

# Nonperturbative effects in inflation

Alexei A. Starobinsky

Landau Institute for Theoretical Physics RAS,  
Moscow, Russia

DESY Theory Workshop 2011  
"Cosmology meets particle physics"  
Hamburg, 27.09.2011

Perturbative results

Stochastic slow-roll inflation

Langevin and Fokker-Planck equations

Application: equilibrium state on the de Sitter background

Probabilities to go to different vacua after inflation

Local duration of inflation

Choice of an initial condition

Conclusions

# Perturbative results

A fixed background - the de Sitter or, more interestingly, quasi-de Sitter space-time (slow roll inflation).

Anomalous growth of light scalar fields occurs for

$0 \leq m^2 \ll H^2$  where  $H \equiv \frac{\dot{a}}{a}$ ,  $a(t)$  is a LFRW scale factor.

The simplest and textbook example  $H = H_0 = \text{const}$ ,  $m = 0$  (Linde, 1982; AS, 1982; Vilenkin and Ford, 1982):

$$\langle \phi^2 \rangle = \frac{H^2 N}{4\pi^2}$$

where  $N = \ln \frac{a}{a(t_0)}$  is the number of e-folds from the beginning of inflation. Straightforward generalization to the quasi-de Sitter case  $|\dot{H}| \ll H^2$ .

For  $0 < m^2 \ll H^2$ , the Bunch-Davies equilibrium value

$$\langle \phi^2 \rangle = \frac{3H_0^4}{8\pi m^2} \gg H_0^2$$

is reached after a large number of e-folds  $N \gg \frac{H_0^2}{m^2}$ . It is almost never reached during slow-roll inflation.

The purely infrared effect - creation of real field fluctuations; renormalization is not important and does not affect it (at least, in the leading approximation).

For the de Sitter inflation (gravitons only) (AS, 1979):

$$P_g(k) = \frac{16GH_0^2}{\pi}; \quad \langle h_{ik}h^{ik} \rangle = \frac{16GH_0^2N}{\pi}.$$

The assumption of small perturbations breaks down for  $N \gtrsim 1/GH_0^2$ . Still ongoing discussion on the final outcome of this effect. My opinion - no screening of the background cosmological constant, instead - stochastic drift through an infinite number of locally de Sitter, but globally non-equivalent vacua.

An always larger effect - due to scalar (adiabatic) fluctuations of an inflaton degree of freedom (Mukhanov and Chibisov, 1981; Hawking, 1982; AS, 1982; Guth and Pi, 1982):

$$P_{\zeta}(k) = \frac{GH_k^4}{\pi|\dot{H}|_k}$$

where the index  $k$  means that the quantity is taken at the moment  $t = t_k$  of the Hubble radius crossing during inflation for each spatial Fourier mode  $k = a(t_k)H(t_k)$ .

The consistency relation

$$r \equiv \frac{P_g}{P_{\zeta}} = \frac{16|\dot{H}_k|}{H_k^2} = 8|n_g| < 0.24$$

(the last inequality - from the most recent observations).

# Beyond small perturbations during slow-roll inflation

Non-perturbative effects  $\implies$  when backreaction is taken into account.

Locally – **around our world-line** – slow-roll inflation has both the beginning and the end.

Globally it has no beginning and no end in the most of interesting cases – **in the sense that inflating patches always exist somewhere in space and time (but outside our past and future light cones) - "eternal inflation"**.

Taking backreaction into account  $\implies$  quantum background.

For sufficiently large  $N_f = \ln\left(\frac{a_f}{a}\right)$ ,  $\langle h^2 \rangle$  becomes larger than 1. Loop corrections proportional to higher powers of  $h$  may become important in this regime only.

Stochastic approach to inflation ("stochastic inflation"):

$$\hat{R}_\mu^\nu - \frac{1}{2}\delta_\mu^\nu \hat{R} = 8\pi G \hat{T}_\mu^\nu(\hat{g}_{\alpha\beta})$$

- not as a function of  $\langle \hat{g}_{\alpha\beta} \rangle$  !

Leads to QFT in a stochastic background.

Stochastic inflation:

- 1) can deal with an arbitrary large (though sufficiently smooth) global inhomogeneity;
  - 2) takes backreaction of created fluctuations into account;
  - 3) goes beyond any finite order of loop corrections.
- Fully developed in Starobinsky (1984,1986) though the first simplified application (but beyond the one-loop approximation) was already in Starobinsky (1982).

