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# The SSC formalism and the collapse hypothesis for inflationary origin of the seeds of cosmic structure

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## Abstract.

This article resumes recent advances in the development of a formalism that allows the incorporation of a hypothetical spontaneous reduction or collapse of wave function of matter fields within the context of semiclassical gravity. The proposal is applied to the inflationary scenario for the emergence of the primordial inhomogeneities and anisotropies out of the initially homogeneous and isotropic quantum state, to which, the early stages of inflation are supposed to drive the universe. In previous works we have argued that a scheme of this kind is required if we want to satisfactorily account for the emergence of the seeds of cosmic structure.

## 1. Introduction

Although the inflationary account for the primordial power spectrum is phenomenologically very successful, there is an obscure part in our understanding of the emergence of the seeds of cosmic structure: How does a universe<sup>1</sup>, which at one point in time, is described by a state that is fully homogeneous and isotropic (H&I), evolve into a state that is not, given that the dynamics does not contain any source for the undoing of such symmetry? The issue has been treated in [1], and further elaborated in [2]. We have been pursuing the proposal that the resolution of such problem is tied to the generic resolution of the measurement problem in quantum theory, based on the ideas about spontaneous (i.e. measurement unrelated) collapse of the wave function. The fact that the situation at hand is one where gravity plays an important role, induces us to be particularly sympathetic to the ideas of L. Diósi and R. Penrose that such collapse is somehow tied to fundamental quantum aspects of gravitation, or more generically, that the issue must have its roots in the quantum gravity interface. Previous studies have concentrated on the phenomenological implementation of such ideas and to the uncovering of the essential features that would be required for the scheme to match the now detailed observations of the cosmic microwave background (CMB) and matter distribution spectra.

This work is devoted, instead, to the formal characterization of the process of collapse within the semiclassical treatment of the problem, an issue that can have applications in dealing with

<sup>1</sup> The argument can be applied also to any relevant and causally connected part of such universe. In fact most inflationary cosmologists would say that the Bunch Davies (BD) vacuum, or a state closely related to it, characterizes only a region that, as a result of the relatively propitious initial conditions, undergoes a successful period of inflation. The other regions, being causally disconnected with the former, play no role in the discussion, and, according to the inflationary paradigm, do not, in any way, prevent the characterization of the relevant region in terms of the pure BD state.



other problems related to the quantum gravity interface.

Before entering the main discussion it seems appropriate to point out the connection of the work with several other various lines of work motivated by closely related issues. Those issues are discussed in Sec. 2 to 5 while Sec. 6 is devoted to the conceptual cosmological problem. Section 7 describes our proposed formalism and Sec. 8 illustrates its application to a baby version of the cosmological problem. We end in Sec. 8 with a brief discussion of the work, its problematic aspects and the prospects for its future development.

## 2. The measurement problem in quantum theory

The extent of the existing work on the subject [11] clearly puts it beyond what we can even begin to address here. However, we should just point out that the great majority of said work can be divided into two great lines: I) works on interpretation of the theory, and II) works based on modifications of the theory.

The basic issue is the fact that the standard presentation of quantum mechanics involves two distinct evolution processes for the quantum state of a system. In Penrose's words, these are the U-Process (characterized by the unitary and deterministic Schrödinger equation), and the R-Process (characterized by the non-deterministic change of the state into one of the eigenstates of a "measured" observable with probabilities dictated by Born's rule).

With regards to I), we will just characterize the situation concerning the Interpretation-Measurement Problem with the following two quotes:

*"Either the wave function as given by Schrödinger equation is not everything, or it is not right"* [6] in J. Bell's Are There Quantum Jumps? and *"There is now in my opinion no entirely satisfactory interpretation of quantum mechanics"* by S. Weinberg [7].

Regarding II) the main proposals trace back to the works [4], [5] and related developments [31, 32, 33], which we will generically call *Collapse Theories*. These are based on the unification of  $R$  and  $U$  evolution processes into a single evolution equation, which is supposed to hold universally. Moreover, the proposals remove the special role for the concept of measurement which becomes some particular case of generic interactions between subsystems of larger systems. In this work, we will not adhere to the early suggestions of the collapse occurring at discrete jumps, or the more modern continuous versions where the collapse is controlled by a stochastic function, but will often present the discussion in terms of the former simply because it is easier to handle mathematically. The basic aspect we will take from these proposals is that the evolution of a quantum state is controlled by a modified Schrödinger's equation that includes some spontaneous and unpredictable jumps

$$|\Psi\rangle \rightarrow |\Psi'\rangle \quad (1)$$

which are, moreover, independent of any external notion of measurement. This removes the anthropomorphization usually associated with the measurement concept, and makes the theory applicable to closed systems. The ultimate closed system being, of course, the subject of study of cosmology. The study of any regime where quantum and gravity come together becomes even more alluring given the intriguing suggestion by Diósi [8] and R. Penrose [9] that the fundamental modification of quantum theory envisioned by the collapse theories might trace its origin to the gravity/quantum interface.

**The argument:** The stability of a quantum superposition of two different stationary mass distributions is examined, where the perturbing effect of each distribution on the space-time structure is taken into account, in accordance with the principles of general relativity. It is argued that the definition of the time-translation operator for the superposed space-times involves an inherent ill-definedness, leading to an essential uncertainty in the energy of the superposed state which, in the Newtonian limit, is proportional to the gravitational self-energy  $E_\Delta$  of the

difference between the two mass distributions. This is consistent with a suggested finite lifetime of the order of  $t \approx \hbar/E_\Delta$  for the superposed state, in agreement with a certain proposal made by the author for a gravitationally induced spontaneous quantum state reduction, and with closely related earlier suggestions by Diósi and by Ghirardi et al.

This is one of the general subjects where our work could be relevant: If gravity is tied to the origin of the modifications in quantum theory envisioned by the collapse theories, it becomes of interest how to incorporate gravitation in any such scheme. One possibility, that can not be discarded, is that the whole enterprise would have to wait for a fully satisfactory and workable theory of quantum gravity. The other possibility, based on the work [3] and which we want to describe here, is to do so in terms of an effective semiclassical scheme where matter is treated quantum mechanically and gravitation is treated in terms of the space-time metric and the semiclassical Einstein's equations.

