Primordial Fluctuations and Non-Gaussianities in Multifield Dirac-Born-Infeld Inflation

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We study Dirac-Born-Infeld inflation models with multiple scalar fields. We show that the adiabatic and entropy modes propagate with a common effective sound speed and are thus amplified at the sound horizon crossing. In the small sound speed limit, we find that the amplitude of the entropy modes is much higher than that of the adiabatic modes. We show that this could strongly affect the observable curvature power spectrum as well as the amplitude of non-Gaussianities, although their shape remains as in the single-field Dirac-Born-Infeld case.

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The past decade has seen an accumulation of cosmological data of increasing precision. Together with future experiments planned to measure the cosmic microwave background (CMB) fluctuations with yet further accuracy, we may be able to piece together more clues about early Universe physics. In parallel with this observational effort, there has been tremendous progress in recent years in the construction of early Universe models in the framework of high energy physics and string theory.

A particularly interesting class of models based on string theory is known as Dirac-Born-Infeld (DBI) inflation [1,2], associated with the motion of a D3-brane in a higher-dimensional background spacetime. The characteristic of DBI inflation, and that which gives it its name, is that the action is of the Dirac-Born-Infeld type and thus contains nontrivial kinetic terms. Most studies of DBI inflation models (or even of string-based inflationary models) have so far concentrated on a single-field description meaning, in the DBI case, that the inflaton corresponds to a radial coordinate of the brane in the extra dimensions. Taking into account the "angular" coordinates of the brane naturally leads to a multifield description since each brane coordinate in the extra dimensions gives rise to a scalar field from the effective four-dimensional point of view. This setup has started to be explored only very recently [3,4].

In this Letter, we show that the multifield DBI action contains some terms, higher order in spacetime gradients and vanishing in the homogeneous case, which have been overlooked. The inclusion of these terms leads to drastic consequences on the primordial fluctuations generated in these types of models. The scalar-type perturbations in multifield models can be divided into (instantaneous) adiabatic modes, fluctuations along the trajectory in field space, and entropy modes which are orthogonal to the former [5]. In contrast with previous expectations, we show that, in DBI models, these two classes of modes propagate with the same speed, namely, an effective speed of sound c_s smaller than the speed of light. As a consequence, the amplification of quantum fluctuations occurs at the sound horizon crossing for both types of modes. Moreover, when $c_s \ll 1$, this leads to an enhancement of the amplitude of the entropy modes with respect to that of the usual adiabatic modes. As primordial non-Gaussianities—potentially detectable in forthcoming experiments if strong enough—discriminate between various models, we also study the impact of the entropy modes on non-Gaussianity in the DBI case.

Our starting point is the DBI Lagrangian governing the dynamics of a D3-brane:

$$L_{\rm DBI} = -\frac{1}{f}\sqrt{-\det(g_{\mu\nu} + fG_{IJ}\partial_\mu\phi^I\partial_\nu\phi^J)}, \qquad (1)$$

where $f = f(\phi^I)$ is a function of the scalar fields ϕ^I ($I = 1, 2, \ldots$) and $G_{IJ}(\phi^K)$ is a metric in field space. From a higher-dimensional point of view, (1) is proportional to the square root of the determinant of the induced metric on the brane, meaning that the ϕ^I correspond to the brane coordinates in the extra dimensions, f embodies the warp factor, and G_{IJ} is (up to a rescaling) the metric in the extra dimensions. We also allow for the presence of a potential and hence consider a full action of the form

$$S = \int d^4x\sqrt{-g}\left[\frac{\mathcal{R}}{2} + P\right], \qquad P = -\frac{1}{f(\phi^I)}(\sqrt{\mathcal{D}} - 1) - V(\phi^I), \qquad (2)$$

where we have set $8\pi G = 1$. The determinant $\mathcal{D} = \det(\delta^\mu_\nu + fG_{IJ}\partial^\mu\phi^I\partial_\nu\phi^J)$ coming from Eq. (1) can be rewritten as

$$\begin{aligned}\mathcal{D} &= \det(\delta^J_I - 2fX^J_I)\\ &= 1 - 2fG_{IJ}X^{IJ} + 4f^2X^{[I}_I X^{J]}_J - 8f^3X^{[I}_I X^J_J X^{K]}_K\\ &\quad + 16f^4X^{[I}_I X^J_J X^K_K X^{L]}_L,\end{aligned} \qquad (3)$$

