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Signature of primordial non-Gaussianity of ϕ^3 -type in the mass function and bias of dark matter haloes

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We explore the effect of a cubic correction $g_{\text{NL}}\phi^3$ on the mass function and bias of dark matter haloes extracted from a series of large N-body simulations and compare it to theoretical predictions. Such cubic terms can be motivated in scenarios like the curvaton model, in which a large cubic correction can be produced while simultaneously keeping the quadratic $f_{\text{NL}}\phi^2$ correction small. The deviation from the Gaussian halo mass function is in reasonable agreement with the theoretical predictions. The scale-dependent bias correction $\Delta b_\kappa(k, g_{\text{NL}})$ measured from the auto- and cross-power spectrum of haloes, is similar to the correction in f_{NL} models, but the amplitude is lower than theoretical expectations. Using the compilation of LSS data in Slosar et al. [JCAP, **08**, 031 (2008)], we obtain for the first time a limit on g_{NL} of $-3.5 \times 10^5 < g_{\text{NL}} < +8.2 \times 10^5$ (at 95% CL). This limit will improve with the future LSS data by 1-2 orders of magnitude, which should test many of the scenarios of this type.

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I. INTRODUCTION

In standard single field inflation, primordial curvature perturbations are produced by the inflaton field as it slowly rolls down its potential. Most of these models predict a nearly scale-invariant spectrum of adiabatic curvature fluctuations in agreement with cosmological observations. In addition, very small deviations from Gaussianity are expected. Therefore, any evidence for or against the detection of primordial non-Gaussianity would strongly constrain inflationary scenarios.

Non-Gaussianity can be produced by nonlinearities in the relation between the primordial curvature perturbation Φ (Here and henceforth, the usual Bardeen potential in matter-dominated era) and the inflaton field, interactions of scalar fields, a modified dispersion relation or a departure from the natural adiabatic vacuum state (see [1] for a review). Any non-Gaussianity that is generated outside the horizon induces a three-point function (or bispectrum) $B_\Phi(\mathbf{k}_1, \mathbf{k}_2, \mathbf{k}_3)$ that is peaked on squeezed triangles (i.e. $k_1 \ll k_2 \sim k_3$) for realistic values of the scalar spectral index. The resulting non-Gaussianity depends only on the local value $\Phi(\mathbf{x})$ of the Bardeen's curvature potential and can thus be conveniently parametrised up to third order by

$$\Phi(\mathbf{x}) = \phi(\mathbf{x}) + f_{\text{NL}} [\phi^2(\mathbf{x}) - \langle \phi^2 \rangle] + g_{\text{NL}} \phi^3(\mathbf{x}), \quad (1)$$

where $\phi(\mathbf{x})$ is an isotropic Gaussian random field and f_{NL} , g_{NL} are dimensionless, phenomenological parameters. While the quadratic term generates the irreducible three-point function or bispectrum at leading order, the cubic term does so for the irreducible four-point function or trispectrum. These statistics can be computed straightforwardly from a perturbative expansion of the homogeneous Robertson-Walker background

[2, 3]. Convolved with the appropriate transfer function (e.g. the radiation transfer function for the CMB temperature anisotropy), they can be used to constrain the value of the coupling parameters f_{NL} and g_{NL} . No significant detection of primordial non-Gaussianity has been reported from measurements of the three-point correlation function of the cosmic microwave background (CMB) anisotropies [4, 5, 6, 7, 8]. The tightest limits are $-4 < f_{\text{NL}} < 80$ at 95% confidence level [7].

If $\mathcal{O}(f_{\text{NL}}) \sim \mathcal{O}(g_{\text{NL}})$ then the cubic correction should always be negligibly small compared to the quadratic one since curvature perturbations are typically $\mathcal{O}(10^{-5})$. However, this condition is not satisfied by some multi-field inflationary models such as the curvaton scenario, in which a large g_{NL} and a small f_{NL} can be simultaneously produced. In this model, curvature perturbations are generated by an additional scalar field, the curvaton, whose energy density is negligible during inflation [9, 10, 11]. Non-Gaussianity is generated by curvaton self-interactions which effectively contribute a non-quadratic term to the curvaton potential [12, 13, 14, 15]. While the value and the sign of g_{NL} depend upon the exact form of the self-interaction term (which can dominate the mass term if the curvaton mass is small enough and the curvaton vacuum expectation value during inflation is large enough [16]), it is generically of magnitude $|g_{\text{NL}}| \sim 10^4 - 10^5$ for realistic models in which the ratio of the energy density of the curvaton to the total energy density at time of decay is small. There are other realizations where one can have large g_{NL} and small f_{NL} [17]. If f_{NL} were small, then the imprint of non-Gaussianity would be detected only in four-point statistics such as the CMB trispectrum [18, 19, 20]. Thus far, no observational limits have been set on g_{NL} by measuring the CMB trispectrum [21, 22]. Nevertheless, since the cur-

rent bound $|f_{\text{NL}}| \lesssim 100$ implies a relative contribution for the quadratic term of ~ 0.1 per cent, a third order coupling parameter $|g_{\text{NL}}| \sim 10^6$ should also be consistent with the data.

Large-scale structures offer another venue to test for the presence of primordial non-Gaussianity. Deviation from Gaussianity can significantly affect the high mass end of the mass function [23, 24], the large-scale two-point correlation [25, 26], the bispectrum [27, 28, 29, 30] of dark matter haloes hosting the observed galaxies as well as void abundances [31, 32] and topological measures of the cosmic web [33, 34]. Recently, references [35, 36, 37] showed that the local quadratic coupling $f_{\text{NL}}\phi^2$ induces a scale-dependent bias $\Delta b_\kappa(k, f_{\text{NL}})$ in the large-scale power spectrum of biased tracers,

$$\Delta b_\kappa(k, f_{\text{NL}}) = 3f_{\text{NL}} [b(M) - 1] \delta_c \frac{\Omega_m H_0^2}{k^2 T(k) D(z)}, \quad (2)$$

where $b(M)$ is the linear bias parameter, H_0 is the Hubble parameter, $T(k)$ is the matter transfer function normalised to unity as $k \rightarrow 0$, $D(z)$ is the growth factor normalised to $(1+z)^{-1}$ in the matter era and $\delta_c \sim 1.68$ is the present-day (linear) critical density threshold. Reference [37] applied eq. (2) to constrain the value of f_{NL} using a compilation of large-scale structure data and found $-29 < f_{\text{NL}} < +69$ at 95% confidence. These limits are comparable with those from the CMB, demonstrating the competitiveness of the method. Forthcoming all sky surveys should achieve constraints of the order of $f_{\text{NL}} \sim 1$ [35, 38, 39, 40] and should be sensitive to a possible scale-dependence of f_{NL} [41]. On the numerical side however, while simulations of structure formation have confirmed the scaling $\Delta b_\kappa(k, f_{\text{NL}})$ with k [35, 42, 43, 44], the exact amplitude of the non-Gaussian bias correction remains somewhat debatable.

All numerical studies to date have only implemented the quadratic term $f_{\text{NL}}\phi^2$. The purpose of this work is to quantify the impact of the cubic term $g_{\text{NL}}\phi^3$ on the mass function and bias of dark matter haloes extracted from cosmological simulations and assess the ability of forthcoming measurements of the large-scale bias of galaxies/quasars to constrain the size of a local cubic correction. This paper is organized as follows. We begin with a brief description of the N-body simulations and illustrate the extent to which the coupling $g_{\text{NL}}\phi^3$ affects the matter power spectrum and the halo mass function (Sec. II). We pursue with the non-Gaussian halo bias (Sec. III), to which we derive analytically the scale-dependent and scale-independent contribution, Δb_κ and Δb_{I} , and demonstrates the large suppression of the simulated Δb_κ relative to theory. We then place limits on the coupling parameter g_{NL} and forecast constraints from future large-scale surveys and CMB experiments (Sec. IV). We also show that our findings consistently apply to more general models with non-zero f_{NL} and g_{NL} (Sec. V). We conclude with a discussion of the results in Sec. VI.

II. THE NON-GAUSSIAN SIMULATIONS

A. Characteristics of the N-body runs

We utilize a series of large N-body simulations of the Λ CDM cosmology seeded with Gaussian and non-Gaussian initial conditions. The (dimensionless) power spectrum of the Gaussian part $\phi(\mathbf{x})$ of the Bardeen potential is the usual power-law $\Delta_\phi^2(k) \equiv k^3 P_\phi(k)/(2\pi^2) = A_\phi (k/k_0)^{n_s-1}$. The non-Gaussianity is of the ‘‘local’’ form $\Phi = \phi + g_{\text{NL}}\phi^3$. We adopt the standard (CMB) convention in which $\Phi(\mathbf{x})$ is primordial, and not extrapolated to present epoch. It is important to note that the local transformation is performed before multiplication by the matter transfer function. $T(k)$ is computed with CMBFAST [45] for the WMAP5 best-fitting parameters [6]: $h = 0.7$, $\Omega_m = 0.279$, $\Omega_b = 0.0462$, $n_s = 0.96$ and a normalisation of the Gaussian curvature perturbations $A_\phi = 7.96 \times 10^{-10}$ at the pivot point $k_0 = 0.02 \text{Mpc}^{-1}$. This yields a density fluctuations amplitude $\sigma_8 \approx 0.81$ when the initial conditions are Gaussian. Five sets of three 1024^3 simulations, each of which has $g_{\text{NL}} = 0, \pm 10^6$, were run with the N-body code GADGET2 [46]. We used the same Gaussian random seed field ϕ in each set of runs so as to minimise the sampling variance. We also explored scenarios with non-zero f_{NL} and g_{NL} and ran 2 realisations for each of the non-Gaussian models characterized by $(f_{\text{NL}}, g_{\text{NL}}) = (\pm 100, -3 \times 10^5)$. In all cases, the box size is $1600 h^{-1} \text{Mpc}$ with a force resolution of 0.04 times the mean interparticle distance. The particle mass of these simulations thus is $3.0 \times 10^{11} M_\odot/h$, enough to resolve haloes down to $10^{13} M_\odot/h$.

In the curvaton scenario, generic polynomial interaction terms of the form $\lambda m_\sigma^4 (\sigma/m_\sigma)^n$ (where λ is some coupling strength) to the quadratic potential of the curvaton field σ yield $|g_{\text{NL}}| \gg 1$ even when the non-linearity parameter f_{NL} is very small [13, 14]. One typically finds $|g_{\text{NL}}| \sim \mathcal{O}(10^4) - \mathcal{O}(10^5)$ when f_{NL} varies in the range $-100 < f_{\text{NL}} < 100$. For practical reasons however, the values of g_{NL} adopted in our simulations are about an order of magnitude larger so as to produce an effect strong enough to be unambiguously measured despite the small simulated volume. Furthermore, we have also considered positive and negative values of g_{NL} so as to assess the sensitivity of the non-Gaussian bias to the sign of the coupling parameter. The simulations with $(f_{\text{NL}}, g_{\text{NL}}) = (-100, -3 \times 10^5)$ may be seen as a particular realisation of the curvaton model in which the coupling constant λ is positive, and the non-quadratic term is very steep ($n \sim 5 - 10$) but contributes little to the total curvaton potential.

B. The matter power spectrum

Non-Gaussian corrections to the primordial curvature perturbation can renormalise the input power spectrum

of fluctuations used to seed the simulations. Since our simulations implement its unrenormalised version $\Delta_\phi^2(k) = A_\phi(k/k_0)^{n_s-1}$, it is desirable to ascertain the effect of the local coupling term on the simulated density power spectrum before discussing the halo mass function and bias. For f_{NL} models with $|f_{\text{NL}}| \lesssim 100$, renormalisation effects are unlikely to be noticeable due to the limited dynamical range of current cosmological simulations [47]. As we will see shortly however, they can be significant in simulations of g_{NL} models with similar level of non-Gaussianity.

The cubic order term $g_{\text{NL}}\phi^3$ renormalises the amplitude A_ϕ of the power spectrum of initial curvature perturbations to $A_\phi \rightarrow A_\phi + 6g_{\text{NL}}\langle\phi^2\rangle$, where

$$\langle\phi^2\rangle = \int \frac{d^3k}{(2\pi)^3} P_\phi(k). \quad (3)$$

For scale invariant initial conditions, $\langle\phi^2\rangle$ has a logarithmic divergence at large and small scales (see [47] for a more detailed discussion of perturbative corrections in non-Gaussian cosmologies). In practice, a low- and high- k cutoff are naturally provided by the finite box size and the resolution of the simulations. Therefore, the effective amplitude of density fluctuations in non-Gaussian simulations with cubic coupling is $\sigma_8 + \delta\sigma_8$ with

$$\begin{aligned} \delta\sigma_8 &= 3g_{\text{NL}}\langle\phi^2\rangle \\ &= 3g_{\text{NL}} \left(\frac{k_0}{k_{\text{min}}}\right)^{1-n_s} \left[1 - \left(\frac{k_{\text{min}}}{k_{\text{max}}}\right)^{1-n_s}\right] \frac{A_\phi}{1-n_s}. \end{aligned} \quad (4)$$

Recall that $k_0 = 0.02 \text{ hMpc}^{-1}$ is our choice of normalisation point, and k_{min} and k_{max} are the integration limits set by the fundamental mode and the Nyquist frequency of the periodic cubical box over which the initial conditions are generated. Equivalently,

$$\delta\sigma_8 = 3g_{\text{NL}} \left(\frac{Lk_0}{2\pi}\right)^{1-n_s} [1 - N^{n_s-1}] \frac{A_\phi}{1-n_s}, \quad (5)$$

where $N = 1024$ is the number of mesh points along one dimension. This result becomes $\delta\sigma_8 = 3g_{\text{NL}} \ln(N)A_\phi$ in the scale-invariant limit $n_s \rightarrow 1$. For the cosmological setup considered here, the absolute deviation is

$$\delta\sigma_8 \approx 0.015 \left(\frac{g_{\text{NL}}}{10^6}\right). \quad (6)$$

This correction is fairly large for the values of g_{NL} adopted here and, therefore, must be taken into account in the comparison between the theory and the simulations. As we will see below, this is especially important when studying the high mass tail of the halo mass function which is exponentially sensitive to the amplitude of density fluctuations.

