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Supersymmetric DBI inflation

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ABSTRACT

We discuss a supersymmetric version of DBI (Dirac–Born–Infeld) inflation, which is a typical inflation model in string cosmology. The supersymmetric DBI action together with a superpotential always leads to correction terms associated with the potential into the kinetic term, which drastically change the dynamics of DBI inflation. We find two significant features of supersymmetric DBI inflation. The first one is that ultra-relativistic motion is prohibited to cause inflation, which leads to order of unity sound velocity squared and hence small non-Gaussianities of primordial curvature perturbations. The second one is that the relation between the tensor-to-scalar ratio and the field variation is modified. Then, significant tensor-to-scalar ratio $r \gtrsim 0.01$ is possible because the variation of the canonically normalized inflaton can be beyond the reduced Planck scale. These new features are in sharp contrast with those of the standard non-supersymmetric DBI inflation and hence have a lot of interest implications on upcoming observations of cosmic microwave background (CMB) anisotropies by the Planck satellite as well as direct detection experiments of gravitational waves like DECIGO and BBO.

Recent observations of CMB anisotropies like the Wilkinson Microwave Anisotropy Probe (WMAP) satellite strongly suggest the presence of accelerated expansion called inflation in the early Universe [1]. During inflation, primordial curvature [2] and tensor [3] perturbations are generated and stretched out to the cosmological scales, which become seeds of the large scale structure formation and the CMB anisotropies. The properties of primordial curvature fluctuations are well known and are almost scale-invariant, adiabatic, and Gaussian. Though tensor perturbations have not yet been found unfortunately, they are expected to be detectable in upcoming CMB experiments like the Planck [4] and the CMBPol [5]. However, we do not know the origin of the inflaton at all, except that it is an effective scalar field. (See Refs. [6] for recent review of inflation model building.)

String theory is the most powerful candidate to unify all of the fundamental interactions. Then, it is natural to pursue the candidate of an inflaton in string theory. In fact, inflation models in the brane setting were proposed [7,8] and have been investigated intensively. Among them, a particularly interesting class of inflation models is DBI inflation [9], which is associated with the relativistic motion of a D-brane in the warped flux compactification. This model has distinctive predictions for primordial perturbations: (i) it can naturally generate large non-Gaussianities of primordial curvature perturbations thanks to the ultra-relativistic motion [9], (ii) it is quite difficult to produce detectable tensor perturbations because the maximal field variation of the inflaton is constrained to be less than the reduced Planck scale $M_{\rm pl}$ [10]. Current observations like the WMAP satellite are precise enough to rule out simple UV models of DBI inflation [11] though more elaborated models are still compatible with the present observations [12,13].

Almost all of these studies of DBI inflation, however, have been based on the non-supersymmetric setup. Supersymmetry is one of the most promising solutions to the hierarchy problem of the Standard Model as well as the unification of the fundamental interactions. Once a probe D-brane is placed on supersymmetric backgrounds, one expects that the world-volume effective theory of the probe brane becomes supersymmetric. Therefore, it is quite important to consider DBI inflation in the supersymmetric framework. Recently, some attempts to supersymmetrize non-canonical kinetic terms have been done [14,15]. However, in order to incorporate a potential term, one needs to introduce a superpotential and solve the equation of motion for the auxiliary field consistently. This is a difficult task when the non-canonical kinetic terms are present.

In this Letter, we discuss the supersymmetric version of DBI inflation. First of all, the supersymmetric DBI action with the superpotential is studied. By solving the equation of motion for the auxiliary field consistently, we show that correction terms associated with the potential always appear in the kinetic term, which drastically changes the dynamics of DBI inflation. Then, using the newly

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obtained action, we investigate the dynamics of supersymmetric DBI inflation in detail. It is revealed that the predictions of primordial perturbations are completely different from those of non-supersymmetric DBI inflation, which may require us to reanalyze all of the DBI inflation models including elaborated models in preparation for the upcoming experiments.