### 3. Semi-classical gravity

Semiclassical gravity is based on the notion that the energy-momentum tensor expectation value is what should play the source term in Einstein's equations when the quantum nature of matter is taken into account. Thus, one writes:

$$G_{\mu\nu} = 8\pi G \langle \hat{T}_{\mu\nu} \rangle. \quad (2)$$

The resulting theory, and its potential problems, have been considered for instance in [29], where it is argued that it is, in fact, far from clear that the theory should be regarded as being unviable. One of the works that is taken as most damaging is that of [30], describing an actual experiment involving mass distributions controlled by a quantum mechanical device capable, in principle, of leading to a macroscopic superposition of the mass distribution, a Schrödinger's cat state. The results are in contradiction with the above formula, if one assumes there is nothing like an effective collapse of the wave function. The connection with our problem lies in the question: how would we treat the situation if there was one such process?

#### 3.1. Relation to the stochastic gravity formalism

One of the criticisms often presented against the idea of semiclassical gravity is that Eq. (2) fails to take into account the so called "quantum fluctuations". Dealing with such criticisms is the objective of the stochastic gravity formalism [35].

The basic idea behind such approach is to add to the Eq. (2) a term characterizing those "quantum fluctuations":

$$G_{\mu\nu}(x) = 8\pi G (\langle \hat{T}_{\mu\nu}(x) \rangle + \xi_{\mu\nu}(x)) \quad (3)$$

where  $\xi_{\mu\nu}$  is a stochastic classical field, corresponding to a distribution with ensemble averages (which we denote by an overline, to distinguish them from the standard quantum expectation values) such that,

$$\overline{\xi_{\mu\nu}(x)} = 0, \quad \overline{\xi_{\mu\nu}(x)\xi_{\rho\sigma}(y)} = \langle \hat{T}_{\mu\nu}(x)\hat{T}_{\rho\sigma}(y) \rangle. \quad (4)$$

Those properties characterize the ensemble from where the individual representatives of the field  $\xi_{\mu\nu}(x)$  are drawn. That is, they characterize the random field representing the "quantum fluctuations". What is implicitly assumed in this approach is that the quantum uncertainties might be well characterized by these stochastic fields. The world out there would correspond to one such realization of the stochastic variable and, as the selection of the particular one that becomes realized, is not determined by the setting, we have a certain degree of unpredictability. That unpredictability would somehow correspond to the one which is often associated with the  $R$  process of standard quantum theory. Therefore, it should not be surprising that one can

relate this approach to the collapse theories, which, as we explained before, do incorporate that  $R$ -process within the modified evolution equations.

In order to see this, let us consider one of such jumps or collapses in the state of the matter fields. To simplify the discussion we will be assuming that the space-time is foliated by hyper surfaces labeled by the time function “ $t$ ” and that the quantum fields are described in the Heisenberg picture. The states would therefore, not evolve with time except in those special instances associated with a collapse<sup>2</sup>. Then collapse is described by an instantaneous jump:  $|\psi(t)\rangle = \theta(t_0 - t)|0\rangle + \theta(t - t_0)|\xi\rangle$ . Now, let us consider its gravitational effects. In this case Einstein’s semiclassical equations read:

$$G_{\mu\nu} = 8\pi G \langle \psi(t) | \hat{T}_{\mu\nu} | \psi(t) \rangle \quad (5)$$

which we can write as

$$G_{\mu\nu} = 8\pi G \langle 0 | \hat{T}_{\mu\nu} | 0 \rangle + 8\pi G \xi_{\mu\nu} \quad (6)$$

where

$$\xi_{\mu\nu} \equiv \theta(t - t_0) (\langle \xi | \hat{T}_{\mu\nu} | \xi \rangle - \langle 0 | \hat{T}_{\mu\nu} | 0 \rangle) \quad (7)$$

might be seen as corresponding to an individual stochastic step. Stochastic gravity might correspond, according to this, to a continuous version of dynamical collapses (like CSL [5]).

#### 4. Emergent gravity

There are many arguments suggesting that gravitation might be an emergent phenomenon. The arguments range from ideas based on string theory developments [36, 37] to those motivated directly by the thermodynamical characteristics of the laws of black hole mechanics [38]; see also the discussion in [39].

It has been argued that in that case the quantization of variables, directly connected with the space-time metric, would be as inappropriate as the quantization of the Navier-Stokes equations. Moreover, if one must take into account the interaction of gravitation with some matter fields described in a quantum mechanical language (as in inflationary cosmology for instance), the semiclassical treatment would be the most appropriate. However if one wants to further consider the modifications of the quantum theory associated with any sort of collapse theory, one finds that one lacks a formalism to do so.

The fact is that, while the semiclassical gravity equation (2) might be valid in the absence of quantum jumps, it is clear that it can not be valid on the jump itself. We take the view, motivated in part by the ideas about the resolution of black hole singularities in Loop Quantum Gravity [40], that during the jump the degrees of freedom of the quantum space-time are excited. In the fluid analogy, this might be thought as corresponding to some chemical reaction or phase transition occurring in the fluid. It is clear that during such processes, which generally involve energy flux between the atomic or molecular degrees of freedom to the macroscopic degrees of freedom characterized in terms of the fluid variables, the Navier-Stokes’ equations can not be valid. If, however, the phase transition takes place in a time period that is very short compared with the other temporal scales characteristic of the situation, one can assume that such effective equation is valid before and after the chemical reaction or phase transition.

<sup>2</sup> The idea is similar to the juncture conditions characterizing thin shells in general relativity [23]. If one knows the energy-momentum tensor outside and inside the shell, then Einstein’s equations can be taken to hold on both sides. At the juncture, corresponding to the infinitely thin shell’s location, one has a singular energy-momentum tensor, and the usual, smooth version of Einstein’s equations must be replaced by something else which takes into account the singular nature of such term.

## 5. Recovering space-time notions from a quantum gravity theory

The recovering of the classical notions of space-time, and Einstein's general relativity in terms of the corresponding variables from a fundamental theory of quantum gravity is far from trivial [41]. One of the most notorious difficulties is the dealing with the so called "Problem of Time" in theories of quantum gravity. The fact is that, generically, in canonical theories of quantum gravity time disappears from the framework and the theory becomes, essentially, a timeless theory [26]. The problem is then how to recover an appropriate characterization of the state of gravitational degrees of freedom, and to do so in a way that respects the standard covariant space-time notions associated with general relativity. The most promising paths in this respect involve the notion of relational time [27], i.e., the selection of a certain variable constructed out of the matter and gravity degrees of freedom, which plays the role of a time variable. Although very interesting proposals of this sort have been considered, those are tied to specific examples, and the general resolution of the problem through this path is still lacking. Moreover, it is not clear if the recovery of time, and more importantly of a unified space-time can be achieved while retaining a fully quantized description.