where we have defined

$$X^{IJ} \equiv -\tfrac{1}{2}\partial^\mu\phi^I\partial_\mu\phi^J, \qquad X^J_I = G_{IK}X^{KJ}, \qquad (4)$$

and where the brackets denote antisymmetrization of the field indices. In the *single-field* case ($I = 1$), the terms in

f^2, f^3, and f^4 in (3) vanish. This is also true for multiple homogeneous scalar fields for which $X^{IJ} = \frac{1}{2}\dot{\phi}^I\dot{\phi}^J$. However, for multiple inhomogeneous scalar fields, these terms, which are higher order in gradients and have not been considered in previous works, do not vanish. We now show that they change drastically the behavior of perturbations.

In order to study the dynamics of linear perturbations about a homogeneous cosmological solution, we expand the initial action (2) to second order in the linear perturbations, including both metric and scalar field perturbations. This is a constrained system, and the number of (scalar) degrees of freedom is the same as the number of scalar fields. It is convenient to express these degrees of freedom in terms of the scalar field perturbations defined in the flat gauge, usually denoted Q^I. To obtain the second-order action, we follow the procedure outlined in Ref. [6] for a Lagrangian of the form $P(X, \phi^J)$, with $X = G_{IJ}X^{IJ}$. As we have stressed above, the multifield DBI Lagrangian is not of this form, but despite that it can be rewritten as

$$\tilde{P}(\tilde{X}, \phi^K) = -\frac{1}{f}(\sqrt{1 - 2f\tilde{X}} - 1) - V, \tag{5}$$

where $\tilde{X} = (1 - \mathcal{D})/(2f)$ [see Eq. (3)]. Although in the homogeneous background $\tilde{X}$ and X coincide, their perturbed values differ. Taking into account the corresponding extra terms, one can show [7] that the second-order action can be written in the compact form

$$S_{(2)} = \frac{1}{2}\int dt\, d^3x\, a^3\Big[\tilde{P}_{,\tilde{X}}\tilde{G}_{IJ}\mathcal{D}_tQ^I\mathcal{D}_tQ^J - \frac{c_s^2}{a^2}\tilde{P}_{,\tilde{X}}\tilde{G}_{IJ}\partial_iQ^I\partial^iQ^J - \mathcal{M}_{IJ}Q^IQ^J + 2\tilde{P}_{,\tilde{X}J}\dot{\phi}_IQ^J\mathcal{D}_tQ^I\Big]. \tag{6}$$

Here a is the cosmological scale factor; the effective (squared) mass matrix is

$$\mathcal{M}_{IJ} = -\mathcal{D}_I\mathcal{D}_J\tilde{P} - \tilde{P}_{,\tilde{X}}\mathcal{R}_{IKLJ}\dot{\phi}^K\dot{\phi}^L + \frac{X\tilde{P}_{,\tilde{X}}}{H}(\tilde{P}_{,\tilde{X}J}\dot{\phi}_I + \tilde{P}_{,\tilde{X}I}\dot{\phi}_J) + \frac{\tilde{X}\tilde{P}^3_{,\tilde{X}}}{2H^2}\Big(1 - \frac{1}{c_s^2}\Big)\dot{\phi}_I\dot{\phi}_J - \frac{1}{a^3}\mathcal{D}_t\Big[\frac{a^3}{2H}\tilde{P}^2_{,\tilde{X}}\Big(1 + \frac{1}{c_s^2}\Big)\dot{\phi}_I\dot{\phi}_J\Big], \tag{7}$$

and we have introduced covariant derivatives $\mathcal{D}_I$ defined with respect to the field space metric G_{IJ}, as well as the time covariant derivative $\mathcal{D}_tQ^I \equiv \dot{Q}^I + \Gamma^I_{JK}\dot{\phi}^JQ^K$, where Γ^I_{JK} is the Christoffel symbol constructed from G_{IJ} and $\mathcal{R}_{IKLJ}$ is the corresponding Riemann tensor. Finally, we have defined the (background) matrix

