

Cosmological non-Gaussianity from the neutrino seesaw mechanism

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The neutrino mass generation via conventional seesaw mechanism is realized at high scales around $O(10^{14})$ GeV with natural Yukawa couplings of $O(1)$, making the test of the neutrino seesaw a great challenge. It is intriguing to note that the neutrino seesaw scale is typically around the upper range of the cosmological inflation scale. In this work, we propose a new framework incorporating inflation and the neutrino seesaw in which the inflaton primarily decays into right-handed neutrinos after inflation. This decay process is governed by the inflaton interaction with the right-handed neutrinos that respects the shift symmetry. Through the neutrino seesaw mechanism, fluctuations of the Higgs field can modulate the inflaton decays and contribute to the curvature perturbation. We investigate the induced non-Gaussian signatures and demonstrate, for the first time, that such signatures provide an important means to directly probe the high scale of natural neutrino seesaw. We further analyze the interplay of the non-Gaussianity signatures with the low-energy neutrino experiments, and their interplay with the Higgs self-coupling measurements at the LHC and future colliders.

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I. INTRODUCTION

The discovery of neutrino oscillations has pointed to tiny but nonzero neutrino masses of $O(0.1)$ eV, which can be naturally generated by including the right-handed neutrinos within the established structure of the Standard Model (SM) of particle physics. These right-handed neutrinos are the chiral partners of the left-handed neutrinos and they join together the Yukawa interactions with the Higgs doublet (just like any other leptons and quarks in the SM), but the right-handed neutrinos are pure singlets of the SM gauge group. As such they can naturally acquire large Majorana masses (M_R) and realize the seesaw mechanism [1,2] to naturally generate the tiny neutrino masses $m_\nu \sim v^2/M_R$, where $v = O(100)$ GeV denotes the vacuum expectation

value of the SM Higgs doublet and the neutrino-Higgs Yukawa couplings (y_ν) are set to their natural values of $O(1)$. This generally predicts a high scale for neutrino seesaw, $M_R \sim v^2/m_\nu = O(10^{14})$ GeV.¹ Hence, given the

¹For the conventional SM setup before 1998, the neutrinos were assumed for simplicity to be massless and have only left-handed components because the SM is structured to have all the right-handed fermions be weak singlets for each fermion family, where the right-handed neutrinos (N_R) are pure gauge singlets and their absence does not affect the gauge anomaly cancellation of the SM. Weinberg realized [3] that without N_R , the left-handed neutrinos can acquire small Majorana masses from a gauge-invariant dimension-five operator ($LLHH$) which is suppressed by a large UV cutoff scale $\Lambda_\nu \sim v^2/m_\nu$, far beyond the weak scale. This dimension-five operator is, however, nonrenormalizable and its “minimal UV completion” is given by the conventional seesaw [1,2] with $\Lambda_\nu = M_R$ after adding back N_R for each fermion family. The existence of the right-handed neutrinos is predicted by the SM structure and provides the minimal UV completion for the dimension-5 Weinberg operator [3] through the seesaw mechanism (naturally generating the light neutrino masses), yet, the right-handed neutrinos point to a brand-new seesaw scale $\Lambda_\nu \sim v^2/m_\nu$ that is beyond the SM. Hence, it is extremely important to probe the right-handed neutrinos as the last missing piece of the SM and test the neutrino mass generation via the seesaw mechanism.

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current capabilities of particle physics experiments, probing the natural neutrino seesaw mechanism at such high scales (M_R) poses a great challenge.

On the other hand, it is believed that the early Universe underwent an inflationary epoch, during which the Universe expanded exponentially over a very short period. Inflation not only resolves the flatness and horizon problems, but also seeds the primordial fluctuations that form the large-scale structure of the Universe. The energy scale of inflation could be as high as $O(10^{16})$ GeV, characterized by the nearly constant Hubble parameter H_{inf} which is typically around 10^{14} GeV, providing an important window for probing new physics at high-energy scales. It is intriguing to observe that both the neutrino seesaw scale and the Hubble parameter during inflation can be realized around the same scale of $O(10^{14})$ GeV [4,5]. In the minimal setup, the inflation is triggered by a scalar inflaton field. The primordial fluctuations are generated by quantum fluctuations of the inflaton and can be directly measured through the cosmic microwave background (CMB). The current CMB data indicate that these fluctuations are adiabatic and Gaussian [5–8].

However, inflation could also generate non-Gaussianity (NG) in the primordial perturbations [9,10], as characterized by n -point ($n \geq 3$) correlation functions of the comoving curvature perturbation ζ . The primordial non-Gaussianity can arise in models of multifield inflation or single-field inflation with interactions, offering an ideal opportunity to probe the relevant new physics at high-energy scales. One notable example is the “cosmological collider” method [11–14], which aims to explore high-scale particle physics by studying the non-Gaussian properties of large-scale structures. Given the present nonobservation of non-Gaussianity, the existing CMB measurements have already set constraints on the non-Gaussian parameter $f_{\text{NL}} \lesssim O(10)$ depending on the non-Gaussian shapes under consideration. The future detection of the 21 cm tomography could eventually reach the sensitivity down to the level of $f_{\text{NL}} = O(0.01)$ [15–17].

At the end of inflation, the inflaton would oscillate at the bottom of its potential and eventually transfer its energy to the SM particles, thereby reheating the Universe. This sets the stage for the transition of the Universe from the inflationary epoch to a radiation-dominated period, preceding the onset of big bang nucleosynthesis. While observations of large-scale structures provide information about the inflation, the dynamics of the reheating process remain unclear. It is natural to expect that the inflaton couples directly to the right-handed neutrinos and predominantly decays into them after inflation. Then, these right-handed neutrinos could further decay into the SM particles through Yukawa interactions and thus complete the reheating process. It is appealing that this approach also naturally provides an initial setup for the leptogenesis [18], which generates sufficient right-handed neutrinos after reheating.

On the other hand, during inflation, not only does the inflaton fluctuate, but other light scalar fields also undergo fluctuations. In particular, the Higgs boson would acquire a field value near the Hubble scale, which varies across different horizon patches. This variation leads to discrimination on the right-handed neutrino masses in local regions of the Universe via seesaw mechanism. In consequence, the inflaton’s decay rate into right-handed neutrinos is modulated by the Higgs field value. This modulated reheating scenario provides a source of primordial curvature perturbations [19]. The associated non-Gaussian signatures open up a new window for probing the neutrino seesaw scale.

In this work, we propose a new framework incorporating inflation and neutrino seesaw in which the inflaton primarily decays into right-handed neutrinos after inflation. This decay process is governed by the inflaton interaction with the right-handed neutrinos that respects the shift symmetry. Through the neutrino seesaw mechanism, fluctuations of the Higgs field can modulate the inflaton decays, contributing to the curvature perturbations. We investigate the effects of Higgs-modulated reheating and the associated non-Gaussianity (bispectrum). We demonstrate its potential to probe the high-scale neutrino seesaw mechanism. We further analyze the interplay of the non-Gaussianity signatures with the low-energy neutrino experiments, and their interplay with the Higgs self-coupling measurements at the LHC and future colliders. In passing, this approach also provides a new framework of the cosmological Higgs collider, in which the Higgs-modulated reheating is naturally realized by the inflaton decays into right-handed neutrinos within the neutrino seesaw. Thus, particles that couple to the Higgs field would induce cosmological collider signatures, which we may call the neutrino-assisted cosmological Higgs collider (NCHC).²

This paper is organized as follows. In Sec. II, we discuss the dynamics and evolution of the Higgs field during and after inflation. In Sec. III, we present a minimal framework incorporating inflation and neutrino seesaw, in which the inflaton decay is modulated by the Higgs boson through right-handed neutrinos. Then, we give the model realization and setup in Sec. III A and analyze the curvature perturbation from the Higgs-modulated reheating through right-handed neutrinos in Sec. III B. For Sec. IV, we first study the comoving curvature perturbation from Higgs-modulated reheating in Sec. IV A. Then, we present the systematic analysis on the three-point correlation function (bispectrum) of the comoving curvature perturbation in Sec. IV B. With these, we study the probe of neutrino

²This NCHC scenario differs from the previous cosmological Higgs collider study in the literature [20] in which the inflaton is assumed to couple to certain newly added singlet scalar fields and predominantly decay into these scalars.

seesaw parameter space from the non-Gaussianity in Sec. IV C and the dependence of non-Gaussianity on the Higgs self-coupling in Sec. IV D. Finally, we conclude in Sec. V. Appendixes A–E provide the necessary formulas and technical derivations to support the analyses in the main text.

II. DYNAMICS OF HIGGS FIELD IN THE EARLY UNIVERSE

In this section, we discuss the physics of the Higgs field during and after inflation. The Lagrangian density of the SM Higgs kinetic term and potential term is given by

$$\mathcal{L} = \sqrt{-g} \left[-g_{\mu\nu} D^\mu \mathbb{H}^\dagger D^\nu \mathbb{H} + \mu^2 \mathbb{H}^\dagger \mathbb{H} - \lambda (\mathbb{H}^\dagger \mathbb{H})^2 \right], \quad (2.1)$$

where $\mathbb{H}$ is the SM Higgs doublet containing four independent scalar components. Note that the Higgs self-coupling might become negative at a scale around 10^{11} GeV depending on the precise value of the measured top quark mass [21]. Given that the Hubble parameter during inflation could be as high as 10^{14} GeV, the Higgs vacuum might become unstable during inflation, but the running of Higgs self-coupling is very sensitive to the measured top quark mass. Within the 3σ range of the current top mass measurement [22], it is still possible to keep the Higgs coupling positive and have a value of $O(0.01)$ at the inflation scale.³

In the following discussions, we denote the quantities at different epochs by using the corresponding subscripts, such as $A_{\text{inf}} \equiv A(t = t_{\text{inf}})$ and $A_{\text{reh}} \equiv A(t = t_{\text{reh}})$. Here t_{inf} is the physical time at the end of inflation, and t_{reh} is the physical time at the completion of reheating. As the effects of slow-roll parameters are fairly small, the Hubble parameter remains constant throughout the entire epoch of inflation. Thus, H_{inf} could be used to represent the Hubble parameter during inflation.

A. Dynamics of Higgs field during inflation

During inflation, the Universe exponentially expands and can be described by the de Sitter spacetime if the slow roll of the inflaton is neglected. In contrast to the inflaton, a massless spectator scalar field with self-interaction will exhibit infrared (IR) divergences in de Sitter spacetime [24–26]. The framework of stochastic inflation [27,28] provides a systematic approach to deal with the IR behavior for the superhorizon mode of the massless spectator field.

Although the SM Higgs doublet $\mathbb{H}$ contains four real scalar components, three of them correspond to the Goldstone modes that become the longitudinal components of the SU(2) weak gauge bosons. During inflation, the

fluctuation of the Higgs field is on the order of the Hubble parameter, which means the masses of the weak gauge bosons are also of the order of the Hubble scale and thus rather large. On the other hand, the CP -even component h of the Higgs doublet $\mathbb{H}$ is much lighter due to the smallness of Higgs self-coupling in comparison with the weak gauge coupling, $\lambda \ll g^2$. Hence for the present study we only need to deal with the light Higgs field h . In the unitary gauge, the Higgs doublet takes the following form:

$$\mathbb{H} = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ h \end{pmatrix}. \quad (2.2)$$

Thus, the Lagrangian density of the pure Higgs sector can be expressed as follows:

$$\mathcal{L} = \sqrt{-g} \left[-\frac{1}{2} g_{\mu\nu} (\partial^\mu h \partial^\nu h) - V(h) \right], \quad (2.3a)$$

$$V(h) = -\frac{1}{2} \mu^2 h^2 + \frac{\lambda}{4} h^4. \quad (2.3b)$$

In the above, $V(h)$ is the Higgs potential, in which the quadratic mass term $\frac{1}{2} \mu^2 h^2$ could be omitted for a large value of h during the inflation.

The Higgs field h can be decomposed into a long-wavelength mode (h_L) and the short-wavelength modes [which include contributions above a physical cutoff scale $\epsilon a(t)H$]. They are both generated by quantum fluctuations,

$$h(\mathbf{x}, t) = h_L(\mathbf{x}, t) + \int \frac{d^3 k}{(2\pi)^3} \theta(k - \epsilon a(t)H_{\text{inf}}) \times \left[a_{\mathbf{k}} h_{\mathbf{k}}(t) e^{-i\mathbf{k} \cdot \mathbf{x}} + a_{\mathbf{k}}^\dagger h_{\mathbf{k}}^*(t) e^{i\mathbf{k} \cdot \mathbf{x}} \right], \quad (2.4)$$

where ϵ is a small parameter such that the short-wavelength modes satisfy the massless Klein-Gordon equation in the de Sitter space, which has the following solution:

$$h_{\mathbf{k}} = \frac{H_{\text{inf}}}{\sqrt{2k^3}} (1 + i k \tau) e^{-i k \tau}, \quad (2.5)$$

where $\tau = -1/(aH)$ is the conformal time.

In this framework, the short-wavelength modes $h_{\mathbf{k}}(t)$ are initially subhorizon and correspond to the normalized modes of a massless scalar field in the de Sitter spacetime. As the Universe expands, these modes are stretched and eventually cross the physical cutoff $\epsilon a(t)H$, transitioning into superhorizon modes h_L . The long-wavelength, superhorizon modes h_L can be effectively treated as a classical stochastic field, following the Langevin equation

$$\dot{h}_L(\mathbf{x}, t) = -\frac{1}{3H_{\text{inf}}} \frac{\partial V}{\partial h_L} + f(\mathbf{x}, t). \quad (2.6)$$

It shows that the evolution of the long-wavelength modes is driven by an effective stochastic “force” $f(\mathbf{x}, t)$, which is

³Adding additional light scalar particle(s) to the Higgs sector at weak scale could lift the Higgs self-coupling to the level of $O(0.1)$ at the inflation scale [23]. For the current study, we will choose the minimal SM Higgs sector.

generated by the “freezing out” of short-wavelength modes,

$$f(\mathbf{x}, t) = \int \frac{d^3k}{(2\pi)^3} \delta(k - \epsilon a(t) H_{\text{inf}}) \epsilon a(t) \times H_{\text{inf}}^2 \left(a_{\mathbf{k}} h_{\mathbf{k}} e^{-i\mathbf{k} \cdot \mathbf{x}} + a_{\mathbf{k}}^\dagger h_{\mathbf{k}}^* e^{i\mathbf{k} \cdot \mathbf{x}} \right), \quad (2.7)$$

where we have used the equation $da(t)/dt = H_{\text{inf}} a(t)$, and the two-point correlation function of the stochastic noise $f(\mathbf{x}, t)$ is given by

$$\langle f(\mathbf{x}_1, t_1) f(\mathbf{x}_2, t_2) \rangle = \frac{H_{\text{inf}}^3}{4\pi^2} \delta(t_1 - t_2) j_0(\epsilon a(t_1) H_{\text{inf}} |\mathbf{x}_1 - \mathbf{x}_2|), \quad (2.8)$$

where $j_0(z) = (\sin z)/z$. Hence, the behavior of the Higgs field on superhorizon scales can be described as a classical stochastic process with a probability distribution ρ satisfying a Fokker-Planck equation,

$$\frac{\partial \rho[h(\mathbf{x}, t)]}{\partial t} = \frac{1}{3H_{\text{inf}}} \frac{\partial}{\partial h} \left\{ \rho[h(\mathbf{x}, t)] \frac{\partial}{\partial h} V[h(\mathbf{x}, t)] \right\} + \frac{H_{\text{inf}}^3}{8\pi^2} \frac{\partial^2}{\partial h^2} \rho[h(\mathbf{x}, t)]. \quad (2.9)$$

The last term on the right-hand side represents the effect of stochastic noise, originating from the subhorizon modes of the Higgs field as they cross the horizon. This term encapsulates the quantum nature of the fluctuations.

To obtain the probability distribution $\rho(h, t)$, we expand it in terms of a set of eigenfunctions $\Psi_n(h)$ as follows:

$$\rho(h, t) = \Psi_0(h) \sum_{n=0}^{\infty} a_n \Psi_n(h) \exp(-\Lambda_n t), \quad (2.10)$$

where Λ_n and $\Psi_n(h)$ denote the corresponding eigenvalues and eigenfunction of the following differential equation:

$$\tilde{D}_h \Psi_n(h) = -\frac{8\pi^2 \Lambda_n}{H_{\text{inf}}^3} \Psi_n(h). \quad (2.11)$$

In the above, the operator $\tilde{D}_h$ is defined as follows:

$$\tilde{D}_h = \frac{\partial^2}{\partial h^2} - \left[\left(\frac{\partial v(h)}{\partial h} \right)^2 - \frac{\partial^2 v(h)}{\partial h^2} \right], \quad (2.12)$$

where $v(h) \equiv [4\pi^2/(3H_{\text{inf}}^4)]V(h)$ and $V(h) = \frac{\lambda}{4} h^4$ is the SM Higgs potential. The eigenfunctions are orthonormalized as follows:

$$\int_{-\infty}^{+\infty} dh \Psi_n(h) \Psi_{n'}(h) = \delta_{n,n'}. \quad (2.13)$$

The derivation of the above equations is provided in Appendix E. We note that all the eigenvalues are non-negative and increase with the index n , and the lowest eigenvalue vanishes ($\Lambda_0 = 0$).

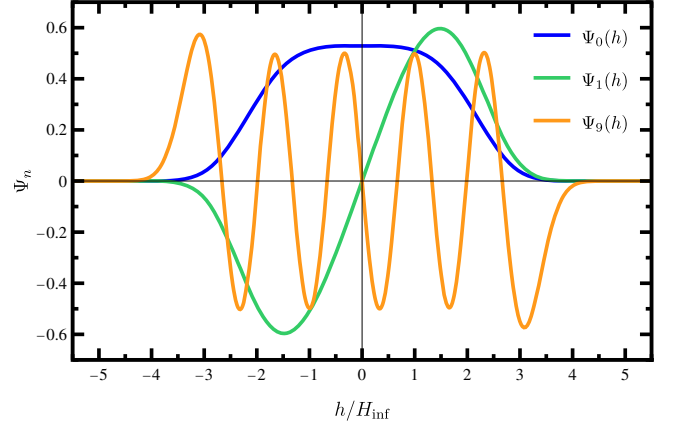


FIG. 1. Shape of the eigenfunction Ψ_n as a function of Higgs field h , where we choose the eigenvalue index $n = 0, 1, 9$ for illustration. The SM Higgs self-coupling constant is set as $\lambda = 0.01$.

In Fig. 1, we plot the shape of the eigenfunction $\Psi_n(h)$ as a function of the Higgs field h , and we choose the eigenvalue index $n = 0, 1, 9$ for illustration. The eigenvalue equation of Eqs. (2.11) and (2.12) is a Sturm-Liouville problem [29] and the eigenfunction $\Psi_n(h)$ is proved to have n zero points and $n + 1$ extreme points.

