Letter

Cosmological signatures of neutrino seesaw mechanism

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The tiny neutrino masses are most naturally explained by seesaw mechanism through singlet right-handed neutrinos, which can further explain the matter-antimatter asymmetry in the Universe. In this work, we propose a new approach to study cosmological signatures of neutrino seesaw through the interaction between inflaton and right-handed neutrinos, which respects the shift symmetry. In our framework, after inflation the inflaton predominantly decays into right-handed neutrinos and its decay rate is modulated by the fluctuations of the Higgs field, which act as the source of curvature perturbations. This gives a new realization of Higgs modulated reheating, and it produces primordial non-Gaussian signatures that can be measured by the forthcoming large-scale structure surveys. We demonstrate, for the first time, that the measurements of non-Gaussianity by these surveys can probe the neutrino seesaw scale and have the potential to constrain a large portion of the seesaw parameter space, opening up a new window for directly testing the high-scale seesaw mechanism.

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Introduction. Understanding the origin of tiny neutrino masses of $O(0.1\ \mathrm{eV})$ poses a major challenge to the standard model (SM) of particle physics. Neutrino seesaw provides the most natural explanation of tiny neutrino masses by including the singlet right-handed neutrinos [1,2], and it further explains the matter-antimatter asymmetry (baryon asymmetry) in the Universe via leptogenesis [3]. But, the natural scale of the seesaw mechanism is around 10^{14} GeV for the Higgs-neutrino Yukawa couplings of $O(1)$. Probing such a high scale of neutrino seesaw is extremely difficult and far beyond the reach of current particle experiments.

In contrast, inflation provides the most appealing mechanism for dynamics of the early Universe, during which the Universe underwent a short period of rapid exponential expansion that resolves the flatness and horizon problems as well as simultaneously generating the primordial fluctuations for seeding the large-scale structures of the Universe. The energy scale of inflation could be as high as 10^{16} GeV, characterized by a nearly constant Hubble parameter H_{inf} around 10^{14} GeV, which coincides with the scale of the neutrino seesaw. Inflation is typically driven by a scalar field known as inflaton. The primordial fluctuations arise from the inflaton's quantum fluctuations and can be directly measured through their contributions to the cosmic microwave background (CMB). The current CMB data indicate that these fluctuations are predominantly adiabatic and Gaussian. However, during inflation the primordial perturbations could also exhibit non-Gaussianity (NG) [4,5]. The NG is sensitive to new physics effects at high energy scales. Although the current CMB observations only set a weak limit on the NG parameter $f_{\mathrm{NL}} = O(10)$ [6], the upcoming experiments will improve detection sensitivity to the level of $f_{\mathrm{NL}} = O(0.01)$ [5,7,8], opening up an important window for probing the high-scale new physics.

We note that the neutrino seesaw scale M with natural Yukawa couplings $[y_\nu = O(1)]$ is around 10^{14} GeV, which coincides with the upper range of inflation scale. Thus, the seesaw mechanism could leave distinctive imprints in the cosmological evolution. It is natural to expect that

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the inflaton couples directly to the right-handed neutrinos and decays predominantly into them after inflation. Then the right-handed neutrinos further decay into the SM particles via Yukawa interactions, completing the reheating process. Moreover, during inflation the Higgs field acquires a value near the Hubble scale, varying across different horizon patches. This variation leads to space-dependent right-handed neutrino masses via the seesaw mechanism, which modulate the inflaton decay rate into right-handed neutrinos. With these, we give a *new realization of Higgs modulated reheating* and provide a source of primordial curvature perturbations [9]. In this letter, we construct the inflaton coupling to right-handed neutrinos through a dimension-5 operator respecting shift symmetry. We investigate the effects of Higgs-modulated reheating and the associated NG signatures, with which we demonstrate the potential to probe the high-scale neutrino seesaw within our framework. We map the non-Gaussianity $f_{\rm NL}^{\rm local}$ onto the (y_ν, M) plane, which shows sensitivity to probing the light neutrino mass and interplays with the low energy neutrino oscillation experiments. We will further show the sensitivity of the NG measurement to the SM Higgs self-coupling at the inflation scale, which is quantitatively connected to the Higgs self-coupling at the TeV scale (through the renormalization group evolution). Hence we may establish the interplay between the Higgs self-coupling constraints at the inflation scale and the Higgs self-coupling measurements at the TeV scale of the LHC.

