Classicalization of inflationary perturbations by collapse models in the light of BICEP2

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Inflationary and hence quantum origin of primordial perturbations is on a firmer ground than ever post the BICEP2 observations of primordial gravitational waves. One crucial ingredient of success of this paradigm rests on explaining the observed classicality of cosmological inhomogeneities despite their quantum origin. Although decoherence provides a partial understanding of this issue, the question of single outcome motivates the analysis of quantum collapse models in cosmological context which generically modify the dynamics of primordial perturbations and hence can leave their imprints on observables. We revisit one such recently proposed working model of classicalization by spontaneous collapse [1] in the light of BICEP2 observations to look for possible modifications to tensor power spectra and their implications. We show that it can potentially change the consistency relation of single-field models and a precise measurement of $n_T$ and its running could serve as a test of such dynamics in the early universe.

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I. INTRODUCTION

The recent observation of $B$–mode polarization of Cosmic Microwave Background Radiation (CMBR) by BICEP2 [2] has put the cosmological inflationary paradigm [3, 4] on stronger footing than ever by confirming one of its many observational predictions, namely the generation of primordial tensor modes or primordial gravitational waves. Inflationary dynamics gives rise to both scalar and tensor fluctuations during inflation which then ‘redshift’ out of the horizon and freeze. These primordial scalar fluctuations, up on re-entering the horizon at later stages, give rise to $TT$ anisotropy spectrum of the CMBR that has been measured by many observations such as WMAP [5] and PLANCK [6]. On the other hand, the tensor modes give rise to the $BB$ spectrum of CMBR which has recently been detected by BICEP2 [2].

Detection of $B$–modes of CMBR is of profound importance as according to the simplest model of inflation it sets the scale of inflation in a unique way, which turns out to be surprisingly close to the Grand Unification (GUT) scale ($\sim 10^{16}$ GeV). But this is not all what the detection of $B$–modes by BICEP2 can provide. BICEP2 measurement has further reinforced the quantum nature of gravity as it indicates that the tensor perturbations generated during inflation are of quantum nature. BICEP2 has measured the power of these tensor modes over the scalar ones on large cosmological scales, called the tensor-to-scalar ratio $r$ and the measured value of $r$ by BICEP2 is [2]

$$r = 0.2^{+0.07}_{-0.05}.$$ (1)

It has been argued in [7] that if these primordial gravitational waves were of classical nature, it would have produced negligible tensor amplitude compared to the scalar one and the value of $r$ would have at least been suppressed by an extra power of slow-roll parameter $\epsilon$ in comparison to the standard scenario. That these primordial tensor modes are indeed of quantum nature has also been argued in other literature such as [8–10].

Primordial scalar perturbations are also generated quantum mechanically and this scenario is well supported by observations.

From this discussion it is evident that the primordial modes, both scalar and tensor, are being generated quantum mechanically at very early times. But these quantum fluctuations are the ones which give rise to CMBR temperature fluctuations and its $B$–mode polarization on large angular scales which are classical in nature. This leads us to the problem of quantum-to-classical transition in the cosmological context, which is a more serious form of the so-called “quantum measurement problem” which also prevails in laboratory systems.

Heuristically, the classical nature of primordial perturbations is argued by the large occupation number of superhorizon modes and effective irrelevance of the commutator of the field variable and its conjugate momentum on superhorizon scales. However, the perturbations generically evolve into highly squeezed states [8, 13, 14] on superhorizon scales, which are highly non-classical states. But the quantum nature is not directly evident in observations because it can be shown that quantum expecta-

1 However, it has been pointed out that large enough primordial gravitational wave signal can be produced via purely classical mechanisms, such as gravitational bremsstrahlung, in a multifield scenario [11] or non-perturbatively generated gravitational waves via electromagnetic fields amplified by an axion-like inflaton [12] which can overshadow the quantum tensor signal.
tions of highly squeezed states are indistinguishable from average of a classical stochastic field. This is so-called ‘decoherence without decoherence’ [13]. Furthermore decoherence, despite the ambiguity in system and environment split for a cosmological scenario, selects the field amplitude basis as the pointer basis of the system and justifies the standard calculation. At this point, what still remains unresolved is the issue of single outcome. This problem is also present in laboratory systems and only gets more intriguing in cosmological context [15]. One possible way out is to appeal to the many-worlds interpretation. The alternative is the so-called collapse models which we explore in this paper. Just as in laboratory systems, it is important to investigate whether and how collapse models can be distinguished from the standard quantum mechanical setup in a cosmological context.