# Langevin equation for the large-scale field

The first main idea: splitting of the inflaton field  $\phi$  into a large-scale and a small-scale parts with respect to  $H$ . More exactly, the border is assumed to lie at  $k = \epsilon aH$  with

$$\exp\left(-\frac{H^2}{|H|}\right) \ll \epsilon \ll 1.$$

The second main idea: a non-commutative part of the large-scale field is very small (it is composed from decaying modes), so we may neglect it. Then the remaining part is equivalent (not equal!) to a stochastic c-number (classical) field with some distribution function.

$$\frac{d\phi}{d\tau^{(n)}} = -\frac{1}{3H^{n+1}} \frac{dV}{d\phi} + f,$$

$$\langle f(\tau_1^{(n)}) f(\tau_2^{(n)}) \rangle = \frac{H^{3-n}}{4\pi^2} \delta(\tau_1^{(n)} - \tau_2^{(n)}).$$

Applicability conditions – the standard slow-roll ones:

$$V'^2 \ll 48\pi G V^2, \quad |V''| \ll 8\pi G V/3.$$



The Gaussian white noise  $f$  describes the flow of small-scale linear field modes through the border  $k = \epsilon a H$  to the large-scale region in the course of the universe expansion.

The time-like variables:  $\tau^{(n)} = \int H^n(t, \mathbf{r}) dt$ , where  $H^2 = 8\pi G V(\phi)/3$ .

This is **not** a time reparametrization  $t \rightarrow f(t)$  in GR. Different  $\tau^{(n)}$  describe different stochastic processes and even have different dimensionality. Different "clocks" are needed to measure them:

- 1)  $n = 0$ : phase of a wave function of a massive particle ( $m \gg H$ );
- 2)  $n = 1$ : scalar metric perturbations ( $\delta N$  formalism);
- 3)  $n = 3$ : dispersion of a light scalar field generated during inflation

$$\langle \chi^2 \rangle = \frac{1}{4\pi^2} \langle \int H^3 dt \rangle = \frac{\langle \tau^{(3)} \rangle}{4\pi^2} .$$

See F. Finelli *et al.*, Phys. Rev. D **79**, 044007 (2009) for more details.

# Einstein-Smoluhovsky (Fokker-Planck) equation

$$\frac{\partial \rho}{\partial \tau} = \frac{\partial}{\partial \phi} \left( \frac{V'}{3H^{n+1}} \rho \right) + \frac{1}{8\pi^2} \frac{\partial^2}{\partial \phi^2} (H^{3-n} \rho) .$$

Probability conservation:  $\int \rho d\phi = 1$ .

Remarks.

- ▶ More generally, the last term can be written the form

$$\frac{1}{8\pi^2} \frac{\partial}{\partial \phi} \left( H^{(3-n)\alpha} \frac{\partial}{\partial \phi} (H^{(3-n)(1-\alpha)} \rho) \right)$$

with  $0 \leq \alpha \leq 1$ .

$\alpha = 0$  – Ito calculus.

$\alpha = 1/2$  – Stratonovich calculus.

However, keeping terms explicitly depending on  $\alpha$  exceeds the accuracy of the stochastic approach. Thus,  $\alpha$  may put 0.

- ▶ All results are independent of the form of a cutoff in the momentum space as far as it occurs for  $k \ll aH$  ( $\epsilon \ll 1$ ).
- ▶ Backreaction is taken into account:  $\delta T_{\mu}^{\nu} = (V - V_{clas}) \delta_{\mu}^{\nu}$ .
- ▶ No necessity in any infrared cutoff. Problems with the so called "volume weighting" arise because quantities like  $a^3 \rho$  are considered which are not normalizable, thus, they may not be considered as probabilities of anything from the mathematical point of view ("unitarity breaking"). Their physical justification is also flawed since it based on the wrong assumption that all Hubble physical volumes ("observers") at given  $\tau$  are clones of each other while it is not so.

- ▶ Another possible source of apparent infrared divergences: use of "gauge invariant" (with respect to a background space-time metric) variables like  $\zeta(\mathbf{r}, t)$  which are not generally covariant with respect to the full metric and, therefore, not directly observable. In contrast, quantities like  $\zeta(\mathbf{r}, t) - \zeta(0, t_0)$  are generally covariant and observable though non-local.
- ▶ The accuracy of the stochastic approach is not sufficient for calculating quantities  $\sim H^2$  in  $\langle \phi^2 \rangle$  and  $\sim H^4$  in EMT average values because of the omission of a contribution from the small-scale part (including the conformal anomaly). However, all larger quantities (if exist) can be calculated quantitatively correctly.

# Transition to predictions for the post-inflationary evolution

From  $\rho(\phi, \tau)$  during inflation to the distribution  $w(\tau)$  over the total local duration of inflation:

$$w(\tau) = \lim_{\phi \rightarrow \phi_{end}} j = \lim_{\phi \rightarrow \phi_{end}} \frac{|V'|}{3H^{n+1}} \rho(\phi, \tau).$$

For the graceful exit to a post-inflationary epoch, the stochastic force should be much less than the classical one during last e-folds of inflation.

