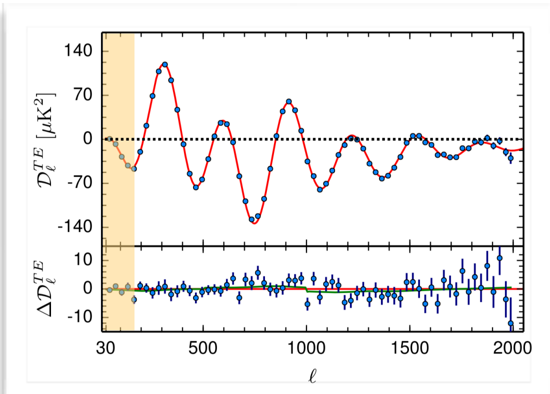
Luca Santoni

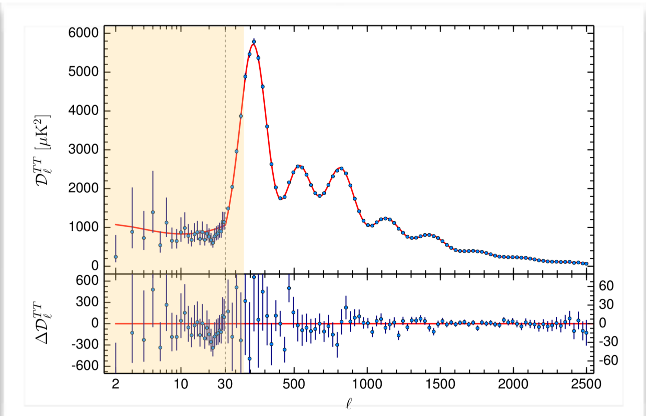
Dissipative inflation via scalar particle production

based on [arXiv:2305.07695] with
Paolo Creminelli, Soubhik Kumar and Borna Salehian

Cosmic Inflation

- Inflation is the most compelling scenario for the early universe.

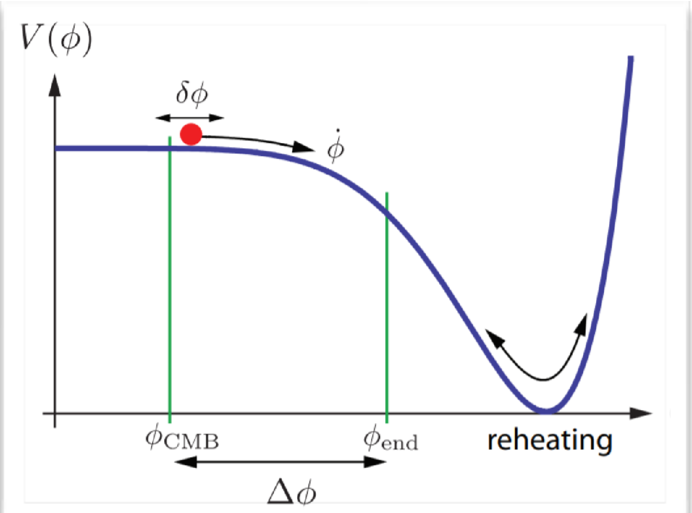




- *Coherence* + (approximate) *scale invariance* ⇒ phase of (quasi-)de Sitter expansion in the early Universe:

$$\mathrm{d}s^2 = -\,\mathrm{d}t^2 + e^{2Ht}\mathrm{d}\mathbf{x}^2.$$

- The simplest model of inflation based on a slowly-rolling scalar field ϕ:



$$S = \int \mathrm{d}^4x\sqrt{-g}\left[\frac{M_{\rm Pl}^2}{2}R - \frac{1}{2}(\partial\phi)^2 - V(\phi)\right]$$

$$\ddot\phi + 3H\dot\phi + V'(\phi) = 0$$

The future threshold $f_{\rm NL} \sim 1$

- The minimal slow-roll model $\mathscr{L} = -\frac{1}{2}(\partial\phi)^2 - V(\phi)$ of inflation is very weakly coupled, with slow-roll suppressed non-Gaussianity.

- Reaching the threshold $f_{\rm NL} \sim 1$ in the future will be very informative.
  The non-observation of any $f_{\rm NL} \sim 1$ would rule out large classes of models; these can be distinguished according to:
  1. perturbations generated by a second field; [Bernardeau and Uzan '02] +…
  2. subluminal propagation: $\mathscr{L} = P((\partial\phi)^2, \phi)$; [Cheung et al '07] +…
  3. different symmetry breaking patterns for inflation: solid, super-solid, gaugid…; [Endlich, Nicolis and Wang '12] +…
  4. dissipative effects: *this talk* and $\Delta\mathscr{L} = \frac{\alpha}{f}\phi F\tilde{F}$: $\ddot{\varphi}_k + (3H+\gamma)\dot{\varphi}_k + \left(\frac{k^2}{a^2} + V''\right)\varphi_k = \ldots$; [Creminelli, Kumar Salehian and LS '23], [Anber and Sorbo '09] +…
  5. warm inflation: $\Delta\mathscr{L} = \frac{\alpha}{16\pi f}\phi \mathrm{Tr}[G_{\mu\nu}\tilde{G}^{\mu\nu}]$; [Berghaus, Graham and Kaplan '19] +…
  6. alternative models to inflation (genesis…).

Outline and summary

- I will describe a new mechanism that gives rise to dissipation during cosmic inflation.
- As opposed to $\Delta\mathscr{L} = \frac{\alpha}{f}\phi F\tilde{F}$, dissipation takes place on parametrically sub-horizon scales, allowing to describe the dynamics in a local manner.
- The presence of dissipation leads to primordial non-Gaussianity with strength $f_{\rm NL}^{\rm eq} \simeq O(10)$.

The model

The model

[Creminelli, Kumar, Salehian and LS '23]

- Inflaton ϕ couples to complex scalar field χ charged under (softly-broken) global U(1):

$$S = \int \mathrm{d}^4x\sqrt{-g}\left[\frac{1}{2}M_{\rm Pl}^2 R - \frac{1}{2}(\partial\phi)^2 - V(\phi) - |\partial\chi|^2 + M^2|\chi|^2 - i\frac{\partial_\mu\phi}{f}\left(\chi\partial^\mu\chi^* - \chi^*\partial^\mu\chi\right) - \frac{1}{2}m^2(\chi^2 + \chi^{*2})\right]$$

- $V(\phi)$ is a slow-roll potential and only source of breaking of shift symmetry, $\phi \to \phi +$ const.
- Last term is the only one breaking U(1): $\chi \to \mathrm{e}^{i\alpha}\chi \quad \Rightarrow \quad$ the hierarchy $m^2 \ll M^2$ is radiatively stable.