In a generic collapse model Schrödinger equation is modified by adding stochastic and non-linear terms. The stochastic nature of the equation helps explaining the probabilistic outcome of quantum measurements without allowing for superluminal communication and the presence of non-linear terms breaks down the underlying superposition principle of quantum mechanics. For a detailed review on collapse models refer to [16]. Though a proper field theoretic treatment of collapse dynamics is not yet known, a few attempts have been made to apply such collapse models, especially Continuous Spontaneous Localization (CSL) model [16], into inflationary dynamics to resolve the problem of quantum to classical transition in cosmological context, such as [12, 14, 17]. As these collapse models explicitly modify the standard dynamics, we can anticipate that its observational implications would diverge from the standard dynamics and the aim of this brief paper is to determine how and where these collapse mechanisms differ from the standard scenario observationally within the context of an illustrative example in the light of recent BICEP observation.

II. GENERIC SINGLE-FIELD SLOW-ROLL INFLATION AND ITS OBSERVATIONAL IMPLICATIONS

The most economical model of inflation is the so-called slow-roll single field model. The power spectra for scalar and tensor perturbations in such a model can be given as [18]

\[ P_R = \frac{1}{8\pi^2} \left( \frac{H^2}{\epsilon M_{Pl}^2} \right) \left( \frac{k}{k_*} \right)^{n_s-1} A_s(k_*) \left( \frac{k}{k_*} \right)^{n_s-1}, \]

\[ P_T = \frac{2}{\pi^2} \left( \frac{H^2}{M_{Pl}^2} \right) \left( \frac{k}{k_*} \right)^{n_T} A_T(k_*) \left( \frac{k}{k_*} \right)^{n_T} \]

respectively, where \( k_* \) is called the pivot scale and for PLANCK it is chosen to be \( k_* = 0.05 \text{ Mpc}^{-1} \). It is to be noted that as the comoving curvature perturbations, denoted by \( R \), and the tensor modes, denoted by \( h \), freeze on superhorizon scales, it is customary to derive the power spectra given above at horizon crossing of each mode \( (k = aH) \) during inflation. In the above equations \( H \) is the Hubble parameter during inflation, \( M_{Pl} = 2.4 \times 10^{18} \text{ GeV} \) is the reduced Planck mass, \( \epsilon \) is the first Hubble slow-roll parameter defined as

\[ \epsilon \equiv -\frac{\ddot{H}}{H^2}. \]

\( n_s \) and \( n_T \) are the scalar spectral index and the tensor spectral index respectively which are the measures of the scale dependence of the respective spectrum and for this generic slow-roll single-field model turn out to be

\[ n_s - 1 = 2\eta - 4\epsilon, \]

\[ n_T = -2\epsilon, \]

where \( \eta \) is the second Hubble slow-roll parameter defined as

\[ \eta \equiv -\frac{\dot{\phi}_0}{H\dot{\phi}_0}, \]

where \( \dot{\phi}_0 \) is the inflaton field.

Each of the tensor modes, which are the traceless and transverse part of the metric fluctuations, is associated with two helicity states, often denoted as \( + \) and \( \times \) polarization. As has been first pointed out by Grishchuk in [19], the evolution of the Fourier modes of each of the helicity states of these tensor modes is identical to that of a massless scalar in de Sitter space with the correspondence

\[ h_{a}^s = \frac{2}{a M_{Pl} v_0^s}, \]

where \( s = +, \times \) is the helicity states and \( v_0^s \) is the re-defined tensor modes. Hence the evolution of quantum tensor modes can be reduced to that of two decoupled massless scalar modes. We will make use of this fact while dealing with CSL-modified inflationary dynamics.