## From $\delta N$ - to $N$ -formalism

The same way to obtain the joint distribution  $w(0, \tau_1; |\mathbf{r}|, \tau_2)$  from the 2-point joint probability distribution  $\rho(\phi_1, 0, \tau_1; \phi_2, |\mathbf{r}|, \tau_2)$  during inflation.

Let  $n = 1$ . When

$$ds^2 = dt^2 - a^2(t) e^{2\zeta(\mathbf{r})} d\mathbf{r}^2 + \text{small terms}$$

after inflation and complete thermalization where

$$\zeta(\mathbf{r}) = N(\mathbf{r}) \equiv \tau^{(1)}(\mathbf{r}) .$$

# QFT of a self-interacting scalar field in the de Sitter background

A.A. Starobinsky and J. Yokoyama, Phys. Rev. D **50**, 6357 (1994).

The equilibrium (static) solution for the 1-point distribution:

$$\rho_{\text{eq}}(\phi) = \text{const } e^{-2v}, \quad v = \frac{4\pi^2 V(\phi)}{3H_0^4}.$$

Arbitrary Green functions and n-point distributions can be constructed, too, using solutions of the same Fokker-Planck equation.

$$V(\phi) = V_0 + \frac{1}{2}m^2\phi^2 + \frac{1}{4}\lambda\phi^4, \quad 0 < \lambda \ll 1, \quad H_0^2 = \frac{8\pi GV_0}{3}.$$

Three regimes:

1. Perturbative regime  $\sqrt{\lambda}H_0^2 \ll m^2 \ll H_0^2$ .

$$\langle \phi^2 \rangle = \frac{3H_0^4}{8\pi m^2} \left( 1 - 3\beta + \frac{429}{16}\beta^2 + \dots \right), \quad \beta = \frac{3\lambda H_0^4}{8\pi^2 m^4}.$$

Compare to the same result in the one-loop (Gaussian) approximation:

$$\langle \phi_G^2 \rangle = \frac{3H_0^4}{8\pi m^2} (1 - 3\beta + 18\beta^2 + \dots).$$



2. Massless self-interacting regime  $|m^2| \ll \sqrt{\lambda} H_0^2$ .

$$\langle \phi^2 \rangle = \sqrt{\frac{3}{2\pi^2}} \frac{\Gamma(0.75)}{\Gamma(0.25)} \frac{H_0^2}{\sqrt{\lambda}} \approx 0.132 \frac{H_0^2}{\sqrt{\lambda}}$$

$$\langle \phi_G^2 \rangle = \frac{1}{\pi\sqrt{8}} \frac{H_0^2}{\sqrt{\lambda}} \approx 0.113 \frac{H_0^2}{\sqrt{\lambda}}$$

The scale  $m^2 \sim \lambda H_0^2$  proposed recently in arXiv:1005.3551 is not critical at all!

3. Symmetry breaking regime  $m^2 < 0$ ,  $\sqrt{\lambda} H_0^2 \ll |m^2| \ll H_0^2$ .

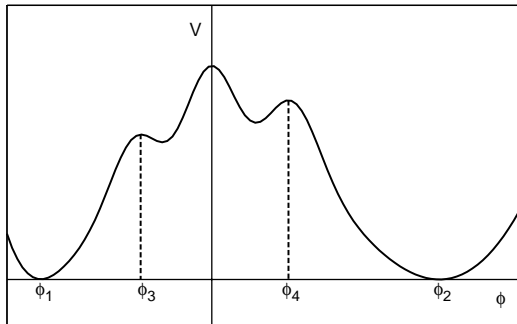
$$\langle \phi^2 \rangle = \frac{|m^2|}{\lambda} + \frac{3H_0^4}{16\pi^2|m^2|} + \mathcal{O}(e^{-1/(4\beta)})$$

The (modulus of) exponent is the action for the Hawking-Moss instanton.

See also F. Finelli *et al.*, Phys. Rev. D **82**, 064020 (2010).

# Probabilities to go to different vacua after inflation

Let inflation may end in two vacua:  $\phi = \phi_1$  and  $\phi = \phi_2$  with  $V(\phi_1) = V(\phi_2) = 0$  (to consider a larger number of post-inflationary vacua,  $\phi$  should have more than one-dimensional internal space).