$$\tilde{G}_{IJ} = G_{IJ} + \frac{2fX}{1 - 2fX}e_{\sigma I}e_{\sigma J} = \perp_{IJ} + \frac{1}{c_s^2}e_{\sigma I}e_{\sigma J}, \tag{8}$$

where $e^I_\sigma = \dot{\phi}^I/\sqrt{2X}$ ($\dot{\sigma} \equiv \sqrt{2X}$ is also used in the following) is the unit vector pointing along the trajectory in field space, $\perp_{IJ} \equiv G_{IJ} - e_{\sigma I}e_{\sigma J}$ is the projector orthogonal to the vector e^I_σ, and

$$c_s^2 \equiv \frac{\tilde{P}_{,\tilde{X}}}{\tilde{P}_{,\tilde{X}} + 2\tilde{X}\tilde{P}_{,\tilde{X}\tilde{X}}} = 1 - f\dot{\sigma}^2. \tag{9}$$

Let us stress that the only difference between action (6) and the corresponding expression in Ref. [6] is the term in spatial gradients, with coefficient $c_s^2\tilde{P}_{,\tilde{X}}\tilde{G}_{IJ}$ instead of $\tilde{P}_{,X}G_{IJ}$. This crucial difference implies that all perturbations, both adiabatic and entropic, propagate with the same speed of sound in multifield DBI inflation, in contrast with Refs. [3,4,6], where they have different speeds. Finally, one should recall that the above expressions apply to the DBI context where $\tilde{P}$ is given in (5) so that $\tilde{P}_{,\tilde{X}} = 1/c_s$.

For simplicity, let us now restrict our attention to two fields ($I = 1, 2$). The perturbations can then be decomposed into $Q^I = Q_\sigma e^I_\sigma + Q_s e^I_s$, where e^I_s, the unit vector orthogonal to e^I_σ, characterizes the entropy direction. [For N fields, the entropy modes would span an $(N-1)$-dimensional subspace in field space.] As in standard inflation, it is more convenient, after going to conformal time $\tau = \int dt/a(t)$, to work in terms of the canonically normalized fields

$$v_\sigma \equiv \frac{a}{c_s}\sqrt{\tilde{P}_{,\tilde{X}}}Q_\sigma, \qquad v_s \equiv a\sqrt{\tilde{P}_{,\tilde{X}}}Q_s. \tag{10}$$

Note that the adiabatic and entropy coefficients differ because $\tilde{G}_{IJ}$ is anisotropic. The equations of motion for v_σ and v_s then follow from the action (6), reexpressed in terms of the rescaled quantities (10). One finds

$$v''_\sigma - \xi v'_s + \Big(c_s^2k^2 - \frac{z''}{z}\Big)v_\sigma - \frac{(z\xi)'}{z}v_s = 0, \tag{11}$$

$$v''_s + \xi v'_\sigma + \Big(c_s^2k^2 - \frac{\alpha''}{\alpha} + a^2\mu_s^2\Big)v_s - \frac{z'}{z}\xi v_\sigma = 0, \tag{12}$$

where

$$\xi = \frac{a}{\dot{\sigma}\tilde{P}_{,\tilde{X}}c_s}[(1 + c_s^2)\tilde{P}_{,s} - c_s^2\dot{\sigma}^2\tilde{P}_{,\tilde{X}s}], \tag{13}$$

$$z = \frac{a\dot{\sigma}}{c_sH}\sqrt{\tilde{P}_{,\tilde{X}}}, \qquad \alpha = a\sqrt{\tilde{P}_{,\tilde{X}}}, \tag{14}$$

and μ_s^2 follows from the mass matrix (7) (see [6,7] for details). We will assume that the effect of the coupling ξ can be neglected when the scales of interest cross out the sound horizon, so that the two degrees of freedom are decoupled and the system can easily be quantized. In the slow-varying regime, where the time evolution of H, c_s, and $\dot{\sigma}$ is slow with respect to that of the scale factor, one gets $z''/z \simeq 2/\tau^2$ and $\alpha''/\alpha \simeq 2/\tau^2$. The solutions of (11) and (12) corresponding to the Minkowski-like vacuum on small scales are thus