Dynamics of higgs field during and after inflation. During inflation, the Universe is effectively a de Sitter spacetime. The dynamics of a spectator Higgs field in this de Sitter spacetime can be described through the stochastic inflation approach [10,11]. In the unitary gauge, the Higgs field is given by $\mathbb{H} = \frac{1}{\sqrt{2}}(0, h)^T$. The potential of the SM Higgs field during inflation is $V(h) = \frac{1}{4}\lambda h^4$, where its mass term can be neglected and the Higgs self-coupling λ could have a value of $O(0.01)$ at the inflation scale within the 3σ range of the current top mass measurement [12,13]. During inflation, the long-wave mode of the Higgs field value h can be described as a classical motion with stochastic noise,

$$\dot{h}(\mathbf{x}, t) = -\frac{1}{3H_{\rm inf}}\frac{\partial V}{\partial h} + f(\mathbf{x}, t), \tag{1}$$

where $H_{\rm inf}$ is the Hubble parameter during inflation, and $f(\mathbf{x}, t)$ is a stochastic background and has the two-point correlation function,

$$\langle f(\mathbf{x}_1, t_1) f(\mathbf{x}_2, t_2)\rangle = \frac{H_{\rm inf}^3}{4\pi^2} j_0(\epsilon a(t_1) H_{\rm inf}|\mathbf{x}_{12}|)\delta(t_1 - t_2), \tag{2}$$

where $j_0(z) = (\sin z)/z$ and $\mathbf{x}_{12} = \mathbf{x}_1 - \mathbf{x}_2$. If inflation lasts long enough, the distribution of the Higgs field would eventually reach an equilibrium with a probability function,

$$\rho_{\rm eq}(h) = \frac{2\lambda^{1/4}}{\Gamma(1/4)}\left(\frac{2\pi^2}{3}\right)^{1/4}\exp\left(\frac{-2\pi^2\lambda h^4}{3H_{\rm inf}^4}\right). \tag{3}$$

The root-mean-square value of the Higgs field $\bar{h} = \sqrt{\langle h^2\rangle}$ is derived as follows:

$$\bar{h} = \left[\int_{-\infty}^{+\infty} {\rm d}h h^2 \rho_{\rm eq}(h)\right]^{1/2} \simeq 0.363\left(\frac{H_{\rm inf}}{\lambda^{1/4}}\right). \tag{4}$$

After inflation, we consider the inflaton potential as quadratic near its minimum, the inflaton oscillates and behaves like a matter component ($w = 0$). Consequently, the Universe expands as $a \sim t^{2/3}$, with the Hubble parameter given by $H = 2/(3t)$. The evolution of the super-horizon mode of the Higgs field h after inflation is governed by the Klein-Gordon equation,

$$\ddot{h}(t) + \frac{2}{t}\dot{h}(t) + \lambda h^3(t) = 0. \tag{5}$$

Thus, for $t \gg (\sqrt{\lambda} h_{\rm inf})^{-1}$ and $h_{\rm inf} > 0$, we derive the evolution of $h(t)$ as follows [15]:

$$h(t) = A\left(\frac{h_{\rm inf}}{\lambda}\right)^{\frac{1}{3}} t^{-\frac{2}{3}}\cos\left(\lambda^{\frac{1}{6}} h_{\rm inf}^{\frac{1}{3}}\omega t^{\frac{1}{3}} + \theta\right), \tag{6}$$

where $h_{\rm inf}$ is the Higgs field value at the end of inflation, which varies in different Hubble patches, and the parameters (A, ω, θ) are given by

$$A = \left(\frac{2}{9}\right)^{\frac{1}{3}} 5^{\frac{1}{4}} \simeq 0.9, \qquad \omega = \frac{\Gamma^2(3/4)}{\sqrt{\pi}} 6^{\frac{1}{3}} 5^{\frac{1}{4}} \simeq 2.3,$$
$$\theta = -3^{-\frac{1}{3}} 2^{\frac{1}{6}}\,\omega - \arctan 2 \simeq -2.9. \tag{7}$$

The solution can be readily generalized to the case of $h_{\rm inf} < 0$. Equation (6) shows that after inflation, the Higgs field oscillates in its quartic potential $\frac{1}{4}\lambda h^4$, but its oscillation amplitude will gradually decrease over time [16].

