

The effective field theory of inflation

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The effective field theory of inflation

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ABSTRACT: We study the effective field theory of inflation, i.e. the most general theory describing the fluctuations around a quasi de Sitter background, in the case of single field models. The scalar mode can be eaten by the metric by going to unitary gauge. In this gauge, the most general theory is built with the lowest dimension operators invariant under spatial diffeomorphisms, like g^{00} and $K_{\mu\nu}$, the extrinsic curvature of constant time surfaces. This approach allows us to characterize all the possible high energy corrections to simple slow-roll inflation, whose sizes are constrained by experiments. Also, it describes in a common language all single field models, including those with a small speed of sound and Ghost Inflation, and it makes explicit the implications of having a quasi de Sitter background. The non-linear realization of time diffeomorphisms forces correlation among different observables, like a reduced speed of sound and an enhanced level of non-Gaussianity.

KEYWORDS: Spontaneous Symmetry Breaking, Space-Time Symmetries, Gauge Symmetry, Cosmology of Theories beyond the SM.

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1. Introduction

The effective field theory approach, i.e. the description of a system through the lowest dimension operators compatible with the underlying symmetries, has been very fruitful in many areas, from particle physics to condensed matter. The purpose of this paper is to apply this methodology to describe the theory of fluctuations around an inflating cosmological background.

The usual way to study a single field inflationary model is to start from a Lagrangian for a scalar field ϕ and solve the equation of motion for ϕ together with the Friedmann equations for the FRW metric. We are interested in an inflating solution, i.e. an accelerated expansion with a slowly varying Hubble parameter, with the scalar following an homogeneous time-dependent solution $\phi_0(t)$. At this point one studies perturbations around this background solution to work out the predictions for the various cosmological observables.

The theory of perturbations around the time evolving solution is quite different from the theory of ϕ we started with: while ϕ is a scalar under all diffeomorphisms (diffs), the perturbation $\delta\phi$ is a scalar only under spatial diffs while it transforms non-linearly with respect to time diffs:

$$t \rightarrow t + \xi^0(t, \vec{x}) \quad \delta\phi \rightarrow \delta\phi + \dot{\phi}_0(t)\xi^0 . \quad (1.1)$$

In particular one can choose a gauge $\phi(t, \vec{x}) = \phi_0(t)$ where there are no inflaton perturbations, but all degrees of freedom are in the metric. The scalar variable $\delta\phi$ has been eaten by the graviton, which has now three degrees of freedom: the scalar mode and the two tensor helicities. This phenomenon is analogous to what happens in a spontaneously broken gauge theory. A Goldstone mode, which transforms non-linearly under the gauge symmetry, can be eaten by the gauge boson (unitary gauge) to give a massive spin 1 particle. The non-linear sigma model of the Goldstone can be embedded and UV completed into a linear representation of the gauge symmetry like in the Higgs sector of the Standard Model. This is analogous to the standard formulation of inflation, where we start from a Lagrangian for ϕ with a linear representation of diffs. In this paper we want to stress the alternative point of view, describing the theory of perturbations during inflation directly around the time evolving vacuum where time diffs are non-linearly realized. This formalism has been firstly introduced, for a generic FRW background, in [1] to study the possibility of violating the Null Energy Condition; here we will extend this formalism focusing on an inflationary solution.

We will show that in unitary gauge the most generic Lagrangian with broken time diffeomorphisms (but unbroken spatial diffs) describing perturbations around a flat FRW with Hubble rate $H(t)$ is given by

$$S = \int d^4x \sqrt{-g} \left[\frac{1}{2} M_{\text{Pl}}^2 R + M_{\text{Pl}}^2 \dot{H} g^{00} - M_{\text{Pl}}^2 (3H^2 + \dot{H}) + \frac{M_2(t)^4}{2!} (g^{00} + 1)^2 \right. \\ \left. + \frac{M_3(t)^4}{3!} (g^{00} + 1)^3 + \dots - \frac{\bar{M}_2(t)^2}{2} \delta K^\mu{}_\mu{}^2 + \dots \right]. \quad (1.2)$$

The first two operators after the Einstein-Hilbert term are fixed by the requirement of having a given unperturbed solution $H(t)$, while all the others are free and parametrize all the possible different theories of perturbations with the same background solution. As time diffs are broken one is allowed to write any term that respects spatial diffs, including for example g^{00} and the extrinsic curvature $K^\mu{}_\nu$ of the surfaces at constant time. The coefficients of the operators will be in general time dependent. The reader may be worried by the use of a Lagrangian that is not invariant under diffeomorphisms. But clearly diff. invariance can be restored as in a standard gauge theory. One performs a time-diffeomorphism with parameter $\xi^0(t, \vec{x})$ and promotes the parameter to a field $\pi(t, \vec{x})$ which shifts under time diffs: $\pi(t, \vec{x}) \rightarrow \pi(t, \vec{x}) - \xi^0(t, \vec{x})$. The scalar π is the Goldstone mode which non linearly realizes the time diffs and it describes the scalar perturbations around the FRW solution.

It is well known that the physics of the longitudinal components of massive gauge bosons can be studied, at sufficiently high energy, concentrating on the scalar Goldstone mode (equivalence theorem). The same is true in our case: for sufficiently high energy the mixing with gravity is irrelevant and we can concentrate on the Goldstone mode. In this regime the physics is very transparent and most of the information about cosmological perturbations can be obtained. Performing the broken diff transformation on the

Lagrangian (1.2) and concentrating on the Goldstone mode π one gets

$$S_\pi = \int d^4x \sqrt{-g} \left[M_{\text{Pl}}^2 \dot{H} (\partial_\mu \pi)^2 + 2M_2^4 \left(\dot{\pi}^2 + \dot{\pi}^3 - \dot{\pi} \frac{1}{a^2} (\partial_i \pi)^2 \right) - \frac{4}{3} M_3^4 \dot{\pi}^3 - \frac{\bar{M}^2}{2} \frac{1}{a^4} (\partial_i^2 \pi)^2 + \dots \right]. \quad (1.3)$$

Every invariant operator in unitary gauge is promoted to a (non-linear) operator for the Goldstone: the non-linear realization of diff invariance forces the relation among various terms.

Let us briefly point out what are the advantages of this approach before moving to a systematic construction of the theory.

- Starting from a “vanilla” scenario of inflation with a scalar field with minimal kinetic term and slow-roll potential, we have parameterized our ignorance about all the possible high energy effects in terms of the leading invariant operators. Experiments will put bounds on the various operators, for example with measurements of the non-Gaussianity of perturbations and studying the deviation from the consistency relation for the gravitational wave tilt. In some sense this is similar to what one does in particle physics, where one puts constraints on the size of the operators that describe deviations from the Standard Model and thus encode the effect of new physics.
- It is explicit what is forced by the symmetries and by the requirement of an inflating background and what is free. For example eq. (1.3) shows that the spatial kinetic term $(\nabla\pi)^2$ is proportional to \dot{H} , while the time kinetic term $\dot{\pi}^2$ is free. Another example is the unitary gauge operator $(g^{00} + 1)^2$. Once written in terms of the Goldstone π , this gives a quadratic term $\dot{\pi}^2$, which reduces the speed of sound of π excitations, and a cubic term $\dot{\pi}(\nabla\pi)^2$, which increases the interaction among modes, i.e. the non-Gaussianity. Therefore, barring cancellations with other operators, a reduced speed of sound is related by symmetry to an enhanced non-Gaussianity. Notice moreover that the coefficient of this operator is constrained to be positive, to avoid propagation of π excitations out of the lightcone.
- One knows all the possible operators. For example, at the leading order in derivatives, the interaction among three π modes can be changed by $(g^{00} + 1)^2$ and $(g^{00} + 1)^3$. This will correspond to two different shapes of the 3-point function which can be in principle experimentally distinguished to fix the size of each operator.
- All the possible single field models are now unified. For example there has been interest in models with a modified Lagrangian $L((\partial\phi)^2, \phi)$, like DBI inflation [2–6] which have rather peculiar predictions. In our language these correspond to the case in which the operators $(g^{00} + 1)^n$ are large. Another interesting limit is when $\dot{H} \rightarrow 0$; in this case the leading spatial kinetic term is coming from the operator proportional to \bar{M}^2 and it is of the form $(\nabla^2\pi)^2$. This limit describes Ghost Inflation [7].

- In the ϕ language one can perform a field redefinition $\phi \rightarrow \tilde{\phi}(\phi)$. It is true that the resulting Lagrangian will describe the same physics, but this is not obvious. A simple example is given by the Lagrangian

$$f(\phi)^2(\partial\phi)^2 - V(\phi), \quad (1.4)$$

where f is a generic function. This is equivalent to a Lagrangian with minimal kinetic term and a different potential through the field redefinition $\tilde{\phi}(\phi)$, $d\tilde{\phi}/d\phi = f(\phi)$. However the equivalence among different Lagrangians becomes more complicated when we consider more general terms. On the other hand this ambiguity is absent at the level of π , which realizes a sort of standard non-linear representation of time diffs.

- In the ϕ language is it not obvious how to assess the importance of an operator for the study of perturbations, because some of the legs of an operator may be evaluated on the background solution. For example in a theory with all operators of the form $(\partial\phi)^{2n}$, all of them may have the same importance if the background velocity $\dot{\phi}_0$ is large enough, as it happens in DBI inflation. On the other hand the usual way of estimating the importance of an operator works in the π language. Even more clear is the case of Ghost Inflation where, given the non-relativistic dispersion relation for π the scaling of operators is clear only in the π language.
- The parametrization of the operators directly around the solution is crucial if one calculates loop corrections of cosmological perturbations. A diagram with a given number of external legs will in general contain a UV divergence. This is easy to renormalize in the Lagrangian (1.2), because there is only a finite number of terms which describe the interaction among n perturbations. On the other hand at the level of the ϕ Lagrangian, there is an *infinite* number of operators contributing to the interaction among n perturbations. For each operator in fact one can put many of its legs on the background, so that the relation among an operator and a diagram for perturbations is rather obscure.

2. Construction of the action in unitary gauge

Inflation is a period of accelerated cosmic expansion with an approximately constant Hubble parameter. This quasi de Sitter background has a privileged spatial slicing, given by a physical clock which allows to smoothly connect to a decelerated hot Big Bang evolution. The slicing is usually realized by a time evolving scalar $\phi(t)$. Another example one may keep in mind is given by a perfect fluid.¹ To describe perturbations around this solution one can choose a gauge where the privileged slicing coincides with surfaces of constant t , i.e. $\delta\phi(\vec{x}, t) = 0$. In this gauge there are no explicit scalar perturbations, but only metric fluctuations. As time diffeomorphisms have been fixed and are not a gauge symmetry

¹Indeed, as shown for example in [8], non-vorticious excitations of a perfect fluid may be described by a derivatively coupled scalar.

anymore, the graviton now describes three degrees of freedom: the scalar perturbation has been eaten by the metric.