$$v_{\sigma k} \simeq v_{sk} \simeq \frac{1}{\sqrt{2kc_s}}e^{-ikc_s\tau}\Big(1 - \frac{i}{kc_s\tau}\Big), \tag{15}$$

when μ_s^2/H^2 is negligible for the entropic modes (if

μ_s^2/H^2 is large, the entropic modes are suppressed). The power spectra for v_σ and v_s after sound horizon crossing therefore have the same amplitude $\mathcal{P}_v = (k^3/2\pi^2)|v_k|^2$. The power spectra for Q_σ and Q_s are thus

$$\mathcal{P}_{Q_\sigma} \simeq \frac{H^2}{4\pi^2 c_s \tilde{P}_{,\tilde{X}}}, \qquad \mathcal{P}_{Q_s} \simeq \frac{H^2}{4\pi^2 c_s^3 \tilde{P}_{,\tilde{X}}}, \tag{16}$$

evaluated at the sound horizon crossing. One recognizes the familiar result of k inflation for the adiabatic part [8,9], while for small c_s, the entropic modes are amplified with respect to the adiabatic modes: $Q_s \simeq Q_\sigma/c_s$.

These results can be reexpressed in terms of the comoving curvature perturbation $\mathcal{R} = (H/\dot{\sigma})Q_\sigma$ which is useful to relate the perturbations during inflation to the primordial fluctuations during the standard radiation era. We recover the usual single-field result for the power spectrum of $\mathcal{R}$ at the sound horizon crossing:

$$\mathcal{P}_{\mathcal{R}_*} \simeq \frac{H^4}{8\pi^2 c_s \tilde{X}\tilde{P}_{,\tilde{X}}} = \frac{H^2}{8\pi^2 \epsilon c_s}, \tag{17}$$

where $\epsilon = -\dot{H}/H^2 = \tilde{X}\tilde{P}_{,\tilde{X}}/H^2$ (the subscript $*$ indicates that the corresponding quantity is evaluated at the sound horizon crossing). It is then convenient to define an entropy perturbation $\mathcal{S} = c_s \frac{H}{\dot{\sigma}} Q_s$ such that $\mathcal{P}_{\mathcal{S}_*} \simeq \mathcal{P}_{\mathcal{R}_*}$. The power spectrum for the tensor modes is, as usual, governed by the transition at Hubble radius, and its amplitude $\mathcal{P}_\mathcal{T} = (2H^2/\pi^2)_{k=aH}$ is much smaller than the curvature amplitude for $c_s \ll 1$.

Leaving aside the possibility that the entropy modes during inflation lead directly to primordial entropy fluctuations that could be detectable in the CMB fluctuations (potentially correlated with adiabatic modes as discussed in Ref. [10]), we consider here only the influence of the entropy modes on the final curvature perturbation. Indeed, on large scales, the curvature perturbation can evolve in time in the multifield case, because of the entropy modes. This transfer from the entropic to the adiabatic modes depends on the details of the scenario and on the background trajectory in field space, but it can be parametrized by a transfer coefficient [11] which appears in the formal solution $\mathcal{R} = \mathcal{R}_* + T_{\mathcal{RS}}\mathcal{S}_*$ of the first-order evolution equations for $\mathcal{R}$ and $\mathcal{S}$ which follow from (11) and (12) in the slow-varying regime on large scales. Equivalently, one can define the "transfer angle" Θ by $\sin\Theta = T_{\mathcal{RS}}/\sqrt{1+T_{\mathcal{RS}}^2}$ (so that $\Theta = 0$ if there is no transfer and $|\Theta| = \pi/2$ if the final curvature perturbation is mostly of entropic origin). This implies in particular that the final curvature power spectrum can be formally expressed as $\mathcal{P}_\mathcal{R} = (1+T_{\mathcal{RS}}^2)\mathcal{P}_{\mathcal{R}_*} = \mathcal{P}_{\mathcal{R}_*}/\cos^2\Theta$. Therefore, the tensor to scalar ratio is given by

$$r \equiv \frac{\mathcal{P}_\mathcal{T}}{\mathcal{P}_\mathcal{R}} = 16\epsilon c_s \cos^2\Theta. \tag{18}$$

This expression combines the result of k inflation [9], where the ratio is suppressed by the sound speed c_s and that of standard multifield inflation [5].