The cubic coupling term $g_{\text{NL}}\phi^3$ can also induce a scale-dependent correction to the matter power spectrum which can be quantified by the fractional change $\beta_m(k, g_{\text{NL}}) = P_{\text{mm}}(k, g_{\text{NL}})/P_{\text{mm}}(k, g_{\text{NL}} = 0) - 1$. In

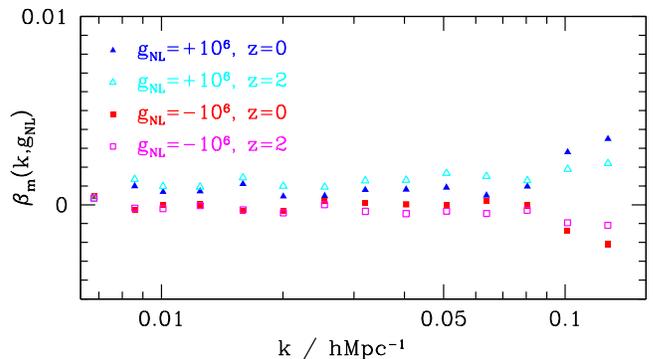


FIG. 1: Non-Gaussian fractional correction $\beta_m(k, g_{\text{NL}}) = P_{\text{mm}}(k, g_{\text{NL}})/P_{\text{mm}}(k, 0) - 1$ to the matter power spectrum after subtracting a scale-independent normalisation shift $6g_{\text{NL}}\langle\phi^2\rangle$ induced by the cubic coupling $g_{\text{NL}}\phi^3$.

Fig. 1, symbols show the result of measuring $\beta_m(k, g_{\text{NL}})$ from the snapshots at $z = 0$ and 2 after correction of the normalisation shift $|2\delta\sigma_8/\sigma_8| = 0.037$. There is some evidence for a scale-dependent correction at wavenumber $k \gtrsim 0.1 \text{ hMpc}^{-1}$ but the resulting deviation is broadly consistent with zero. We will thus neglect $\beta_m(k, g_{\text{NL}})$ henceforth.

C. The halo multiplicity function

Halo was identified using the MPI parallelised version of the AHF halo finder [48] which is based on the spherical overdensity (SO) finder developed by [49]. The main reason for using a SO finder is that it is more closely connected to the predictions of the spherical collapse model, on which most of the analytic formulae presented in this paper are based. Namely, the virial mass M of a halo is defined by the radius at which the inner overdensity exceeds $\Delta_{\text{vir}}(z)$ times the background density $\bar{\rho}(z)$ [50, 51]. The value of the overdensity threshold $\Delta_{\text{vir}}(z)$ is obtained from the collapse of a spherical tophat perturbation and has a dependence on redshift through the matter density $\Omega_m(z)$ [52, 53]. We discard poorly resolved haloes and only study those containing at least 34 particles or, equivalently, with a mass larger than $M = 10^{13} M_\odot/h$.

Analytic arguments based on the Press-Schechter theory [54, 55] predict that the halo mass function $n(M, z)$ is entirely specified by the distribution $\nu f(\nu)$ of first-crossings, or multiplicity function

$$\nu f(\nu) = M^2 \frac{n(M, z)}{\bar{\rho}} \frac{d \ln M}{d \ln \nu}. \quad (7)$$

The peak height $\nu(M, z) = \delta_c(z)/\sigma(M)$ is the typical amplitude of fluctuations that produce haloes of mass M by redshift z . Here, $\sigma(M)$ is the variance of the density field δ_M linearly extrapolated to present epoch and smoothed

on mass scale M , whereas $\delta_c(z) \approx 1.68D(0)/D(z)$ is the critical linear overdensity for (spherical) collapse at redshift z .

Despite the lack of a compelling theoretical description of the multiplicity function for Gaussian initial conditions, the fractional deviation from Gaussianity can be modelled accurately using the Press-Schechter formalism. In this approach, the halo mass function $n(M, z)$ is related to the probability $P(> \delta_c, M)$ that a region of mass M exceeds the critical density for collapse $\delta_c(z)$ through the relation $n(M, z) = -2(\bar{\rho}/M) dP/dM$. The non-Gaussian fractional correction to the multiplicity function then is $R(\nu, g_{\text{NL}}) \equiv f(\nu, g_{\text{NL}})/f(\nu, 0) = (dP/dM)(> \delta_c, M, g_{\text{NL}})/(dP/dM)(> \delta_c, M, 0)$. The level excursion probability $P(> \delta_c, M, g_{\text{NL}})$ can be computed once the probability distribution function (PDF) of the smoothed density field δ_M , $P(\delta_M)$, is known. Here, we will consider the simple extensions proposed by [56] and [57], in which $P(\delta_M)$ is generically expressed as the inverse transform of a cumulant generating function. Both extensions have been shown to give reasonable agreement with numerical simulations of non-Gaussian cosmologies [42, 44, 58].

In [56], the saddle-point technique is applied directly to $P(\delta_M)$. The resulting Edgeworth expansion is then used to obtain $P(> \delta_c, M, g_{\text{NL}})$. For f_{NL} non-Gaussianity, reference [42] found that the resulting non-Gaussian mass function agrees well with the simulations. For g_{NL} non-Gaussianity, neglecting cumulants other than the kurtosis $S_4(M) = \langle \delta_M^4 \rangle / \sigma^6$ and truncating the series expansion at S_4 , the non-Gaussian fractional correction reads

$$R_{\text{LV}}(\nu, g_{\text{NL}}) \approx \left\{ 1 + \frac{1}{4!} \sigma^2 S_4 (\nu^4 - 4\nu^2 - 3) \right. \quad (8)$$

$$\left. - \frac{1}{4!} \sigma^2 \frac{dS_4}{d \ln \nu} (\nu^2 - 3) \right\} \exp[\nu^2 \delta \sigma_8]$$

$$= \left\{ 1 + \frac{1}{4!} \sigma^2 S_4 (\nu^4 - 6\nu^2 + 3) \quad (9)$$

$$\left. - \frac{1}{4!} \frac{d(\sigma^2 S_4)}{d \ln \nu} (\nu^2 - 3) \right\} \exp[\nu^2 \delta \sigma_8]$$

after integration over regions above $\delta_c(z)$. Note that we have omitted writing the redshift dependence explicitly. Strictly speaking however, $R(\nu, g_{\text{NL}})$ depends distinctly upon the variables M (or ν) and z due to the presence of $\sigma^2 S_4(M)$. Our notation is motivated by the fact that the measured non-Gaussian correction, as plotted in Fig.3, appears to depend mostly on the peak height. The exponential factor in the right-hand side is the correction induced by the renormalisation of the amplitude of linear density fluctuations, eq. (6). For consistency, we have also used the Press-Schechter multiplicity function to derive this last term although a Sheth-Tormen mass function [59] may be more appropriate.

In [57], it is the level excursion probability $P(> \delta_c, M)$ that is calculated within the saddle-point approximation. Including only a cubic coupling $g_{\text{NL}} \phi^3$ and truncating the

resulting expression at the kurtosis, we find

$$P(> \delta_c, M, g_{\text{NL}}) \quad (10)$$

$$\approx \frac{1}{\sqrt{2\pi}} \frac{\sigma}{\delta_c} \left(1 + 3g_{\text{NL}} \sigma^2 \langle \phi^2 \rangle - \frac{S_4}{12} \sigma^2 \delta_c^2 \right)$$

$$\times \exp \left\{ -\frac{\delta_c^2}{2\sigma^2} \left[1 - 6g_{\text{NL}} \langle \phi^2 \rangle - \frac{S_4}{12} \delta_c^2 \right] \right\}$$

at first order in g_{NL} . Note that we have already included the renormalisation of the fluctuation amplitude. For rare events, $\sigma \ll 1$ and the first parenthesis in the right-hand side can be neglected. To ensure that the resulting mass function is properly normalised, we follow [60] and use

$$\nu_* f(\nu_*) = M^2 \frac{n(M, z, g_{\text{NL}})}{\bar{\rho}} \frac{d \ln M}{d \ln \nu_*}. \quad (11)$$

for the non-Gaussian mass function, where $\nu_* = \delta_*/\sigma$, $\delta_* = \delta_c \sqrt{1 - 2\delta\sigma_8 - S_4 \delta_c^2 / 12}$ ¹ and f is the same multiplicity function as in the Gaussian case. Taking the derivative of the level excursion probability then gives

$$\frac{(dP/dM)(> \delta_c, M, g_{\text{NL}})}{(dP/dM)(> \delta_c, M, 0)} \approx \exp \left[\frac{S_4 \delta_c^4}{4! \sigma^2} + \nu^2 \delta \sigma_8 \right] \quad (12)$$

$$\times \left(\frac{\delta_*}{\delta_c} + \frac{1}{4!} \frac{\delta_c^3}{\delta_*} \frac{dS_4}{d \ln \sigma} \right).$$

The fractional change in the multiplicity function eventually reads as

$$R_{\text{MVJ}}(\nu, g_{\text{NL}}) \approx \exp \left[\frac{\nu^4}{4!} \sigma^2 S_4 + \nu^2 \delta \sigma_8 \right] \quad (13)$$

$$\times \left\{ 1 - \frac{\nu^2}{8} \sigma^2 S_4 - \frac{\nu^2}{4!} \frac{d(\sigma^2 S_4)}{d \ln \nu} \right\}$$

after expanding δ_* at the first order and ignoring the shift in the normalisation amplitude, i.e. $\delta_*^{\pm 1} \approx 1 \mp S_4 \delta_c^2 / 24$. In the limit $\sigma^2 S_4 \ll 1$ and $\nu \gg 1$, the two theoretical expectations reduce to $1 + \nu^4 \sigma^2 S_4 / 24$. However, they differ in the coefficient of the $\nu^2 \sigma^2 S_4$ term, which is $-1/4$ and $-1/6$ for the LV and MVJ formula, respectively. There-

¹ In local f_{NL} models, this formula only involves the skewness S_3 . As it is incorrectly quoted in some of the literature on non-Gaussian halo mass functions, let us write down its explicit expression:

$$R_{\text{MVJ}}(\nu, f_{\text{NL}}) = \exp \left[\frac{S_3 \delta_c^3}{6\sigma^2} \right] \left[\frac{\delta_c^2}{6\delta_*} \frac{dS_3}{d \ln \sigma} + \frac{\delta_*}{\delta_c} \right],$$

or, in terms of the peak height $\nu = \delta_c/\sigma$,

$$R_{\text{MVJ}}(\nu, f_{\text{NL}}) \approx \exp \left[\frac{\nu^3}{6} \sigma S_3 \right] \left[1 - \frac{\nu}{3} \sigma S_3 - \frac{\nu}{6} \frac{d(\sigma S_3)}{d \ln \nu} \right],$$

after expanding $\delta_* = \sqrt{1 - S_3 \delta_c / 3}$ at the first order.

fore, we shall also consider the approximation

$$R(\nu, g_{\text{NL}}) = \exp \left[\frac{\nu^4}{4!} \sigma^2 S_4 + \nu^2 \delta\sigma_8 \right] \times \left\{ 1 - \frac{\nu^2}{4} \sigma^2 S_4 - \frac{\nu^2}{4!} \frac{d(\sigma^2 S_4)}{d \ln \nu} \right\}, \quad (14)$$

which is designed to match better the Edgeworth expansion of [56] when the peak height is $\nu \sim 1$.

Calculating the fractional change in the mass function requires knowledge of the kurtosis $S_4(M) \equiv g_{\text{NL}} S_4^{(1)}(M)$ of the smoothed density field δ_M , which we compute analytically using the general formula (valid for $n \geq 3$)

$$\sigma^{2n-2} S_n^{(1)}(M) = n! \left(\prod_{i=1}^{n-1} \int \frac{d^3 k_i}{(2\pi)^3} \alpha_M(k_i) P_\phi(k_i) \right) \times \alpha_M(|\mathbf{k}_1 + \dots + \mathbf{k}_{n-1}|), \quad (15)$$

where

$$\alpha_M(k) = \frac{2}{3\Omega_m H_0^2} D(0) k^2 T(k) W_M(k) \quad (16)$$

is evaluated at $z = 0$ and with a (spherically symmetric) window $W_M(k)$ function of characteristic mass scale M . A tophat filter is adopted throughout this paper. Over the mass range probed by our simulations, $10^{13} \lesssim M \lesssim 5 \times 10^{15} M_\odot/h$, the normalised kurtosis $\sigma^2 S_4^{(1)}(M)$ is a monotonic decreasing function of M that varies in the narrow range $\sim 4 - 6 \times 10^{-7}$ for the top-hat filter assumed here (see Fig. 2). Note also that the $\sigma^2 S_4$ term dominates the total contribution to the non-Gaussian correction eqs. (9), (13) and (14) when the peak height is $\nu \gtrsim 2$ (One finds $|d(\sigma^2 S_4)/d \ln \nu| \lesssim 0.1 \sigma^2 S_4$).

The fractional correction is plotted in Fig. 3 for the haloes extracted from the simulations at redshift $z = 0.3, 0.5, 1$ and 2 . In the top panel, the data are compared to the theoretical predictions eqs. (9), (13) and (14) evaluated at $z = 0$. As we can see, the level of non-Gaussianity in the halo multiplicity function is consistent with the theory. Our approximation (14) performs equally well regardless of the sign of g_{NL} . It agrees better with the measurements than the formulae of [57] which significantly overestimates the data for $g_{\text{NL}} = 10^6$, and that of [56] which is not always positive definite for $g_{\text{NL}} = -10^6$. The bottom panel shows that the discrepancy somewhat worsens at higher redshift, especially in the case $g_{\text{NL}} = 10^6$. However, it is possible the agreement may be improved by adding higher order powers of $\sigma^2 S_4$ and higher order cumulants.

To conclude this section, one should keep in mind that all these extensions are based on Press-Schechter and, therefore, provide a bad fit to the Gaussian mass function of haloes. In this respect, excursion set approaches may be more promising since they seem to reproduce both the Gaussian halo counts and the dependence on f_{NL} [61, 62].

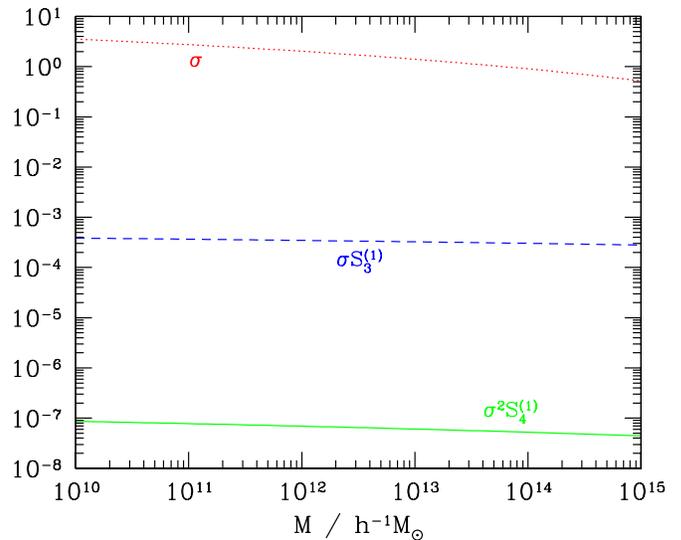


FIG. 2: Variance σ (dotted), skewness $\sigma S_3^{(1)}$ (dashed) and kurtosis $\sigma^2 S_4^{(1)}$ (solid) of the smoothed linear density field δ_M as a function of mass scale M .