What is the most general Lagrangian in this gauge? One must write down operators that are functions of the metric $g_{\mu\nu}$, and that are invariant under the (linearly realized) time dependent spatial diffeomorphisms $x^i \rightarrow x^i + \xi^i(t, \vec{x})$. Spatial diffeomorphisms are in fact unbroken. Besides the usual terms with the Riemann tensor, which are invariant under all diffs, many extra terms are now allowed, because of the reduced symmetry of the system. They describe the additional degree of freedom eaten by the graviton. For example it is easy to realize that g^{00} is a scalar under spatial diffs, so that it can appear freely in the unitary gauge Lagrangian. Polynomials of g^{00} are the only terms without derivatives. Given that there is a preferred slicing of the spacetime, one is also allowed to write geometric objects describing this slicing. For instance the extrinsic curvature $K_{\mu\nu}$ of surfaces at constant time is a tensor under spatial diffs and it can be used in the action. Notice that generic functions of time can multiply any term in the action. In appendix A we prove that the most generic Lagrangian can be written as

$$S = \int d^4x \sqrt{-g} \left[\frac{1}{2} M_{\text{Pl}}^2 R - c(t) g^{00} - \Lambda(t) + \frac{1}{2!} M_2(t)^4 (g^{00} + 1)^2 + \frac{1}{3!} M_3(t)^4 (g^{00} + 1)^3 + \right. \\ \left. - \frac{\bar{M}_1(t)^3}{2} (g^{00} + 1) \delta K^\mu{}_\mu - \frac{\bar{M}_2(t)^2}{2} \delta K^\mu{}_\mu{}^2 - \frac{\bar{M}_3(t)^2}{2} \delta K^\mu{}_\nu \delta K^\nu{}_\mu + \dots \right], \quad (2.1)$$

where the dots stand for terms which are of higher order in the fluctuations or with more derivatives. We denote by $\delta K_{\mu\nu}$ the variation of the extrinsic curvature of constant time surfaces with respect to the unperturbed FRW: $\delta K_{\mu\nu} = K_{\mu\nu} - a^2 H h_{\mu\nu}$ with $h_{\mu\nu}$ is the induced spatial metric. Notice that only the first three terms in the action above contain linear perturbations around the chosen FRW solution, all the others are explicitly quadratic or higher. Therefore the coefficients $c(t)$ and $\Lambda(t)$ will be fixed by the requirement of having a given FRW evolution $H(t)$, i.e. requiring that tadpole terms cancel around this solution. Before fixing these coefficients, it is important to realize that this simplification is not trivial. One would expect that there are an infinite number of operators which give a contribution at first order around the background solution. However one can write the action as a polynomial of linear terms like $\delta K_{\mu\nu}$ and $g^{00} + 1$, so that it is evident whether an operator starts at linear, quadratic or higher order. All the linear terms besides the ones in eq. (2.1) will contain derivatives and they can be integrated by parts to give a combination of the three linear terms we considered plus covariant terms of higher order. This construction is explicitly carried out in appendix B. We conclude that *the unperturbed history fixes $c(t)$ and $\Lambda(t)$, while the difference among different models will be encoded into higher order terms.*

We can now fix the linear terms imposing that a given FRW evolution is a solution. As we discussed, the terms proportional to c and Λ are the only ones that give a stress energy tensor

$$T_{\mu\nu} = - \frac{2}{\sqrt{-g}} \frac{\delta S_{\text{matter}}}{\delta g^{\mu\nu}} \quad (2.2)$$

which does not vanish at zeroth order in the perturbations and therefore contributes to the right hand side of the Einstein equations. During inflation we are mostly interested in

a flat FRW Universe (see appendix B for the general case)

$$ds^2 = -dt^2 + a^2(t)d\vec{x}^2 \tag{2.3}$$

so that Friedmann equations are given by

$$H^2 = \frac{1}{3M_{\text{Pl}}^2} [c(t) + \Lambda(t)] \tag{2.4}$$

$$\frac{\ddot{a}}{a} = \dot{H} + H^2 = -\frac{1}{3M_{\text{Pl}}^2} [2c(t) - \Lambda(t)] . \tag{2.5}$$

Solving for c and Λ we can rewrite the action (2.1) as

$$\begin{aligned} S = \int d^4x \sqrt{-g} & \left[\frac{1}{2} M_{\text{Pl}}^2 R + M_{\text{Pl}}^2 \dot{H} g^{00} - M_{\text{Pl}}^2 (3H^2 + \dot{H}) + \frac{1}{2!} M_2(t)^4 (g^{00} + 1)^2 \right. \\ & + \frac{1}{3!} M_3(t)^4 (g^{00} + 1)^3 + \left. -\frac{\bar{M}_1(t)^3}{2} (g^{00} + 1) \delta K^\mu{}_\mu - \frac{\bar{M}_2(t)^2}{2} \delta K^\mu{}_\mu{}^2 \right. \\ & \left. - \frac{\bar{M}_3(t)^2}{2} \delta K^\mu{}_\nu \delta K^\nu{}_\mu + \dots \right]. \tag{2.6} \end{aligned}$$

As we said all the coefficients of the operators in the action above may have a generic time dependence. However we are interested in solutions where H and \dot{H} do not vary significantly in one Hubble time. Therefore it is natural to assume that the same holds for all the other operators. With this assumption the Lagrangian is approximately time translation invariant.² Therefore the time dependence generated by loop effects will be suppressed by a small breaking parameter.³ This assumption is particularly convenient since the rapid time dependence of the coefficients can win against the friction created by the exponential expansion, so that inflation may cease to be a dynamical attractor, which is necessary to solve the homogeneity problem of standard FRW cosmology.

It is important to stress that this approach does describe the most generic Lagrangian not only for the scalar mode, but also for gravity. High energy effects will be encoded for example in operators containing the perturbations in the Riemann tensor $\delta R_{\mu\nu\rho\sigma}$. As these corrections are of higher order in derivatives, we will not explicitly talk about them below.

Let us give some examples of how to write simple models of inflation in this language. A model with minimal kinetic term and a slow-roll potential $V(\phi)$ can be written in unitary gauge as

$$\int d^4x \sqrt{-g} \left[-\frac{1}{2} (\partial\phi)^2 - V(\phi) \right] \rightarrow \int d^4x \sqrt{-g} \left[-\frac{\dot{\phi}_0(t)^2}{2} g^{00} - V(\phi_0(t)) \right] . \tag{2.7}$$

²The limit in which the time shift is an exact symmetry must be taken with care because $\dot{H} \rightarrow 0$. This implies that the spatial kinetic term for the Goldstone vanishes, as we will see in the discussion of Ghost Inflation.

³Notice that this symmetry has nothing to do with the breaking of time diffeomorphisms. To see how this symmetry appears in the ϕ language notice that, after a proper field redefinition, one can always assume that $\dot{\phi} = \text{const}$. With this choice, invariance under time translation in the unitary gauge Lagrangian is implied by the shift symmetry $\phi \rightarrow \phi + \text{const}$. This symmetry and the time translation symmetry of the ϕ Lagrangian are broken down to the diagonal subgroup by the background. This residual symmetry is the time shift in the unitary gauge Lagrangian.

As the Friedmann equations give $\dot{\phi}_0(t)^2 = -2M_P^2\dot{H}$ and $V(\phi(t)) = M_{\text{pl}}^2(3H^2 + \dot{H})$ we see that the action is of the form (2.6) with all but the first three terms set to zero. Clearly this cannot be true exactly as all the other terms will be generated by loop corrections: they encode all the possible effects of high energy physics on this simple slow-roll model of inflation.

A more general case includes all the possible Lagrangians with at most one derivative acting on each ϕ : $L = P(X, \phi)$, with $X = g^{\mu\nu}\partial_\mu\phi\partial_\nu\phi$. Around an unperturbed solution $\phi_0(t)$ we have

$$S = \int d^4x \sqrt{-g} P(\dot{\phi}_0(t)^2 g^{00}, \phi(t)) \tag{2.8}$$

which is clearly of the form above with $M_n^4(t) = \dot{\phi}_0(t)^{2n}\partial^n P/\partial X^n$ evaluated at $\phi_0(t)$. Terms containing the extrinsic curvature contain more than one derivative acting on a single scalar and will be crucial in the limit of exact de Sitter, $\dot{H} \rightarrow 0$.

3. Action for the Goldstone boson

The unitary gauge Lagrangian describes three degrees of freedom: the two graviton helicities and a scalar mode. This mode will become explicit after one performs a broken time diffeomorphism (Stückelberg trick) as the Goldstone boson which non-linearly realizes this symmetry. In analogy with the equivalence theorem for the longitudinal components of a massive gauge boson [9], we expect that the physics of the Goldstone decouples from the two graviton helicities at short distance, when the mixing can be neglected. Let us review briefly what happens in a non-Abelian gauge theory before applying the same method in our case.

The unitary gauge action for a non-Abelian gauge group A_μ^a is

$$S = \int d^4x -\frac{1}{4}\text{Tr} F_{\mu\nu}F^{\mu\nu} - \frac{1}{2}m^2\text{Tr} A_\mu A^\mu, \tag{3.1}$$

where $A_\mu = A_\mu^a T^a$. Under a gauge transformation we have

$$A_\mu \rightarrow U A_\mu U^\dagger + \frac{i}{g} U \partial_\mu U^\dagger \equiv \frac{i}{g} U D_\mu U^\dagger. \tag{3.2}$$

The action therefore becomes

$$S = \int d^4x -\frac{1}{4}\text{Tr} F_{\mu\nu}F^{\mu\nu} - \frac{1}{2}\frac{m^2}{g^2}\text{Tr} D_\mu U^\dagger D_\mu U. \tag{3.3}$$

The gauge invariance can be “restored” writing $U = \exp [iT^a \pi^a(t, \vec{x})]$, where π^a are scalars (the Goldstones) which transform non-linearly under a gauge transformation Λ as

$$e^{iT^a \tilde{\pi}^a(t, \vec{x})} = \Lambda(t, \vec{x}) e^{iT^a \pi^a(t, \vec{x})} \tag{3.4}$$

Going to canonical normalization $\pi_c \equiv m/g \cdot \pi$, we see that the Goldstone boson self-interactions become strongly coupled at the scale $4\pi m/g$, which is parametrically higher

than the mass of the gauge bosons. The advantage of reintroducing the Goldstones is that for energies $E \gg m$ the mixing between them and the transverse components of the gauge field becomes irrelevant, so that the two sectors decouple. Mixing terms in eq. (3.2) are in fact of the form

$$\frac{m^2}{g} A_\mu^a \partial^\mu \pi^a = m A_\mu^a \partial^\mu \pi_c^a \quad (3.5)$$

which are irrelevant with respect to the canonical kinetic term $(\partial\pi_c)^2$ for $E \gg m$. In the window $m \ll E \ll 4\pi m/g$ the physics of the Goldstone π is weakly coupled and it can be studied neglecting the mixing with transverse components.