The model

[Creminelli, Kumar, Salehian and LS '23]

- Let's set $\phi_0(t) = \rho f t$.
- The equations of motion for $\chi \to a^{3/2}\chi$ are (to leading order in slow-roll):

$$\ddot{\chi}_1 - \frac{\vec{\nabla}^2 \chi_1}{a^2} - \left(M^2 - m^2 + \frac{9H^2}{4} \right) \chi_1 - 2\rho\dot{\chi}_2 = 0\,,$$

$$\ddot{\chi}_2 - \frac{\vec{\nabla}^2 \chi_2}{a^2} - \left(M^2 + m^2 + \frac{9H^2}{4} \right) \chi_2 + 2\rho\dot{\chi}_1 = 0\,.$$

- We can decompose χ_1 and χ_2 as:

$$\chi_i(t,x) = \int \frac{\mathrm{d}^3k}{(2\pi)^3} \mathrm{e}^{i\vec{k}\cdot\vec{x}} \left[(F_k(t))_{ij}\hat{a}_j(\vec{k}) + (F_k^*(t))_{ij}\hat{a}_j^\dagger(-\vec{k}) \right] \,.$$

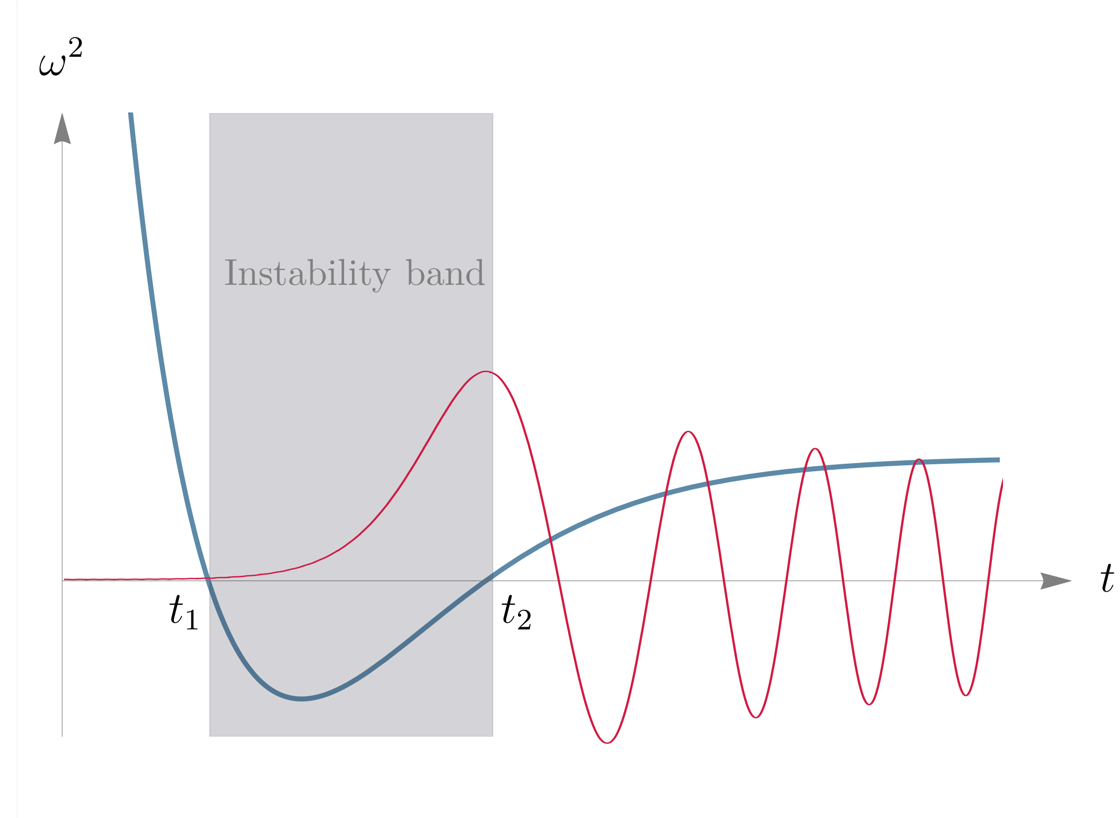
- The solution for $F_{ij}(t)$ can be obtained in WKB approximation: i.e., we assume that the coefficients in the e.o.m. depend so weakly on time that all the time dependence in $F_{ij}(t)$ is encoded in a common phase factor $\mathrm{e}^{i\int \omega_\pm \mathrm{d}t}$. [Weinberg '62]

The model

[Creminelli, Kumar, Salehian and LS '23]

- Two dispersion relations $\omega_\pm$, one for each additional degree of freedom.

- For given choices of the parameters, satisfying
$$m^2 < M^2 + \frac{9H^2}{4} < \rho^2 + \frac{m^4}{4\rho^2} < 2\rho^2 \,,$$
there is an instability ($\omega_-^2 < 0$) for physical momenta in the range:
$$-m^2 + M^2 + \frac{9}{4}H^2 < \frac{k^2}{a^2} < m^2 + M^2 + \frac{9}{4}H^2 \,.$$



- At t_2, the amplitude of the ω_- mode is enhanced by a factor $e^{\pi\xi}$, where
$$\xi \simeq \frac{m^4}{8H\rho M^2} \,.$$

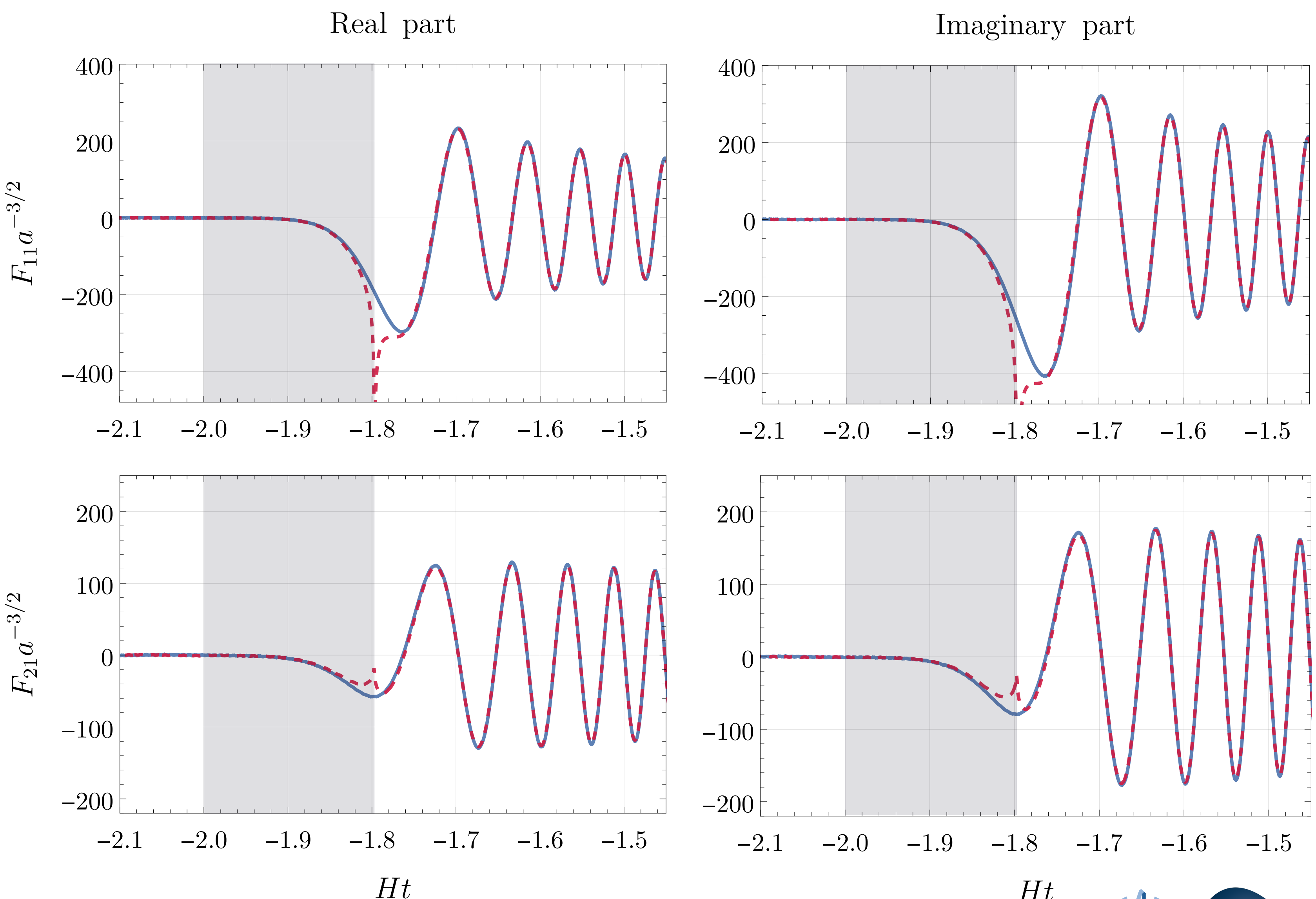
- The WKB approximation breaks down near the turning points t_1 and t_2.

