Letter Editors' Suggestion

Large-field inflation and the cosmological collider

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Large-field inflation is a major class of inflation models featuring a near- or super-Planckian excursion of the inflaton field. We point out that the large excursion generically introduces significant scale dependence to spectator fields through inflaton couplings, which in turn induces characteristic distortions to the oscillatory shape dependence in the primordial bispectrum mediated by a spectator field. This so-called cosmological collider signal can thus be a useful indicator of large field excursions. We show an explicit example with signals from the "tower states" motivated by the swampland distance conjecture.

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I. INTRODUCTION

Despite the success of cosmic inflation as a widely accepted scenario for the primordial universe, its microscopic model largely remains unclear, due to the large degeneracy when confronting models with data. While it is virtually impossible to pin down the model of inflation, it is nevertheless possible to probe some general properties of inflation models by cosmological observations [1].

A notable property is the excursion range $\Delta\phi$ of the inflaton field from the horizon exit of CMB-scale modes to the end of inflation. Roughly, models with $\Delta\phi \gtrsim M_{\rm Pl} \simeq 2.4 \times 10^{18}$ GeV are called large-field inflation, while models with $\Delta\phi \ll M_{\rm Pl}$ are called small-field inflation. It turns out that large and small field models have drastically different properties and consequences, making $\Delta\phi$ a useful classifier of inflation models.

One particular feature of large-field inflation is its noticeable scale dependence, as shown by the well-known Lyth bound $\Delta\phi \sim \sqrt{2\epsilon} N_e M_{\rm Pl}$ in single-field slow-roll (SFSR) models [2]. Here ϵ is the first slow-roll parameter and $40 \lesssim N_e \lesssim 60$ is the e-folding number. A large $\Delta\phi$ ($\gtrsim M_{\rm Pl}$) implies a not-so-small ϵ, and consequently a relatively high inflation scale and a potentially observable tensor mode, since the tensor-to-scalar ratio $r = 16\epsilon$ in SFSR models. This conclusion survives in a much wider range of models beyond SFSR [3,4].

Recently there has been revived interest in large-field inflation, and in particular in whether a scalar field can consistently traverse a super-Planckian excursion within quantum gravity. In particular, the swampland distance conjecture (SDC) holds that a super-Planckian distance $\Delta\phi \gg M_{\rm Pl}$ always corresponds to a breakdown of low-energy effective field theory, with ϕ acting as a modulus controlling the mass of an infinite tower of light states (henceforth "tower states") with masses behaving as $m_{\rm tower} \sim \exp(-O(1) \times \Delta\phi/M_{\rm Pl})$ asymptotically [5,6]. This characteristic exponential behavior can be understood as a feedback effect from the loops of the tower states modifying the kinetic term of ϕ [7,8]. A controlled model of inflation should exist well away from this asymptotic region, which poses a familiar challenge for constructing realistic examples in string theory [9].

While large super-Planckian field excursions may be in tension with consistent quantum gravity, modest excursions $\Delta\phi \sim O(1) \times M_{\rm Pl}$ are less constrained [10], and arise in concrete examples such as axion monodromy [11–13]. Although most of the tower states are heavy in such a scenario, it is still expected that the masses of some states can change significantly as ϕ evolves over a Planckian range (which can even improve the behavior of the potential [14]). Thus, mild but visible scale dependence is a very generic feature of large-field inflation. It may well trigger the onset of dramatic events such as particle production [15] and phase transitions [16,17]. It is therefore desirable to have a more clear and direct probe of such scale dependence, as a generic indicator of large-field inflation.