Let us follow the same steps for our case of broken time diffeomorphisms. Let us concentrate for instance on the two operators:

$$\int d^4x \sqrt{-g} [A(t) + B(t)g^{00}(x)] . \quad (3.6)$$

Under a broken time diff. $t \rightarrow \tilde{t} = t + \xi^0(x)$, $\vec{x} \rightarrow \vec{\tilde{x}} = \vec{x}$, g^{00} transforms as:

$$g^{00}(x) \rightarrow \tilde{g}^{00}(\tilde{x}(x)) = \frac{\partial \tilde{x}^0(x)}{\partial x^\mu} \frac{\partial \tilde{x}^0(x)}{\partial x^\nu} g^{\mu\nu}(x) . \quad (3.7)$$

The action written in terms of the transformed fields is given by:

$$\int d^4\tilde{x} \sqrt{-\tilde{g}(\tilde{x}(x))} \left| \frac{\partial \tilde{x}}{\partial x} \right| \left[A(t) + B(t) \frac{\partial x^0}{\partial \tilde{x}^\mu} \frac{\partial x^0}{\partial \tilde{x}^\nu} \tilde{g}^{\mu\nu}(\tilde{x}(x)) \right] . \quad (3.8)$$

Changing integration variables to \tilde{x} , we get:

$$\int d^4\tilde{x} \sqrt{-\tilde{g}(\tilde{x})} \left[A(\tilde{t} - \xi^0(x(\tilde{x}))) + B(\tilde{t} - \xi^0(x(\tilde{x}))) \frac{\partial(\tilde{t} - \xi^0(x(\tilde{x})))}{\partial \tilde{x}^\mu} \frac{\partial(\tilde{t} - \xi^0(x(\tilde{x})))}{\partial \tilde{x}^\nu} \tilde{g}^{\mu\nu}(\tilde{x}) \right] . \quad (3.9)$$

The procedure to reintroduce the Goldstone is now similar to the gauge theory case. Whenever ξ^0 appears in the action above, we make the substitution

$$\xi^0(x(\tilde{x})) \rightarrow -\tilde{\pi}(\tilde{x}) . \quad (3.10)$$

This gives, dropping the tildes for simplicity:

$$\int d^4x \sqrt{-g(x)} \left[A(t + \pi(x)) + B(t + \pi(x)) \frac{\partial(t + \pi(x))}{\partial x^\mu} \frac{\partial(t + \pi(x))}{\partial x^\nu} g^{\mu\nu}(x) \right] . \quad (3.11)$$

One can check that the action above is invariant under diffs at all orders (and not only for infinitesimal transformations) upon assigning to π the transformation rule

$$\pi(x) \rightarrow \tilde{\pi}(\tilde{x}(x)) = \pi(x) - \xi^0(x) . \quad (3.12)$$

With this definition π transforms as a scalar field plus an additional shift under time diffs.

Applying this procedure to the unitary gauge action (2.6) we obtain

$$\begin{aligned}
 S = \int d^4x \sqrt{-g} & \left[\frac{1}{2} M_{\text{Pl}}^2 R - M_{\text{Pl}}^2 \left(3H^2(t + \pi) + \dot{H}(t + \pi) \right) + \right. \\
 & + M_{\text{Pl}}^2 \dot{H}(t + \pi) \left((1 + \dot{\pi})^2 g^{00} + 2(1 + \dot{\pi}) \partial_i \pi g^{0i} + g^{ij} \partial_i \pi \partial_j \pi \right) + \\
 & \frac{M_2(t + \pi)^4}{2!} \left((1 + \dot{\pi})^2 g^{00} + 2(1 + \dot{\pi}) \partial_i \pi g^{0i} + g^{ij} \partial_i \pi \partial_j \pi + 1 \right)^2 + \\
 & \left. \frac{M_3(t + \pi)^4}{3!} \left((1 + \dot{\pi})^2 g^{00} + 2(1 + \dot{\pi}) \partial_i \pi g^{0i} + g^{ij} \partial_i \pi \partial_j \pi + 1 \right)^3 + \dots \right], \tag{3.13}
 \end{aligned}$$

where for the moment we have neglected for simplicity terms that involve the extrinsic curvature.

This action is rather complicated, and at this point it is not clear what is the advantage of reintroducing the Goldstone π from the unitary gauge Lagrangian. In analogy with the gauge theory case, the simplification occurs because, at sufficiently short distances, the physics of the Goldstone can be studied neglecting metric fluctuations. As for the gauge theory case, the regime for which this is possible can be estimated just looking at the mixing terms in the Lagrangian above. In eq. (3.13) we see in fact that quadratic terms which mix π and $g_{\mu\nu}$ contain fewer derivatives than the kinetic term of π so that they can be neglected above some high energy scale. In general the answer will depend on which operators are present. Let us start with the simplest case in which only the tadpole terms are relevant ($M_2 = M_3 = \dots = 0$). This corresponds to the standard slow-roll inflation case. The leading mixing with gravity will come from a term of the form

$$\sim M_{\text{Pl}}^2 \dot{H} \dot{\pi} \delta g^{00}. \tag{3.14}$$

After canonical normalization ($\pi_c \sim M_{\text{Pl}} \dot{H}^{1/2} \pi$, $\delta g_c^{00} \sim M_{\text{Pl}} \delta g^{00}$), we see that the mixing terms can be neglected for energies above $E_{\text{mix}} \sim \epsilon^{1/2} H$, where ϵ is the usual slow-roll parameter $\epsilon \equiv -\dot{H}/H^2$. Another case which will be of interest is when the operator M_2 gets large. In this case we have mixing terms of the form

$$\sim M_2^4 \dot{\pi} \delta g^{00} \tag{3.15}$$

which, upon canonical normalization (notice that now $\pi_c \sim M_2^2 \pi$), becomes negligible at energies larger than $E_{\text{mix}} \sim M_2^2/M_{\text{Pl}}$.⁴

In the regime $E \gg E_{\text{mix}}$ the action dramatically simplifies to

$$\begin{aligned}
 S_\pi = \int d^4x \sqrt{-g} & \left[\frac{1}{2} M_{\text{Pl}}^2 R - M_{\text{Pl}}^2 \dot{H} \left(\dot{\pi}^2 - \frac{(\partial_i \pi)^2}{a^2} \right) + 2M_2^4 \left(\dot{\pi}^2 + \dot{\pi}^3 - \dot{\pi} \frac{(\partial_i \pi)^2}{a^2} \right) - \frac{4}{3} M_3^4 \dot{\pi}^3 + \dots \right]. \tag{3.16}
 \end{aligned}$$

Given an inflationary model, one is interested in computing predictions for present cosmological observations. From this point of view, it seems that the decoupling limit (3.16)

⁴In the theories we are studying Lorentz symmetry is spontaneously broken, so one should define a separate regime of energies and momenta for which the mixing can be neglected. For cosmological perturbations, we will be only interested in the energy range.

is completely irrelevant for these extremely infrared scales. However, as for standard single field slow-roll inflation, one can prove that there exists a quantity, the usual ζ variable, which is constant out of the horizon at any order in perturbation theory [10, 11] (see appendix D of [12] for a generalization including terms with higher spatial derivatives). The intuitive reason for the existence of a conserved quantity is that after exiting the horizon different regions evolve exactly in the same way. The only difference is how much one has expanded with respect to another and it is this difference that remains constant.

Therefore the problem is reduced to calculating correlation functions just after horizon crossing. We are therefore interested in studying our Lagrangian with an IR energy cutoff of order H . If the decoupling scale E_{mix} is smaller than H , the Lagrangian for π (3.16) will give the correct predictions up to terms suppressed by E_{mix}/H .

As we discussed, we are assuming that the time dependence of the coefficients in the unitary gauge Lagrangian is slow compared to the Hubble time, that is, suppressed by some generalized slow roll parameters. This implies that the additional π terms coming from the Taylor expansion of the coefficients are small. In particular, the relevant operators, i.e. the ones which dominate moving towards the infrared, like the cubic term, are unimportant at the scale H and have therefore been neglected in the Lagrangian (3.16).

In conclusion, with the Lagrangian (3.16) one is able to compute all the observables which are not dominated by the mixing with gravity, like for example the non-Gaussianities in standard slow-roll inflation [13, 14]. Notice however that the tilt of the spectrum can be calculated, at leading order, with the Lagrangian (3.16). As we will see later, its value can in fact be deduced simply by the power spectrum at horizon crossing computed neglecting the mixing terms. It is important to stress that our approach does not lose its validity when the mixing with gravity is important so that the Goldstone action is not sufficient for predictions. The action (2.6) contains all the information about the model and can be used to calculate all predictions even when the mixing with gravity is large.

