Constraints on holographic multi-field inflation

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In holographic inflation, the four-dimensional cosmological dynamics is postulated to be dual to the renormalization group flow of a three-dimensional Euclidean conformal field theory with marginally relevant operators. The scalar potential of the 4D theory—in which inflation is realized—is highly constrained, with use of the Hamilton–Jacobi equations. We show that in multi-field holographic realizations of inflation, the dynamical masses of all fields additional to the inflaton must respect an upper bound of the form \( \mu \leq 3H/2 \) up to slow roll corrections. Indeed this bound applies to any multi-field model of inflation which uses the Hamilton–Jacobi equations. This upper bound reduces to the analytic continuation of the well known Breitenlohner–Freedman bound found in AdS spacetimes in the case when the masses are approximately constant, and it is found to be independent of the number of fields, the field space geometry and/or the shape of the inflationary trajectory followed in multi-field space. We infer that such models do not allow fields sufficiently heavy to be integrated out, and may present a number of interesting phenomenological consequences that could be confirmed or constrained by future surveys. For instance, a detection of “cosmological collider” oscillatory patterns in the non-Gaussian bispectrum due to massive fields, would uncover the existence of fields with dynamical masses larger than 3H/2 during inflation, therefore ruling out holographic inflation or any inflationary models based on the Hamilton–Jacobi equations.

Introduction: The observation of departures from a perfectly Gaussian distribution of primordial curvature perturbations would allow us to infer rather fundamental information about cosmic inflation \cite{1–5}. It is by now well understood that single field models of inflation cannot account for the generation of local non-Gaussianity unless a nontrivial self-interaction—together with a non-attractor background evolution—plays a role in inducing it \cite{6–10}. This is mostly due to the fact that the dynamics of curvature perturbations is highly constrained by the diffeomorphism invariance of the gravitational theory within which inflation is realised. On the other hand, multi-field models of inflation can accommodate non-gravitational interactions affecting the dynamics of curvature perturbations: isocurvature fields orthogonal to the inflationary trajectory can efficiently transfer their non-Gaussian statistics—resulting from their own self-interactions—to curvature perturbations \cite{11–20}. As such, the detection of non-Gaussianity could reveal signatures only attributable to the presence of additional degrees of freedom interacting with curvature perturbations \cite{21–25}, and lead to the discovery of new energy scales playing a role during the epoch of inflation.

Understanding the theoretical restrictions on the various classes of interactions coupling together fields in multi-field systems would allow us to interpret future observations related to non-Gaussianity. For example, multi-field inflationary systems derived from supergravity models, characterized by non-flat Kähler geometries, are severely restricted due to the way in which the gravitational interaction couples chiral fields together. As a consequence, it is not easy to spontaneously break supersymmetry and keep, at the same time, every chiral field stabilized while sustaining inflation. A similar situation holds in string theory compactifications, where various fields have a geometrical origin that determines how they couple to other fields at energies much lower than the compactification scale, making it hard to build a quasi-de Sitter stage where all the moduli are stabilized (for a recent up to date discussion see Ref. \cite{26} and references therein). These restrictions do not only impose a challenge to the construction of realistic models of inflation in fundamental theories, but they also have consequences for the prediction of observable primordial spectra \cite{27}.

Other classes of potentially well motivated multi-field constructions, enjoying a constrained structure, have received less attention. In particular, holographic models of inflation have self-interactions constrained by certain requirements on the “holographic” correspondence connecting the 4D inflationary bulk cosmology and the 3D Euclidean conformal field theory (CFT) \cite{28–33}. More precisely, in this class of theories, the inflationary dynamics in 4D is dual to a renormalization group flow realized in a 3D Euclidean CFT with marginally relevant scalar operators deforming the conformal symmetry. In these constructions, the action determining the homogeneous cosmological dynamics is conjectured to be

\begin{equation}
S = \int d^4x \sqrt{-g} \left( \frac{1}{2} R - \frac{1}{2} \gamma_{ab}(\phi) \partial_a \phi^a \partial_b \phi^b - V(\phi) \right),
\end{equation}

where the first term corresponds to the Einstein-Hilbert
action for the spacetime metric $g_{\mu\nu}$ (in units where the reduced Planck mass is unity), and $\gamma_{ab}$ is a sigma-model metric characterizing the geometry of the multi-scalar field target space spanned by $\phi^a$ (with $a = 1, \cdots, N$). The inflationary potential $V$ is determined by a “fake” superpotential $W(\phi)$ as

