Loop quantum cosmology, non-Gaussianity, and CMB power asymmetry

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We argue that the anomalous power asymmetry observed in the cosmic microwave background (CMB) may have originated in a cosmic bounce preceding inflation. In loop quantum cosmology (LQC), the big bang singularity is generically replaced by a bounce due to quantum gravitational effects. We compute the spectrum of inflationary non-Gaussianity and show that strong correlation between observable scales and modes with longer (superhorizon) wavelength arise as a consequence of the evolution of perturbations across the LQC bounce. These correlations are strongly scale dependent and induce a dipole-dominated modulation on large angular scales in the CMB, in agreement with observations.

I. INTRODUCTION

The Planck team has reported evidence of anomalous features in the large-scale cosmic microwave background (CMB) which point towards violations of statistical isotropy [1]. Some of these anomalies were already observed by WMAP [2–4], and have been now reconfirmed [5], hence reducing the possibility of instrumental origin or systematics. Although the associated statistical significances are still inconclusive, the intriguing possibility that these features are traces of fundamental physics beyond the simplest inflationary models has attracted considerable attention in the theoretical community.

Among the different anomalies detected, the *power asymmetry*—unequal power in different regions of the sky—has received significant attention, and it is also the main target of this paper. This asymmetry was first modeled by adding a modulating factor to an otherwise statistically isotropic temperature distribution [6]

$$\delta T(\hat{n}) = (1 + A(\hat{n}))\delta T_{\mathrm{iso}}(\hat{n}), \tag{1}$$

where $A(\hat{n}) = \sum_{LM} A_{LM} Y_{LM}(\hat{n})$ is the modulating function. Planck has reported no evidence for A_{LM} different from zero for $L \geq 2$, but a statistically significant (about 3σ) $L = 1$ *dipolar* modulation has been found [1]. Furthermore, the effect of this dipole on the angular power spectrum C_ℓ has been detected *only* for low multipoles with $\ell < 64$. More precisely, after separating the ℓ range into bins of size $\Delta\ell = 64$, only the first bin shows a non-vanishing signal, with average amplitude $A_{L=1} = 0.07 \pm 0.02$. Observations are compatible with $A_{L=1} = 0$ for larger values of ℓ. This scale dependence implies that the simple parametrization (1) is insufficient to account for the observed modulation. The theoretical challenge is therefore to find a mechanism able to produce a dipolar modulation present only on large scales, with negligible contribution to

the quadrupole, octupole, etc., and furthermore respecting the existing constraints on the CMB: a remarkably Gaussian, almost scale-invariant spectrum of adiabatic perturbations. Not surprisingly, it has been difficult to find a completely satisfactory model [7].

One of the first and most compelling ideas to generate such a power asymmetry was introduced by Erickcek, Carroll, and Kamionkowski in Ref. [8]. It relies on the fact that the presence of a very long-wavelength, super-Hubble perturbation of a curvaton field will induce a dipole modulation in the observed spectrum, provided its wavelength is still short enough that its amplitude shows a gradient across the sky. This long mode could originate, for instance, as a remnant of the preinflationary epoch. This idea has been implemented by several authors and improved in different directions (see e.g. Refs. [9–11]).

Another interesting proposal, introduced in Ref. [12], works instead in a single-field inflationary model, and generates the power modulation through a non-Gaussian *coupling* between observable scales in the CMB and even larger, super-Hubble scales. The reason why non-Gaussianity can induce anisotropies in the observed spectrum, even if the underlying statistics is isotropic, is simply because a typical realization looks significantly more anisotropic if the underlying distribution is non-Gaussian. Or in other words, because observable modes couple to the *particular realization* of the long-wavelength modes in our Universe, which is generically anisotropic. This model, therefore, requires a mechanism to generate large correlations between very different scales. The consistency relation proposed in Ref. [13] tells us that it is difficult to find a realistic single-field model producing that type of non-Gaussianity.