We, thus, take the conservative view that the characterization of the gravitational degrees of freedom in terms of space-time notions should be done in the context of approximated, and classical level, even if the correct underlying theory is a theory of quantum gravity.

Moreover, it appears very likely that in attempts to incorporate a notion of time by selecting a variable to act as a clock, one ends with a theory where the remaining degrees of freedom are characterized by an evolution equation that involves subtle deviations from the expected Schrödinger's equation. In fact, as discussed in [28], the attempts to recover time in this fashion seem to lead to equations involving loss of unitarity.

## 6. Cosmology

Now, let us consider the problem we have mentioned concerning the emergence of the seeds of structure during inflation.

The problem is that, according to the inflationary ideas, the very early stages of inflation drive our universe into a situation that is described in terms of a background metric and scalar field configuration that is perfectly H&I, and by a perturbation field constructed out of metric and scalar field perturbations in a state closely connected to the Bunch Davies vacuum, and which is also exactly H&I. The issue is the inability of our physical theories to account for the emergence, out of this situation, of the primordial inhomogeneities and anisotropies. We note that several arguments have been put forward in this direction, but as discussed in [2], they all fail to be fully satisfactory. The difficulties are also alluded to in page 476 of S. Weinberg's book on cosmology [14] where we find "*... the field configurations must become locked into one of an ensemble of classical configurations with ensemble averages given by quantum expectation values... It is not apparent just how this happens...*".

The most popular of those arguments rely on the ideas of "decoherence", but that alone is insufficient as reflected in the following quote: "*.. However decoherence is not enough to explain the breakdown of translational invariance..*" found in page 348 of the well known book [15]. Moreover, even W. Zurek says [16]: "*The interpretation based on the ideas of decoherence and ein-selection has not really been spelled out to date in any detail. I have made a few half-hearted attempts in this direction, but, frankly, I was hoping to postpone this task, since the ultimate questions tend to involve such anthropic attributes of the observership as perception, awareness, or consciousness, which, at present, cannot be modeled with a desirable degree of rigor.*"

In the cosmological situation at hand, we can not appeal to "observers" or "measuring apparatuses", as the emergence of the seeds of structure is not only prior to observers or apparatuses, but is actually a prerequisite for the conditions where such things are possible.

This might lead us to the general concern: Is quantum theory applicable to cosmology? If not, where is the boundary of applicability of Quantum theory? Is this an issue that is resolved with some suitable reformulation or interpretation of quantum theory?

Such approach was considered in [42], and led to the development of the Consistent Decohering Histories approach as the only one applicable to cosmology. As we have pointed out in [2] and [17], that approach does not offer a fully satisfactory resolution of the issues. As we said, we need to account for the breaking of the H&I character of the initial state (i.e. the situation corresponding to the early states of inflation), in a context where the dynamics does not contain anything capable of generating such breaking<sup>3</sup>. As already pointed out, our approach is based on the idea that such breaking is to find an explanation within the so called **collapse theories**.

## 7. The proposal

Based on the discussions above, we will look at Einstein's general relativity as an emergent phenomenon, or as "hydrodynamical level of description" of some underlying theory of quantum gravity. We will assume that such description validity requires  $R \ll 1/l_{Plank}^2$ . However, we will not assume that this condition is sufficient. For instance, the situations considered in Penrose's experimental proposals, which clearly satisfy this latter condition, might be ones where important modifications associated with the quantum gravity interface become relevant. We will assume, however, that such situation includes, with the appropriate caveats to be discussed below, the inflationary regime. Of course, that situation is one where the matter fields (particularly the inflation) still require a full quantum treatment.

The setting will thus naturally be semiclassical Einstein's gravity with the **collapse** reflecting some remanent signature from quantum aspects of the gravitation (here, we are following general ideas by Penrose and Diósi): i.e., besides  $U$ , we have, sometimes, spontaneous jumps of the quantum state:

$$....|0\rangle_{k_1} \otimes |0\rangle_{k_2} \otimes |0\rangle_{k_3} \otimes .... \rightarrow ....|\Xi\rangle_{k_1} \otimes |0\rangle_{k_2} \otimes |0\rangle_{k_3} \otimes ....$$

More precisely, we will rely on the notion of *Semiclassical Self-consistent Configuration* (SSC).

**DEFINITION:** The set  $g_{\mu\nu}(x), \hat{\varphi}(x), \hat{\pi}(x), \hat{H}, |\xi\rangle \in \hat{H}$  represents a SSC if and only if  $\hat{\varphi}(x)$ ,  $\hat{\pi}(x)$  and  $\hat{H}$  correspond to a quantum field theory constructed over a space-time with metric  $g_{\mu\nu}(x)$  (as described in, say, [10]), and the state  $|\xi\rangle$  in  $\mathcal{H}$  is such that

$$G_{\mu\nu}[g(x)] = 8\pi G \langle \xi | \hat{T}_{\mu\nu}[g(x), \hat{\varphi}(x), \hat{\pi}(x)] | \xi \rangle.$$

It is, in a sense, the general relativistic version of Schrödinger-Newton equation [34].

This, however, can not describe the transition from a H&I SSC to one that is not. For that we need to add a collapse. We must describe that as a transition from one SSC to another, not simply from one state to another. So, we need the collapse to be represented by a full transition of the type SSC-I  $\rightarrow$  SSC-II. In particular, they will describe a transition from an H&I SSC to one that is not H&I SSC. That involves changing the state, and thus the space-time, and thus the Hilbert space where the state "lives".

We take the view that, at the jump, the degrees of freedom of the underlying theory of quantum gravity, which are not susceptible of a metric description, are excited. In the fluid analogy, this might be thought as corresponding to a chemical reaction or phase transition occurring in underlying matter constituents of the fluid.

<sup>3</sup> People sometimes find attractive to relate this issue with that of spontaneous symmetry breaking. In this regard, we turn the reader to [44] where some of the myths connected with those ideas are clarified.