$$V = 3W^2 - 2\gamma^{ab}W_aW_b,$$

(2)

where $\gamma^{ab}$ is the inverse of $\gamma_{ab}$, and $W_a = \partial W / \partial \phi^a$. The inflationary solutions admitted by (2) that are dual to the renormalization group flow are given by Hamilton–Jacobi equations, of the form

$$\dot{\phi}^a = -2\gamma^{ab}W_b,$$

(3)

$$H = W,$$

(4)

where $H = \dot{a}/a$ is the Hubble parameter. The trajectory described by this solution is dual to the renormalization group flow of the boundary operators with fixed points representing static de Sitter configurations of the cosmological bulk. In this sense, the entire cosmological history, starting from a static de Sitter universe (inflation), and ending in another static de Sitter universe (our dark energy dominated universe) may be understood as the consequence of renormalization group flow from the UV-fixed point (late universe) to the IR-fixed point (early universe).

The purpose of this article is to study some of the consequences on the dynamics of multi-field fluctuations coming from the constrained structure of the potential of Eq. (2). We are interested in how the structure of (2), together with the multi-field trajectory (3), constrains the interactions between the primordial curvature perturbation and other (isocurvature) fields in multi-field holographic models of inflation, and how this affects the observational signatures. Our results apply to any model described by the Hamilton–Jacobi equations (2) and (3), regardless of the holographic interpretation.

In particular, there is a known upper mass bound on all fields additional to the inflaton

$$m \leq m_{\text{max}} \equiv \frac{3}{2}H.$$

(5)

This bound was derived in [34] in the single-field case, and argued to be valid in the multifield case in [35] under the implicit assumption that all masses are constant. It coincides with the analytic continuation of the Breitenlohner–Freedman bound encountered in scalar field theories in AdS spacetimes [36].

As we shall see, on a time-dependent background such as inflation, both sides of (5) receive corrections. First of all, for fields orthogonal to the trajectory, the bound applies to the dynamical, “entropy” mass, which is different from that obtained from the Hessian of the potential. Secondly, the upper bound receives corrections if the masses evolve in time, which is the generic situation during inflation, but these corrections are small if inflation is slow-roll. This bound happens to be independent of the strength of the interaction between the fields and the inflaton, of the number of fields, of the field space geometry and even of the shape of the inflationary trajectory followed in multi-field space. As a consequence, in holographic models of inflation, degrees of freedom additional to the inflaton could lead to certain distinguishable non-Gaussian features that could be observed in future surveys. In particular, if future observations reveal evidence of spin-0 fields with entropy masses larger than $3H/2$, then holographic versions of inflation would be essentially ruled out.

**Derivation of the bound:** We can motivate the bound by considering the simplest case: a straight trajectory in a model with $N$ fields with canonical kinetic terms $\gamma_{ab} = \delta_{ab}$. In this case the dynamical “entropy” mass coincides with the naive mass (we will later discuss the most general situation). Without loss of generality we can take the inflationary trajectory to be along the $\phi^i \equiv \phi$ direction, with all other fields stabilized: $\phi^i \equiv \sigma^i = \sigma^i_0$ for $i = 2, \cdots, N$. Note that the equations (3) imply $W_{\sigma^i} = 0$ on the inflationary trajectory, and we can expand the superpotential as

$$W(\phi, \sigma^i) = w(\phi) + \frac{1}{2} \sum_{i=2}^{N} h_i(\phi)(\sigma^i - \sigma^i_0)^2 + \cdots,$$

(6)

where $w(\phi)$ and $h_i(\phi)$ are given functions of $\phi$. Inserting this expression into (2) gives

$$V = 3w^2 - 2(w')^2 + \frac{1}{2} \sum_i m_i^2(\phi)(\sigma^i - \sigma^i_0)^2,$$

(7)

where the masses $m_i(\phi)$ of the fields $\sigma^i$ at each point on the trajectory are found to be given by

$$m_i^2(\phi) = 6wh_i - 4h_i^2 - 4w'h_i,$$

(8)

We can rewrite this expression in a more useful way by noticing from Eq. (3) that $w = H$ and $w' = -\phi/2$. Then, we obtain $m_i^2 = 6Hh_i(1 + \delta_i/3) - 4h_i^2$, where we have defined $\delta_i = h_i/Hh_i$. From the previous expression, it follows that $m_i^2$ has a maximum value given by

$$m_{\text{max}} = \frac{3H}{2}(1 + \delta/3).$$

(9)

Notice that $\delta_i$ measures the running of $h_i$, which near the maximum satisfies $h_i \sim H$. If background quantities evolve slowly, then we expect $\delta \sim O(\epsilon)$, implying that masses stay almost constant during slow roll. On the other hand, unless $h_i \ll H$, a large value of $\delta_i$ could make the field $\sigma^i$ tachyonic.