If inflation lasts for enough time, the only remaining eigenfunction is the ground state eigenfunction $\Psi_0(h)$ corresponding to the eigenvalue $\Lambda_0 = 0$. Consequently, the probability distribution of the long-wavelength modes asymptotically approaches that of an equilibrium state,

$$\rho_{\text{eq}}(h) = \frac{2\lambda^{1/4}}{\Gamma(1/4)} \left(\frac{2\pi^2}{3} \right)^{1/4} \exp\left(\frac{-2\pi^2 \lambda h^4}{3H_{\text{inf}}^4} \right), \quad (2.14)$$

where the normalization is imposed,

$$\int_{-\infty}^{+\infty} dh \rho_{\text{eq}}(h) = 1. \quad (2.15)$$

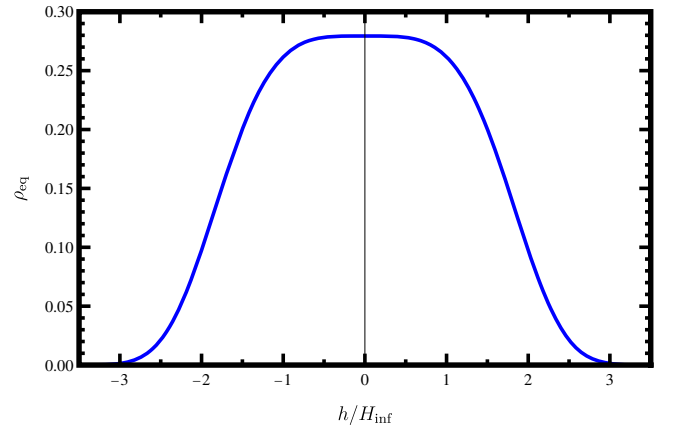


FIG. 2. Probability density distribution ρ_{eq} as a function of the Higgs field h at the end of inflation. The SM Higgs self-coupling constant is set as $\lambda = 0.01$.

In Fig. 2, we present the equilibrium probability distribution $\rho_{\text{eq}}(h)$ as a function of the Higgs field h in the unitary gauge.

The root-mean-square value of the Higgs field h can be derived as follows:

$$\bar{h} = \sqrt{\langle h^2 \rangle} = \left[\int_{-\infty}^{+\infty} dh h^2 \rho_{\text{eq}}(h) \right]^{1/2} \simeq 0.363 \left(\frac{H_{\text{inf}}}{\lambda^{1/4}} \right). \quad (2.16)$$

For the present analysis, the SM Higgs self-coupling constant is $\lambda = 0.01$, corresponding to $\bar{h} \simeq 1.15 H_{\text{inf}}$. Thus, the Higgs field in our Universe can be approximated by a uniform background $\bar{h} \simeq H_{\text{inf}}$ combined with Gaussian quantum fluctuations $\delta h(\mathbf{x}, t)$ around the background $\bar{h}$, namely,

$$h(\mathbf{x}, t) = \bar{h}(t) + \delta h(\mathbf{x}, t), \quad (2.17)$$

which enables the application of the mean-field (MF) approximation in the subsequent analysis. In order to obtain the mode functions of the quantum fluctuation $\delta h(\mathbf{x}, t)$, we expand Eq. (2.1) by substituting $h(\mathbf{x}, t) = \bar{h}(t) + \delta h(\mathbf{x}, t)$ to obtain the Lagrangian for δh ,

$$\mathcal{L}(\delta h) \rightarrow \sqrt{-g} \left[-\frac{1}{2} \partial_\mu \delta h \partial^\mu \delta h + \frac{1}{2} \mu^2 (\delta h)^2 - \frac{3}{2} \lambda \bar{h}^2 (\delta h)^2 - (\lambda \bar{h}) \delta h^3 - \frac{\lambda}{4} \delta h^4 \right]. \quad (2.18)$$

In the MF approximation, the field δh is taken to be the linear solution to the equation of motion. The mass of the Higgs fluctuation is $m_h^2 = 3\lambda \bar{h}^2 - \mu^2 \simeq 0.4\lambda^{1/2} H_{\text{inf}}^2$, which means the mass of the fluctuation δh is suppressed by the SM Higgs self-coupling constant λ . In this case, the mass of the Higgs fluctuation in m_h could easily satisfy $m_h^2 \ll H_{\text{inf}}^2$ with a small λ .

For convenience, we use the approximation that the fluctuation δh is massless; thus, the mode functions of $\delta h(\mathbf{k})$ are also the solution for the massless Klein-Gordon equation in the de Sitter spacetime, as shown in Eq. (2.5).

B. Evolution of Higgs field after inflation

After inflation, if the inflaton potential is quadratic around the bottom of the potential, the inflaton would oscillate and behave like the cold matter ($w = 0$) having a mass $m_\phi \sim O(1-10)H_{\text{inf}}$. Consequently, the Universe will expand as $a(t) \sim t^{2/3}$, from which the Hubble parameter is given by $H = 2/(3t)$. Using the Lagrangian (2.1) and considering only the quartic term of the Higgs potential, we find that the evolution of the superhorizon mode of the Higgs field after inflation is described by the following Klein-Gordon equation:

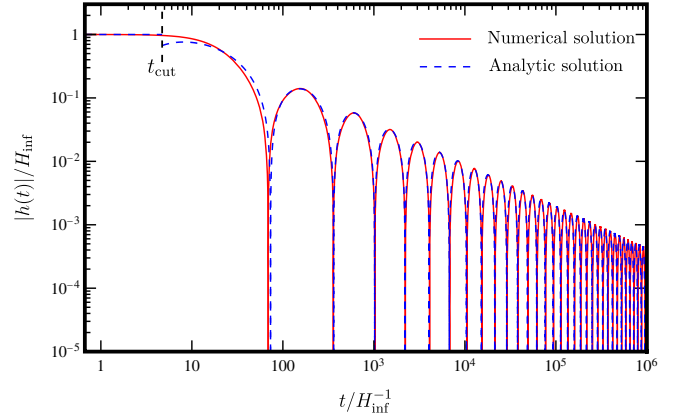


FIG. 3. Evolution of background value of the Higgs field $|h(t)|$ after inflation, where we set an initial value $h_{\text{inf}} = H_{\text{inf}}$. In this plot, the red solid curve represents the numerical solution and the blue dashed curve denotes the analytic solution.

$$\ddot{h}(t) + \frac{2}{t} \dot{h}(t) + \lambda h^3(t) = 0. \quad (2.19)$$

We solve Eq. (2.19) numerically to determine the evolution of the Higgs field h . The results are presented in Fig. 3, where for illustration we set the Higgs self-coupling constant $\lambda = 0.01$ and choose an initial value $h_{\text{inf}} = H_{\text{inf}}$. In this plot, the red solid curve represents the numerical solution, whereas the blue dashed curve denotes the analytical solution.

Additionally, we present a semianalytical solution to Eq. (2.19) which is given in Appendix A. In the following, we derive the analytical formulas for the evolution of the Higgs field $h(t)$ in the case of $h_{\text{inf}} > 0$,

$$h(t) = \begin{cases} h_{\text{inf}}, & t \leq t_{\text{cut}}, \\ AH_{\text{inf}} \left(\frac{h_{\text{inf}}}{H_{\text{inf}} \lambda} \right)^{1/3} (H_{\text{inf}} t)^{-2/3} \cos(\lambda^{1/6} h_{\text{inf}}^{1/3} \omega t^{1/3} + \theta), & t > t_{\text{cut}}, \end{cases} \quad (2.20)$$

where the relevant parameters are given as follows:

$$t_{\text{cut}} = \frac{\sqrt{2}}{3\sqrt{\lambda} h_{\text{inf}}}, \quad A = \left(\frac{2}{9} \right)^{1/3} 5^{1/4} \simeq 0.9, \quad (2.21a)$$

$$\omega = \frac{\Gamma^2(3/4)}{\sqrt{\pi}} 6^{1/3} 5^{1/4} \simeq 2.3, \quad \theta = -3^{-1/3} 2^{1/6} \omega - \arctan 2 \simeq -2.9. \quad (2.21b)$$

One can readily derive the solution for the case $h_{\text{inf}} < 0$. In Fig. 3, we compare the numerical solutions with our analytic solution. We find that for $t \gg t_{\text{cut}}$ the analytic solution agrees with the numerical results very well. This result indicates that after inflation the Higgs field would oscillate in its quartic potential $\frac{1}{4}\lambda h^4$ and the amplitude of these oscillations will decrease due to the Hubble friction term in Eq. (2.19).

III. HIGGS-MODULATED REHEATING USING RIGHT-HANDED NEUTRINOS

In this section, we present a minimal realization incorporating inflation and neutrino seesaw, in which the inflaton decay is modulated by the Higgs boson through right-handed neutrinos. We will give the model realization and setup in Sec. III A and study the Higgs-modulated reheating through right-handed neutrinos in Sec. III B.

A. Model realization and setup

In addition to the particle content of the SM, we introduce the scalar field inflaton ϕ and right-handed neutrinos N_R . The relevant Lagrangian is given as follows:

$$\begin{aligned} \Delta\mathcal{L} = & \sqrt{-g} \left[-\frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) \right. \\ & + \bar{N}_R i \not{\partial} N_R + \frac{1}{\Lambda} \partial_\mu \phi \bar{N}_R \gamma^\mu \gamma^5 N_R \\ & \left. + \left(-\frac{1}{2} M \bar{N}_R^c N_R - y_\nu \bar{L}_L \tilde{\mathbb{H}} N_R + \text{H.c.} \right) \right], \quad (3.1) \end{aligned}$$

where $\mathbb{H}$ is the SM Higgs doublet and $\tilde{\mathbb{H}} = i\sigma_2 \mathbb{H}^*$, with σ_2 as the second Pauli matrix. In Eq. (3.1), $V(\phi)$ is the inflaton potential and its concrete form is irrelevant to the following discussion. After inflation, the potential $V(\phi)$ is assumed to be dominated by the inflaton mass term under which the inflaton ϕ will oscillate. In the above, we have suppressed the flavor indices for the SM leptons and right-handed neutrinos. Each left-handed lepton doublet $L_L = (\nu_L, e_L)^T$ interacts with a right-handed neutrino N_R through the Yukawa coupling y_ν , which is generally a complex matrix. Because of the shift symmetry, the inflaton ϕ couples to the right-handed neutrinos through a dimension-five effective operator with cutoff Λ .⁴ To maintain the perturbative unitarity of the theory during inflation requires Λ to be no less than $60H_{\text{inf}}$, namely, $\Lambda \geq (\dot{\phi})^{1/2} \simeq 60H_{\text{inf}}$. The

⁴As an illustration, this dimension-5 operator in Eq. (3.1) may be induced from a UV model having an approximate global $U(1)_{B-L}$ symmetry which is spontaneously broken by a new scalar field Φ with a $U(1)_{B-L}$ charge -2 and having a vacuum expectation value (VEV), $\langle \Phi \rangle = f$. After spontaneous symmetry breaking (SSB), the Yukawa interaction $(y_x \bar{N}_R^c N_R \Phi + \text{H.c.})$ will generate a Majorana mass $M = y_x f$ for the right-handed neutrino N_R , where $f = M/y_x = O(10)M$ for a natural Yukawa coupling $y_x = O(0.1)$. The inflaton ϕ emerges as a pseudo-Nambu-Goldstone boson from Φ and has its mass generated by soft breaking of the $U(1)_{B-L}$ symmetry, and its interactions with fermions are dictated by their $U(1)$ charges. The residual shift symmetry enforces that the inflaton-fermion coupling (such as the N_R - N_R - ϕ coupling) takes the form of a dimension-5 operator. Its cutoff scale $\Lambda = f$ is given by the $U(1)_{B-L}$ breaking scale f , which is about a factor of $O(10)$ of the N_R mass scale as shown above since the N_R mass originates from the SSB of $U(1)_{B-L}$.

shift symmetry plays a key role for maintaining the flatness of the inflaton potential throughout inflation and is widely realized in models, such as the natural inflation [30] and axion monodromy inflation [31,32], or other models with inflation driven by pseudo-Nambu-Goldstone bosons. The inflaton may couple to SM fermions or gauge bosons through higher-dimensional operators, allowing it to decay into SM particles which are usually suppressed. Consequently, the inflaton is expected to primarily decay into right-handed neutrinos.

For simplicity, we will focus on the case of one generation of fermions for the present study. In this case, the neutrino mass matrix is shown as follows:

$$\mathbf{M}_\nu = \begin{pmatrix} 0 & \frac{y_\nu h}{\sqrt{2}} \\ \frac{y_\nu h}{\sqrt{2}} & M \end{pmatrix}. \quad (3.2)$$

By diagonalizing the neutrino seesaw mass matrix, we derive the neutrino mass eigenstates of ν and N as follows⁵:

$$m_\nu \simeq -\frac{y_\nu^2 h^2}{2M}, \quad M_N \simeq M + \frac{y_\nu^2 h^2}{2M}, \quad (3.3)$$

for $M \gg |y_\nu h|$. The mixing angle θ for diagonalizing $\mathbf{M}_\nu$ is given by

$$\tan \theta = \frac{\sqrt{2} y_\nu h}{\sqrt{M^2 + 2y_\nu^2 h^2} + M} \simeq \frac{y_\nu h}{\sqrt{2} M}. \quad (3.4)$$

In the above, we see that the heavy neutrino mass eigenvalue is shifted by an amount of $\frac{y_\nu^2 h^2}{2M}$ relative to M . This lifting effect is crucial for our mechanism to work, as we are actually probing the seesaw scale of the heavy neutrino mass eigenvalue.

In the Lagrangian (3.1), the inflaton is coupled to the right-handed neutrino N_R through a dimension-five operator $\frac{1}{\Lambda} \partial_\mu \phi \bar{N}_R \gamma^\mu \gamma^5 N_R$. After inflation, the inflaton decays through this operator until the reheating completes. The inflaton should decay into the mass eigenstates (ν and N) instead of chiral eigenstates (ν_L and N_R). So, in terms of mass eigenstates, the relevant interaction vertices from this dimension-five operator take the following form:

$$\begin{aligned} & \frac{\cos^2 \theta}{\Lambda} \partial_\mu \phi \bar{N} \gamma^\mu \gamma^5 N + \frac{\sin^2 \theta}{\Lambda} \partial_\mu \phi \bar{\nu} \gamma^\mu \gamma^5 \nu \\ & - \left(\frac{\sin 2\theta}{2\Lambda} \partial_\mu \phi \bar{N} \gamma^\mu \gamma^5 \nu + \text{H.c.} \right). \quad (3.5) \end{aligned}$$

Hence, there are three decay channels $\phi \rightarrow NN, N\nu, \nu\nu$. Thus, we compute the decay rates of the inflaton as

⁵For the parameter space, we consider that the value of $y_\nu h/M$ always satisfies the condition $|y_\nu h/M| \ll 1$ at the time of reheating.

follows⁶:

$$\Gamma(\phi \rightarrow NN) = \frac{m_\phi M_N^2 \cos^4 \theta}{4\pi\Lambda^2} \left(1 - \frac{4M_N^2}{m_\phi^2}\right)^{1/2}, \quad (3.6a)$$

$$\Gamma(\phi \rightarrow N\nu) = \frac{m_\phi M_N^2 (\sin 2\theta)^2}{32\pi\Lambda^2} \left(1 - \frac{M_N^2}{m_\phi^2}\right)^2, \quad (3.6b)$$

$$\Gamma(\phi \rightarrow \nu\nu) = \frac{m_\phi m_\nu^2 \sin^4 \theta}{4\pi\Lambda^2} \left(1 - \frac{4m_\nu^2}{m_\phi^2}\right)^{1/2}, \quad (3.6c)$$

where the last decay rate $\Gamma(\phi \rightarrow \nu\nu)$ is suppressed by light neutrino mass factor m_ν^2 and is thus fully negligible. Since the mixing angle $\theta \simeq \frac{y_\nu h}{\sqrt{2}M} \ll 1$, we see that the inflaton decay is dominated by the channel $\phi \rightarrow NN$. Consequently, if we neglect the kinematic factors in the above formula, the total decay rate of the inflaton can be approximated as follows:

$$\Gamma \simeq \frac{m_\phi M^2}{4\pi\Lambda^2} \left[1 + \frac{1}{4} \left(\frac{y_\nu h}{M}\right)^2\right]. \quad (3.7)$$

In our setup, the reheating occurs instantaneously at the time $\Gamma = H(t_{\text{reh}}) = 2/(3t_{\text{reh}})$. Equation (3.7) shows that the decay of the inflaton through the right-handed neutrino is modulated by the Higgs field, and then the Higgs fluctuation would induce the curvature perturbation.

Note that our scenario differs from the inflaton decays through the SM fermion channel, where $\Gamma \propto m_f^2 \propto (y_f h)^2$. Since the Higgs field value decreases after inflation, the decay width of inflaton $\Gamma \propto h^2$ would decrease even faster than the Hubble parameter $H(t)$, preventing the completion of reheating [20]. However, unlike the SM fermions, the mass of the right-handed neutrino is mainly contributed by the Majorana mass M instead of the Higgs field value h . This feature prevents the inflaton decay rate Γ from fast decreasing with h . Thus, a viable Higgs-modulated reheating can be realized. We note that the conventional seesaw mechanism has the seesaw scale M typically around 10^{14} GeV, which is comparable to the Hubble scale H_{inf} during inflation. We consider the parameter space of $M_N < m_\phi/2$, and thus the inflaton decaying into two heavy neutrinos is generally kinetically allowed. In our setup, the dimension-five operator discussed above causes the inflaton to decay predominantly into right-handed neutrinos after inflation. If the inflaton couples to the SM fermions via dimension-five operators and under the shift symmetry, the corresponding decay rates are suppressed by the fermion masses, which depend on the Higgs field value (that decreases quickly after inflation). Couplings between the inflaton and SM gauge bosons (via operators such as

$\phi F^{\mu\nu} \tilde{F}_{\mu\nu}$) can be forbidden if the shift symmetry is anomaly-free with respect to the SM gauge group, i.e., the sum of the anomaly parts of fermion triangle loops (containing ϕ and two SM gauge bosons as external lines) vanishes. For the inflaton coupling with the SM Higgs, it may induce additional decay of the inflaton into pairs of the SM Higgs boson, which could become a dominant channel if this coupling would be large enough (and its consequence was discussed in [20]). In this work, we consider a different scenario, in which the inflaton-Higgs derivative coupling is negligibly small.

In passing, we note that our model differs from the literature [33], where a right-handed neutrino is introduced and its mass is modulated solely due to the Dirac mass term and the curvature perturbation is from the kinematic blocking of inflaton decays. In our model, the modulation arises from the neutrino seesaw mechanism. Additionally, the model of [33] has the inflaton couple to the right-handed neutrino via a dimension-four operator without shift symmetry, which would induce a large Planck-scale Majorana mass term for the right-handed neutrino and could make the inflaton decay difficult. In our model, the inflaton has derivative coupling with the right-handed neutrinos under the shift symmetry, so it does not directly contribute to the Majorana mass of the right-handed neutrinos.

