

Inflation without an inflaton

Daniele Bertacca^{1,2,3}, Raul Jimenez^{4,5}, Sabino Matarrese^{1,2,3,6} and Angelo Ricciardone^{7,8}¹*Dipartimento di Fisica e Astronomia Galileo Galilei, Università degli Studi di Padova, via Marzolo 8, I-35131, Padova, Italy*²*INFN, Sezione di Padova, via Marzolo 8, I-35131, Padova, Italy*³*INAF-Osservatorio Astronomico di Padova, Vicolo dell'Osservatorio 5, I-35122 Padova, Italy*⁴*ICC, University of Barcelona, Martí i Franquès, 1, E08028 Barcelona, Spain*⁵*ICREA, Pg. Lluís Companys 23, Barcelona, 08010, Spain*⁶*Gran Sasso Science Institute, Viale F. Crispi 7, I-67100 L'Aquila, Italy*⁷*Dipartimento di Fisica "Enrico Fermi", Università di Pisa, Pisa I-56127, Italy*⁸*INFN sezione di Pisa, Pisa I-56127, Italy*

(Received 21 December 2024; accepted 11 June 2025; published 8 July 2025)

We propose a novel scenario in which scalar perturbations, which seed the large-scale structure of the universe, are generated without relying on a scalar field (the inflaton). In this framework, inflation is driven by a de Sitter space time, where tensor metric fluctuations (i.e., gravitational waves) naturally arise from quantum vacuum oscillations, and scalar fluctuations are generated via second-order tensor effects. We compute the power spectrum of such scalar fluctuations and show it to be consistent with near scale invariance. We derive the necessary conditions under which scalar perturbations become significant and much larger than the tensor modes, and we identify a natural mechanism to end inflation via a transition to a radiation-dominated phase. Our proposed mechanism could remove the need for a model-dependent scenario: the choice of a scalar field, as the inflaton, to drive inflation.

DOI: [10.1103/vfny-pgc2](https://doi.org/10.1103/vfny-pgc2)

Introduction. The inflationary paradigm [1–9] remains the most successful candidate to explain the origin and evolution of the universe and its large-scale structure [4]. This has led to the search among fundamental particle physics for an inflaton scalar field. So far there has not been any definitive theoretical argument that leads to a single inflaton candidate; in fact, it is possible to build any inflationary potential, even in fundamental theories, such as string theory, that fits the observational data: namely the observed tilt of the spectrum of scalar fluctuations, the upper limit on the energy scale of inflation (the yet to be discovered tensor modes) [10,11], and the non-Gaussian $1 - n_s$ signature (e.g., Ref. [12]). This is perhaps the main weakness of the inflation paradigm, as it depends on a model-dependent construction. It is therefore interesting to search for scenarios that are fully model independent, so that inflation becomes a theory rather than a model. Along this line of reasoning, an attempt is done in Refs. [13–15], in a framework termed quantum Fisher cosmology (QFC). It has been shown that the tilt of the primordial scalar power spectrum can be predicted to be $n_s = 0.9672$ just by considering the Heisenberg uncertainty principle in measuring time in de Sitter spacetime (dS).

Here, we focus on how scalar perturbations are generated in a model-independent fashion, within a purely quantum physics framework. We show that scalar perturbations arise

as a second-order effect from tensor perturbations and can become significantly enhanced, allowing them to dominate over the linear tensor modes, which are inherently present in dS. The generation of second-order scalar modes from tensor perturbations was first studied in Refs. [16–18]. Recently, a quantitative analysis of such tensor-induced scalar perturbations was done in Refs. [19,20] for postinflationary epochs. Our scenario relies on a similar mechanism to generate the scalar perturbations. In addition, we discuss the instability of the dS space [21–25], which provides both a natural way for a graceful exit from inflation and also the means to end into a radiation-dominated epoch.

If one describes dS as a quantum coherent state composite of gravitons, then the self-coupling of gravitons—as well as their coupling to other relativistic particle species, such as those in the standard model (SM), which must always be present—leads to quantum scattering and decay of the constituent gravitons of dS. In this case, the final quantum state cannot be described as a coherent state, and there will no longer be the dispersion relations of the free quanta propagating on a classical dS background.

