



## Geometric inflation

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### ABSTRACT

We argue that the presence of an inflationary epoch is a natural, almost unavoidable, consequence of the existence of a sensible action involving an infinite tower of higher-curvature corrections to the Einstein-Hilbert action. No additional fields besides the metric are required. We show that a family of such corrections giving rise to a well-posed cosmological evolution exists and automatically replaces the radiation-dominated early-universe Big Bang by a period of exponential growth of the scale factor, which is gracefully connected with standard late-time  $\Lambda$ CDM cosmology. The class of higher-curvature theories giving rise to sensible cosmological evolution share additional remarkable properties such as the existence of Schwarzschild-like non-hairy black holes, or the fact that, just like for Einstein gravity, the only degrees of freedom propagated on the vacuum are those of the standard graviton.

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The Einstein-Hilbert action is expected to be the first in an infinite series of higher-curvature terms. This is, for instance, what String Theory predicts, corrections being weighted by powers of  $\alpha'$  or the Planck length, e.g., [1–6]. Naively, one would expect the terms appearing in the corresponding stringy four-dimensional effective actions to give rise to sensible modifications of Einstein's gravity predictions, for example, yielding second-order equations for cosmological evolution. However, it is well-known that this is not the case in general—for instance, one can get the standard  $R^2$  Starobinsky model [7] corrected by a Weyl<sup>2</sup> term [8], which would spoil the nice behavior of the former. The idea is that such “inconsistencies” would be healed if one considers the full-fledged UV-complete theory; they would be artifacts resulting from truncating the series at some particular order in the parameter expansion. We do not control the exact mechanism by which, say, String Theory (or any other putative UV-complete description of gravity) deals with the infinite tower of effective higher-curvature terms at suf-

ficiently high energies. Whatever it is, though, the net result must be to produce sensible dynamics.

It is not always the case that we have at our disposal a UV-complete theory from which we can derive the low energy dynamics in a top-down approach. Hence, an alternative attempt at capturing such higher energy effects would consist in considering actions which already satisfy the requirement of producing sensible dynamics by themselves at all orders. This is in line with recent approaches to scrutinize candidates for low energy effective actions by demanding the absence of physical inconsistencies such as unitarity or causality violation order by order [9,10]. We shall follow this bottom-up approach with a broad question in mind: if any meaningful description of the cosmological dynamics can be expressed as an action for the metric tensor, order by order, what would be its predictions?

In this letter we analyze this question by constructing an infinite family of higher-curvature corrections to the Einstein-Hilbert action, selected by the criterium that they give rise to a well-posed cosmological problem. Interestingly, we find that the standard radiation-dominated decelerating early-universe predicted by Einstein's gravity (in the absence of additional mechanisms) is generically replaced by an inflationary era of exponential growth of the scale factor. For these “geometric inflation” models, the problem would be explaining a universe without inflation rather than

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the opposite. The effect of the higher-curvature terms becomes increasingly irrelevant as  $a(t)$  grows and, eventually, the evolution smoothly transits to standard post-inflationary  $\Lambda$ CDM cosmology. We stress that, in this setup, inflation takes place regardless of any extra (scalar) fields. This seems also relevant in the context of the recent discussions about inflationary models and the swampland conjectures [11,12].

The class of theories satisfying the well-posed cosmology criterion –whose first representative was constructed in [13]– is a subclass of a broader, recently identified, family [14–20], characterized by additional remarkable properties such as possessing second-order equations of motion on maximally symmetric backgrounds, as well as the existence of non-hairy generalizations of the Schwarzschild black hole (see below).

Before presenting the explicit details of our construction, some comments are in order. First, we emphasize that the strategy followed here differs from the usual effective field theory one, where one would include all possible higher-curvature invariants (up to field redefinitions) [21,22]. Ours should be regarded as a bottom-up approach in which the terms included are chosen in such a way that when the infinite tower of terms is considered, this is guaranteed to yield sensible cosmological dynamics in an explicit and computationally accessible way. That said, it is in fact also possible that the class of terms considered here actually provides a basis of higher-curvature operators for a general effective action. Indeed, after the first version of this paper was released, a proof was presented in [23,24] that any higher-curvature effective action can be mapped via field redefinitions to a “Generalized quasi-topological” (GQT) gravity [17]. The terms considered in the present paper represent a subset of all GQT theories. Hence, it could still be that one requires the full set of terms. We do suspect, however, that the class of terms belonging to the subset considered here does suffice to provide such a basis. Finally, regarding how well our approach may capture the physics of putative gravitational UV completions, we would like to point out Ref. [25], which appeared soon after the first version of our paper. There, the authors showed that using T-duality invariance at all orders in  $\alpha'$  as a guiding principle in an effective description of the bosonic sector of String Theory, the modified Friedmann equations obtained from the corresponding higher-curvature corrections turn out to be strikingly similar to the ones found within our framework –see (5) and (6) below.