Boundary conditions at the end of inflation:

$$\rho(\phi_1, \tau) = \rho(\phi_2, \tau) = 0.$$

Method of calculation (Starobinsky (1984,1986)): consider the quantities

$$Q_m(\phi) = \int_0^\infty \tau^m \rho(\phi, \tau) d\tau$$

where  $\tau = 0$  corresponds to the local beginning of inflation.

$$Q_m(\phi_1) = Q_m(\phi_2) = 0.$$

By integrating the Fokker-Planck equation over  $\tau$ , we get for  $m = 0$ :

$$Q_0(\phi) = \frac{8\pi^2}{H^{3-n}} \exp\left(\frac{\pi}{GH^2(\phi)}\right) \int_{\phi_1}^{\phi} d\psi \exp\left(-\frac{\pi}{GH^2(\psi)}\right) \times$$

$$\left(C_0 - \int_{\phi_1}^{\psi} \rho_0(\psi_1) d\psi_1\right),$$

$$C_0 = \frac{\int_{\phi_1}^{\phi_2} d\phi \exp\left(-\frac{\pi}{GH^2(\phi)}\right) \int_{\phi_1}^{\phi} \rho_0(\psi) d\psi}{\int_{\phi_1}^{\phi_2} d\phi \exp\left(-\frac{\pi}{GH^2(\phi)}\right)}.$$

$P_1 = C_0$  – the absolute probability to go to the vacuum

$\phi = \phi_1$ ;

$P_2 = 1 - C_0$  – the absolute probability to go to the vacuum

$\phi = \phi_2$ .

No  $n$  dependence in  $C$  !

## Local duration of inflation

$$Q_1(\phi) = \frac{8\pi^2}{H^{3-n}} \exp\left(\frac{\pi}{GH^2(\phi)}\right) \int_{\phi_1}^{\phi} d\psi \exp\left(-\frac{\pi}{GH^2(\psi)}\right) \times \\ \left(C_1 - \int_{\phi_1}^{\psi} Q_0(\psi_1) d\psi_1\right),$$

$$C_1 = \frac{\int_{\phi_1}^{\phi_2} d\phi \exp\left(-\frac{\pi}{GH^2(\phi)}\right) \int_{\phi_1}^{\phi} Q_0(\psi) d\psi}{\int_{\phi_1}^{\phi_2} d\phi \exp\left(-\frac{\pi}{GH^2(\phi)}\right)}.$$

$$\langle \tau_1 \rangle = \frac{C_1}{C_0}, \quad \langle \tau_2 \rangle = \frac{\tilde{C}_1}{1 - C_0},$$

$$\langle \tau \rangle_{tot} = C_0 \langle \tau_1 \rangle + (1 - C_0) \langle \tau_2 \rangle = \int_{\phi_1}^{\phi_2} Q_0(\phi) d\phi.$$

$\tilde{C}_1$  is  $C_1$  with  $\phi_1$  and  $\phi_2$  interchanged.

# Choice of an initial condition

- ▶ Static solutions – not normalizable in the inflationary (i.e. unstable) case.
  - ▶  $\rho_0(\phi) = \delta(\phi - \phi_0)$  – why?
  - ▶ "Eternal inflation as an initial condition":  $\rho_0(\phi) \propto \rho_{E_1}(\phi)$ 
    - the wave function of the lowest energy level of the Schrodinger equation arising through the separation of variables in the Fokker-Planck equation ( $E_0 = 0$  due to hidden supersymmetry of the former).
    - 1) Not possible in the continuum spectrum case.
    - 2) In the discrete spectrum case, generically  $E_2 - E_1 \sim E_1$ 
      - not enough time for relaxation.
- As a whole, "eternal" inflation seems not be eternal enough to fix the initial condition uniquely.

However, if inflation had occurred at all, the dependence of predictions on  $\rho_0(\phi)$  is comparatively weak: for almost all  $\rho_0(\phi)$  except from the HH-like one  $\rho_0(\phi) \propto \exp\left(\frac{\pi}{GH^2(\phi)}\right)$ , the main contribution comes from the highest maximum of  $V(\phi)$  without any necessity of a "tunneling" initial condition. On the other hand, if  $\rho_0(\phi) \propto \exp\left(\frac{\pi}{GH^2(\phi)}\right)$ , there is practically no inflation at all, and final probabilities  $P_1$  and  $P_2$  are equal to the initial ones.

# Conclusions

- ▶ During slow-roll inflation, backreaction of created inflaton fluctuations has to be taken into account for sufficiently long inflation that leads to QFT in a stochastic background.
- ▶ No problems of principle in predicting all joint probability distributions for light scalar fields, including the inflaton itself, during and after inflation ( $N$ -formalism) in the original (probability conserving) stochastic approach, once an initial condition  $\rho_0(\phi)$  is given. No necessity to refer to other universes outside our light cone.



- ▶ No satisfactory principle to fix  $\rho_0(\phi)$  uniquely.
- ▶ Some dependence on  $\rho_0(\phi)$  remains in final answers, so a possibility to get some knowledge on it from observational data does not seem hopeless. However, if inflation had occurred at all, the dependence of predictions on  $\rho_0(\phi)$  is weak and mainly produced by the region around the highest maximum of  $V(\phi)$ . For this, no specific "tunneling" initial condition is needed.