WKB solution

[Creminelli, Kumar, Salehian and LS '23]

- The final WKB result is:

$$F_k(t) = \begin{cases} \left(\vec{Q}_+ \,,\, i\vec{Q}_+\right) e^{-i\int^t \omega_+} + \left(\vec{Q}_- \,,\, -i\vec{Q}_-\right) e^{-i\int^t \omega_-}, & t < t_1 \\ e^{-i\theta_1}\left(\vec{Q}_- \,,\, -i\vec{Q}_-\right) e^{\int_{t_1}^t |\omega_-|} & t_1 < t < t_2 \\ e^{-i\theta_1 + \pi\xi}\left[\left(\vec{Q}_- \,,\, -i\vec{Q}_-\right) e^{-i\int_{t_2}^t \omega_-} + i\left(\vec{Q}_-^* \,,\, -i\vec{Q}_-^*\right) e^{+i\int_{t_2}^t \omega_-}\right] & t > t_2 \end{cases}$$



$$\vec{Q}_\pm \equiv \frac{1}{2\sqrt{2}\left(\frac{k^2}{a^2} + \mu^2\right)^{1/4}} \begin{pmatrix} u_\pm \\ \mp \frac{i}{u_\pm} \end{pmatrix}$$

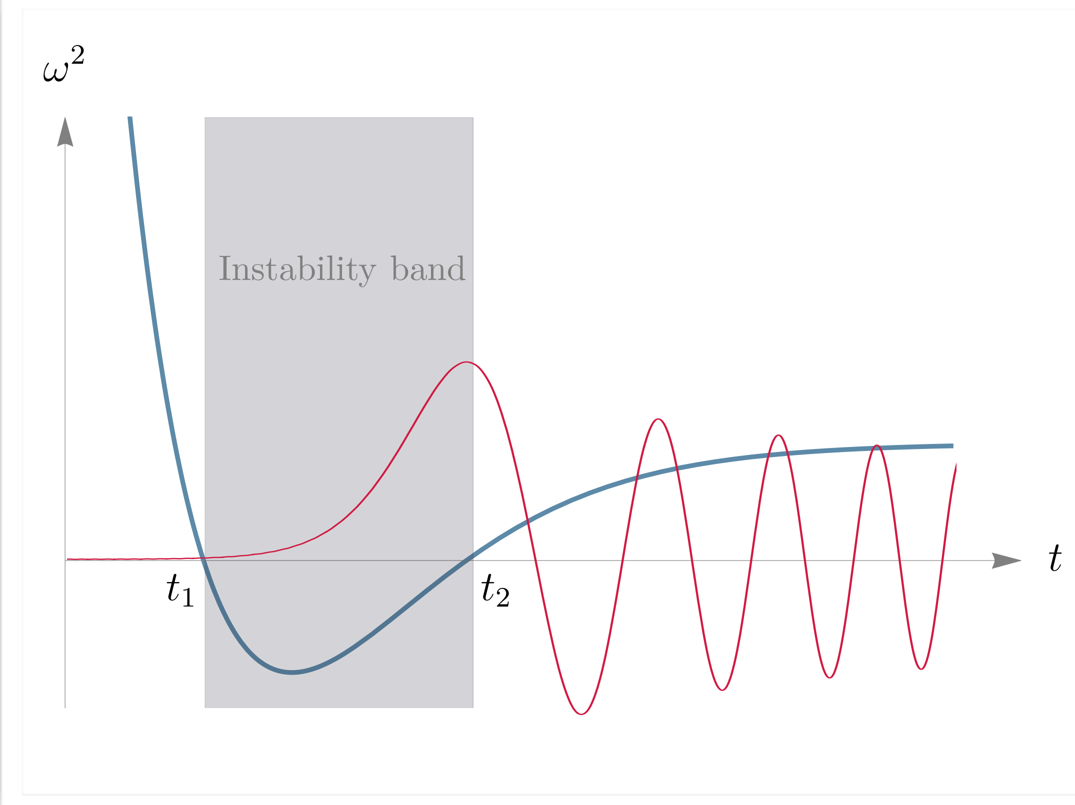
$$u_\pm \equiv \left[\frac{\omega_\pm}{\sqrt{\frac{k^2}{a^2} + \mu^2} \pm \left(\rho - \frac{m^2}{2\rho}\right)}\right]^{1/2}$$

$$\mu^2 \equiv \rho^2 - M^2 - \frac{9H^2}{4} + \frac{m^4}{4\rho^2}$$

Local response and hierarchies

[Creminelli, Kumar, Salehian and LS '23]

- Let us focus on the regime in which the χ-particle production takes place on parametrically sub-horizon scales.



- If:

$$H \ll m \ll M \lesssim \rho \ll f \ ,$$

then the instability window is narrow and localized on scales much shorter than Hubble i.e.,

$$\frac{k}{a} \sim M \gg H \ , \qquad H(t_2 - t_1) \sim \frac{m^2}{M^2} \ll 1 \ .$$

- This is a necessary condition in order to avoid non-local responses.

- Compare instead with $\Delta \mathscr{L} = \frac{\alpha}{f} \phi F \tilde{F}$ [Anber and Sorbo '09]: the instability occurs on scales comparable to H and additional resonances and instabilities are found. [Domcke et al '20], [Caravano et al '22], [Peloso and Sorbo '22]

- As a byproduct, the *local approximation* enables an analytic control over the dynamics of the perturbations.

Backreaction and slow-roll background evolution

[Creminelli, Kumar, Salehian and LS ’23]

- The exponentially amplified χ modes will eventually backreact on the inflationary evolution.

- At the background level, two main types of backreaction effects:
_ the large production of χ particles extracts energy from the inflaton, providing a new source of dissipation that can potentially overcome the Hubble friction in the ϕ_0 equation of motion;
_ the energy density of the produced particles contributes to the Friedmann equations.

- I will focus on the *large-backreaction* regime: dissipation due to χ production is comparable to, or larger than, the standard Hubble friction.
⇒ the evolution is dominated by dissipation and deviations from standard slow-roll are at least $O(1)$.

- This can allow for instance to have inflation on potentials that would otherwise be too steep to support slow-roll.

Background equations

[Creminelli, Kumar, Salehian and LS '23]

- To be as general as possible, we will allow m^2 and M^2 to depend on $X \equiv (\partial\phi)^2/(2\rho^2 f^2)$:

$$M^2(X) = M_0^2\left(c_0 + c_1\frac{(\partial\phi)^2}{\rho^2 f^2} + \dots\right), \qquad m^2(X) = m_0^2\left(c_0' + c_1'\frac{(\partial\phi)^2}{\rho^2 f^2} + \dots\right).$$

- The Friedmann equations are:

$$3M_{\rm Pl}^2H^2 = \frac{\dot\phi_0^2}{2} + V + \langle|\dot\chi|^2\rangle + \frac{1}{a^2}\langle|\partial_i\chi|^2\rangle - \left(M^2 - XM_X^2\right)\langle|\chi|^2\rangle + i\rho\langle\chi\dot\chi^* - \dot\chi\chi^*\rangle + \frac{1}{2}\left(m^2 - Xm_X^2\right)\langle\chi^2 + \chi^{*2}\rangle$$

$$-M_{\rm Pl}^2\dot H = \frac{\dot\phi_0^2}{2} + \langle|\dot\chi|^2\rangle + \frac{1}{3a^2}\langle|\partial_i\chi|^2\rangle + XM_X^2\langle|\chi|^2\rangle + i\rho\langle\chi\dot\chi^* - \dot\chi\chi^*\rangle - \frac{1}{2}Xm_X^2\langle\chi^2 + \chi^{*2}\rangle$$