We note that the derivative coupling between the inflaton and heavy neutrino can also induce cosmological collider signals during the inflation [34,35]. For the present study, our primary focus is on the predictions for the local type f_{NL} which is generated from the Higgs-modulated reheating through neutrino seesaw. This differs from the conventional cosmological collider signals generated by the inflaton correlation functions during inflation (which does not invoke Higgs-modulated reheating).

B. Higgs-modulated reheating

In our model, the decay rate of the inflaton is influenced by the SM Higgs field. Fluctuations of the Higgs field value (as generated during inflation) cause variations in the inflaton's decay rate across different Hubble patches. These variations perturb the local expansion history, seeding large-scale inhomogeneity and anisotropy in the Universe through these Higgs fluctuations. The δN formalism [36–44] can be used to compute these fluctuations. The number of e-folds of the cosmic expansion after inflation can be derived as follows:

$$\begin{aligned} N(\mathbf{x}) &= \int_{t_{\text{inf}}}^{t_{\text{reh}}(\mathbf{x})} d \ln a(t) = \int_{t_{\text{inf}}}^{t_{\text{reh}}(\mathbf{x})} dt H(t) + \int_{t_{\text{reh}}(\mathbf{x})}^{t_{\text{f}}} dt H(t) \\ &= \int_{\rho_{\text{inf}}}^{\rho_{\text{reh}}(h(\mathbf{x}))} d\rho \frac{H}{\dot{\rho}} + \int_{\rho_{\text{reh}}(h(\mathbf{x}))}^{\rho_{\text{f}}} d\rho \frac{H}{\dot{\rho}}, \end{aligned} \quad (3.8)$$

⁶In our practical calculation, we include all the kinematic factors.

where $a(t)$ is the scale factor, $\rho(t)$ is the total energy density of the universe at the time t , t_{inf} is the physical time at the end of inflation, t_{reh} is the physical time at which reheating occurs, ρ_f is a reference energy density, and t_f is the reference time where the energy density $\rho = \rho_f$ after the completion of reheating.

For the present study, we consider the Universe as a perfect fluid, both before and after the completion of reheating. At the end of inflation (t_{inf}), we assume the inflaton's decay rate Γ is significantly smaller than the Hubble scale. Reheating is completed at a subsequent time t_{reh} when the Hubble parameter satisfies $H(t_{\text{reh}}) = \Gamma_{\text{reh}}$. As discussed in the previous section, during the period before reheating completion ($t_{\text{inf}} < t < t_{\text{reh}}$), the inflaton oscillates near the minimum of a quadratic potential. This corresponds to the matter-dominated Universe, where the pressure $p = 0$ and thus the equation of state parameter $w = p/\rho = 0$. Throughout this stage, the Universe expands as $a \sim t^{2/3}$, with the Hubble parameter $H = 2/(3t)$. After reheating completes ($t > t_{\text{reh}}$), the Universe becomes radiation dominated, which means that the equation of state parameter is $w = 1/3$ and the scale factor behaves as $a \sim t^{1/2}$, with the Hubble parameter given by $H = 1/(2t)$. Here, the right-handed neutrinos decay fast enough after being produced.⁷

As a result, the state of the Universe changes before and after the completion of reheating. For a fluctuation $\delta\Gamma_{\text{reh}}(\mathbf{x})$ in the decay rate, there will be variations in the local reheating time $t_{\text{reh}}(\mathbf{x})$ across different Hubble patches. These variations translate into differences in the expansion history among these patches, which can be quantitatively expressed as fluctuations in the number of e-folds $\delta N(\mathbf{x}, t)$ of local expansion after inflation.

On the other hand, the comoving curvature perturbation after reheating, $\zeta_h(\mathbf{x}, t)$, is equal to the $\delta N(\mathbf{x}, t)$ of cosmic expansion among different Hubble patches under the uniform energy density gauge,

$$\zeta_h(\mathbf{x}, t) = \delta N(\mathbf{x}, t) = N(\mathbf{x}, t) - \langle N(\mathbf{x}, t) \rangle. \quad (3.9)$$

Since the Higgs fluctuation causes the fluctuation of the inflaton decay rate $\delta\Gamma_{\text{reh}}(\mathbf{x})$ in the Higgs-modulated reheating, we can express the comoving curvature perturbation as a function of the local Higgs fluctuation $h(\mathbf{x}, t_{\text{reh}})$. By utilizing the continuity equation

$$\dot{\rho} + 3H(1+w)\rho = 0, \quad (3.10)$$

we integrate Eq. (3.8) and derive the local e-folding number,

⁷We note that the decay width of the right-handed neutrino is $\Gamma_N = \frac{y_N^2 M_N}{8\pi}$. Comparing Γ_N with the Hubble parameter at reheating (H_{reh}), we find that, in most of the parameter space that can be probed in the near future from non-Gaussianity, this relation ($\Gamma_N > H_{\text{reh}}$) can be satisfied.

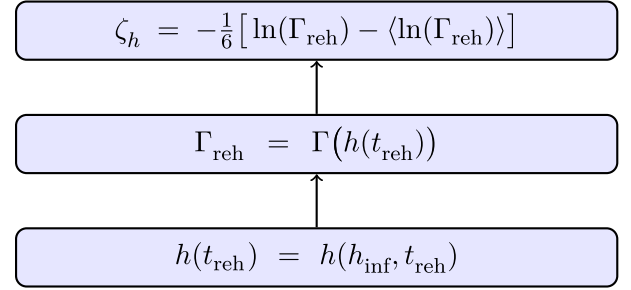


FIG. 4. Schematic plot showing how the comoving curvature perturbation ζ_h sourced from the Higgs-modulated reheating is a function of the Higgs field h_{inf} during the inflation, $\zeta_h = \zeta_h(h_{\text{inf}})$.

$$N(\mathbf{x}) = -\frac{1}{3(1+w_1)} \ln \frac{\rho_{\text{reh}}(h(\mathbf{x}))}{\rho_{\text{inf}}} - \frac{1}{3(1+w_2)} \ln \frac{\rho_f}{\rho_{\text{reh}}(h(\mathbf{x}))}, \quad (3.11)$$

where $w_1 = 0$ is the equation of state parameter before reheating is completed,⁸ and the equation of state parameter after reheating is $w_2 = 1/3$. Utilizing the first Friedmann equation $3H^2 M_p^2 = \rho$ and noting that reheating completes when $H(t_{\text{reh}}) = \Gamma_{\text{reh}}$ is reached, we can derive the following comoving curvature perturbation after reheating ($t > t_{\text{reh}}$):

$$\begin{aligned} \zeta_h(\mathbf{x}, t > t_{\text{reh}}) &= \delta N(\mathbf{x}) = N(\mathbf{x}) - \langle N(\mathbf{x}) \rangle \\ &= -\frac{1}{12} [\ln \rho_{\text{reh}}(\mathbf{x}) - \langle \ln \rho_{\text{reh}}(\mathbf{x}) \rangle] \\ &= -\frac{1}{6} [\ln(H_{\text{reh}}) - \langle \ln(H_{\text{reh}}) \rangle] \\ &= -\frac{1}{6} [\ln(\Gamma_{\text{reh}}) - \langle \ln(\Gamma_{\text{reh}}) \rangle]. \end{aligned} \quad (3.12)$$

In this scenario, we derive a relationship between curvature perturbation $\zeta_h(\mathbf{x})$ from Higgs-modulated reheating and the Higgs field h_{inf} during inflation, as illustrated in Fig. 4.

In the case where the decay of the right-handed neutrino N is significantly delayed (corresponding to a much smaller Yukawa coupling between the Higgs and N), the right-handed neutrino begins to dominate the Universe shortly after the inflaton decays. Since the mass of N is close to that of the inflaton, the Universe quickly becomes matter dominated and remains so until N decays into Standard Model particles. In this regime, the decay rate of N is largely independent of the Higgs vacuum expectation value, and thus the modulation effect is suppressed. We note here that, in most of the parameter space accessible to near-future experiments, the decay rate of N remains larger than the Hubble rate at reheating. Therefore, we neglect this effect in our present analysis.

⁸Our approach also applies to the general case of $w_1 \neq 1/3$.

In Sec. II, we have demonstrated that the value of the Higgs field at the completion of reheating, $h(t_{\text{reh}})$, is determined by the initial value of h_{inf} . Given that the inflaton's decay rate Γ is a function of the Higgs value $h(t_{\text{reh}})$ from Eqs. (3.3) and (3.6), and that the comoving curvature perturbation ζ depends on the decay rate Γ , we can thus establish a relationship between the comoving curvature perturbation ζ and the Higgs field h_{inf} during inflation.

Note that, at the time of the reheating occurring, the value of the Higgs field after inflation becomes an oscillatory function of its initial value,

$$h(t_{\text{reh}}, h_{\text{inf}}) \sim (h_{\text{inf}})^{\frac{1}{3}} \cos(\omega_{\text{reh}} h_{\text{inf}}^{\frac{1}{3}} + \theta), \quad (3.13)$$

with the oscillating frequency estimated as

$$\omega_{\text{reh}} = \lambda_s^{\frac{1}{6}} t_{\text{reh}}^{\frac{1}{3}} \omega. \quad (3.14)$$

When t_{reh} is large, the oscillation frequency can become very high. Note that ζ_h is a function of h^2 and can be expanded into the form $A + Bh^2/M^2 + O(h^4/M^4)$, which includes a factor $\cos^2(\omega_{\text{reh}} h_{\text{inf}}^{1/3} + \theta)$. Since h_{inf} varies across different Hubble volumes and ζ_h is highly sensitive to h_{inf} , averaging over a sufficiently large volume allows the factor $\cos^2(\omega_{\text{reh}} h_{\text{inf}}^{1/3} + \theta)$ to be effectively treated as $1/2$ [45]. Consequently, in the subsequent calculations, we directly set $\cos^2(\omega_{\text{reh}} h_{\text{inf}}^{1/3} + \theta) \rightarrow 1/2$.

Combined with the inflaton fluctuation $\delta\phi$ during inflation, the total comoving curvature perturbation can be written as follows:

$$\zeta = \zeta_\phi + \zeta_h, \quad (3.15)$$

where ζ_ϕ is generated by the inflaton fluctuation $\delta\phi$,

$$\zeta_\phi \simeq -\frac{H_{\text{inf}}}{\dot{\phi}_0} \delta\phi(\mathbf{x}), \quad (3.16)$$

and ζ_h originates from the Higgs-modulated reheating. Because these two components are generated at different times and are independent of each other, the power spectrum of ζ contains both contributions,

$$\mathcal{P}_\zeta = \mathcal{P}_\zeta^{(\phi)} + \mathcal{P}_\zeta^{(h)}, \quad (3.17)$$

where $\mathcal{P}_\zeta^{(\phi)}$ denotes the contribution induced by inflaton fluctuations,

$$\mathcal{P}_\zeta^{(\phi)} = \left(\frac{H}{\dot{\phi}}\right)^2 \mathcal{P}_\phi = \left(\frac{H_{\text{inf}}}{\dot{\phi}}\right)^2 \frac{H_{\text{inf}}^2}{4\pi^2}. \quad (3.18)$$

For convenience, we define R as the square root of the ratio between the power spectrum of the Higgs-modulated reheating and that of the comoving curvature perturbation ζ ,

$$R \equiv \left(\frac{\mathcal{P}_\zeta^{(h)}}{\mathcal{P}_\zeta^{(\phi)}}\right)^{1/2}, \quad (3.19)$$

where $\mathcal{P}_\zeta^{(o)} \simeq 2.1 \times 10^{-9}$ is the observed curvature perturbation [6,7]. To be consistent with observation, we should require $R < 1$.

In the literature, various modulated reheating models were studied, which often assume that all primordial perturbations originate from Higgs-modulated reheating ($R = 1$) [33,46–49]. The existing cosmological observations require $R < 1$. In this work, we compute $\mathcal{P}_\zeta^{(h)}$ through the mean-field method and impose $R < 1$ for our parameter space.

Moreover, the modulated reheating can provide a source of primordial NG. Primordial non-Gaussianity is characterized by the three-point correlation function of ζ , which is also called the bispectrum $\langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \zeta_{\mathbf{k}_3} \rangle$. This primordial non-Gaussianity leaves an imprint on the CMB anisotropy and large-scale structure observations. Different physics effects involved during inflation lead to different shape functions of the primordial NG, and several templates of shape functions are measured by using the Planck-2018 data to identify signals of potential new physics [5,6,8]. For instance, the local non-Gaussianity of the Bardeen potential Φ is given by

$$\langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle'_{\text{local}} = 2A^2 f_{\text{NL}}^{\text{local}} \left\{ \frac{1}{k_1^3 k_2^3} + \frac{1}{k_2^3 k_3^3} + \frac{1}{k_3^3 k_1^3} \right\}. \quad (3.20)$$

In the above, $\langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle'$ is defined as the three-point correlation function excluding the δ function of momentum conservation, $\langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle = (2\pi)^3 \delta^3(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) \times \langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle'$.

For studying the modulated reheating, the comoving curvature perturbation ζ is expanded as a function of the fluctuation of the Higgs field during the inflation δh_{inf} around its mean value $\bar{h}$,

$$\begin{aligned} \zeta_h(\delta h_{\text{inf}}(\mathbf{x})) &= \delta N(\delta h_{\text{inf}}(\mathbf{x})) \\ &= N' \delta h_{\text{inf}} + \frac{1}{2} N'' (\delta h_{\text{inf}})^2 + \dots, \end{aligned} \quad (3.21)$$

where N' and N'' denote the first and second derivatives of the e-folding number N with respect to the Higgs field δh_{inf} , evaluated at its mean value,

$$N' = \left. \frac{dN}{dh_{\text{inf}}} \right|_{\bar{h}}, \quad N'' = \left. \frac{d^2 N}{dh_{\text{inf}}^2} \right|_{\bar{h}}. \quad (3.22)$$

This perturbative approach is also referred to as the mean-field method and has been commonly used to calculate the n -point correlation functions of curvature perturbation. Using this expansion, we can determine the amplitude of

curvature perturbations $\mathcal{P}_\zeta = N'^2 \mathcal{P}_{\delta h}$ and the primordial local non-Gaussianity [46–50].

IV. PROBING NEUTRINO SEESAW USING PRIMORDIAL NON-GAUSSIANITY

As discussed in Sec. III B, for the scenario of Higgs-modulated reheating, the fluctuation of the Higgs field can contribute to the comoving curvature perturbation ζ with a fraction R . Moreover, the Higgs-modulated reheating will also generate primordial non-Gaussianities, which could be detected by the CMB or large-scale structure observations. In this section, we investigate the primordial local non-Gaussianity arising from the Higgs-modulated reheating. We will demonstrate that the primordial non-Gaussianity provides a viable approach to probe the neutrino seesaw scale in this framework. For this purpose, we compute the two- and three-point correlation functions of the comoving curvature perturbation from the Higgs-modulated reheating ζ_h . In this section, the Hubble scale during inflation H_{inf} is abbreviated as H , and the notation of the Higgs field value during inflation h_{inf} is simplified as h , which differs from the Higgs field value $h(t)$ after inflation.

A. Comoving curvature perturbation from Higgs-modulated reheating

As discussed in the Sec. III B, the comoving curvature perturbation ζ from the Higgs-modulated reheating depends on the logarithm of decay rate at the completion of reheating $\ln(\Gamma_{\text{reh}})$, as described by Eq. (3.12).

The relation between Γ_{reh} and the Higgs field h_{inf} during inflation is shown in Fig. 4, namely, $\Gamma_{\text{reh}} = \Gamma(h_{\text{reh}}(t_{\text{reh}}, h_{\text{inf}}))$. In the mean-field approximation, we expand the comoving curvature perturbation to the order of δh_{inf}^2 ,

$$\zeta_h(\mathbf{x}) = -\frac{1-3w}{6(1+w)} \left[\frac{\Gamma'_0}{\Gamma_0} \delta h_{\text{inf}}(\mathbf{x}) + \frac{\Gamma_0 \Gamma''_0 - \Gamma_0' \Gamma_0'}{2\Gamma_0^2} \delta h_{\text{inf}}^2(\mathbf{x}) \right], \quad (4.1)$$

where we have defined

$$\Gamma_0 = \Gamma_{\text{reh}}|_{h_{\text{inf}}(\mathbf{x})=\bar{h}}, \quad (4.2a)$$

$$\Gamma'_0 = \left. \frac{d\Gamma_{\text{reh}}}{dh_{\text{inf}}} \right|_{h_{\text{inf}}(\mathbf{x})=\bar{h}} = \left. \frac{d\Gamma_{\text{reh}}}{dh_{\text{reh}}} \frac{\partial h_{\text{reh}}}{\partial h_{\text{inf}}} \right|_{h_{\text{inf}}(\mathbf{x})=\bar{h}}, \quad (4.2b)$$

$$\begin{aligned} \Gamma''_0 &= \left. \frac{d^2\Gamma_{\text{reh}}}{dh_{\text{inf}}^2} \right|_{h_{\text{inf}}(\mathbf{x})=\bar{h}} \\ &= \left. \frac{d^2\Gamma_{\text{reh}}}{dh_{\text{reh}}^2} \left(\frac{\partial h_{\text{reh}}}{\partial h_{\text{inf}}} \right)^2 \right|_{h_{\text{inf}}(\mathbf{x})=\bar{h}} + \left. \frac{d\Gamma_{\text{reh}}}{dh_{\text{reh}}} \left(\frac{\partial^2 h_{\text{reh}}}{\partial h_{\text{inf}}^2} \right) \right|_{h_{\text{inf}}(\mathbf{x})=\bar{h}}. \end{aligned} \quad (4.2c)$$

After inflation, the inflaton potential is quadratic, so the Universe is matter dominated, implying $w = 0$. We can establish the relation between the curvature perturbation and the Higgs fluctuation as follows:

$$\begin{aligned} \zeta_h(\mathbf{x}) &= -\frac{1}{6} \left[\frac{\Gamma'_0}{\Gamma_0} \delta h_{\text{inf}}(\mathbf{x}) + \frac{\Gamma_0 \Gamma''_0 - \Gamma_0' \Gamma_0'}{2\Gamma_0^2} \delta h_{\text{inf}}^2(\mathbf{x}) \right] \\ &\equiv z_1 \delta h_{\text{inf}}(\mathbf{x}) + \frac{1}{2} z_2 \delta h_{\text{inf}}^2(\mathbf{x}), \end{aligned} \quad (4.3)$$

where z_1 and z_2 are the linear and second-order coefficients,

$$z_1 = -\frac{1}{6} \frac{\Gamma'_0}{\Gamma_0}, \quad z_2 = -\frac{1}{6} \left[\frac{\Gamma''_0}{\Gamma_0} - \left(\frac{\Gamma'_0}{\Gamma_0} \right)^2 \right]. \quad (4.4)$$

In the following, since we only deal with Higgs fluctuations δh_{inf} during the inflation, we will omit the subscript “inf” and directly use the notation δh for convenience, i.e., $\delta h \equiv \delta h_{\text{inf}}$.