## 8. Example: a baby version of inflation

As an example of the application of the formalism, we consider the collapse associated with the generation of a primordial perturbation characterizing the breakdown of homogeneity and isotropy (H&I) at a single wave number  $k_0$ . The discussion here is a shorthand version of that presented in [3] in which some inessential complications have been omitted<sup>4</sup>.

As we indicated in the SSC formalism, space-time is thus treated as classical, and, in our case (working in a specific gauge and ignoring the tensor perturbations), the relevant metric is

$$ds^2 = a^2(\eta) \left[ -(1 + 2\psi)d\eta^2 + (1 - 2\psi)\delta_{ij}dx^i dx^j \right], \quad \psi(\eta, \vec{x}) \ll 1 \quad (8)$$

with  $a(\eta)$  the scale factor and  $\psi(\eta, x)$  the Newtonian potential. Later we will limit ourselves to the case where this metric is perfectly H&I, and to the transition to the case where it develops inhomogeneity and anisotropy characterized by  $\vec{k}_0$ . The scalar field with action  $S = 1/2 \int d^4x (\nabla_\mu \phi \nabla^\mu \phi - m^2 \phi^2)$  is treated at the quantum field theory on curved space-times level, so we write:

$$\hat{\phi}(x) = \sum_{\alpha} \left( \hat{a}_{\alpha} u_{\alpha}(x) + \hat{a}_{\alpha}^{\dagger} u_{\alpha}^*(x) \right), \quad (9)$$

with the functions  $u_{\alpha}(x)$  a complete set of normal modes, satisfying the field equation on the corresponding space-time:

$$(g^{\mu\nu} \nabla_{\mu} \nabla_{\nu} + m^2) u_{\alpha}(x) = 0, \quad (10)$$

and orthonormal w.r.t. the symplectic product:

$$((\phi_1, \pi_1), (\phi_2, \pi_2))_{\text{Sympl}} \equiv -i \int_{\Sigma} [\phi_1 \pi_2^* - \pi_1 \phi_2^*] d^3x. \quad (11)$$

For simplicity, we'll set everything in a co-moving coordinate box of size  $L$ .

Finally, one constructs the state such that Einstein's semiclassical equations hold. This is nontrivial, but is a well defined problem. We will, in fact, consider two such constructions and their matching at the collapse time.

### 8.1. The homogeneous and isotropic case: SSC-I

This situation corresponds to choosing  $\psi(\eta, \vec{x}) \equiv 0$ . We assume an almost de Sitter expansion (characterized by the parameters  $H_0^{(I)}$  and  $\epsilon^{(I)}$ ). That is,  $a^{(I)}(\eta) = -1/(H_0 \eta)^{1+\epsilon^{(I)}}$ . More precisely, the conformal expansion factor is defined as  $\mathcal{H}^{(I)} \equiv \frac{\dot{a}^{(I)}(\eta)}{a^{(I)}(\eta)}$  and the slow-roll parameter  $\epsilon^{(I)} \equiv 1 - \dot{\mathcal{H}}^{(I)}/\mathcal{H}^2^{(I)}$ .

The quantum field construction will be based on a complete set of modes of the form  $u_{\vec{k}}^{(I)}(x) = v_{\vec{k}}^{(I)}(\eta) e^{i\vec{k} \cdot \vec{x}}/L^{3/2}$ . These must then satisfy the evolution equation:

$$\ddot{v}_{\vec{k}}^{(I)} + 2\mathcal{H}^{(I)} \dot{v}_{\vec{k}}^{(I)} + \left( k^2 + a^{2(I)} m^2 \right) v_{\vec{k}}^{(I)} = 0, \quad (12)$$

and be normalized:

$$v_{\vec{k}}^{(I)} \dot{v}_{\vec{k}}^{(I)*} - \dot{v}_{\vec{k}}^{(I)} v_{\vec{k}}^{(I)*} = i\hbar a^{-2(I)}. \quad (13)$$

For the modes with  $k \neq 0$ , the most general solution to the evolution equation is a linear combination of the Hankel functions of first and second kind,  $\eta^{3/2} H_{\nu}^{(1)}(-k\eta)$  and  $\eta^{3/2} H_{\nu}^{(2)}(-k\eta)$

<sup>4</sup> One such simplification is the avoidance of the separation of the modes of the quantum field into symmetric and antisymmetric, or real and imaginary, and denoted by *R&I* in [3].



with  $(\nu^{(I)})^2 = (9/4) - (m/H_0^{(I)})^2$ . We are only limited by the requirement that the norm be positive, which is easily achieved.

However, the Hankel functions are not well behaved at the origin, and thus the zero mode is not included. For  $k = 0$ , the general solution to the equation is a linear combination of the functions  $\eta^{(3-2\nu)/2}$  and  $\eta^{(3+2\nu)/2}$ . The choice is again arbitrary, provided it has positive symplectic norm. In this case, we take

$$v_0^{(I)}(\eta) = \sqrt{\frac{\hbar}{H_0^{(I)}}} \left[ 1 - \frac{i}{6} \left( -H_0^{(I)} \eta \right)^3 \right] \left( -H_0^{(I)} \eta \right)^{m^2/3H_0^{2(I)}}. \quad (14)$$

For the  $k \neq 0$ , we make the Bunch-Davis choice, i.e., we use modes which, in the asymptotic past, behave as purely “positive frequency solutions”. This fixes  $\hat{H}^{(I)}$  as the Fock space of that construction.