- The inflaton's background equation is:

$$\frac{1}{a^3}\partial_t\left[\left(1 + \frac{\langle|\chi|^2\rangle}{\rho^2 f^2}(M_X^2 - 2\rho^2) - \frac{1}{2\rho^2 f^2}\langle\chi^2 + \chi^{*2}\rangle m_X^2\right)a^3\dot\phi_0\right] + V'(\phi) - \frac{im^2}{f}\langle\chi^2 - \chi^{*2}\rangle = 0$$

Background equations

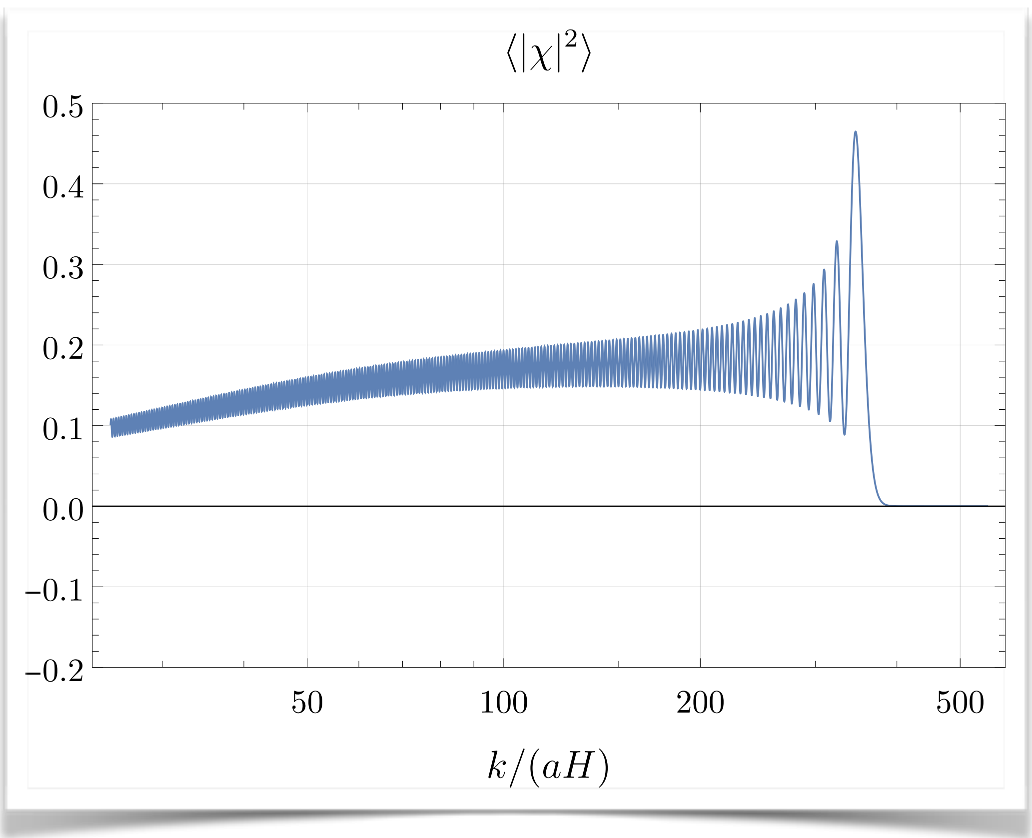
[Creminelli, Kumar, Salehian and LS '23]

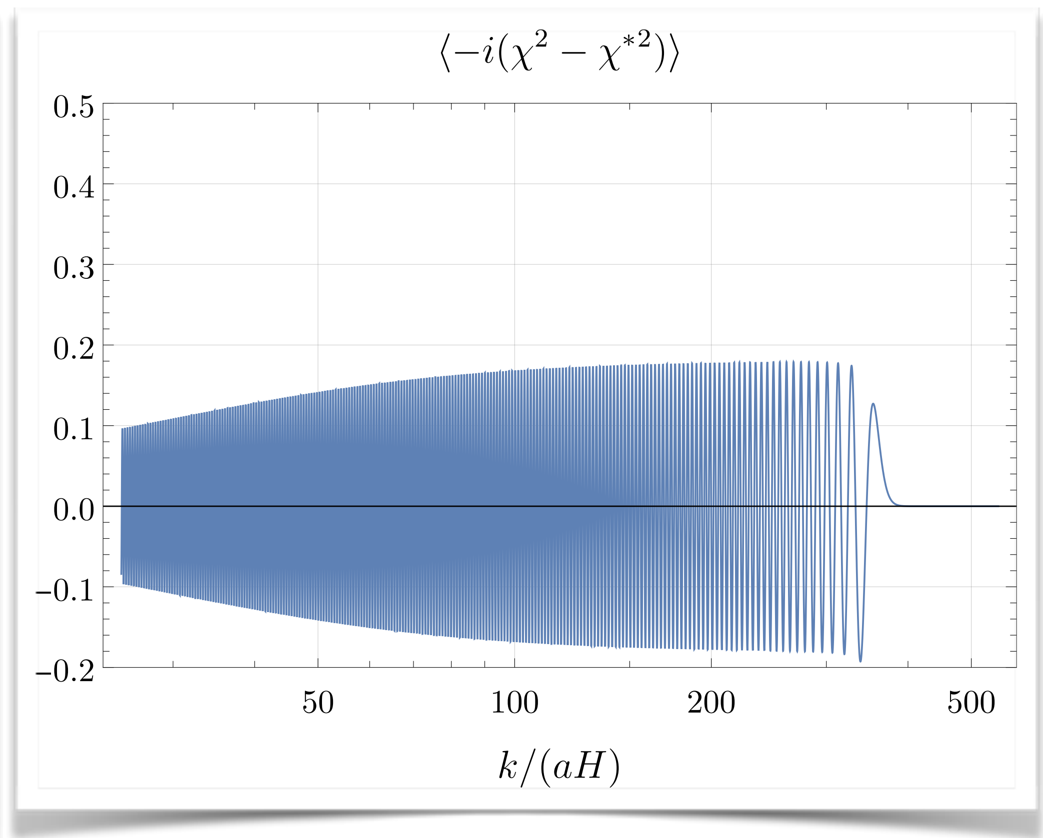
- The field expectation values can be computed using the WKB solution. For instance:

$$\langle |\chi|^2 \rangle = \frac{1}{2}\langle \chi_1^2 + \chi_2^2 \rangle = \frac{1}{4\pi^2 a^3}\int \mathrm{d}k\, k^2 \left([F_k(t)\cdot F_k^\dagger(t)]_{11} + [F_k(t)\cdot F_k^\dagger(t)]_{22} \right) .$$

- Using simple dimensional analysis:

$$\langle |\chi|^2 \rangle \simeq \frac{\mathrm{e}^{2\pi\xi}}{4\pi^2} M^2, \qquad \langle \chi^2 + \chi^{*2} \rangle \simeq \frac{\mathrm{e}^{2\pi\xi}}{4\pi^2} m^2, \qquad -i\langle \chi^2 - \chi^{*2} \rangle \simeq \frac{\mathrm{e}^{2\pi\xi}}{2\pi^2} m^2 .$$





- In the limit $\xi \gg 1$, $\langle \chi^2 - \chi^{*2} \rangle$ becomes the dominant term in the background equations.