From Eq. (4.3), we further derive the corresponding form in Fourier space,

$$\begin{aligned} \zeta_h(\mathbf{k}) &= \int d^3\mathbf{x} \zeta(\mathbf{x}) e^{-i\mathbf{k}\cdot\mathbf{x}} \\ &= z_1 \delta h(\mathbf{k}) + \frac{z_2}{2} \int d^3\mathbf{x} \delta h^2(\mathbf{x}) e^{-i\mathbf{k}\cdot\mathbf{x}} \\ &= z_1 \delta h(\mathbf{k}) + \frac{z_2}{2} \int d^3\mathbf{x} \frac{d^3\mathbf{k}_1}{(2\pi)^3} \frac{d^3\mathbf{k}_2}{(2\pi)^3} \delta h(\mathbf{k}_1) \\ &\quad \times \delta h(\mathbf{k}_2) e^{i(\mathbf{k}_1+\mathbf{k}_2-\mathbf{k})\cdot\mathbf{x}} \\ &= z_1 \delta h(\mathbf{k}) + \frac{z_2}{2} \int \frac{d^3\mathbf{k}_1}{(2\pi)^3} \delta h(\mathbf{k}_1) \delta h(\mathbf{k}-\mathbf{k}_1). \end{aligned} \quad (4.5)$$

Thus, we derive the two-point correlation function of ζ from the Higgs-modulated reheating to the leading order,

$$\langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \rangle_h = z_1^2 \langle \delta h_{\mathbf{k}_1} \delta h_{\mathbf{k}_2} \rangle, \quad (4.6)$$

from which we obtain the power spectrum of ζ ,

$$\mathcal{P}_\zeta^{(h)} = z_1^2 \mathcal{P}_{\delta h} = \frac{z_1^2 H^2}{4\pi^2}. \quad (4.7)$$

From Eq. (3.19), we use R^2 to reflect the ratio between the contribution from Higgs-modulated reheating and the comoving curvature perturbation ζ . So, in this case, we have

$$R = \left(\frac{\mathcal{P}_\zeta^{(h)}}{\mathcal{P}_\zeta^{(o)}} \right)^{1/2} = |z_1| \left(\frac{\mathcal{P}_{\delta h}}{\mathcal{P}_\zeta^{(o)}} \right)^{1/2}, \quad (4.8)$$

where $\mathcal{P}_\zeta^{(o)} \simeq 2.1 \times 10^{-9}$ is the observed curvature perturbation [6,7]. Thus, we require that the ratio R^2 be less than unity, i.e., $R = (\mathcal{P}_\zeta^{(h)}/\mathcal{P}_\zeta^{(o)})^{1/2} < 1$.

B. Three-point correlation function of curvature perturbation

In this subsection, we derive the three-point correlation function of the comoving curvature perturbation ζ contributed by the Higgs-modulated reheating, $\zeta_h = z_1 \delta h + \frac{1}{2} z_2 \delta h^2$. The three-point correlation function of ζ from modulated reheating $\langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \zeta_{\mathbf{k}_3} \rangle_h$ consists of two parts,

$$\langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \zeta_{\mathbf{k}_3} \rangle_h = z_1^3 \langle \delta h_{\mathbf{k}_1} \delta h_{\mathbf{k}_2} \delta h_{\mathbf{k}_3} \rangle + z_1^2 z_2 \langle \delta h^4 \rangle (\mathbf{k}_1, \mathbf{k}_2, \mathbf{k}_3). \quad (4.9)$$

On the right-hand side of the above formula, the first term $z_1^3 \langle \delta h_{\mathbf{k}_1} \delta h_{\mathbf{k}_2} \delta h_{\mathbf{k}_3} \rangle$ is the three-point correlation function of the Higgs fluctuation $\delta h(\mathbf{k})$ generated by the self-interactions of the Higgs field. The second term arises from replacing one $\delta h(\mathbf{k})$ by the nonlinear term $\frac{1}{2} z_2 \delta h^2$, which exists even if the Higgs fluctuation $\delta h(\mathbf{k})$ is purely Gaussian.

As discussed in Sec. II A, the Higgs field could be treated as a massless scalar boson in de Sitter spacetime during inflation. Because of the SM Higgs self-interaction term, $\Delta \mathcal{L} = -\sqrt{-g}[(\lambda \bar{h}) \delta h^3]$, the three-point correlation function of δh is presented as the diagram in Fig. 5.

According to the Schwinger-Keldysh (SK) path integral formalism [51,52], we can compute the three-point correlation function of δh via the following integral:

$$\begin{aligned} \langle \delta h_{\mathbf{k}_1} \delta h_{\mathbf{k}_2} \delta h_{\mathbf{k}_3} \rangle'(\tau_f) &= -i3!\lambda \bar{h} \int_{-\infty}^{\tau_f} d\tau a^4 \left[\prod_{i=1}^3 G_+(\mathbf{k}_i, \tau) - \prod_{i=1}^3 G_-(\mathbf{k}_i, \tau) \right] \\ &= 12\lambda \bar{h} \text{Im} \left(\int_{-\infty}^{\tau_f} d\tau a^4 \prod_{i=1}^3 G_+(\mathbf{k}_i, \tau) \right), \end{aligned} \quad (4.10)$$

where λ is the SM Higgs self-coupling constant, $\bar{h}$ is the uniform Higgs background during inflation, and

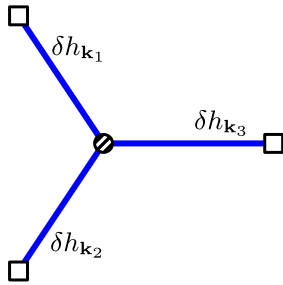


FIG. 5. Three-point correlation function of δh from the Higgs self-interaction $\Delta \mathcal{L} = -\sqrt{-g}[(\lambda \bar{h}) \delta h^3]$. The Higgs propagator is depicted by a blue solid line with a dot and a square at its end points, representing a bulk-to-boundary propagator, which includes one “plus type” and one “minus type.” The square at one end of the propagator indicates its boundary point ($\tau_f \rightarrow 0^-$). The shaded dot at the vertex means that contributions from both plus- and minus-type propagators must be summed.

the $G_{\pm}(\mathbf{k}_i, \tau)$ is the bulk-to-boundary propagator of massless scalar in the SK path integral defined in Appendix D. In Eq. (4.10), $\langle \delta h_{\mathbf{k}_1} \delta h_{\mathbf{k}_2} \delta h_{\mathbf{k}_3} \rangle'$ is defined as the three-point correlation function without including the δ function of momentum conservation, $\langle \delta h_{\mathbf{k}_1} \delta h_{\mathbf{k}_2} \delta h_{\mathbf{k}_3} \rangle = (2\pi)^3 \delta^3(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) \langle \delta h_{\mathbf{k}_1} \delta h_{\mathbf{k}_2} \delta h_{\mathbf{k}_3} \rangle'$. We derive the integral in the last line of Eq. (4.10) to the leading order of τ_f as follows:

$$\begin{aligned} \text{Im} \left(\int_{-\infty}^{\tau_f} d\tau a^4 \prod_{i=1}^3 G_+(\mathbf{k}_i, \tau) \right) &= \text{Im} \int_{-\infty}^{\tau_f} \frac{d\tau}{(H\tau)^4} \frac{H^6}{8k_1^3 k_2^3 k_3^3} \left(\prod_{i=1}^3 (1 - ik_i \tau) \right) e^{i(k_1 + k_2 + k_3)\tau} \\ &= \frac{H^2}{24k_1^3 k_2^3 k_3^3} \left\{ (k_1^3 + k_2^3 + k_3^3) \left[\ln(k_t |\tau_f|) + \gamma - \frac{4}{3} \right] \right. \\ &\quad \left. + k_1 k_2 k_3 - \sum_{a \neq b} k_a^2 k_b \right\}, \end{aligned} \quad (4.11)$$

where $\tau_f \rightarrow 0^-$ is the conformal time when the inflation ends, $\gamma \simeq 0.577$ is the Euler-Mascheroni constant, and wave number $k_t = k_1 + k_2 + k_3$ is around the scale of the present observable Universe. In the above correlation function, the leading contribution is given by the logarithmic term of $\ln(k_t |\tau_f|)$ and there is no inverse power term of $k_t |\tau_f|$ as proved in Ref. [51].

Thus, we further derive the three-point correlation function of δh as follows:

$$\begin{aligned} \langle \delta h_{\mathbf{k}_1} \delta h_{\mathbf{k}_2} \delta h_{\mathbf{k}_3} \rangle' &= \frac{1}{2} \lambda \bar{h} H^2 \left\{ \left(\frac{1}{k_1^3 k_2^3} + 2 \text{ perm.} \right) \left(-N_e + \gamma - \frac{4}{3} \right) \right. \\ &\quad \left. + \frac{1}{k_1^2 k_2^2 k_3^3} - \left(\frac{1}{k_1 k_2^2 k_3^3} + 5 \text{ perm.} \right) \right\}, \end{aligned} \quad (4.12)$$

where N_e is the number of e-folds of the expansion from the time at which the fluctuation mode with momentum k_t first passed outside the horizon [$k_t^{-1} \simeq 1/(a_k H)$] until the end of inflation. The e-folding number N_e is derived as follows:

$$N_e = \ln \frac{a_{\text{end}}}{a_k} = \ln \frac{-(H\tau_f)^{-1}}{k_t/H} = -\ln(k_t |\tau_f|) \sim 60. \quad (4.13)$$

For the second term on the right-hand side of Eq. (4.9), it can be expressed as a four-point correlation function of δh , and to the leading order it is given by the product of two two-point correlation functions,

$$\begin{aligned}
z_1^2 z_2 \langle \delta h^4 \rangle(\mathbf{k}_1, \mathbf{k}_2, \mathbf{k}_3) &= \frac{z_1^2 z_2}{2} \int \frac{d^3 \mathbf{k}_0}{(2\pi)^3} \langle \delta h(\mathbf{k}_1) \delta h(\mathbf{k}_2) \delta h(\mathbf{k}_0) \delta h(\mathbf{k}_3 - \mathbf{k}_0) \rangle + (2 \text{ perm.}) \\
&= \frac{z_1^2 z_2}{2} \left[\int \frac{d^3 \mathbf{k}_0}{(2\pi)^3} \langle \delta h(\mathbf{k}_1) \delta h(\mathbf{k}_0) \rangle \langle \delta h(\mathbf{k}_2) \delta h(\mathbf{k}_3 - \mathbf{k}_0) \rangle + (\mathbf{k}_1 \leftrightarrow \mathbf{k}_2) \right] + (2 \text{ perm.}) \\
&= \frac{z_1^2 z_2}{2} \left[\int d^3 \mathbf{k}_0 (2\pi)^3 \delta^3(\mathbf{k}_1 + \mathbf{k}_0) \delta^3(\mathbf{k}_2 + \mathbf{k}_3 - \mathbf{k}_0) \frac{H^4}{4k_1^3 k_2^3} + (\mathbf{k}_1 \leftrightarrow \mathbf{k}_2) \right] + (2 \text{ perm.}) \\
&= (2\pi)^3 \delta^3(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) z_1^2 z_2 \left[\frac{H^4}{4k_1^3 k_2^3} + (2 \text{ perm.}) \right]. \tag{4.14}
\end{aligned}$$

Combined with the part arising from the nonlinear rate, we derive the three-point correlation function of the comoving curvature perturbation ζ from Higgs-modulated reheating as follows:

$$\begin{aligned}
\langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \zeta_{\mathbf{k}_3} \rangle'_h &= \frac{z_1^3}{2} \lambda \bar{h} H^2 \left\{ \left[\frac{1}{k_1^3 k_2^3} + (2 \text{ perm.}) \right] \left(-N_e + \gamma - \frac{4}{3} \right) + \frac{1}{k_1^2 k_2^2 k_3^2} - \left[\frac{1}{k_1 k_2^2 k_3^3} + (5 \text{ perm.}) \right] \right\} \\
&\quad + \frac{z_1^2 z_2}{4} H^4 \left[\frac{1}{k_1^3 k_2^3} + (2 \text{ perm.}) \right]. \tag{4.15}
\end{aligned}$$

In the above, we have included the non-Gaussianity contributions from both the nonlinear term and the Higgs self-interactions, whereas the previous studies only include the former [46–49].

C. Probing the seesaw scale through local non-Gaussianity

In the previous subsection, we derived the three-point correlation function $\langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \zeta_{\mathbf{k}_3} \rangle'_h$ for the curvature perturbation. For this subsection, we further compute the local non-Gaussianity $f_{\text{NL}}^{\text{local}}$ originating from the Higgs-modulated reheating in our model, with which we study the probe of the seesaw mechanism in this framework.

The three-point correlation function $\langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \zeta_{\mathbf{k}_3} \rangle'_h$ contributes to three distinct classes of non-Gaussian shape templates (local, equilateral, and orthogonal types) according to Refs. [5,8]. We present a systematic analysis of the specific contributions of $\langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \zeta_{\mathbf{k}_3} \rangle'_h$ to these templates as in Appendix B. We evaluate the amplitude of a given non-Gaussian shape template, parametrized by f_{NL}^i , which is expressed as a function of the model parameters

$$f_{\text{NL}}^i = f_{\text{NL}}^i(M, y_\nu, \lambda, H_{\text{inf}}, m_\phi, \Lambda), \tag{4.16}$$

where “ i ” represents the type of non-Gaussian shape template. For the present analysis, we choose a set of

TABLE I. The relevant parameters with benchmark values chosen for the present analysis.

Parameters	$\mathcal{P}_\zeta$	N_e	H_{inf}	m_ϕ	Λ	λ
Values	2.1×10^{-9}	60	$(1, 3) \times 10^{13} \text{ GeV}$	$40 H_{\text{inf}}$	$60 H_{\text{inf}}$	0.01

relevant parameters having benchmark values, as shown in Table I. The amplitude of the comoving curvature perturbation power spectrum $\mathcal{P}_\zeta$ is taken as $\ln(10^{10} \mathcal{P}_\zeta) \simeq 3.047$, according to the Planck-2018 data [6,7]. The SM Higgs self-coupling is set to be $\lambda = 0.01$.

We find that the non-Gaussianity predicted by our model mainly belongs to the local type. This is due to the following reasons. First, in the three-point correlation function arising from the Higgs self-coupling, as described in Eq. (4.12), the local type is amplified by the number of e-folds N_e , whereas the other two shapes do not receive such enhancement. Second, as shown in Eq. (4.14), the contribution from the nonlinear term exclusively generates local-type non-Gaussianity. Hence, for the present study, we primarily focus on the magnitude of local NG as predicted by our model, which can be approximately expressed by the following⁹:

$$f_{\text{NL}}^{\text{local}} \simeq -\frac{10}{3} \frac{z_1^3 H^3}{(2\pi)^4 \mathcal{P}_\zeta^2} \left(\frac{\lambda \bar{h}}{2H} N_e - \frac{z_2 H}{4z_1} \right). \tag{4.17}$$

It is found that the local-type non-Gaussianity f_{NL} arising from the nonlinear term [46–48,50] is given by

$$f_{\text{NL}}^{\text{local}}(\text{NT}) = \frac{5z_2}{6z_1^2} = 5 \left(1 - \frac{\Gamma_0'' \Gamma_0}{\Gamma_0^2} \right). \tag{4.18}$$

We note that this result is compatible with the scenario where the Higgs fluctuation is the only source of primordial fluctuation, i.e., $R = 1$. Additionally, these studies assumed the absence of intrinsic non-Gaussianity in the Higgs field

⁹The accurate formula is given in Eq. (B5) of Appendix B.

TABLE II. Comparison of major contributions to the non-Gaussianity arising from the Higgs self-interaction and the nonlinear term for three sets of benchmark points with specific neutrino seesaw scale M and the Higgs-neutrino Yukawa coupling y_ν . The ratio R of Eq. (3.19) is presented in the fourth row. The predicted values of $f_{\text{NL}}^{\text{local}}$ (HSC), $f_{\text{NL}}^{\text{local}}$ (NT), and $f_{\text{NL}}^{\text{local}}$ (total) are shown in the fifth, sixth, and seventh rows, respectively. In each entry, the number outside (inside) the parentheses corresponds to the input of Higgs self-coupling $\lambda = 0.01(0.02)$.

Benchmarks	A_1	A_2	B_1	B_2	C_1	C_2
M/H_{inf}	3	3	7	7	14	14
y_ν	0.3	0.6	0.3	0.6	0.3	0.6
R	0.02 (0.01)	0.08 (0.06)	0.03 (0.02)	0.10 (0.07)	0.03 (0.02)	0.13 (0.09)
$f_{\text{NL}}^{\text{local}}$ (HSC)	0.037 (0.019)	2.38 (1.19)	0.07 (0.036)	4.55 (2.28)	-0.15 (-0.08)	-9.8 (-4.9)
$f_{\text{NL}}^{\text{local}}$ (NT)	0.008 (0.003)	0.49 (0.17)	0.01 (0.005)	0.93 (0.33)	-0.03 (-0.01)	-2.0 (-0.7)
$f_{\text{NL}}^{\text{local}}$ (total)	0.045 (0.021)	2.86 (1.36)	0.09 (0.041)	5.49 (2.61)	-0.18 (-0.09)	-11.8 (-5.6)

due to its self-coupling $\Delta\mathcal{L} = -\sqrt{-g}(\lambda\bar{h})\delta h^3$, whose contribution, however, could be significant as will be shown below. For the present analysis, we incorporate both the intrinsic non-Gaussianity of the Higgs self-interactions and the nonlinear term. By assuming $R = 1$ and purely Gaussian δh , our result will reduce to Eq. (4.18) as shown in Appendix B. However, in the general case, Eq. (4.17) includes two distinct contributions: one originating from the Higgs three-point correlation function induced by Higgs self-interactions and another from the nonlinear term.