To complete the SSC, we still need a state  $|\xi^{(I)}\rangle$  in  $\hat{H}^{(I)}$  such that its expectation value for the energy-momentum tensor leads to the desired nearly de Sitter expansion. Consider a state in which all the modes with  $k \neq 0$  are in their vacuum state, while the zero mode is excited in a coherent state:

$$|\xi^{(I)}\rangle = c e^{\xi_0^{(I)} \hat{a}_0^{(I)\dagger}} |0^{(I)}\rangle. \quad (15)$$

The semiclassical Einstein’s equation now reads

$$3\mathcal{H}^{2(I)} = 4\pi G \left( (\dot{\phi}^{2(I)})_{\xi,0} + a^{2(I)} m^2 (\phi^{2(I)})_{\xi,0} \right), \quad (16)$$

$$\mathcal{H}^{2(I)} + 2\dot{\mathcal{H}}^{(I)} = -4\pi G \left( (\dot{\phi}^{2(I)})_{\xi,0} - a^{2(I)} m^2 (\phi^{2(I)})_{\xi,0} \right) \quad (17)$$

where we have used a normal ordering recipe for the energy-momentum tensor, and defined:

$$(\dot{\phi}^{2(I)})_{\xi,0} \equiv \langle \xi^{(I)} | : (\partial_\eta \hat{\phi}^{(I)})^2 : | \xi^{(I)} \rangle, \quad (\hat{\phi}^{2(I)})_{\xi,0} \equiv \langle \xi^{(I)} | : (\phi^{(I)})^2 : | \xi^{(I)} \rangle. \quad (18)$$

In general, the “classical relations”  $(\dot{\phi}^{2(I)})_{\xi,0} = (\dot{\phi}_{\xi,0}^{(I)})^2$  and  $(\phi^{2(I)})_{\xi,0} = (\phi_{\xi,0}^{(I)})^2$  will not hold; they, however, do for our simple coherent a state  $|\xi^{(I)}\rangle$ .

Writing scale factor as  $a(\eta) = (-1/H_0\eta)^{1+\epsilon_1}$  and substituting in Einstein’s equation up to  $\epsilon^{(I)}$ , we obtain  $3(\dot{\phi}_{\xi,0}^{(I)})^2 \approx \epsilon^{(I)} a^{2(I)} m^2 (\phi_{\xi,0}^{(I)})^2$ . This leads to

$$\phi_{\xi,0}^{(I)}(\eta) \propto \eta \sqrt{\epsilon^{(I)} m^2 / 3 (H_0^{(I)})^2}. \quad (19)$$

On the other hand, taking the parameter  $\xi_0^{(I)}$  as real, we find

$$\langle \xi^{(I)} | \hat{\phi}^{(I)}(x) | \xi^{(I)} \rangle = \frac{2\xi_0^{(I)}}{L^{3/2}} \sqrt{\frac{\hbar}{H_0^{(I)}}} \left( -H_0^{(I)} \eta \right)^{m^2/3H_0^{2(I)}}, \quad (20)$$

That is, we have compatibility if we set

$$\epsilon^{(I)} = \frac{m^2}{3H_0^{2(I)}}, \quad H_0^{(I)} = 16\pi G \hbar \epsilon^{(I)} \frac{(\xi_0^{(I)})^2}{L^3}. \quad (21)$$

This was the explicit SSC-I construction representing a H&I state, corresponding to the early stages of inflation.

### 8.2. An inhomogeneous and anisotropic case: SSC-II

Next, we want to consider a situation where the universe is no longer H&I, but has been excited in the  $\vec{k}_0$  mode: we will denote this new SSC by SSC-II. It will be characterized by the parameters  $H_0^{(\text{II})}$  and  $\epsilon^{(\text{II})}$  (which might, in principle, differ slightly from those corresponding to the SSC-I discussed in the previous section), and a Newtonian potential described by  $\psi(\eta, \vec{x}) = \epsilon P(\eta) \cos(\vec{k}_0 \cdot \vec{x})$ , where  $P(\eta)$  is an (in principle) arbitrary function of  $\eta$ , and  $\epsilon$  is a small (expansion) parameter (not to be confused with the slow roll parameter  $\epsilon$ ).

**The strategy:** We first construct the “generic” Hilbert space assuming that  $P(\eta)$  is given. Then make an “educated” guess for the form of the quantum state, and then, by requiring that our construction be an SSC (i.e., that Einstein’s semiclassical equations hold) we will find what the function  $P(\eta)$  ought to be.

The first step is to find the complete set of modes, which we write as

$$u_{\vec{k}}^{(\text{II})}(x) = \frac{1}{L^{3/2}} [v_{\vec{k}}^{(\text{II})0}(\eta) e^{i\vec{k} \cdot \vec{x}} + \epsilon (\delta v_{\vec{k}}^{(\text{II})-}(\eta) e^{i(\vec{k}-\vec{k}_0) \cdot \vec{x}} + \delta v_{\vec{k}}^{(\text{II})+}(\eta) e^{i(\vec{k}+\vec{k}_0) \cdot \vec{x}})] \quad (22)$$

to the zeroth order in  $\epsilon$ . The evolution equation is given, at the lowest order, by

$$\ddot{v}_{\vec{k}}^{(\text{II})0} + 2\mathcal{H}^{(\text{II})} \dot{v}_{\vec{k}}^{(\text{II})0} + \left(k^2 + a^{2(\text{II})} m^2\right) v_{\vec{k}}^{(\text{II})0} = 0, \quad (23)$$

with normalization condition

$$v_{\vec{k}}^{(\text{II})0} \dot{v}_{\vec{k}}^{(\text{II})0*} - \dot{v}_{\vec{k}}^{(\text{II})0} v_{\vec{k}}^{(\text{II})0*} = i\hbar a^{-2(\text{I})}, \quad (24)$$

which is identical to the construction we have already done. Thus, we take  $v_{\vec{k}}^{(\text{II})0}(\eta)$  as before. At first order in  $\epsilon$  the corresponding evolution equation takes the form

$$\delta \ddot{v}_{\vec{k}}^{(\text{II})\pm} + 2\mathcal{H}^{(\text{II})} \delta \dot{v}_{\vec{k}}^{(\text{II})\pm} + \left[(\vec{k} \pm \vec{k}_0)^2 + a^{2(\text{II})} m^2\right] \delta v_{\vec{k}}^{(\text{II})\pm} = F_{\vec{k}}(\eta) \quad (25)$$

where

$$F_{\vec{k}}(\eta) \equiv 4\dot{v}_{\vec{k}}^{(\text{II})0} \dot{P} - 2 \left(2k^2 + a^{2(\text{II})} m^2\right) v_{\vec{k}}^{(\text{II})0} P. \quad (26)$$