## 1. The theory

At every order in curvature, and up to multiplicities, there exists a density  $\mathcal{R}_{(n)}$  constructed from contractions of the metric and the Riemann tensor such that the theory defined by the action

$$S = \int \frac{d^4x \sqrt{|g|}}{16\pi G} \left\{ -2\Lambda + R + \sum_{n=3}^{\infty} \lambda_n L^{2n-2} \mathcal{R}_{(n)} \right\}, \quad (1)$$

satisfies the following properties: i) it possesses second-order linearized equations around any maximally symmetric background; ii) it admits bona-fide non-hairy generalizations of the Schwarzschild(-AdS) black hole (and, more generally, of the Einstein gravity Taub-NUT/bolt solutions) characterized by a single function,  $g_{tt}g_{rr} = -1$ , and whose thermodynamic properties can be accessed in a fully analytic fashion; iii) it possesses a well-posed cosmological initial-value problem, namely, it admits cosmological Friedmann-Lemaître-Robertson-Walker (FLRW) solutions

$$ds^2 = -dt^2 + a(t)^2 \left( \frac{dr^2}{1-kr^2} + r^2 d\Omega^2 \right), \quad (2)$$

where the associated generalized Friedmann equations for the scale factor  $a(t)$  are second-order.

Notice that in (1) we have introduced dimensionless couplings  $\lambda_n$ , as well as a new energy scale  $\sim L^{-1}$ , which we will tacitly assume to be below the Planck scale,  $L_{\text{pl}}^{-1}$ . If  $L^{-1} \ll L_{\text{pl}}^{-1}$  and  $\lambda_3 \neq 0$ , the theory is afflicted by causality issues, which can be seen as a fingerprint of the presence of an infinite tower of massive higher spin particles [10]. We will assume that  $L^{-1} \lesssim L_{\text{pl}}^{-1}$  and will be more explicit on the values of  $\lambda_n$  below.

The discovery of four-dimensional theories satisfying requirements i) and ii) was ignited by the construction of Einsteinian cubic gravity [14–16]. By now, the existence of a family of theories satisfying such properties –including terms with arbitrarily high orders in curvature– and their connection to previously known (and some unknown) higher-dimensional theories such as Lovelock [26,27] or Quasi-topological gravities [28–31] is well established [17–20,32–37]. The observation that a subclass of such Generalized quasi-topological theories also satisfies iii) –at least at cubic order– was put forward in [13], and subsequently extended to the quartic and quintic orders in [37].<sup>1</sup> The cubic representative turns out to involve a simple linear combination of the Einsteinian cubic gravity density  $\mathcal{P} = 12R_a^c d R_c^e f R_e^a b + R_{ab}^{cd} R_{cd}^{ef} R_{ef}^{ab} - 12R_{abcd} R^{ac} R^{bd} + 8R_a^b R_c^d R_c^a R_b^d$  and a previously characterized [17] invariant  $\mathcal{C} = R_{abcd} R^{abc} R^{de} - \frac{1}{4} R_{abcd} R^{abcd} R - 2R_{abcd} R^{ac} R^{bd} + \frac{1}{2} R_{ab} R^{ab} R$ , which is trivial when evaluated on a static and spherically symmetric ansatz; the exact combination reads  $\mathcal{R}_{(3)} \propto \mathcal{P} - 8\mathcal{C}$ . In the Supplement we present the explicit form of  $\mathcal{R}_{(n)}$  up to  $n = 8$  as well as a rationale to construct these terms in general.

## 2. Generalized Friedmann equations

When evaluated on a FLRW ansatz of the form (2), the full non-linear equations of (1) reduce to a couple of second-order differential equations for the scale factor. Focusing on flat spatial metrics,  $k = 0$ , the generalized Friedmann equations read<sup>2</sup>

$$3F(H) = 8\pi G \rho + \Lambda, \quad (5)$$

$$-\frac{\dot{H}}{H} F'(H) = 8\pi G (\rho + P), \quad (6)$$

where

$$F(H) \equiv H^2 + L^{-2} \sum_{n=3}^{\infty} (-1)^n \lambda_n (LH)^{2n}, \quad (7)$$

and  $F'(H) \equiv dF(H)/dH$ . Here,  $H \equiv \dot{a}/a$  is the usual Hubble parameter, and  $\rho$  and  $P$  are the density and pressure of a perfect fluid, whose energy-momentum tensor appears in the right-hand-side of the modified Einstein's equations. Observe that if we set all the higher-curvature couplings to zero,  $F(H) = (\dot{a}/a)^2$ ,  $\dot{H}F'(H)/H = 2(\ddot{a} - \dot{a}^2)/a^2$ , and (5) and (6) reduce to the usual Einstein-gravity versions. Note also that using both equations one obtains the conservation equation

$$\frac{d\rho}{dt} + 3H(P + \rho) = 0.$$

<sup>1</sup> See also [47] for a previous proposal of a ghost-free higher-curvature modification of gravity giving rise to an inflationary epoch.