In the following, we will demonstrate that these contributions are significant for sample inputs of the Higgs self-coupling $\lambda = 0.01$ and the e-folding $N_e = 60$. For illustration, we provide three sets of representative benchmark points where the seesaw scale is chosen as $M = (3, 7, 14)H_{\text{inf}}$, respectively. For each given seesaw scale M , we input the sample values of the Higgs-neutrino Yukawa coupling as $y_\nu = (0.3, 0.6)$, respectively. With these, we evaluate the contributions to the local non-Gaussianity $f_{\text{NL}}^{\text{local}}$ by the Higgs self-interaction and nonlinear term, which are presented and compared in Table II for the three sets of benchmarks of the neutrino seesaw scale M and Yukawa coupling y_ν . For comparison, we also vary the Higgs self-coupling as $\lambda = 0.02$ and present the corresponding contributions in the parentheses of each entry for the ratio R and the non-Gaussianities $f_{\text{NL}}^{\text{local}}$ (HSC), $f_{\text{NL}}^{\text{local}}$ (NT), and $f_{\text{NL}}^{\text{local}}$ (total), as shown in Table II, where HSC, NT, and total stand for the contributions by the Higgs self-coupling, the nonlinear term, and their sum, respectively. This demonstrates that, in the neutrino seesaw parameter space, the Higgs self-interaction gives the dominant contribution to non-Gaussianity, whereas the nonlinear term provides a sizable but subdominant contribution. It also shows that the non-Gaussianity measurements are sensitive to the size of the Higgs self-coupling λ at the seesaw scale.

In the seesaw mechanism, the mass of the light neutrino ν is determined by the Majorana mass M and the neutrino-Higgs Yukawa coupling y_ν through the formula

$$m_\nu = \frac{y_\nu^2 v^2}{2M} \sim O(0.05) \text{ eV}, \quad (4.19)$$

where $v \simeq 246$ GeV is the vacuum expectation value of the Higgs field after electroweak symmetry breaking.¹⁰ The neutrino oscillation data give $\Delta m_{31}^2 \simeq 2.5 \times 10^{-3}$ and $\Delta m_{21}^2 \simeq 7.5 \times 10^{-5} \text{ eV}^2$ [55], which require at least one of the light neutrino masses to be $m_\nu \gtrsim 0.05 \text{ eV}$. With this, we can estimate the neutrino seesaw scale to be around 10^{14} GeV for the Yukawa coupling $y_\nu = O(1)$. On the other hand, cosmological measurements based on the CMB alone already set upper bounds on the sum of the light neutrino masses, approximately $\sum m_\nu \lesssim 0.26 \text{ eV}$ [7]. When combined with observations of large-scale structures, this bound can be tightened to $O(0.1) \text{ eV}$. For instance, the eBOSS Collaboration [56] placed a 95% upper bound $\sum m_\nu \lesssim 0.10 \text{ eV}$ and the DES Collaboration [57] set a constraint $\sum m_\nu \lesssim 0.13 \text{ eV}$ at 95% confidence level. Given the existing cosmological upper bounds on the neutrino mass sum $\sum m_\nu$ and two mass-squared differences $(\Delta m_{3\ell}^2, \Delta m_{21}^2)$ measured by oscillation data, we can determine the largest light neutrino mass for the normal mass ordering (NO) and inverted mass ordering (IO) from the following conditions:

$$\text{NO: } \sum m_\nu = m_3 + \sqrt{m_3^2 - \Delta m_{31}^2} + \sqrt{m_3^2 - \Delta m_{31}^2 + \Delta m_{21}^2}, \quad (4.20a)$$

$$\text{IO: } \sum m_\nu = m_2 + \sqrt{m_2^2 - \Delta m_{21}^2} + \sqrt{m_2^2 - |\Delta m_{32}|^2}, \quad (4.20b)$$

¹⁰In Eq. (4.19), m_ν is the mass of the light neutrino ν at the electroweak scale. During and after the inflation, we can include the renormalization-group running effect for this neutrino mass at the corresponding Hubble scale of $O(10^{13}) \text{ GeV}$, which is about 30% larger than its low-energy value at the electroweak scale [53,54].

where m_3 is the largest mass for the normal ordering and m_2 is the largest mass for the inverted ordering.

If we choose the weaker bound on the light neutrino mass sum $\sum m_\nu \lesssim 0.26$ eV by CMB alone [7], we can solve the conditions (4.20) to obtain the largest light neutrino mass to be around 0.1 eV for either NO or IO. In comparison, choosing the stronger bound $\sum m_\nu \lesssim 0.13$ eV [57], we find from the conditions (4.20) that the largest light neutrino mass is $m_3 \simeq 0.06$ eV for the NO and $m_2 \simeq 0.05$ eV for the IO.

To probe the seesaw parameter space in our model, we compute the non-Gaussianity $f_{\text{NL}}^{\text{local}}$ for different values of the seesaw scale M and neutrino-Higgs Yukawa coupling y_ν and set the SM Higgs self-coupling $\lambda = 0.01$. We present our numerical findings in Figs. 6 and 7. The colored region

satisfies the requirement $R < 1$, whereas the white region in the upper-right corner of each plot corresponds to $R > 1$ and is excluded. The region with blue color corresponds to $f_{\text{NL}}^{\text{local}} > 0$, and the red regions represent $f_{\text{NL}}^{\text{local}} < 0$.

In Figs. 6 and 7, the green contours describe the existing 2σ bounds, $-11.1 \leq f_{\text{NL}}^{\text{local}} \leq 9.3$, as given by the Planck-2018 data [8]. Moreover, we display contours for $f_{\text{NL}} = \pm 1, \pm 0.1, \pm 0.01$, plotted as orange, yellow, and white curves, respectively. These contours represent the potential sensitivity reaches of the ongoing and future observations, such as those from DESI [58], CMB-S4 [59], Euclid [60], SPHEREx [61], LSST [62], and SKA [63] experiments. We present the seesaw predictions for the light neutrino mass of $m_\nu = 0.1(0.06)$ eV (pink curves) and $m_\nu = 0.05$ eV (purple curves), and for the Hubble

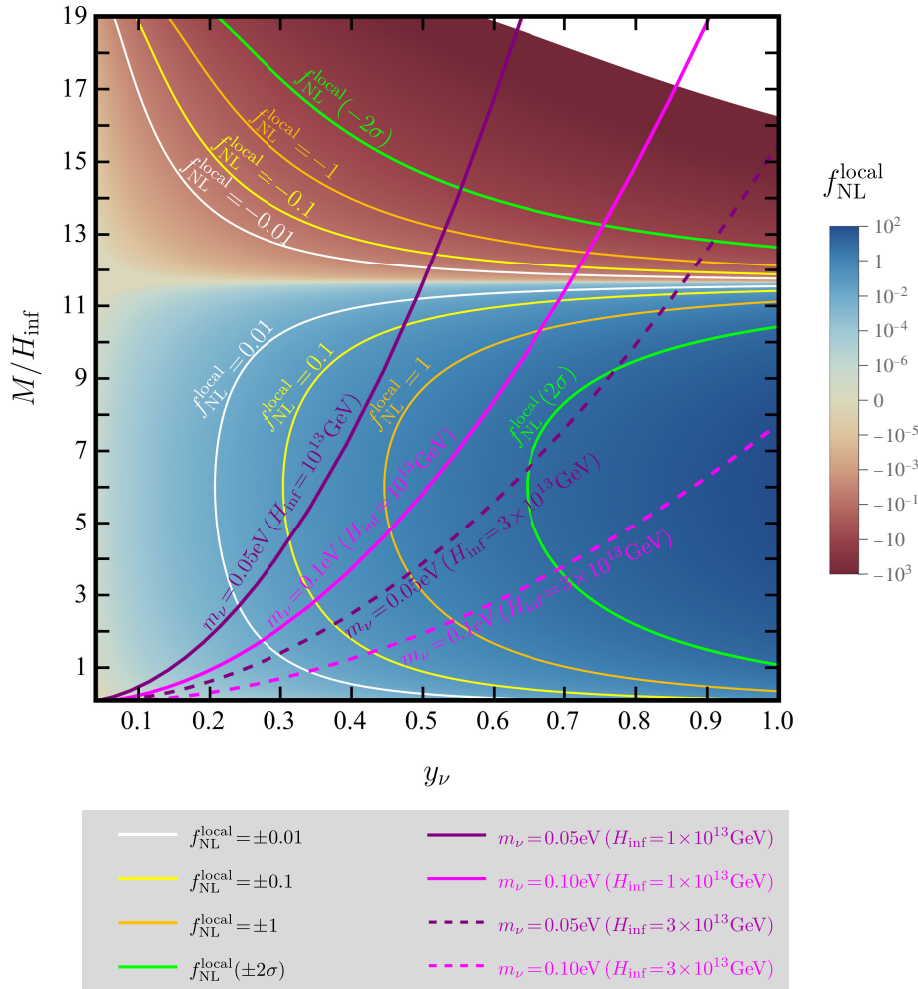


FIG. 6. Prediction of the non-Gaussianity $f_{\text{NL}}^{\text{local}}$ from the seesaw parameter space of heavy neutrino mass scale M versus Yukawa coupling y_ν , where the SM Higgs self-coupling is input as $\lambda = 0.01$, and the Hubble parameter during inflation is set as $H_{\text{inf}} = 10^{13}$ and 3×10^{13} GeV, respectively. The blue region represents the parameter space with positive $f_{\text{NL}}^{\text{local}}$, whereas the red region represents the parameter space with negative $f_{\text{NL}}^{\text{local}}$. The green contour depicts the 2σ bound based on Planck-2018 data, corresponding to $-11.1 \leq f_{\text{NL}}^{\text{local}} \leq 9.3$. The colored region satisfies the requirement of $R < 1$ and the uncolored region in the upper-right corner corresponds to $R \geq 1$. The seesaw predictions for the light neutrino mass of $m_\nu = 0.1$ eV (pink curves) and $m_\nu = 0.05$ eV (purple curves) and for the Hubble parameter $H = 10^{13}$ GeV (solid curves) and $H = 3 \times 10^{13}$ GeV (dashed curves) are given.

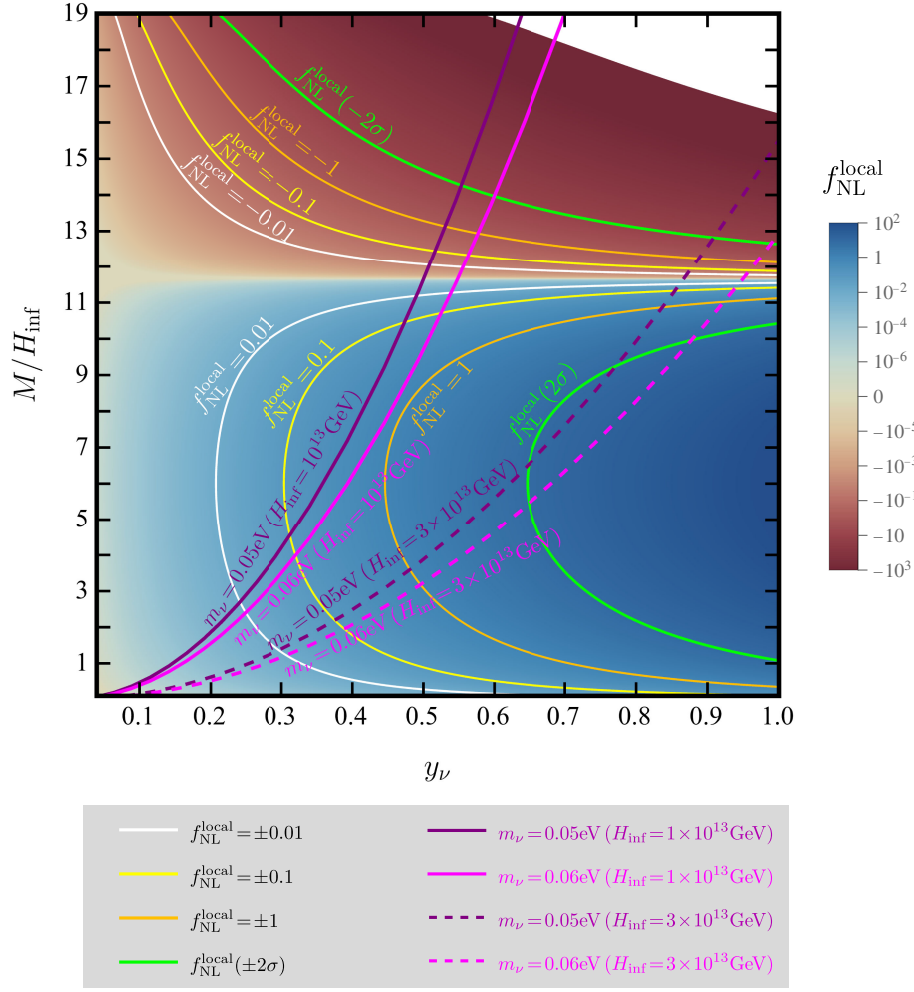


FIG. 7. Prediction of the non-Gaussianity $f_{\text{NL}}^{\text{local}}$ from the seesaw parameter space of heavy neutrino mass scale M versus Yukawa coupling y_ν , where the SM Higgs self-coupling is input as $\lambda = 0.01$, and the Hubble parameter during inflation is set as $H_{\text{inf}} = 10^{13}$ and 3×10^{13} GeV, respectively. The seesaw predictions for the light neutrino mass of $m_\nu = 0.06$ (pink curves) and $m_\nu = 0.05$ eV (purple curves) and for the Hubble parameter $H = 10^{13}$ GeV (solid curves) and $H = 3 \times 10^{13}$ GeV (dashed curves) are given.

parameter $H = 10^{13}$ GeV (solid curves) and $H = 3 \times 10^{13}$ GeV (dashed curves), where we have included the renormalization-group running effects for the light neutrino mass at the Hubble scale. It shows that a larger value of the Hubble parameter will shift the pink and purple curves toward the right-hand-side region with larger Yukawa coupling y_ν .

For the local-type NG $f_{\text{NL}}^{\text{local}} > 0$, we see from Fig. 6 that the existing measurements of Planck-2018 (shown as the 2σ green contours) already have sensitivity to probe the case of a light neutrino mass around $m_\nu = 0.05$ eV for Hubble parameter $H = 3 \times 10^{13}$ GeV, as shown by the purple dashed curve; whereas for a larger light neutrino mass around $m_\nu = 0.1$ eV with $H = 3 \times 10^{13}$ GeV, a large portion of the parameter region in our model is already excluded by the Planck-2018 data as shown by the pink dashed curve. Then, for a smaller Hubble parameter $H = 10^{13}$ GeV, Fig. 6 shows that for inputting the light

neutrino mass of range $m_\nu = (0.05-0.1)$ eV, our seesaw predictions (purple and pink solid curves) have significant parameter space consistent with the current bound of Planck-2018 (the 2σ green contour), but they can be further probed by the improved non-Gaussianity measurements of future experiments as shown by the (orange, yellow, white) contours. For instance, probing the case of a light neutrino mass $m_\nu = 0.05$ eV and Hubble parameter $H = 10^{13}$ GeV (purple solid curve) requires a sensitivity to $f_{\text{NL}}^{\text{local}} \sim 0.1$ or even $f_{\text{NL}}^{\text{local}} \sim 0.01$, depending on the seesaw mass scale M . In contrast, the case of $m_\nu = 0.1$ eV (pink solid curve) can be more effectively probed in the near future.

Then, for the case of $f_{\text{NL}}^{\text{local}} < 0$, Fig. 6 shows that the existing bounds of Planck-2018 measurements (shown as the -2σ green contour) have already excluded a large portion of the seesaw parameter space in our model. For instance, for a light neutrino mass $m_\nu = 0.1$ eV (0.05 eV) and Hubble parameter $H = 10^{13}$ GeV, the parameter space

with seesaw scale $M \gtrsim 13H_{\text{inf}}$ ($M \gtrsim 14H_{\text{inf}}$) is excluded by the Planck-2018 data in our model. The future measurements of non-Gaussianity under planning will be able to extensively probe the seesaw parameter space in our model with $M \gtrsim 12H_{\text{inf}}$ in the case of $f_{\text{NL}}^{\text{local}} < 0$.

In parallel, we further present the seesaw predictions in Fig. 7 for choosing the largest light neutrino mass $m_\nu = 0.06$ eV (representing normal ordering of light neutrino masses) or $m_\nu = 0.05$ eV (representing the inverted ordering), as depicted by the pink curves and purple curves, respectively. We see that, for $H = 3 \times 10^{13}$ GeV (dashed curves), the existing constraints of Planck 2018 (green contours) already have sensitivity to probe a part of the seesaw parameter space in our model with the largest light neutrino mass corresponding to the NO versus IO of light neutrinos. Figure 7 further shows that, for a smaller Hubble parameter $H = 10^{13}$ GeV (solid curves), the future non-Gaussianity measurements with sensitivities of $f_{\text{NL}}^{\text{local}} \lesssim 0.1$ can sensitively probe the seesaw predictions in our framework with $m_\nu = 0.06$ eV (pink solid curve) versus $m_\nu = 0.05$ eV (purple solid curve).

From Figs. 6 and 7, we see that the cosmological non-Gaussianity measurements are rather sensitive to the value of light neutrino mass scale m_ν , which is also constrained by the low-energy neutrino experiments, especially the determination of the light neutrino mass ordering (as will be measured by the neutrino oscillation experiments JUNO [64] and DUNE [65]) and the determination of light neutrino mass scale (by the on-going and future neutrinoless double- β decay experiments [66]). Hence, our analyses of Figs. 6 and 7 demonstrate the important interplay between the light neutrino mass determinations by the low-energy experiments and the high-scale cosmological measurements on the non-Gaussianity in our model.

As a final point, we note in Figs. 6 and 7 that the non-Gaussianity $f_{\text{NL}}^{\text{local}}$ has sign flip around the value of $M \simeq 12H_{\text{inf}}$. This can be understood as follows. We note that, in the parameter space considered in this study, the sign of the non-Gaussianity is mainly determined by the contribution from the Higgs self-interaction. As shown in Eq. (4.17), it is linked to the coefficient z_1 , which depends on the derivative of the inflaton decay width. Since the Higgs field during inflation enters the decay width through neutrino seesaw as shown in Eq. (3.3), the derivative $\Gamma'_0 = d\Gamma_{\text{reh}}/dh_{\text{inf}}$ is given by¹¹

$$\Gamma'_0 = \frac{d\Gamma_{\text{reh}}}{dM_N} \frac{dM_N}{dh_{\text{inf}}}, \quad (4.21)$$

¹¹In fact, a quantity (with dependence on the Higgs field h) could be converted into a function of the right-handed neutrino mass M_N , including the left-handed neutrino mass m_ν and the mixing angle θ .

where $\Gamma'_0 \propto z_1$ according to Eq. (4.4) and the derivative $\frac{dM_N}{dh_{\text{inf}}} \simeq \frac{y_\nu^2 h_{\text{reh}}}{M^2} \frac{dh_{\text{reh}}}{dh_{\text{inf}}}$ is nonzero. Hence, the condition $z_1 = 0$ requires $\frac{d\Gamma_{\text{reh}}}{dM_N} = 0$. At the time of reheating, we have $M_N \simeq M$, at which we compute the derivative

$$\begin{aligned} \frac{d\Gamma_{\text{reh}}}{dM_N} &= \frac{m_\phi M}{8\pi\Lambda^2} \left[\left(1 - \frac{M^2}{m_\phi^2}\right)^2 \left(1 - \frac{4M^2}{m_\phi^2}\right)^{\frac{1}{2}} - \frac{8M^2}{m_\phi^2} \right] \\ &\quad \times \left(1 - \frac{4M^2}{m_\phi^2}\right)^{-\frac{1}{2}}. \end{aligned} \quad (4.22)$$

From the condition $\frac{d\Gamma_{\text{reh}}}{dM_N} = 0$, we solve the ratio of the seesaw scale M over the inflaton mass m_ϕ as follows:

$$\frac{M}{m_\phi} \simeq 0.29. \quad (4.23)$$

Thus, for the input of $m_\phi = 40H_{\text{inf}}$, the sign-transition point $z_1 = 0$ corresponds to $M \simeq 11.6H_{\text{inf}}$. This nicely explains that the transition between $f_{\text{NL}}^{\text{local}} > 0$ and $f_{\text{NL}}^{\text{local}} < 0$ in Figs. 6 and 7 happens around the horizontal line of $M \simeq 12H_{\text{inf}}$.