In this paper, we exploit the recently developed cosmological collider (CC) observables to probe large-field inflation. The aim of CC physics is to look for heavy

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states with mass around or heavier than the inflation Hubble parameter H [18–23]. Such states were produced in a fast evolving background during inflation and then source a characteristic oscillatory signal, known as the CC signal, in the 3-point correlator of the curvature perturbation ζ [20]. One normally defines a scaleless and dimensionless shape function $\mathcal{S}(k_1, k_2, k_3)$ to describe the 3-point correlator, with k_i ($i = 1, 2, 3$) the magnitude of the three external momenta. Then, the CC signal appears when $k_1 \simeq k_2 \gg k_3$, the so-called squeezed limit, where we can write, ignoring a factor that might depend on $\mathbf{k}_1 \cdot \mathbf{k}_3$,

$$\mathcal{S}_{\rm signal}(k_1, k_2, k_3) \sim A\left(\frac{k_1}{k_3}\right)^{\alpha} \sin\left[\omega \log\left(\frac{k_1}{k_3}\right) + \delta\right]. \quad (1)$$

Here, A is the overall amplitude of the signal and (α, ω, δ) describe the shape dependence, which are in principle measurable from future CMB/LSS/21 cm observations [24–27] and also calculable from a given model [28,29]. For instance, with a tree-level exchange of a scalar field of mass $m > 3H/2$, we have $\alpha = -1/2$ and $\omega = \sqrt{m^2/H^2 - 9/4}$.

The main point of this paper is that the scale dependence in large-field inflation can be particularly significant for a massive spectator state (such as tower states) and this will make observables such as α and ω further dependent on k_1/k_3. It is not new that slow-roll correction can introduce scale dependence of $\mathcal{O}(\epsilon, \eta)$ for inflation correlators [30]. The new observation is that the direct inflaton coupling, such as the ϕ-dependent mass of the tower states, can introduce scale dependence of $\mathcal{O}(\epsilon^{1/2})$ and this can be a significant effect for large-field inflation. Therefore, measuring the k_1/k_3-dependence in α and ω can be a useful indicator of large-field inflation.

The recently updated constraint $r \lesssim 0.03$ by BICEP/Keck [31] translates to a bound in SFSR models, $16\epsilon \lesssim 0.03$. But we stress that this constraint on ϵ does not directly apply when there are spectator fields. In particular, the spectator fields appearing in our signals can suppress r by introducing more power to the scalar mode than to the tensor [32,33]. We will show below that the spectator field with time-dependent mass could bring a new correction to n_s in either direction, in addition to the n_s correction from constant-mass spectator fields considered in [32]. Thus, in principle, the observational constraints on n_s and r need to be reanalyzed for our model. At the same time, Ref. [32] shows that the relative correction to r is typically of $\mathcal{O}(1)$ in the parameter space we are interested in. Since our aim is to illustrate a physical effect instead of fitting a specific model to data, we will loosely take $16\epsilon < \mathcal{O}(0.1)$.

II. SCALE DEPENDENT INFLATION PARAMETERS

During the inflation, the Hubble parameter $H = H(t)$ has a weak time dependence characterized by the slow-roll parameters, including $\epsilon = -\dot{H}/H^2$ and $\eta = \dot{\epsilon}/(H\epsilon)$. In SFSR models, the inflaton ϕ has a nearly, but not exactly, constant speed $\dot{\phi}_0 = \dot{\phi}_0(t)$, again characterized by the slow-roll parameters. The time dependence in $\phi_0(t)$ and $H(t)$ can easily introduce a time-dependent mass to a spectator field σ. So time-dependent masses in inflation are rather generic.

Therefore, choosing a reference time $t_* = 0$,

$$H(t) = H_*\left[1 - \epsilon_* H_* t + \left(\epsilon_*^2 - \frac{1}{2}\epsilon_*\eta_*\right)(H_* t)^2 + \cdots\right]. \quad (2)$$

Using the relation $\dot{\phi}_0^2 \simeq 2\epsilon M_{\rm Pl}^2 H^2$, we get

$$\dot{\phi}_0(t) = \dot{\phi}_{0*}\left[1 + \left(\frac{1}{2}\eta_* - \epsilon_*\right)H_* t + \cdots\right]. \quad (3)$$

Using these relations we can find the time dependence of the masses induced by some familiar interactions. For example, the spectator can have a nonminimal coupling to the Ricci scalar $-\xi R\sigma^2$. Then, using $R = 12H^2 + 6\dot{H}$, we have $m_\sigma^2 = m_{\sigma*}^2(1 - 2\epsilon_* H_* t + \cdots)$ with $m_{\sigma*}^2 = 24\xi H_*^2$. Take another example of $(\partial\phi)^2\sigma^2/\Lambda^2$; we have $m_\sigma^2 = m_{\sigma*}^2[1 + (\eta_* - 2\epsilon_*)H_* t + \cdots]$ with $m_{\sigma^*}^2 = 2\dot{\phi}_{0*}^2/\Lambda^2$.