<sup>2</sup> In the general case,  $k \neq 0$ , the equations read

$$\Lambda - 3F(h) = -8\pi G \rho, \quad (3)$$

$$\Lambda - \left[ \frac{F'(h)}{H} \frac{dh}{dt} + 3F(h) \right] = +8\pi G P, \quad (4)$$

where we introduced  $h \equiv \sqrt{H^2 + k/a^2}$ .

Assuming that the fluid is composed of matter ( $P_m = 0$ ) and radiation ( $P_r = \rho_r/3$ ), one gets the familiar expression for the energy density

$$\rho = \rho_m + \rho_r = \frac{3F(H_0)}{8\pi G} \left[ \frac{\Omega_m}{a^3} + \frac{\Omega_r}{a^4} \right], \quad (8)$$

where  $\Omega_m = \rho_m/\rho_c$  and  $\Omega_r = \rho_r/\rho_c$ . As in [13], we redefine the critical density as  $\rho_c = 3F(H)/(8\pi G)$ , in order to avoid possible misinterpretations of our models as some kind of dark fluid.

Interestingly enough, the cosmological evolution of this class of theories is entirely dictated by the function  $F(H)$ ,<sup>3</sup> which is in turn controlled by  $L^{-1}$  and the  $\lambda_n$ . Ultimately, one would like to be able to compute those couplings from a putative UV-complete theory. In that respect, notice that there are ambiguities in the definition of the  $\mathcal{R}_{(n)}$  appearing in (1): there exist distinct densities  $\mathcal{R}_{(n)}^A$  and  $\mathcal{R}_{(n)}^B$  differing by a third density  $\mathcal{T}_{(n)}^{AB}$  with the property that  $\mathcal{T}_{(n)}^{AB}$  makes no contribution to the field equations for the classes of metrics considered here (see the Supplement). If we call  $\lambda_n^A$  and  $\lambda_n^B$  the gravitational couplings multiplying these two densities, the only effect in the cosmology will be the appearance of either coupling in (7). That is, it is tantamount an ambiguity in the numerical value of the  $n$ -th gravitational coupling. We expect that this is fixed by either an explicit UV-completed theory or by further consistency-checks that the theory must overcome. In the meanwhile, we will extract generic conclusions from the very fact that there exists such a function and we will focus on two particular models for the sake of definiteness.

### 3. Explicit models

At this level, the gravitational couplings  $\lambda_n$  are free parameters in our Lagrangian. There are, nonetheless, certain reasonable constraints that can be imposed. On the one hand, requiring positive-mass black holes to exist, would fix the sign of the first non-vanishing coupling [20], e.g.,  $\lambda_3 > 0$ . It is also convenient to choose  $F(H)$  to be a bijective function. This avoids pathological situations such as the absence of solutions for high enough energy densities, or the appearance of singularities at points for which  $F(H)$  reaches an extremum. The simplest way of satisfying these requirements consists in setting all odd-order couplings to zero and imposing the even ones to be positive, i.e.,  $\lambda_{2k+1} = 0$ ,  $\lambda_{2k} > 0$ , for all  $k \in \mathbb{Z}^+$ . The odd ones can be safely included if they are negative (except for  $\lambda_3$ ), or positive but sufficiently small. Besides this, we can also choose the relative values of the different couplings. For concreteness, in this letter we will consider two models (which we denote ‘‘GeomInf 1’’ and ‘‘GeomInf 2’’ henceforth) to be compared with the standard  $\Lambda$ CDM one, namely

$$\text{GeomInf 1: } \lambda_{2k+1} = 0, \quad \lambda_{4+2k} = \lambda_4/k!, \quad k \in \mathbb{Z}^+, \quad (9)$$

$$\text{GeomInf 2: } \lambda_3 > 0, \quad \lambda_{n \geq 4} = (-1)^n \lambda_3 / (n-4)!. \quad (10)$$

The above choices simplify the functional form of  $F(H)$ , making it possible to sum the corresponding infinite series (many other summable choices are possible). One finds  $F(H) = H^2 + \lambda_4 H^8 L^6 e^{(HL)^4}$  for the first and  $F(H) = H^2 - \lambda_3 L^4 H^6 [1 - (HL)^2 e^{(HL)^2}]$  for the second. With regards to  $L^{-1}$ , the most reasonable choice seems to be that it corresponds to some new scale below the Planck mass, but high enough to make the higher-curvature effects become negligible at late times.