Inflaton-neutrino interaction and inflaton decay. The right-handed neutrinos N_R can couple to the inflaton ϕ through a unique dimension-5 operator, which has a cutoff scale Λ and respects the inflaton's shift symmetry [17]. Thus, we construct the minimal model incorporating both inflation and neutrino seesaw with the following Lagrangian:

$$\begin{aligned}\Delta\mathcal{L} = \sqrt{-g}\Bigg[&-\frac{1}{2}\partial_\mu\phi\partial^\mu\phi - V(\phi) + \overline{N}_R {\rm i}\partial\!\!\!/ N_R \\ &+ \frac{1}{\Lambda}\partial_\mu\phi\overline{N}_R\gamma^\mu\gamma^5 N_R \\ &+ \left(-\frac{1}{2}M\overline{N_{\rm R}^{\rm c}}N_R - y_\nu\overline{L}_{\rm L}\widetilde{\mathbb{H}}N_R + {\rm H.c.}\right)\Bigg],\end{aligned} \tag{8}$$

where $V(\phi)$ is the inflaton potential and $L_{\rm L} = (\nu_{\rm L}, e_{\rm L})^T$ is the left-handed lepton doublet. After inflation we consider the inflaton mass term dominates the potential $V(\phi)$

under which the inflaton ϕ will oscillate. Due to the shift symmetry, inflaton couples to the right-handed neutrinos through the dimension-5 operator of Eq. (8) [21]. The perturbative unitarity imposes a lower bound on its cutoff scale, $\Lambda \gtrsim 60H_{\rm inf}$. In our setup, this dimension-5 operator causes the inflaton to decay predominantly into right-handed neutrinos after inflation. If the inflaton couples to the SM fermions via dimension-5 operators and with the shift symmetry, the corresponding decay rates are suppressed by the fermion masses, as the Higgs field quickly decreases after inflation. (This also applies to the case of inflaton coupling to top quarks.) 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 and thus ensure inflaton to mainly decay into the right-handed neutrinos.

For simplicity, we focus on analyzing the case of one family of fermions. For the neutrino seesaw with $|y_\nu h| \ll M$, the two mass eigenstates ν and N have masses,

$$m_\nu = -\frac{y_\nu^2 h^2}{2M}, \qquad M_N = M + \frac{y_\nu^2 h^2}{2M}. \tag{9}$$

The rotation angle θ for this mass diagonalization is given by $\tan\theta \simeq y_\nu h/(\sqrt{2}M)$. In Eq. (9), the heavy neutrino mass M_N has a shift $y_\nu^2 h^2/(2M)$ from M, which is crucial for our mechanism as we are really probing the seesaw effect on the heavy neutrino mass eigenvalue.

For $|y_\nu h| \ll M$, the inflaton decay rate into neutrinos is given by

$$\Gamma \simeq \frac{m_\phi M^2}{4\pi\Lambda^2}\left[1+\frac{1}{4}\left(\frac{y_\nu h}{M}\right)^2\right], \tag{10}$$

where kinetic factors are ignored for simplicity, but $m_\phi > 2M_N$ is always required to ensure that the inflaton decay channel $\phi \to NN$ is kinematically open. We see that the inflaton decay rate depends on the Higgs field value h [22]. On the other hand, Refs. [24–26] studied cosmological collider (CC) signals, which manifest as oscillatory features in the primordial NG—a qualitatively different type of NG signal from what we study. In particular, Ref. [25] explored the CC signal arising from a coupling similar to our Eq. (8), this would provide a complementary probe within our model framework.

Curvature perturbation from Higgs modulated reheating. In our approach, the inflaton decay rate is affected by the value of the SM Higgs field. The variation of the Higgs field $h(\mathbf{x}, t_{\rm reh})$ leads to a spatial variation of the decay rate $\Gamma_{\rm reh}(\mathbf{x})$. It perturbs the local expansion history, seeding large-scale inhomogeneity and anisotropy. These fluctuations can be described by the δN formalism [27–35]. The number of e-folds of the cosmic expansion after inflation can be computed as [36],

$$\begin{aligned} N(\mathbf{x}) &= \int \mathrm{d}\ln a(t) = \int_{t_{\rm end}}^{t_{\rm reh}(\mathbf{x})} \mathrm{d}tH(t) + \int_{t_{\rm reh}(\mathbf{x})}^{t_{\rm f}} \mathrm{d}tH(t) \\ &= \int_{\rho_{\rm end}}^{\rho_{\rm reh}(h(\mathbf{x}))} \mathrm{d}\rho\frac{H}{\dot\rho} + \int_{\rho_{\rm reh}(h(\mathbf{x}))}^{\rho_{\rm f}} \mathrm{d}\rho\frac{H}{\dot\rho}, \end{aligned} \tag{11}$$

where $a(t)$ is the scale factor and $\rho(t)$ is the total energy density of the Universe at the time t. The curvature perturbation during reheating, $\zeta(\mathbf{x}, t)$, is equal to the $\delta N(\mathbf{x}, t)$ of cosmic expansion among different Hubble patches in the uniform energy density gauge,

$$\zeta_h(\mathbf{x}, t) = \delta N(\mathbf{x}, t) = N(\mathbf{x}, t) - \langle N(\mathbf{x}, t)\rangle. \tag{12}$$