In all such examples the time dependence is weak, in the sense of being $\mathcal{O}(\epsilon, \eta)$. Therefore, to see an $\mathcal{O}(1)$ effect, we would need the number of e-folds $N \simeq \min\{1/\epsilon, 1/\eta\} \sim \mathcal{O}(30)$. Practically, we have access only to $\mathcal{O}(10)$ e-folds from CMB and LSS, so it is difficult to see this time dependence directly.

Stronger time dependence is possible if we consider nonderivative couplings of the form $m_{\sigma*}^2 f(\phi/M_{\rm Pl})\sigma^2$. Such a coupling breaks the approximate shift symmetry of the inflaton field, which is expected to be broken at least by quantum gravity effects. The evolution of the inflaton field leads to $m_\sigma^2(t) = m_{\sigma*}^2 f'(1 - \sqrt{2\epsilon_*}H_* t + \cdots)$. In this case we only need $\mathcal{O}(\epsilon^{-1/2})$ e-folds to see the effect. For large-field inflation the required e-folds could be $\mathcal{O}(10)$, and thus possible for CMB/LSS observation. We note that this coupling is small, which would not change the inflaton dynamics. For the same reason, it would not generate a sizable non-Gaussian signal. At the same time, there could certainly be other, shift-symmetry preserving, couplings which can generate an observable signal. The effect of this nonderivative coupling would be modulating the shape of the signal, as will be detailed below.

For large-field models, a rather generic interaction exists, namely the "tower states" motivated by the SDC, with mass $\propto e^{-\alpha\phi/M_{\rm Pl}}$, where α is an $\mathcal{O}(1)$ number. With rolling background $\phi(t)$, this introduces a time-dependent mass $m_\sigma^2 = m_{\sigma*}^2 e^{-2\alpha\sqrt{2\epsilon_*}Ht}$ which is again an $\mathcal{O}(\epsilon^{1/2})$ effect. Although this exponential ansatz is best motivated in asymptotic regions of the potential, the scale-dependent frequency and size of the resulting signal is largely independent of the particular ansatz we choose, so long as the spectator couples directly to the inflaton via nonderivative couplings. Hence, we focus on this particular ansatz as an example, and consider alternative choices in the supplemental material [34].

III. QUASISINGLE (LARGE) FIELD INFLATION

As a concrete example we study quasi-single-field inflation [18], where we have a scalar field σ with mass $m_\sigma \sim \mathcal{O}(H)$ coupled to the inflaton through $(\partial_\mu\phi)^2\sigma$. At the same time, σ has a cubic self-coupling $\lambda\sigma^3$. The σ self-coupling is not constrained by the slow-roll condition and can *a priori* be large (namely $\lambda/H \sim 1$ or larger). So the bispectrum from this cubic coupling can be potentially large. This is a well-studied model with large non-Gaussianity. Below, we will consider the effect of inflaton modulation in this model.

As discussed before, we assume that the inflaton modulates the mass m_σ through an exponential factor. Then, the Lagrangian is given by

$$\mathcal{L} = -\frac{1}{2}(\partial_\mu\phi)^2 - V(\phi) - \frac{1}{2}(\partial_\mu\sigma)^2 - \frac{1}{2}e^{-2\alpha\phi/M_{\mathrm{Pl}}}m^2\sigma^2 - \frac{1}{6}\lambda_3\sigma^3 + \frac{1}{2}\lambda_5(\partial_\mu\phi)^2\sigma. \quad (4)$$

This model is certainly not complete, as the σ potential is not stable, and the inflaton rolling introduces a tree-level tadpole to σ. But we can still treat it as an effective description of a more complete model, in which the σ field is stabilized at some background value σ_0, for example, by an approximate cancellation between the terms proportional to λ_3 and λ_5, and we are simply expanding the full action around σ_0. We also note that (4) only includes the operators of lowest dimensions for generating the process (16). Corrections from omitted higher dimensional operators are thus negligible so long as the model remains perturbative. We further elaborate on the validity of this Lagrangian in the supplemental material [34].