<sup>3</sup> This is reminiscent of what happens for black holes in Lovelock gravity [48]; indeed, Lovelock cosmologies display a similar structure [49].

### 4. Early universe cosmology

For small values of the scale factor, radiation dominates over matter and dark energy. In that regime, the standard Einstein gravity scaling is  $a(t) \sim t^{1/2}$ , which follows straightforwardly from (5) and (8) for  $F(H) = H^2$ . Without additional mechanisms, this is obviously unsatisfactory from an inflationary perspective, as it predicts a negative acceleration  $\ddot{a}/a \sim -1/(4t^2)$ . The introduction of higher-order terms changes this behavior dramatically. Consider first the case in which we truncate the series in (1) at a certain order in curvature  $n = n_{\max}$ , as in [13]. Neglecting all contributions but radiation, the generalized Friedmann equations can be solved parametrically as<sup>4</sup>

$$a(p) = a(p_0) \left( \frac{F(p_0)}{F(p)} \right)^{1/4}, \quad t(p) = - \int_{p_0}^p \frac{dq F'(q)}{4q F(q)}. \quad (12)$$

For a given  $F(H)$ , it is straightforward to obtain  $a(t)$  from these equations. The limit  $a \rightarrow 0$  implies  $F(p) \rightarrow \infty$ , which for the kind of models we are considering corresponds to  $p \rightarrow \infty$ . We find  $F(p) \sim L^{2(n_{\max}-1)} p^{2n_{\max}}$  and  $t(p) \sim n_{\max}/(2p)$ , where the proportionality constants depend on the particular model. Putting this together in (12), we find

$$a(t) \sim t^{n_{\max}/2} \quad \text{when } t \rightarrow 0. \quad (13)$$

Of course, this reduces to the Einstein gravity result for  $n_{\max} = 1$ . Crucially, the introduction of higher-curvature terms changes the sign of the scale factor acceleration making it positive, namely  $\ddot{a}/a \sim n_{\max}(n_{\max} - 2)/(4t^2)$ . Hence, positively accelerated expansion occurs provided

$$\ddot{a}(t) > 0 \Leftrightarrow n_{\max} > 2. \quad (14)$$

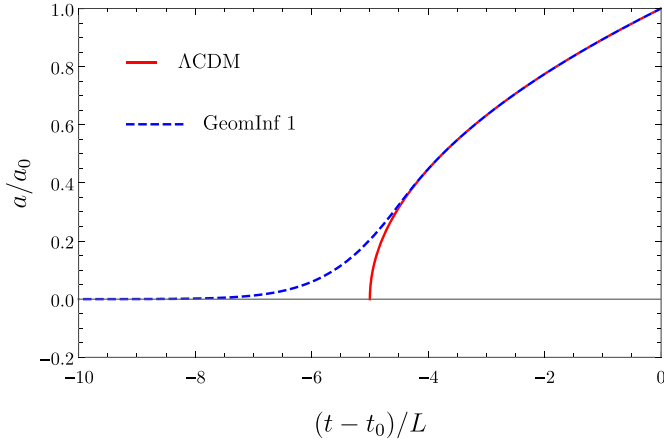
The larger  $n_{\max}$ , the faster the acceleration at early times. This fact can also be expressed in terms of the slow-roll parameter,  $\epsilon \equiv -\dot{H}/H^2 = 1 - \ddot{a}/(aH^2)$ , for which we find  $\epsilon \sim 2/n_{\max}$ . Positive acceleration occurs for  $0 < \epsilon < 1$ , which of course translates into (14).

On the other hand, when we only include a finite number of terms, the expansion is polynomial rather than exponential, and at some finite moment in the past a singularity would be reached when  $a = 0$ , albeit the series truncation certainly ceases to be trustworthy when the characteristic curvature scale becomes of order  $\sim L$ . This suggests that, when dealing with early time cosmology, it is not consistent to truncate the higher-order terms, as they will become relevant before reaching the singularity. One can immediately see that when the full tower of higher-curvature terms is included, the scale factor will grow faster than any polynomial near  $a = 0$ . Therefore, the expansion will be at least exponential. This can be verified explicitly for our two models in (9) and (10). In particular, when  $t \rightarrow -\infty$  we find  $a(t) \sim e^{-(-3t/L)^{4/3}/4}$  for GeomInf 1, and  $a(t) \sim e^{-t^2/L^2}$  for GeomInf 2. More generally, if  $F(H) \sim e^{(HL)^w}$ , one finds (up to multiplicative powers which are irrelevant for the order of magnitude)

<sup>4</sup> In fact, a parametric solution can be written in general. The result is

$$a(p) = J(p), \quad t(p) = \int_{p_0}^p \frac{dq J'(q)}{q J(q)}, \quad (11)$$

where  $J$  is the function  $J = P^{-1} \circ F$  and  $p_0$  is such that  $J(p_0) = 1$ .