D. Dependence of non-Gaussianity on the Higgs self-coupling

In this subsection, we analyze the dependence of NG on the Higgs self-coupling (λ).

As shown in Eq. (4.17), the local-type non-Gaussianity $f_{\text{NL}}^{\text{local}}$ receives contributions from both the HSC term and the NT term. In Table II, we have presented numerically their individual contributions and their sum for comparison. It shows that the Higgs self-coupling produces the dominant contribution to non-Gaussianity (as shown in the fifth row), and the nonlinear term gives the subdominant contribution (as shown in the sixth row). To examine the sensitivity of $f_{\text{NL}}^{\text{local}}$ to the Higgs self-coupling contribution in Table II, we vary the Higgs self-coupling from $\lambda = 0.01$ to $\lambda = 0.02$. It shows that, in each entry of Table II, the number outside (inside) the parentheses corresponds to the input of Higgs self-coupling $\lambda = 0.01(0.02)$. We see that increasing the λ value by a factor 2 generally causes the reduction of both $f_{\text{NL}}^{\text{local}}$ (HSC) and $f_{\text{NL}}^{\text{local}}$ (NT) by about a factor of $\frac{1}{2}$ and $\frac{1}{3}$, respectively. These features can be explained by the following simple scaling behaviors:

$$f_{\text{NL}}^{\text{local}}(\text{HSC}) \propto \lambda^{-1}, \quad f_{\text{NL}}^{\text{local}}(\text{NT}) \propto \lambda^{-\frac{3}{2}}, \quad (4.24)$$

which we have derived in Appendix C. (We note that the above scaling behaviors are derived by using the mean-field approximation, which are valid for the parameter space under consideration, but not for arbitrarily small or large coupling λ .) These features are also reflected in Fig. 8, where the Higgs self-coupling $\lambda = 0.01$ corresponds to the

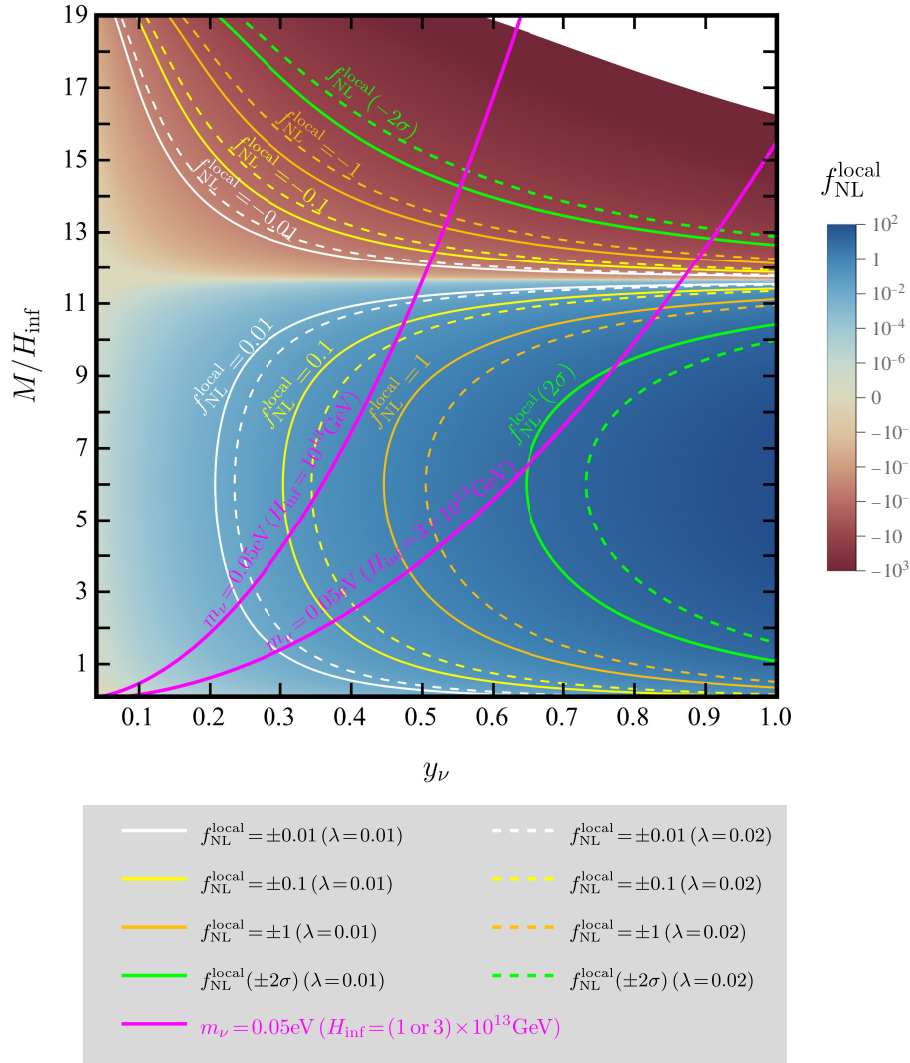


FIG. 8. Prediction of the non-Gaussianity $f_{\text{NL}}^{\text{local}}$ from the seesaw parameter space of heavy neutrino mass scale M versus the Yukawa coupling constant y_ν , where we input the SM Higgs self-coupling $\lambda = 0.01$ (solid curves) and $\lambda = 0.02$ (dashed curves), and the Hubble parameter during inflation is set as $H_{\text{inf}} = 10^{13}$ and 3×10^{13} GeV, respectively. The contours in solid curves represent the given bounds of non-Gaussianity $f_{\text{NL}}^{\text{local}}$ with Higgs self-coupling $\lambda = 0.01$, whereas the contours in dashed curves correspond to the bounds with Higgs self-coupling $\lambda = 0.02$. The seesaw predictions are presented by pink curves for the light neutrino mass of $m_\nu = 0.05$ eV, with the Hubble parameter $H = 10^{13}$ and $H = 3 \times 10^{13}$ GeV, respectively.

$f_{\text{NL}}^{\text{local}}$ contours given by the solid curves, and the value $\lambda = 0.02$ corresponds to the $f_{\text{NL}}^{\text{local}}$ contours given by the dashed curves. (The $f_{\text{NL}}^{\text{local}}$ contours in solid curves are the same as those in Figs. 6 and 7.) In Fig. 8, we show that, by inputting a larger Higgs self-coupling value $\lambda = 0.02$, the non-Gaussianity contours (given by dashed curves) impose weaker bounds on the seesaw parameter space of (M, y_ν) as compared to the contours (in solid curves) with a smaller coupling $\lambda = 0.01$.¹² The above analyses show that the

measurements of non-Gaussianity $f_{\text{NL}}^{\text{local}}$ are sensitive to the probe of the Higgs self-coupling λ at the seesaw scale, which is quantitatively connected to the low-energy values of λ (measured by the LHC and future high-energy colliders [67]) via the renormalization-group evolution. Hence, this also demonstrates the important interplay on probing the Higgs self-coupling λ between the high-scale cosmological non-Gaussianity measurements and the TeV-scale collider measurements.

In passing, we comment on the effect of the cutoff scale Λ for the ϕ - N_R interaction of Eq. (3.1). It can affect the time of reheating completion t_{reh} . A smaller Λ leads to a larger inflaton decay rate and thus smaller t_{reh} , allowing less time for the Higgs field value to decrease [cf. Eq. (2.20)] and

¹²This is in contrast to the conventional collider probe of the Higgs self-coupling λ , where a larger λ value always produces stronger signals of the di-Higgs production [67].

TABLE III. Comparison of the ratio R and the non-Gaussianity in our framework with different cutoff Λ for three sets of benchmark points with specific neutrino seesaw scale M and the Higgs-neutrino Yukawa coupling y_ν . The ratio R of Eq. (3.19) with three benchmark cutoffs is presented in the fourth, fifth, and sixth rows. The predicted non-Gaussianity $f_{\text{NL}}^{\text{local}}$ with different cutoffs Λ are shown in the last three rows.

Benchmarks	A_1	A_2	B_1	B_2	C_1	C_2
M/H_{inf}	3	3	7	7	14	14
y_ν	0.3	0.6	0.3	0.6	0.3	0.6
$R(\Lambda = 60H_{\text{inf}})$	0.021	0.083	0.026	0.10	0.033	0.13
$R(\Lambda = 80H_{\text{inf}})$	0.010	0.039	0.012	0.048	0.015	0.062
$R(\Lambda = 100H_{\text{inf}})$	0.005	0.021	0.007	0.027	0.009	0.034
$f_{\text{NL}}^{\text{local}}(\Lambda = 60H_{\text{inf}})$	0.045	2.9	0.086	5.5	-0.18	-12
$f_{\text{NL}}^{\text{local}}(\Lambda = 80H_{\text{inf}})$	0.004	0.29	0.009	0.55	-0.018	-1.2
$f_{\text{NL}}^{\text{local}}(\Lambda = 100H_{\text{inf}})$	0.001	0.048	0.001	0.092	-0.003	-0.20

causing larger fluctuations, thereby increasing the non-Gaussianity. These are shown in Table III. Hence, further understanding of the UV dynamics of inflation for the determinations of H_{inf} , m_ϕ , and Λ would be beneficial for a definitive probe of the seesaw parameter space.

V. CONCLUSIONS

The seesaw mechanism provides the most appealing resolution to the origin of tiny masses of active neutrinos. However, the natural seesaw scale is as high as 10^{14} GeV, posing a great challenge for experimental tests at particle colliders. Considering the fact that the natural neutrino seesaw scale can be around the upper range of the inflation scale, we proposed a new framework incorporating inflation and neutrino seesaw in which the inflaton primarily decays into heavy right-handed neutrinos. The inflaton couples to the right-handed neutrinos through an effective interaction that respects the shift symmetry. Through the neutrino seesaw, fluctuations of the Higgs field can modulate the inflaton decays, leading to non-Gaussian signatures at large scales. This provides, for the first time, an important means to directly probe the seesaw mechanism in the early Universe by measuring the non-Gaussian signatures. Moreover, it is appealing that this scenario also naturally provides an initial setup for the leptogenesis of matter-antimatter asymmetry where sufficient right-handed neutrinos are generated after reheating. This approach further provides a new framework of the cosmological Higgs collider, in which the Higgs-modulated reheating is naturally realized by the inflaton decays into right-handed neutrinos through neutrino seesaw.

In Sec. II, we studied the dynamics and evolution of the Higgs field during and after inflation. During the inflation, large quantum fluctuations of the Higgs field are generated due to the high scale of inflation $H_{\text{inf}} = O(10^{13}-10^{14})$ GeV. If the inflation lasts long enough, the distribution of the Higgs field would finally reach an equilibrium state and the average of the Higgs field takes a value around the Hubble scale, as shown in Eq. (2.16). After inflation, the value of Higgs field $h(t)$ would oscillate and decrease. We also gave a

semianalytic formula (2.20) for the evolution of the Higgs field after inflation, which fits well with the exact numerical calculations as demonstrated in Fig. 3.

In Sec. III, we presented an approach that incorporates both the inflation and neutrino seesaw mechanism. The inflaton ϕ and the right-handed neutrino N_R can couple together through a unique dimension-five operator as in Eq. (3.1) that respects the shift symmetry. In this approach, the inflaton primarily decays into right-handed neutrinos after inflation. With the neutrino seesaw mechanism, the Higgs field can influence the reheating process through the modulation of inflaton decays into right-handed neutrinos. In consequence, the curvature perturbations are generated by fluctuations in the Higgs field. We established a relation (3.12) between the curvature perturbations and the Higgs fluctuations, as illustrated in Fig. 4. A further expanded formula for this relation up to $O(\delta h_{\text{inf}}^2)$ is given in Eq. (4.1).

In Sec. IV, we investigated the primordial local non-Gaussianity arising from the Higgs-modulated reheating through the inflaton decays into right-handed neutrinos within the seesaw mechanism. We demonstrated that our method provides observable signatures of primordial non-Gaussianity, which can be used to probe the neutrino seesaw mechanism in our framework. We performed a full analysis for the two- and three-point correlation functions of the comoving curvature perturbation ζ_h as contributed by the Higgs fluctuations in our model. We found that the primordial non-Gaussianity is contributed by both the Higgs self-interaction term and the nonlinear term, where the former can be dominant and was not considered in the literature. We presented the predictions for the local non-Gaussianity $f_{\text{NL}}^{\text{local}}$ from the parameter space of neutrino seesaw, as shown in Figs. 6 and 7. We found that, for fairly modest values of Higgs self-coupling, the sensitivities to the non-Gaussianity of $f_{\text{NL}}^{\text{local}} = O(1)$ could probe the seesaw scale $M \sim 10^{13}$ GeV and can also have important interplay with probing the light neutrino mass scale and mass ordering in the low-energy neutrino experiments. We further studied the dependence of the non-Gaussianity on Higgs self-coupling λ (as shown in Fig. 8). We found that

the non-Gaussianity measurements are also sensitive to the SM Higgs self-coupling λ at the neutrino seesaw scale and thus provide complementary probes of λ to the on-going LHC collider experiment.

In the near future, combining the neutrino data (such as those from the oscillation experiments JUNO [64] and DUNE [65] and the neutrinoless double- β decay experiments [66]) with the improved cosmological non-Gaussianity measurements (such as those from DESI [58], CMB-S4 [59], Euclid [60], SPHEREx [61], LSST [62], and SKA [63] experiments) is expected to provide a more sensitive probe of the high-scale neutrino seesaw mechanism through our approach. The main idea and results of this work are summarized in the companion Letter [68].

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DATA AVAILABILITY

The data supporting this study's findings are available within the article.

APPENDIX A: EVOLUTION OF HIGGS FIELD AFTER INFLATION

In this appendix, we provide further technical derivations for the evolution of the Higgs field after inflation to support the analysis of Sec. II. For the sake of convenience, we rescale all physical quantities according to their dimensions by the Hubble parameter H_{inf} during the inflation. Thus, they become dimensionless under the following rescaling¹³:

¹³In this rescaling, we have also included a numerical factor $\frac{3}{2}$ for convenience, such that the initial time t_0 defined below Eq. (A2) will simply become $t_0 = 1$.

$$h \rightarrow \left(\frac{3}{2}H_{\text{inf}}\right)h, \quad t \rightarrow \left(\frac{3}{2}H_{\text{inf}}\right)^{-1}t, \quad \dot{h} \rightarrow \left(\frac{3}{2}H_{\text{inf}}\right)^2\dot{h}. \quad (\text{A1})$$

After inflation and before the completion of reheating, we assume the Universe is matter dominated. Hence, the scale factor $a(t)$ and the Hubble parameter $H(t)$ take the following forms:

$$a(t) = \left(\frac{t}{t_0}\right)^{2/3}, \quad H(t) = \frac{2}{3t}, \quad (\text{A2})$$

where t_0 represents the initial time of matter domination, and the Hubble parameter at t_0 satisfies $H(t_0) = \frac{2}{3t_0} = H_{\text{inf}}$. The equation of motion for the Higgs field $h(t)$ is given by

$$\ddot{h}(t) + \frac{2}{t}\dot{h}(t) + \lambda h^3(t) = 0. \quad (\text{A3})$$

We can define the conformal time τ during the matter-dominated stage as follows:

$$\tau = \int \frac{dt}{a(t)} = 3t^{1/3}. \quad (\text{A4})$$

The definition of conformal time used here differs from that used in the calculation of correlation functions. During inflation, conformal time is always negative, with values approaching zero, which corresponds to an infinitely distant future. However, in the present context of a different cosmic expansion, conformal time is positive. These two definitions can be related by a simple constant shift. Thus, we can express the scale factor a in terms of conformal time τ as $a(\tau) = \frac{1}{9}\tau^2$. Based on this conformal time, we define

$$\begin{aligned} \varphi &\equiv ah = t^{2/3}h = \frac{1}{9}\tau^2h, \\ \varphi' &\equiv \frac{d\varphi}{d\tau}, \quad \varphi'' \equiv \frac{d^2\varphi}{d\tau^2}, \\ a' &\equiv \frac{da}{d\tau}, \quad a'' \equiv \frac{d^2a}{d\tau^2}. \end{aligned} \quad (\text{A5})$$

For the following derivations, we will use the two relations involving derivatives of the scale factor a ,

$$\frac{a'}{a} = \frac{2}{\tau}, \quad \frac{a''}{a} = \frac{2}{\tau^2}. \quad (\text{A6})$$

We then obtain the equation of motion for the redefined field $\varphi(\tau)$,

$$\varphi''(\tau) - \frac{a''}{a}\varphi(\tau) + \lambda\varphi(\tau)^3 = 0. \quad (\text{A7})$$

For a quantity A , we denote A_0 as $A(t = t_0) = A(\tau = \tau_0)$, where t_0 and τ_0 denote the initial physical time and its corresponding conformal time, respectively. We set the

initial conditions as follows:

$$t_0 = 1, \quad \tau_0 = 3, \quad a_0 = 1, \\ h_0 = h_{\text{inf}}, \quad \dot{h}_0 = 0, \quad \varphi_0 = h_0, \quad \varphi'_0 = \frac{2}{3}h_0. \quad (\text{A8})$$

In the above, $t_0 = 1$ implies $H(t_0) = H_{\text{inf}}$ and h_0 is the initial value of $h(t)$ which is determined at the end of inflation. We take $h_0 > 0$ as an example for the following calculation, but the results can be readily generalized to the case of $h_0 < 0$. When the conformal time is small, the term $-\frac{2}{\tau^2}\varphi$ dominates, and the other term $\lambda\varphi^3$ can be neglected for small λ . Hence, we could first solve the following equation:

$$\varphi'' - \frac{2}{\tau^2}\varphi = 0, \quad (\text{A9})$$

with the initial conditions $\varphi_0 = h_0$ and $\varphi'_0 = \frac{2}{3}h_0$. We thus derive the following solution:

$$\varphi(\tau) = \frac{1}{9}h_0\tau^2. \quad (\text{A10})$$

As the conformal time increases, the term $\lambda\varphi^3$ becomes increasingly significant. When $\lambda\varphi(\tau)^3 = \frac{2}{\tau^2}\varphi(\tau)$, we identify the conformal time cutoff τ_{cut} as follows:

$$\tau_{\text{cut}} = 3^{2/3}2^{1/6}\lambda^{-1/6}h_0^{-1/3}. \quad (\text{A11})$$

Then, we deduce $\varphi_{\text{cut}} \equiv \varphi(\tau_{\text{cut}}) = \frac{1}{9}h_0\tau_{\text{cut}}^2$ and $\varphi'_{\text{cut}} \equiv \varphi'(\tau_{\text{cut}}) = \frac{2}{9}h_0\tau_{\text{cut}}$. Note that the solution (A10) is valid when $\tau \leq \tau_{\text{cut}}$.