In this work we introduce a scenario which has features in common with both of the previous ideas. We consider a single-field inflationary model preceded by a bounce described by loop quantum cosmology (LQC). This scenario has been analyzed in great detail, both for the background space-time and for perturbations (see

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Refs. [14–17] for reviews). In short, the evolution of perturbations across the bounce excites quanta out of an initial vacuum, and as a consequence the onset of inflation is reached in an excited state, rather than in the Bunch-Davies vacuum. The presence of those perturbations, remnants of the preinflationary phase, has an important impact on the non-Gaussianity generated during inflation. We compute those non-Gaussianities and the modulation they produce on observable scales, and show that the observed power asymmetry can have an origin on the quantum bounce preceding inflation, while still respecting all observational constraints. We set $c = 1$ but keep G and $\hbar$ explicitly. Numerical values are given in Planck units.

II. LQC AND THE POWER SPECTRUM

In LQC, the mean effective space-time geometry is described by equations which incorporate the leading quantum corrections to general relativity. For instance, the Friedmann equation reads [14,18]

$$H^2 = \frac{8\pi G}{3} \rho \left(1 - \frac{\rho}{\rho_{\max}} \right), \qquad (2)$$

where H is the Hubble rate, ρ the energy density, and $\rho_{\max} \approx 0.4\rho_{\rm Pl}$ its *upper bound*, which is a fraction of the Planck energy density $\rho_{\rm Pl}$. In this paper the matter sector is assumed to be a single scalar field with mass m. More complicated potentials can be incorporated. However, the effects we are seeking have a quantum gravity origin, and the form inflaton potential produces subleading contributions. At low energies $\rho \ll \rho_{\max}$, Eq. (2) becomes indistinguishable from the general relativistic Friedmann equation. However, close to the Planck scale, quantum gravitational effects break the linear relation between H^2 and ρ, and bring H to zero while ρ attains its maximum value. This is a quantum bounce that "bridges" a contracting and an expanding phase of the Universe.

The theory of cosmological perturbations in LQC that we use in this paper was developed in Ref. [19], and has been used in Refs. [19,20] to study the evolution of perturbations across the bounce through the end of inflation, and to compute the primordial power spectrum of tensor and scalar perturbations. This scenario is based on first principles, and trans-Planckian issues can be addressed squarely. The free parameters relevant for phenomenological studies are the value of the inflaton field at some arbitrary reference time, e.g. the bounce, $\phi(t_B) := \phi_B$, its mass m, and the initial state of perturbations at some initial time. References [19,20] have explored the predictions for the power spectrum across the parameter space and contrasted the results with observations.

From the viewpoint of inflation, the effect of the preceding bounce translates to an excited state $|\beta\rangle$ for perturbations at the onset of slow roll. Therefore, the LQC preinflationary evolution can be conveniently encoded in the Bogoliubov coefficients α_k and β_k relating $|\beta\rangle$ and the Bunch-Davies vacuum, i.e., relating the mode functions of curvature perturbations $\mathcal{R}_k(t)$ that result from the preinflationary evolution and the Bunch-Davies modes $\mathcal{R}_k^{\rm BD}(t)$. These coefficients can computed numerically for a choice of the free parameters. For $\phi_B = 1.22$ and $m = 1.10 \times 10^{-6}$, both in Planck units, and "Minkowski-like" vacuum initial condition[1] for perturbations at one Planck second before the bounce, we obtain that the average number of quanta present at the onset of inflation, given by $|\beta_k|^2$, is approximately 10^{-3} for the reference mode $k_\star$ which today corresponds to 0.002 Mpc^{-1}. For the longest-wavelength mode we can directly measure, $k_{\min} \approx k_\star/8.9$, we obtain $|\beta_{k_{\min}}|^2 = 1.2$, and $|\beta_k|^2 \sim 1/k$ for lower values of k [20]. Other choices of initial data give similar results (see Refs. [19,20] for details).