4. The various limits of single field inflation

4.1 Slow-roll inflation and high energy corrections

The simplest example of the general Lagrangian (2.6) is obtained by keeping only the first three terms, which are fixed once we know the background Hubble parameter $H(t)$, and setting to zero all the other operators of higher order: $M_2 = M_3 = \bar{M}_1 = \bar{M}_2 \dots = 0$. In the ϕ language, this corresponds to standard slow-roll inflation, with no higher order terms. In this case, as discussed in the last section, predictions at the scale H can be made neglecting the mixing with gravity and concentrating on the Goldstone Lagrangian (3.16). One is interested in calculating, soon after horizon crossing, the conserved quantity ζ . This is defined, at linear order, by choosing the gauge $\pi = 0$ (unitary gauge in our language) and the spatial part of the metric to be

$$g_{ij} = a^2(t) [(1 + 2\zeta(t, \vec{x}))\delta_{ij} + \gamma_{ij}] \tag{4.1}$$

where γ is transverse and traceless and it describes the two tensor degrees of freedom. The relation between π and ζ is very simple. As we are neglecting the mixing with gravity,

the metric is unperturbed in the π language; to set $\pi = 0$ one has to perform a time diffeomorphism $t \rightarrow t - \pi(t, \vec{x})$ which gives a spatial metric of the form (4.1) with

$$\zeta(t, \vec{x}) = -H\pi(t, \vec{x}). \quad (4.2)$$

For each mode k , one is only interested in the dynamics around horizon crossing $\omega(k) = k/a \sim H$. During this period the background can be approximated as de Sitter up to slow-roll corrections. Therefore, the 2-point function of the canonically normalized scalar π_c is given by the de Sitter result

$$\langle \pi_c(\vec{k}_1) \pi_c(\vec{k}_2) \rangle = (2\pi)^3 \delta(\vec{k}_1 + \vec{k}_2) \frac{H_*^2}{2k_1^3}, \quad (4.3)$$

where here and below $*$ means the value of a quantity at horizon crossing. This implies that the 2-point function of ζ is given by

$$\langle \zeta(\vec{k}_1) \zeta(\vec{k}_2) \rangle = (2\pi)^3 \delta(\vec{k}_1 + \vec{k}_2) \frac{H_*^4}{4M_{\text{Pl}}^2 |\dot{H}_*|} \frac{1}{k_1^3} = (2\pi)^3 \delta(\vec{k}_1 + \vec{k}_2) \frac{H_*^2}{4\epsilon_* M_{\text{Pl}}^2} \frac{1}{k_1^3}. \quad (4.4)$$

As the variable ζ is constant outside the horizon, this equation is exact for all k up to slow-roll corrections. In particular it allows us to calculate the tilt of the spectrum at leading order in slow-roll

$$n_s - 1 = \frac{d}{d \log k} \log \frac{H_*^4}{|\dot{H}_*|} = \frac{1}{H_*} \frac{d}{dt_*} \log \frac{H_*^4}{|\dot{H}_*|} = 4 \frac{\dot{H}_*}{H_*^2} - \frac{\ddot{H}_*}{H_* \dot{H}_*}. \quad (4.5)$$

Notice however that not all observables can be calculated from the π Lagrangian (3.16): this happens when the leading result comes from the mixing with gravity or is of higher order in the slow-roll expansion. For example, as the first two terms of eq. (3.16) do not contain self-interactions of π , the 3-point function $\langle \zeta(\vec{k}_1) \zeta(\vec{k}_2) \zeta(\vec{k}_3) \rangle$ would be zero. One is therefore forced to look at subleading corrections, taking into account the mixing with gravity in eq. (3.13).

Obviously our choice of setting to zero all the higher order terms cannot be exactly true. At the very least they will be radiatively generated even if we put them to zero at tree level. The theory is non-renormalizable and all interactions will be generated with divergent coefficients at sufficiently high order in the perturbative expansion. As additional terms are generated by graviton loops, they may be very small. For example it is straightforward to check that starting from the unitary gauge interaction $M_{\text{Pl}}^2 \dot{H} g^{00}$ a term of the form $(g^{00} + 1)^2$ will be generated with a logarithmically divergent coefficient $M_2^4 \sim \dot{H}^2 \log \Lambda$. This implies that one should assume $M_2^4 \gtrsim \dot{H}^2$ ⁽⁵⁾. This lower limit is however very small. For example the dispersion relation of π will be changed by the additional contribution to the time kinetic term: this implies, as we will discuss thoroughly below, that the speed of π excitations deviates slightly from the speed of light, by a relative amount $1 - c_s \sim M_2^4 / (|\dot{H}| M_{\text{Pl}}^2) \sim |\dot{H}| / M_{\text{Pl}}^2$. Using the normalization of the scalar spectrum

⁵The explicit calculation of logarithmic divergences in a theory of a massless scalar coupled to gravity has been carried out a long time ago in [15].

eq. (4.4), we see that the deviation from the speed of light is $\gtrsim \epsilon^2 \cdot 10^{-10}$. A not very interesting lower limit.

The size of the additional operators will be much larger if additional physics enters below the Planck scale. In general our approach gives the correct parametrization of all possible effects of new physics. As usual in an effective field theory approach, the details of the UV completion of the model are encoded in the higher dimension operators. This is very similar to what happens in physics beyond the Standard Model. At low energy the possible effects of new physics are encoded in a series of higher dimensional operators compatible with the symmetries [16]. The detailed experimental study of the Standard model allows us to put severe limits on the size of these higher dimensional operators. The same can be done in our case, although the set of conceivable observations is unfortunately much more limited. One example of a possible experimental limit on higher dimension operators is the consistency relation for the gravitational wave tilt. As is well known, the gravity wave spectrum from the Einstein-Hilbert action is given by

$$\langle \gamma^s(\vec{k}_1) \gamma^{s'}(\vec{k}_2) \rangle = (2\pi)^3 \delta(\vec{k}_1 + \vec{k}_2) \frac{H_*^2}{M_{\text{Pl}}^2} \frac{1}{k_1^3} \delta_{ss'} \quad (4.6)$$

where γ^s denotes the two possible polarizations of the gravity wave. The ratio between this contribution and the scalar one (4.4) is given by ϵ_* . The gravitational wave tilt, $n_g = -2\epsilon_*$, is thus fixed once the ratio between tensor and scalar modes is known.

This prediction is valid if one assumes $M_2 = 0$, i.e. $c_s = 1$. As we will see in fact, the scalar spectrum goes as c_s^{-1} , while predictions for gravitational waves are not changed by M_2 . The experimental verification of the consistency relation, even with large errors, would tell us that c_s cannot deviate substantially from 1 which implies

$$M_2^4 \lesssim M_{\text{Pl}}^2 |\dot{H}|. \quad (4.7)$$

Notice that the higher dimension operators will not only influence scalar fluctuations, but also the tensor modes, although these corrections are arguably much harder to test. For example the unitary gauge operator $-\bar{M}_3(t)/2 \cdot \delta K^\mu{}_\nu \delta K^\nu{}_\mu$, whose relevance for scalar fluctuations will be discussed later on, contains terms of the form \dot{g}_{ij}^2 . This will change the gravity wave dispersion relation. It is in fact straightforward to obtain the action for the tensor modes γ_{ij} in the presence of this operator. One gets

$$S_\gamma = \frac{M_{\text{Pl}}^2}{8} \int d^4x \sqrt{-g} \left[\left(1 - \frac{\bar{M}_3^2}{M_{\text{Pl}}^2} \right) \dot{\gamma}_{ij} \dot{\gamma}_{ij} - \frac{1}{a^2} \partial_t \gamma_{ij} \partial_t \gamma_{ij} \right]. \quad (4.8)$$

Therefore the spectrum of gravity waves (4.6) will get corrections of order $\bar{M}_3^2/M_{\text{Pl}}^2$. This correction is small unless we push \bar{M}_3^2 up to the Planck scale. It is easy to realize that operators of the form $(g^{00} + 1)^n$ do not influence tensor modes as they do not affect the transverse-traceless components of the metric.

Other examples of experimental limits on various operators will be discussed in the following sections.

4.2 Small speed of sound and large non-Gaussianities

The Goldstone action (3.16) shows that the spatial kinetic term $(\partial_i\pi)^2$ is completely fixed by the background evolution to be $M_{\text{Pl}}^2\dot{H}(\partial_i\pi)^2$. In particular only for $\dot{H} < 0$, it has the “healthy” negative sign. This is an example of the well studied relationship between violation of the null energy condition, which in a FRW Universe is equivalent to $\dot{H} < 0$, and the presence of instabilities in the system [17, 8]. Notice however that the wrong sign of the operator $(\partial_i\pi)^2$ is not enough to conclude that the system is pathological: higher order terms like $\delta K^\mu{}_\mu{}^2$ may become important in particular regimes, as we will discuss thoroughly below. Reference [1] studies examples in which $\dot{H} > 0$ can be obtained without pathologies.

The coefficient of the time kinetic term $\dot{\pi}^2$ is, on the other hand, not completely fixed by the background evolution, as it receives a contribution also from the quadratic operator $(g^{00} + 1)^2$. In eq. (3.16) we have

$$\left(-M_{\text{Pl}}^2\dot{H} + 2M_2^4\right)\dot{\pi}^2. \quad (4.9)$$

To avoid instabilities we must have $-M_{\text{Pl}}^2\dot{H} + 2M_2^4 > 0$. As time and spatial kinetic terms have different coefficients, π waves will have a “speed of sound” $c_s \neq 1$. This is expected as the background spontaneously breaks Lorentz invariance, so that $c_s = 1$ is not protected by any symmetry. As we discussed in the last section, deviation from $c_s = 1$ will be induced at the very least by graviton loops.⁶ The speed of sound is given by

$$c_s^{-2} = 1 - \frac{2M_2^4}{M_{\text{Pl}}^2\dot{H}}. \quad (4.10)$$

This implies that in order to avoid superluminal propagation we must have $M_2^4 > 0$ (assuming $\dot{H} < 0$). Superluminal propagation would imply that the theory has no Lorentz invariant UV completion [18]. In the following we will concentrate on the case $c_s \leq 1$, see [19] for a phenomenological discussion of models with $c_s > 1$.

Using the equation above for c_s^2 the Goldstone action can be written at cubic order as

$$S_\pi = \int d^4x \sqrt{-g} \left[-\frac{M_{\text{Pl}}^2\dot{H}}{c_s^2} \left(\dot{\pi}^2 - c_s^2 \frac{(\partial_i\pi)^2}{a^2} \right) + M_{\text{Pl}}^2\dot{H} \left(1 - \frac{1}{c_s^2} \right) \left(\dot{\pi}^3 - \dot{\pi} \frac{(\partial_i\pi)^2}{a^2} \right) - \frac{4}{3} M_3^4 \dot{\pi}^3 \dots \right]. \quad (4.11)$$

From the discussion in section (3) we know that the mixing with gravity can be neglected at energies $E \gg E_{\text{mix}} \simeq M_2^2/M_{\text{Pl}}$. This implies that predictions for cosmological observables, which are done at energies of order H , are captured at leading order by the

⁶If we neglect the coupling with gravity and the time dependence of the operators in the unitary gauge Lagrangian (so that $\pi \rightarrow \pi + \text{const}$ is a symmetry), $c_s = 1$ can be protected by a symmetry $\partial_\mu\pi \rightarrow \partial_\mu\pi + v_\mu$, where v_μ is a constant vector. Under this symmetry the Lorentz invariant kinetic term of π changes by a total derivative, while the operator proportional to M_2^4 in eq. (3.16) is clearly not invariant, so that $c_s = 1$. Notice that the theory is not free as we are allowed to write interactions with more derivatives acting on π . This symmetry appears in the study of the brane bending mode of the DGP model [18].