In fact, as pointed out in Ref. [22], based on particle creation [21] and the fluctuation-dissipation theorem [26] or, equivalently, on thermodynamic considerations [27], the dS spacetime is unstable and the timescale of this instability can be exponentially large given any initial perturbation. This occurs because the geometry of the gravitational field is coupled to the energy-momentum tensor and this quantity, in turn, governs the dynamics of the gravitational field. Consequently, the quantum fluctuations of the vacuum or, equivalently, the effects of particle creation are linked to a background

gravitational field. If we use thermodynamic considerations, see Ref. [27], we note that the existence of such a maximally symmetric state cannot guarantee stability against small fluctuations.

Recent works [28,29], using an approach related to quantum thermodynamics, provide a natural mechanism for the irreversible relaxation of the cosmological constant, and dS decay, in order to escape the inflationary epoch, without the need for subsequent reheating, and suggest an alternative way to generate the primordial perturbations, which could arise from the thermal fluctuations described by the power spectrum (see also Ref. [30]).

In this Letter, we derive exact expressions for the second-order scalar potentials and the scalar power spectrum resulting from second-order tensor perturbations. We demonstrate that the latter agrees with the expected nearly scale-invariance from observations, opening the way for numerous potential follow-up studies and extensions.

Scalar perturbations from tensor modes. We work on a pure dS metric $ds^2 = -dt^2 + e^{2t/\alpha}(dx^2 + dy^2 + dz^2)$ where $\alpha \equiv (3/\Lambda)^{1/2}$ and Λ is the vacuum energy. We assume Einstein gravity.

For generality, on the right-hand side (RHS) of Einstein's equations, we allow for the presence of a stress-energy tensor, which accounts for the sum of the cosmological constant driving our dS expansion plus a generic fluid, with energy density ρ , isotropic pressure p , four-velocity u^μ and anisotropic stress tensor π^μ_ν , namely

$$T^\mu_\nu = -\frac{\Lambda}{8\pi G}\delta^\mu_\nu + (\rho + p)u^\mu u_\nu + p\delta^\mu_\nu + \pi^\mu_\nu. \quad (1)$$

As we will see below, the considered fluid unavoidably arises from the vacuum expectation value of the second-order contribution to the Einstein's tensor from gravitational waves (GW), which on subhorizon scales leads to nonvanishing energy, pressure and anisotropic stress. We consider a background metric perturbed up to second order in the Poisson gauge (see, e.g., Ref. [18])

$$ds^2 = a^2(\eta)[-(1 + \psi_2(\mathbf{x}, \eta))d\eta^2 + \omega_{2i}(\mathbf{x}, \eta)d\eta dx^i + \gamma_{ij}(\mathbf{x}, \eta)dx^i dx^j], \quad (2)$$

where η is the conformal time, $a(\eta)$ the scale factor and the metric perturbations contain only tensor modes at first order and scalar, vector and tensor ones at second order. Here, ω_{2i} is a solenoidal (divergence-free or transverse) vector while the conformal spatial metric is $\gamma_{ij} = (1 - \phi_2)\delta_{ij} + \chi_{1ij} + \chi_{2ij}/2$ where χ_{1ij} and χ_{2ij} are the linear and second-order tensor modes (GW). In this case, considering only the tensor contribution we have, in the comoving frame of such a GW fluid,

$$u^\mu = \frac{1}{a}\left(\delta^\mu_0 + \frac{1}{2}v^\mu_2\right), \quad (3)$$

from which we can compute u_μ as:

$$u_0 = g_{0\rho}u^\rho = a(-1 - \frac{1}{2}\psi_2); \quad u_i = \frac{1}{2}a\omega_{2i} + a(\frac{1}{2}v_{2i}). \quad (4)$$

For the 00 component at second order we find

$$\frac{1}{2}\delta^2 T^0_0 = -\frac{1}{2}\delta^2 \rho. \quad (5)$$

The other components of T^μ_ν are given by

$$\frac{1}{2}\delta^2 T^i_0 = -(1 + w)\bar{\rho}\left[\frac{1}{2}v^i_2\right], \quad (6)$$

$$\frac{1}{2}\delta^2 T^i_j = +\frac{1}{2}\delta^2 p\delta^i_j + \frac{1}{a^2}\frac{1}{2}\delta^2 \pi^i_j, \quad (7)$$

where w is the equation of state at the background level and the anisotropic stress tensor can be decomposed into a trace-free scalar part, Π , a vector part, Π_i , and a tensor part, Π_{ij} , at each order according to

$$\pi_{ij} = a^2\left[\Pi_{,ij} - \frac{1}{3}\nabla^2\Pi\delta_{ij} + \frac{1}{2}(\Pi_{,i,j} + \Pi_{,j,i}) + \Pi_{ij}\right], \quad (8)$$

$$\frac{1}{a^2}\delta^2 \pi_{ij} = \left[\Pi_{2,i,j} - \frac{1}{3}\nabla^2\Pi_2\delta_{ij} + \frac{1}{3}(\Pi_{2,i,j} + \Pi_{2,j,i}) + \Pi_{2ij}\right]. \quad (9)$$

We can split the peculiar velocity v^i_2 as $v_{2i} = v_{2,i} + v_{2i}^\perp$, where $v_{2i}^\perp$ is a solenoidal (divergence-free or transverse) vector.