The inflationary scalar power spectrum in the presence of an excited state $|\beta\rangle$ is given by $P_{\mathcal{R}}(k) = P_{\mathcal{R}}^{\rm BD}(k)|\alpha_k + \beta_k|^2$, where $P_{\mathcal{R}}^{\rm BD}(k)$ is the Bunch-Davies result. Figure 1 shows the numerically computed $P_{\mathcal{R}}(k)$ for the choice of parameters mentioned above. Two energy scales play an important role in the power spectrum. First, LQC introduces a *new energy scale* $k_{\rm LQC}/a(t_B) := \sqrt{R_B/6} \approx 3.21$ that is directly related to the space-time scalar curvature at the bounce $R_B = 48\pi\rho_{\max} \approx 62$. A second scale is provided by the value of curvature at the onset of accelerated expansion $k_I/a(t_I) := \sqrt{R_I/6} \approx 10^{-5}$, which occurs at time t_I–notice that t_I is not the onset of *slow roll*, which happens at later times. Modes $k > k_{\rm LQC}$ are "inside" the curvature radius—i.e. its wavelength is shorter than the curvature radius—during the bounce, and exit during the slow-roll era. These perturbations reach the onset of slow roll in the Bunch-Davies vacuum, and their power spectrum is indistinguishable from standard results. Modes $k_I < k < k_{\rm LQC}$ exit the curvature radius soon before the bounce and reenter after it, to exit again during slow roll. The crossing process around the bounce time amplifies the amplitude of those modes, which then reach inflation in an excited state. Finally, modes $k < k_I$ are outside the Hubble radius during the entire evolution, even at the onset of inflation, and their spectrum is significantly suppressed. The reference mode $k_\star$, that today corresponds to 0.002 Mpc^{-1}, is approximately one third of $k_{\rm LQC}$, $k_\star = 0.36\, k_{\rm LQC}$. Consequently, the LQC corrections to the observable power spectrum only appear for the largest angular scales in the sky, that correspond to multiples

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[1]As explained in Ref. [19], although Minkowski-like initial conditions do not produce a state with the desired ultraviolet behavior—i.e. a Hadamard or fourth adiabatic order state—one can always modify the initial data for sufficiently ultraviolet modes to make the state ultraviolet regular. Since such modifications do not affect observable predictions, we do not describe the details in this paper, which can be found in Refs. [19,21].
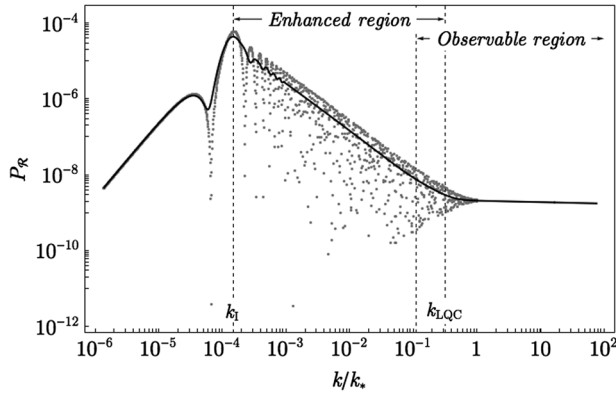
FIG. 1. Scalar power spectrum at the end of inflation for $\phi_B = 1.22$ and $m = 1.10 \times 10^{-6}$, and "Minkowski-like" vacuum initial conditions at one Planck second before the bounce [20]. Gray points show the numerical result for individual modes. The spectrum is rapidly oscillatory, and its average is shown in black. The spectrum is amplified for low wave numbers $k_I < k < k_{\mathrm{LQC}}$, but the enhancement is only significant for the very low multipoles in the CMB, for which the observational error bars are large because of cosmic variance. Therefore, this power spectrum is compatible with current observations.

$\ell \lesssim 30$, and are only significant for the lowest multiples. Furthermore, the LQC corrections significantly increase the power for modes with wavelengths larger than the Hubble radius today. This power enhancement conveniently reaches a maximum around k_I, at which perturbation theory is still well under control.