Now let us begin with the supersymmetric DBI action in the warped throat. Assuming that a probe D3-brane is moving in the supersymmetric ten-dimensional geometry of the form

$$ds_{10}^2 = H^{-\frac{1}{2}}(y)\,ds_4^2 + H^{\frac{1}{2}}(y)\,ds_6^2, \tag{1}$$

where ds_4^2, ds_6^2 are four-dimensional spacetime and six-dimensional internal space respectively, the supersymmetric DBI Lagrangian in the flat spacetime [16] is generalized as follows:

$$\mathcal{L}_{\rm DBI} = \int d^4\theta \left[\Phi\Phi^\dagger + \frac{1}{16T}\left(D^\alpha\Phi D_\alpha\Phi\right)\left(\overline{D}^{\dot\alpha}\Phi^\dagger\overline{D}_{\dot\alpha}\Phi^\dagger\right) \times \frac{1}{1+A+\sqrt{(1+A)^2-B}}\right], \tag{2}$$

where we have employed the static gauge and normalized the D3-brane tension and the string slope parameter $2\pi\alpha'$ to unity. The chiral and anti-chiral superfields are denoted by Φ and $\Phi^\dagger$, D_α and $\overline{D}_{\dot\alpha}$ are the supercovariant derivatives, $T = T(\Phi, \Phi^\dagger)$ is a function of Φ, $\Phi^\dagger$ corresponding to the warp factor $T = H^{-1}$, and A, B are given by

$$A \equiv \frac{\partial_\mu\Phi\partial^\mu\Phi^\dagger}{T}, \qquad B \equiv \frac{\partial_\mu\Phi\partial^\mu\Phi\partial_\nu\Phi^\dagger\partial^\nu\Phi^\dagger}{T^2}. \tag{3}$$

Here we have turned on one complex scalar field associated with two independent fluctuations along the throat direction y. In order to incorporate the potential, we add the superpotential term to the Lagrangian,

$$\mathcal{L}_{\rm pot} = \int d^2\theta\, W(\Phi) + \text{h.c.} \tag{4}$$

Several kinds of superpotentials are induced by the background fluxes. For example, we can introduce the superpotential of a mass term $W = \frac{1}{2}m\Phi^2$ on the D3-brane in the presence of a constant Ramond–Ramond 3-form background [17,18].

The component Lagrangian is

$$\mathcal{L} = -T\big(1 + 2T^{-1}\partial_\mu\varphi\partial^\mu\overline{\varphi} + T^{-2}\left(\partial_\mu\varphi\partial^\mu\overline{\varphi}\right)^2 - T^{-2}\left(\partial_\mu\varphi\partial^\mu\varphi\right)\left(\partial_\nu\overline{\varphi}\partial^\nu\overline{\varphi}\right)\big)^{1/2} + T + \overline{F}F + \frac{\partial W}{\partial\varphi}F + \frac{\partial\overline{W}}{\partial\overline{\varphi}}\overline{F} + G(\varphi)\left(-2F\overline{F}\partial_\mu\varphi\partial^\mu\overline{\varphi} + F^2\overline{F}^2\right), \tag{5}$$

where we have dropped the fermions since they do not contribute to the dynamics of the inflation. The function $G(\varphi)$ is defined as

$$G(\varphi) = \frac{1}{T}\frac{1}{1+A+\sqrt{(1+A)^2-B}}, \tag{6}$$

with the replacement of Φ by φ in A and B.

The Lagrangian for the scalar component φ in the chiral superfield Φ is obtained by solving the equation of motion for the auxiliary field F in Φ. This is in general a simultaneous equation for F and $\bar{F}$. After eliminating $\bar{F}$, we find the equation for F is given by

$$2G(\varphi)\frac{\partial W}{\partial\varphi}F^3 + \frac{\partial\overline{W}}{\partial\overline{\varphi}}\left(1 - 2G(\varphi)\partial_\mu\varphi\partial^\mu\overline{\varphi}\right)F + \left(\frac{\partial\overline{W}}{\partial\overline{\varphi}}\right)^2 = 0. \tag{7}$$

Unlike the standard (quasi-)canonical case, a salient feature of the supersymmetric DBI model is that the equation for F is *cubic* and can be solved analytically by Cardano's method,[1]

$$F = \omega^k\sqrt[3]{-\frac{q}{2} + \sqrt{\left(\frac{q}{2}\right)^2 + \left(\frac{p}{3}\right)^3}} + \omega^{3-k}\sqrt[3]{-\frac{q}{2} - \sqrt{\left(\frac{q}{2}\right)^2 + \left(\frac{p}{3}\right)^3}}. \tag{8}$$