For this study, we describe the Universe as a perfect fluid both before and after the completion of reheating. During the period $t_{\rm end} < t < t_{\rm reh}$, we consider the inflaton potential is dominated by its mass term. Thus, when the inflaton oscillates near the minimum of the potential, the Universe is matter dominated ($w = 0$). (Our approach also applies to the general case of $w \neq 1/3$.) For the period $t > t_{\rm reh}$, consider the right-handed neutrinos decay fast enough after being produced, so the Universe transitions to a radiation-dominated phase ($w = 1/3$). Using the equation of state $\dot\rho + 3H(1+w)\rho = 0$, the locally expanded e-folding number can be expressed as follows:

$$N(\mathbf{x}) = -\frac{1}{3}\ln\frac{\rho_{\rm reh}(h(\mathbf{x}))}{\rho_{\rm inf}} - \frac{1}{4}\ln\frac{\rho_f}{\rho_{\rm reh}(h(\mathbf{x}))}. \tag{13}$$

Using first Friedmann equation $3H^2M_p^2 = \rho$, and noting that reheating completes when $H(t_{\rm reh}) = \Gamma_{\rm reh}$ (where we take the sudden reheating approximation), we determine the curvature perturbation after reheating ($t > t_{\rm reh}$),

$$\begin{aligned} \zeta_h(\mathbf{x}, t > t_{\rm reh}) &= \delta N(\mathbf{x}) = N(\mathbf{x}) - \langle N(\mathbf{x})\rangle \\ &= -\frac{1}{12}[\ln\rho_{\rm reh}(\mathbf{x}) - \langle\ln\rho_{\rm reh}(\mathbf{x})\rangle] \\ &= -\frac{1}{6}[\ln(\Gamma_{\rm reh}) - \langle\ln(\Gamma_{\rm reh})\rangle]. \end{aligned} \tag{14}$$

Combined with the inflaton fluctuation $\delta\phi$ during inflation, the total comoving curvature perturbation is given by $\zeta = \zeta_\phi + \zeta_h$, where ζ_ϕ is generated by the inflaton fluctuation $\delta\phi$,

$$\zeta_\phi \simeq -\frac{H_{\rm inf}}{\dot\phi_0}\delta\phi(\mathbf{x}), \tag{15}$$

and ζ_h originates from the effect of 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,

$$\mathcal{P}_\zeta = \mathcal{P}_\zeta^{(\phi)} + \mathcal{P}_\zeta^{(h)}, \tag{16}$$

where $\mathcal{P}_\zeta^{(\phi)}$ is the contribution of inflaton fluctuations,

$$\mathcal{P}_\zeta^{(\phi)} = \left(\frac{H_{\rm inf}}{\dot\phi}\right)^2 \mathcal{P}_\phi = \left(\frac{H_{\rm inf}}{\dot\phi}\right)^2 \frac{H_{\rm inf}^2}{4\pi^2}. \tag{17}$$

We further define R as square root of the ratio between the power spectra of Higgs-modulated reheating and of the comoving curvature perturbation ζ,

$$R \equiv \left(\frac{\mathcal{P}_\zeta^{(h)}}{\mathcal{P}_\zeta^{(o)}}\right)^{1/2}, \tag{18}$$

where $\mathcal{P}_\zeta^{(o)} \simeq 2.1 \times 10^{-9}$ is the observed curvature perturbation [37,38]. To agree with the observation, we should require $R < 1$.

Modulated reheating can also provide a source of primordial NG. The primordial NG from the three-point correlation function of ζ is known as the bispectrum $\langle \zeta_{\mathbf{k}_1}\zeta_{\mathbf{k}_2}\zeta_{\mathbf{k}_3}\rangle$. To compute the n-point correlation function of ζ_h, we expand the curvature perturbation,

$$\zeta_h = \delta N = N'\delta h_{\rm inf} + \frac{1}{2}N''(\delta h_{\rm inf})^2 + \cdots, \tag{19}$$

where N' and N'' denote the first and second derivatives of the e-folding number N with respect to $h_{\rm inf}$. The expansion allows us to determine the amplitude of the curvature perturbations as $\mathcal{P}_\zeta^{(h)} = N'^2\mathcal{P}_{h_{\rm inf}}$ and the primordial local NG $f_{\rm NL}^{\rm local}$ [39–43].