**Fig. 1.** Scale factor  $a(t)$  as a function of time as predicted by the standard  $\Lambda$ CDM model and our geometric inflation model GeomInf 1 defined in (9) (for clarity reasons, we omit the curve corresponding to GeomInf 2 in this plot). The normalization of  $a$  is arbitrary. Here, we have normalized it by  $a_0$  corresponding to its value at some later time  $t_0$  at which the higher-curvature corrections become negligible. In particular, we have fixed this value so that  $H(t_0) = 0.1L^{-1}$ . While for the  $\Lambda$ CDM model the Big Bang singularity is reached at  $t - t_0 = -5L$ , the value  $a = 0$  is never reached at finite times for the geometric inflation model (only approached asymptotically as  $t \rightarrow -\infty$ ).

$$a(t) \sim \exp \left[ -\frac{1}{4} \left( \frac{4(1-w)t}{wL} \right)^{\frac{w}{w-1}} \right] \quad \text{if } w > 1, \quad (15)$$

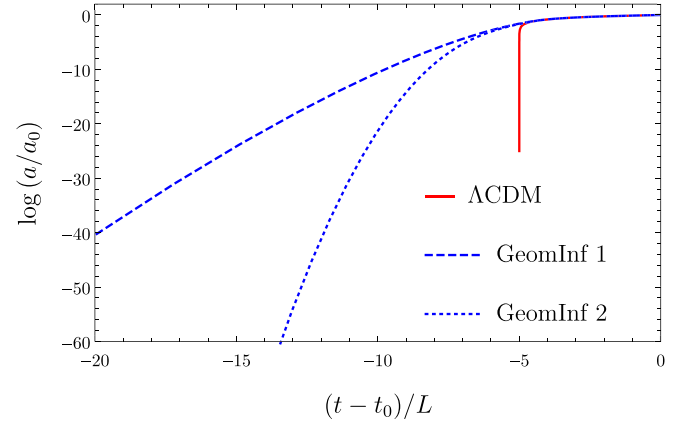
and  $a(t) \sim \exp[-e^{-4t/L}/4]$  if  $w = 1$ . Observe that for these models, there is no notion of *origin of time* at which we can set  $t = 0$ . Rather,  $a \rightarrow 0$  is only reached asymptotically as  $t \rightarrow -\infty$ .<sup>5</sup>

## 5. Geometric inflation

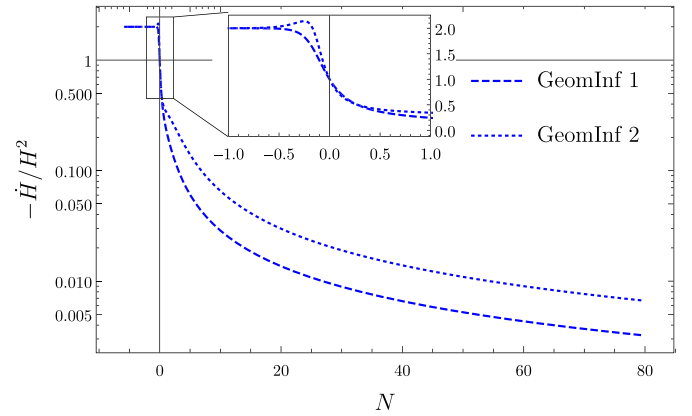
Let us now analyze our two models in more detail. In order to do so, we solve the generalized Friedmann equations in each case. In Fig. 1, we plot the scale factor as a function of time for the standard  $\Lambda$ CDM model and GeomInf 1 as defined in (9). The general intuitions observed in the previous section become manifest. In particular, we observe that the Big Bang singularity followed by a decelerating expansion predicted by Einstein's gravity for a radiation-dominated universe is replaced by an exponential growth of the scale factor which eventually connects with the Einstein gravity prediction at sufficiently late times, where the higher-curvature terms become increasingly irrelevant.