III. BACKREACTION ON THE INFLATIONARY GEOMETRY

An important question is whether the energy in the quanta present at the onset of slow roll produces significant backreaction on the inflationary space-time geometry which cannot be neglected. The energy density on the perturbations can be computed and compared to the vacuum result. The *difference* in energy density between the state $|\beta\rangle$ and the Bunch-Davies vacuum can be written as

$$\Delta\rho(t) = \frac{\hbar\epsilon}{4\pi G} \int \frac{d^3 k}{(2\pi)^3} \left\{ |\beta_k|^2 \left(|\dot{\mathcal{R}}_k^{\mathrm{BD}}|^2 + \frac{k^2}{a^2}|\mathcal{R}_k^{\mathrm{BD}}|^2 \right) \right.$$
$$\left. + \mathrm{Re}\left[\alpha_k \beta_k^\star \left[(\dot{\mathcal{R}}_k^{\mathrm{BD}})^2 + \frac{k^2}{a^2}(\mathcal{R}_k^{\mathrm{BD}})^2 \right] \right] \right\}, \quad (3)$$

where the integral in k is extended from $-\infty$ to ∞. It is well known that the difference of expectation values of the energy-momentum tensor between states that are at least of fourth adiabatic order is always finite, hence no renormalization is required in (3). Numerical evaluation shows that the ratio of $\Delta\rho$ with the background energy density ρ_0, is $\Delta\rho/\rho_0 \approx 10^{-3}$ at the onset of slow roll and decreases

exponentially with cosmic time. Therefore, since the energy density in the Bunch-Davies vacuum is known to be negligibly small, the backreaction of the state $|\beta\rangle$ on the inflationary geometry can be neglected.

IV. THE BISPECTRUM

At leading order in perturbation theory, non-Gaussianity is characterized by the bispectrum $B_{\mathcal{R}}(k_1, k_2, k_3)$, defined in terms of the three-point function

$$\langle \mathcal{R}_{\vec{k}_1} \mathcal{R}_{\vec{k}_2} \mathcal{R}_{\vec{k}_3} \rangle = (2\pi)^3 \delta^3(\vec{k}_1 + \vec{k}_2 + \vec{k}_3) B_{\mathcal{R}}(\vec{k}_1, \vec{k}_2, \vec{k}_3).$$

The non-Gaussianity generated from excited states in inflation has been analyzed by several authors (see e.g. Refs. [22–25]). It was pointed out in Ref. [23] that the main characteristic of the associated bispectrum is an enhancement in "squeezed" configurations which involve very different scales, $k_1 \approx k_2 \gg k_3$. These non-Gaussian correlations originate from quantum interactions between particles that are present at the onset of slow roll. In realistic models, including the scenario presented in this paper, the conclusions of the consistency relation proposed in Ref. [13] are still satisfied in the limit $k_1/k_3 \to 0$, but nevertheless there are important effects for small, but finite, values of k_1/k_3.

The full expression for the bispectrum $B_{\mathcal{R}}(k_1, k_2, k_3)$ as a function of the state $|\beta\rangle$ was obtained in Refs. [23–25], and is given in the Appendix. We have used those results to numerically compute $B_{\mathcal{R}}(k_1, k_2, k_3)$ at the end of inflation. The main difference from previous analysis is that the state $|\beta\rangle$ is now *computed* from a quantum gravity framework, rather than postulated. The resulting bispectrum, therefore, carries information about the preinflationary evolution. Those effects are more clearly displayed by plotting the ratio with the standard Bunch-Davies bispectrum $B_{\mathcal{R}}^{\mathrm{BD}}(k_1, k_2, k_3)$ (see Fig. 2). As expected, the bispectrum is strongly peaked in squeezed triangles, where the ratio $B_{\mathcal{R}}/B_{\mathcal{R}}^{\mathrm{BD}}$ shows a growth proportional to $\frac{k_{\mathrm{LQC}}^2}{k_1 k_3}$ —i.e., proportional to the product of the number of quanta in the modes k_1 and k_3, as expected for an enhancement originated from particle interactions. As for the power spectrum, the enhancement *only* appears for modes $k_I \lesssim k \lesssim k_{\mathrm{LQC}}$. Therefore, there are strong correlations between the largest wavelengths we can directly observe ($k \approx k_{\mathrm{LQC}}$) and super-Hubble modes with k values in the range $k_I \lesssim k \lesssim k_{\mathrm{LQC}}$. The bispectrum becomes negligibly small when the three momenta are in observable scales, hence respecting current constraints on non-Gaussianity. This scale dependence of the bispectrum will play an important role in the next section.