With a rolling background $\phi_0(t) \simeq \phi_{0*} - |\dot\phi_{0*}|(t - t_*) = \sqrt{2\epsilon}M_{\mathrm{Pl}}H_*(t - t_*)$ (we assume $\dot\phi_0 < 0$ without loss of generality), the ϕ-dependence gets translated to a time dependence,

$$m_{\mathrm{eff}}(t) = m_* e^{-\alpha\sqrt{2\epsilon}Ht}, \quad (5)$$

where t is the physical time. Here and below we remove the $*$ in ϵ and H. Decomposing the σ field into Fourier modes with fixed 3D momentum $\mathbf{k}$, we can find the equation of motion for a single mode $\sigma_{\mathbf{k}}$ as

$$\sigma_{\mathbf{k}}'' - \frac{2}{\tau}\sigma_{\mathbf{k}}' + (k^2 + m_*^2|H\tau|^{2\alpha\sqrt{2\epsilon}-2})\sigma_{\mathbf{k}} = 0, \quad (6)$$

where a prime denotes the derivative with respect to conformal time τ, which is related to t via $e^{Ht} = a = -1/(H\tau)$. With the usual Bunch-Davies initial condition for the mode, the above equation has a unique solution up to an irrelevant overall phase.

We are not aware of any named special functions that solve this equation, but it is possible to proceed analytically by making approximations. For example, when the mode leaves the horizon at late times ($|k\tau| \ll 1$) and when the mass $m(\tau)$ is still much greater than Hubble, we can apply the Wentzel-Kramers-Brillouin (WKB) approximation. Rewriting $\sigma_{\mathbf{k}}(\tau) = \tau\chi_{\mathbf{k}}(\tau)$, the equation for χ_k is

$$\chi_{\mathbf{k}}'' + (k^2 + m_*^2|H\tau|^{2\alpha\sqrt{2\epsilon}-2} - 2\tau^{-2})\chi_{\mathbf{k}} = 0. \quad (7)$$

At late times, we can neglect the momentum, and the WKB solution is

$$\chi_{\mathbf{k}}(\tau) \simeq Ae^{+\mathrm{i}\vartheta(\tau)} + Be^{-\mathrm{i}\vartheta(\tau)}. \quad (8)$$

When the effective mass $m_{\mathrm{eff}} = m_*|H\tau|^{\alpha\sqrt{2\epsilon}} \lesssim H$, the WKB solution stops oscillating and becomes invalid. Therefore, the above approximation works only when $m_{\mathrm{eff}} \gg H$, in which case the phase $\vartheta(\tau)$ can be approximated by

$$\vartheta(\tau) \simeq \frac{1}{\alpha\sqrt{2\epsilon}}\frac{m_*}{H}(|H\tau|^{\alpha\sqrt{2\epsilon}} - 1) = \frac{m_*}{H}\log|H\tau|\left(1 + \alpha\sqrt{\frac{\epsilon}{2}}\log|H\tau| + \mathcal{O}(\epsilon)\right). \quad (9)$$

Therefore, the oscillation frequency of the mode function in the late-time limit is itself time dependent, just as expected. We also solve the mode equation (6) numerically, and show the solutions with several choices of parameters in Fig. 1, in

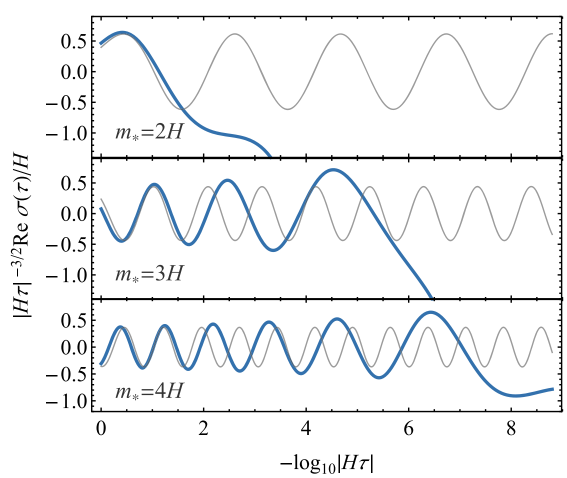


FIG. 1. The real part of the mode functions with time dependent masses (5) (dark blue). In all plots we take $\alpha = 1$ and $\epsilon_* = 0.02$. For comparison, the corresponding constant-mass mode functions are shown in gray curves.

which the mild time dependence of the oscillation frequency is evident.