**Origin of the bound:** In the particular case of holographic inflation, because the inflationary trajectory is dual to the renormalization group flow in the CFT side of the duality, the potential driving inflation must always
admit monotonic solutions of the form (3), regardless of the initial conditions. This is satisfied for flows that are solutions of the Hamilton–Jacobi equations. This restricts the value of the masses of the fields, simply because a field with mass larger than 3H/2 has non-monotonic trajectories. To appreciate this, let us disregard the motion of the inflation $\phi$ and focus on the background evolution of one of the massive fields $\sigma$ with a mass $m$. Its background equation of motion is given by

$$\dot{\sigma} + 3H\sigma + m^2(\sigma - \sigma_0) = 0.$$  \hspace{1cm} (10)

The general solution is of the form $\sigma(t) = \sigma_0 + A_+ e^{\omega_+ t} + A_- e^{\omega_- t}$ with:

$$\omega_{\pm} = -\frac{3}{2} H \pm \frac{3}{2} H \sqrt{1 - \frac{4m^2}{9H^2}}.$$  \hspace{1cm} (11)

If $m < m_{\text{max}}$, the solutions are overdamped, and the field $\sigma$ reaches $\sigma = \sigma_0$ monotonically at a time $t \gg H^{-1}$. On the other hand, if $m > m_{\text{max}}$ the underdamped solutions are oscillatory, and not of the desired form $\dot{\sigma} = f(\sigma)$.

**Bending trajectories in arbitrary field geometries:** We can see that the upper bound primarily comes from the negative sign of the second term in (2), which affects equally any field in the system. One could wonder whether the bound changes in more general situations where $\gamma_{ab}$ is non-canonical, and the inflationary trajectory in multi-field space is not associated with a single field (that is, it is not a straight line in multi-field space). To analyze this situation let us simplify our strategy by reducing the problem to a two field system, but allowing arbitrary sigma-model metrics $\gamma_{ab}$ and arbitrary inflationary trajectories. Of course, we will demand that the system continues to respect the background equations (3) and (4). To proceed, we introduce a local basis, consisting of unit-vectors $T^a$ and $N^a$, tangential and normal to the trajectory $\phi^a = \phi^a(t)$, in the following way [37, 38]:

$$T^a = \frac{\dot{\phi}_0}{\dot{\phi}}, \quad N^a = -\sqrt{\epsilon_0} \epsilon_{ab} T^b,$$  \hspace{1cm} (12)

where $\dot{\phi}_0 = \sqrt{\gamma_{ab}\phi^a\phi^b}$. This basis can be used to decompose the metric as $\gamma_{ab} = T_aT_b + N_aN_b$. To describe the local dynamics around the trajectory, it is useful to introduce the rate of turning $\Omega$ at which the inflationary trajectory bends in field space, defined by

$$\frac{D}{dt} T^a = -\Omega N_a, \quad \frac{D}{dt} N^a = \Omega T^a,$$  \hspace{1cm} (13)

where $D_t$ is a covariant derivative defined to act on vectors as $D_t A^a = \dot{A}^a + \Gamma_{bc}^a \dot{\phi}^b A^c$, where $\Gamma_{bc}^a$ represents Christoffel symbols.

The basis $(T^a, N^b)$ may be used to project gradients of various quantities, allowing us to find useful expressions for them. For instance, using the equations of motion (3) and (4), it is straightforward to find that the projection of the Hessian of $W$ gives

$$N^a T^b \nabla_a W_b = \frac{1}{2} \Omega,$$  \hspace{1cm} (14)

$$T^a T^b \nabla_a W_b = -\frac{1}{4} H(2\epsilon - \eta),$$  \hspace{1cm} (15)

$$N^a N^b \nabla_a W_b = W_{NN},$$  \hspace{1cm} (16)

where $\epsilon \equiv -\dot{H}/H^2$ and $\eta \equiv \dot{\epsilon}/H\epsilon$ are the usual slow-roll parameters, required to be small in order to achieve inflation with the right properties. In Eq. (16), the quantity $W_{NN}$ is a model dependent parameter, on a par with $\Omega$ and/or slow-roll parameters. We can also use $(T^a, N^b)$ to project gradients of $V$. For instance, one can write $\nabla_a V = T_a V_T + N_a V_N$. With the help of the equations of motion, one finds $V_T = \sqrt{\epsilon/2}H^2(6-2\epsilon+\eta)$ and $V_N = \Omega \dot{\phi}$. One also finds expressions for the projections of second gradients of $V$ along $T^a$ and $N^a$ [27, 38]:

$$T^a T^b \nabla_a \nabla_b V = 3H^2(2\epsilon - \eta/2) + \Omega^2$$  \hspace{1cm} (17)

$$N^a T^b \nabla_a \nabla_b V = \Omega H \left( 3 + \frac{\dot{\Omega}}{H \Omega} + \eta - 2\epsilon \right)$$  \hspace{1cm} (18)

$$N^a N^b \nabla_a \nabla_b V = \mu^2 - 3\Omega^2 - \epsilon H^2 \mathbb{R},$$  \hspace{1cm} (19)

In the previous expressions, $\mathbb{R}$ is the Ricci scalar computed from the sigma model metric $\gamma_{ab}$. It can be shown that $\mathbb{R} = 2\epsilon_{abcd} T^a T^b N^c N^d$, where $\epsilon_{abcd}$ is the Riemann tensor. We have introduced the dynamical entropy mass $\mu$ of the isocurvature fluctuation, orthogonal to the inflationary trajectory [39]. Note that $\mu^2$ differs from $m^2 = N^a N^b \nabla_a \nabla_b V$. While on short wavelengths the isocurvature field interacts with the curvature perturbations thanks to $\Omega$, on long wavelengths the isocurvature fluctuation $\sigma$ satisfies the same equation of motion given in (10), with $\sigma_0 = 0$ and $m^2$ replaced by $\mu^2$ (from where one sees that indeed $\mu$ may be identified as the mass of $\sigma$, independently of the value of $\Omega$). We will now show that in the general case the bound (9) becomes an upper bound on $\mu$.

Now, Eqs. (17)-(19) are completely general, and assume nothing about the origin of $V$. Since $V$ is given by Eq. (2), the first and second gradients of $V$ are found to be given by:

$$\nabla_a V = 6W_a \nabla_a W_b - 4W_b \nabla_a W_a,$$  \hspace{1cm} (20)

$$\nabla_a \nabla_b V = 6W_a \nabla_b W_c + 6W_b \nabla_a W_c - 4\nabla_b W_c \gamma^{cd} \nabla_a W_d - 4\nabla_c W_c \nabla_a W_b,$$  \hspace{1cm} (21)

These equations relate field-derivatives of $V$ with field-derivatives of $W$. We may now project both sides of these equations along different combinations of $T^a$ and $N^a$, and obtain relations between second derivatives of $W$ and the quantities parametrizing the inflationary trajectory. In the end, one only finds new information coming from the projection of $N^a N^b \nabla_a \nabla_b V$. In fact, one finds a non-trivial expression for the entropy mass $\mu^2$ in terms of
the superpotential. To obtain it, one may first insert
\[ \gamma_{ab} = T_a T_b + N_a N_b \] in (21), then project along \( N^a \) to obtain \( N^a N^b \nabla_a \nabla_b V \), and use Eqs. (14)-(16) to simplify the expression. One finds
\[ N^a N^b \nabla_a \nabla_b V = 6 W W_{NN} - 4(W_{NN})^2 - \epsilon H^2 R + 2 \dot{W}_{NN} - 3 \Omega^2, \]
where \( R = 2 R_{abcd} T^a N^b T^c N^d \) appears from simplifying the term containing \( \nabla_a \nabla_b W \) in Eq. (21). Then, by comparing this result with Eq. (19), we finally arrive at:
\[ \mu^2 = 6 W W_{NN} - 4(W_{NN})^2 + 2 \dot{W}_{NN}. \]
This result has exactly the same form of (8), except that \( h_i \) is replaced by \( W_{NN} \). Consequently, we conclude that the entropy mass is constrained by \( \mu \leq \mu_{\text{max}} \), where
\[ \mu_{\text{max}} = \frac{3}{2} H \left( 1 + \frac{\delta}{3} \right), \]
with \( \delta \equiv \dot{W}_{NN} / HW_{NN} \). Notice that Eq. (23) is completely model independent, and only relies on the requirement that the background equations of motion satisfy the first order differential equations (3) and (4). In particular, it does not depend on the value of the turning rate \( \Omega \). On the other hand, notice that \( \mu^2 \) is not bounded from below, and it may become negative (situations where the non-trivial field geometry can play a role in making \( \mu^2 \) negative has been analyzed in Ref. [40, 41]).