Recall that the three modes involved in the bispectrum must form a triangle, so it is really a function of two momenta, e.g. $B_{\mathcal{R}}(\vec{k}_1, \vec{k}_3)$. The shape of the bispectrum,
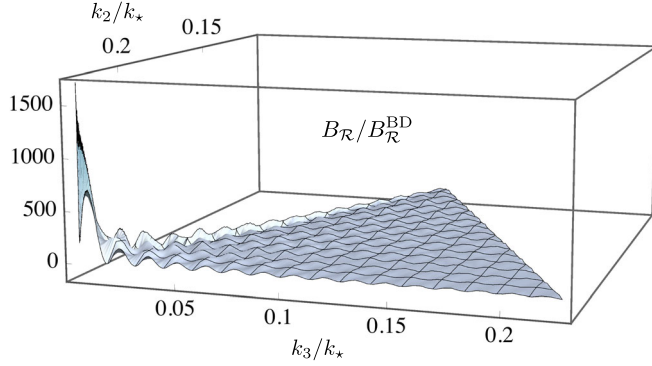
FIG. 2 (color online).   Ratio of the inflationary bispectrum for the excited initial state arising from the LQC preinflationary evolution versus the Bunch-Davies bispectrum, as a function of k_2 and k_3 for a fixed $k_1 = 0.22k_\star$. The plot shows the range of k_2 and k_3 allowed by the triangle condition $\vec{k}_1 + \vec{k}_2 + \vec{k}_3 = 0$. The bispectrum $B_\mathcal{R}$ is highly scale dependent and shows a prominent enhancement for squeezed configurations $k_3 \ll k_2 \approx k_1$.

shown in Fig. 2, can be understood by writing its dominant contribution in the squeezed limit $x := k_3/k_1 \ll 1$:

$$B_\mathcal{R} \approx 4\epsilon \Delta_\mathcal{R}(k_1)\Delta_\mathcal{R}(k_3) \times \mathrm{Re}\left[f_t \frac{1 - e^{i\tilde{k}_t\eta_0}}{1 + (1+\mu)x}\right.$$
$$\left. + f_1 \frac{1 - e^{i\tilde{k}_1\eta_0}}{(1+\mu)x} + f_2 \frac{1 - e^{i\tilde{k}_2\eta_0}}{(1-\mu)x} + f_3 \frac{1 - e^{i\tilde{k}_3\eta_0}}{1 + (-1+\mu)x}\right],$$
$$(4)$$

where $\mu = \hat{k}_1 \cdot \hat{k}_3$, $k_t = k_1 + k_2 + k_3$, $\tilde{k}_i = k_t - 2k_i$, and $\Delta_\mathcal{R}(k) = \hbar \frac{2\pi^2}{k^3}\left(\frac{H}{\dot{\phi}}\right)^2\left(\frac{H}{2\pi}\right)^2$ evaluated at Hubble crossing during inflation. The functions f_i contain the information of the state $|\beta\rangle$ and are defined in the Appendix. The parameter η_0 is the value of the conformal time at the onset of slow roll. For the example $\phi_B = 1.22$ and $m = 1.1 \times 10^{-6}$ we obtain $|\eta_0 k_\star| \approx 10^3$, so all observable modes are deeply inside the Hubble radius at the onset of slow roll.

The largest contribution to (4) comes from very squeezed configurations for which $|\tilde{k}_3\eta_0| \ll 1$. In that limit $B_\mathcal{R}$ is dominated by the terms proportional to f_t and f_3. This contrasts with the example considered in Refs. [23,25] of a scale-invariant excited state $|\beta\rangle$, where additionally the analysis was restricted to observable scales for which $|\eta_0 k_i| \gg 1$ for all i. In that case the leading contributions come from the terms proportional to f_1 and f_2.

It is important to note that the non-Gaussianity computed here is generated during slow-roll inflation. Extra contributions to the bispectrum will certainly arise from the evolution across the bounce. However, since the state of perturbations is very close to vacuum at the bounce time, those contributions are expected to be subdominant. This expectation is indeed borne out in explicit computations in bouncing models [26,27], where it is shown that the

non-Gaussianities generated across the bounce are much smaller than the bispectrum shown in Fig. 2, particularly for squeezed triangles, which are the relevant configurations for this paper. It is therefore reasonable to assume that the leading contributions to the bispectrum in squeezed configurations are dominated by inflationary non-Gaussianity.