We finally turn to primordial non-Gaussianities, whose detection would provide an additional window on the very early Universe. This aspect is especially important for DBI models since it is well known that (single-field) DBI inflation produces a (relatively) high level of non-Gaussianity for small c_s [2]. How, therefore, do the entropic modes, whose amplitude is much larger than that of the adiabatic fluctuations, affect the primordial non-Gaussianity? In the small c_s limit, one can estimate the dominant contribution by extracting from the third-order Lagrangian the analogue of the terms giving the dominant contribution in the single-field case, but including now the entropy components. These terms are [7]

$$S_{(3)}^{(\mathrm{main})} = \int dt d^3x \frac{a^3}{2c_s^5 \dot{\sigma}}[\dot{Q}_\sigma^3 + c_s^2 \dot{Q}_\sigma \dot{Q}_s^2] - \frac{a}{2c_s^3\dot{\sigma}}[\dot{Q}_\sigma(\nabla Q_\sigma)^2 - c_s^2 \dot{Q}_\sigma(\nabla Q_s)^2 + 2c_s^2 \dot{Q}_s \nabla Q_\sigma \nabla Q_s],$$

where we have replaced f by $1/\dot{\sigma}^2$ since, for $c_s \ll 1$, $f\dot{\sigma}^2 \simeq 1$. Following the standard procedure [12–14], one can compute the 3-point functions involving adiabatic and entropy fields. The purely adiabatic 3-point function is naturally the same as in single-field DBI [15,16]. The new contribution is

$$\begin{aligned}\langle Q_\sigma(\boldsymbol{k}_1)Q_s(\boldsymbol{k}_2)Q_s(\boldsymbol{k}_3)\rangle = {}& -(2\pi)^3\delta\Big(\sum_i \boldsymbol{k}_i\Big)\frac{H^4}{4\sqrt{2c_s\epsilon}c_s^4}\frac{1}{(\prod_i k_i^3)K^3}[-2k_1^2k_2^2k_3^2 - k_1^2(\boldsymbol{k}_2\cdot\boldsymbol{k}_3)(2k_2k_3 - k_1K + 2K^2)\\ & + k_3^2(\boldsymbol{k}_1\cdot\boldsymbol{k}_2)(2k_1k_2 - k_3K + 2K^2) + k_2^2(\boldsymbol{k}_1\cdot\boldsymbol{k}_3)(2k_1k_3 - k_2K + 2K^2)],\end{aligned} \tag{19}$$

where $K = \sum_i k_i$. By using $\mathcal{R} \approx \mathcal{A}_\sigma Q_{\sigma*} + \mathcal{A}_s Q_{s*}$, with $\mathcal{A}_\sigma = (H/\dot{\sigma})_*$ and $\mathcal{A}_s = T_{\mathcal{RS}} c_{s*} \mathcal{A}_\sigma$, one can express the 3-point function of the curvature perturbation $\mathcal{R}$, which is the observable quantity, in terms of the correlation functions of the scalar fields. We find

$$\begin{aligned}\langle \mathcal{R}(\boldsymbol{k}_1)\mathcal{R}(\boldsymbol{k}_2)\mathcal{R}(\boldsymbol{k}_3)\rangle &= (\mathcal{A}_\sigma)^3\langle Q_\sigma(\boldsymbol{k}_1)Q_\sigma(\boldsymbol{k}_2)Q_\sigma(\boldsymbol{k}_3)\rangle + \mathcal{A}_\sigma(\mathcal{A}_s)^2[\langle Q_\sigma(\boldsymbol{k}_1)Q_s(\boldsymbol{k}_2)Q_s(\boldsymbol{k}_3)\rangle + \text{perm.}]\\ &= (\mathcal{A}_\sigma)^3\langle Q_\sigma(\boldsymbol{k}_1)Q_\sigma(\boldsymbol{k}_2)Q_\sigma(\boldsymbol{k}_3)\rangle(1 + T_{\mathcal{RS}}^2),\end{aligned} \tag{20}$$

where we have implicitly assumed that the vectors $\boldsymbol{k}_i$ are of the same order of magnitude (so that the slowly varying background parameters are evaluated at about the same time). As we see, the above quantity depends on the symmetrized version of the 3-point function (19), which has exactly the same shape as in single-field DBI. Note that the enhancement of the mixed correlation $\langle Q_\sigma Q_s Q_s \rangle$ by a factor of $1/c_s^2$ is compensated by the ratio between $\mathcal{A}_\sigma$ and $\mathcal{A}_s$ so that the adiabatic and mixed contributions in (20) are exactly of the same order. In principle, there are other contributions to the observable 3-point function, in particular, those coming from the 4-point function of the scalar fields, which can be reexpressed in terms of the power spectrum via Wick's theorem [17]. The amplitude of this contribution will depend on the specific models. We implicitly ignore them in the following.