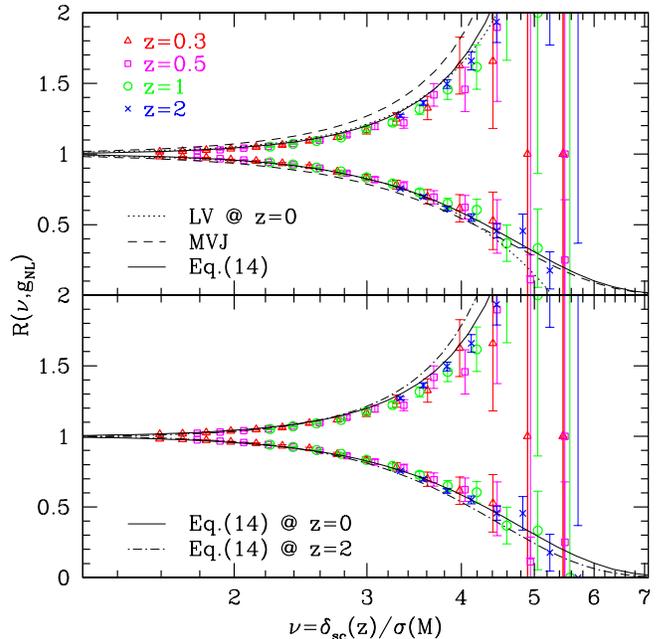


FIG. 3: *Top panel*: Fractional correction to the Gaussian multiplicity function of dark matter haloes as a function of the peak height $\nu(M, z)$ for a coupling parameter $g_{\text{NL}} = \pm 10^6$. The dotted, dashed and solid curves show the theoretical predictions eqs. (9), (13) and (14) at $z = 0$, respectively. Error bars denote Poisson errors. For illustration, $M = 10^{15} M_\odot/h$ corresponds to $\nu = 3.2, 5.2, 7.7$ at redshift $z = 0, 1$ and 2 , respectively. Similarly, $M = 10^{14} M_\odot/h$ and $10^{13} M_\odot/h$ correspond to $\nu = 1.9, 3, 4.5$ and $1.2, 1.9, 2.9$ respectively. *Bottom panel*: a comparison with equation (14) evaluated at $z = 0$ and 2 .

III. THE NON-GAUSSIAN BIAS SHIFT

A. Theoretical considerations

In order to calculate the scale-dependent bias correction induced by the g_{NL} coupling term to the correlation of haloes of mass M collapsing at redshift z , we follow [36] and consider the two-point correlation $\xi_{\text{hh}}(\mathbf{r})$ of regions of the smoothed density field δ_M above a threshold $\delta_c(z) = \nu(z)\sigma$. The two-point correlation function of this level excursion set, which was originally derived by [26], can be expressed in the high threshold approximation as

$$\xi_{\text{hh}}(\mathbf{r}) = -1 + \exp \left\{ \sum_{n=2}^{\infty} \sum_{j=1}^{n-1} \frac{\nu^n \sigma^{-n}}{j!(n-j)!} \right. \\ \left. \times \xi^{(n)} \left(\begin{array}{cc} \mathbf{x}_1, \dots, \mathbf{x}_1, & \mathbf{x}_2, \dots, \mathbf{x}_2 \\ j \text{ times} & (n-j) \text{ times} \end{array} \right) \right\}, \quad (17)$$

where $\mathbf{r} = \mathbf{x}_1 - \mathbf{x}_2$. For the non-Gaussian model considered here, the leading-order correction induced by non-zero three-point and four-point correlations of the density field reads

$$\Delta \xi_{\text{hh}} = \frac{\nu^3}{\sigma^3} \xi^{(3)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2) + \frac{\nu^4}{\sigma^4} \left[\frac{1}{3} \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2) \right. \\ \left. + \frac{1}{4} \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2, \mathbf{x}_2) \right]. \quad (18)$$

In the case $f_{\text{NL}} = 0$ and $g_{\text{NL}} \neq 0$, only the four-point functions contribute at first order. It should also be noted that, at linear order, $\xi^{(2)}(\mathbf{x}_1, \mathbf{x}_2)$ amounts to a renormalisation of the linear bias and, therefore, does not contribute to the scale-dependent correction. The Fourier mode of the curvature perturbation $\Phi(\mathbf{x})$ is given by

$$\Phi(\mathbf{k}) = \phi(\mathbf{k}) + g_{\text{NL}} \int \frac{d^3 k_1}{(2\pi)^3} \int \frac{d^3 k_2}{(2\pi)^3} \phi(\mathbf{k}_1) \phi(\mathbf{k}_2) \phi(\mathbf{k} - \mathbf{k}_1 - \mathbf{k}_2) \quad (19)$$

and its trispectrum is

$$T_{\Phi}(\mathbf{k}_1, \mathbf{k}_2, \mathbf{k}_3, \mathbf{k}_4) = 6g_{\text{NL}} [P_{\phi}(k_1)P_{\phi}(k_2)P_{\phi}(k_3) + (\text{cyclic})] \quad (20)$$

Therefore, the trispectrum of the density field linearly extrapolated to present epoch is simply

$$T_{\delta}(\mathbf{k}_1, \mathbf{k}_2, \mathbf{k}_3, \mathbf{k}_4) = \left(\prod_{i=1}^4 \alpha_M(k_i) \right) T_{\Phi}(\mathbf{k}_1, \mathbf{k}_2, \mathbf{k}_3, \mathbf{k}_4). \quad (21)$$

The four-point correlation $\xi^{(4)}$ of the density field then is the Fourier transform of $T_{\delta}(\mathbf{k}_1, \mathbf{k}_2, \mathbf{k}_3, \mathbf{k}_4)$ with conservation of momentum enforced. Hence,

$$\xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2) = 6g_{\text{NL}} \left(\prod_{i=1}^3 \int \frac{d^3 k_i}{(2\pi)^3} \alpha_M(k_i) P_{\phi}(k_i) \right) \\ \times \left[1 + 3 \frac{P_{\phi}(k_{123})}{P_{\phi}(k_3)} \right] \alpha_M(k_{123}) \\ \times e^{i\mathbf{k}_{123} \cdot \mathbf{r}}, \quad (22)$$

and

$$\xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2, \mathbf{x}_2) = 6g_{\text{NL}} \left(\prod_{i=1}^3 \int \frac{d^3 k_i}{(2\pi)^3} \alpha_M(k_i) P_{\phi}(k_i) \right) \\ \times \left[1 + 2 \frac{P_{\phi}(k_{123})}{P_{\phi}(k_1)} + \frac{P_{\phi}(k_{123})}{P_{\phi}(k_3)} \right] \\ \times \alpha_M(k_{123}) e^{i\mathbf{k}_{12} \cdot \mathbf{r}}, \quad (23)$$

where $\mathbf{k}_{ij\dots l} = \mathbf{k}_i + \mathbf{k}_j + \dots + \mathbf{k}_l$ for shorthand convenience. Since we will examine the effect of non-Gaussianity on Fourier space statistics only, we take the Fourier transform of the four-point functions. After some simplification, we arrive at

$$\int d^3 r \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2) e^{-i\mathbf{k} \cdot \mathbf{r}} = 6g_{\text{NL}} \alpha_M(k) P_{\phi}(k) \int \frac{d^3 k_1}{(2\pi)^3} \alpha_M(k_1) P_{\phi}(k_1) \int \frac{d^3 k_2}{(2\pi)^3} \alpha_M(k_2) P_{\phi}(k_2) \\ \times \alpha_M(|\mathbf{k} + \mathbf{k}_{12}|) \left[3 + \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_{12}|)}{P_{\phi}(k)} \right] \quad (24)$$

$$\int d^3 r \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2, \mathbf{x}_2) e^{-i\mathbf{k} \cdot \mathbf{r}} = 6g_{\text{NL}} \int \frac{d^3 k_1}{(2\pi)^3} \alpha_M(k_1) \alpha_M(|\mathbf{k} + \mathbf{k}_1|) P_{\phi}(k_1) P_{\phi}(|\mathbf{k} + \mathbf{k}_1|) \int \frac{d^3 k_2}{(2\pi)^3} \alpha_M(k_2) \\ \times \alpha_M(|\mathbf{k} + \mathbf{k}_2|) P_{\phi}(k_2) \left[1 + 2 \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_2|)}{P_{\phi}(|\mathbf{k} + \mathbf{k}_1|)} + \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_2|)}{P_{\phi}(k_2)} \right]. \quad (25)$$

For realistic values of the spectral index ($n_s \sim 1$), the products $\alpha_M(|\mathbf{k} + \mathbf{k}_i|) P_{\phi}(|\mathbf{k} + \mathbf{k}_i|)$ appearing in the right-hand side of the above equalities formally diverges whenever $\mathbf{k} + \mathbf{k}_1 = 0$ due to the ultraviolet divergence of $P_{\phi}(k)$. To cure this problem, one can set $P_{\phi}(k) = 0$ for sufficiently small wavenumbers or excise a thin shell centred at wavenumber k_i from the integral. In the large-scale limit $k \ll k_i$, the ratio $P_{\phi}(|\mathbf{k} + \mathbf{k}_i|)/P_{\phi}(k)$ vanishes whereas $P_{\phi}(|\mathbf{k} + \mathbf{k}_i|)/P_{\phi}(k_i)$

tends towards unity. In this case, the above expressions reduce to

$$\int d^3r \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2) e^{-i\mathbf{k}\cdot\mathbf{r}} \approx 3g_{\text{NL}} \sigma^4 S_3^{(1)}(M) \alpha_M(k) P_\phi(k) \quad (26)$$

$$+ 6g_{\text{NL}} \alpha_M(k) \int \frac{d^3k_1}{(2\pi)^3} \alpha_M(k_1) P_\phi(k_1) \int \frac{d^3k_2}{(2\pi)^3} \alpha_M(k_2) P_\phi(k_2) \alpha_M(k_{12}) P_\phi(k_{12}),$$

and

$$\int d^3r \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2, \mathbf{x}_2) e^{-i\mathbf{k}\cdot\mathbf{r}} \quad (27)$$

$$\approx 24g_{\text{NL}} \sigma^2 \int \frac{d^3k_1}{(2\pi)^3} \alpha_M^2(k_1) P_\phi^2(k_1).$$

Only the first term in the right-hand side of eq.(26) is not well behaved in the limit $k \rightarrow 0$ where it becomes proportional to k^{n_s-2} . The second scales as k^2 , while the Fourier transform of $\xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2, \mathbf{x}_2)$ asymptotes to a constant. A similar decomposition also arises in f_{NL} models. For this type of non-Gaussianity, the first order correction is furnished by the three-point function $\xi^{(3)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2)$, whose Fourier transform can be split into the familiar term $2f_{\text{NL}}\sigma^2\alpha_M(k)P_\phi(k)$, and a second piece given by

$$\frac{1}{2}f_{\text{NL}}\alpha_M(k) \int \frac{d^3k_1}{(2\pi)^3} \alpha_M^2(k_1) P_\phi^2(k_1). \quad (28)$$

In both quadratic and cubic local non-Gaussianity, the term proportional to $\alpha_M(k)$ can be neglected since, at the pivot point $k = k_0$, its magnitude relative to the term involving $\alpha_M(k)P_\phi(k)$ is only $\mathcal{O}(0.01)$ and $\mathcal{O}(10^{-6})$, respectively. Moreover, it decreases as one goes to larger scales. By contrast, it is not so obvious at first sight how to deal with the term (28). In the non-Gaussian halo power spectrum, this term would appear multiplied by $\nu^4/(4\sigma^4)$,

$$6g_{\text{NL}}b_{\text{L}}^2 \frac{\delta_c^2(z)}{\sigma^2} \int \frac{d^3k_1}{(2\pi)^3} \alpha_M^2(k_1) P_\phi^2(k_1) \quad (29)$$

$$\simeq 6 \times 10^{-4} g_{\text{NL}} b_{\text{L}}^2 \frac{D^2(0)}{D^2(z)} \left(\frac{M}{10^{13} M_\odot/h} \right)^{0.375}.$$

The approximation (second line) holds for $10^{13} \leq M \leq 10^{14} M_\odot/h$. For $g_{\text{NL}} = 10^6$ and $b_{\text{L}} \gtrsim 3$, this can be much larger than the typical shot-noise correction applied to the halo power spectra we measure in the simulations (see below). For $g_{\text{NL}} = -10^6$, this will certainly produce a halo power spectrum which is negative at sufficiently low wavenumber. It is plausible that higher order counter-terms in the expansion eq.(18) renormalises its value. However, such a calculation is beyond the scope of this paper, so the best option is to ignore this term in the following of the analysis.

The non-Gaussian correction ΔP_{hh} to the halo power spectrum P_{hh} thus is the Fourier transform of

$\nu^4 \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2, \mathbf{x}_2)/3\sigma^4$ which, in the limit of long wavelength $k \ll 1$, becomes

$$\Delta P_{\text{hh}}(k) = g_{\text{NL}} \nu^4 S_3^{(1)}(M) \alpha_M(k) P_\phi(k) \quad (30)$$

$$= g_{\text{NL}} b_{\text{L}}^2(z) \delta_c^2(z) S_3^{(1)}(M) \alpha_M(k) P_\phi(k),$$

where we have used $b_{\text{L}}(z) = \nu^2/\delta_c(z)$ as is appropriate for high density peaks. The smoothing window that appears in $\alpha_M(k)$ effectively makes little difference because we are considering the limit where k^{-1} is much larger than the smoothing radius, so we will omit it in the following. For small non-Gaussianity, we can also write $\Delta P_{\text{hh}} \approx 2b_{\text{L}}\Delta b_\kappa P_\delta(k)$ where $P_\delta(k) = \alpha_M^2 P_\phi(k)$ is the power spectrum of the smoothed density field. The scale-dependent bias correction $\Delta b_\kappa(k, g_{\text{NL}})$ can eventually be recast into the form

$$\Delta b_\kappa(k, g_{\text{NL}}) = \frac{1}{2} g_{\text{NL}} b_{\text{L}}(z) \delta_c^2(z) S_3^{(1)}(M) \alpha_M^{-1}(k) \quad (31)$$

$$= \frac{3}{4} g_{\text{NL}} b_{\text{L}}(z) \delta_c^2(0) \frac{D(0)}{D(z)^2} S_3^{(1)}(M) \frac{\Omega_{\text{m}} H_0^2}{k^2 T(k)}$$

$$= \frac{1}{4} g_{\text{NL}} \delta_c(z) S_3^{(1)}(M) \Delta b_\kappa(k, f_{\text{NL}} = 1),$$

where $\Delta b_\kappa(k, f_{\text{NL}})$ is the scale-independent bias induced by the quadratic coupling $f_{\text{NL}}\phi^2$, eq. (2). We have also assumed the Eulerian bias prescription $b(M) = 1 + b_{\text{L}}(M)$.