IV. POWER SPECTRUM IN LARGE-FIELD MODELS

It is important to check the scale dependence of the power spectrum induced by the σ-φ coupling. This coupling arises from the λ_5 term in (4) evaluated with the inflation background,

$$\Delta\mathcal{L} = a^3\mu\varphi'\sigma, \tag{10}$$

where $\mu \equiv -\lambda_5\dot{\phi}_0$. A full calculation of the power spectrum should include slow-roll corrections everywhere consistently, including the coefficients of $\sigma'_{\mathbf{k}}$ and $\varphi'_{\mathbf{k}}$ in the mode equation, together with the slow-roll induced-mass correction to $\sigma_{\mathbf{k}}$. However, in order to isolate the effect of a time-dependent mass m_{eff}, we only retain the time dependence in m_{eff}, and assume a de Sitter limit for the background evolution. The resulting power spectrum then reflects the scale dependence from $m_{\mathrm{eff}}(\tau)$.

We numerically evolve the following set of equations, with the usual Bunch-Davies initial condition.

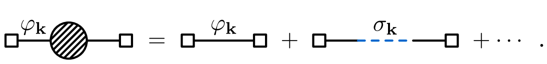
$$\sigma''_{\mathbf{k}} - \frac{2}{\tau}\sigma'_{\mathbf{k}} + (k^2 + m_*^2|H\tau|^{2\alpha\sqrt{2\epsilon}-2})\sigma_{\mathbf{k}} = \frac{-\mu}{H\tau}\varphi_{\mathbf{k}}, \tag{11}$$

$$\varphi''_{\mathbf{k}} - \frac{2}{\tau}\varphi'_{\mathbf{k}} + k^2\varphi_{\mathbf{k}} = -\frac{3\mu}{H\tau^2}\sigma_{\mathbf{k}} + \frac{\mu}{H\tau}\sigma'_{\mathbf{k}}. \tag{12}$$

The power spectrum P_ζ is then computed by

$$P_\zeta(k) = (H/\dot{\phi}_0)^2(k^3/2\pi^2)\langle\varphi_{-\mathbf{k}}(\tau_f)\varphi_{\mathbf{k}}(\tau_f)\rangle', \tag{13}$$

where the prime $\langle\cdots\rangle'$ means the momentum-conserving δ-function is removed, and τ_f is chosen so that $|k\tau_f| \ll 1$ for all relevant k. This is effectively computing the following set of diagrams:



(14)

In our numerical result we do not observe any oscillatory signal, but only a smooth scale dependence, as expected. So, the scale dependence from the time-dependent mass is degenerate with the scale dependence from the slow-roll potential, and is not an independent observable. However, we do need to check that the new scale dependence from the heavy-field mixing is not much larger than the observed value, to avoid significant tuning between the background contribution and the massive-state correction. The scalar tilt $n_s - 1 \equiv \mathrm{d}\log P_\zeta/\mathrm{d}\log k$ induced by the time-dependent mass is shown in Fig. 2. This result is to be added with the slow-roll contributions, like those considered in [32], to get the observed value $n_s^{(\mathrm{CMB})} \simeq 0.967$ [1]. We see that the correction Δn_s in Fig. 2 stays at the same order as the slow-roll contribution $n_s^{(\mathrm{CMB})} - 1 \simeq -0.033$ for most of the parameter space. Therefore, the scale dependence from the time-dependent mass is compatible with the current observations for most of the parameter space we are interested in.