This behavior is even more manifest in the logarithmic plot of Fig. 2, where we also include GeomInf 2 in the comparison. Naturally, the exact details depend on the exact choice of parameters  $\lambda_n$  and energy scale  $L^{-1}$ , but the general message is that an inflationary epoch appears to be unavoidable, regardless of the specific model. It is also worth emphasizing that late-time cosmology can be essentially identical to the one predicted by Einstein gravity, as long as the new energy scale is high enough. Again, the exact point at which the crossover from the geometric inflation era to the standard  $\Lambda$ CDM-like phase occurs, depends on the exact details of each model, but a *graceful exit* is automatically implemented in all models. This smooth transition is manifest in Figs. 1 and 2.

In Fig. 3, we plot the slow-roll parameter  $\epsilon$  as a function of the number of  $e$ -folds. This is defined as  $N = \log(a/a_f)$ , where the reference value  $a_f$  denotes the value at which inflation ends, defined



**Fig. 2.** We plot  $\log(a(t)/a_0)$ . In geometric inflation models, the Big Bang is replaced by a period of exponential growth, which is made evident.



**Fig. 3.** Slow-roll parameter  $\epsilon = -\dot{H}/H^2$  as a function of the number of  $e$ -folds  $N = \log(a/a_f)$ . The reference value  $a_f$  is taken such that  $\epsilon(a_f) = 1$ , meaning that inflation ends. (Inset) Zoom of the previous plot where we observe that there is a smooth transition between the Einstein gravity value for radiation-dominated universes,  $\epsilon = 2$ , and the inflationary period, in which  $\epsilon \rightarrow 0$ . Geometric inflation provides a graceful exit to the inflationary period.

so that  $\epsilon(a_f) = 1$ . For our class of theories, these two quantities can be written in terms of the function  $F(H)$  defined in (7) as

$$N = \frac{1}{4} \log \left( \frac{F(H)}{F(H_f)} \right), \quad \epsilon = \frac{4F(H)}{HF'(H)}. \quad (16)$$

When  $\epsilon = 0$ ,  $H$  remains constant and the expansion is exponential—corresponding to a pure de Sitter space. When  $|\epsilon| \ll 1$ , which is the case for most values of  $N$  shown in the plot for our two models—and continues to be the case as  $N \rightarrow \infty$ —the growth is quasi-exponential. This is the typical behavior of standard slow-roll inflationary models (see, for example, [38]). After the end of inflation is reached, the slow-roll parameter smoothly connects with the radiation-dominated Einstein gravity result, corresponding to  $\epsilon = 2$ —see inset in Fig. 3.

As the  $\epsilon(N)$  curves make manifest, our models implement all nice features of slow-roll models without introducing any additional fields and/or potentials. The behavior of  $\epsilon$  for large  $N$  is  $\epsilon \sim 1/(4N)$  for GeomInf 1, and  $\epsilon \sim 1/(2N)$  for GeomInf 2. More generally, if  $F(H) \sim e^{(HL)^w}$ , we get  $\epsilon \sim 1/(wN)$ . However, observational quantities such as the spectral index will have a different

<sup>5</sup> Let us note that in the generic case a singularity remains in the limit  $t \rightarrow -\infty$ . However, we note that it is possible to achieve an explicit de Sitter phase at early

times through a judicious choice of function  $F(H)$ . For example, this will be the case if  $F(H) \rightarrow +\infty$  at some (finite)  $H = H_*$ .

dependence on  $\epsilon$  with respect to the slow-roll models driven by scalars. For that, one would need to perform a thorough analysis of perturbations.

## 6. Discussion

Under the present proposal, the evolution of the Universe from inflation to late time acceleration can be explained by the bottom-up construction of an effective geometric field theory of gravity. We solely demand consistence of the theory order-by-order in the series of higher-curvature terms. We showed that if the full series is taken into consideration, and only in that case, then the expansion of the Universe at early times has no other choice than becoming exponential. As proved in [13] for the cubic case, the evolution of the scale factor smoothly transits to standard post-inflationary  $\Lambda$ CDM cosmology.

The surprising performance of our theory in describing the accelerated expansion periods of the Universe is nothing but a first step towards building up a new cosmological framework. Further checks will represent an increasingly demanding challenge; for instance, the study of cosmological perturbations. Remarkably, a preliminary analysis of linear perturbations shows that those also satisfy second-order dynamics<sup>6</sup>—this was anticipated in [37] for the cubic theory— which opens up the window for a more thorough study of observational tests such as the signatures in the Cosmic Microwave Background (CMB) [39].