Now, let us consider the observational implications of this simplest model of inflation. PLANCK measures the amplitude \( A_s \) and the scalar spectral index \( n_s \) as [2]

\[ A_s = 2.215 \times 10^{-9}, \]

\[ n_s = 0.9603 \pm 0.0073, \]

from the \( TT \) anisotropy spectrum of the CMBR. BICEP2, by detecting the \( B \)-polarization of CMBR measured tensor-to-scalar ratio \( r \), defined as

\[ r \equiv \frac{A_T}{A_s}, \]

to be 0.2. According to the simplest single field model the tensor-to-scalar ratio turns out to be of the order of slow-roll parameter:

\[ r = 16\epsilon, \]
Hence the BICEP2 measurement of \( r \) as 

\[ r = -2n_T. \] 

Thus independent measurements of \( r \) and \( n_T \) can unambiguously determine whether the inflationary dynamics was indeed simple or not.

The recent observation of \( r \) by BICEP2 and the measurement of scalar amplitude \( A_s \) by PLANCK indirectly provides the amplitude of the tensor perturbations. But one can see from Eq. (2) that the amplitude of the tensor power solely depends on the Hubble parameter during inflation and considering the central values of observed \( r \) and \( A_s \) one gets the Hubble parameter during inflation as

\[ H \simeq 1.1 \times 10^{14} \text{GeV}. \] 

Now, during inflationary era the universe is dominated by the potential energy of the inflaton field and thus the Friedmann equation during inflation is written as

\[ H^2 = \frac{1}{3M_{Pl}^2}V(\varphi_0). \] 

Hence the BICEP2 measurement of \( r \) also sets the scale of inflation as

\[ V^{1/4} \sim 2.1 \times 10^{16} \text{GeV}, \]

which is very close to the GUT scale. However, inferring the scale of inflation from tensor amplitude becomes more involved if one invokes large extra dimensions \([20]\) or classical sources of primordial gravitational waves in a multi-field scenario \([11]\).

Another implication of the BICEP2 result comes from the Lyth bound \([21]\). Writing the first slow-roll parameters \( \epsilon \) as

\[ \epsilon = M_{Pl}^2 \frac{H^2}{V_0}, \]

one can determine the field excursion during the inflation as

\[ \Delta \varphi_0 = O(1) \left( \frac{r}{0.01} \right)^{1/2} M_{Pl}, \]

which shows that the inflaton field excursion is super-Planckian during inflation if \( r < 0.2 \), as has been observed by BICEP2. This leads to some tension with effective field theory description of inflation as the field excursion becomes of the same order of the natural cutoff scale \(^2\).

\(^2\) However there is a recent debate in literature whether sub-Planckian field excursion can be made consistent with recent BICEP2 observation in a single-field scenario \([22]\).

### III. CSL-MODIFIED SINGLE-FIELD DYNAMICS AND ITS OBSERVATIONAL IMPLICATIONS

#### A. Scalar perturbations

For this analysis we would keep our focus on the derivations done in \([1, 15]\). The collapse models modify the dynamics of particles in the Schrödinger picture. Hence while applying the CSL modifications to the inflationary dynamics one requires to analyze the field dynamics in the Schrödinger picture. An elaborate analysis of Schrödinger picture dynamics of scalar perturbations during inflation is given in \([13]\) (see also \([24]\)).

We would study the primordial scalar perturbations during inflation in terms of the so-called Mukhanov-Sakasi variable, a gauge-invariant quantity denoted by \( \zeta \), which is related to the comoving curvature perturbation as

\[ \zeta(\tau, x) = \frac{a\varphi_0}{H} R(\tau, x). \] 

In Schrödinger picture, the standard scalar perturbation is analyzed in terms of its wave-functional \([13]\) defined as

\[ \Psi[\zeta(\tau, x)] = \prod_k \Psi^R_k[\zeta^R_k(\tau)] \Psi^I_k[\zeta^I_k(\tau)] \] 

where we have

\[ \zeta^I_k(\tau) = \frac{1}{\sqrt{2}} \left( \zeta^R_k(\tau) + i\zeta^I_k(\tau) \right). \]