Background equations

[Creminelli, Kumar, Salehian and LS '23]

- The background equations boil down to:

$$3M_{\rm Pl}^2H^2 \simeq \frac{\dot\phi_0^2}{2} + V + \frac{e^{2\pi\xi}}{4\pi^2}M^4\,, \qquad -M_{\rm Pl}^2\dot H \simeq \frac{\dot\phi_0^2}{2} + \frac{e^{2\pi\xi}}{4\pi^2}M^4\,,$$

$$3H\dot\phi_0 + V' + \frac{e^{2\pi\xi}}{2\pi^2}\frac{m^4}{f} \simeq 0\ .$$

has the same sign of $3H\dot\phi_0$ (friction)

- In the large-backreaction regime, $H\dot\phi_0 \lesssim e^{2\pi\xi}m^4/(2\pi^2 f)$:

$$2\pi\xi \simeq \ln\left|\frac{2\pi^2 f V'}{m^4}\right|\ .$$

- If the energy density is dominated by the inflaton potential:

$$M_{\rm Pl}^2H^2 \simeq V\,, \qquad \varepsilon \equiv -\frac{\dot H}{H^2} \ll 1 \qquad \text{(slow roll).}$$

- In order for the solution to be an attractor, we require ξ to be a monotonic function of $\dot\phi_0$ (this can be generically obtained by choosing $M^2(X)$ and $m^2(X)$).

Perturbations

[Creminelli, Kumar, Salehian and LS '23]

- Let us define $\varphi(t, \vec{x}) \equiv \phi(t, \vec{x}) - \phi_0(t)$ and expand the inflaton's equation of motion (in the decoupling limit):

$$\nabla_\mu \left[\left(1 + \frac{|\chi|^2}{\rho^2 f^2}(M_X^2 - 2\rho^2) - \frac{1}{2\rho^2 f^2}(\chi^2 + \chi^{*2}) m_X^2 \right) \nabla^\mu \phi \right] - V'(\phi) + \frac{i m^2}{f}(\chi^2 - \chi^{*2}) = 0 \, .$$

- The equation for φ is schematically of the form:

$$\mathcal{D}\varphi(t, \vec{x}) = \delta\mathcal{O}[\chi; t] \, ,$$

where $\mathcal{O} \supseteq |\chi|^2, (\chi^2 + \chi^{*2}), \ \partial_t |\chi|^2, \ \partial_t(\chi^2 + \chi^{*2}), \ (\chi^2 - \chi^{*2})$.

- For each operator, we distinguish:

$$\delta\mathcal{O} \equiv \mathcal{O} - \langle \mathcal{O} \rangle_{\varphi=0} = \delta\mathcal{O}_S + \delta\mathcal{O}_R \, .$$

$\delta\mathcal{O}_S$: represents intrinsic inhomogeneities in $\mathcal{O}$ (stochastic): $\delta\mathcal{O}_S \equiv [\mathcal{O} - \langle \mathcal{O} \rangle]_{\varphi=0}$.

$\delta\mathcal{O}_R$: is the response induced by the coupling to φ and is a functional of φ. At linear order: $\delta\mathcal{O}_R = \dfrac{\partial\langle \mathcal{O} \rangle}{\partial \dot{\phi}_0}\dot{\varphi} + \ldots$

Perturbations

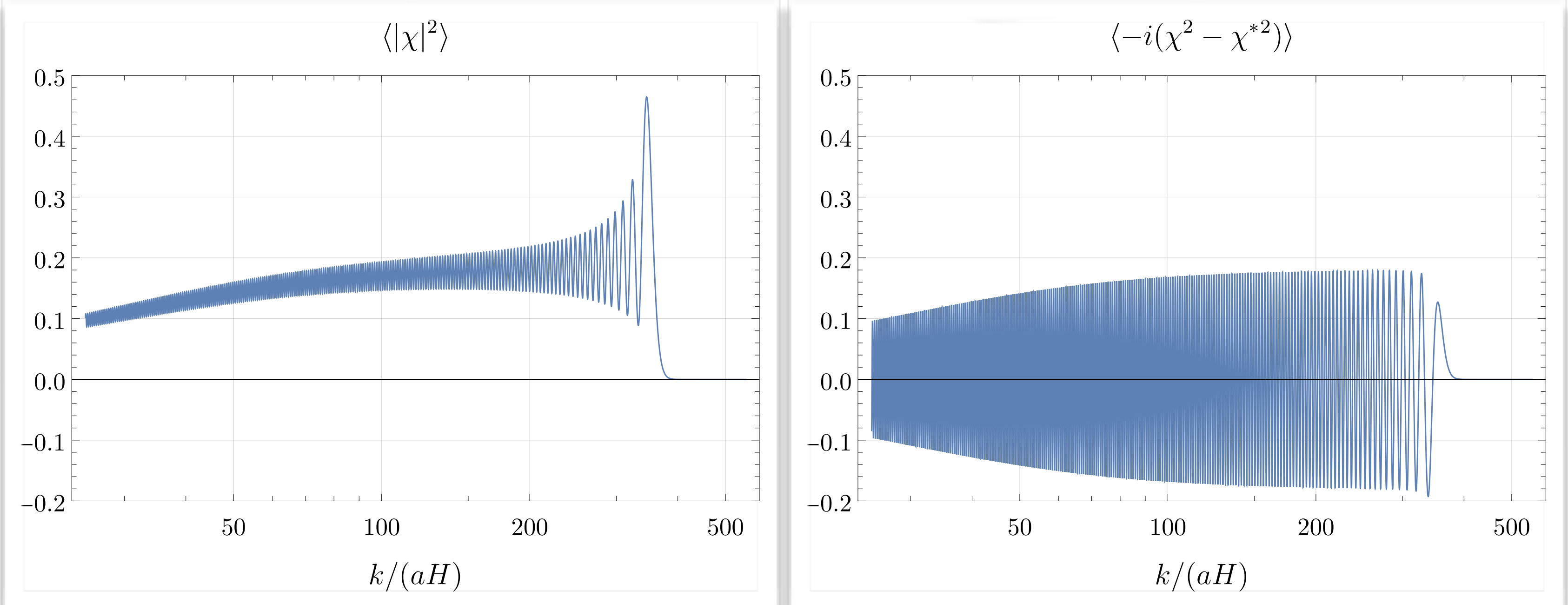
[Creminelli, Kumar, Salehian and LS '23]

- The type of response $\delta\mathcal{O}_R$ depends on the operator.
- Let us focus on:

$$|\chi|^2\,, \qquad (\chi^2 - \chi^{*2})\,.$$

$\langle |\chi|^2 \rangle$ is sensitive to a wide range of scales i.e., the response is *non-local*.

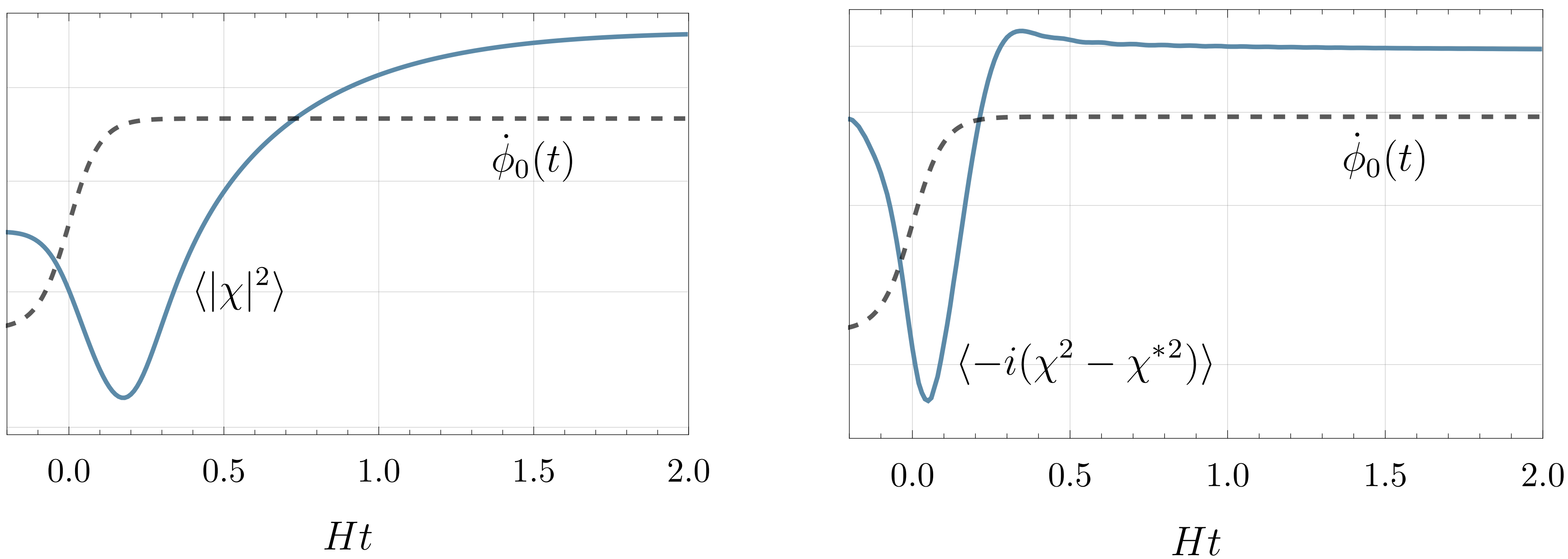
$\langle \chi^2 - \chi^{*2} \rangle$ gets its leading contribution from modes in a narrow range of scales around the instability window: the response is *local*.