Here ω is the complex cubic root, $k = 0, 1, 2$, and p and q are given by

$$p = \left(\frac{\partial W}{\partial\varphi}\right)^{-1}\frac{\partial\overline{W}}{\partial\overline{\varphi}}\frac{1 - 2G\partial_\mu\varphi\partial^\mu\overline{\varphi}}{2G},$$
$$q = \frac{1}{2G}\left(\frac{\partial W}{\partial\varphi}\right)^{-1}\left(\frac{\partial\overline{W}}{\partial\overline{\varphi}}\right)^2. \tag{9}$$

We note that if $W = 0$, the unique solution is given by $F = 0$ and the bosonic part of the DBI Lagrangian (2) is not changed compared with the non-supersymmetric case. A remarkable fact in the case $W \neq 0$ is that there are three different on-shell actions associated with the $k = 0, 1, 2$ solutions in Eq. (8). In the following, we concentrate on the $k = 0$ branch since it is continuously connected to the ordinary solution $F = -\partial\overline{W}/\partial\bar{\varphi}$ in the canonical limit. The other solutions with $k = 1, 2$ do not have any definite limit and will yield essentially inequivalent theories.

Now we denote the phase factor of $\partial W/\partial\varphi$ as α. Since the functions A, B and G are real, the phase of p and q in the solutions (9) are -2α and -3α respectively. Then from the $k = 0$ solution in (8), the phase of F is given by $\pi - \alpha$. As a result, the phase factor of the product of $\partial W/\partial\varphi$ and F becomes π and does not depend on α in the on-shell Lagrangian. Therefore only the absolute values of $\partial W/\partial\varphi$ and F contribute to the Lagrangian.

We further impose the global $U(1)_R$ symmetry on the superpotential $W(\Phi)$ and the warp factor T. This is always possible when the geometry (1) has a $U(1)$ isometry in the y direction. A typical example of this kind of geometry is the near horizon limit of N coincident D3-branes [9]. Since the supersymmetric DBI Lagrangian given in Eq. (2) is invariant under the $U(1)_R$ symmetry, the dynamics of the scalar field φ depends only on its radial component f. In this case, f is identified with the fluctuation along the radial direction in $AdS_5 \times S^5$.

Under these circumstances, the full on-shell action for the scalar field f in curved spacetime is given by

$$S = \int d^4x\sqrt{-g}\left(\frac{1}{2}M_{\rm pl}^2R + \mathcal{L}_f\right), \tag{10}$$

$$\mathcal{L}_f = \mathcal{L}_{\rm DBI} + \mathcal{L}_{\rm aux}, \tag{11}$$

$$\mathcal{L}_{\rm DBI} = T\frac{\gamma - 1}{\gamma}, \tag{12}$$

$$\mathcal{L}_{\rm aux} = F^2 - 2\sqrt{2}\frac{dW}{df}F + GF^2\left(2X + F^2\right), \tag{13}$$

where F and $\partial W/\partial f$ are real and positive, and

[1] In the case where the fermions are present, a perturbative solution to the equation for F was discussed in [14].

$$X \equiv -\frac{1}{2}\partial_\mu f \partial^\mu f, \qquad \gamma = \frac{1}{\sqrt{1-2\frac{X}{T}}},$$

$$G = \frac{1}{T}\frac{2\gamma^2}{(1+\gamma)^2}. \tag{14}$$

The equation for the auxiliary field can be rewritten as

$$2GF^3 + (1+2GX)F - \sqrt{2}\frac{dW}{df} = 0. \tag{15}$$

Although we have the analytic solutions to the above equation, since its complexity would make it difficult to capture the essence of the physical properties, we look for approximate solutions under the assumption that one term in the left-hand side of Eq. (15) is subdominant. This will provide a valuable intuition for the clear characteristics of our model. Later, we will mention the case that all the terms in Eq. (15) are comparable.