V. NON-GAUSSIAN MODULATION

We follow some of the ideas presented in Ref. [12] to compute the CMB power modulation arising from coupling with super-Hubble scales. The statistics of the classical Bardeen potential Φ for observable scales is modified in the presence of a given long-wavelength perturbation $\Phi(\vec{k}_l)$ if the spectrum is non-Gaussian. Concretely, the two-point function acquires off-diagonal contributions of the form [28,29]

$$\langle\Phi_{\vec{k}}\Phi_{\vec{k}'}\rangle = P_\Phi(k)[(2\pi)^3\delta(\vec{k}+\vec{k}') + G(\vec{k},\vec{k}_l)\Phi(\vec{k}_l)], \quad (5)$$

where $\vec{k}_l = -(\vec{k}+\vec{k}')$ (i.e. it closes a triangle with $\vec{k}$ and $\vec{k}'$), $P_\Phi(k)$ is the Φ-power spectrum, and we take $G(\vec{k},\vec{k}_l) = \frac{5}{3}B_\mathcal{R}(\vec{k},\vec{k}_l)[\Delta_\mathcal{R}(k)\Delta_\mathcal{R}(k_l)]^{-1}$. For bispectra peaked in squeezed configurations, the leading contributions to $G(\vec{k},\vec{k}_l)$ come from the regime $k_l \ll k$, for which $k \approx k'$. Therefore, the correlation function is close to being diagonal. The off-diagonal terms break both homogeneity and isotropy. These terms vanish if we average over $\Phi_{\vec{k}_l}$, as expected, since our model respects these symmetries at the fundamental level. But in the particular realization of $\Phi_{\vec{k}_l}$ chosen by our Universe, the modulation term may be important, and may produce deviations from homogeneity and isotropy in the CMB much larger than would be expected from a typical realization of a Gaussian spectrum.

The off-diagonal terms in (5) source analogous terms in the covariance matrix of temperature spherical harmonic coefficients

$$\langle a_{\ell m} a^\star_{\ell'm'}\rangle = \delta_{\ell\ell'}\delta_{mm'}C_\ell + \sum_{LM}A_{LM}\mathcal{G}^{\ell\ell'L}_{-mm'M}(C_\ell + C_{\ell'}),$$

where $\mathcal{G}^{\ell_1\ell_2\ell_3}_{m_1m_2m_3}$ is a product of Wigner 3-j symbols, and C_ℓ is the standard angular power spectrum. The momentum dependence of the kernel $G(\vec{k},\vec{k}_l)$ factorizes for the dominant contributions in a power law for k (except for typical oscillations which are unimportant for our computations and can be averaged). The kernel can be expanded in Legendre polynomials

$$G(\vec{k},\vec{k}_l) = \sum_L g_L(k_l)\left(\frac{k_\star}{k}\right)^{\alpha_L} P_L(\mu),$$

where $\mu = \hat{k}\cdot\hat{k}_l$, and $g_L(k_l)$ encodes the k_l dependence. Then, the mean value of the modulation amplitude,

averaged over M, for a typical realization of the long-wavelength modes can be computed as (see Ref. [12] for further details of the computation)

$$\langle A_L(\ell)\rangle = \frac{C_\ell(-\alpha_L)}{2C_\ell(0)}\left[\int \frac{dk_l k_l^2}{(2\pi)^3}|g_L(k_l)|^2 P_\Phi(k_l)\right]^{1/2}, \quad (6)$$

where the coefficients $C_\ell(\alpha_L)$ have been defined as the CMB temperature power spectrum computed by replacing $n_s \rightarrow n_s + \alpha_L$. Therefore, if $\alpha_L = 0$ for all L, i.e. if $G(\vec{k}, \vec{k}_l)$ does not depend on k, the modulation amplitude is ℓ independent.