We note that when reheating occurs, the value of the Higgs field is an oscillatory function of its initial value,

$$h(t_{\rm reh}, h_{\rm inf}) \propto h_{\rm inf}^{\frac{1}{3}} \cos\left(\omega_{\rm reh} h_{\rm inf}^{\frac{1}{3}} + \theta\right), \tag{20}$$

where the oscillating frequency is given by

$$\omega_{\rm reh} = \lambda^{\frac{1}{6}} t_{\rm reh}^{\frac{1}{3}} \omega. \tag{21}$$

When $t_{\rm reh}$ is large, the oscillation frequency can become very high. Note that ζ_h can be expanded into the form $A + Bh^2/M^2 + O(h^4/M^4)$, which includes a factor $\cos^2(\omega_{\rm reh} h_{\rm inf}^{1/3} + \theta)$. Since $h_{\rm inf}$ varies across different Hubble volumes and ζ_h is highly sensitive to $h_{\rm inf}$, averaging over a sufficiently large volume makes the factor $\cos^2(\omega_{\rm reh} h_{\rm inf}^{1/3} + \theta)$ be effectively $1/2$ [44]. Thus, for the following analysis we set $\cos^2(\omega_{\rm reh} h_{\rm inf}^{1/3} + \theta) \to 1/2$ in the expression of ζ_h.

Bispectrum from Higgs fluctuations. We expand the curvature perturbation in terms of the Higgs fluctuation as follows:

$$\zeta_h(\mathbf{x}) = z_1 \delta h_{\rm inf}(\mathbf{x}) + \frac{1}{2} z_2 \delta h_{\rm inf}^2(\mathbf{x}), \tag{22}$$

where the coefficients z_1 and z_2 are given by

$$z_1 = -\left.\frac{\Gamma'}{6\Gamma}\right|_{h_{\rm inf}=\bar h}, \qquad z_2 = \left.\frac{\Gamma'\Gamma' - \Gamma\Gamma''}{6\Gamma^2}\right|_{h_{\rm inf}=\bar h}, \tag{23}$$

with Γ' (Γ'') being the first (second) derivative of Γ respect to $h_{\rm inf}$. In the following, we abbreviate the Hubble scale during inflation $H_{\rm inf}$ as H and the Higgs field value during inflation $h_{\rm inf}$ as h, which should differ from the Higgs field value $h(t)$ after inflation. 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,

$$\langle \zeta_{\mathbf{k}_1}\zeta_{\mathbf{k}_2}\zeta_{\mathbf{k}_3}\rangle_h = z_1^3 \langle \delta h_{\mathbf{k}_1}\delta h_{\mathbf{k}_2}\delta h_{\mathbf{k}_3}\rangle + z_1^2 z_2 \langle \delta h^4\rangle(\mathbf{k}_1, \mathbf{k}_2, \mathbf{k}_3). \tag{24}$$

On the right-hand side of the equality in Eq. (24), the first term is the three-point correlation function of the Higgs fluctuation $\delta h(\mathbf{k})$ generated by the self-interactions of the Higgs field, whereas 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.

During inflation, the Higgs fluctuation δh could be treated as a nearly massless scalar. Due to the SM Higgs self-coupling term $\Delta\mathcal{L} = -\sqrt{-g}[(\lambda\bar h)\delta h^3]$ and according to the Schwinger-Keldysh (SK) path integral formalism [45,46], the three-point correlation function of the Higgs fluctuation δh, $\langle \delta h_{\mathbf{k}_1}\delta h_{\mathbf{k}_2}\delta h_{\mathbf{k}_3}\rangle'$ [47], can be computed through the following integral:

$$\langle \delta h_{\mathbf{k}_1}\delta h_{\mathbf{k}_2}\delta h_{\mathbf{k}_3}\rangle' = 12\lambda\bar h\, {\rm Im}\left(\int_{-\infty}^{\tau_f} {\rm d}\tau a^4 \prod_{i=1}^{3} G_+(\mathbf{k}_i, \tau)\right), \tag{25}$$