Considering a graviton perturbation, with generic w and c_s , where $\delta^2 p = c_s^2 \delta^2 \rho$ and manipulating the Einstein equations, we find the potential equation for ϕ_2 , i.e.,

$$\begin{aligned} & \phi_2'' + 3(1 + c_s^2)\mathcal{H}\phi_2' + [2\mathcal{H}' + (1 + 3c_s^2)\mathcal{H}^2]\phi_2 - c_s^2\nabla^2\phi_2 \\ &= 8\pi Ga^2[2\mathcal{H}' + 3(1 + c_s^2)\mathcal{H}^2]\Pi_2 + \frac{8\pi G}{3}a^2\nabla^2\Pi_2 + 8\pi Ga^2\mathcal{H}\Pi_2 - \frac{3}{2}[2\mathcal{H}' + (1 + 3c_s^2)\mathcal{H}^2]\nabla^{-4}\partial_i\partial^j\mathcal{A}^i_j \\ & - \frac{3}{2}\mathcal{H}\nabla^{-4}\partial_i\partial^j\mathcal{A}^{i'}_{j'} - \frac{1}{2}\nabla^{-2}\partial_i\partial^j\mathcal{A}^i_j - [2\mathcal{H}' + (1 + 3c_s^2)\mathcal{H}^2]\nabla^{-2} \\ & \times \left[-\frac{3}{4}\chi_1^{lk,m}\chi_{1kl,m} - \frac{1}{2}\chi_1^{kl}\nabla^2\chi_{1lk} + \frac{1}{2}\chi_1^{km}\chi_{1m,k} + \frac{1}{2}\chi'_{1km}\chi'_{1km} - \frac{3}{2}\nabla^{-2}\partial_i\partial_j(\chi_{1k}^{i'}\chi_{1kj}^{j'}) \right] \\ & + \frac{3}{8}\chi_{1,k}^{ml}\chi_{1lm}^k - \frac{3}{8}\chi_1^{kl'}\chi'_{1kl} - \frac{1}{4}\chi_{1k,l}^m\chi_{1,m}^{lk} - \mathcal{H}\nabla^{-2}\left[-\frac{3}{2}\chi_1^{lk,m'}\chi_{1kl,m} + \frac{1}{2}\chi_1^{kl'}\nabla^2\chi_{1lk} - \frac{1}{2}\chi_1^{kl}\nabla^2\chi_{1lk}' + \frac{1}{2}\chi_{1,l}^{km'}\chi_{1m,k}^l \right. \\ & \left. + \frac{1}{2}\chi_{1,l}^{km}\chi_{1m,k}' - 2\mathcal{H}\chi_{1km}'\chi'_{1km} \right] + 3\mathcal{H}\nabla^{-4}(-2\mathcal{H}\chi_{1k,j}^{i'}\chi_{1k,i}^{j'} + \chi_{1k,i}^{j'}\nabla^2\chi_{1k,j}^{i'}) + \frac{1}{2}\nabla^{-2}\partial_i\partial_j(\chi_{1k}^{i'}\chi_{1kj}^{j'}) \\ & + c_s^2\left[-\frac{1}{8}\chi_1^{ik'}\chi'_{1ki} - \mathcal{H}\chi_1^{ik}\chi'_{1ki} + \frac{1}{2}\chi_1^{mk}\nabla^2\chi_{1mk} + \frac{3}{8}\chi_{1,k}^{ml}\chi_{1lm}^k - \frac{1}{4}\chi_{1k,l}^m\chi_{1,m}^{lk} \right], \end{aligned} \quad (10)$$

where

$$\mathcal{A}_j^i = \frac{1}{2} \chi_1^{lk,i} \chi_{1kl,j} + \chi_1^{kl} \chi_{1lk,j}^i - \chi_1^{kl} \chi_{1l,jk}^i - \chi_1^{kl} \chi_{1lj,k}^i + \chi_1^{kl} \chi_{1j,kl}^i + \chi_1^{ik} \chi_{1l,jk}^l - \chi_{1,l}^{ki} \chi_{1j,k}^l. \quad (11)$$