These wave functionals satisfy the functional Schrödinger equation as

\[ i \frac{\partial \Psi^I_k}{\partial \tau} = \hat{H}_k \Psi^I_k, \]

where the Hamiltonian \( \hat{H}_k \equiv \hat{H}^R_k + \hat{H}^I_k \) and the ground state solution of the functional Schrödinger equation is written as

\[ \Psi^R_k[\tau, \zeta^R_k] = \sqrt{N_k(\tau)} \exp \left( -\frac{\Omega_k(\tau)}{2} \left( \zeta^R_k \right)^2 \right). \]

Here \( \Omega_k \) is related to the mode functions \( f_k \) in the Heisenberg picture as

\[ \Omega_k = \frac{f_k^*}{f_k}. \]

Thus, in Schrödinger picture the power spectrum of \( \zeta \) turns out to be

\[ \mathcal{P}_\zeta(k) = \frac{k^3}{2\pi^2} |f_k|^2 = \frac{k^3}{2\pi^2 \text{Re} \Omega_k}. \]

and thus the power spectrum of the comoving curvature perturbations can be obtained as

\[ \mathcal{P}_\zeta(k) = \frac{k^3}{8\pi^2 \epsilon M_{Pl}^2 \alpha^2 \text{Re} \Omega_k}. \]
It has been proposed in [15] that CSL-like modifications can be applied to the inflationary perturbations directly in the Fourier space by adding CSL-like non-linear and stochastic terms to the functional Schrödinger equation of \( \zeta \) which then looks like

\[
dW_{k}^{R,1} = -i\gamma_{k}^{R,1} d\tau + \sqrt{\gamma} \left( \zeta_{k}^{R,1} - \langle \zeta_{k}^{R,1} \rangle \right) dW_{\tau} \tag{24}
\]

where the stochastic behavior due to CSL mechanism is encoded in the Wiener process \( W_{\tau} \) and \( \gamma \) is called the collapse parameter. The most general stochastic wave-functional which satisfies this stochastic functional Schrödinger equation can be written as

\[
\Psi_{k}^{R,1} \left( \tau, \zeta_{k}^{R,1} \right) = \sqrt{N_{k} \left( \tau \right)} \exp \left\{ i\sigma_{k}^{R,1} \left( \tau \right) + i\chi_{k}^{R,1} \left( \tau \right) \right\} \frac{\sqrt{\gamma}}{2} \left( \zeta_{k}^{R,1} - \langle \zeta_{k}^{R,1} \rangle \right)^{2} \frac{d\tau}{2} - \frac{\text{Re} \Omega_{k} \left( \tau \right)}{2} \left( \zeta_{k}^{R,1} - \langle \zeta_{k}^{R,1} \rangle \right)^{2} - i\frac{\text{Im} \Omega_{k} \left( \tau \right)}{2} \left( \langle \zeta_{k}^{R,1} \rangle \right)^{2}, \tag{25}
\]

where \( \zeta_{k}^{R,1}, \sigma_{k}^{R,1} \) and \( \chi_{k}^{R,1} \) are real numbers.

In [13], the collapse parameter \( \gamma \) was taken to be constant and it was inferred that such a case is not capable of explaining the quantum-to-classical transition of primordial modes. Then it was argued in [1] that taking \( \gamma \) to be constant the scenario loses one of the crucial features of CSL dynamics, known as amplification mechanism. In CSL-modified quantum mechanics, the collapse parameter is taken to be directly proportional to the mass of the system and its number density reflecting the fact that heavier objects become classical faster than the lighter ones. Similarly, in tune with the expectation that superhorizon modes behave classically one could assume in the cosmological context that the collapse parameter would become stronger as a generic mode starts to cross the horizon facilitating its collapse to one of its field eigenstates. Hence a phenomenological form of the collapse parameter was proposed in [1] as

\[
\gamma = \frac{\gamma_{0} \left( k \right)}{\left( -k\tau \right)^{\alpha}}, \tag{26}
\]

where \( 0 < \alpha < 2 \). It was also shown in [1] that with \( 1 < \alpha < 2 \) the CSL-modified scalar dynamics successfully explains the quantum-to-classical transition of primordial scalar modes without destroying the phase coherence of the superhorizon modes essential to explain the peaks and troughs of the CMBR anisotropy spectrum.