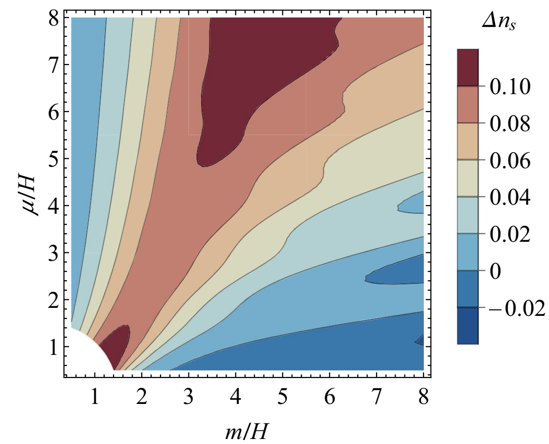
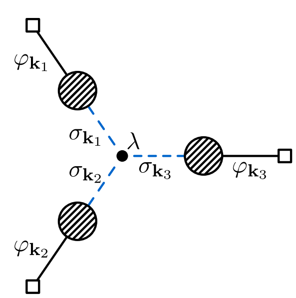


FIG. 2. The correction to scalar tilt, Δn_s, from an intermediate state with time-dependent mass given in (5), plotted on the plane of σ mass m at a reference scale and the mixing parameter μ between σ and $\delta\phi$. In this plot we fix $\epsilon = 1.25\times 10^{-3}$, which corresponds to $r = 0.02$ (< 0.02) when $\mu = 0$ (> 0). The lower-left corner with $\sqrt{m^2+\mu^2} < 3H/2$, where no oscillation signals occur, is excluded.

V. THE SQUEEZED BISPECTRUM

Now we study the consequence of the scale-dependence σ mass in the 3-point correlator mediated by σ. This is a well-studied "discovery channel" for the CC, and it was found that the largest signal comes from the following diagram, where the blobs again include arbitrary numbers of two-point mixing insertions [35,36].



(15)

The diagram (15) can be computed in the following way

$$\langle\varphi_{\mathbf{k}_1}\varphi_{\mathbf{k}_2}\varphi_{\mathbf{k}_3}\rangle' = 2\lambda_3\mathrm{Im}\int_{-\infty}^{0}\mathrm{d}\tau\, a^4\prod_{i=1}^{3}\langle\sigma_{\mathbf{k}_i}(\tau)\varphi_{-\mathbf{k}_i}(\tau_f)\rangle, \tag{16}$$

from which we can get the shape function $\mathcal{S}$ by

$$\mathcal{S}(k_1,k_2,k_3) = \frac{(k_1k_2k_3)^2}{2\pi P_\zeta^{1/2}H^3}\langle\varphi_{\mathbf{k}_1}\varphi_{\mathbf{k}_2}\varphi_{\mathbf{k}_3}\rangle'. \tag{17}$$

We will first try to understand the qualitative features of this process by taking analytical approximations, and then present full numerical results.

For a crude analytical understanding, we focus on the oscillatory signal in the squeezed limit $k_3 \ll k_1 \simeq k_2$ and consider the perturbative regime $\mu \ll m$ and $m/H \gg 1$. The signal with constant mass can be estimated as [36,37]

$$\mathcal{S}_{\rm signal} \simeq \frac{\lambda_3\mu^3}{\nu^{5/2}H^4}e^{-\pi\nu}\sqrt{\frac{k_3}{k_1}}\sin\left(\nu\log\frac{k_3}{k_1}+\vartheta\right), \tag{18}$$

where $\nu \equiv \sqrt{m^2/H^2-9/4}$. We emphasize that, for constant m, the signal has a fixed oscillation frequency ν and its size has a fixed scaling with the momentum ratio, characterized by the power law $(k_3/k_1)^{1/2}$.

When the mass has a weak time dependence, we can still use (18), but this time evaluate $m = m(\tau)$ at the saddle point $\tau \simeq -\nu/(2k_1)$ of the integrand $\sim e^{\pm 2ik_1\tau}(-\tau)^{\pm i\nu}$. Consequently, the time dependence of m is translated to a scale dependence of ν in (18), which appears in both the signal frequency and in the overall size, and this immediately leads to a main result of this work. The CC signal from a time-dependent mass has a scale-dependent oscillation frequency, and its overall size scales with k_3/k_1 differently from the power law $(k_3/k_1)^{1/2}$. The deviation in the scaling with k_3/k_1 is quite significant due to the exponential factor $e^{-\pi\nu}$, which makes the signal very sensitive to the mass m.