The features above described are achieved without the introduction of any additional fields besides the metric. As opposed to other alternatives studied in the cosmological context—see e.g., [40,41]—, our theories cannot be thought of as scalar-tensor models. In fact, on maximally symmetric backgrounds they only propagate the usual transverse and traceless graviton of General Relativity so, in a sense, they are as “pure-metric” as it gets amongst higher-curvature theories. On this respect, our proposal sheds a different light on discussions such as the role of reheating and the status of the string swampland conjecture [11,12] under the present geometrical scheme—perhaps scalar fields play a less important role in the early universe than our current expectations suggest.

A better understanding on the multiplicities of the  $\mathcal{R}_{(n)}$  densities and the related ambiguities in the actual numerical values of the gravitational couplings  $\lambda_n$  seems necessary. It was shown in [37] that the cubic density  $\mathcal{R}_{(3)}$  constructed in [13] is the dimensional reduction of so-called five dimensional quasi-topological gravity [28,29]. It is intriguing to explore whether there is a network connecting these classes of theories in different spacetime dimensions that may ultimately provide hints of a hidden mathematical structure. A complementary viewpoint would be to further scrutinize the  $\mathcal{R}_{(n)}$  densities to better understand whether the triviality of the ambiguities—when evaluated on spherically symmetric and cosmological ansätze—is due to some sort of integrable structure. The value of the effective couplings,  $\lambda_n$ , could be otherwise constrained by using observations (SNIa [42], [43],  $H_0$  [44], BAO's [45], [46]).

The very fact that a purely geometric mechanism triggers inflation is tantalizing. We expect to contribute further to put together more pieces of this exciting puzzle.

<sup>6</sup> On general grounds, it is expected that higher-order time derivatives still appear at non-linear level in perturbation theory. One should study in that case whether there is a strong coupling problem, or if, on the contrary, the non-linear terms can be safely neglected.

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## Appendix A. Supplementary material

Supplementary material related to this article can be found online at <https://doi.org/10.1016/j.physletb.2020.135242>.

## References

- [1] D.J. Gross, J.H. Sloan, Nucl. Phys. B 291 (1987) 41.
- [2] D.J. Gross, E. Witten, Nucl. Phys. B 277 (1986) 1.
- [3] M.T. Grisaru, D. Zanon, Phys. Lett. B 177 (1986) 347.
- [4] E.A. Bergshoeff, M. de Roo, Nucl. Phys. B 328 (1989) 439.
- [5] M.B. Green, C. Stahn, J. High Energy Phys. 09 (2003) 052, arXiv:hep-th/0308061 [hep-th].
- [6] S. Frolov, I.R. Klebanov, A.A. Tseytlin, Nucl. Phys. B 620 (2002) 84, arXiv:hep-th/0108106 [hep-th].
- [7] A.A. Starobinsky, Phys. Lett. B 91 (99) (1980).
- [8] L. Alvarez-Gaume, A. Kehagias, C. Kounnas, D. Lust, A. Riotto, Fortschr. Phys. 64 (2016) 176, arXiv:1505.07657 [hep-th].
- [9] A. Adams, N. Arkani-Hamed, S. Dubovsky, A. Nicolis, R. Rattazzi, J. High Energy Phys. 10 (2006) 014, arXiv:hep-th/0602178 [hep-th].
- [10] X.O. Camanho, J.D. Edelstein, J. Maldacena, A. Zhiboedov, J. High Energy Phys. 02 (2016) 020, arXiv:1407.5597 [hep-th].
- [11] G. Obied, H. Ooguri, L. Spodyneiko, C. Vafa, arXiv:1806.08362 [hep-th], 2018.
- [12] P. Agrawal, G. Obied, P.J. Steinhardt, C. Vafa, Phys. Lett. B 784 (2018) 271, arXiv:1806.09718 [hep-th].
- [13] G. Arciniega, J.D. Edelstein, L.G. Jaime, arXiv:1810.08166 [gr-qc], 2018.
- [14] P. Bueno, P.A. Cano, Phys. Rev. D 94 (2016) 104005, arXiv:1607.06463 [hep-th].
- [15] R.A. Hennigar, R.B. Mann, Phys. Rev. D 95 (2017) 064055, arXiv:1610.06675 [hep-th].
- [16] P. Bueno, P.A. Cano, Phys. Rev. D 94 (2016) 124051, arXiv:1610.08019 [hep-th].
- [17] R.A. Hennigar, D. Kubizňák, R.B. Mann, Phys. Rev. D 95 (2017) 104042, arXiv:1703.01631 [hep-th].
- [18] P. Bueno, P.A. Cano, Class. Quantum Gravity 34 (2017) 175008, arXiv:1703.04625 [hep-th].
- [19] J. Ahmed, R.A. Hennigar, R.B. Mann, M. Mir, J. High Energy Phys. 05 (2017) 134, arXiv:1703.11007 [hep-th].
- [20] P. Bueno, P.A. Cano, Phys. Rev. D 96 (2017) 024034, arXiv:1704.02967 [hep-th].
- [21] C.P. Burgess, Living Rev. Relativ. 7 (2004) 5, arXiv:gr-qc/0311082 [gr-qc].
- [22] C.P. Burgess, M. Williams, J. High Energy Phys. 08 (2014) 074, arXiv:1404.2236 [gr-qc].
- [23] P. Bueno, P.A. Cano, J. Moreno, A. Murcia, arXiv:1906.00987 [hep-th], 2019.
- [24] P. Bueno, P.A. Cano, R.A. Hennigar, arXiv:1909.07983 [hep-th], 2019.
- [25] O. Hohm, B. Zwiebach, arXiv:1905.06963 [hep-th], 2019.
- [26] D. Lovelock, Aequ. Math. 4 (1970) 127.
- [27] D. Lovelock, J. Math. Phys. 12 (1971) 498.
- [28] J. Oliva, S. Ray, Class. Quantum Gravity 27 (2010) 225002, arXiv:1003.4773 [gr-qc].
- [29] R.C. Myers, B. Robinson, J. High Energy Phys. 08 (2010) 067, arXiv:1003.5357 [gr-qc].
- [30] M.H. Dehghani, A. Bazrafshan, R.B. Mann, M.R. Mehdizadeh, M. Ghanaatian, M.H. Vahidinia, Phys. Rev. D 85 (2012) 104009, arXiv:1109.4708 [hep-th].
- [31] A. Cisterna, L. Guajardo, M. Hassaine, J. Oliva, J. High Energy Phys. 04 (2017) 066, arXiv:1702.04676 [hep-th].
- [32] X.-H. Feng, H. Huang, Z.-F. Mai, H. Lu, Phys. Rev. D 96 (2017) 104034, arXiv:1707.06308 [hep-th].
- [33] R.A. Hennigar, J. High Energy Phys. 09 (2017) 082, arXiv:1705.07094 [hep-th].