Case (i): subdominance of the first term. The auxiliary field F is given by

$$F \simeq \sqrt{2}\frac{\gamma+1}{3\gamma-1}\frac{dW}{df}. \tag{16}$$

The condition that the first term is negligible is satisfied for

$$\frac{8\gamma^2(\gamma+1)}{(3\gamma-1)^3}\frac{1}{T}\left(\frac{dW}{df}\right)^2 \ll 1. \tag{17}$$

Since the prefactor in the left-hand side of the above inequality is of the order unity for $\gamma \geqslant 1$, $(dW/df)^2 \ll T$. Under this condition, the Lagrangian $\mathcal{L}_f$ for the scalar field f is dominated by the DBI kinetic term $\mathcal{L}_{\rm DBI}$ only except the case $\gamma - 1 \ll 1$ when the DBI kinetic term is significantly suppressed. Hence, only inflation with a usual (almost) canonical kinetic term can happen. In this case, large tensor perturbations are prohibited due to the Lyth bound and the constrained field variation.

Case (ii): subdominance of the last term. The solution is obtained by taking the limit $q \to 0$ in Eq. (8) and is found to be $F = 0$, which leads to no potential and hence no inflation.

Case (iii): subdominance of the middle term. The solution is given by taking the limit $p \to 0$ in Eq. (8). We obtain

$$F \simeq \frac{1}{\sqrt{2}}\left(\frac{1+\gamma}{\gamma}\right)^{\frac{2}{3}}\left(T\frac{dW}{df}\right)^{\frac{1}{3}}. \tag{18}$$

The subdominance of the middle term is satisfied for

$$\frac{8\gamma^2(\gamma+1)}{(3\gamma-1)^3}\frac{1}{T}\left(\frac{dW}{df}\right)^2 \gg 1. \tag{19}$$

Note that this condition is just the opposite inequality of Eq. (17) and equivalent to $(dW/df)^2 \gg T$. Substituting this solution of F in Eq. (13) yields

$$\mathcal{L}_{\rm aux} = -\frac{1}{2^{\frac{2}{3}}}\left(\frac{\gamma+1}{\gamma}\right)^{\frac{2}{3}} V(f), \tag{20}$$

where the potential $V(f)$ is defined so that $\mathcal{L}_{\rm aux} \to -V(f)$ for $\gamma \to 1$ (i.e. no kinetic term limit $X \to 0$),

$$V(f) \equiv \left(\frac{27T}{2}\right)^{\frac{1}{3}}\left(\frac{dW}{df}\right)^{\frac{4}{3}}. \tag{21}$$

The condition (19) can be recast into $V \gg T$. Note that this is not a sufficient condition for inflation because the kinetic terms (γ) depending on the potential appears in $\mathcal{L}_{\rm aux}$. In fact, the slow-roll parameter ϵ is given by

$$\epsilon = -\frac{\dot{H}}{H^2} \simeq \frac{3(\gamma-1)}{2\gamma+1}, \tag{22}$$

where we have used $V \gg T$. Thus, inflation can happen only for $\gamma \simeq 1$, that is, the ultra-relativistic motion of the D-brane is prohibited in the supersymmetric DBI inflation, which is in marked contrast to the standard non-supersymmetric case. For k-inflation type Lagrangian ($\mathcal{L}_f = K(f, X)$) [19] including the DBI inflation as a special case, the non-Gaussianities of the curvature perturbations are enhanced by $1/c_s^2$ [20]. Then, the standard non-supersymmetric DBI inflation predicts large non-Gaussianities for ultra-relativistic motion because of $c_s^2 = 1/\gamma^2$ [9]. On the other hand, in our case, the sound velocity squared are estimated as

$$c_s^2 \simeq 3/\left(3\gamma^2 + \gamma - 1\right) \simeq 1, \tag{23}$$

for $V \gg T$ and $\gamma \simeq 1$. Thus, c_s^2 becomes almost unity, and hence negligible non-Gaussianity is predicted for the supersymmetric DBI inflation. Next, we discuss tensor perturbations and comment on the generalized Lyth bound [10,21]. The field variation of f can be related to the e-folding number N for $\mathcal{L}_f = K(f, X)$ as,