where $G_\pm(\mathbf{k}_i, \tau)$ is the bulk-to-boundary propagator of massless scalar in the SK path integral [46]. In Eq. (25), we denote the integral part ${\rm Im}(\cdots) \equiv A$ and compute it to the leading order of $k_t\tau_f$,

$$A = \frac{H^2}{24 k_1^3 k_2^3 k_3^3}\left\{(k_1^3 + k_2^3 + k_3^3)\left[\ln(k_t|\tau_f|) + \gamma - \frac{4}{3}\right] + k_1k_2k_3 - \sum_{a\neq b} k_a^2 k_b\right\}, \tag{26}$$

where $\gamma \simeq 0.577$ is the Euler-Mascheroni constant, and the wave number $k_t = k_1 + k_2 + k_3$ is around the scale of the present observable Universe. Here $H \equiv H_{\rm inf}$ for simplicity. For the second term on the right-hand side of Eq. (24), 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,

$$z_1^2 z_2 \langle \delta h^4\rangle(\mathbf{k_1}, \mathbf{k_2}, \mathbf{k_3}) = (2\pi)^3\delta^3(\mathbf{k_1} + \mathbf{k_2} + \mathbf{k_3}) z_1^2 z_2 \left(\frac{H^4}{4k_1^3 k_2^3} + 2\,{\rm perm.}\right). \tag{27}$$

For this study, we mainly focus on the magnitude of local NG predicted by our model, which can be approximated as follows:

$$f_{\rm NL}^{\rm local} \simeq -\frac{10}{3}\frac{z_1^3 H^3}{(2\pi)^4 \mathcal{P}_\zeta^2}\left(\frac{\lambda \bar{h}}{2H}N_e - \frac{z_2 H}{4 z_1}\right), \quad (28)$$

where N_e is the e-folding number corresponding to the present Universe,

$$N_e = \ln\frac{a_{\rm end}}{a_k} = \ln\frac{-(H\tau_f)^{-1}}{k_t/H} = -\ln(k_t|\tau_f|) \sim 60. \quad (29)$$

We find that the contribution from the Higgs self-coupling can be dominant (not studied before), whereas nonlinear term contribution is non-negligible.

Probing neutrino seesaw using non-Gaussianity. In our analysis, 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 [37,38]. We set the SM Higgs self-coupling $\lambda = 0.01$, and the Hubble parameter $H_{\rm inf} = 10^{13}$ GeV (or, 3×10^{13} GeV). We also set the inflaton mass $m_\phi = 40 H_{\rm inf}$ and the cutoff scale $\Lambda = 60 H_{\rm inf}$. (The effects due to the variation of inputs will be shown in Table I.) With these inputs, we present our findings in Figs. 1 and 2.

In Fig. 1, the colored region obeys the condition $R < 1$ and the white region in the upper-right corner corresponds to $R \geq 1$. The region with blue color corresponds to $f_{\rm NL}^{\rm local} > 0$, whereas the red regions represent $f_{\rm NL}^{\rm local} < 0$. The green contours describe the 2σ bounds on $f_{\rm NL}^{\rm local}$ from Planck-2018 data, $-11.1 \leq f_{\rm NL}^{\rm local} \leq 9.3$ [6]. We further present contours for $f_{\rm NL}^{\rm local} = \pm 1, \pm 0.1, \pm 0.01$, which are plotted as orange, yellow, and white curves, respectively. These contours represent sensitivity reaches by the future experiments. We see that the local-type NG measurements for $f_{\rm NL}^{\rm local} > 0$ and $f_{\rm NL}^{\rm local} < 0$ can probe different seesaw parameter space of the (y_ν, M) plane, so their probes are complementary.

In Fig. 1, we set two benchmarks for the light neutrino mass $m_\nu = 0.06$ eV and 0.05 eV [48], shown as the pink and purple curves, respectively, for the Hubble parameter $H_{\rm inf} = 10^{13}$ GeV (solid curves) and 3×10^{13} GeV (dashed curves). We see that a larger Hubble parameter shifts the pink and purple curves toward the regions with larger Yukawa coupling y_ν. For the case with $f_{\rm NL}^{\rm local} > 0$, we see that for a light neutrino mass $m_\nu = 0.06(0.05)$ eV and $H_{\rm inf} = 3 \times 10^{13}$ GeV, the existing Planck-2018 data already excluded part of the parameter space as shown by the green contour. For a smaller Hubble parameter $H_{\rm inf} = 10^{13}$ GeV, our predictions of the local-type NG are beyond the reach of Planck-2018, but will be largely probed by the future measurements with improved sensitivities of $f_{\rm NL}^{\rm local} = \pm 1, \pm 0.1, \pm 0.01$ (shown by the orange, yellow, and white contours).