- Non-local responses can lead to memory effects and unwanted large oscillations in the background solution.

Perturbations

[Creminelli, Kumar, Salehian and LS '23]



Local vs. non-local response: numerical solutions.

- Various ways to suppress the non-local operator $|\chi|^2$:

_ $2\pi\xi \gg 1$, in such a way that: $\dfrac{\frac{H\dot\phi_0}{f^2}\langle|\chi|^2\rangle}{\frac{im^2}{f}\langle\chi^2-\chi^{*2}\rangle} \simeq \dfrac{H\rho^3}{m^4} \simeq \dfrac{1}{8\xi} \ll 1$;

_ moderate `fine tuning': $M_X^2 = 2\rho^2$ removes $|\chi|^2$ (and $\partial_t|\chi|^2$) from the equation;

_ χ decays into an extra sector on time scales $\ll H^{-1}$.

Linearized equation

[Creminelli, Kumar, Salehian and LS '23]

- In the local approximation the φ equation reads:

$$\ddot{\varphi}_k + (3H+\gamma)\dot{\varphi}_k + \left(\frac{k^2}{a^2} + V''\right)\varphi_k = -\frac{m^2}{f}\delta\mathcal{O}_S(k)\,,$$

where $\mathcal{O} \equiv -i(\chi^2 - \chi^{*2})$ and

$$\gamma \sim \frac{m^2}{f}\frac{\partial\langle\mathcal{O}\rangle}{\partial\dot{\phi}_0} \sim \frac{\xi m^4}{\pi M f^2}\mathrm{e}^{2\pi\xi}\,.$$

- Two scales in the problem:
  _ $k/a \sim \gamma$ (friction becomes dominant);
  _ $k/a \sim \sqrt{\gamma H} \gg H$ (freeze-out).

- The solution for φ_k is the superposition of homogeneous (vacuum fluctuations) and particular solutions.

- Vacuum fluctuations are subleading in the large-friction limit $\gamma \gg H$. [Lopez Nacir et al '11]
  I will focus on the particular solution.

Power spectrum

Stochastic noise and induced power spectrum

[Creminelli, Kumar, Salehian and LS '23]

- The particular solution to the inhomogeneous equation can be derived using standard Green's function methods:

$$\varphi(\tau,k) = -\frac{m^2}{f}\int \mathrm{d}\tau' G_k(\tau,\tau')a(\tau')^2\delta\mathcal{O}_S(\tau',k)\ .$$

- The inflaton two-point function is:

$$\langle\varphi_k(\tau)\varphi_{k'}(\tau)\rangle = \frac{m^4}{f^2}\int \mathrm{d}\tau'\mathrm{d}\tau'' a(\tau')^2 a(\tau'')^2\, G_k(\tau,\tau')G_{k'}(\tau,\tau'')\langle\delta\mathcal{O}_S(\tau',k)\delta\mathcal{O}_S(\tau'',k')\rangle\ .$$

- In the local approximation, the noise two-point function $\langle\delta\mathcal{O}_S(\tau',k)\delta\mathcal{O}_S(\tau'',k')\rangle$ is proportional to a delta-function $\delta(\tau-\tau')$ (*locality in time*) and it is independent of the spatial momentum (*locality in space*):

$$\langle\delta\mathcal{O}_S(\tau',k)\delta\mathcal{O}_S(\tau'',k')\rangle \simeq (2\pi)^3\delta(\vec{k}+\vec{k}')\,\delta(\tau'-\tau'')\,H^4\tau^4\nu_{\mathcal{O}}\ , \qquad \nu_{\mathcal{O}} \simeq \frac{\mathrm{e}^{4\pi\xi}}{4\pi^2}\frac{M}{m}.$$

Stochastic noise and induced power spectrum

[Creminelli, Kumar, Salehian and LS '23]

- The power spectrum for $\zeta = -H\varphi/\dot{\phi}_0$ is:

$$\Delta_\zeta^2 \equiv \frac{k^3}{2\pi^2}\frac{H^2}{\dot{\phi}_0^2}\langle\varphi_k\varphi_{-k}\rangle' \simeq \frac{1}{32\xi^2}\left(\frac{\gamma}{\pi H}\right)^{3/2}\frac{H^4 M}{m^5} \ .$$

- The smallness of the power spectrum is due to the presence of many independent sources that contribute incoherently on sub-Hubble scales.
[In contrast, the case of axion inflation in the strong-backreaction regime predicts large power for the fluctuations simply because the instability continues up to large scales.]

- For the same reason, the large number also makes the statistics of perturbations close to Gaussian (by the central limit theorem) and thus compatible with observations.

Non-Gaussianity

Non-Gaussianity

[Creminelli, Kumar, Salehian and LS '23]

- In the local approximation, expanding the φ equation to quadratic order in perturbations:

$$-\nabla_\mu \nabla^\mu \varphi + V''\varphi + \frac{1}{2}V'''\varphi^2 + \frac{1}{\rho f^2}\left[m_X^2\dot{\varphi} + \left(m_X^2 + m_{XX}^2\right)\frac{\dot{\varphi}^2}{2\rho f} - m_X^2\frac{(\partial_i\varphi)^2}{2\rho f a^2}\right]\langle\mathcal{O}\rangle + \frac{1}{f}\left(m^2 + m_X^2\frac{\dot{\varphi}}{\rho f}\right)\left(\delta\mathcal{O}_R + \delta\mathcal{O}_S\right) = 0$$

- Three different sources of non-Gaussianity:
  1. inflaton's self interactions, $V'''\varphi^2$ (slow-roll suppressed);
  2. non-Gaussianity induced by the statistics of the noise fluctuation $\delta\mathcal{O}_S$:

$$\langle\varphi_{k_1}\varphi_{k_2}\varphi_{k_3}\rangle = -\left(\frac{m^2}{f}\right)^3\int \mathrm{d}\tau_1\mathrm{d}\tau_2\mathrm{d}\tau_3 a(\tau_1)^2 a(\tau_2)^2 a(\tau_3)^2 G_{k_1}(0,\tau_1)G_{k_2}(0,\tau_2)G_{k_3}(0,\tau_3)\langle\delta\mathcal{O}_S(\tau_1,k_1)\delta\mathcal{O}_S(\tau_2,k_2)\delta\mathcal{O}_S(\tau_3,k_3)\rangle$$

$$\Rightarrow \qquad f_{\rm NL}^{\rm eq} \simeq \frac{40\pi}{9}\xi\frac{m^2}{M^2} \lesssim 1$$

  3. non-Gaussianity sourced by non-linear coupling between φ and χ:
     (i) by direct coupling between φ and $\delta\mathcal{O}_S$;
     (ii) from non-linear dependence of $\delta\mathcal{O}_R$ on φ.