$$\frac{df}{M_{\rm pl}} = \sqrt{\frac{r}{8c_s K_X}}\, dN, \tag{24}$$

where r is the tensor-to-scalar ratio and K_X is the partial derivative of K with respect to X. Here, you should notice that $c_s K_X = 1$ both for the canonical kinetic term ($c_s = K_X = 1$) and for the standard DBI case $c_s^{-1} = \gamma = K_X$, which leads to the so-called Lyth bound, namely, significant tensor-to-scalar ratio $r \gtrsim 0.01$ is possible only for $\Delta f \gtrsim M_{\rm pl}$. However, the relation $c_s K_X = 1$ does not hold true in our case. Instead, the following relation is obtained for $\gamma \sim 1$ and $V \gg T$,

$$c_s K_X \sim \frac{V}{3T} \gg 1. \tag{25}$$

Therefore, the tensor-to-scalar ratio r is enhanced by the factor $c_s K_X$ in comparison to the standard non-supersymmetric DBI inflation, which leads to significant tensor-to-scalar ratio $r \gtrsim 0.01$ even for apparent sub-Planck variation of the field. This can be easily understood by expanding the Lagrangian around $\gamma = 1$ and taking the leading terms for $V \gg T$,

$$\mathcal{L}_f \simeq \frac{V}{3T}X - V. \tag{26}$$

Thus, the kinetic term is enhanced by $V/3T$. If we take the canonical kinetic term by redefining the field f as $f_{\rm can} \sim f\sqrt{V/(3T)}$, the Lyth bound applies for $f_{\rm can}$.[2] Therefore, the observable tensor perturbations are predicted because the variation of the canonically normalized inflaton $f_{\rm can}$ can be beyond the reduced Planck scale. Finally, we would like to mention the case that all of the terms in the left-hand side of Eq. (15) are comparable. Under this condition, $V \sim (dW/df)^2$ and T are comparable. Then, by comparing the kinetic part (12) and the auxiliary part (13) in the Lagrangian, it is easy to verify that inflation is possible only for $\gamma - 1 \ll 1$ in this case as well.

In summary, we have discussed the supersymmetric DBI inflation. In order to accommodate the potential term in addition to the

[2] Even in our case, Planck-suppressed operators for the canonically normalized field $f_{\rm can}$ must be controlled to guarantee large tensor perturbations [22]. One of such methods is to introduce (approximate) shift symmetry [23]. It is manifest from the chiralities of Φ and $\Phi^\dagger$ that our DBI action given in Eq. (2) can be easily modified to respect it approximately. However, it should be notice that we have to abandon a global $U(1)_R$ symmetry in this case, though the analysis runs almost parallel.

DBI kinetic term consistently, the equation of motion for the auxiliary field F is derived and solved. Inserting its solution into the Lagrangian, we obtain the effective Lagrangian for the supersymmetric DBI inflation, in which the kinetic term related to the potential always appears in addition to the DBI kinetic term. We find that ultra-relativistic motion of the D-brane is forbidden to cause inflation, which has the significant implications on the prediction of the primordial perturbations. Firstly, the non-Gaussianities of the primordial curvature perturbations are negligible because the sound velocity squared are almost unity. Second, the significant tensor-to-scalar ratio is possible in our model, especially in Case (iii), because of the enhancement of the kinetic term. These two features are totally different from those of the standard non-supersymmetric DBI inflation. Provided that our model be realized, upcoming observations such as Planck and CMBPol experiments will detect such tensor perturbations though the non-Gaussianities of the curvature perturbations will, unfortunately, not be observed.

These new predictions are based on the fact that one always encounters kinetic (derivative) terms accompanied by the potential in supersymmetric models with non-canonical kinetic terms. This feature is not confined to DBI inflation but quite generic to non-trivial kinetic terms appearing in inflation models such as k-inflation [19] and G-inflation [24], which must also be supersymmetrized once supersymmetry would be found as fundamental symmetry. For example, similar structures, such as a cubic equation of the auxiliary field and potential-induced kinetic terms, appear in the k-inflation models with superpotentials. We will discuss elsewhere the supersymmetrization of these models and its implications for cosmology by solving the equation for an auxiliary field adequately.

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