Let us highlight that the master Eq. (10) is one of the main results of the paper and describes the dynamic of the gravitational potential ϕ_2 on all scales. Regarding ψ_2 , using the traceless part of the ij components we find

$$\psi_2 = \phi_2 - 8\pi G a^2 \Pi_2 - \frac{\mathcal{F}_\chi}{4}, \quad (12)$$

where

$$\mathcal{F}_\chi = 4\nabla^{-2} \left(\frac{3}{4} \chi_1^{lk,m} \chi_{1kl,m} + \frac{1}{2} \chi_1^{kl} \nabla^2 \chi_{1lk} - \frac{1}{2} \chi_{1,l}^{km} \chi_{1lm,k}^l \right) - 6\nabla^{-4} \partial_i \partial^j \mathcal{A}_j^i. \quad (13)$$

Considering only the tensor contribution, the gauge-invariant curvature perturbation on uniform density hypersurfaces is

$$\zeta_2 = -\phi_2 - \frac{\mathcal{H}}{\bar{\rho}'} \delta^2 \rho. \quad (14)$$

As a first approximation, let's assume that ϕ_2 is constant on very large scales. In this case looking for only the particular solution for ϕ_2 and using the energy-momentum constraint equation we get

$$\frac{8\pi G}{H_\Lambda^2} \Pi_2 = \eta_{\text{in}}^2 \left[\frac{8\pi G}{\eta_{\text{in}}^2 H_\Lambda^2} \Pi_{2\text{in}} - \left(\phi_2 - \frac{1}{4} \mathcal{F}_\chi \right) \right] \left(\frac{\eta}{\eta_{\text{in}}} \right)^{(5+3c_s^2)} + \left(\phi_2 - \frac{1}{4} \mathcal{F}_\chi \right) \eta^2 \quad (15)$$

$$\psi_2 = \left[-\frac{8\pi G}{\eta_{\text{in}}^2 H_\Lambda^2} \Pi_{2\text{in}} + \left(\phi_2 - \frac{1}{4} \mathcal{F}_\chi \right) \right] \left(\frac{\eta}{\eta_{\text{in}}} \right)^{3(1+c_s^2)}. \quad (16)$$

We can quickly verify that the previous two equations can be combined again to recover the usual relation

$$\frac{8\pi G}{\eta_{\text{in}}^2 H_\Lambda^2} \Pi_2 + \psi_2 \left(\frac{\eta}{\eta_{\text{in}}} \right)^2 = \left(\phi_2 - \frac{1}{4} \mathcal{F}_\chi \right) \left(\frac{\eta}{\eta_{\text{in}}} \right)^2. \quad (17)$$

Note that if

$$\left(\phi_2 - \frac{1}{4} \mathcal{F}_\chi \right) = \frac{8\pi G}{\eta_{\text{in}}^2 H_\Lambda^2} \Pi_{2\text{in}}, \quad (18)$$

we will always have $\psi_2 = 0$ and $\phi_2 = \mathcal{F}_\chi/4$; this could also be a solution in our scenario. Here, η_{in} is a suitable initial condition. Precisely, η_{in} is that initial moment in which the hypothesis that ϕ_2 is constant over time is still correct. Now, considering just a mode with a given comoving k , we know that this assumption is reasonably correct just after horizon crossing. η_{in} must depend on comoving k (i.e., for different k we should have a suitable η_{in}). In this case, as a first approximation, we can set $|\eta_{\text{in}}(k)| \simeq 1/c_s k$. If we Fourier transform all these variables we obtain

$$\tilde{\psi}_2 \simeq \left[-\frac{8\pi G}{H_\Lambda^2} c_s^2 k^2 \tilde{\Pi}_{2\text{in}} + \left(\tilde{\phi}_2 - \frac{1}{4} \tilde{\mathcal{F}}_\chi \right) \right] (c_s k |\eta|)^{3(1+c_s^2)}. \quad (19)$$