Goldstone action (4.11) if $H \gg M_2^2/M_{\text{Pl}}$, or equivalently for $\epsilon/c_s^2 \ll 1$. If this is not the case one is not assured that the Goldstone action contains the leading effects.

The calculation of the 2-point function follows closely the case $c_s = 1$ if we use a rescaled momentum $\bar{k} = c_s k$ and take into account the additional factor c_s^{-2} in front of the time kinetic term. We obtain

$$\langle \zeta(\vec{k}_1) \zeta(\vec{k}_2) \rangle = (2\pi)^3 \delta(\vec{k}_1 + \vec{k}_2) \frac{1}{c_{s*}} \cdot \frac{H_*^4}{4M_{\text{Pl}}^2 |\dot{H}_*| k_1^3} = (2\pi)^3 \delta(\vec{k}_1 + \vec{k}_2) \frac{1}{c_{s*}} \cdot \frac{H_*^2}{4\epsilon_* M_{\text{Pl}}^2 k_1^3}. \quad (4.12)$$

The variation with time of the speed of sound introduces an additional contribution to the tilt

$$n_s = \frac{d}{d \log k} \log \frac{H_*^4}{|\dot{H}_*| c_{s*}} = \frac{1}{H_*} \frac{d}{dt_*} \log \frac{H_*^4}{|\dot{H}_*| c_{s*}} = 4 \frac{\dot{H}_*}{H_*^2} - \frac{\ddot{H}_*}{\dot{H}_* H_*} - \frac{\dot{c}_{s*}}{c_{s*} H_*}. \quad (4.13)$$

The result agrees with the one found in [20].

From the action (4.11) we clearly see that the same operator giving a reduced speed of sound induces cubic couplings of the Goldstones of the form $\dot{\pi}(\nabla\pi)^2$ and $\dot{\pi}^3$. The non-linear realization of time diffeomorphisms forces a relation between a reduced speed of sound and an enhanced level of the 3-point function correlator, i.e. non-Gaussianities. This relationship was stressed in the explicit calculation of the 3-point function in [21].

To estimate the size of non-Gaussianities, one has to compare the non-linear corrections with the quadratic terms around freezing, $\omega \sim H$. In the limit $c_s \ll 1$, the operator $\dot{\pi}(\nabla\pi)^2$ gives the leading contribution, as the quadratic action shows that a mode freezes with $k \sim H/c_s$, so that spatial derivatives are enhanced with respect to time derivatives. The level of non-Gaussianity will thus be given by the ratio:

$$\frac{\mathcal{L}_{\dot{\pi}(\nabla\pi)^2}}{\mathcal{L}_2} \sim \frac{H\pi \left(\frac{H}{c_s}\pi\right)^2}{H^2\pi^2} \sim \frac{H}{c_s^2}\pi \sim \frac{1}{c_s^2}\zeta, \quad (4.14)$$

where in the last step we have used the linear relationship between π and ζ , eq. (4.2). Taking $\zeta \sim 10^{-5}$ we have an estimate of the size of the non-linear correction.⁷ Usually the magnitude of non-Gaussianities is given in terms of the parameters f_{NL} , which are parametrically of the form: $\mathcal{L}_{\dot{\pi}(\nabla\pi)^2}/\mathcal{L}_2 \sim f_{\text{NL}}\zeta$. The leading contribution will thus give

$$f_{\text{NL}, \dot{\pi}(\nabla\pi)^2}^{\text{equil.}} \sim \frac{1}{c_s^2}. \quad (4.15)$$

The superscript “equil.” refers to the momentum dependence of the 3-point function, which in these models is of the so called equilateral form [22]. This is physically clear in the Goldstone language as the relevant π interactions contain derivatives, so that they

⁷The size of the non-linear corrections depend on the specific value of ζ . Even if the typical value of ζ is small, one may be interested in very large (and therefore very unlikely) fluctuations, for example to study the production of primordial black holes. For sufficiently large values of ζ , $\zeta \gtrsim c_s^2$, non-linear corrections become of order 1 and the perturbative expansion breaks down. Therefore, predictions which depend on very large values of ζ may lie out of the regime of validity of the effective field theory.

die out quickly out of the horizon; the correlation is only among modes with comparable wavelength.

In the Goldstone Lagrangian (4.11) there is an additional independent operator, $-\frac{4}{3}M_3^4\dot{\pi}^3$, contributing to the 3-point function, coming from the unitary gauge operator $(g^{00} + 1)^3$. We thus have two contributions of the form $\dot{\pi}^3$ which give

$$f_{\text{NL}, \dot{\pi}^3}^{\text{equil.}} \sim 1 - \frac{4}{3} \frac{M_3^4}{M_{\text{Pl}}^2 |\dot{H}| c_s^{-2}} . \quad (4.16)$$

The size of the operator $-\frac{4}{3}M_3^4\dot{\pi}^3$ is not constrained by the non-linear realization of time diffeomorphisms: it is a free parameter. In DBI inflation [2] we have $M_3^4 \sim M_{\text{Pl}}^2 |\dot{H}| c_s^{-4}$, so that its contribution to non-Gaussianities is of the same order as the one of eq. (4.15). The same approximate size of the M_3^4 is obtained if we assume that both the unitary gauge operators $M_2^4(g^{00} + 1)^2$ and $M_3^4(g^{00} + 1)^3$ become strongly coupled at the same energy scale.

It is interesting to look at the experimental limits on non-Gaussianities as a constraint on the size of the unitary gauge operator $(g^{00} + 1)^2$ and therefore on the speed of sound. The explicit calculation [21] gives the contribution of the operator $\dot{\pi}(\nabla\pi)^2$ to the experimentally constrained parameter $f_{\text{NL}}^{\text{equil.}}$; at leading in order in c_s^{-1} we have⁸

$$f_{\text{NL}}^{\text{equil.}} = \frac{85}{324} \cdot \frac{1}{c_s^2} . \quad (4.17)$$

The experimentally allowed window [23]

$$-256 < f_{\text{NL}}^{\text{equil.}} < 332 \quad \text{at 95\% C.L.} \quad (4.18)$$

translates into the constraint

$$c_s > 0.028 \quad \text{at 95\% C.L.} \quad (4.19)$$

Notice however that, although in principle the operators $\dot{\pi}(\nabla\pi)^2$ and $\dot{\pi}^3$ give a different momentum dependence of the 3-point function, this difference is not experimentally appreciable at present, so that the constraint (4.18) is on the joint effect of the two operators. The constraint on the speed of sound will hold barring a cancellation between the two operators. In the case of DBI inflation for example the effect of the operator $M_3^4(g^{00} + 1)^3$ is sizeable as we discussed. However there is no cancellation and the constraint on the speed of sound is only slightly changed to

$$\text{DBI : } c_s > 0.031 \quad \text{at 95\% C.L.} \quad (4.20)$$

Although we concentrated so far on the Goldstone Lagrangian, it is important to stress that this general approach is useful also when one is interested in taking into account the full mixing with gravity. For example, going back to the unitary gauge Lagrangian (2.6), we can easily see how many coefficients will be relevant in calculating the 3-point function.

⁸This is obtained setting $P_{,XXX} = 0$ in the notation of [21].

At leading order in slow-roll and in derivatives there are 2 coefficients as we discussed: M_2 and M_3 . At first order in slow-roll, there will be 4 new parameters describing the slow variation of the coefficients: the conventional ϵ and η slow-roll parameters and two additional ones for the coefficients of the operators $(g^{00} + 1)^2$ and $(g^{00} + 1)^3$. This in fact is what one finds in the explicit calculation [21].⁹

All the discussion can be straightforwardly extended to the 4-point function (and higher order correlators). In the Goldstone Lagrangian we have 3 operators contributing to the 4-point function (again at leading order in slow-roll and derivatives): $(g^{00} + 1)^2$, which is fixed by the speed of sound c_s , $(g^{00} + 1)^3$ and $(g^{00} + 1)^4$. Let us estimate the effect of the operator which is fixed by the speed of sound. As we did for the 3-point function, it is easy to see that the effect will be dominated by the operator $(\nabla\pi)^4$ and that the level of non-Gaussianity induced by it can be estimated as

$$\frac{\mathcal{L}_{(\nabla\pi)^4}}{\mathcal{L}_2} \sim \frac{\left(\frac{H}{c_s}\pi\right)^4}{H^2\pi^2} \sim \frac{H^2}{c_s^4}\pi^2 \sim \frac{1}{c_s^4}\zeta^2. \quad (4.21)$$

This matches with the explicit calculation done in [24].

4.2.1 Cutoff and naturalness

As discussed, for $c_s < 1$ the Goldstone action contains non-renormalizable interactions. Therefore the self-interactions among the Goldstones will become strongly coupled at a certain energy scale, which sets the cutoff of our theory. This cutoff can be estimated looking at tree level partial wave unitarity, i.e. finding the maximum energy at which the tree level scattering of π s is unitary. The calculation is straightforward, the only complication coming from the non-relativistic dispersion relation. The cutoff scale Λ turns out to be

$$\Lambda^4 \simeq 16\pi^2 M_2^4 \frac{c_s^7}{(1 - c_s^2)^2} \simeq 16\pi^2 M_{\text{Pl}}^2 |\dot{H}| \frac{c_s^5}{1 - c_s^2}. \quad (4.22)$$

The same result can be obtained looking at the energy scale where loop corrections to the $\pi\pi$ scattering amplitude become relevant. As expected the theory becomes more and more strongly coupled for small c_s , so that the cutoff scale decreases. On the other hand, for $c_s \rightarrow 1$ the cutoff becomes higher and higher. This makes sense as there are no non-renormalizable interactions in this limit and the cutoff can be extended up to the Planck scale. This cutoff scale is obtained just looking at the unitary gauge operator $(g^{00} + 1)^2$; depending on their size the other independent operators may give an even lower energy cutoff. Notice that the scale Λ indicates the maximum energy at which our theory is weakly coupled and make sense; below this scale new physics must come into the game. However new physics can appear even much below Λ .