The normalization condition (needed only at one time) is

$$\begin{aligned} \dot{v}_{\vec{k}+\vec{k}_0}^{(\text{II})0*} \delta v_{\vec{k}}^{(\text{II})+} - v_{\vec{k}+\vec{k}_0}^{(\text{II})0*} \delta \dot{v}_{\vec{k}}^{(\text{II})+} - \dot{v}_{\vec{k}}^{(\text{II})0} \delta v_{\vec{k}+\vec{k}_0}^{(\text{II})-*} + v_{\vec{k}}^{(\text{II})0} \delta \dot{v}_{\vec{k}+\vec{k}_0}^{(\text{II})-*} \\ = 4 \left( v_{\vec{k}}^{(\text{II})0} \dot{v}_{\vec{k}+\vec{k}_0}^{(\text{II})0*} - \dot{v}_{\vec{k}}^{(\text{II})0} v_{\vec{k}+\vec{k}_0}^{(\text{II})0*} \right) P. \end{aligned} \quad (27)$$

If we had  $P(\eta)$  and the initial conditions satisfying (27), the Eq. (25) would define a unique solution for the functions  $\delta v_{\vec{k}}^{(\text{II})\pm}$ . As we said, we will assume that  $P(\eta)$  is given and take the initial conditions to be

$$\delta \dot{v}_{\vec{k}}^{(\text{II})\pm}(\eta_c) = 0, \quad \delta v_{\vec{k}}^{(\text{II})\pm}(\eta_c) = 4v_{\vec{k}}^{(\text{II})0}(\eta_c) P(\eta_c). \quad (28)$$

This completes the construction of the generic Hilbert space representation  $\hat{H}^{(\text{II})}$  of the quantum field. Next, we need to find the state  $|\zeta^{(\text{II})}\rangle \in \mathcal{H}^{(\text{II})}$  that completes the SSC construction.

The symmetries of the space-time lead us to consider the multi coherent state “ansatz” :

$$|\zeta^{(\text{II})}\rangle = \dots |\zeta_{-2\vec{k}_0}^{(\text{II})}\rangle \otimes |\zeta_{-\vec{k}_0}^{(\text{II})}\rangle \otimes |\zeta_0^{(\text{II})}\rangle \otimes |\zeta_{\vec{k}_0}^{(\text{II})}\rangle \otimes |\zeta_{2\vec{k}_0}^{(\text{II})}\rangle \dots \quad (29)$$

The vector in Fock space is characterized by the specific modes that are excited (all other modes are assumed to be in the vacuum of the corresponding oscillator) and the parameters  $\zeta_{\vec{k}}^{(\text{II})}$  indicate the coherent state for the mode  $\vec{k}$ . The expectation value of the field operator in such a state is given by

$$\phi_{\zeta}^{(\text{II})}(x) = \phi_{\zeta,0}^{(\text{II})}(\eta) + \left( \delta\phi_{\zeta,\vec{k}_0}^{(\text{II})}(\eta) e^{i\vec{k}_0 \cdot \vec{x}} \right) + \left( \delta\phi_{\zeta,2\vec{k}_0}^{(\text{II})}(\eta) e^{i2\vec{k}_0 \cdot \vec{x}} \right) + \dots \quad (30)$$

In fact, for a general state of the form considered above, we have

$$\begin{aligned} \phi_{\zeta,0}^{(\text{II})}(\eta) &= \zeta_0^{(\text{II})} v_0^{(\text{II})0}(\eta) + \varepsilon [\zeta_{-\vec{k}_0}^{(\text{II})} \delta v_{-\vec{k}_0}^{(\text{II})+}(\eta) + \zeta_{\vec{k}_0}^{(\text{II})} \delta v_{\vec{k}_0}^{(\text{II})-}(\eta)] + c.c., \\ \delta\phi_{\zeta,\vec{k}_0}^{(\text{II})}(\eta) &= \zeta_{\vec{k}_0}^{(\text{II})} v_{\vec{k}_0}^{(\text{II})0}(\eta) + \varepsilon [\zeta_0^{(\text{II})} \delta v_0^{(\text{II})+}(\eta) + \zeta_{2\vec{k}_0}^{(\text{II})} \delta v_{2\vec{k}_0}^{(\text{II})-}(\eta)] + \dots, \\ \delta\phi_{\zeta,2\vec{k}_0}^{(\text{II})}(\eta) &= \zeta_{2\vec{k}_0}^{(\text{II})} v_{2\vec{k}_0}^{(\text{II})0}(\eta) + \varepsilon [\zeta_{\vec{k}_0}^{(\text{II})} \delta v_{\vec{k}_0}^{(\text{II})+}(\eta) + \zeta_{3\vec{k}_0}^{(\text{II})} \delta v_{3\vec{k}_0}^{(\text{II})-}(\eta)] + \dots, \end{aligned}$$

We set  $\delta\phi_{\zeta,n\vec{k}_0}^{(\text{II})}(\eta) = 0$  for all  $n \geq 2$ , simply by imposing the required relations between the parameters  $\zeta_{\pm\vec{k}_0}^{(\text{II})}$ ,  $\zeta_{\pm 2\vec{k}_0}^{(\text{II})}$ ,  $\zeta_{\pm 3\vec{k}_0}^{(\text{II})}$ , etc. It is easy to see that  $|\zeta_{\pm n\vec{k}_0}^{(\text{II})}| \sim \varepsilon^n |\zeta_0^{(\text{II})}|$ .

The conditions above ensure that there are no terms in  $e^{\pm i n \vec{k}_0 \cdot \vec{x}}$  (with  $n \geq 2$ ) appearing in the expectation value of the energy-momentum tensor. That is necessary for compatibility of our state ansatz with the semiclassical Einstein’s equations. Let us consider these in detail up to the first order in  $\varepsilon$ . The zero order equations are just the same as the equations we found in SSC-I, and are used to fix the construction of SSC up to that order, i.e., to determine the relation between  $a^{(\text{II})}$  and  $\zeta_0^{(\text{II})}$ .

The key result, and the aspect that enables us to carry out the construction in a complete manner, is the following fact (see [3] for the details): the first order semiclassical Einstein’s equations for the metric and the above state can be combined into a simple dynamical equation for the Newtonian potential, which is independent of the first order quantities  $\delta\tilde{\phi}_{\zeta,\vec{k}_0}^{(\text{II})}$  and  $\delta\dot{\phi}_{\zeta,\vec{k}_0}^{(\text{II})}$ . To the level of precision we are working, the equation above becomes simply

$$\ddot{P} + \epsilon^{(\text{II})} \mathcal{H}^{(\text{II})} \dot{P} + \left[ k_0^2 - \epsilon^{(\text{II})} \mathcal{H}^2(\text{II}) \right] P = 0. \quad (31)$$

The general solution reads:

$$P(\eta) = C_1 \eta^{\frac{1}{2}[1+\epsilon^{(\text{II})}]} J_{\alpha}(-k\eta) + C_2 \eta^{\frac{1}{2}[1+\epsilon^{(\text{II})}]} Y_{\alpha}(-k\eta), \quad (32)$$

where  $J_{\alpha}(-k\eta)$  and  $Y_{\alpha}(-k\eta)$  are the Bessel functions of first and second kind,  $\alpha = [1 + 3\epsilon^{(\text{II})}]/2$ .