Now, on superhorizon scales, still assuming ϕ_2 constant, ζ_2 becomes

$$\zeta_2 = -\phi_2 + \frac{3\mathcal{H}^3}{4\pi G a^2 \bar{\rho}'} \psi_2 = -\phi_2 - \frac{H_\Lambda^2}{4\pi G(1+w)\bar{\rho}} \psi_2, \quad (20)$$

where we left w generic and, let us stress, that is different from c_s^2 . Including Eq. (19) and

$$\bar{\rho}(\eta) = \bar{\rho}_{\text{in}} \left(\frac{\eta}{\eta_{\text{in}}} \right)^{3(1+w)} = \bar{\rho}_{\text{in}} (c_s k |\eta|)^{3(1+w)}$$

in Eq. (20), it becomes

$$\begin{aligned} \tilde{\zeta}_2 - \tilde{\phi}_2 &= -\frac{H_\Lambda^2}{4\pi G(1+w)\bar{\rho}} \tilde{\psi}_2 \\ &= -\frac{H_\Lambda^2}{4\pi G(1+w)\bar{\rho}_{\text{in}}} \left[-\frac{8\pi G}{H_\Lambda^2} c_s^2 k^2 \tilde{\Pi}_{2\text{in}} + \left(\tilde{\phi}_2 - \frac{1}{4} \tilde{\mathcal{F}}_\chi \right) \right] (c_s k |\eta|)^{3(c_s^2-w)}. \end{aligned} \quad (21)$$

Now we need to know the value of c_s and ρ_{in} at a given mode k . It is transparent that $w - c_s^2 > 0$ in order for the scalar fluctuations to be larger than the tensor ones. This can be achieved by the same gravitons produced during de Sitter phase [31]. For pure radiation with equation of state 1/3, the scalar perturbations will no longer grow. This provides a natural route to end inflation.

We conclude this section with a further comment. As can be seen from the metrics we have considered, with our approach we can analytically study vectors and tensors at the second order. These contributions will also be analyzed in detail in future work.

Power spectrum. As a simple analytical example, let us consider the case where $\psi_2 = 0$, i.e., when $\phi_2 = \mathcal{F}_\chi/4$ and we focus only on the first terms in (13); the second term will give a similar dependence. Then the primordial power spectrum of scalar fluctuations ϕ is

$$\langle \phi(\mathbf{k})\phi(\mathbf{k}') \rangle = \frac{1}{4} \langle \phi_2(\mathbf{k})\phi_2(\mathbf{k}') \rangle = (2\pi)^3 \delta^{(3)}(\mathbf{k} + \mathbf{k}') \mathcal{P}_\phi(k) \quad (22)$$

and

$$\mathcal{P}_\phi(k) = \frac{1}{64(2\pi)^3} \frac{1}{k^4} \int d^3k_1 d^3k_2 \delta^{(3)}[\mathbf{k} - (\mathbf{k}_1 + \mathbf{k}_2)] \mathcal{K}_h(\mathbf{k}_1, \mathbf{k}_2, k^2) P_h(k_1) P_h(k_2), \quad (23)$$

where the kernel is given by

$$\begin{aligned} \mathcal{K}_h(\mathbf{k}_1, \mathbf{k}_2, k^2) = & \{ (k_1^2 + k_2^2 + 3\mathbf{k}_1 \cdot \mathbf{k}_2)^2 [(1 - \hat{\mathbf{k}}_1 \cdot \hat{\mathbf{k}}_2)^4 + (1 + \hat{\mathbf{k}}_1 \cdot \hat{\mathbf{k}}_2)^4] \\ & + 8(\mathbf{k}_1 \cdot \mathbf{k}_2)(k_1^2 + k_2^2 + 3\mathbf{k}_1 \cdot \mathbf{k}_2)(\hat{\mathbf{k}}_1 \times \hat{\mathbf{k}}_2)^2 [3 + (\hat{\mathbf{k}}_1 \cdot \hat{\mathbf{k}}_2)^2] + 8k_1^2 k_2^2 (\hat{\mathbf{k}}_1 \times \hat{\mathbf{k}}_2)^4 [1 + (\hat{\mathbf{k}}_1 \cdot \hat{\mathbf{k}}_2)^2] \}. \end{aligned} \quad (24)$$

Defining $P_h(k) = 2\pi^2 \Delta_h^2(k)/k^3$, the tensor power spectrum in dS reads [32]

$$\Delta_h^2(k) = \frac{16}{\pi} \left(\frac{H_{\text{inf}}}{m_{\text{pl}}} \right)^2, \quad (25)$$

where $H_{\text{inf}} = \sqrt{\Lambda/3}$ is the Hubble constant during inflation.