If we are interested in using our Lagrangian for making predictions for cosmological correlation functions, then we need to use it at a scale of order the Hubble parameter H

⁹The explicit calculation shows that one of the coefficients does not give rise to an independent momentum dependence of the 3-point function, so that it cannot be disentangled from the other parameters.

during inflation. We therefore need that this energy scale is below the cutoff, $H \ll \Lambda$. Using the explicit expression for the cutoff (4.22) in the case $c_s \ll 1$ one gets

$$H^4 \ll M_{\text{Pl}}^2 |\dot{H}| c_s^5 \tag{4.23}$$

which can be rewritten using the spectrum normalization (4.12) as an inequality for the speed of sound

$$c_s \gg P_\zeta^{1/4} \simeq 0.003 . \tag{4.24}$$

A theory with a lower speed of sound is strongly coupled at $E \simeq H$. Not surprisingly this value of the speed of sound also corresponds to the value at which non-Gaussianity are of order one: the theory is strongly coupled at the energy scale H relevant for cosmological predictions.

Let us comment on the naturalness of the theory. One may wonder whether the limit of small c_s is natural or instead loop corrections will induce a larger value. The Goldstone self-interactions, $\dot{\pi}(\nabla\pi)^2$ and $(\nabla\pi)^4$ for example, will induce a radiative contribution to $(\nabla\pi)^2$. It is easy to estimate that these contributions are of order $c_s^{-5} \Lambda^4 / (16\pi^2 M_2^4)$, where Λ is the UV cutoff, i.e. the energy scale at which new physics enters in the game. We can see that it is impossible to have large radiative contribution; even if we take Λ at the unitarity limit (4.22), the effect is of the same order as the tree level value. This makes sense as the unitarity cutoff is indeed the energy scale at which loop corrections become of order one.

We would like also to notice that the action (3.16) is natural from an effective field theory point of view [28]. The relevant operators are in fact protected from large renormalizations if we assume an approximate shift symmetry of π . In this case the coefficients of the relevant operators are sufficiently small and they will never become important for observations as cosmological correlation functions probe the theory at a fixed energy scale of order H : we never go to lower energy. Clearly here we are only looking at the period of inflation, where an approximate shift symmetry is enough to make the theory technically natural; providing a graceful exit from inflation and an efficient reheating are additional requirements for a working model which are not discussed in our formalism.

4.3 De-Sitter limit and the Ghost condensate

In the previous section we saw that the limit $c_s \rightarrow 0$ is pathological as the theory becomes more and more strongly coupled. However we have neglected in our discussion the higher derivative operators in the unitary gauge Lagrangian (2.6)

$$\int d^4x \sqrt{-g} \left(-\frac{\bar{M}_2(t)^2}{2} \delta K^\mu{}_\mu{}^2 - \frac{\bar{M}_3(t)^2}{2} \delta K^\mu{}_\nu \delta K^\nu{}_\mu \right) . \tag{4.25}$$

These operators give rise in the Goldstone action to a spatial kinetic term of the form

$$\int d^4x \sqrt{-g} \left[-\frac{\bar{M}^2}{2} \frac{1}{a^4} (\partial_i^2 \pi)^2 \right] , \tag{4.26}$$

where $\bar{M}^2 = \bar{M}_2^2 + \bar{M}_3^2$. This spatial kinetic term will make the Goldstone propagate even in the limit $c_s \rightarrow 0$. It is therefore interesting to consider our general Lagrangian in the limit $\dot{H} = 0$, when the gravitational background is exactly de Sitter space which implies $c_s = 0$. As H is now time independent, it is possible to impose an additional symmetry to the theory: the time independence of all the coefficients in the unitary gauge Lagrangian. Looking back at the procedure (3.13) to reintroduce the Goldstone π , we realize that this symmetry forbids any dependence on π without derivatives. The Goldstone action is thus invariant under shift of π

$$\pi(\vec{x}, t) \rightarrow \pi(\vec{x}, t) + \text{const.} \tag{4.27}$$

This is the limit of Ghost Condensation [25], where the Goldstone has a non-relativistic dispersion relation $\omega \propto k^2$. More generally one can consider intermediate situations where both the spatial kinetic term $c_s^2(\nabla\pi)^2$ and the higher derivative one $(\nabla^2\pi)^2$ are present. The predictions of the theory will change significantly depending on which term dominates at the energy of freezing $\omega \sim H$ [7, 27].

As with the previous models, one must find the energy regime for which the mixing of the Goldstone with gravity can be neglected. One simple way to estimate this range is to look at the δK operators which contain terms like

$$\delta K_{ij} \supset (\partial_i\partial_j\pi + \partial_i g_{0j}) . \tag{4.28}$$

Going to canonical normalization this shows that the mixing with gravity can be neglected for $k \gtrsim M_2^2/M_{\text{Pl}}$. As the dispersion relation of the Goldstone is of the form $\omega^2 = \bar{M}^2/M_2^4 \cdot k^4$, we see that the energy E_{mix} under which the mixing is relevant is $E_{\text{mix}} \simeq \bar{M}M_2^2/M_{\text{Pl}}^2$ [25]. Notice that this scale has nothing to do with the curvature of the background. This is a quite remarkable feature of this example, as usually the mixing with gravity is related to the background stress energy tensor and therefore to the curvature of spacetime: the more a system curves space, the more it mixes with gravity. In this case on the other hand, the mixing will be relevant even on a flat Minkowski background. This is what one calls a proper modification of gravity: gravity, for example the Newtonian potential generated by a source, is modified at scales much smaller than the curvature. This model of modification of gravity and its rich phenomenology has been studied in [25].¹⁰

As we are interested in inflation, we concentrate on the opposite limit $H \gg E_{\text{mix}}$, when the mixing can be neglected and one can focus on the π Lagrangian. Let us briefly describe the main features of Ghost Inflation, referring for details to [25, 7, 27], where the theory is studied with an approach very close to the one presented in this paper. Most of the interesting features can be understood looking at the scaling with energy of the various operators. Given the non-relativistic dispersion relation, $\omega \propto k^2$, the way an operator scales with energy does not coincide with its mass dimension as in the Lorentz invariant

¹⁰Also in the case of models with a reduced speed of sound, the scale of mixing with gravity can become parametrically smaller than the horizon; it is enough to have $\epsilon/c_s^2 \ll 1$. In this case the model can be considered a way of modifying gravity. Notice however that one can not take the limit $\dot{H} = 0$ without considering the spatial higher derivative terms: the scalar mode would not propagate otherwise.

case. A rescaling of the energy by a factor s , $E \rightarrow sE$, (equivalent to a time rescaling $t \rightarrow s^{-1}t$), must go together with a momentum transformation $k \rightarrow s^{1/2}k$ ($x \rightarrow s^{-1/2}x$ on the spatial coordinates). As the quadratic action for the Goldstones is of the form

$$\int d^4x \left[2M_2^4 \dot{\pi}^2 - \frac{\bar{M}^2}{2} \frac{1}{a^4} (\partial_i^2 \pi)^2 \right] \quad (4.29)$$

we must assign to π the scaling dimension $1/4$

$$\pi \rightarrow s^{1/4} \pi. \quad (4.30)$$

to keep the quadratic action invariant. With this rule it is easy to check that all the allowed Goldstone operators, besides the kinetic term (4.29), are irrelevant, i.e. they have positive scaling dimension and they become less and less relevant going down in energy. This shows that the theory makes sense as an effective field theory. In particular the higher derivative time kinetic operator $\ddot{\pi}^2$, which would naively seem as important as $(\nabla^2 \pi)^2$ and would describe the presence of a ghost in the theory, has dimension 2 and it can be neglected at low energies. If one assumes that there is a single scale M in the problem, $M \simeq M_2 \simeq \bar{M}$, this will also set the energy cutoff of the effective field theory description.

The scaling dimension of π also allows us to estimate the spectrum of perturbations produced in Ghost Inflation. The dimension of π tells us how the amplitude of quantum fluctuations changes with energy. At the scale of the cutoff $\Lambda \simeq M$, the quantum fluctuations of the canonically normalized Goldstone field $\pi_c \simeq M^2 \pi$ are of the order of the cutoff $\delta\pi_c(M) \sim M$. Going down in energy we can estimate the quantum fluctuations at freezing $E \sim H$. In the standard case the scalar would have dimension 1 and its fluctuations at freezing would be of order H ; in this case on the other hand we have

$$\delta\pi_c(H) \sim \delta\pi_c(M) \left(\frac{H}{M} \right)^{1/4} \sim (HM^3)^{1/4}. \quad (4.31)$$

Quantum fluctuations at the scale H are much enhanced with respect to a scalar with a Lorentz invariant dispersion relation. The spectrum of ζ will thus be given by [7]

$$\langle \zeta(\vec{k}_1) \zeta(\vec{k}_2) \rangle \sim (2\pi)^3 \delta(\vec{k}_1 + \vec{k}_2) \frac{H^2}{M^4} (HM^3)^{1/2} \frac{1}{k_1^3}. \quad (4.32)$$

The correct normalization of the spectrum requires $(H/M)^{5/4} \simeq 10^{-5}$.

The non-Gaussianity will be dominated by the operator with the lowest dimension. It is straightforward to see that the operator $M_2^4 \dot{\pi} (\nabla \pi)^2$ coming from the unitary gauge operator $M_2^4 (g^{00} + 1)^2$ has dimension $1/4$ and it is the least irrelevant operator. At the cutoff scale M the theory is strongly coupled. As the cubic operator has dimension $1/4$, at energies of order H it will give a level of non-Gaussianity of order $(H/M)^{1/4}$, which is parametrically of order $P_\zeta^{1/5}$. The same result can be obtained with the approach used in the last section, i.e. comparing the interaction term with the free action at freezing

$$\frac{\mathcal{L}_{\dot{\pi}(\nabla\pi)^2}}{\mathcal{L}_2} \sim \frac{M^4 H \pi (HM) \pi^2}{M^4 H^2 \pi^2} = M \pi = \frac{M}{H} \zeta \sim \left(\frac{H}{M} \right)^{1/4}. \quad (4.33)$$

The level of non-Gaussianity is extremely high compared to standard slow-roll as a consequence of the very low dimension of the most relevant operators. The explicit calculation [7] gives an effect which is somewhat smaller than the naive estimate and comparable to the existing experimental bound [23].

In our discussion we have neglected so far the unitary gauge operator

$$\int d^4x \sqrt{-g} \left(-\frac{\bar{M}_1(t)^3}{2} (g^{00} + 1) \delta K^\mu{}_\mu \right). \quad (4.34)$$

This operator is odd under time reversal, so that it is consistent to set it to zero. If this term is present, there is a second operator with dimension 1/4 in the Goldstone Lagrangian, of the form $\nabla^2\pi(\nabla\pi)^2$. Its contribution to the 3-point function would be comparable with $\dot{\pi}(\nabla\pi)^2$. The unitary gauge operator (4.34) also contributes to the π quadratic Lagrangian as we are now going to discuss.

