the shift symmetry (4). On the other hand, for a small field value of $|\phi| \ll \sqrt{\kappa}$, the canonically normalized field is $\sqrt{\kappa}\phi$. Thus, ϕ^2 and ϕ are the dynamical variables for high and low scales, respectively.

We can interpret the above phenomenon in the following way. If we go to high energy scales, namely the large field value ϕ, the self-interaction in the Kähler potential becomes strong, and the scalar field forms a bound state ϕ^2. On the other hand, as ϕ becomes small, the self-coupling becomes smaller and the symmetry-breaking term becomes more relevant. Thus ϕ^2 breaks up and ϕ becomes the suitable variable. Such a phenomenon of forming a bound state seems quite generic if one considers a large-scale inflation in which the inflaton takes a very large field value during inflation. Because of the strong self-interactions, the inflaton kinetic term runs with scales, and the inflaton dynamics is significantly changed. The novelty here is the existence of the shift symmetry, without which we cannot control the effect of the higher order terms on the inflationary dynamics. We will come back to this point later.

In order to construct a realistic inflation model, we consider the following Kähler and super-potentials[1]

$$K = \kappa|\phi|^2 + ic\left(\phi^2 - \phi^{\dagger 2}\right) - \frac{1}{4}\left(\phi^2 - \phi^{\dagger 2}\right)^2 + |X|^2, \tag{7}$$

$$W = mX\phi, \tag{8}$$

where both κ and m break the symmetry, and so we assume $\kappa \ll 1$ and $m \ll 1$.[2] These small parameters are naturally understood in 't Hooft's sense. The superpotential produces the inflaton potential. We assume that X and ϕ have $U(1)_R$ charges 2 and 0, respectively. We assign a Z_2 symmetry under which both X and ϕ flip the sign.

The Lagrangian is given by

$$\mathcal{L} = \left(\kappa + 2|\phi|^2\right)\partial_\mu \phi^\dagger \partial^\mu \phi + \partial_\mu X^\dagger \partial^\mu X - V \tag{9}$$

with

$$V = e^K\left(|D_X W|^2 K^{X\bar{X}} + |D_\phi W|^2 K^{\phi\bar{\phi}} - 3|W|^2\right). \tag{10}$$

The scalar potential looks complicated, but it can be reduced to a simple form during inflation. One can show that, during inflation, X acquires a mass of the order of the Hubble scale and is stabilized at the origin, if $|c| \gtrsim 1$.[3] Then the scalar potential becomes

$$V \approx \frac{1}{2} e^{\frac{\kappa}{2}(\phi_R^2 + \phi_I^2) - 2c\phi_R\phi_I + \phi_R^2\phi_I^2} m^2\left(\phi_R^2 + \phi_I^2\right). \tag{11}$$

The flat direction is given by $\phi_R\phi_I = \text{constant}$ because of the symmetry (4). Therefore if ϕ_R has a very large value, ϕ_I is stabilized at a point where the Kähler potential is minimal. For $\phi_R > 1$ and $\kappa \ll 1$, ϕ_I is stabilized at

$$\phi_I \approx \frac{c}{\phi_R}. \tag{12}$$

Here and in what follows we focus on the case of $\phi_R > 0$ and $\phi_I > 0$ without loss of generality. The scalar potential is then reduced to the following form:

$$V \approx \frac{1}{2} e^{\frac{\kappa}{2}(\phi_R^2 + \frac{c^2}{\phi_R^2}) - c^2} m^2\left(\phi_R^2 + \frac{c^2}{\phi_R^2}\right), \tag{13}$$

for $|\phi_R| > 1$. Since we explicitly breaks the shift symmetry (4) by the κ term, there appears a non-vanishing exponential prefactor. However, for $|\phi_R| < 1/\sqrt{\kappa}$, the exponential prefactor is of $O(1)$, and therefore can be dropped.[4] The Lagrangian for the inflaton ϕ_R is summarized by

$$\mathcal{L} \approx \frac{1}{2}\phi_R^2(\partial\phi_R)^2 - \frac{1}{2}m^2\phi_R^2, \tag{14}$$

for $1 \ll \phi_R \ll 1/\sqrt{\kappa}$. In terms of the canonically normalized field, $\varphi \equiv \phi_R^2/2$, we have

$$\mathcal{L} \approx \frac{1}{2}(\partial\varphi)^2 - m^2\varphi, \tag{15}$$

for $1 \ll \varphi \ll 1/\kappa$. Thus our model is equivalent to the linear term inflation model.

After inflation ends, the inflaton field will oscillate about the origin. Then the ϕ_I is no longer negligible, and the inflaton is expressed by a complex scalar field. As the amplitude decreases, the κ-term becomes more important, and in the end, the kinetic term arises mainly from the κ-term. The Lagrangian is

[1] One can add an additional breaking term in the superpotential which induces a periodic potential for φ [5]. The interesting feature may be found in the non-Gaussianity in this case [7].

[2] The breaking of the shift symmetry in the superpotential could produce radiative corrections to the Kähler potential. In particular, $\kappa = O(m^2)$ is induced [4]. Here we consider a more general case that κ and m are not related to each other.

[3] A self-coupling $\sim |X|^4$ in the Kähler potential produces a Hubble-induced mass term for X about the origin.

[4] Actually, the inflaton does slow-roll if the exponential prefactor gives a main contribution to the tilt of the potential.