**Non-Gaussianity:** Up to now, we have shown that the bound on the mass of the additional fields is independent of the number of fields, and independent of the path followed by the inflationary trajectory in field space. We now address the observational consequences of this bound. Isocurvature fields can enhance the amplitude of non-Gaussianity up to levels that can be distinguished from single field models. Consequently, the universal bound (24) gives a window of opportunity to test holographic models of inflation.

To make the following discussion simple, we restrict ourselves to a 2-field model setup and work in the comoving gauge, where the inflaton fluctuation along the trajectory vanishes. In this gauge, the two scalar degrees of freedom are the isocurvature field \( \sigma \), and the curvature perturbation \( \zeta \), that is defined to perturb the metric as \( ds^2 = -dt^2 + a^2 e^{2 \zeta} dx^2 \). As soon as \( \Omega \neq 0 \) the field \( \sigma \) will have an effect on the 3-point function of \( \zeta \). The 3-point function can be computed using the in-in formalism, in which case the interaction picture Hamiltonian induced by a non-vanishing \( \Omega \) is given by
\[ H_I(t) = -\int d^3x \left[ {\mathcal{L}}^\text{int}_2 + {\mathcal{L}}^\text{int}_3 \right] \]
where
\[ {\mathcal{L}}^\text{int}_2 \propto \Omega \times \dot{\zeta} \sigma, \]
\[ {\mathcal{L}}^\text{int}_3 \propto \Omega \times \dot{\zeta}^2 \sigma. \]

Correlation functions for single field inflation are highly constrained by dilations and special conformal transformations, which are non linearly realised by \( \zeta \) at horizon crossing [42–46]. In particular the squeezed limit of the 3-point function \( \langle \zeta^3 \rangle \), which is when one of the momenta is taken to be soft, \( i.e. \) much smaller in magnitude than the other two. Now from Eq. (26) note that the vertex \( \langle \zeta^3 \rangle \) induces an interaction between the curvature mode \( \zeta \) and the massive field \( \sigma \). This is fixed to be proportional to the power spectrum of the other two long wavelengths modes, times a slow roll factor. Any deviations from this will be a signal of new physics appearing at horizon crossing. This implies for (26) that the squeezed limit has a dependence on the mass of \( \sigma \) [16, 21–25]. For the case when \( \mu < 3H/2 \) we have
\[ \langle \zeta q k_1 \zeta k_2 \rangle \sigma \simeq P_\zeta(q) P_\zeta(k) \left( \frac{q}{k} \right)^{3/2 - \nu} \]
where \( q \) is the soft momentum, and \( k_{1,2} \) are \( 1/k_1 \sim 1/k_2 \ll 1/q \). Also we have defined \( \nu \equiv \sqrt{\frac{9}{4} - \frac{\mu^2}{H^2}} > 0 \). The correlation functions contain a small non-analytic dependence due to interactions with the massive field. On the other hand for masses, \( \mu > 3H/2 \)
\[ \langle \zeta q k_1 \zeta k_2 \rangle \sigma \simeq P_\zeta(q) P_\zeta(k) \left( \frac{q}{k} \right)^{3/2} \cos(-i \nu \log \frac{q}{k} - \phi_0), \]
where the phase \( \phi_0 \) is fixed in terms of \( \nu \). The modulations appear because fields whose masses are \( \mu > 3/2H \) oscillate after horizon crossing, creating an interference pattern [21, 23]. These results imply a sharp difference between the masses above and below \( \frac{3}{2} H \). Detecting a signal such as (28) would disfavour holographic models based on (2). This pattern is part of what is know as cosmological collider [21]. These signatures could be observed with future surveys by looking, for example, at the dark matter distribution or the 21 cm line [47–49].

**Concluding remarks:** Future cosmological surveys, aimed at characterizing the distribution of primordial curvature perturbations, will be able to constrain holographic realizations of inflation. As we have seen, the value of the mass \( \mu \) of fields additional to the inflaton crucially determines the shape of non-Gaussian imprints in the primordial distribution of curvature perturbations. If future observations reveal the existence of massive scalar fields with approximately constant masses above the bound \( 3H/2 \), holographic constructions, or indeed any multi-field model which uses the Hamilton–Jacobi equations, would be ruled out.

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1 This variable is often called \( \mathcal{R} \) but here we follow the notation of [6].
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