The value of $\langle A_L(\ell)\rangle$ for different multipoles L can be computed numerically from the bispectrum shown in Fig. 2, and some of its properties can be qualitatively understood already from Eq. (4). The leading contribution comes from the terms proportional to f_t and f_3, which are dominated by a *dipole* $L = 1$. The terms proportional to f_1 and f_2 contribute to a subdominant quadrupole. The octopole $L = 3$ is subdominant compared to the quadrupole, and this hierarchy continues for higher multipoles. Since the factors f_i in (4) scale approximately as $(k_l k)^{-1}$ for every i, we have $\alpha_L \approx 1$ for all L, and therefore the amplitude $A(\ell)$ is expected to *decrease* with ℓ.

Figure 3 shows $\langle A_L(\ell)\rangle$ for $L = 1, 2, 3$ as a function of ℓ. The amplitudes $\langle A_L(\ell)\rangle$ are all scale dependent, and such that $\langle A_1\rangle \gtrsim 10\langle A_2\rangle \gtrsim 10\langle A_3\rangle$. The average value of $\langle A_1\rangle$ for $\ell \lesssim 64$ is in agreement with the observed value $A_1^{\text{obs}} = 0.07 \pm 0.02$. Other values of our free parameters ϕ_B and m, and of the initial conditions for perturbations, can decrease or increase the value of $\langle A_1\rangle$. For instance, we find that choosing vacuum initial data for perturbations far into the past of the bounce significantly decreases the amplitudes of all multipoles.

There has been some debate based on symmetry arguments about whether a statistically homogeneous and isotropic bispectrum can generate a dipole modulation. Our computation is an example where the answer is in the affirmative. Furthermore, as shown in Ref. [30], a careful analysis reveals that symmetry arguments restrict A_{L+1} to be suppressed with respect to A_L for even L. The hierarchy of multipoles we find here, $\langle A_0\rangle > \langle A_1\rangle > \langle A_2\rangle > \cdots$, is therefore in agreement with those restrictions.

VI. DISCUSSION

We have presented an inflationary scenario in which perturbations start the slow-roll phase in an excited state, rather than the Bunch-Davies vacuum. This state arises from the preinflationary evolution provided by loop quantum cosmology, in which the big bang singularity is replaced by a bounce. Two new scales appear in the problem, k_I and k_{LQC}, related to the onset of the exponential expansion and the bounce, respectively. The number density of quanta at the onset of slow roll is significant only for the range $k_I \lesssim k \lesssim k_{\text{LQC}}$. During inflation these excitations induce non-Gaussian correlations, which we have computed. To the best of our knowledge, this is the first computation of non-Gaussianity in LQC. The result is compatible with existing observational constraints. Furthermore, large correlations arise between the longest modes we can observe, with $k \approx k_{\text{LQC}}$, and super-Hubble modes with $k \gtrsim k_I$. We have shown that those non-Gaussian correlations, which involve super-Hubble modes, are able to modify the observed power spectrum at large scales, inducing correlations between CMB angular multipoles $\ell \lesssim 30$ that differ in $\Delta\ell \approx 1$. These correlations are strongly scale dependent, and produce a power asymmetry in agreement with observations.

Other observed anomalies at large scales [1]—parity violation, power suppression for a bunch of multipoles around $\ell = 20$, multipole alignments, etc.—seem also to require correlations between low multipoles qualitatively similar to the ones obtained in this paper. Whether the agreement is also quantitative will be analyzed in future work.

We emphasize that the fact that the power spectrum shown in Fig. 1 is amplified for low wave numbers $k_I < k < k_{\text{LQC}}$ is not necessarily in conflict with observations, since i) the enhancement is only significant for the very low multipoles in the CMB, for which the observational error bars are large because of cosmic variance; ii) the observed power suppression is only significant for a few multipoles around $\ell = 20$, indicating that the effect is more likely to originate from correlations between multipoles, rather than a suppression of the primordial two-point function.