At the late-time stage $\tau \gg \tau_{\text{cut}}$, the term $\lambda\varphi^3$ dominates and the term $-\frac{2}{\tau^2}\varphi$ may be neglected. In consequence, we simplify Eq. (A7) as follows:

$$\varphi''(\tau) + \lambda\varphi^3(\tau) = 0. \quad (\text{A12})$$

We can define a conserved quantity as the conformal “energy”:

$$\frac{1}{2}\varphi'^2 + \frac{\lambda}{4}\varphi^4 = E, \quad (\text{A13})$$

where we could define $E_K = \frac{1}{2}\varphi'^2$ as the conformal kinematic energy and $E_V = \frac{\lambda}{4}\varphi^4$ as the conformal potential energy. This equation implies that the field φ oscillates within the potential $\frac{\lambda}{4}\varphi^4$ without attenuation or damping. The solution to Eq. (A12) is given by an elliptic sine function, which can be approximated as follows:

$$\varphi(\tau) = \varphi_{\text{max}} \cos \left[\frac{\Gamma^2(3/4)}{\sqrt{\pi}} \sqrt{\lambda} \varphi_{\text{max}} (\tau - \tau_{\text{cut}}) + \theta \right], \quad (\text{A14})$$

where φ_{max} is the oscillation amplitude and θ is the phase, both of which are determined by initial conditions. Matching the energy at the time τ_{cut} , we could determine the phase θ and the oscillation amplitude φ_{max} ,

$$\begin{cases} \frac{E_K}{E_V} = 4 = \tan^2\theta, \\ E = E_K + E_V = \frac{5}{4}\varphi_{\text{cut}}^4 = \frac{\lambda}{4}\varphi_{\text{max}}^4, \end{cases} \Rightarrow \begin{cases} \theta = -\arctan 2 \simeq -1.1, \\ \varphi_{\text{max}} = 2^{1/3}3^{-2/3}5^{1/4} \left(\frac{h_0}{\lambda} \right)^{1/3}. \end{cases} \quad (\text{A15})$$

Then, we derive an analytic solution as a piecewise function,

$$\varphi(\tau) = \begin{cases} \frac{1}{9}h_0\tau^2, & \tau \leq \tau_{\text{cut}}, \\ \varphi_{\text{max}} \cos \left[\frac{\Gamma^2(3/4)}{\sqrt{\pi}} \sqrt{\lambda} \varphi_{\text{max}} (\tau - \tau_{\text{cut}}) + \theta \right], & \tau > \tau_{\text{cut}}. \end{cases} \quad (\text{A16})$$

With these, we derive the solution of Higgs field $h(t)$ in terms of physical time t as follows:

$$\begin{aligned} h(t) &= \frac{\varphi(3t^{1/3})}{a(t)} \\ &= \begin{cases} h_0, & t \leq t_{\text{cut}}, \\ A \left(\frac{h_0}{\lambda} \right)^{1/3} t^{-2/3} \cos \left[\lambda^{1/6} h_0^{1/3} \omega (t^{1/3} - t_{\text{cut}}^{1/3}) + \theta \right], & t > t_{\text{cut}}, \end{cases} \end{aligned} \quad (\text{A17})$$

where $t_{\text{cut}} = (\frac{1}{3}\tau_{\text{cut}})^3$, and the coefficients A and ω are given by

$$A = 2^{1/3}3^{-2/3}5^{1/4} \simeq 0.9, \quad (\text{A18a})$$

$$\omega = \frac{\Gamma^2(3/4)}{\sqrt{\pi}} 2^{1/3}3^{1/3}5^{1/4} \simeq 2.3. \quad (\text{A18b})$$

Using the definition of τ_{cut} in Eq. (A11), we could convert the term $\lambda^{1/6}h_0^{1/3}\omega t_{\text{cut}}^{1/3}$ of Eq. (A17) into a constant phase,

$$\lambda^{1/6}h_0^{1/3}\omega t_{\text{cut}}^{1/3} = \frac{1}{3}\lambda^{1/6}h_0^{1/3}\omega\tau_{\text{cut}} = 3^{-1/3}2^{1/6}\omega. \quad (\text{A19})$$

Finally, we further derive Eq. (A17) as follows:

$$h(t) = \begin{cases} h_0, & t \leq t_{\text{cut}}, \\ A \left(\frac{h_0}{\lambda} \right)^{1/3} t^{-2/3} \cos \left[\lambda^{1/6} h_0^{1/3} \omega t^{1/3} + \theta' \right], & t > t_{\text{cut}}, \end{cases} \quad (\text{A20})$$

where the phase $\theta' = -\lambda^{1/6}h_0^{1/3}\omega t_{\text{cut}}^{1/3} + \theta \simeq -2.9$. Recall that we employ the factor $\frac{3}{2}H_{\text{inf}}$ to define dimensionless physical quantities in Eq. (A1). We can recover the

dimensions of relevant physical variables by the following rescaling:

$$\begin{aligned} t &\rightarrow \left(\frac{3}{2}H_{\text{inf}}\right)t, & t_0 &\rightarrow \left(\frac{3}{2}H_{\text{inf}}\right)t_0, & t_{\text{cut}} &\rightarrow \left(\frac{3}{2}H_{\text{inf}}\right)t_{\text{cut}}, \\ h(t) &\rightarrow h(t)/\left(\frac{3}{2}H_{\text{inf}}\right), & h_0 &\rightarrow h_0/\left(\frac{3}{2}H_{\text{inf}}\right), \end{aligned} \quad (\text{A21})$$

under which the form of the formula (A20) remains unchanged.

APPENDIX B: NON-GAUSSIANITY FROM THREE-POINT CORRELATION FUNCTIONS

The Planck Collaboration gives a list of primordial non-Gaussianity measurements with various shapes [8]. For the present study, we will use the shape templates of local-, equilateral-, and orthogonal-type non-Gaussianity measurements.

For the local-type non-Gaussianity, we compute the three-point correlation function,

$$\begin{aligned} \langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle'_{\text{local}} &= B_{\Phi}^{\text{local}}(k_1, k_2, k_3) \\ &= 2A^2 f_{\text{NL}}^{\text{local}} \left(\frac{1}{k_1^3 k_2^3} + \frac{1}{k_2^3 k_3^3} + \frac{1}{k_3^3 k_1^3} \right) \\ &\equiv 2A^2 f_{\text{NL}}^{\text{local}} F^{\text{local}}(k_1, k_2, k_3), \end{aligned} \quad (\text{B1})$$

where Φ denotes the Bardeen gravitational potential, which is related to the comoving curvature perturbation ζ via $\Phi \equiv \frac{3}{5}\zeta$ on superhorizon scales. The power spectrum $P_{\Phi} = \langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \rangle'$ is given by $P_{\Phi} = A/k^3$, where A is the normalization constant. Given $\Phi = \frac{3}{5}\zeta$ and $P_{\Phi} = \langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \rangle' = A/k^3$ (under the approximation $n_s \simeq 1$), we derive the normalization constant $A = 2\pi^2 (\frac{3}{5})^2 P_{\zeta}$. In the above formula and hereafter, we have defined $\langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle'$ as the three-point correlation function without including the δ function of momentum conservation, $\langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle = (2\pi)^3 \delta^3(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) \langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle'$.

For the equilateral-type non-Gaussianity, the three-point correlation function is computed as follows:

$$\begin{aligned} \langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle'_{\text{equil}} &= B_{\Phi}^{\text{equil}}(k_1, k_2, k_3) \\ &= 6A^2 f_{\text{NL}}^{\text{equil}} \left[-\frac{1}{k_1^3 k_2^3} - \frac{1}{k_2^3 k_3^3} - \frac{1}{k_3^3 k_1^3} \right. \\ &\quad \left. - \frac{2}{(k_1 k_2 k_3)^2} + \left(\frac{1}{k_1 k_2^2 k_3^2} + 5 \text{ perm.} \right) \right] \\ &\equiv 6A^2 f_{\text{NL}}^{\text{equil}} F^{\text{equil}}(k_1, k_2, k_3). \end{aligned} \quad (\text{B2})$$

For the orthogonal-type non-Gaussianity, we can derive the three-point correlation function as follows:

$$\begin{aligned} \langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle'_{\text{ortho}} &= B_{\Phi}^{\text{ortho}}(k_1, k_2, k_3) \\ &= 6A^2 f_{\text{NL}}^{\text{ortho}} \left[-\frac{3}{k_1^3 k_2^3} - \frac{3}{k_2^3 k_3^3} - \frac{3}{k_3^3 k_1^3} \right. \\ &\quad \left. - \frac{8}{(k_1 k_2 k_3)^2} + \left(\frac{3}{k_1 k_2^2 k_3^2} + 5 \text{ perm.} \right) \right] \\ &\equiv 6A^2 f_{\text{NL}}^{\text{ortho}} F^{\text{ortho}}(k_1, k_2, k_3). \end{aligned} \quad (\text{B3})$$

Next, we can use the above templates to analyze a three-point correlation function of the curvature perturbation $\langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \zeta_{\mathbf{k}_3} \rangle'$ through

$$\begin{aligned} \langle \Phi_{\mathbf{k}_1} \Phi_{\mathbf{k}_2} \Phi_{\mathbf{k}_3} \rangle' &= \left(\frac{3}{5} \right)^3 \langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \zeta_{\mathbf{k}_3} \rangle' \\ &= 2A^2 [f_{\text{NL}}^{\text{local}} F^{\text{local}}(k_1, k_2, k_3) \\ &\quad + 3f_{\text{NL}}^{\text{equil}} F^{\text{equil}}(k_1, k_2, k_3) \\ &\quad + 3f_{\text{NL}}^{\text{ortho}} F^{\text{ortho}}(k_1, k_2, k_3)], \end{aligned} \quad (\text{B4})$$

where $F^{\text{local}}(k_1, k_2, k_3)$, $F^{\text{equil}}(k_1, k_2, k_3)$, and $F^{\text{ortho}}(k_1, k_2, k_3)$ are functions of the external momenta (k_1, k_2, k_3) . In the present analysis, we find that the three-point correlation function $\langle \zeta_{\mathbf{k}_1} \zeta_{\mathbf{k}_2} \zeta_{\mathbf{k}_3} \rangle'$ can be expressed as the sum of these three shape templates with relevant coefficients f_{NL} .

We can directly solve the coefficients f_{NL} as follows:

$$\begin{aligned} f_{\text{NL}}^{\text{equil}} &= \frac{5}{12} P, & f_{\text{NL}}^{\text{ortho}} &= -\frac{1}{12} P, \\ f_{\text{NL}}^{\text{local}} &= \frac{1}{2} \left[P \left(N_e - \gamma + \frac{7}{3} \right) + Q \right], \end{aligned} \quad (\text{B5})$$

where the quantities P and Q are given by

$$P = -\frac{20z_1^3 H^3}{3(2\pi)^4 P_{\zeta}^2} \frac{\lambda \bar{h}}{2H}, \quad (\text{B6a})$$

$$Q = \frac{20z_1^3 H^3}{3(2\pi)^4 P_{\zeta}^2} \frac{z_2 H}{4z_1}. \quad (\text{B6b})$$

We may also compare our results with the literature [46–48,50], which considered the special case of $R = 1$ and ignored the contribution from Higgs self-coupling. In this case, given $R = |z_1|(P_{\delta h}/P_{\zeta})^{1/2}$, we obtain $(P_{\delta h}/P_{\zeta})^{1/2} = R/|z_1|$. Under the assumption that Higgs fluctuation δh is scale invariant, i.e., $P_{\delta h}^{1/2} = H/(2\pi)$, we deduce the following result:

$$f_{\text{NL}}^{\text{local}} = \frac{Q}{2} = \frac{5z_1^3 H^3}{6(2\pi)^4 P_{\zeta}^2} \frac{z_2 H}{z_1} = \frac{5P_{\delta h}^2}{6P_{\zeta}^2} z_1^2 z_2 = R^4 \frac{5z_2}{6z_1^2}. \quad (\text{B7})$$

Using $z_1 = -\Gamma'_0/(6\Gamma_0)$ and $z_2 = -(\Gamma_0\Gamma''_0 - \Gamma'_0\Gamma'_0)/(6\Gamma_0^2)$ as defined in Sec. IV A, and the condition that $R = 1$, we thus derive the local non-Gaussianity,

$$f_{\text{NL}}^{\text{local}} = \frac{5z_2}{6z_1^2} = 5 \left(1 - \frac{\Gamma''_0\Gamma_0}{(\Gamma'_0)^2} \right), \quad (\text{B8})$$

which agrees with the literature [46–48,50].

APPENDIX C: ANALYSIS OF NG DEPENDENCE ON THE HIGGS SELF-COUPLING

The local-type non-Gaussianity $f_{\text{NL}}^{\text{local}}$ in Eq. (B5) consists of two distinct contributions: the Higgs self-coupling term $f_{\text{NL}}^{\text{local}}(\text{HSC})$ generated by the three-point correlation function of Higgs self-interactions in Eq. (4.10) and the nonlinear term $f_{\text{NL}}^{\text{local}}(\text{NT})$ arising from the substitution of $\delta h(\mathbf{k})$ with the nonlinear component $\frac{1}{2}z_2\delta h^2$ in Eq. (4.14). These contributions are expressed as follows:

$$f_{\text{NL}}^{\text{local}}(\text{HSC}) = -\frac{5H_{\text{inf}}^2}{3(2\pi)^4 P_\zeta^2} \left(N_e - \gamma + \frac{7}{3} \right) z_1^3 \lambda \bar{h}, \quad (\text{C1a})$$

$$f_{\text{NL}}^{\text{local}}(\text{NT}) = \frac{5H_{\text{inf}}^4}{6(2\pi)^4 P_\zeta^2} z_1^2 z_2. \quad (\text{C1b})$$

In this appendix, we give an analysis to estimate the λ dependence of the above non-Gaussianity contributions $f_{\text{NL}}^{\text{local}}(\text{HSC})$ and $f_{\text{NL}}^{\text{local}}(\text{NT})$.

1. Dependence of $f_{\text{NL}}^{\text{local}}(\text{HSC})$ on Higgs self-coupling

The λ dependence of $f_{\text{NL}}^{\text{local}}(\text{HSC})$ in Eq. (C1a) arises through the three parameters $z_1(\lambda)$, λ , and $\bar{h}(\lambda)$,

$$f_{\text{NL}}^{\text{local}}(\text{HSC}) \propto \lambda z_1^3(\lambda) \bar{h}(\lambda). \quad (\text{C2})$$

In the above, the root-mean-square value of the Higgs field $\bar{h}$ is given by Eq. (2.16), from which we deduce

$$\bar{h} = \left[\int_{-\infty}^{+\infty} dh h^2 \rho_{\text{eq}}(h) \right]^{\frac{1}{2}} \simeq 0.363 \left(\frac{H_{\text{inf}}}{\lambda^{1/4}} \right) \propto \lambda^{-\frac{1}{4}}. \quad (\text{C3})$$

The parameter z_1 is defined in Eq. (4.4) as

$$z_1 = -\frac{1}{6} \frac{\Gamma'_0}{\Gamma_0}, \quad (\text{C4})$$

where Γ_0 is defined in Eq. (4.2) with $h_{\text{inf}} = \bar{h}$. Neglecting kinematic factors, we can approximate the total decay rate as in Eq. (3.7),

$$\Gamma_{\text{reh}} \simeq \frac{m_\phi M^2}{4\pi\Lambda^2} \left[1 + \frac{1}{4} \left(\frac{y_\nu h_{\text{reh}}}{M} \right)^2 \right]. \quad (\text{C5})$$

From the semianalytical solution for the Higgs field $h(t)$ in Eq. (2.20), and after averaging over a sufficiently large volume such that the factor $\cos^2(\omega_{\text{reh}} h_{\text{inf}}^{1/3} + \theta)$ is treated as $\frac{1}{2}$, we find that h_{reh} scales as follows:

$$h_{\text{reh}}(\lambda) \propto h_{\text{inf}}^{\frac{1}{3}} \lambda^{-\frac{1}{3}}. \quad (\text{C6})$$

Given the condition $y_\nu h_{\text{reh}}/M \ll 1$ in the parameter space under consideration, the λ dependence of Γ_0 in Eq. (C5) can be neglected for this estimate. The remaining λ dependence in z_1 arises from Γ'_0 , which is the derivative of Γ_{reh} with respect to h_{inf} . Using Eqs. (C5) and (C6) for computing Γ'_0 , we obtain the λ dependence of Γ'_0 and z_1 ,

$$\begin{aligned} z_1 \propto \Gamma'_0(\lambda) &\propto \frac{d\Gamma_{\text{reh}}}{dh_{\text{reh}}} \frac{\partial h_{\text{reh}}}{\partial h_{\text{inf}}} \Big|_{h_{\text{inf}}=\bar{h}} \\ &\propto h_{\text{reh}}(\lambda) \bar{h}^{-\frac{2}{3}} \lambda^{-\frac{1}{3}} \propto \lambda^{-\frac{7}{2}}, \end{aligned} \quad (\text{C7})$$

where in the last step we have used the mean-field approximation $h_{\text{inf}} = \bar{h}$ with Eq. (C3).

Given the λ dependence of $z_1(\lambda)$ in Eq. (C7) and $\bar{h}(\lambda)$ in Eq. (C3), we can thus deduce the λ dependence of the non-Gaussianity $f_{\text{NL}}^{\text{local}}(\text{HSC})$,

$$f_{\text{NL}}^{\text{local}}(\text{HSC}) \propto \lambda z_1^3(\lambda) \bar{h}(\lambda) \propto \lambda^{-1}. \quad (\text{C8})$$

This explains why the predicted values of $f_{\text{NL}}^{\text{local}}(\text{HSC})$ reduce by about a factor $\frac{1}{2}$ as the Higgs self-coupling varies from $\lambda = 0.01$ to $\lambda = 0.02$, as shown in the fifth row of Table II.