The non-Gaussianity parameter $f_{\rm NL}$ is defined by

$$\langle \mathcal{R}(\boldsymbol{k}_1)\mathcal{R}(\boldsymbol{k}_2)\mathcal{R}(\boldsymbol{k}_3)\rangle = -(2\pi)^7 \delta\Big(\sum_i \boldsymbol{k}_i\Big)\Big[\frac{3}{10} f_{\rm NL}(\mathcal{P}_{\mathcal{R}})^2\Big] \times \frac{\sum_i k_i^3}{\prod_i k_i^3}, \tag{21}$$

from which we obtain, for the equilateral configuration,

$$f_{\rm NL}^{(3)} = -\frac{35}{108}\frac{1}{c_s^2}\frac{1}{1+T_{\mathcal{RS}}^2} = -\frac{35}{108}\frac{1}{c_s^2}\cos^2\Theta. \tag{22}$$

One can easily understand this result. The curvature power spectrum is amplified by a factor of $(1 + T_{\mathcal{RS}}^2)$ due to the feeding of curvature by entropy modes. Similarly, the 3-point correlation function for $\mathcal{R}$ resulting from the 3-point correlation functions of the adiabatic and entropy modes is enhanced by the same factor $(1 + T_{\mathcal{RS}}^2)$. However, since $f_{\rm NL}$ is roughly the ratio of the 3-point function with respect to the square of the power spectrum, one sees that $f_{\rm NL}$ is now reduced by the factor $(1 + T_{\mathcal{RS}}^2)$. The so-called UV model of DBI inflation is under strong observational pressure because it generates a high level of non-Gaussianities that exceed the experimental bound [18,19]. We stress that their reduction by multiple-field effects may be very important for model building.

We end by revisiting the consistency condition relating the non-Gaussianity of the curvature perturbation, the tensor to scalar ratio r, and the tensor spectral index $n_{\mathcal{T}} = -2\epsilon$, given in Ref. [20] for single-field DBI. In our case, substituting $f_{\rm NL}^{(3)} \simeq -\frac{1}{3}\frac{1}{c_s^2}\cos^2\Theta$ in (18) gives

$$r + 8n_{\mathcal{T}} = -r\big(\sqrt{-3f_{\rm NL}^{(3)}}\cos^{-3}\Theta - 1\big). \tag{23}$$

As we can can see from (22) and (23), violation of the standard inflation consistency relation [corresponding to a vanishing right-hand side in (23)] would be stronger in multifield DBI than in single-field DBI and thus easier to detect. In the multifield case, the consistency condition is only an inequality (unless Θ is observable when the entropy modes survive after inflation).

To summarize, we have shown that both adiabatic and entropy modes propagate with the same speed of sound c_s, in multifield DBI models. Both modes are thus amplified at the sound horizon crossing, with an enhancement of the entropy modes with respect to the adiabatic ones in the small c_s limit. The amplitude of the non-Gaussianities, which are important in DBI models, is also strongly affected by the entropy modes, although their shape remains as in the single-field case. All of these features are generic in any model governed by the multifield DBI action. The model-specific quantity (depending on the field metric, the warp factor, and the potential) is the transfer coefficient between the initial entropy modes and the final curvature perturbation between the time when the fluctuations cross out the sound horizon and the end of inflation. Recent analyses (see, e.g., [21]) in slightly different contexts show that this transfer can be very efficient, leading to a final curvature perturbation of entropic origin (as in the curvaton scenario). More generally, our results show that multifield effects, common in string theory-motivated inflation models, deserve close attention as the entropy modes produced could significantly affect the cosmological observable quantities.

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