The change in the mean number density of haloes also creates a scale-independent shift which we denote by $\Delta b_{\text{I}}(g_{\text{NL}})$. As shown in [42] for f_{NL} models, the inclusion of this correction noticeably improves the agreement with the simulations at wavenumber $k \lesssim 0.1 h\text{Mpc}^{-1}$. Using a peak-background split and considering the limit of small non-Gaussianity, this contribution reads

$$\Delta b_{\text{I}}(g_{\text{NL}}) = -\frac{1}{\sigma} \frac{\partial}{\partial \nu} \ln [R(\nu, g_{\text{NL}})]$$

$$\approx -\frac{1}{\sigma} \left[\frac{1}{3!} (\nu^3 - 3\nu) \sigma^2 S_4 + 2\nu \delta \sigma_8 \quad (32)$$

$$+ \frac{1}{4!} (\nu^3 - 8\nu) \frac{d(\sigma^2 S_4)}{d \ln \nu} - \frac{\nu}{4!} \frac{d^2(\sigma^2 S_4)}{d \ln \nu^2} \right]$$

after truncating Δb_{I} at first order in g_{NL} . This approximation should perform reasonably well for moderate values of the peak height, $\nu \lesssim 4$, for which the fractional change in the mass function, equation (14), matches well the numerical data. It is worth noticing that $\Delta b_{\text{I}}(g_{\text{NL}})$ has a sign opposite to that of g_{NL} because the bias decreases

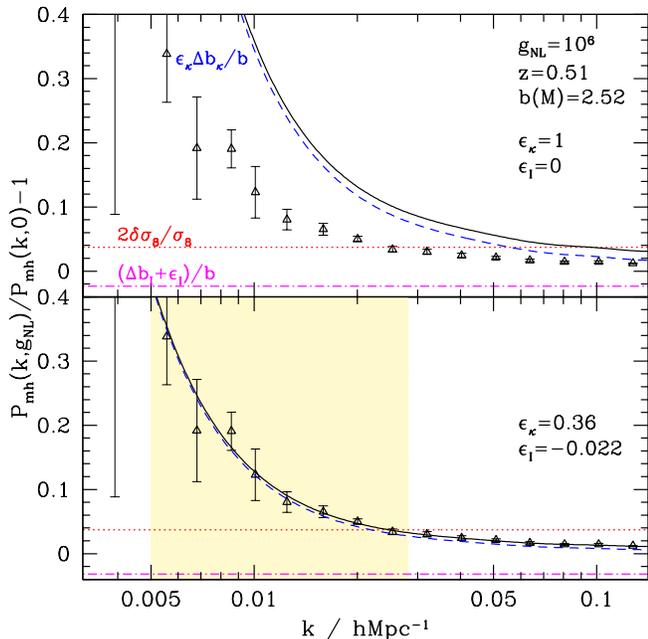


FIG. 4: *Top panel* : Non-Gaussian bias correction computed from the halo-matter power spectrum of haloes of mass $M > 2 \times 10^{13} M_{\odot}/h$ extracted from the snapshot at $z = 0.5$ (filled symbols). The solid curve represents $P_{\text{mh}}(k, g_{\text{NL}})/P_{\text{mh}}(k, 0) - 1$ with a non-Gaussian bias shift $\Delta b(k, g_{\text{NL}})$ given by eq. (34). The dashed, dotted and dotted-dashed curves show three separate contributions that arise at first order in g_{NL} . *Bottom panel* : $\Delta b(k, g_{\text{NL}})$ is replaced by the theoretical model eq. (35). The shaded region indicates the data points used to fit the parameters ϵ_{κ} and ϵ_{I} . Error bars indicate the scatter among 5 realisations.

when the mass function goes up. Δb_{I} also includes a correction induced by the renormalisation of σ_8 . In practice, we estimate $\Delta b_{\text{I}}(g_{\text{NL}})$ for a given halo sample by evaluating $\sigma^2 S_4$ and ν at the scale corresponding to the average halo mass \bar{M} of the sample, as it is unclear to which extent $R(\nu, g_{\text{NL}})$ agrees with the data in the limit $M \gg M_{\star}$.

B. Comparison with the simulations

To assess the effect of primordial non-Gaussianity on the halo bias, we will consider the ratios

$$\frac{P_{\text{mh}}(k, g_{\text{NL}})}{P_{\text{mh}}(k, 0)} - 1 = \frac{\Delta b(k, g_{\text{NL}})}{b(M)} + 2 \frac{\delta\sigma_8}{\sigma_8} \quad (33)$$

$$\frac{P_{\text{hh}}(k, g_{\text{NL}})}{P_{\text{hh}}(k, 0)} - 1 = \left(1 + \frac{\Delta b(k, g_{\text{NL}})}{b(M)}\right)^2 + 2 \frac{\delta\sigma_8}{\sigma_8} - 1,$$

where $\Delta b(k, g_{\text{NL}})$ is generally the sum of a scale-dependent and a scale-independent term. One should bear in mind that the scale-independent shift $2\delta\sigma_8/\sigma_8$ arises from the matter power spectrum and, therefore, is

distinct from the term $-2\nu\delta\sigma_8/\sigma$ appearing in Δb_{I} . Following [42], we shall also quantify the departure from the theoretical scaling as a function of wavemode amplitude with the ratio $\Delta b^s/\Delta b^t$, where Δb^s is the non-Gaussian bias correction measured from the simulation and Δb^t is eq. (35).

We interpolate the dark matter particles and halo centres onto a regular cubical mesh. The resulting dark matter and halo fluctuation fields are then Fourier transformed to yield the matter-matter, halo-matter and halo-halo power spectra $P_{\text{mm}}(k)$, $P_{\text{mh}}(k)$ and $P_{\text{hh}}(k)$, respectively. These power spectra are measured for a range of halo masses and redshifts, covering the relevant range of statistical properties corresponding to the available data sets of galaxy or quasar populations with different luminosities and bias. Note that these quantities are computed on a 512^3 grid, whose Nyquist wavenumber is sufficiently large ($\approx 1 h\text{Mpc}^{-1}$) to allow for an accurate measurement of the power in wavemodes of amplitude $k \lesssim 0.1 h\text{Mpc}^{-1}$. The halo power spectrum is corrected for the shot-noise due to the discrete nature of dark matter haloes, which we assume to be the standard Poisson term $1/\bar{n}_{\text{h}}$. This discreteness correction is negligible for $P_{\text{mm}}(k)$ due to the large number of dark matter particles. Yet another important quantity is the linear halo bias $b(M)$ which must be measured accurately from the Gaussian simulations as it controls the magnitude of the scale-dependent shift. Here, we shall use the ratio $P_{\text{mh}}(k)/P_{\text{mm}}(k)$ as a proxy for the halo bias since it is less sensitive to shot-noise.

1. An effective non-Gaussian bias correction

Summarizing the analytical considerations of Sec. III A, non-Gaussianity of the g_{NL} type add a correction $\Delta b(k, g_{\text{NL}})$ to the bias $b(k)$ of dark matter haloes which is at leading order

$$\Delta b(k, g_{\text{NL}}) = \Delta b_{\kappa}(k, g_{\text{NL}}) + \Delta b_{\text{I}}(g_{\text{NL}}), \quad (34)$$

We found however that this theoretical expectation significantly overestimates the magnitude of the non-Gaussian bias shift measured from the simulations. This is exemplified in the top panel of Fig. 4, where $P_{\text{mh}}(k, g_{\text{NL}})/P_{\text{mh}}(k, 0) - 1$ is plotted for haloes of mass $M > 2 \times 10^{13} M_{\odot}/h$ identified at $z = 0.5$. Clearly, the predicted scale-dependent correction Δb_{κ} is much steeper than measured from the halo samples. In order to improve the agreement with the numerical data, we modify the above relation as follows :

$$\Delta b(k, g_{\text{NL}}) = \epsilon_{\kappa} \Delta b_{\kappa}(k, g_{\text{NL}}) + [\Delta b_{\text{I}}(g_{\text{NL}}) + \epsilon_{\text{I}}], \quad (35)$$

and treat ϵ_{κ} and ϵ_{I} as free parameters that we fit to our measurements of the cross-power spectrum in the range $0.005 \leq k \leq 0.03 h\text{Mpc}^{-1}$ (weighted by the scatter among 5 realisations) where the scale-dependent effect is largest. The bottom panel of Fig. 4 shows the resulting

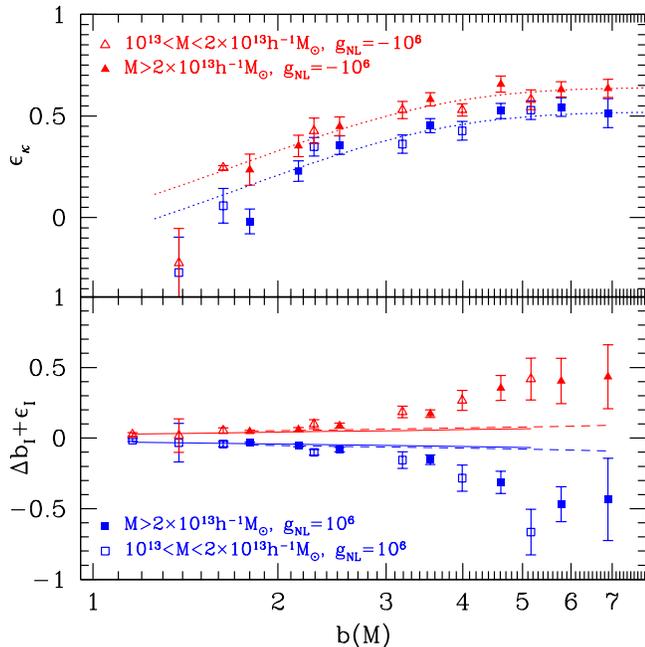


FIG. 5: Best-fitting ϵ_κ and $\Delta b_I + \epsilon_I$ as a function of halo bias and $g_{\text{NL}} = \pm 10^6$, for two different mass cuts as indicated in the Figure. In the top panel, the dotted curve is our best fit to $\epsilon_\kappa(b, g_{\text{NL}})$. In the bottom panel, the solid and dashed lines show the scale-independent shift Δb_I predicted by a peak-background split, eq. (32), for the low and high mass samples, respectively.

best-fit contributions $\epsilon_\kappa \Delta b_\kappa$ and $\Delta b_I + \epsilon_I$ for the halo sample mentioned above. Overall, as seen in Fig. 5, ϵ_κ and ϵ_I appear to depend mainly upon the linear halo bias $b(M)$ and the coupling parameter g_{NL} , although dependencies on redshift or other halo observables are not excluded (The data is too noisy for a reliable estimate of these). The most striking feature of Fig. 5 is the functional dependence of ϵ_κ on $b(M)$ and g_{NL} . Firstly, ϵ_κ is a monotonically increasing function of the bias and never reaches unity, even for the most biased samples for which the high peak approximation should be valid. Secondly, ϵ_κ is noticeably larger for $g_{\text{NL}} = -10^6$, suggesting thereby that second (and higher) order contributions to the scale-dependent bias may be important. Furthermore, for $b \lesssim 1.5$ where the high peak approximation breaks down, there is evidence that the effect reverses sign. The bottom panel of Fig. 5 shows that the scale-independent correction has sign opposite to that of g_{NL} , in agreement with theoretical expectations from the peak-background split (see Sec.III A). However, whereas for $b \lesssim 3$ the magnitude of the correction is comparable to that predicted by eq. (32), it is considerably larger for $b \gtrsim 3$, reaching up to 5-10 per cent of the linear halo bias.

Assuming that ϵ_κ is a function of $b(M)$ and g_{NL} only and asymptotes to a constant in the highly biased limit,

we find that the following parametrised form

$$\epsilon_\kappa(b, g_{\text{NL}}) = c_1 - c_2 g_{\text{NL}} - \frac{c_3}{1 + \exp(c_4 b)} \quad (36)$$

captures reasonably well the increase of ϵ_κ with halo bias for $1.5 < b(M) < 7$. The best-fit values of the parameters are

$$\begin{aligned} c_1 &= 0.59 \pm 0.03, & c_2 &= (6.0 \pm 0.9) \times 10^{-8} \\ c_3 &= 2.1 \pm 0.5, & c_4 &= 0.88 \pm 0.13. \end{aligned} \quad (37)$$

We do not provide a fitting formula for ϵ_I (or $\Delta b_I + \epsilon_I$) since it is not directly measurable in real data.

Explaining these findings clearly requires a better theoretical understanding, which we leave for a future investigation. Notwithstanding this, the phenomenological prescription (35) with parameters ϵ_κ and ϵ_I fitted to the data provides, as we will see below, a good description of the large-scale halo power spectrum in simulations of g_{NL} models.

2. Non-Gaussian bias from auto- and cross-power spectra

We have measured auto- and cross-power spectra for a range of halo masses and redshifts spanning the range $0 < z < 2$. The ratios defined in eq.(33) are shown in Figs 6 and 7 as a function of wavenumber for the mass threshold $M > 2 \times 10^{13} M_\odot/h$ and the mass bin $10^{13} < M < 2 \times 10^{13} M_\odot/h$, respectively. The fractional deviation $\Delta b^s / \Delta b^t$ is also shown at the bottom of each panel. The shaded region indicates a departure less than 20 per cent. Error bars denote the scatter around the mean and, therefore, may underestimate the true errors as they are computed from a small number of realisations. Note that, in order to reduce the impact of sampling variance, we first compute the ratios $P_{\text{mh}}(k, g_{\text{NL}}) / P_{\text{mh}}(k, 0)$ and $P_{\text{hh}}(k, g_{\text{NL}}) / P_{\text{hh}}(k, 0)$ for each realisation before calculating the average.