Note that for the scalar perturbations

$$\Delta_\phi(k) = \frac{k^3}{2\pi^2} \frac{1}{64(2\pi)^3} \frac{1}{k^4} \int d^3k_1 d^3k_2 \delta^{(3)}[\mathbf{k} - (\mathbf{k}_1 + \mathbf{k}_2)] \mathcal{K}_h(\mathbf{k}_1, \mathbf{k}_2) \left[\frac{2\pi^2}{k_1^3} \frac{16}{\pi} \left(\frac{H_{\text{inf}}}{m_{\text{pl}}} \right)^2 \right] \left[\frac{2\pi^2}{k_2^3} \frac{16}{\pi} \left(\frac{H_{\text{inf}}}{m_{\text{pl}}} \right)^2 \right], \quad (26)$$

which shows that the CMB measured level of fluctuations determines the scale of inflation. In particular, performing the above integral yields $H_{\text{inf}} \sim 5 \times 10^{-4} m_{\text{pl}}$; this is not far from the GUT scale. This corresponds to a tensor-to-scalar ratio of $r \sim 0.0006$, which lies at the detection threshold of Stage IV CMB experiments [33], but distinctly different from predictions of models such as Starobinsky's inflation.

Note that the kernel differs distinctly from that of standard single-clock inflation (which for this case is in the bispectrum). This difference will also affect the bispectrum of scalar perturbations, thereby imprinting a unique signature in the squeezed non-Gaussian features of the large-scale structure [12]. As a consequence, this provides a means to observationally test our scenario. The power spectrum is nearly scale invariant; this can be seen easily as no uncompensated power of 4 terms of $\mathbf{k}$ appear in the scaling of the momenta in (24). This raises the issue of how we obtain the required red tilt observed by CMB experiments. To obtain this we have to apply the argument developed in Refs. [13–15], where the idea is that in any de Sitter spacetime, quantum uncertainty will always generate a (red) tilt of the scale-invariant power spectrum.

Starting from Eq. (10) we note that at smaller scales, more precisely starting from the horizon $1/H_\Lambda$, the first-order tensor modes give a fundamental contribution because they are the sources (via the particular solution for the ϕ_2). Consequently they give a contribution to the scalars ϕ_2 and ψ_2 that we are studying at second order. At even smaller scales, i.e., for $k\eta \gg 1$ we have $\chi_{1ij} \propto e^{\pm k\eta}/a$; therefore we infer that there will be a free streaming effect (the tensor modes will oscillate and decay) and at even smaller scales we have to treat them as quantized gravitons that are continuously generated. In other

words, on scales smaller than the horizon we can treat them as Bose-Einstein graviton condensates [24,25]. In this case an open quantum system approach is needed in a quantum thermodynamic system and this will be studied and explored in further work.

Discussion. We have shown that it is possible to generate nearly scale-invariant scalar adiabatic perturbations in pure de Sitter. This is a scenario where the inflaton does not exist, and thus opens up the possibility to provide a picture of inflation that is model independent. There are some interesting features in our calculation. Starting from (10) we need an extra relation to close the equations and find a general solution. As we have pointed out in the previous section, a possible way is to look at the particular solution which is constant, i.e., nondynamical, on very large scales. (Let us emphasize that this is perfectly in line with the approaches used when analyzing the behavior of the gravitational potential on very large scales.) However, this is not true for ψ_2 and consequently for ζ_2 when the effective/generic fluid has $w - c_s^2 > 0$. Finally, let us highlight that, as Refs. [24,25] have also pointed out, this additional contribution (in our case this extra fluid) is necessary to regularize dS. In particular, this other kind of contribution (precisely, quantum at scales much smaller than the horizon, and classical on very large scales) necessarily also plays the role of a cosmic clock. In this case, in addition to the self-coupling of gravitons, we have their coupling to other species of relativistic particles, such as those of the SM, which must always be present, leading to quantum dispersion and decay of the gravitons constituting dS on quantum scales, far below the horizon.

First, note that the scalar fluctuations are generated outside the horizon from the tensor perturbations. Assuming these