Since $\dot{\phi}_0(t)$ is the only source of breaking of Lorentz symmetries, non-linearly realized general covariance and shift symmetry ensure that $\delta\mathcal{O}_R$ depends on φ only through the Lorentz scalar $\partial_\mu\phi\partial^\mu\phi$.

Non-Gaussianity from non-linear coupling

[Creminelli, Kumar, Salehian and LS '23]

- The quadratic equation of motion for φ in the local approximation:

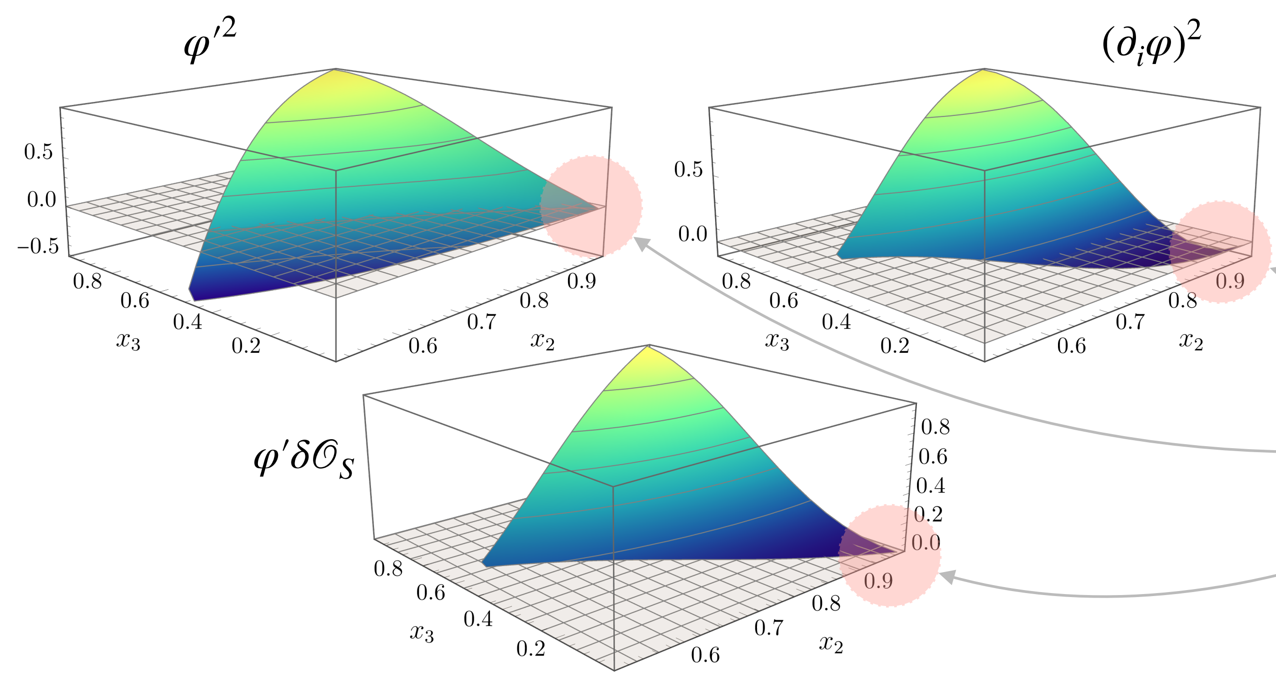
$$\varphi'' + (2H+\gamma)a\varphi' - \vec{\nabla}^2\varphi + a^2 V''\varphi \simeq -a^2\frac{m^2}{f}\delta\mathcal{O}_S + \frac{\gamma}{2\rho f}\left[(\partial_i\varphi)^2 - 2\pi\xi\varphi'^2\right] - a\frac{2\pi\xi m^2}{\rho f^2}\varphi'\delta\mathcal{O}_S\,.$$

- Neglecting the homogeneous (i.e., vacuum) solution,

$$\varphi = \varphi^{(1)} + \varphi^{(2)}$$

$$\varphi_k^{(2)}(\tau) = -\frac{1}{\rho f}\int \mathrm{d}\tau' G_k(\tau,\tau')\int\frac{\mathrm{d}^3p}{(2\pi)^3}\left[\frac{\gamma}{2}\,\vec{p}\cdot(\vec{k}-\vec{p})\,\varphi_p^{(1)}(\tau')\varphi_{k-p}^{(1)}(\tau') + \pi\xi\gamma\,\partial_{\tau'}\varphi_p^{(1)}(\tau')\partial_{\tau'}\varphi_{k-p}^{(1)}(\tau') + a(\tau')\frac{2\pi\xi m^2}{f}\partial_{\tau'}\varphi_{k-p}^{(1)}(\tau')\delta\mathcal{O}_S(\tau',p)\right]$$

- Non-gaussianity shapes peak in the *equilateral* configuration:



Vanishing in the squeezed limit $k_3 \ll k_1, k_2$, consistently with the fact that there is an *attractor* solution for the dynamics which is controlled by a *single clock*, $\phi_0(t)$.

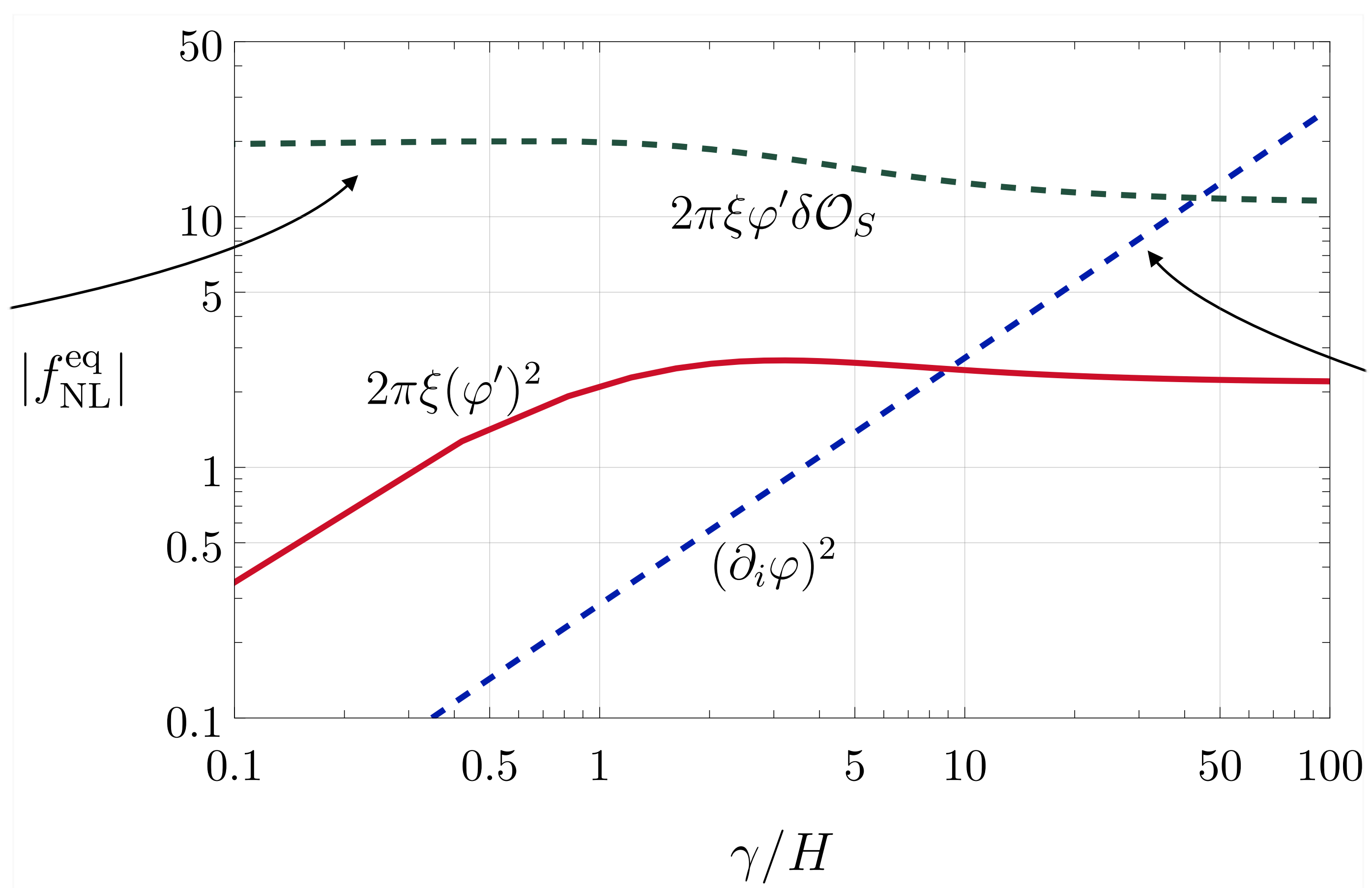
Non-Gaussianity from non-linear coupling