The remaining Einstein’s equations to this order take the form of two constraints (again, see [3] for the details). Using these, we express the initial values that would determine the specific solution  $P(\eta)$ :

$$\begin{pmatrix} P \\ \dot{P} \end{pmatrix} = \frac{\sqrt{4\pi G \epsilon^{(\text{II})} \mathcal{H}^{(\text{II})}}}{k_0^2 - \mathcal{H}^2(\text{II}) \epsilon^{(\text{II})}} \begin{pmatrix} (3\mathcal{H}^{(\text{II})} - a m \sqrt{3/\epsilon^{(\text{II})}}) & 1 \\ (a m \sqrt{3/\epsilon^{(\text{II})}} \mathcal{H}^{(\text{II})} - k_0^2 + (\epsilon^{(\text{II})} - 3) \mathcal{H}^2(\text{II})) & -\mathcal{H}^{(\text{II})} \end{pmatrix} \begin{pmatrix} \delta\tilde{\phi}_{\zeta,\vec{k}_0}^{(\text{II})} \\ \delta\dot{\phi}_{\zeta,\vec{k}_0}^{(\text{II})} \end{pmatrix}. \quad (33)$$

Therefore, given the two numbers  $\delta\tilde{\phi}_{\zeta,\vec{k}_0}^{(\text{II})}(\eta_c)$  and  $\dot{\delta\tilde{\phi}}_{\zeta,\vec{k}_0}^{(\text{II})}(\eta_c)$ , we would have a completely determined space-time metric. That is the condition above would determine  $P$  and  $\dot{P}$  at  $\eta_c$ , and these would determine the constants  $C_1$  and  $C_2$  in (32). And then, the Eq. (25), together with the initial conditions (27) (which depend on  $P(\eta_c)$  which, as we saw was fixed), would fix the mode functions for all  $\eta$ . Furthermore, those two numbers would determine the state parameters  $\zeta_{\vec{k}_0}$  (and thus the rest as well). Thus, we have a complete SSC-II (to this order in  $\varepsilon$ ).

### 8.3. The actual collapse process or the matching of SSC-I and SSC-II

Next, we want to consider the process of collapse. That is represented by the transition from the SSC-I to the SSC-II, and, in particular, the **matching** of the H&I region of space-time with the region where such symmetry has disappeared. We will be considering here that the transition corresponds to the hypersurfaces  $\eta = \eta_c$  of SSC-I and SSC-II. Note that this gives such hypersurface  $\Sigma_c$  a preferred status in the resulting space-time, and is not something to be thought as related to a gauge freedom: to the past of that hypersurface  $\Sigma_c$ , the space-time is H&I, and to the future it is not. We will assume here the induced metric is continuous on  $\Sigma_c$ . As shown in [3], the continuity can be obtained as a conclusion of a natural assumption regarding the energy-momentum tensor. In any event, continuity requires  $P(\eta_c) = 0$  and thus

$$(3\mathcal{H}^{(\text{II})} - am\sqrt{3/\epsilon^{(\text{II})}})\delta\tilde{\phi}_{\zeta,\vec{k}_0}^{(\text{II})}(\eta_c) + \dot{\delta\tilde{\phi}}_{\zeta,\vec{k}_0}^{(\text{II})}(\eta_c) = 0, \quad (34)$$

and therefore

$$\dot{P} = -\sqrt{4\pi G\epsilon^{(\text{II})}}\mathcal{H}^{(\text{II})}\delta\tilde{\phi}_{\zeta,\vec{k}_0}^{(\text{II})}(\eta_c). \quad (35)$$

Finally, we want to use the conditions prior to the collapse to help determine the situation just after it. Following the early work in [1], we will assume that the collapse is characterized by a loose analogy with something like an imprecise measurement (of the operators  $\hat{\phi}_{\vec{k}_0}^{(\text{I})}(\eta)$ ) in standard quantum theory: Before the collapse the operator had an uncertainty  $\Delta\hat{\phi}_{\vec{k}_0}^{(\text{I})}(\eta_c)$ , and after the collapse, the new state will have an expectation value for that operator given by a random number within the range of uncertainty prior to collapse. The final result is then:

$$\varepsilon\delta\tilde{\phi}_{\zeta,\vec{k}_0}^{(\text{II})}(\eta_c) = x_{\vec{k}_0}\sqrt{\langle 0_{\vec{k}_0}^{(\text{I})} | \left[ \Delta\hat{\phi}_{\vec{k}_0}^{(\text{I})}(\eta_c) \right]^2 | 0_{\vec{k}_0}^{(\text{I})} \rangle} \approx x_{\vec{k}_0} a(\eta_c)^{-1} \sqrt{\frac{\hbar}{2k}} \quad (36)$$

with  $x_{\vec{k}_0}$  taken to be a random variable distributed according to a Gaussian function centered at zero with unit-spread. Thus, a specific choice of the random number  $x_{\vec{k}_0}$  fully determines the SSC-II.

This finishes the analysis of the emergence of an actual anisotropic and inhomogeneous perturbation associated with a single mode, through a single discrete collapse in that mode. We note that the scheme leads to a description of space-time in terms of two regions, one corresponding to the homogeneous and isotropic metric of SSC-I, and one of the slightly inhomogeneous and anisotropic metric of SSC-II. The matching conditions on the hypersurface  $\Sigma_c$  are such that the metric is continuous, but the extrinsic curvature is not. This issue is closely tied to the jump of the energy-momentum tensor across such surface. This in turn is closely connected to the issue of conservation laws in any collapse theory (see however [43]).