4.3.1 De-Sitter limit without the Ghost condensate

In this section we want to study the effect of the operator (4.34) on the quadratic π action. We will see that, if the coefficient of this operator is sufficiently large, we obtain a new de Sitter limit, where the dispersion relation at freezing is of the form $\omega^2 \propto k^2$, instead of the Ghost Condensate behavior $\omega^2 \propto k^4$.

For simplicity we can take \bar{M}_1 to be time independent. Reintroducing the Goldstone we get a 3-derivative term of the form $-\bar{M}_1^3 \dot{\pi} \nabla^2 \pi / a^2$ ⁽¹¹⁾. This would be a total time derivative without the time dependence of the scale factor $a(t)$ and of the metric determinant. Integrating by parts we get a standard 2-derivative spatial kinetic term

$$- \int d^4x \sqrt{-g} \frac{\bar{M}_1^3 H}{2} \left(\frac{\partial_i \pi}{a} \right)^2. \quad (4.35)$$

In the exact de Sitter limit, $\dot{H} = 0$, and taking $M_2 \sim \bar{M}_1 \sim M$, this operator gives a dispersion relation of the form $\omega^2 = c_s^2 k^2$, with a small speed of sound¹²

$$c_s^2 = \frac{H}{M} \ll 1. \quad (4.36)$$

This will hold only if the higher derivative operators $\delta K^\mu{}_\mu{}^2$ and $\delta K^\mu{}_\nu \delta K^\nu{}_\mu$ are subdominant. If we assume that they are characterized by the same mass scale, $\bar{M}_2 \sim \bar{M}_3 \sim M$, the dispersion relation will get two contributions

$$\omega^2 \sim \frac{H}{M} k^2 + \frac{k^4}{M^2}. \quad (4.37)$$

¹¹The operator gives also a contribution to $\dot{\pi}^2$ proportional to H . We will assume that this is small compared to $M_2^4 \dot{\pi}^2$. In Minkowski space the operator we are studying can be forbidden by a $\phi \rightarrow -\phi$ symmetry, which is equivalent to time reversal in unitary gauge [25]. In a de Sitter background this symmetry is broken by the metric, so that this operator cannot be set to zero.

¹²In this model the mixing with gravity is rather different from the previous cases. The reason is that a time derivative is integrated by parts to get to eq. (4.35), so that the Goldstone terms contain the same number of derivatives as the terms describing the mixing with gravity. This implies that the mixing does not become less and less relevant going to high energy. On the other hand one can choose the model parameters in such a way that the mixing is always irrelevant. See [23] for the explicit calculations.

The two spatial kinetic term are comparable at freezing $\omega \sim H$. On the other hand, if the k^4 contribution is somewhat suppressed, it becomes irrelevant at freezing and therefore for inflationary predictions. In this limit we have a new kind of Ghost Inflation with an exactly de Sitter background, but with a $\omega^2 \propto k^2$ dispersion relationship at freezing.

Following what we did for the other models it is straightforward to obtain the spectrum normalization and an estimate of the 3-point function non-Gaussianity.

$$\langle \zeta(\vec{k}_1) \zeta(\vec{k}_2) \rangle \sim (2\pi)^3 \delta(\vec{k}_1 + \vec{k}_2) \frac{H^4}{M^4} \frac{1}{c_s^3} \frac{1}{k_1^3} \sim (2\pi)^3 \delta(\vec{k}_1 + \vec{k}_2) \left(\frac{H}{M} \right)^{5/2} \frac{1}{k_1^3}. \quad (4.38)$$

$$\frac{\mathcal{L}_{\dot{\pi}(\nabla\pi)^2}}{\mathcal{L}_2} \sim \frac{1}{c_s^2} H \pi \sim \left(\frac{H}{M} \right)^{1/4}. \quad (4.39)$$

A comparable contribution will come from the Goldstone operator $(\nabla\pi)^2 \nabla^2 \pi$. Not surprisingly the estimates above are the same as the ones we obtained in the Ghost Condensate case eqs (4.32) and (4.33). As we discussed in fact, taking all the operators at the same scale one gets a comparable contribution at freezing from the k^2 and k^4 spatial kinetic terms. We thus expect similar predictions when we assume that only one of the two contributions is present.

Now that we have found two different de Sitter limits, one dominated at freezing by $(g^{00} + 1)\delta K^\mu{}_\mu$ and the other by $\delta K^\mu{}_\mu{}^2$ and $\delta K^\mu{}_\nu \delta K^\nu{}_\mu$, one may wonder if there are other possibilities. One could imagine that both these spatial kinetic terms are suppressed for some reason and the leading operators come at higher order. In this case one would end up with a dispersion relation of the form

$$\omega^2 \sim k^{2n} \quad n \geq 3. \quad (4.40)$$

However it is easy to realize that this cannot be the case, because the theory would not make sense as an effective field theory. Following the same logic we used for Ghost Condensation, we find that the scaling dimension of the operator π would be

$$\pi \rightarrow s^{-\frac{1}{2} + \frac{3}{2n}} \pi. \quad (4.41)$$

This implies that the operator $\dot{\pi}(\nabla\pi)^2$, which is linked by symmetry to the time kinetic term $\dot{\pi}^2$, has dimension $(7 - 3n)/(2n)$. For $n \geq 3$ this operator is strong at low energy, so that the effective field theory does not make sense.

5. Conclusions

Given the ongoing experimental effort to test inflation and the proliferation of different models, it is quite important to characterize the most general theory of inflation. In this paper we took a novel point of view: instead of writing down a general Lagrangian and study the fluctuations around an inflating solution, we directly describe the effective theory of fluctuations around a quasi de Sitter background, where spatial diffeomorphisms are explicit and the time ones are non-linearly realized. We showed that the most generic

action can be written at leading order in derivatives in the form

$$\begin{aligned}
 S = \int d^4x \sqrt{-g} & \left[\frac{1}{2} M_{\text{Pl}}^2 R + M_{\text{Pl}}^2 \dot{H} g^{00} - M_{\text{Pl}}^2 (3H^2 + \dot{H}) + \frac{1}{2!} M_2(t)^4 (g^{00} + 1)^2 \right. \\
 & + \frac{1}{3!} M_3(t)^4 (g^{00} + 1)^3 + -\frac{\bar{M}_1(t)^3}{2} (g^{00} + 1) \delta K^\mu{}_\mu - \frac{\bar{M}_2(t)^2}{2} \delta K^\mu{}_\mu{}^2 \\
 & \left. - \frac{\bar{M}_3(t)^2}{2} \delta K^\mu{}_\nu \delta K^\nu{}_\mu + \dots \right]. \tag{5.1}
 \end{aligned}$$

Cosmological correlation functions test this effective field theory at a scale of order the Hubble parameter H . In this approach the role of symmetries is made much more transparent. One can see explicitly which features are implied by the inflating background solution and in particular the quite different behavior in the cases $\dot{H} < 0$, $\dot{H} = 0$ and $\dot{H} > 0$ as the coefficient of the operator g^{00} is fixed by \dot{H} . From this point of view, our approach makes clearer the relationship among inflation, theories of modification of gravity and theories which violate the Null Energy Condition (equivalent to $\dot{H} > 0$ in the cosmological context) like the bouncing models [1, 29, 30]. Another example of the role of symmetries is given by the link between a reduced speed of sound and an enhanced level of non-Gaussianity as both come from the same operator $M_2(t)^4 (g^{00} + 1)^2$ and are thus related by the non-linear realization of time diffeomorphisms.

All the possible deviations from a vanilla slow-roll scenario are systematically encoded in the size of higher order operators, similarly to what happens in the study of the Standard Model of particle physics. Moreover all single field models are unified in a common framework and this allows us to draw general conclusions which are independent of the specific realization, as done in [12, 31] for example.

It is easy to think about possible extensions of our formalism. Along the same lines it would be interesting to study the most general theory of (single field) quintessence and to work out its phenomenological consequences. Differently from inflation, which probes the effective theory at a scale of order H , we would be interested in this case to the subhorizon dynamics of perturbations. It would also be interesting to use our approach for the study of fluctuations in fluids like in radiation or matter dominance [8]. Finally it should be straightforward to introduce additional fields into the game and study multi-field inflationary models.

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A. The most general Lagrangian in unitary gauge

Let us study what are the rules for writing down the most general Lagrangian in unitary gauge. In a theory which is only invariant under spatial diffeomorphisms there is a preferred

slicing of spacetime given by a function $\tilde{t}(x)$ (with time-like gradient), which non-linearly realizes time diffeomorphisms. For example if the breaking is given by a time evolving scalar, surfaces of constant \tilde{t} are also of constant value of the scalar. Unitary gauge is the one in which the time coordinate t is chosen to coincide with \tilde{t} , so that the additional degree of freedom \tilde{t} does not explicitly appear in the action. One can therefore build various terms:

1. Terms which are invariant under all diffeomorphisms: these are just polynomials of the Riemann tensor $R_{\mu\nu\rho\sigma}$ and of its covariant derivatives, contracted to give a scalar.¹³
2. A generic function of \tilde{t} becomes $f(t)$ in unitary gauge. We are therefore free to use generic functions of time in front of any terms in the Lagrangian.
3. The gradient $\partial_\mu\tilde{t}$ becomes δ_μ^0 in unitary gauge. Thus in every tensor we can always leave free an upper 0 index. For example we can use g^{00} (and functions of it) in the unitary gauge Lagrangian, or the component of the Ricci tensor R^{00} .
4. It is useful to define a unit vector perpendicular to surfaces of constant \tilde{t}

$$n_\mu = \frac{\partial_\mu\tilde{t}}{\sqrt{-g^{\mu\nu}\partial_\mu\tilde{t}\partial_\nu\tilde{t}}} . \tag{A.1}$$

This allows to define the induced spatial metric on surfaces of constant \tilde{t} : $h_{\mu\nu} \equiv g_{\mu\nu} + n_\mu n_\nu$. Every tensor can be projected on the surfaces using $h_{\mu\nu}$. In particular we can use in our action the Riemann tensor of the induced 3d metric ${}^{(3)}R_{\alpha\beta\gamma\delta}$ and covariant derivatives with respect to the 3d metric.