TABLE I. Comparison of the non-Gaussianity predictions in our framework for three sets of benchmark points with specific cutoff scales Λ and inflaton masses m_ϕ.

Benchmarks	A_1	A_2	B_1	B_2	C_1	C_2	D_1	D_2
$\Lambda/H_{\rm inf}$	60	60	70	70	80	80	100	100
$m_\phi/H_{\rm inf}$	30	40	30	40	30	40	30	40
$f_{\rm NL}^{\rm local}$	.34	1.8	.10	.54	.034	.18	.01	.031

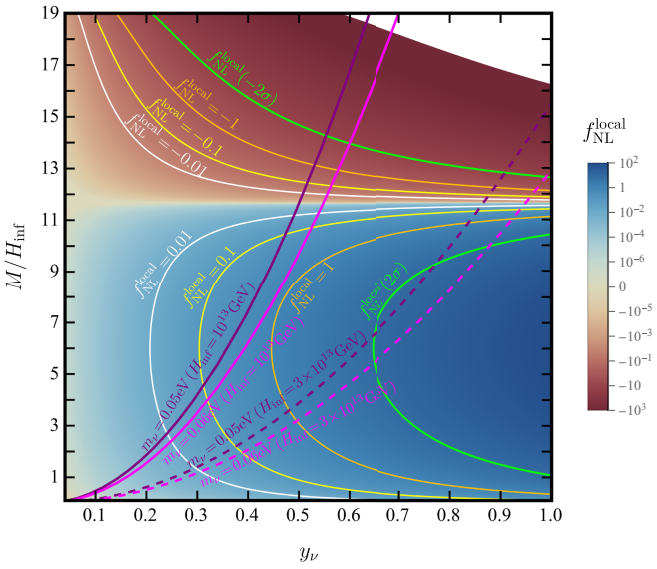


FIG. 1. Prediction of the non-Gaussianity (NG) $f_{\rm NL}^{\rm local}$ from the seesaw parameter space of the heavy neutrino mass scale M versus Yukawa coupling y_ν.

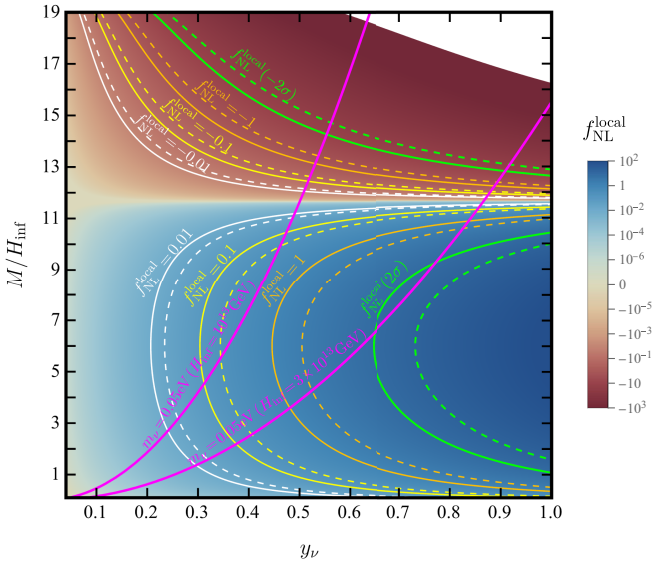


FIG. 2. Prediction of the non-Gaussianity $f_{\rm NL}^{\rm local}$ from the seesaw parameter space of heavy neutrino mass scale M versus the Yukawa coupling y_ν, where we input the SM Higgs self-coupling $\lambda = 0.01$ (solid curves) and $\lambda = 0.02$ (dashed curves), and the Hubble parameter during inflation is set as $H_{\rm inf} = 10^{13}$ GeV and 3×10^{13} GeV respectively.