- [34] Y.-Z. Li, H.-S. Liu, H. Lu, J. High Energy Phys. 02 (2018) 166, arXiv:1708.07198 [hep-th].
- [35] P. Bueno, P.A. Cano, A. Ruiperez, J. High Energy Phys. 03 (2018) 150, arXiv:1802.00018 [hep-th].
- [36] P. Bueno, P.A. Cano, R.A. Hennigar, R.B. Mann, J. High Energy Phys. 10 (2018) 095, arXiv:1808.01671 [hep-th].
- [37] A. Cisterna, N. Grandi, J. Oliva, arXiv:1811.06523 [hep-th], 2018.
- [38] D. Baumann, in: Physics of the Large and the Small, TASI 09, Proceedings of the Theoretical Advanced Study Institute in Elementary Particle Physics, Boulder, Colorado, USA, 1-26 June 2009 (2011), pp. 523–686, arXiv:0907.5424 [hep-th].
- [39] P.A.R. Ade, et al., Planck, Astron. Astrophys. 594 (2016) A13, arXiv:1502.01589 [astro-ph.CO].
- [40] T.P. Sotiriou, V. Faraoni, Rev. Mod. Phys. 82 (2010) 451, arXiv:0805.1726 [gr-qc].
- [41] G.W. Horndeski, Int. J. Theor. Phys. 10 (1974) 363.
- [42] A.G. Riess, et al., Supernova Search Team, Astron. J. 116 (1998) 1009, arXiv:astro-ph/9805201 [astro-ph].
- [43] S. Perlmutter, et al., Supernova Cosmology Project, Astrophys. J. 517 (1999) 565, arXiv:astro-ph/9812133 [astro-ph].
- [44] O. Farooq, B. Ratra, Astrophys. J. 766 (2013) L7, arXiv:1301.5243 [astro-ph.CO].
- [45] Astron. Astrophys. 574 (2015) A59, arXiv:1404.1801 [astro-ph.CO].
- [46] S. Alam, et al., BOSS, Mon. Not. R. Astron. Soc. 470 (2017) 2617, arXiv:1607.03155 [astro-ph.CO].
- [47] B. Tekin, arXiv:1602.08949 [gr-qc], 2016.
- [48] X.O. Camanho, J.D. Edelstein, Class. Quantum Gravity 30 (2013) 035009, arXiv:1103.3669 [hep-th].
- [49] X.O. Camanho, Lovelock gravity, black holes and holography, Ph.D. thesis, Santiago de Compostela U., 2013, arXiv:1509.08129 [hep-th].