As we can see, once ϵ_κ and ϵ_I are fitted to the ratio of cross-power spectra, the theoretical prediction eq. (35) provides a reasonable description of the non-Gaussian bias in the halo power spectrum P_{hh} , indicating that non-Gaussianity does not generate much stochasticity and the predicted scaling $\Delta b_\kappa(k, g_{\text{NL}}) \propto k^{-2} T(k)^{-1}$ applies equally well for the auto- and cross-power spectrum. This was also found to be true in f_{NL} models [42]. The inclusion of a scale-independent correction $\Delta b_I + \epsilon_I$ significantly improve the agreement at $k \lesssim 0.03 h\text{Mpc}^{-1}$. For the highly biased samples $b > 4$ however, this correction is so large that the non-Gaussian bias shift reverses sign at wavenumber $k \gtrsim 0.05 h\text{Mpc}^{-1}$. Such an effect is not seen in the simulations, but we expect large deviations from the relation (35) in that range of wavenumber, where second- and higher-order corrections induced by the cubic coupling $g_{\text{NL}} \phi^3$ together with the nonlinear bias created by the gravitational evolution of matter density fluctuations may become important. Even though the data is noisier due to the low

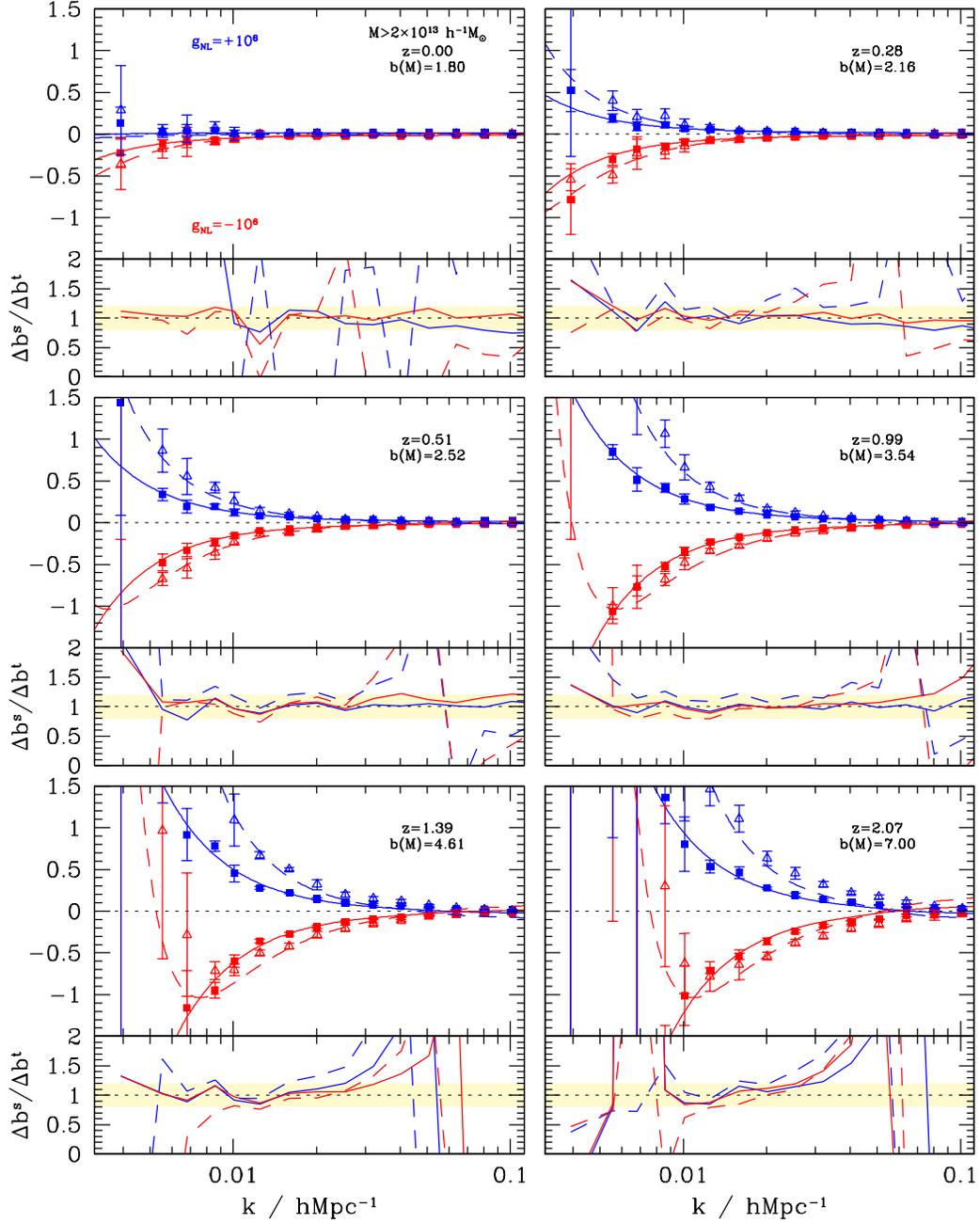


FIG. 6: Non-Gaussian bias correction measured in the simulation outputs at redshift $0 < z < 2$ for haloes of mass $M > 2 \times 10^{13} M_{\odot}/h$. In each panel, the upper plot shows the ratio $P_{\text{bh}}(k, g_{\text{NL}})/P_{\text{bh}}(k, 0) - 1$ (dashed curves, empty symbols) and $P_{\text{mh}}(k, g_{\text{NL}})/P_{\text{mh}}(k, 0) - 1$ (solid curves, filled symbols). The error bars represent the scatter among 5 realisations. The respective output redshift and linear halo bias are also quoted. The bottom of each panel displays the departure from the theoretical prediction, $\Delta b^s/\Delta b^t$. The shaded area indicates the domain where the deviation is less than 20 per cent. The parameters ϵ_{κ} and ϵ_1 are fitted individually to each sample. For illustration, ϵ_{κ} takes the best-fit values 0.06 and 0.60 for the lowest and highest biased samples, respectively.

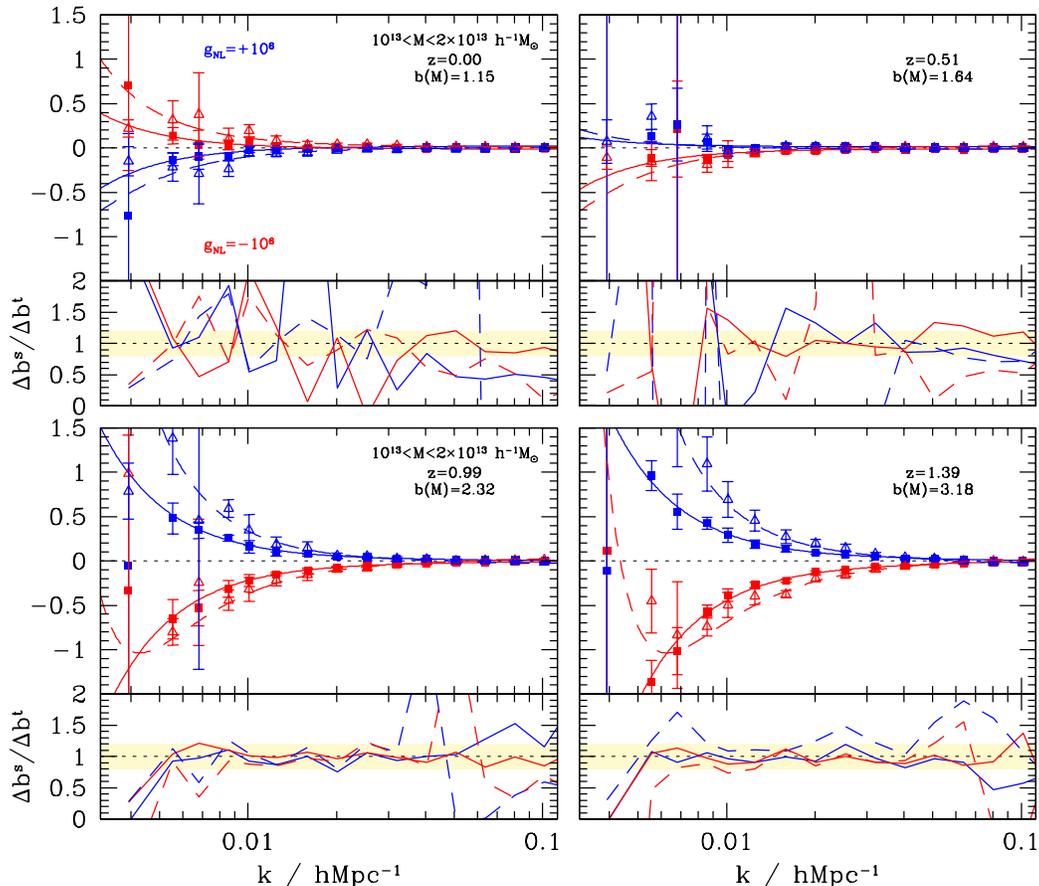


FIG. 7: Same as Fig. 6 but for haloes in the mass range $10^{13} < M < 2 \times 10^{13} M_{\odot}/h$.

number density of haloes, it is worth noting that, for the highly biased samples at $k \lesssim 0.01 \text{ hMpc}^{-1}$, the cross-power spectrum $P_{\text{mh}}(k, g_{\text{NL}} = -10^6)$ goes negative while $P_{\text{hh}}(k, g_{\text{NL}} = -10^6)$ remains positive and increases sharply, in agreement with the analytic prediction. Still, there is some evidence that the ratio $P_{\text{hh}}(k, g_{\text{NL}} = -10^6)/P_{\text{hh}}(k, 0) - 1$ saturates at a value noticeably larger than -1 before the sharp upturn, whereas our model predicts $2\delta\sigma_8/\sigma_8 - 1 \approx -0.96$ at the minimum.

Fig. 7 further explore the effect in the low mass samples, for which the $z = 0$ haloes with $b(M) \approx 1.15$ constitute an almost unbiased sample of the density field. In this case, the sign of the scale-dependent contribution is reversed, namely, the large-scale halo power spectrum in simulations of $g_{\text{NL}} = -10^6$ is enhanced relative to that of the Gaussian ones. This is in rough agreement with the theory, which predicts a similar effect for $b(M) < 1$. Again, haloes with a similar bias also have a comparable scale-dependent and scale-independent bias regardless of mass or redshift. Finally, note that the sample at $z = 1.39$ shown in Fig. 7 corresponds closely to the quasar sample used by [37], for which $z = 1.8$ and $b = 2.7$.

IV. CONSTRAINTS ON THE COUPLING PARAMETER g_{NL}

A. Constraints on g_{NL} from current large-scale structure data

Reference [37] took advantage of the scale-dependence of the bias to constrain f_{NL} from a sample of highly biased luminous red galaxies (LRGs) and quasars (QSOs). It is straightforward to translate their $2\text{-}\sigma$ limit $-29 < f_{\text{NL}} < +69$ into a constraint on g_{NL} since the non-Gaussian scale-dependent bias $\Delta b_{\kappa}(k, g_{\text{NL}})$ has the same functional form as $\Delta b_{\kappa}(k, f_{\text{NL}})$.

Constraints will arise mostly from the QSO sample at median redshift $z = 1.8$, which covers a large comoving volume and is highly biased, $b = 2.7$. In light of our results (see Fig. 5), we expect the parameter $\epsilon_{\kappa}(b, g_{\text{NL}})$ to vary with g_{NL} . However, in order to simplify the analysis, we will assume that, at fixed b , $\epsilon_{\kappa}(b, g_{\text{NL}})$ is given by the mean of $\epsilon_{\kappa}(b, g_{\text{NL}} = \pm 10^6)$. For a sample with bias $b \sim 2.7$, this implies $\epsilon_{\kappa} \simeq 0.4$. Furthermore, assuming that the typical mass of QSO-hosting haloes

is $\sim 10^{13} M_\odot/h$ yields $S_3^{(1)}(M) \simeq 2.3 \times 10^{-4}$. Hence, the multiplicative factor $(1/4) \delta_c(z) \epsilon_\kappa S_3^{(1)}(M)$ is approximately $\simeq 8.4 \times 10^{-5}$. Our limits on g_{NL} thus are

$$-3.5 \times 10^5 < g_{\text{NL}} < +8.2 \times 10^5 \quad (38)$$

at 95% confidence level. The scale-independent correction $\Delta b_{\text{I}} + \epsilon_{\text{I}}$ is not directly measured as it adds up to the bias b which is fitted to the data. For the limits obtained here, $|\Delta b_{\text{I}} + \epsilon_{\text{I}}|$ should be much smaller than b and can thus be ignored. Note also that, whereas the non-Gaussian bias scales as $D(z)^{-1}$ in f_{NL} models, we have $\Delta b(k, g_{\text{NL}}) \propto D(z)^{-2}$ for g_{NL} non-Gaussianity, so one can achieve relatively larger gains from measurements of high redshift tracers. In fact, the extent to which one can improve the observational bounds will strongly depend on our ability to minimize the impact of sampling variance caused by the random nature of the wavemodes, and the shot-noise caused by the discrete nature of the tracers. By comparing differently biased tracers of the same surveyed volume [40] and suitably weighting galaxies (e.g. by the mass of their host halo) [63, 64], it should be possible to circumvent these problems and considerably improve the detection level.