fluctuations are adiabatic, as in the standard picture of inflation, the perturbations induced in the potential ϕ persist after the transition from the dS phase to a radiation dominated epoch through the decay of dS. Similarly to scalar modes, vector perturbations are also generated from second-order tensor fluctuations. As expected, the fluctuations deviate from Gaussianity and exhibit a certain intrinsic level of non-Gaussian features, the quantification of which will be addressed in future work. However, we note the following: for any non-conformally invariant field (with mass $m \ll H$) the quantum generation of particles from the vacuum state will add extra contributions to the second-order scalar perturbation mode, hence contributing to a Gaussianization by the central limit theorem. In our scenario, the magnitude of fluctuations can be determined from Eqs. (23) and (24), which rely exclusively on dS features and potential physical mechanisms that enhance scalar perturbations. These factors suggest a possible explanation for why the fluctuations are at the 10^{-5} level. Particle production (reheating) is generated by de Sitter decaying into radiation as postulated in Refs. [21–23,25]. This out-of-equilibrium transition to a radiation phase is what provides a replenishment of the patch with SM particles. For out-of-equilibrium decay $T_{\text{RH}} = 10^{-10}(E/\text{GeV})^{3/2} \text{ GeV}$ (see Ref. [34]), where E is the energy scale of inflation. Using the value derived from the level of CMB fluctuations above (10^{14-15} GeV) leads to $10^{11} < T_{\text{RH}}/\text{GeV} < 10^{12}$, which allows for baryogenesis. This will result in a particle density of $10^{90} < n_\chi < 10^{100}$, which is the right number of particles in our patch and entropy generated out of equilibrium between

$1 - 3 (S_f/S_i = 10^7(E/\text{GeV})^{-1/2})$, which guarantees no extra injection of entropy.

In our scenario, we assume the existence of dS. It is worth commenting if this is plausible. The first argument one can make is a phenomenological one; currently we leave in an almost de Sitter spacetime driven by dark energy. While the scale of the current de Sitter phase is meV and not relevant for inflation, our mechanism would, without doubt, apply today as the type of fluctuations we are talking about will be present. The difference is that as we are inside the horizon they are still quantum in nature. However, as they cross the horizon and classicalize there might be some observable effects. Further, if the DESI results [35] are to stand the test of time, and it is confirmed that dark energy is weakening (decaying), this will be further support that de Sitter may be decaying. On a more theoretical front, one can argue that dS is the most natural scenario in quantum gravity theories that try to describe physics beyond the Planck scale, see, e.g., Ref. [36].

Acknowledgments. Funding for R.J. was partially provided by PID2022-141125NB-I00, and the “Center of Excellence Maria de Maeztu 2020-2023” award to the ICCUB (CEX2019- 000918-M) funded by MCIN/AEI/10.13039/501100011033. D.B. and S.M. acknowledge support from the COSMOS network through ASI (Italian Space Agency) Grants No. 2016-24-H.0, No. 2016-24-H.1-2018, and No. 2020-9-HH.0.

Data availability. No data were created or analyzed in this study.

-
- [1] A. A. Starobinsky, Spectrum of relict gravitational radiation and the early state of the universe, *JETP Lett.* **30**, 682 (1979).
 - [2] A. A. Starobinsky, A new type of isotropic cosmological models without singularity, *Phys. Lett. B* **91**, 99 (1980) [Adv. Ser. Astrophys. Cosmol. **3**, 130 (1987)].
 - [3] A. A. Starobinsky, Dynamics of phase transition in the new inflationary universe scenario and generation of perturbations, *Phys. Lett. B* **117**, 175 (1982).
 - [4] V. F. Mukhanov and G. V. Chibisov, Quantum fluctuations and a nonsingular universe, *JETP Lett.* **33**, 532 (1981) [*Pisma Zh. Eksp. Teor. Fiz.* **33**, 549 (1981)].
 - [5] A. H. Guth, The inflationary universe: A possible solution to the horizon and flatness problems, *Phys. Rev. D* **23**, 347 (1981) [Adv. Ser. Astrophys. Cosmol. **3**, 139 (1987)].
 - [6] A. D. Linde, A new inflationary universe scenario: A possible solution of the horizon, flatness, homogeneity, isotropy and primordial monopole problems, Adv. Ser. Astrophys. Cosmol. **3**, 149 (1987).
 - [7] A. D. Linde, Chaotic inflation, *Phys. Lett. B* **129**, 177 (1983).
 - [8] A. Albrecht and P. J. Steinhardt, Cosmology for grand unified theories with radiatively induced symmetry breaking, *Phys. Rev. Lett.* **48**, 1220 (1982).
 - [9] L. A. Kofman, A. D. Linde and A. A. Starobinsky, Inflationary universe generated by the combined action of a scalar field and gravitational vacuum polarization, *Phys. Lett. B* **157**, 361 (1985).
 - [10] N. Aghanim *et al.* (Planck Collaboration), Planck 2018 results. I. Overview and the cosmological legacy of Planck, *Astron. Astrophys.* **641**, A1 (2020).
 - [11] G. Galloni, N. Bartolo, S. Matarrese, M. Migliaccio, A. Ricciardone and N. Vittorio, Updated constraints on amplitude and tilt of the tensor primordial spectrum, *J. Cosmol. Astropart. Phys.* **04** (2023) 062.
 - [12] D. Bertacca, R. Jimenez, S. Matarrese, and L. Verde, Breaking the single clock symmetry: Measuring single-field inflation non-Gaussian features, *Phys. Rev. D* **105**, 043503 (2022).
 - [13] C. Gomez and R. Jimenez, Cosmology from quantum information, *Phys. Rev. D* **102**, 063511 (2020).
 - [14] C. Gómez and R. Jimenez, Model independent prediction of the spectral index of primordial quantum fluctuations, *J. Cosmol. Astropart. Phys.* **10** (2021) 052.
 - [15] C. Gomez and R. Jimenez, Quantum fisher cosmology: Confronting observations and the trans-planckian problem, *J. Cosmol. Astropart. Phys.* **09** (2021) 016.
 - [16] K. Tomita, Non-linear theory of gravitational instability in the expanding universe. II, *Prog. Theor. Phys.* **47**, 416 (1972).
 - [17] K. Tomita, Non-linear theory of gravitational instability in the expanding universe. II, *Prog. Theor. Phys.* **45**, 1747 (1971).
 - [18] S. Matarrese, S. Mollerach and M. Bruni, Second order perturbations of the Einstein-de Sitter universe, *Phys. Rev. D* **58**, 043504 (1998).