Now, let us calculate the power spectrum in this scenario. In such a case it has been calculated in [1] that

\[
\frac{a^{2} \text{Re} \Omega_{k}}{H^{2}} = \frac{k^{3} \gamma_{0} \left( k \right)}{H^{2} \left( -k\tau \right)^{-1 + \alpha}}, \tag{27}
\]

which indicates that the power spectrum would not be time-invariant on superhorizon scale unlike the standard scenario. This feature has been observed in both [1,12] and as a cure to it the power was calculated at the end of inflation and not at the horizon-crossing of each mode by using

\[
- k\tau = \frac{k}{k_{0}} e^{-\Delta N}, \tag{28}
\]

where \( k_{0} \) is the comoving wavenumber of the mode which is at the horizon today \( k_{0} = a_{0}H_{0} \) and \( \Delta N \) is the number of e-folds the mode has spent outside the horizon after its exit and thus for observationally relevant modes \( \Delta N \sim 50 - 60 \). We would henceforth consider \( \Delta N \sim 60 \). Also to make the power spectrum nearly scale-invariant we chose the scale-dependence of \( \gamma_{0} \) as

\[
\gamma_{0} \left( k \right) = \tilde{\gamma}_{0} \left( \frac{k}{k_{0}} \right)^{\beta}. \tag{29}
\]

Then the power spectrum of comoving curvature perturbations turns out to be

\[
P_{R} = \frac{1}{8\pi G_{\text{pl}}} \frac{k_{s}^{3} H^{2}}{\tilde{\gamma}_{0}} e^{-\left( 1 + \alpha \right) \Delta N} \left( \frac{k_{s}}{k} \right)^{3 + \alpha - \beta} \left( \frac{k}{k_{s}} \right)^{3 + \alpha - \beta} \equiv A_{s} \left( k_{s} \right) \left( \frac{k}{k_{s}} \right)^{3 + \alpha - \beta}. \tag{30}
\]

Incorporating the correction to the tilt in the spectrum due to quasi-de Sitter evolution of the background, the power spectrum gets modified to

\[
P_{R} = A_{s} \left( k_{s} \right) \left( \frac{k}{k_{s}} \right)^{3 + \alpha - \beta + 2\eta - 4\epsilon} = A_{s} \left( k_{s} \right) \left( \frac{k}{k_{s}} \right)^{n_{s} - 1}. \tag{31}
\]

where the power is now calculated at the end of inflation.

### B. Tensor perturbations

The tensor perturbations would now be straightforward to calculate once the scalar analysis is done as each of the helicity components of the Fourier tensor mode behaves like massless scalar perturbations, as discussed before. It should be noted that tensor perturbations are gauge-invariant by construction and the redefined tensor perturbations \( v_{\mu}^{\nu} \) of Eq. (31) can be identified as the Mukhanov-Sasaki variable of scalar perturbations defined in the previous section. Hence in the Schrödinger picture
each helicity component $v^s_k$ of the tensor modes can be expressed as functionals given in Eq. (17) following similar functional Schrödinger equation given in Eq. (19) in the standard scenario. The ground state solutions of the functional Schrödinger equation would also be Gaussian as given in Eq. (20). Similarly the power spectrum of $v^s_k$ would be same as of $\zeta$ given in Eq. (22) and can be written as

$$P_{v^s}(k) = \frac{k^3}{2\pi^2 \Re \Omega_k}, \quad (32)$$

and following Eq. (3) one can write down the power spectrum for the tensor modes as

$$P_h = \sum_s \frac{4}{a^2 M^2_{Pl}} P_{v^s} = \frac{2}{\pi^2 M^2_{Pl} a^2 \Re \Omega_k} k^3. \quad (33)$$

At this point, we assume CSL collapse mechanism affects each helicity mode of the gravitons the same way as it affects the inflatons. This is the simplest scenario to imagine and is in conformity with the philosophy that the collapse mechanism should be universal in nature. Since in our way of implementing the collapse mechanism we essentially use the fact that each mode is an independent harmonic oscillator to modify the equation of motion, there is no reason to expect that this modification should be sensitive to details of the underlying nature of the field. Hence the CSL-modified dynamics of each helicity mode of the gravitons would be same as that of the massless inflatons or the gauge-invariant Mukhanov-Sasaki variable which has been considered in the previous section. For redefined tensor modes $v_s$, $\Re \Omega_k$ would have the same form as given in Eq. (27), following which the power spectrum of the tensor modes can be determined as