B. Predictions from future LSS surveys

References [39, 40, 47, 65] applied the Fisher matrix formalism to forecast constraints on f_{NL} from forthcoming galaxy redshift surveys. Here, we will simply try to estimate the detection limit for g_{NL} . Following [39, 47], we consider a (nearly spherical) survey of volume V . Assuming the Fourier modes are still uncorrelated and Gaussian distributed, the total signal-to-noise squared reads

$$\left(\frac{S}{N}\right)^2 \approx \frac{V}{4\pi^2} \int_{k_{\text{min}}}^{k_{\text{max}}} dk k^2 \left[\left(1 + \frac{\Delta b_\kappa}{b}\right)^2 - 1 \right]^2 \quad (39)$$

in the limit where sampling variance dominates the errors. Here, $k_{\text{min}} \sim \pi/V^{1/3}$ is the smallest wavemode accessible and k_{max} is not necessarily finite since the integral does converge as one takes k_{max} to infinity. Substituting the expression eq. (31) for the scale-dependent bias $\Delta b_\kappa(k, g_{\text{NL}})$ and setting $T(k) \equiv 1$ over the wavenumber range across which the integral is performed, we arrive at

$$\left(\frac{S}{N}\right)^2 \approx \frac{V}{\pi^2} (k_\star^2)^2 \left(\frac{1}{k_{\text{min}}} - \frac{1}{k_{\text{max}}} \right), \quad (40)$$

where

$$k_\star^2 \simeq 5.0 \times 10^{-12} g_{\text{NL}} \epsilon_\kappa \frac{(1-1/b)}{D^2(z)} \left(\frac{S_3^{(1)}}{10^{-4}} \right) h^2 \text{Mpc}^{-2}. \quad (41)$$

We have also assumed $|g_{\text{NL}}| \lesssim 10^5$, such that $|k_\star^2|$ is at most of the order of k_{min}^2 and the term linear in $\Delta b_\kappa/b$

dominates the signal. When $k_{\text{min}} \ll k_{\text{max}}$, we can further simplify $(S/N)^2$ to

$$\left(\frac{S}{N}\right)^2 \approx 8.1 \times 10^{-13} g_{\text{NL}}^2 \epsilon_\kappa^2 \left(1 - \frac{1}{b}\right)^2 D(z)^{-4} \times \left(\frac{S_3^{(1)}}{10^{-4}}\right)^2 \left(\frac{V}{h^{-3} \text{Gpc}^3}\right)^{4/3}. \quad (42)$$

Note the strong sensitivity of the signal-to-noise squared to the growth factor $D(z)$ (For f_{NL} non-Gaussianity, this dependence is only $D(z)^{-2}$).

To highlight the improvement one could achieve with future galaxy surveys, it is useful to first calculate the detection limit for the SDSS LRG sample centred at $z \sim 0.3$ and covering a volume $v \approx 2 h^{-3} \text{Gpc}^3$. Assuming a linear bias $b = 2$ and a skewness parameter $S_3^{(1)} \sim 2 \times 10^{-4}$ appropriate for haloes of mass $M \sim 10^{12} - 10^{13} M_\odot/h$, the minimum g_{NL} detectable at the 1- σ level is $\simeq 10^6$ for a correction factor $\epsilon_\kappa = 0.3$ which we read off from Fig. 5. For a survey configuration analogous to SDSS-III/BOSS², with central redshift $z = 0.5$ and a comoving volume $V = 6 h^{-3} \text{Gpc}^3$, the minimum g_{NL} would be $\sim 4 \times 10^5$ for galaxies tracing haloes of similar mass and bias. Finally, for a configuration like EUCLID³ with a $V = 100 h^{-3} \text{Gpc}^3$ survey centred at $z = 1.4$, the detection limit would be $\sim 2.1 \times 10^4$. Clearly, these limits are only indicative: they may be significantly improved by selecting highly biased, high redshift (single- or multi-)tracers. Nevertheless, this shows that future galaxy surveys should furnish interesting constraints on the size of the cubic coupling $g_{\text{NL}} \phi^3$.

C. Predictions for CMB temperature anisotropies

The CMB trispectrum provides an alternative probe of local, non-quadratic correction to the Gaussian curvature perturbations, so it is interesting to assess the sensitivity of this statistics to the nonlinear parameter g_{NL} .

Before calculating the detection limit, let us briefly review the formalism established for the CMB trispectrum [18, 19]. The temperature anisotropy field is conveniently decomposed into spherical harmonics, $\Delta T(\hat{\mathbf{n}})/T = \sum_{lm} a_l^m Y_l^m(\hat{\mathbf{n}})$. Statistical isotropy and invariance under parity transformation $\hat{\mathbf{n}} \rightarrow -\hat{\mathbf{n}}$ implies that the 4-point correlation of the spherical harmonic coefficients a_l^m takes the form

$$\langle a_{l_1}^{m_1} a_{l_2}^{m_2} a_{l_3}^{m_3} a_{l_4}^{m_4} \rangle = \sum_{LM} (-1)^M \begin{pmatrix} l_1 & l_2 & L \\ m_1 & m_2 & -M \end{pmatrix} \begin{pmatrix} l_3 & l_4 & L \\ m_3 & m_4 & M \end{pmatrix} Q_{l_3 l_4}^{l_1 l_2}(L). \quad (43)$$

² www.sdss3.org

³ http://sci.esa.int/science-e/www/object/index.cfm?fobjectid=42266

Here, $Q_{l_3 l_4}^{l_1 l_2}(L)$ is the angular average trispectrum and brackets are Wigner-3j symbols. Statistical homogeneity also implies that $Q_{l_3 l_4}^{l_1 l_2}(L)$ is independent of position. The connected part of the trispectrum, $T_{l_3 l_4}^{l_1 l_2}(L)$, encodes information about non-Gaussianity and is obtained by subtracting a Gaussian piece constructed from the power spectra C_l . The invariance of the 4-point harmonic function under the $4!$ permutations of the coefficients $a_{l_i}^{m_i}$ imposes constraints on $T_{l_3 l_4}^{l_1 l_2}(L)$ which we summarize in Appendix A. Eq.(43) can be inverted with the aid of the orthogonality of the Wigner-3j symbols to form an estimator for the CMB trispectrum.

The signal-to-noise for the CMB trispectrum $T_{l_3 l_4}^{l_1 l_2}(L)$ summed up to a certain l_{\max} is [19]

$$\left(\frac{S}{N}\right)^2(< l_{\max}) \approx \sum_{l_1 > l_2 > l_3 > l_4}^{l_{\max}} \sum_L \frac{|T_{l_3 l_4}^{l_1 l_2}(L)|^2}{(2L+1)C_{l_1}C_{l_2}C_{l_3}C_{l_4}} \quad (44)$$

when cosmic variance dominates the errors. Otherwise, one shall include a contribution from the power spectrum of the detector noise to the C_l . Galactic foreground subtraction on a fraction $1 - f_{\text{sky}}$ would further reduce $(S/N)^2$ by a factor of f_{sky} .

Neglecting the ISW effect, the Sachs-Wolfe provides a useful order-of-magnitude estimate of the signal-to-noise as long as l_{\max} does not exceed $\lesssim 100$ [19, 20, 66]. Details of the calculation are given in Appendix A. In short, the CMB trispectrum simplifies to

$$T_{l_3 l_4}^{l_1 l_2}(L) = \frac{27}{2\pi} g_{\text{NL}} (2L+1) \pi_{l_1 l_2 l_3 l_4}^2 \quad (45)$$

$$\times \begin{pmatrix} l_1 & l_2 & L \\ 0 & 0 & 0 \end{pmatrix} \begin{pmatrix} l_3 & l_4 & L \\ 0 & 0 & 0 \end{pmatrix}$$

$$\times [C_{l_1}^{\text{SW}} C_{l_2}^{\text{SW}} C_{l_3}^{\text{SW}} + (\text{cyclic})].$$

where

$$C_l^{\text{SW}} = \frac{2}{9\pi} \int_0^\infty dk k^2 P_\phi(k) j_l^2(kr_*) \approx \frac{2\pi A_\phi}{9l(l+1)}. \quad (46)$$

The last equality assumes a nearly scale-invariant spectrum $n_s \approx 1$. We also use the notation

$$\pi_{l_1 l_2 l_3 l_4} = \sqrt{(2l_1+1) \times \cdots \times (2l_4+1)}. \quad (47)$$

The signal-to-noise can then be recast into the compact form

$$\left(\frac{S}{N}\right)^2(< l_{\max}) = \frac{9}{2} A_\phi^2 \left\{ \frac{1}{6} \int_{-1}^{+1} dx s_{l_{\max}}^3(x) t_{l_{\max}}(x) + \frac{1}{2} \int_{-1}^{+1} dx r_{l_{\max}}^2(x) s_{l_{\max}}^2(x) \right\} \quad (48)$$

where the auxiliary functions $r_l(x)$, $s_l(x)$ and $t_l(x)$ are

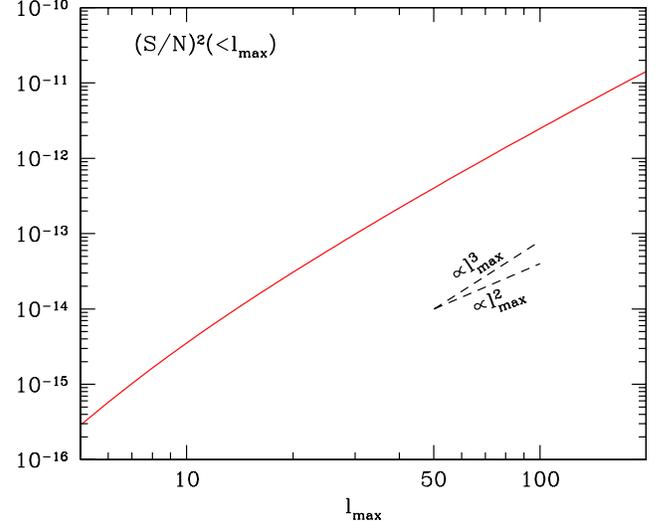


FIG. 8: Signal-to-noise ratio squared for the CMB trispectrum as a function of the maximum multipole l_{\max} . We have assumed $g_{\text{NL}} = 1$.

defined as

$$r_l(x) = \sum_{k=2}^l (2k+1) P_k(x) \quad (49)$$

$$s_l(x) = \sum_{k=2}^l \frac{(2k+1)}{k(k+1)} P_k(x) \quad (50)$$

$$t_l(x) = \sum_{k=2}^l (2k+1)k(k+1) P_k(x). \quad (51)$$

Here, $P_l(x)$ are Legendre polynomials. Note that we have excluded the monopole and dipole from the summation since these modes are unobservable.

Figure. 8 shows the signal-to-noise ratio squared in the Sachs-Wolfe approximation for the CMB trispectrum assuming $g_{\text{NL}} = 1$. Although this approximation breaks down for $l_{\max} \gtrsim 100$, we have extended the calculation up to $l_{\max} = 200$ so as to extrapolate more robustly the l_{\max} -dependence to small angular resolution. A power-law fit to $(S/N)^2$ in the range $50 \leq l_{\max} \leq 200$ gives

$$\left(\frac{S}{N}\right)^2(< l_{\max}) \simeq 2.43 \times 10^{-17} g_{\text{NL}}^2 \left(\frac{A_\phi}{10^{-9}}\right)^2 l_{\max}^{2.6}. \quad (52)$$

Our results appear consistent with the findings of [19] shown in their Fig.2. However, our constant of proportionality is about 20-30 times larger, presumably because they adopted a lower fluctuation amplitude (compare also their prediction for the f_{NL} model with that of [20]). Note that adding the information encoded in temperature-polarization trispectra may enhance $(S/N)^2$ by a factor of a few [18].

Assuming that the scaling eq.(52) persists well beyond the range over which the Sachs-Wolfe effect dominates,

the minimum g_{NL} detectable at $1\text{-}\sigma$ level is $g_{\text{NL}} \simeq 20, 7.9, 3.2, 1.9$ and 1.3×10^4 for $l_{\text{max}} = 250, 500, 1000, 1500$ and 2000 . A more realistic calculation should include the full radiation transfer function, detector noise etc. In this respect, detailed calculations have shown that, for the quadratic coupling $f_{\text{NL}}\phi^2$, $(S/N)^2$ of the CMB bispectrum and trispectrum closely follows the behaviour obtained in the Sachs-Wolfe approximation [19, 20]. It seems reasonable, then, to expect that this is also true for $g_{\text{NL}}\phi^3$.

While our predictions are qualitative they show that, for the WMAP CMB temperature measurement⁴ ($l_{\text{max}} \sim 250$), no detection of a significant trispectrum implies $|g_{\text{NL}}| \leq 2 \times 10^5$ at the $1\text{-}\sigma$ level. This is of the same order as the limit we derived from the QSO sample analyzed by [37]. For a PLANCK-like experiment⁵ ($l_{\text{max}} \sim 1500$), no evidence for a trispectrum would imply $|g_{\text{NL}}| \leq 1.3 \times 10^4$ at the $1\text{-}\sigma$ level. This is comparable to the detection limit that could be achieved with an all-sky survey such as EUCLID.

V. EFFECT OF NON-GAUSSIANITY WITH NON-ZERO f_{NL} AND g_{NL}

In this Section, we examine the halo multiplicity function and large-scale bias in numerical simulations of structure formation with non-zero coupling parameters $(f_{\text{NL}}, g_{\text{NL}}) = (\pm 100, -3 \times 10^5)$. We show that the results are consistent with those obtained from the simulations with non-vanishing g_{NL} solely.

A. Mass function

It is straightforward to calculate the fractional deviation from the Gaussian mass function, eq.(14), to non-zero f_{NL} and g_{NL} . Again, we start with the MVJ formula and neglect second order corrections such as $(\sigma S_3)^2$ etc. Adjusting the coefficient of the terms $\nu\sigma S_3$ and $\nu^2\sigma^2 S_4$ to that of the small ν expansion obtained by [56], we arrive at

$$R(\nu, \text{NL}) = \exp \left[\frac{\nu^3}{3!} \sigma S_3 + \frac{\nu^4}{4!} \sigma^2 S_4 + \nu^2 \delta\sigma_8 \right] \times \left\{ 1 - \frac{\nu}{2} \sigma S_3 - \frac{\nu}{6} \frac{d(\sigma S_3)}{d \ln \nu} - \frac{\nu^2}{4} \sigma^2 S_4 - \frac{\nu^2}{4!} \frac{d(\sigma^2 S_4)}{d \ln \nu} \right\}, \quad (53)$$

where the shorthand notation NL designates the combination $(f_{\text{NL}}, g_{\text{NL}})$. In Fig. 9, this theoretical prediction

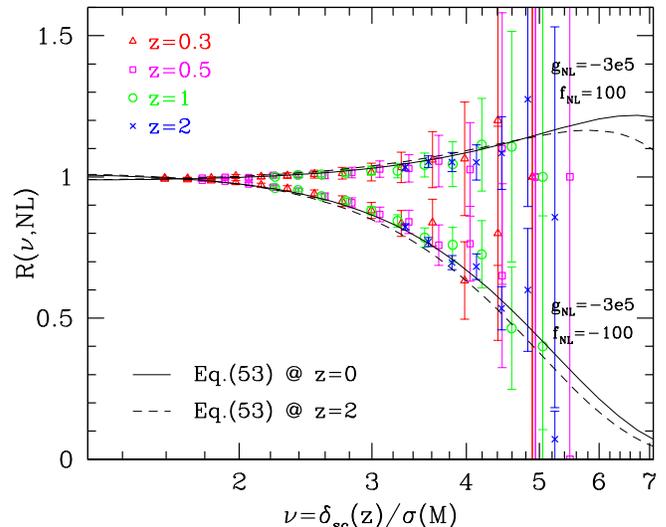


FIG. 9: Fractional correction to the Gaussian multiplicity function of dark matter haloes as a function of the peak height $\nu(M, z)$ for $f_{\text{NL}} = \pm 100$ and a cubic coupling parameter $g_{\text{NL}} = -3 \times 10^5$. The solid and dashed curves show the theoretical prediction eq.(53) at $z = 0$ and 2 , respectively. Error bars denote Poisson errors.

is compared $R(\nu, \text{NL})$ measured in non-Gaussian simulations of $(f_{\text{NL}}, g_{\text{NL}}) = (\pm 100, -3 \times 10^5)$. We account for the fact that the amplitude of density fluctuations is renormalised by $\delta\sigma_8 \approx 0.0045$. Fig. 9 demonstrates that our approximation is in good agreement with the data, although it slightly overestimates the effect at $z = 2$ when $f_{\text{NL}} = -100$. For $f_{\text{NL}} = 100$, the positive and negative contributions from the quadratic and cubic coupling, respectively, almost cancel each other and flatten the fractional deviation over most of the mass range probed by the simulations.