2. Dependence of $f_{\text{NL}}^{\text{local}}(\text{NT})$ on Higgs self-coupling

Next, we study the λ dependence of $f_{\text{NL}}^{\text{local}}(\text{NT})$ from the nonlinear term contribution. From Eq. (C1b), we see that $f_{\text{NL}}^{\text{local}}(\text{NT})$ depends on λ through product $z_1^2 z_2$. Since λ dependence of z_1 is already given by Eq. (C7), we will focus on the analysis of the λ dependence of z_2 . Equation (4.4) gives the following expression of z_2 :

$$z_2 = -\frac{1}{6} \left[\frac{\Gamma''_0}{\Gamma_0} - \left(\frac{\Gamma'_0}{\Gamma_0} \right)^2 \right], \quad (\text{C9})$$

which shows that the λ dependence of z_2 mainly comes from $\Gamma''_0(\lambda)$ and $\Gamma'_0(\lambda)$. The λ dependence of Γ'_0 is given in Eq. (C7). Thus, we analyze the λ dependence of Γ''_0 , which can be derived as follows:

$$\begin{aligned}
\Gamma_0'' &= \left. \frac{d^2 \Gamma_{\text{reh}}}{dh_{\text{inf}}^2} \right|_{h_{\text{inf}}=\bar{h}} \\
&= \left. \frac{d^2 \Gamma_{\text{reh}}}{dh_{\text{reh}}^2} \left(\frac{\partial h_{\text{reh}}}{\partial h_{\text{inf}}} \right)^2 \right|_{h_{\text{inf}}=\bar{h}} + \left. \frac{d\Gamma_{\text{reh}}}{dh_{\text{reh}}} \left(\frac{\partial^2 h_{\text{reh}}}{\partial h_{\text{inf}}^2} \right) \right|_{h_{\text{inf}}=\bar{h}} \\
&= \left[C_1 \left(h_{\text{inf}}^{-\frac{2}{3}} \lambda^{-\frac{1}{3}} \right)^2 + C_2 h_{\text{reh}} h_{\text{inf}}^{-\frac{5}{3}} \lambda^{-\frac{1}{3}} \right] \Big|_{h_{\text{inf}}=\bar{h}} \propto \lambda^{-\frac{1}{3}}, \quad (\text{C10})
\end{aligned}$$

where the coefficients C_1 and C_2 are independent of the Higgs self-coupling λ , and Eqs. (C6), (C5), and (C3) are used.

From Eq. (C5), we can deduce the scaling behaviors

$$\frac{\Gamma_0''}{\Gamma_0} \sim \frac{1}{M^2}, \quad \left(\frac{\Gamma_0'}{\Gamma_0} \right)^2 \sim \frac{1}{M^2} \frac{h_{\text{reh}}^2}{M^2}. \quad (\text{C11})$$

Since $h_{\text{reh}}/M \ll 1$, we can neglect the λ dependence from the second term $(\Gamma_0'/\Gamma_0)^2$ of Eq. (C9) in comparison to its first term Γ_0''/Γ_0 . Hence, we can extract the λ dependence of z_2 as follows:

$$z_2(\lambda) \propto \Gamma_0'' \propto \lambda^{-\frac{1}{3}}. \quad (\text{C12})$$

Using Eqs. (C7) and (C12), we can derive the λ dependence of $f_{\text{NL}}^{\text{local}}(\text{NT})$ as follows:

$$f_{\text{NL}}^{\text{local}}(\text{NT}) \propto z_1^2 z_2 \propto \lambda^{-\frac{3}{2}}. \quad (\text{C13})$$

In summary, we have estimated the λ dependence of both contributions $f_{\text{NL}}^{\text{local}}(\text{HSC})$ and $f_{\text{NL}}^{\text{local}}(\text{NT})$ as follows:

$$f_{\text{NL}}^{\text{local}}(\text{HSC}) \propto \lambda^{-1}, \quad f_{\text{NL}}^{\text{local}}(\text{NT}) \propto \lambda^{-\frac{3}{2}}. \quad (\text{C14})$$

where we have used the condition $y_\nu h_{\text{reh}}/M \ll 1$ and have neglected kinematic factors in the inflaton decay rate Γ_{reh} . Hence, summed contribution to the non-Gaussianity $f_{\text{NL}}^{\text{local}} = f_{\text{NL}}^{\text{local}}(\text{HSC}) + f_{\text{NL}}^{\text{local}}(\text{NT})$ decreases monotonically as the Higgs self-coupling λ increases. These analytical behaviors agree well with the numerical results presented in Table II and Fig. 8. We also note that the above scaling behaviors of Eq. (C14) are derived by using the mean-field approximation, which are valid for the parameter space under consideration, but not for arbitrarily large or small coupling λ .

APPENDIX D: SCHWINGER-KELDYSH PROPAGATORS OF HIGGS FIELD

The Klein-Gordon equation in the de Sitter spacetime for the massless scalar field is given as follows:

$$\square u_{\mathbf{k}} = \ddot{u}_{\mathbf{k}} + 3H\dot{u}_{\mathbf{k}} + \frac{\mathbf{k}^2}{a^2(t)} u_{\mathbf{k}} = 0, \quad (\text{D1})$$

where $H = \dot{a}/a$ is constant and $a(t) = e^{Ht}$. We may consider the conformal coordinates with $d\tau = dt/a(\tau)$ and $a = -(H\tau)^{-1}$, and reexpress the Klein-Gordon equation as follows:

$$u_{\mathbf{k}}'' - \frac{2}{\tau} u_{\mathbf{k}}' + \mathbf{k}^2 u_{\mathbf{k}} = 0. \quad (\text{D2})$$

The normalized mode function of the massless scalar field in de Sitter spacetime with the Bunch-Davies vacuum can be solved as

$$u_{\mathbf{k}}(\tau) = \frac{H}{\sqrt{2k^3}} (1 + ik\tau) e^{-ik\tau}, \quad (\text{D3})$$

where we denote $u' = du/d\tau$ and $u'' = d^2u/d\tau^2$.

In the Schwinger-Keldysh path integral formalism [51,52], the bulk-to-bulk propagators are defined as follows:

$$\begin{aligned}
G_{++}(\mathbf{k}; \tau_1, \tau_2) &= G_{>}(\mathbf{k}; \tau_1, \tau_2) \theta(\tau_1 - \tau_2) \\
&\quad + G_{<}(\mathbf{k}; \tau_1, \tau_2) \theta(\tau_2 - \tau_1), \quad (\text{D4a})
\end{aligned}$$

$$G_{+-}(\mathbf{k}; \tau_1, \tau_2) = G_{<}(\mathbf{k}; \tau_1, \tau_2), \quad (\text{D4b})$$

$$G_{-+}(\mathbf{k}; \tau_1, \tau_2) = G_{>}(\mathbf{k}; \tau_1, \tau_2), \quad (\text{D4c})$$

$$\begin{aligned}
G_{--}(\mathbf{k}; \tau_1, \tau_2) &= G_{<}(\mathbf{k}; \tau_1, \tau_2) \theta(\tau_1 - \tau_2) \\
&\quad + G_{>}(\mathbf{k}; \tau_1, \tau_2) \theta(\tau_2 - \tau_1), \quad (\text{D4d})
\end{aligned}$$

where we have defined

$$\begin{aligned}
G_{>}(\mathbf{k}; \tau_1, \tau_2) &= u_{\mathbf{k}}(\tau_1) u_{\mathbf{k}}^*(\tau_2) \\
&= \frac{H^2}{2k^3} [1 + ik(\tau_1 - \tau_2) + k^2 \tau_1 \tau_2] e^{-ik(\tau_1 - \tau_2)}, \quad (\text{D5a})
\end{aligned}$$

$$G_{<}(\mathbf{k}; \tau_1, \tau_2) = u_{\mathbf{k}}^*(\tau_1) u_{\mathbf{k}}(\tau_2) = G_{>}^*(\mathbf{k}; \tau_1, \tau_2). \quad (\text{D5b})$$

Additionally, the bulk-to-boundary propagator is a special propagator in which external legs terminated at the final slice, i.e., $\tau = \tau_f \rightarrow 0^-$,

$$G_{\pm}(\mathbf{k}, \tau) = G_{\pm+}(\mathbf{k}; \tau, \tau_f), \quad (\text{D6})$$

which means that the bulk-to-boundary propagator ($\tau = \tau_1 < \tau_2 = \tau_f \rightarrow 0_-$) is defined based on the above bulk-to-bulk propagator. The bulk-to-boundary propagator includes one plus-type $G_{+}(\mathbf{k}, \tau)$ and one minus-type $G_{-}(\mathbf{k}, \tau)$,

$$\tau \bullet \text{---} \square = G_{+}(\mathbf{k}, \tau), \quad (\text{D7a})$$

$$\tau \circ \text{---} \square = G_{-}(\mathbf{k}, \tau), \quad (\text{D7b})$$

where the square at one end of the propagator indicates its boundary point ($\tau_f \rightarrow 0^-$), and a black dot and a white dot denote plus and minus types, respectively. In addition, a shaded dot (often representing a vertex) indicates that the possibilities from both the plus- and minus-type propagators should be summed up,

$$\tau \text{---} \text{---} \square = G_+(\mathbf{k}, \tau) + G_-(\mathbf{k}, \tau). \quad (\text{D8})$$

The plus- and minus-type bulk-to-boundary propagators take the following forms:

$$G_+(\mathbf{k}, \tau) = \frac{H^2}{2k^3} [1 - ik(\tau - \tau_f) + k^2 \tau \tau_f] e^{ik(\tau - \tau_f)} \\ \simeq \frac{H^2}{2k^3} [1 - ik\tau] e^{ik\tau}, \quad (\text{D9a})$$

$$G_-(\mathbf{k}, \tau) = \frac{H^2}{2k^3} [1 + ik(\tau - \tau_f) + k^2 \tau \tau_f] e^{-ik(\tau - \tau_f)} \\ \simeq \frac{H^2}{2k^3} [1 + ik\tau] e^{-ik\tau}. \quad (\text{D9b})$$

APPENDIX E: PROBABILITY DISTRIBUTION FUNCTION OF HIGGS FIELD

For the sake of convenience, we abbreviate the Hubble parameter during the inflation H_{inf} as H , namely, $H_{\text{inf}} = H$. A scalar field $\phi(\mathbf{x}, t)$ with only 1 degree of freedom can be represented as the mode functions of long-wavelength part $\phi_L(\mathbf{x}, t)$ and short-wavelength part $\phi_S(\mathbf{x}, t)$ [27,28],

$$\phi(\mathbf{x}, t) = \phi_L(\mathbf{x}, t) + \phi_S(\mathbf{x}, t) \\ = \phi_L(\mathbf{x}, t) + \int \frac{d^3k}{(2\pi)^3} \theta(k - \epsilon a(t)H) \\ \times [a_{\mathbf{k}} \phi_{\mathbf{k}}(t) e^{-i\mathbf{k} \cdot \mathbf{x}} + a_{\mathbf{k}}^\dagger \phi_{\mathbf{k}}^*(t) e^{i\mathbf{k} \cdot \mathbf{x}}]. \quad (\text{E1})$$

The short-wavelength modes $\phi_{\mathbf{k}}(t)$ are initially subhorizon. Over the time, as the Universe expands, these modes are stretched and eventually cross the physical cutoff $\epsilon a(t)H$, transitioning into the superhorizon modes ϕ_L . The superhorizon modes ϕ_L can be effectively treated as a classical stochastic field, obeying the Langevin equation

$$\dot{\phi}_L(\mathbf{x}, t) = -\frac{1}{3H} \frac{\partial V}{\partial \phi_L} + f(\mathbf{x}, t). \quad (\text{E2})$$

In addition to the force from the potential of the field, the long-wavelength modes are also driven by an effective stochastic force f , which is generated by freezing out the short-wavelength modes,

$$f(\mathbf{x}, t) = \int \frac{d^3k}{(2\pi)^3} \delta(k - \epsilon a(t)H) \epsilon a(t) \\ \times H^2 [a_{\mathbf{k}} \phi_{\mathbf{k}}(t) e^{-i\mathbf{k} \cdot \mathbf{x}} + a_{\mathbf{k}}^\dagger \phi_{\mathbf{k}}^*(t) e^{i\mathbf{k} \cdot \mathbf{x}}]. \quad (\text{E3})$$

In order to obtain the Fokker-Planck equation for the one-point probability distribution function, we should derive the two-point correlation function of the stochastic force f . For this, we could derive the following relation for a canonical massless field $\phi_{\mathbf{k}}$ in the late-time limit:

$$\left\langle \left(a_{\mathbf{k}_1} \phi_{\mathbf{k}_1}(t_1) e^{-i\mathbf{k}_1 \cdot \mathbf{x}_1} + a_{\mathbf{k}_1}^\dagger \phi_{\mathbf{k}_1}^*(t_1) e^{i\mathbf{k}_1 \cdot \mathbf{x}_1} \right) \right. \\ \left. \times \left(a_{\mathbf{k}_2} \phi_{\mathbf{k}_2}(t_2) e^{-i\mathbf{k}_2 \cdot \mathbf{x}_2} + a_{\mathbf{k}_2}^\dagger \phi_{\mathbf{k}_2}^*(t_2) e^{i\mathbf{k}_2 \cdot \mathbf{x}_2} \right) \right\rangle \\ = (2\pi)^3 \delta^3(\mathbf{k}_1 - \mathbf{k}_2) e^{i\mathbf{k}_1 \cdot \mathbf{x}_{12}} \frac{H^2}{2k_1^3}. \quad (\text{E4})$$

Then, we can derive the two-point correlation function of f as follows:

$$\langle f(\mathbf{x}_1, t_1) f(\mathbf{x}_2, t_2) \rangle = \int \frac{d^3k_1}{(2\pi)^3} \left(\prod_{i=1}^2 \delta(k_i - \epsilon a(t_i)H) \epsilon a(t_i) H^2 \right) \frac{H^2}{2k_1^3} e^{i\mathbf{k}_1 \cdot \mathbf{x}_{12}} \\ = \frac{1}{4\pi^2} \int_0^{+\infty} dk_1 \left(\prod_{i=1}^2 \delta(k_i - \epsilon a(t_i)H) \epsilon a(t_i) H^2 \right) \frac{H^2}{k_1} j_0(k_1 x_{12}) \\ = \frac{H^5 \epsilon j_0(\epsilon a(t_1)H x_{12}) \delta(\epsilon a(t_1)H - \epsilon a(t_2)H)}{4\pi^2 a(t_1)} \prod_{i=1}^2 a(t_i), \quad (\text{E5})$$

where $\mathbf{x}_{12} = \mathbf{x}_1 - \mathbf{x}_2$, $x_{12} = |\mathbf{x}_{12}|$, and $j_0(z) = (\sin z)/z$ is the spherical Bessel function of zeroth order. Using the property of the δ function, we can simplify the above formula for the two-point correlation function of the stochastic force:

$$\langle f(\mathbf{x}_1, t_1) f(\mathbf{x}_2, t_2) \rangle = \frac{H^3}{4\pi^2} j_0(\epsilon a(t_1)H x_{12}) \delta(t_1 - t_2). \quad (\text{E6})$$

If we take the one-point limit $\mathbf{x}_2 \rightarrow \mathbf{x}_1$ in the position space with $j_0(\epsilon a H x_{12}) \rightarrow 1$, the one-point probability

distribution function for the field ϕ obeys the Fokker-Planck equation

$$\frac{\partial \rho(\phi, t)}{\partial t} = \frac{1}{3H} \left\{ \rho(\phi, t) \frac{\partial^2 V(\phi)}{\partial \phi^2} + \frac{\partial V(\phi)}{\partial \phi} \frac{\partial \rho(\phi, t)}{\partial \phi} \right\} + \frac{H^3}{8\pi^2} \frac{\partial^2 \rho(\phi, t)}{\partial \phi^2}. \quad (\text{E7})$$

To solve Eq. (E7), we decompose the solution $\rho(\phi, t)$ in terms of the eigenfunctions $\{\Psi_n\}$ with eigenvalues $\{\Lambda_n\}$,

$$\rho(\phi, t) = e^{-v} \sum_{n=0}^{\infty} a_n \Psi_n(\phi) e^{-\Lambda_n(t-t_0)}, \quad (\text{E8})$$

where v is defined as

$$v(\phi) \equiv \frac{4\pi^2 V(\phi)}{3H^4}, \quad (\text{E9})$$

and the coefficients a_n could be given by the initial condition of $\rho(t=t_0)$,

$$a_n = \int d\phi \rho(\phi, t_0) e^{v(\phi)} \Psi_n(\phi). \quad (\text{E10})$$

Then, we derive the corresponding equation for the eigenfunctions Ψ_n and eigenvalues Λ_n by substituting Eq. (E8) into Eq. (E7):

$$\left[\frac{\partial^2}{\partial \phi^2} + \frac{\partial^2 v}{\partial \phi^2} - \left(\frac{\partial v}{\partial \phi} \right)^2 \right] \Psi_n(\phi) = -\frac{8\pi^2}{H^3} \Lambda_n \Psi_n(\phi), \quad (\text{E11})$$

which is a typical Sturm-Liouville equation. Thus, all of the eigenvalues Λ_n are non-negative and the eigenfunctions $\Psi_n(\phi)$ are orthonormal functions,

$$\int_{-\infty}^{+\infty} d\phi \Psi_n(\phi) \Psi_{n'}(\phi) = \delta_{n,n'}. \quad (\text{E12})$$

They satisfy the completeness condition,

$$\sum_n \Psi_n(\phi) \Psi_n(\phi_0) = \delta(\phi - \phi_0). \quad (\text{E13})$$

There is a special eigenfunction with the eigenvalue $\Lambda_0 = 0$,

$$\Psi_0(\phi) = N_0 e^{-v(\phi)}, \quad (\text{E14})$$

where N_0 is the normalization constant. Thus, Eq. (E8) can be reexpressed as follows:

$$\rho(\phi, t) = \Psi_0(\phi) \sum_{n=0}^{\infty} b_n \Psi_n(\phi) e^{-\Lambda_n(t-t_0)}, \quad (\text{E15})$$

where coefficients b_n are also the normalization factors analogous to a_n , defined as $b_n \equiv a_n/N_0$. If the inflation lasts long enough, the system quickly approaches its equilibrium state. The only term that remains in the above summation is $\Psi_0(\phi)$ with $\Lambda_0 = 0$. Thus, the equilibrium probability distribution ρ_{eq} is derived as

$$\rho_{\text{eq}}(\phi) = \Psi_0(\phi)^2 = N_0^2 \exp\left(-\frac{8\pi^2 V(\phi)}{3H^4}\right), \quad (\text{E16})$$

where the equilibrium probability distribution satisfies the normalization condition,

$$\int d\phi \rho_{\text{eq}}(\phi) = 1. \quad (\text{E17})$$

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