$$P_h = \frac{2}{\pi^2 M^2_{Pl}} \frac{k^3 H^2}{\tilde{\gamma}_0} e^{-(1+\alpha)\Delta N} \left(\frac{k}{k_0}\right)^{3+\alpha-\beta} \left(\frac{k}{k_*}\right)^{3+\alpha-\beta} \equiv A_T(k_*), \quad (34)$$

and considering the tilt in the power due to quasi-de Sitter background evolution one gets

$$P_h = A_T(k_*) \left(\frac{k}{k_*}\right)^{3+\alpha-\beta-2\epsilon} = A_T(k_*) \left(\frac{k}{k_*}\right)^{n_T}. \quad (35)$$

C. Observables

Let us now illustrate how the CSL-modified primordial dynamics differ from the standard one observationally. The first thing to note from Eq. (30) and Eq. (34) is that the tensor-to-scalar ratio remain the same in the modified dynamics:

$$r = 16\epsilon, \quad (36)$$

which indicates from the Lyth bound that the field excursion during inflation would be super-Planckian even in this case. The scalar spectral index and the tensor spectral index would now become (from Eq. (31) and Eq. (35))

$$n_s - 1 = \delta + 2\eta - 4\epsilon, \quad n_T = \delta - 2\epsilon, \quad (37)$$

where we have defined $\delta = 3 + \alpha - \beta$. We note here that, although the model has three free parameters $\tilde{\gamma}_0$, $\alpha$ and $\beta$ to begin with, the modification to spectral indices can be captured in one effective free parameter $\delta$. The observation of the scalar spectral index by PLANCK indicates that the quantity $\delta$ can at best be of the order of slow-roll parameters so that the comoving curvature power spectrum remains to be nearly scale-invariant.

We also note from Eq. (31) that the tensor amplitude in such a case does not remain to be a sole function of Hubble parameter, the prime feature which is used to determine the scale of inflation using the BICEP2 observations. Even though, if we consider that inflation has indeed taken place at that high scale, then that would help estimating the collapse parameter as

$$\tilde{\gamma}_0 \sim k^2_{T_0} e^{-(1+\alpha)\Delta N}, \quad (38)$$

which turns out to be extremely small. A stronger collapse mechanism can then bring down the scale of inflation even though the field excursions would remain to be super-Planckian.

Most interestingly what this modified dynamics does is to change the consistency relation of the single-field models. In such a scenario the consistency relation turns out to be

$$r = -8n_T + 8\delta. \quad (39)$$

Hence independent accurate measurements of $r$ and $n_T$ would give us a direct handle on $\delta$ in this model.

IV. DISCUSSION AND CONCLUSION

Belief in quantum nature of primordial perturbations makes it essential to understand the apparent classicality of inhomogeneity it gives rise to. The squeezing of the superhorizon modes and decoherence partially explain this conundrum. But the single outcome problem cannot be addressed without appealing to either many-worlds interpretation or collapse mechanisms. Collapse mechanisms generically modify the dynamics of the primordial perturbations and so are expected to leave their imprints on cosmological observables. In [1, 15] some toy models of CSL applied to inflationary dynamics were explored. In [1], within the context of an illustrative example, it was shown that classicality of perturbations can be achieved while still preserving scale-invariance and phase coherence for a certain parameter range. In the light of BICEP results which report high enough tensor-to-scalar
ratio, inconsistent with purely classical dynamics of the primordial tensor modes [7–10], we revisit this model to investigate its observational consequences.