The scale-independent bias shift which arises from the change in the mean number density of haloes can again be estimated using the peak-background split. We find

$$\Delta b_1(\text{NL}) \approx -\frac{1}{\sigma} \left[\frac{1}{3!} (\nu^3 - 3\nu) \sigma^2 S_4 + \frac{1}{2} (\nu^2 - 1) \sigma S_3 + \frac{1}{4!} (\nu^3 - 8\nu) \frac{d(\sigma^2 S_4)}{d \ln \nu} - \frac{\nu}{4!} \frac{d^2(\sigma^2 S_4)}{d \ln \nu^2} + \frac{1}{6} (\nu^2 - 4) \frac{d(\sigma S_3)}{d \ln \nu} - \frac{1}{6} \frac{d^2(\sigma S_3)}{d \ln \nu^2} + 2\nu \delta\sigma_8 \right] \quad (54)$$

at the first order in the nonlinear parameters f_{NL} and g_{NL} .

B. Bias

Having checked that the amount of non-Gaussianity in the mass function is also consistent with our simple

⁴ <http://map.gsfc.nasa.gov/>

⁵ <http://sci.esa.int/science-e/www/area/index.cfm?fareaid=17>

theoretical expectation when both f_{NL} and g_{NL} are non-zero, we now turn to the clustering of dark matter haloes. The quadratic coupling $f_{\text{NL}}\phi^2$ generates the bispectrum at leading order, and a second order contribution to the trispectrum which reads [67]

$$T_{\Phi}(\mathbf{k}_1, \mathbf{k}_2, \mathbf{k}_3, \mathbf{k}_4) = 4f_{\text{NL}}^2 [P_{\phi}(k_{13})P_{\phi}(k_3)P_{\phi}(k_4) + 11 \text{ permutations}] . \quad (55)$$

Following the steps outlined in Sec.III A, the Fourier transform of the four-point correlation of the density field can eventually be written down as

$$\begin{aligned} \int d^3r \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2) e^{-i\mathbf{k}\cdot\mathbf{r}} &= 8f_{\text{NL}}^2 \alpha_M(k) P_{\phi}(k) \int \frac{d^3k_1}{(2\pi)^3} \alpha_M(k_1) P_{\phi}(k_1) \int \frac{d^3k_2}{(2\pi)^3} \alpha_M(k_2) P_{\phi}(k_2) \alpha_M(|\mathbf{k} + \mathbf{k}_{12}|) \\ &\times \left\{ \frac{P_{\phi}(k_{12})}{P_{\phi}(k_2)} \left[1 + \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_{12}|)}{P_{\phi}(k)} \right] + \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_1|)}{P_{\phi}(k)} + \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_1|)}{P_{\phi}(k_1)} \right. \\ &\left. + \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_1|)P_{\phi}(|\mathbf{k} + \mathbf{k}_{12}|)}{P_{\phi}(k_1)P_{\phi}(k_2)} + \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_1|)P_{\phi}(|\mathbf{k} + \mathbf{k}_{12}|)}{P_{\phi}(k)P_{\phi}(k_2)} \right\} \end{aligned} \quad (56)$$

$$\begin{aligned} \int d^3r \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2, \mathbf{x}_2) e^{-i\mathbf{k}\cdot\mathbf{r}} &= 4f_{\text{NL}}^2 \int \frac{d^3k_1}{(2\pi)^3} \alpha_M(k_1) \alpha_M(|\mathbf{k} + \mathbf{k}_1|) P_{\phi}(k_1) \int \frac{d^3k_2}{(2\pi)^3} \alpha_M(k_2) \alpha_M(|\mathbf{k} + \mathbf{k}_2|) P_{\phi}(k_2) \\ &\times \left\{ P_{\phi}(|\mathbf{k} + \mathbf{k}_{12}|) \left[1 + 2 \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_1|)}{P_{\phi}(k_2)} + \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_1|)P_{\phi}(|\mathbf{k} + \mathbf{k}_2|)}{P_{\phi}(k_1)P_{\phi}(k_2)} \right] \right. \\ &\left. + P_{\phi}(k) \left[1 + 2 \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_1|)}{P_{\phi}(k_1)} + \frac{P_{\phi}(|\mathbf{k} + \mathbf{k}_1|)P_{\phi}(|\mathbf{k} + \mathbf{k}_2|)}{P_{\phi}(k_1)P_{\phi}(k_2)} \right] \right. \\ &\left. + 2P_{\phi}(|\mathbf{k} + \mathbf{k}_1|) \left[\frac{P_{\phi}(k_{12})}{P_{\phi}(k_1)} + \frac{P_{\phi}(k_{12})}{P_{\phi}(k_2)} \right] \right\} . \end{aligned} \quad (57)$$

In the large-scale limit $k \rightarrow 0$, these expressions asymptote to

$$\begin{aligned} \int d^3r \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2) e^{-i\mathbf{k}\cdot\mathbf{r}} &\approx 4f_{\text{NL}}^2 \sigma^4 S_3^{(1)}(M) \alpha_M(k) P_{\phi}(k) + 8f_{\text{NL}}^2 \alpha_M(k) \int \frac{d^3k_1}{(2\pi)^3} \alpha_M^2(k_1) P_{\phi}(k_1) \\ &\times \int \frac{d^3k_2}{(2\pi)^3} \alpha_M^2(k_2) \alpha_M(k_{12}) [P_{\phi}(k_1)P_{\phi}(k_2) + P_{\phi}(k_1)P_{\phi}(k_{12}) + P_{\phi}^2(k_{12})] \end{aligned} \quad (58)$$

$$\begin{aligned} \int d^3r \xi^{(4)}(\mathbf{x}_1, \mathbf{x}_1, \mathbf{x}_2, \mathbf{x}_2) e^{-i\mathbf{k}\cdot\mathbf{r}} &\approx 16f_{\text{NL}}^2 \sigma^4 P_{\phi}(k) + 16f_{\text{NL}}^2 \int \frac{d^3k_1}{(2\pi)^3} \alpha_M^2(k_1) \int \frac{d^3k_2}{(2\pi)^3} \alpha_M^2(k_2) P_{\phi}(k_2) P_{\phi}(k_{12}) \\ &\times [P_{\phi}(k_1) + P_{\phi}(k_2)] . \end{aligned} \quad (59)$$

Ignoring the second piece in the right-hand side of Eqs (58) and (59), the non-Gaussian correction to the halo power spectrum reads

$$\begin{aligned} \Delta P_{\text{hh}}(k) &= 4f_{\text{NL}} b_{\text{L}}^2 \delta_c(z) \alpha_M(k) P_{\phi}(k) \\ &+ 4f_{\text{NL}}^2 b_{\text{L}}^2 \delta_c^2(z) P_{\phi}(k) + \left(g_{\text{NL}} + \frac{4}{3} f_{\text{NL}}^2 \right) \\ &\times b_{\text{L}}^2 \delta_c^2(z) S_3^{(1)}(M) \alpha_M(k) P_{\phi}(k) . \end{aligned} \quad (60)$$

If we set $g_{\text{NL}} = 0$ and keep only the first two terms in the right-hand side, then the non-Gaussian (Eulerian) halo power spectrum can be cast into the form

$$P_{\text{hh}}(k) = [b(M) + f_{\text{NL}} b_{\phi}(k)]^2 P_{\delta}(k) \quad (61)$$

where the scale-dependent bias parameter $b_{\phi}(k)$ is

$$b_{\phi}(k) = 2 [b(M) - 1] \delta_c(z) \alpha_M^{-1}(k) . \quad (62)$$

Note that reference [47] obtained this relation by considering the halo power spectrum implied by a bias relation that is a local mapping of the density field. In practice, the term proportional to $P_{\phi}(k)$ is negligible as it contributes only at very small wavenumber $k \lesssim 0.001 h^{-1} \text{Mpc}$. The third term in the right-hand side of eq.(60) is derived in this paper for the first time. In the case $g_{\text{NL}} = 0$, its magnitude relative to the term linear in f_{NL} is $(1/3)f_{\text{NL}}\delta_c(z)S_3^{(1)}(M)$, which is approximately ~ 0.03 at redshift $z = 1.8$ and for a mass scale $M = 10^{13} M_{\odot}/h$. Although its contribution becomes increasingly important at higher redshift, it is fairly small

for the values of f_{NL} considered here. Consequently, we shall employ the approximation

$$\Delta b(k, \text{NL}) = \epsilon_\kappa \Delta b_\kappa(k, g_{\text{NL}}) + \Delta b_\kappa(k, f_{\text{NL}}) + [\Delta b_{\text{I}}(\text{NL}) + \epsilon_{\text{I}}] . \quad (63)$$

to describe the non-Gaussian bias of dark matter haloes.

The quadratic coupling $f_{\text{NL}}\phi^2$ also affects the matter power spectrum at leading order [27, 68], positive values of f_{NL} increasing the small scale power. However, the relative size of this k -dependent correction, $\beta_{\text{m}}(k, f_{\text{NL}})$, is at a per cent level only in the weakly nonlinear regime $k \lesssim 0.1 \text{ hMpc}^{-1}$ [42, 69] and fades rapidly as one goes to larger scales. We will thus neglect it in what follows.

In Fig. 10, the result of measuring ratios of auto- and cross-power spectra in the simulations with $(f_{\text{NL}}, g_{\text{NL}}) = (\pm 100, -3 \times 10^5)$ is shown at $0 < z < 1.5$ for the haloes with $M > 2 \times 10^{13} M_\odot/h$. We do not quote error bars since the data points are obtained by averaging over two realisations only. The solid and dashed curves show the theoretical prediction eq. (63). The value of the multiplicative factor $\epsilon_\kappa(b, g_{\text{NL}})$ was obtained from the four-parameter formula eq.(36), while $\Delta b_{\text{I}} + \epsilon_{\text{I}}$ was individually fitted for each halo sample over the wavenumber range $0.005 < k < 0.03 \text{ hMpc}^{-1}$. As can be seen, the theoretical expectation eq. (63) agrees reasonably well with the numerical data. This demonstrates that the range of validity of the non-Gaussian bias formula eq. (35) extends to smaller values of g_{NL} as well as models with non-vanishing f_{NL} and g_{NL} .

VI. DISCUSSION

In this paper we explored the effect of a local cubic coupling $g_{\text{NL}}\phi^3$ on the mass function and bias of dark matter haloes. We derived analytical expressions for these statistics and tested them against the outcome of numerical simulations.

We showed that current theoretical predictions of the non-Gaussian correction to the mass function reasonably agree with the simulations. The LV formula [56] appears to provide a better fit to the data than the MVJ formula [57], in agreement with some of the literature on the subject [35, 70]. The two approximations can be combined to provide an accurate description if one adjusts the low- ν expansion of the latter so as to match that of the former.

We found that the magnitude of the non-Gaussian scale-dependent bias $\Delta b_\kappa(k, g_{\text{NL}})$ is suppressed relative to a theoretical prediction based on the statistics of highly overdense regions, even on linear scales $k \lesssim 0.01 \text{ hMpc}^{-1}$. This suppression is stronger for the low biased samples $b \lesssim 3$ and, at fixed value of b , for positive values of g_{NL} . We were able to fit the measured halo bias at the expense of introducing two free parameters, ϵ_κ and ϵ_{I} , that depend mostly on the halo bias $b(M)$ and the coupling parameter g_{NL} . These parameters quantify the departure from the theoretical scale-dependent and

scale-independent non-Gaussian bias correction, respectively. We provide a simple fitting formula for $\epsilon_\kappa(b, g_{\text{NL}})$, eq.(36), which should be used when analyzing observational data. In non-Gaussian simulations of the f_{NL} type, the data also hint at a (possibly f_{NL} -dependent) suppression of the non-Gaussian scale-dependent bias relative to theory for wavenodes $k \lesssim 0.03 \text{ hMpc}^{-1}$ [42, 43, 44], but the effect is much weaker than seen in our simulations of g_{NL} models. Clearly, these results require a better theoretical modelling of the non-Gaussian halo bias.

Reference [44] argued that both the MVJ and LV analytic formula can be reconciled with measurements of the non-Gaussian fractional correction to the mass function once non-spherical collapse is included. In practice, the critical density for collapse is replaced by $\delta_c \rightarrow \sqrt{q}\delta_c$, where the value $q = 0.75$ is obtained from a fit to the mass function measured in simulations [71]. Reference [72] argued that such a relation is a consequence of the diffusive nature of the critical threshold for collapse. Their model predicts $q \simeq 0.8$, in good agreement with the findings of [44]. However, we found that substituting $\delta_c \rightarrow \sqrt{q}\delta_c$ in eq.(14) only modestly improve the agreement with the data. Regarding the non-Gaussian bias, it is not obvious how one could justify the replacement $\delta_c \rightarrow \sqrt{q}\delta_c$ given that the linear bias of our (Gaussian) halo samples converges towards the spherical collapse prediction ν^2/δ_c for large peak height.