We note that the low-energy neutrino oscillation data provide $\Delta m^2_{13} \simeq 2.5 \times 10^{-3}\,\mathrm{eV}^2$ and $\Delta m^2_{12} \simeq 7.4 \times 10^{-5}\ \mathrm{eV}^2$ [51], requiring at least one of the light neutrinos has mass $m_\nu \gtrsim 0.05$ eV. Moreover, the cosmological measurements can place an upper bound on the sum of light neutrino masses. Combining this with the neutrino oscillation measurements on the mass-squared differences can determine the upper limits of the light neutrino masses for either normal ordering (NO) or inverted ordering (IO). For instance, cosmological measurements based on the CMB alone already set a 95% upper limit, $\sum m_\nu \lesssim 0.26$ eV [38]. Combining this with the observations of large-scale structure, eBOSS Collaboration [52] placed a 95% upper bound $\sum m_\nu \lesssim 0.10$ eV and DES Collaboration [53] set a constraint $\sum m_\nu \lesssim 0.13$ eV at 95% CL. Combining the tighter bound $\sum m_\nu \lesssim 0.13$ eV with the neutrino oscillation data [51], we find the largest light neutrino mass to be $m_3 \simeq 0.06$ eV for the NO and $m_2 \simeq 0.05$ eV for the IO. Figure 1 shows that the current and future measurements of the local-type NG are sensitive to probing the difference between the cases with light neutrino mass $m_\nu = 0.06$ eV (pink curves) versus $m_\nu = 0.05$ eV (purple curves). The forthcoming oscillation experiments such as JUNO [54] and DUNE [55] are expected to determine the neutrino mass ordering, and give stronger constraints on the allowed light neutrino masses. It is encouraging to see that using the NG measurements to probe neutrino seesaw around the inflation scale (in Fig. 1) could also have sensitivity to the light neutrino masses and their ordering. Hence this may interplay with the low-energy oscillation experiments such as JUNO and DUNE.

Additionally, for the Higgs modulated reheating, the local NG also depends on the SM Higgs self-coupling λ, which can be used to probe the Higgs self-coupling. In order to test the sensitivity of $f^{\mathrm{local}}_{\mathrm{NL}}$ to the SM Higgs self-coupling λ, we vary the value of λ and present the $f^{\mathrm{local}}_{\mathrm{NL}}$ contours at 2σ level in Fig. 2. The $f^{\mathrm{local}}_{\mathrm{NL}}$ contours in solid (dashed) curves correspond to the Higgs self-coupling $\lambda = 0.01(0.02)$. The seesaw predictions are presented by pink curves for the light neutrino mass of $m_\nu = 0.05$ eV, with the Hubble parameter $H = 10^{13}$ GeV and $H = 3 \times 10^{13}$ GeV, respectively. In Fig. 2, with a larger Higgs self-coupling value $\lambda = 0.02$, the non-Gaussianity contours (in dashed curves) impose weaker bounds on the seesaw parameter space of (y_ν, M) as compared to the contours (in solid curves) with a smaller coupling $\lambda = 0.01$.[1] This analysis shows that the NG measurements of $f^{\mathrm{local}}_{\mathrm{NL}}$ 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 [56]) via the renormalization group evolution. Hence, this also demonstrates the interplay on probing the Higgs self-coupling λ between the high-scale cosmological NG measurements and the TeV-scale collider measurements.

In Table I, we show the dependence of the NG on different values of the cutoff scale Λ and inflaton mass m_ϕ. For illustration, we choose a sample input of neutrino seesaw scale $M = 5H_{\mathrm{inf}}$ (with $H_{\mathrm{inf}} = 10^{13}$ GeV), Higgs-neutrino Yukawa coupling $y_\nu = 0.5$, and the SM Higgs self-coupling $\lambda = 0.01$. Such sample inputs may indicate certain parameter degeneracy in NG signals. But the Hubble scale during inflation (H_{inf}) could be determined from other measurements in principle (such as the tensor-to-scalar ratio). The mass of N_R may be inferred from the cosmological collider signatures, whereas the light neutrino masses can be measured by the ongoing and future low-energy neutrino experiments. Although Λ and m_ϕ are more closely related to the UV physics and thus more challenging to determine, such degeneracy could be potentially resolved by analyzing the higher-point correlation functions (such as the four-point function). In principle, measuring more detailed information from higher-point correlators would help to further pin down the underlying parameter space and to distinguish our scenario from other models generating local non-Gaussianity.

The ongoing and forthcoming measurements on the non-Gaussianity (such as those from DESI [57], CMB-S4 [58], Euclid [59], SPHEREx [60], LSST [61], and SKA [62] experiments) will further probe the origin of neutrino mass generation through the seesaw mechanism around the inflation scale. We stress that once the local non-Gaussinity is observed, the associate cosmological collider signals [24–26] can be used to discriminate our model from other local non-Gaussinity sources. A systematic expansion of this article is presented in the companion longer paper of Ref. [15].

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

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