## 9. Discussion

We have presented a well defined framework where one could, in principle, carry out all the analysis of the collapse approach to the inflationary origin of the seeds of cosmic structure. However, the fact that it is so complicated has led us, in the past to seek and use some shortcuts in the actually phenomenological analysis such as those in [20, 21]. But having a clear framework opens the path to the investigations of topics of more fundamental nature.

Of course, the fact that the scheme is well defined does not mean that it is free of problems, or that all questions have been successfully addressed. As far as we see, the most serious problems, afflicting this proposal are: a) the fact that during the collapse or jump from one SSC to another the energy-momentum tensor fails to be conserved (i.e. its four-divergence jumps); b) the fact that the collapse, as described here, takes place on a certain space-like hypersurface, and is thus a highly nonlocal process.

Regarding issue a), as we already indicated, the source is the generic violation of energy conservation that afflicts all collapse theories. The first thing to note in this regard is the fact that new developments in such collapse proposals provide hope for the possibility of modifications curing that problem, as discussed for instance in [33]. On the other hand, if we do take the view that space-time is emergent, and thus not a fundamental concept, it seems natural to question whether energy-momentum and its conservation should still be regarded as fundamental. If we take the view that Einstein's equations are a sort of Navier-Stokes effective equations, then their absolute validity would be limited, and of course so should be the conservation of energy-momentum. After all, as we already noted, nobody would be surprised if a violation of Navier-Stokes equations was observed in a liquid undergoing some process where the fundamental microscopic degrees of freedom of the fluid in question were being excited. In such situation, exemplified by one where the fluid was undergoing a phase transition, we would expect the fluid dynamics description to fail or at least to require some correction.

Regarding the issue b) one might be concerned with the nonlocality being tied to violations of causality. The first thing we should put out is that is easy to imagine processes that look highly nonlocal, when viewed in the context of an effective theory, and which are seen to be essentially local when viewed within the context of a fundamental theory. Consider, for instance, a very large lake that freezes essentially instantaneously as a result of a rapid drop in temperature. It is clear that such instantaneous freezing might, at first sight, look problematic in the context of special relativity, but a moment's thought clarifies that there is nothing odd going on. The particular preferential frame in which the freezing is instantaneous is determined globally (on the scale of the lake's size), as the frame where the temperature drops is uniform and homogeneous. The microscopic process is clearly controlled by local physics, and the instantaneous freeze-out is the result of the change in external conditions taking place in a correlated fashion over large distances. Second, when one considers collapse taking place in a stochastic fashion, and for all modes, it is quite possible that the effective collapse can be described in terms of something more local, as it is explicitly shown to occur in the inflationary problem when addressed in terms of the theory known as continuous spontaneous localization as discussed in [25], a work related to that of [24], but taking a slightly different point of view.

Furthermore, non locality and violation of causality, although superficially connected, are two distinct concepts. In fact, any analysis of a B-EPR system (and its experimental confirmation as the breakdown of Bell's inequalities) indicates that the system possesses nonlocal aspects. However, it is also clear that it does not violate causality (one can not use the EPR setups to send information faster than light, see for instance [45]).

In fact, the standard Copenhagen interpretation, where the collapse is the result of an "observer" performing a measurement, is, if anything, more susceptible to causality violations, as discussed for instance in [46], simply because there, the will of the experimenter plays a central role in determining the time of collapse. The current proposal falls within the class of ideas

known as “spontaneous collapse”, where the actions of experimental physicists do not determine the time of collapse. This fact seriously limits the ability of anyone to use the underlying physics to produce causality violations. Of course, as the final theory about the dynamical collapse is not at hand, we can not say that there is a proof that such a theory exists where all problems related to causality will be avoided. In any case, it is our view that that this line of studies which is motivated by the conceptual difficulties facing the inflationary account of the emergence of the seeds of cosmic structure, can offer one of the few sources of guidance in the search for such theory. It is worth pointing out that the approach does have phenomenological implications, as shown by the studies in [20, 21, 22], and the particular development presented here has lead to suggestions to look for novel kinds of statistical anomalies in then large scale structure and in the cosmic microwave background, which go beyond the usual kind of non-gaussianities [47]. Such analysis might be used to constrain the kind of viable collapse theories, if one takes the view that the phenomenological success of inflation must be backed by the corresponding conceptual clarity.

Finally, it is worth noting that, when considering and assuming the inflationary problem (ignoring the difficulties discussed in [2]) one adopts one of the more popular postures regarding the emergence of classicality, or more precisely, the generation of primordial inhomogeneities and anisotropies. The effective characterization of the actual emergence, if one truly attempts it, is bound to rely on a closely related formalism. These popular postures include, for instance, i) the notion that after a given mode exits the horizon (its physical wavelength, as seen in a co-moving frame, becomes larger than the Hubble radius), the fluctuation corresponding to that mode becomes classical, or ii) that, due to some decoherence effect, we can, at a certain point, adopt the “many worlds interpretation” of quantum theory, and consider the state of the quantum field as characterizing not our universe, but an ensemble of universes of which ours is just a typical element. Therefore, let us say that in case i) we want to produce a description (even an approximate one) of our universe concentrating, for simplicity, on a single mode, and want the characterization to be valid both before and after the mode exits the horizon. In that case, the approach presented here seems to be the best one can do, as long as we do not have a workable theory of quantum gravity which allows us to characterize space-time in a full quantum language. In case ii), we might also be interested in putting together the characterization of our space-time before the decoherence is taken to be effective, and the description of the particular branch of the many worlds (the particular element of the ensemble of universes) in which we happen to find ourselves. Again, in that situation the present analysis offers perhaps the furthest one can go in achieving the said goal, given the present stage of the development of candidate theories of quantum gravity.

The inflationary account for the emergence of the seeds of structure, represents the only situation where general relativity and quantum theory come together in the explanation of observable phenomena. We can let the phenomenological success blind us to its shortcomings [2] but that would mean discarding what might be perhaps the only source of empirical clues available to the enterprise of clarifying the way these two pillars of modern physics are to be reconciled. If instead we try to address these shortcomings we will have to face very difficult questions, and in doing so we would be clearly at risk of making tremendous mistakes. However in this case we can take consolation from Sir Francis Bacon’s profound observation<sup>5</sup> about scientific methodology: *“truth emerges more readily from error than from confusion”*.

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<sup>5</sup> As cited in Thomas Kuhn’s “The structure of Scientific Revolutions”

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