[Creminelli, Kumar, Salehian and LS '23]

- Amplitudes of non-Gaussianity from non-linear coupling:

The interaction of the inflaton with the noise can generate $f_{\rm NL}^{\rm eq} \simeq$ few $\times\, O(10)$.

As opposed to the blue curve, this is *model-dependent*.



The linear growth in $f_{\rm NL}^{\rm eq} \simeq -\frac{\gamma}{4H}$ is a consequence of the non-linearly realized Lorentz symmetries and is *model-independent*.

This is similar to what happens when considering models where the inflaton has a reduced speed of propagation, where $f_{\rm NL}^{\rm eq} \propto \frac{1}{c_s^2}$.

Conclusions

Summary

- Model of dissipative inflation via scalar particle production.
- As opposed to previous implementations based on a coupling to gauge fields, in our model particle production takes place on parametrically sub-horizon scales:
  _ avoids non-local response and memory effects;
  _ good analytic control on the dynamics and the predictions;
  _ makes the statistics of the perturbations naturally close to Gaussian, by virtue of the central limit theorem and the large occupation on short scales.
- Robust against radiative corrections because of (approximate) $U(1)$ and shift symmetries.
- First robust explicit realization of the EFT of dissipative inflation. [Lopez Nacir et al '11]
- Proof of concept that dissipative inflation is not an exclusive feature of couplings to gauge fields, but can be realized more in general.

Open directions

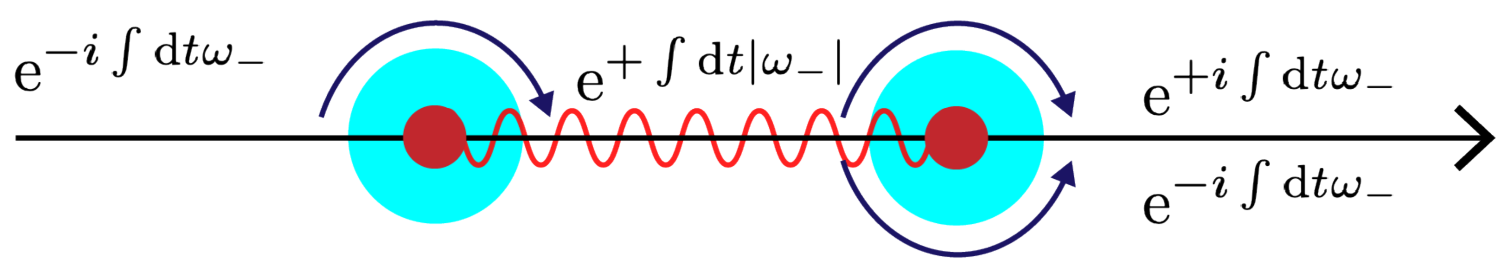
- Tensor modes
- Primordial black holes
- Thermalization and warm inflation
- Fermions

Backup slides

WKB solution

[Creminelli, Kumar, Salehian and LS '23]

- To find the full solution for $F_{ij}(t)$:
  _ we first solve with WKB in the regions far from the turning points ($t \ll t_1$, $t_1 \ll t \ll t_2$ and $t \gg t_2$) where $|\dot{\omega}\_/\omega\_^2| \ll 1$;
  _ we perform an analytical continuation and match the coefficients:



[Landau and Lifshits, vol. 3]
[Dufaux et al '06]

Vacuum fluctuations

- The homogeneous solution to the linearized φ equation:
$$\ddot{\varphi}_k + (3H+\gamma)\dot{\varphi}_k + \left(\frac{k^2}{a^2} + V''\right)\varphi_k = -\frac{m^2}{f}\delta\mathcal{O}_S(k)$$
is
$$\varphi_k \simeq \tau^\alpha \left[A\,\mathrm{H}_\alpha^{(1)}(-k\tau) + B\,\mathrm{H}_\alpha^{(2)}(-k\tau)\right], \qquad \alpha \equiv \frac{3}{2} + \frac{\gamma}{2H}\,.$$
(deviations from scale invariance neglected.)

- Demanding the correct Bunch–Davies initial condition at $k|\tau_0| \gg \gamma/H$ yields
$$A = 0, \qquad B \sim \tau_0^{-\gamma/2H}\,.$$

- The solution at late times ($\tau \to 0$) is
$$\varphi_k(\tau \to 0) \propto \left(\frac{k\tau_0}{\gamma/H}\right)^{-\frac{\gamma}{2H}} \mathrm{e}^{-\frac{\gamma}{2H}}\,.$$

- Increasing the friction will exponentially damp the homogeneous solution.

Non-Gaussianity from non-linear coupling

[Creminelli, Kumar, Salehian and LS '23]

3. non-Gaussianity sourced by non-linear coupling between φ and χ:

(i) by direct coupling between φ and $\delta\mathcal{O}_S$;

(ii) from non-linear dependence of $\delta\mathcal{O}_R$ on φ:

$$\delta\mathcal{O}_R = \mathcal{O}[\sqrt{-\partial_\mu\phi\partial^\mu\phi}] - \mathcal{O}\,|_{\varphi=0}$$

$$= \mathcal{O}[\sqrt{-\partial_\mu\phi\partial^\mu\phi}] - \langle\mathcal{O}\rangle|_{\varphi=0} - \delta\mathcal{O}_S$$

$$= \delta\langle\mathcal{O}\rangle\left(\sqrt{-\partial_\mu\phi\partial^\mu\phi}\right) + \frac{1}{2\langle(\mathcal{O}-\langle\mathcal{O}\rangle)^2\rangle}\frac{\partial\langle(\mathcal{O}-\langle\mathcal{O}\rangle)^2\rangle}{\partial\dot\phi_0}\dot\varphi\delta\mathcal{O}_S + \ldots$$

$$= \frac{\partial\langle\mathcal{O}\rangle}{\partial\dot\phi_0}\left(\dot\varphi - \frac{(\partial_i\varphi)^2}{2\rho f a^2}\right) + \frac{1}{2}\frac{\partial^2\langle\mathcal{O}\rangle}{\partial\dot\phi_0^2}\dot\varphi^2 + \frac{1}{2\nu_\mathcal{O}}\frac{\partial\nu_\mathcal{O}}{\partial\dot\phi_0}\dot\varphi\delta\mathcal{O}_S + \ldots$$

Since $\dot\phi_0(t)$ is the only source of breaking of Lorentz symmetries, non-linearly realized general covariance and shift symmetry ensure that $\delta\mathcal{O}_R$ depends on φ only through the Lorentz scalar $\partial_\mu\phi\partial^\mu\phi$.

At this order, $\mathcal{O}$ can be taken Gaussian. A change in $\dot\phi_0$ induces a variation in the mean and the variance.