- [19] P. Bari, A. Ricciardone, N. Bartolo, D. Bertacca and S. Matarrese, Signatures of primordial gravitational waves on the large-scale structure of the universe, *Phys. Rev. Lett.* **129**, 091301 (2022).
- [20] P. Bari, D. Bertacca, N. Bartolo, A. Ricciardone, S. Giardiello and S. Matarrese, An analytical study of the primordial gravitational-wave-induced contribution to the large-scale structure of the Universe, *J. Cosmol. Astropart. Phys.* **07** (2023) 034.
- [21] E. Mottola, Particle creation in de Sitter space, *Phys. Rev. D* **31**, 754 (1985).
- [22] I. Antoniadis, P. O. Mazur, and E. Mottola, Cosmological dark energy: Prospects for a dynamical theory, *New J. Phys.* **9**, 11 (2007).
- [23] A. M. Polyakov, Phase transitions and the universe, *Sov. Phys. Usp.* **25**, 187 (1982).
- [24] G. Dvali and C. Gomez, Quantum compositeness of gravity: Black holes, ads and inflation, *J. Cosmol. Astropart. Phys.* **01** (2014) 023.
- [25] G. Dvali, C. Gomez, and S. Zell, Quantum break-time of de Sitter, *J. Cosmol. Astropart. Phys.* **06** (2017) 028.
- [26] E. Mottola, A quantum fluctuation dissipation theorem for general relativity, *Phys. Rev. D* **33**, 2136 (1986).
- [27] E. Mottola, Thermodynamic instability of de Sitter space, *Phys. Rev. D* **33**, 1616 (1986).
- [28] R. Alicki, G. Barenboim, and A. Jenkins, Quantum thermodynamics of de Sitter space, *Phys. Rev. D* **108**, 123530 (2023).
- [29] R. Alicki, G. Barenboim, and A. Jenkins, The irreversible relaxation of inflation, *Phys. Lett. B* **866**, 139519 (2025).
- [30] J. Feng and J. J. Zhang, Quantum fisher information as a probe for Unruh thermality, *Phys. Lett. B* **827**, 136992 (2022).
- [31] A. Negro and S. P. Patil, Hadamard regularization of the graviton stress tensor, *arXiv:2403.16806* [hep-th].
- [32] Y. Watanabe and E. Komatsu, Improved calculation of the primordial gravitational wave spectrum in the standard model, *Phys. Rev. D* **73**, 123515 (2006).
- [33] L. Verde, H. Peiris, and R. Jimenez, Optimizing CMB polarization experiments to constrain inflationary physics, *J. Cosmol. Astropart. Phys.* **01** (2006) 019.
- [34] E. W. Kolb and M. S. Turner, The early universe, *Front. Phys.* **69**, 1 (1990).
- [35] M. Abdul Karim *et al.* (DESI Collaboration), DESI DR2 Results II: Measurements of Baryon Acoustic Oscillations and Cosmological Constraints, *arXiv:2503.14738* [astro-ph.CO].
- [36] P. Berglund, T. Hübsch and D. Minic, On de Sitter spacetime and string theory, *Int. J. Mod. Phys. D* **32**, 2330002 (2023).