5. Additional possibilities will come from the covariant derivatives of $\partial_\mu\tilde{t}$. Notice that we can equivalently look at covariant derivatives of n_μ : the derivative acting on the normalization factor just gives terms like $\partial_\mu g^{00}$ which are covariant on their own and can be used in the unitary gauge Lagrangian. The covariant derivative of n_μ projected on the surfaces of constant \tilde{t} gives the extrinsic curvature of these surfaces

$$K_{\mu\nu} \equiv h_\mu^\sigma \nabla_\sigma n_\nu . \tag{A.2}$$

The index ν is already projected on the surface because $n^\nu \nabla_\sigma n_\nu = \frac{1}{2} \nabla_\sigma (n^\nu n_\nu) = 0$. The covariant derivative of n_ν perpendicular to the surface can be rewritten as

$$n^\sigma \nabla_\sigma n_\nu = -\frac{1}{2} (-g^{00})^{-1} h_\nu^\mu \partial_\mu (-g^{00}) \tag{A.3}$$

so that it does not give rise to new terms. Therefore all covariant derivatives of n_μ can be written using the extrinsic curvature $K_{\mu\nu}$ (and its covariant derivatives) and derivatives of g^{00} .

¹³The metric and the completely antisymmetric tensor $(-g)^{-1/2} \epsilon^{\mu\nu\rho\sigma}$ can be used to contract indices.

6. Notice that using at the same time the Riemann tensor of the induced 3d metric and the extrinsic curvature is redundant as ${}^{(3)}R_{\alpha\beta\gamma\delta}$ can be rewritten using the Gauss-Codazzi relation as [32]

$${}^{(3)}R_{\alpha\beta\gamma\delta} = h_{\alpha}^{\mu}h_{\beta}^{\nu}h_{\gamma}^{\rho}h_{\delta}^{\sigma}R_{\mu\nu\rho\sigma} - K_{\alpha\gamma}K_{\beta\delta} + K_{\beta\gamma}K_{\alpha\delta} . \quad (\text{A.4})$$

Thus one can forget about the 3d Riemann tensor altogether. We can also avoid using the induced metric $h_{\alpha\beta}$ explicitly: written in terms of the 4d metric and n_{μ} one gets only terms already discussed above.¹⁴ Finally also the use of covariant derivatives with respect to the induced 3d metric can be avoided: the 3d covariant derivative of a projected tensor can be obtained as the projection of the 4d covariant derivative [32].

We conclude that the most generic action in unitary gauge is given by

$$S = \int d^4x \sqrt{-g} F(R_{\mu\nu\rho\sigma}, g^{00}, K_{\mu\nu}, \nabla_{\mu}, t) , \quad (\text{A.5})$$

where all the free indices inside the function F must be upper 0's.

B. Expanding around a given FRW solution

In this section we want to prove that the most generic theory with broken time diffeomorphisms around a given FRW background (with $k = -1, 0, +1$ depending of the spatial curvature) can be written as

$$S = \int d^4x \sqrt{-g} \left[\frac{1}{2} M_{\text{Pl}}^2 R + M_{\text{Pl}}^2 \left(\dot{H} - \frac{k}{a^2} \right) \cdot g^{00} - M_{\text{Pl}}^2 \left(3H^2 + \dot{H} + 2\frac{k}{a^2} \right) + \dots \right] \quad (\text{B.1})$$

where the dots stand for terms which are invariant under spatial diffeomorphisms and of quadratic (or higher) order in the fluctuations around the given FRW background.¹⁵

As such this statement is trivial. We know that the displayed terms give rise to the wanted FRW evolution so that, if we do not want to move away from it, the additional operators must start quadratic around this solution. What we want to say is that *each one* of the additional invariant terms is quadratic (or of higher order) without cancellation of linear contributions among various operators. These terms will be written as polynomials (quadratic and higher) of linear operators like $g^{00} + 1$, $\delta K_{\mu\nu} = K_{\mu\nu} - K_{\mu\nu}^{(0)}$, $\delta R_{\mu\nu\rho\sigma} \equiv R_{\mu\nu\rho\sigma} - R_{\mu\nu\rho\sigma}^{(0)}$ and so on. Notice that these terms start linear in the perturbations as we have explicitly removed their value evaluated on the given FRW solution. Given the symmetries of a FRW metric, every tensor evaluated on the background ($K_{\mu\nu}^{(0)}$, $R_{\mu\nu\rho\sigma}^{(0)}$, $(\nabla_{\alpha} R_{\mu\nu\rho\sigma})^{(0)}$...) can be written just in terms of $g_{\mu\nu}$, n_{μ} and functions of time. For example

$$K_{\mu\nu}^{(0)} = a^2 H h_{\mu\nu} \quad (\text{B.2})$$

$$R_{\mu\nu\rho\sigma}^{(0)} = 2(H + k)h_{\mu[\rho}h_{\sigma]\nu} + (\dot{H} + H^2)a^2 h_{\mu\sigma}\delta_{\nu}^0\delta_{\rho}^0 + \text{perm.} \quad (\text{B.3})$$

¹⁴Notice that the determinant of the induced metric is related to the one of the full 4d metric by $h = g^{00} \cdot g$ and that the completely antisymmetric 3d tensor can be rewritten in terms of the 4d one as $h^{-1/2} \epsilon^{ijk} = (-g)^{-1/2} (-g^{00})^{-1/2} \epsilon^{0ijk}$.

¹⁵We can always make the coefficient in front of R time independent through an appropriate field redefinition $g_{\mu\nu} \rightarrow g_{\mu\nu} \cdot f(t)$. This corresponds, in the usual formalism, to going to Einstein frame.

where k is a constant which depends on the curvature of the spatial slices and the permutations are acting only on the last term. As such all the operators evaluated on the FRW background are themselves covariant operators, so that operators like $\delta K_{\mu\nu}$ and $\delta R_{\mu\nu\rho\sigma}$ are well defined covariant operators which vanish on the given FRW background and start linear in the perturbations. We stress that this possibility of rewriting the tensors evaluated on the background holds only because of the high degree of symmetry of the FRW background and it would not be true if one were interested in expanding around a less symmetric solution, e.g. a non-homogeneous background.

In equation (B.1) only the displayed operators contain linear terms in the fluctuations, so that the coefficients of $\sqrt{-g} g^{00}$ and $\sqrt{-g}$ are uniquely determined by the background FRW solution.

Let us now see how the Lagrangian can always be cast in the form (B.1). If we take an operator composed by the contraction of two tensors T and G (the generalization with more tensors is straightforward) we can write

$$TG = \delta T \delta G + T^{(0)}G + TG^{(0)} - T^{(0)}G^{(0)}. \quad (\text{B.4})$$

Let us discuss each term of the sum. The first one starts explicitly quadratic in the perturbation as we want. As we said, given the symmetries of the FRW background, the unperturbed tensors $T^{(0)}$ and $G^{(0)}$ can be written as functions of $g_{\mu\nu}$, n_μ and t . Therefore the last term $T^{(0)}G^{(0)}$ is just a polynomial of g^{00} with time dependent coefficients; it contains the terms $\sqrt{-g} g^{00}$ and $\sqrt{-g}$ plus operators which start explicitly quadratic in the perturbations. We are left with tensors of the form $T^{(0)}G$. We want to prove that also these terms can be written as the linear operators in eq. (B.1) plus operators that start quadratic in the fluctuations. By construction G will be linear either in $K_{\mu\nu}$ or $R_{\mu\nu\rho\sigma}$ with covariant derivatives acting on them. Covariant derivatives can be dealt with by successive integration by parts, letting them act on $T^{(0)}$ and the time dependent coefficient of the operator. In doing so we can generate extrinsic curvature terms. In this case we can reiterate eq. (B.4) until no covariant derivatives are left.¹⁶ We are thus left with the only possible scalar linear terms with no covariant derivatives: $K^\mu{}_\mu$ and R^{00} . Both of them can be rewritten in a more useful form. We can integrate by parts the extrinsic curvature term

$$\int d^4x \sqrt{-g} f(t) K^\mu{}_\mu = \int d^4x \sqrt{-g} f \nabla_\mu n^\mu = - \int d^4x \sqrt{-g} n^\mu \partial_\mu f = \int d^4x \sqrt{-g} \sqrt{-g^{00}} \dot{f}. \quad (\text{B.5})$$

While we can deal with R^{00} using the following relationship [32]:

$$(-g^{00})^{-1} R^{00} = R_{\mu\nu} n^\mu n^\nu = K^2 - K_{\mu\nu} K^{\mu\nu} - \nabla_\mu (n^\mu \nabla_\nu n^\nu) + \nabla_\nu (n^\mu \nabla_\mu n^\nu). \quad (\text{B.6})$$

The last two terms can again be integrated by parts:

$$\int d^4x \sqrt{-g} f(t) \nabla_\mu (n^\mu \nabla_\nu n^\nu) = - \int d^4x \sqrt{-g} \partial_\mu f n^\mu K^\nu{}_\nu, \quad (\text{B.7})$$

$$\int d^4x \sqrt{-g} f(t) \nabla_\nu (n^\mu \nabla_\mu n^\nu) = - \int d^4x \sqrt{-g} \partial_\nu f n^\mu \nabla_\mu n^\nu = 0 \quad (\text{B.8})$$

¹⁶There can be also powers of g^{00} from $T^{(0)}$. We can deal with them by writing $g^{00} = -1 + \delta g^{00}$ and thus generating additional contributions to the the g^{00} operator in eq. (B.1) plus terms which are explicitly quadratic or more in the perturbations.

where in the last passage we have used that $\partial_\nu f \propto n_\nu$. This shows that K_μ^μ and R^{00} can be written in terms of the linear operators of eq. (B.1) plus invariant terms that starts quadratically in the fluctuations.

In conclusion, we have shown that the most general Lagrangian of a theory with broken time diffeomorphisms around a given FRW background can be written in the form:

$$S = \int d^4x \sqrt{-g} \left[\frac{1}{2} M_{\text{Pl}}^2 R + M_{\text{Pl}}^2 \left(\dot{H} - \frac{k}{a^2} \right) \cdot g^{00} - M_{\text{Pl}}^2 \left(3H^2 + \dot{H} + 2\frac{k}{a^2} \right) + F^{(2)}(g^{00} + 1, \delta K_{\mu\nu}, \delta R_{\mu\nu\rho\sigma}; \nabla_\mu; t) \right] \quad (\text{B.9})$$

where $F^{(2)}$ starts quadratic in the arguments $g^{00} + 1$, $\delta K_{\mu\nu}$ and $\delta R_{\mu\nu\rho\sigma}$.

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