A important ingredient is the choice of the halo identification algorithm. While we used a spherical overdensity (SO) finder, reference [72] considered a Friends-of-Friends (FoF) finder with a linking length $b = 0.2$. The question of how the spherical overdensity masses can be mapped onto friends-of-friends masses remains a matter of debate [73]. Since the peak height depends on the halo mass M through the variance $\sigma(M)$, any systematic difference will be reflected in the value of ν associated to a specific halo sample. This will in turn affect the size of the fractional deviation from the Gaussian mass function at some specified peak height. The sensitivity of the non-Gaussian mass function and bias to the halo finder will be presented elsewhere.

The observational bound on f_{NL} inferred by [37] from the clustering of a high redshift sample of quasars can be straightforwardly translated into a limit on g_{NL} since $\Delta b_\kappa(k, g_{\text{NL}})$ also scales as $k^{-2}T(k)^{-1}$ at low wavenumber. We have obtained

$$-3.5 \times 10^5 < g_{\text{NL}} < +8.2 \times 10^5 \quad (95\%). \quad (64)$$

These are the first limits derived on g_{NL} . While they are too weak to provide interesting constraints on inflationary scenarios such as the curvaton model, future all-sky redshift surveys should improve them by a factor of ~ 100 . Future CMB observations, including PLANCK, should also improve the limits derived here by an order of magnitude. With these improvements we expect that in the future realistic models of cubic type non-Gaussianity will be tested with real observations.

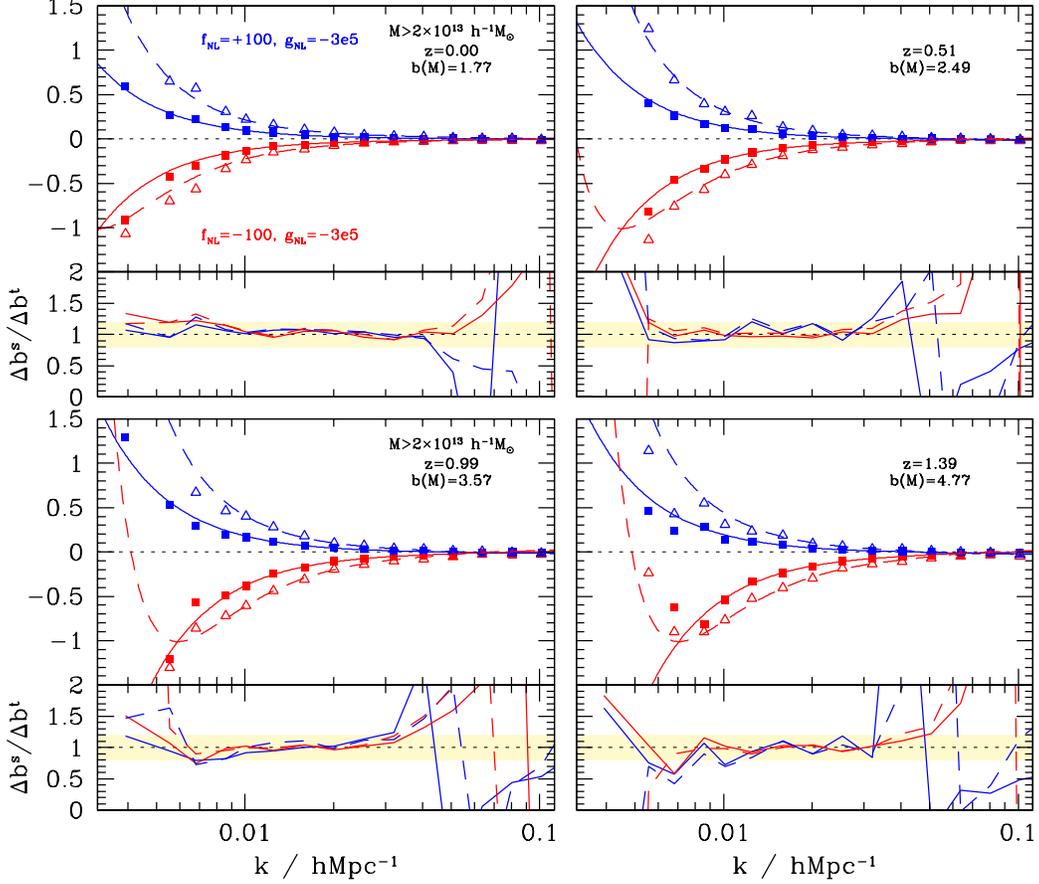


FIG. 10: Non-Gaussian bias correction measured in the simulations with $(f_{\text{NL}}, g_{\text{NL}}) = (\pm 100, -3 \times 10^5)$ for haloes of mass $M > 2 \times 10^{13} M_{\odot}/h$. Error bars are not shown as the data points are averaged over two realisations solely.

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APPENDIX A: SIGNAL-TO-NOISE FOR THE CMB TRISPECTRUM

As shown in [18], the invariance of the 4-point harmonic function of the CMB temperature anisotropy field under the 4! permutations of the coefficients $a_{l_i}^{m_i}$ imposes constraints on the CMB trispectrum $T_{l_3 l_4}^{l_1 l_2}(L)$ which can

be enforced by defining

$$\begin{aligned}
 T_{l_3 l_4}^{l_1 l_2}(L) &= P_{l_3 l_4}^{l_1 l_2}(L) + (2L+1) \sum_{L'} [(-1)^{l_2+l_3} \\
 &\quad \times \left\{ \begin{matrix} l_1 & l_2 & L \\ l_4 & l_3 & L' \end{matrix} \right\} P_{l_1 l_3}^{l_2 l_4}(L') + (-1)^{L+L'} \\
 &\quad \times \left\{ \begin{matrix} l_1 & l_2 & L \\ l_3 & l_4 & L' \end{matrix} \right\} P_{l_1 l_4}^{l_3 l_2}(L')], \quad (\text{A1})
 \end{aligned}$$

where curly brackets are Wigner-6j symbols,

$$\begin{aligned}
 P_{l_3 l_4}^{l_1 l_2}(L) &= t_{l_3 l_4}^{l_1 l_2}(L) + (-1)^{\sum_i l_i} t_{l_3 l_4}^{l_1 l_2}(L) + (-1)^{L+l_3+l_4} \\
 &\quad \times t_{l_4 l_3}^{l_1 l_2}(L) + (-1)^{L+l_1+l_2} t_{l_3 l_4}^{l_2 l_1}(L), \quad (\text{A2})
 \end{aligned}$$

and the reduced trispectrum $t_{l_3 l_4}^{l_1 l_2}(L)$ is symmetric under the exchange of its upper and lower indices and fully characterises the model.

The expansion coefficients a_l^m are related to the primordial curvature perturbation $\Phi(\mathbf{k})$ through

$$a_l^m = 4\pi(-i)^l \int \frac{d^3 k}{(2\pi)^3} \Phi(\mathbf{k}) g_{Tl}(k) Y_l^{m*}(\hat{\mathbf{k}}), \quad (\text{A3})$$

where $g_{\tau l}(k)$ is the radiation transfer function. The reduced trispectrum can be calculated from this relation once the four-point function $T_{\Phi}(\mathbf{k}_1, \mathbf{k}_2, \mathbf{k}_3, \mathbf{k}_4)$ is specified. For a local cubic coupling $g_{\text{NL}}\phi^3$,

$$t_{l_3 l_4}^{l_1 l_2}(L) = \int_0^\infty dr r^2 \beta_{l_2}(r) \beta_{l_4}(r) h_{l_1 L l_2} h_{l_3 L l_4} \times [\mu_{l_1}(r) \beta_{l_3}(r) + \beta_{l_1}(r) \mu_{l_3}(r)] \quad (\text{A4})$$

with

$$\beta_l(r) = \frac{2}{\pi} \int_0^\infty dk k^2 P_\phi(k) g_{\tau l}(k) j_l(kr) \quad (\text{A5})$$

$$\mu_l(r) = \frac{2}{\pi} \int_0^\infty dk k^2 g_{\text{NL}} g_{\tau l}(k) j_l(kr), \quad (\text{A6})$$

and

$$h_{l_1 L l_2} = \frac{1}{\sqrt{4\pi}} \pi_{l_1 L l_2} \begin{pmatrix} l_1 & l_2 & L \\ 0 & 0 & 0 \end{pmatrix}. \quad (\text{A7})$$

Note that most of the contribution to $t_{l_3 l_4}^{l_1 l_2}(L)$ comes from a small volume centred at the comoving distance r_* to the surface of last scattering.

The Sachs-Wolfe approximation $g_{\tau l}(k) \approx -j_l(kr_*)/3$ valid at low multipoles $l \ll 100$ provides a useful order-of-magnitude estimate [19, 20, 66]. In this limit, we can approximate $\mu_l(r)$ as $-g_{\text{NL}}\delta_D(r - r_*)/(3r_*^2)$ since we assume g_{NL} independent of wavenumber. Hence, the reduced trispectrum simplifies to

$$t_{l_3 l_4}^{l_1 l_2}(L) \approx 9g_{\text{NL}} C_{l_2}^{\text{SW}} C_{l_4}^{\text{SW}} (C_{l_1}^{\text{SW}} + C_{l_3}^{\text{SW}}) h_{l_1 L l_2} h_{l_3 L l_4}, \quad (\text{A8})$$

Inserting this expression successively into eqs (A2) and (A1) yields, after some simplification, the CMB trispectrum eq.(45). The following relation between the Wigner-3j and 6j symbols (e.g., Appendix A of [18]),

$$\begin{aligned} & \sum_{l'_3} (2l'_3 + 1) (-1)^{\Sigma + l'_3 - l_3 - m_1 - m'_1} \begin{Bmatrix} l_1 & l_2 & l_3 \\ l'_1 & l'_2 & l'_3 \end{Bmatrix} \\ & \times \begin{pmatrix} l_2 & l'_1 & l'_3 \\ m_2 & m'_1 & -m'_3 \end{pmatrix} \begin{pmatrix} l_1 & l'_3 & l'_2 \\ m_1 & m_3 & -m'_2 \end{pmatrix} \quad (\text{A9}) \\ & = \begin{pmatrix} l_1 & l_2 & l_3 \\ m_1 & m_2 & -m \end{pmatrix} \begin{pmatrix} l_3 & l'_1 & l'_2 \\ m & m'_1 & -m'_2 \end{pmatrix} \end{aligned}$$

where $\Sigma = l_1 + l_2 + l'_1 + l'_2$ and the value of m is set by the triangle condition, can be useful to derive the above result. The signal-to-noise summed up to multipole l_{max} eq. (44) then becomes

$$\begin{aligned} \left(\frac{S}{N}\right)^2 (< l_{\text{max}}) &= \left(\frac{27}{2\pi}\right)^2 g_{\text{NL}}^2 \sum_{l_1 > l_2 > l_3 > l_4}^{l_{\text{max}}} \pi_{l_1 l_2 l_3 l_4}^2 \\ & \times \frac{[C_{l_1}^{\text{SW}} C_{l_2}^{\text{SW}} C_{l_3}^{\text{SW}} + (\text{cyclic})]^2}{C_{l_1}^{\text{SW}} C_{l_2}^{\text{SW}} C_{l_3}^{\text{SW}} C_{l_4}^{\text{SW}}} \quad (\text{A10}) \\ & \times \sum_{L=0}^{2l_{\text{max}}} (2L+1) \begin{pmatrix} l_1 & l_2 & L \\ 0 & 0 & 0 \end{pmatrix}^2 \begin{pmatrix} l_3 & l_4 & L \\ 0 & 0 & 0 \end{pmatrix}^2. \end{aligned}$$

We can recast the sum over the diagonal modes L into a manifestly symmetric form with the aid of the Gaunt integral

$$\frac{1}{2} \int_{-1}^{+1} dx P_{l_1}(x) P_{l_2}(x) P_{l_3}(x) = \begin{pmatrix} l_2 & l_2 & l_3 \\ 0 & 0 & 0 \end{pmatrix}^2 \quad (\text{A11})$$

and the orthogonality relation

$$\sum_{k=0}^{\infty} (2k+1) P_k(x) P_k(y) = 2\delta_D(x-y), \quad (\text{A12})$$

where δ_D is the Dirac delta. We find

$$\begin{aligned} & \sum_{L=0}^{2l_{\text{max}}} (2L+1) \begin{pmatrix} l_1 & l_2 & L \\ 0 & 0 & 0 \end{pmatrix}^2 \begin{pmatrix} l_3 & l_4 & L \\ 0 & 0 & 0 \end{pmatrix}^2 \quad (\text{A13}) \\ & = \frac{1}{2} \int_{-1}^{+1} dx P_{l_1}(x) P_{l_2}(x) P_{l_3}(x) P_{l_4}(x). \end{aligned}$$

There is a strict equality because the Wigner-3j symbols vanish for $L > l_1 + l_2$. Eq. (48) then follows by replacing the above equality into eq.(A10) and summing over all the $4!$ permutations of the quadruplet (l_1, l_2, l_3, l_4) . Although eq. (48) becomes computationally expensive when $l_{\text{max}} \gg 100$ (because we are summing over redundant configurations), we found that it is quite efficient for $l_{\text{max}} \lesssim 200$.

Following [20], we can roughly estimate the dependence of $(S/N)^2$ on l_{max} by considering only the contribution of the $L = 1$ mode in eq.(A10). Consequently, the product of the Wigner-3j symbols squared reduces to $\sim l_1 l_3 \delta_{l_1-1, l_2} \delta_{l_3-1, l_4}$ and yields $(S/N)^2 \propto l_{\text{max}}^2$. However, including all L modes as in eq. (48) gives a steeper dependence, $(S/N)^2 \propto l_{\text{max}}^{2.6}$ (see Fig.8), due to the fact that the Wigner-3j symbols decay slowly with increasing L . This is quite apparent in the classical limit $l_1, l_2, L \gg 1$, where

$$\begin{aligned} & \begin{pmatrix} l_1 & l_2 & L \\ 0 & 0 & 0 \end{pmatrix}^2 \quad (\text{A14}) \\ & \approx \frac{1}{\pi} [(l_1 + l_2)^2 - L^2]^{-1/2} [L^2 - (l_1 - l_2)^2]^{-1/2}. \end{aligned}$$

This implies that the terms in the summation eq.(A10) decay only as $1/L$ for $L \gg l_1, l_2$. By contrast, the second order contribution to the CMB trispectrum induced by the quadratic coupling $f_{\text{NL}}\phi^2$ adds an additional multiplicative factor of $(C_L^{\text{SW}})^2$ in the summation over the L modes which increases the relative contribution of the low- L modes (since these now decay as L^{-5}). This is the reason why considering only $L \leq 10$ modes as done in [20] still provides a good approximation to the signal-to-noise of the CMB trispectrum for the f_{NL} model.

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