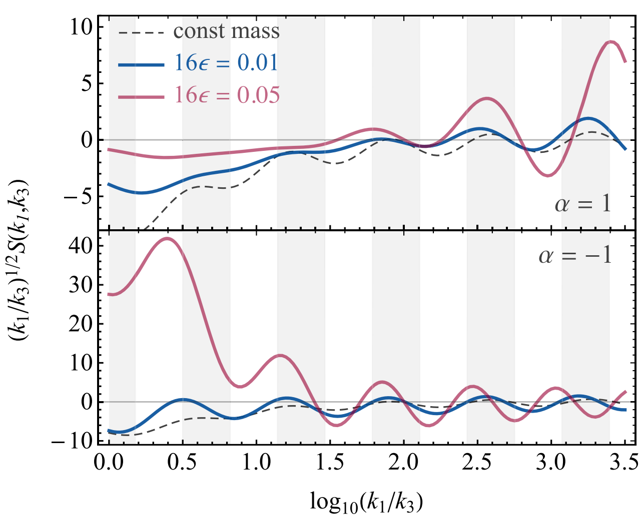


FIG. 3. The shape function (17) from the process (15). In both panels we take $m_* = \mu = 3H$ and $\lambda_3 = H$. The gray shades indicate the expected frequency $\nu_\mu = \sqrt{m_*^2+\mu^2}$ for constant mass in the squeezed limit. The blue and red curves show the signals from time-dependent masses, in which both the amplitude and the frequency of the oscillations change visibly with k_1/k_3.

In the strongly-coupled regime ($\mu \gtrsim m$), no closed analytical estimate like (18) is known. However, it is known that the oscillation frequency of the signal is modified to $\nu \to \nu_\mu \equiv \sqrt{(m^2+\mu^2)/H^2-9/4}$ [32]. We anticipate that the signal size is sensitive to m so long as μ is not too much greater than m. Therefore, our main result, namely, the scale-dependent oscillation frequency and the signal size, still holds.

The above analytical arguments rely on approximations in several limits which are never really reached by realistic parameters. Therefore a numerical approach is indispensable to get the signal shape precisely. We compute (16) numerically and get the full bispectrum. Several examples are shown in Fig. 3. We show the signals with the mass decreasing/increasing in the upper/lower panel, respectively. Compared with the bispectrum of constant mass (dashed curves), the slow changes of the frequency and the amplitude are evident in both cases.

VI. DISCUSSIONS

Large-field inflation models typically show significant scale dependence that can distort the CC signals. In this paper we show that a direct coupling between a massive state and the inflaton can lead to an oscillatory signal in the bispectrum, with significant and nonstandard momentum-ratio dependence in both the size and the frequency. We use quasisingle-field inflation and an exponential ansatz for the mass to illustrate the point, but we stress that the signal distortion depends mainly on the large-field excursion and the direct coupling. Thus, we expect that similar effects should also show up in a broader class of models with these ingredients.

Phenomenologically, our work shows that the momentum-ratio dependence of the signal size and frequency can be informative probes of rich dynamics during inflation. (See Ref. [38] for a related example.) Therefore, our work invites efforts on developing more realistic templates allowing for the possibility that the signal size and the frequency could be momentum-ratio dependent.

There are other sources of scale dependence, including the slow-roll background, as explored in [37,39,40]. We expect the effect to be weaker than ours in large-field inflation, as detailed before. There is also a known slow change of signal frequency in the not-so-squeezed configurations. This change is automatically included in a numerical approach, and can also be resolved analytically by pushing the calculation to higher orders in the momentum ratio.

In this paper we only considered the time-dependent mass. It would be interesting to explore the similar time dependences in other parameters. In particular, a time-dependent two-point mixing parameter μ could induce a similar change in the signal frequency. It would also be interesting to incorporate these nonstandard momentum-ratio dependences in template-based Fisher forecasts for future observations. We leave these topics for future studies.

ACKNOWLEDGMENTS

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