The first point to note is that in this illustrative example the tensor-to-scalar ratio remains unchanged indicating super-Planckian field excursions during inflation as demanded by Lyth bound [21]. Secondly, the tensor amplitude no longer directly yields the scale of inflation as the collapse parameter also enters in this conversion. A stronger collapse parameter would bring down the scale of inflation. Unfortunately, so far there is no other theoretical or observational guide to estimate the scale of collapse parameter. Furthermore, the spectral tilts get modified by the one and the same combination of the free parameters of the model which eventually also change the consistency relation of the standard single-field scenario by capturing the deviation from this in one effective free parameter $\delta$. From PLANCK's observation of the scalar spectral index, the free parameter $\delta$ can at best be of the order of slow-roll parameters. It is also in principle possible that $\delta$ be identically zero in which case the collapse mechanism achieves the required classicality without leaving any imprint on observations. All the same, it is more reasonable to expect $\delta$ to be non-zero which then would make this model testable by observations. One important feature which distinguishes this kind of modification from other scenarios which also modify the consistency relation like curvaton [25] or multifield [26] models is that generically such extensions of the minimal model modify the scalar sector while leaving the tensor sector untouched, however the collapse models modify both and hence are potentially distinguishable by precision measurement of $n_T$. However modification to initial conditions of tensor modes, i.e. deviation from Bunch-Davies vacuum would also modify the tensor spectral tilt [27] reflecting the scale-dependence of Bogoliubov coefficients. It is possible to arrange this scale-dependence in such a way as to shift the tensor spectral index by a constant, thus mimicking the effect of the collapse mechanism considered here. But in principle, the generic scenario of non-Bunch Davies initial condition should be distinguishable from the collapse scenario. This would require precision measurement of running of tensor spectral index which has been argued to be possible in the near future [28].

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[1] S. Das, K. Lochan, S. Sahu, and T. P. Singh, Phys.Rev. D88, 085020 (2013), 1304.5994.
[2] P. Ade et al. (BICEP2 Collaboration) (2014), 1403.3985.
[3] A. H. Guth, Phys.Rev. D23, 347 (1981).
[4] A. D. Linde, Phys.Lett. B129, 177 (1983).
[5] G. Hinshaw et al. (WMAP), Astrophys.J.Suppl. 208, 19 (2013), 1212.5226.
[6] P. Ade et al. (Planck Collaboration) (2013), 1303.5076.
[7] A. Ashoorioon, P. B. Dev, and A. Mazumdar (2012), 1211.4678.
[8] S. Bose and L. Grishchuk, Phys.Rev. D66, 043529 (2002), gr-qc/0111064.
[9] L. M. Krauss and F. Wilczek, Phys.Rev. D89, 047501 (2014), 1309.5343.
[10] L. M. Krauss and F. Wilczek (2014), 1404.0634.
[11] L. Senatore, E. Silverstein, and M. Zaldarriaga (2011), 1109.0542.
[12] J. L. Cook and L. Sorbo, Phys.Rev. D85, 023534 (2012), 1109.0022.
[13] D. Polarski and A. A. Starobinsky, Class.Quant.Grav. 13, 377 (1996), gr-qc/9504030.
[14] A. Albrecht, P. Ferreira, M. Joyce, and T. Prokopec, Phys.Rev. D50, 4807 (1994), astro-ph/9303001.
[15] J. Martin, V. Vennin, and P. Peter, Phys.Rev. D86, 103524 (2012), 1207.2086.
[16] A. Bassi, K. Lochan, S. Satin, T. P. Singh, and H. Ulbricht, Rev.Mod.Phys. 85, 471 (2013), 1204.4325.
[17] P. Canate, P. Pearle, and D. Sudarsky, Phys.Rev. D87, 104024 (2013).
[18] D. Baumann (2009), 0907.5424.
[19] L. P. Grishchuk, Sov.Phys.JETP 40, 409 (1975).
[20] C. M. Ho and S. D. H. Hsu (2014), 1404.0745.
[21] D. H. Lyth, Phys.Rev.Lett. 78, 1861 (1997), hep-ph/9606387.
[22] S. Choudhury and A. Mazumdar (2014), 1403.5549.
[23] S. Antusch and D. Nolde (2014), 1404.1821.
[24] L. Sriramkumar and T. Padmanabhan, Phys.Rev. D71, 103512 (2005), gr-qc/0408034.
[25] T. Fujita, M. Kawasaki, and S. Yokoyama (2014), 1404.0951.
[26] S. A. Kim and A. R. Liddle, Phys.Rev. D74, 023513 (2006), astro-ph/0605604.
[27] A. Ashoorioon, K. Dimopoulos, M. Sheikh-Jabbari, and G. Shiu (2014), 1403.6099.
[28] J. Caliguri and A. Kosowsky (2014), 1403.5324.