Some of the features appearing in our scenario—remnants, large-amplitude perturbations associated with an infrared scale k_I, correlations with super-Hubble modes, etc.—have been identified in previous phenomenological
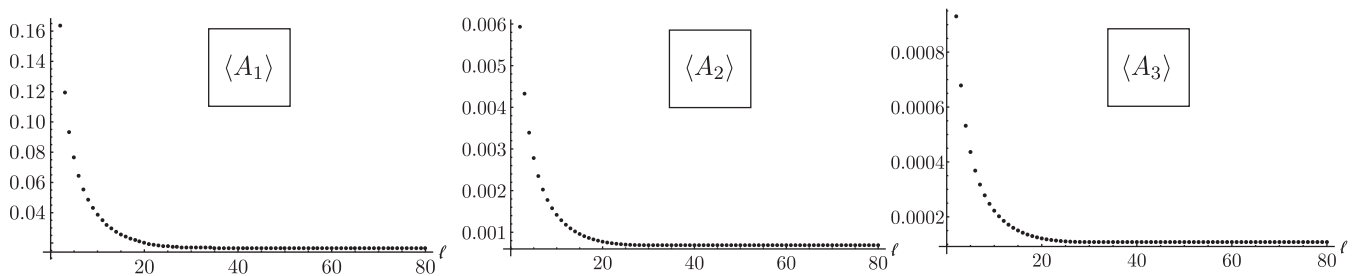

FIG. 3.   Amplitude $\langle A_L(\ell)\rangle$ as a function of ℓ for $L = 1, 2, 3$, for parameter values $\phi_B = 1.22$ and $m = 1.10 \times 10^{-6}$, and "Minkowski-like" vacuum initial conditions at 1 Planck second before the bounce.

analysis (see e.g. Refs. [8,12]) as ingredients needed to account for some of the observed anomalies at large scales. Here, these features arise from a concrete quantum gravity proposal based on first principles. Therefore, our results provide further motivation to consider the observed anomalies as real physical features, which have an origin beyond the simplest inflationary models, rather than statistical flukes or instrumental noise. Future work will be addressed to provide further robustness to the model introduced here and to extend its quantitative predictions.

ACKNOWLEDGMENTS

APPENDIX: INFLATIONARY BISPECTRUM FROM AN EXCITED STATE

The expression for the scalar bispectrum generated during slow-roll inflation when the state of perturbations $|\beta\rangle$ is given by a Bogoliubov transformation of the Bunch-Davies vacuum, with coefficient α_k and β_k, is given by [22–25]

$$B_{\mathcal{R}}(k_1, k_2, k_3) = \Delta_{\mathcal{R}}(k_1)\Delta_{\mathcal{R}}(k_2)\left\{\frac{1}{2}\left(3\epsilon - 2\eta + \epsilon\frac{k_1^2 + k_2^2}{k_3^2}\right)\right.$$
$$+ 4\epsilon\frac{k_1^2 k_2^2}{k_3^3}\text{Re}\left[f_t\frac{1 - e^{ik_t\eta_0}}{k_t} + f_1\frac{1 - e^{i\tilde{k}_1\eta_0}}{\tilde{k}_1}\right.$$
$$\left.\left. + f_2\frac{1 - e^{i\tilde{k}_2\eta_0}}{\tilde{k}_2} + f_3\frac{1 - e^{i\tilde{k}_3\eta_0}}{\tilde{k}_3}\right]\right\}$$
$$+ 2 \text{ cyclic permut. } k_1 \to k_2 \to k_3, \tag{A1}$$

where $k_t = k_1 + k_2 + k_3$, $\tilde{k}_i = k_t - 2k_i$, and the parameter η_0 is the value of the conformal time at the onset of slow roll. The functions f_i contain the information of the state $|\beta\rangle$:

$$f_t = [1 + F(k_1)(1 + F(k_2)) + \text{cyclic perm.}],$$
$$f_1 = F(k_1)^\star[1 + F(k_2) + F(k_3)] - F(k_2)^\star F(k_3),$$

and f_2 and f_3 can be obtained by cyclically permuting the momenta in f_1, and $F(k) = |\beta_k|^2 + \alpha_k^\star \beta_k$. The star indicates complex conjugation. Also, $\Delta_{\mathcal{R}}(k) = \hbar\frac{2\pi^2}{k^3}\left(\frac{H}{\dot{\phi}}\right)^2\left(\frac{H}{2\pi}\right)^2$ evaluated at Hubble exit during inflation.

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