Unity of Cosmological Inflation Attractors

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Recently, several broad classes of inflationary models have been discovered whose cosmological predictions, in excellent agreement with Planck, are stable with respect to significant modifications of the inflaton potential. Some classes of models are based on a nonminimal coupling to gravity. These models, which we call ξ attractors, describe universal cosmological attractors (including Higgs inflation) and induced inflation models. Another class describes conformal attractors (including Starobinsky inflation and T models) and their generalization to α attractors. The aim of this Letter is to elucidate the common denominator of these attractors: their robust predictions stem from a joint pole of order 2 in the kinetic term of the inflaton field in the Einstein frame formulation prior to switching to the canonical variables. Model-dependent differences only arise at subleading level in the kinetic term. As a final step towards the unification of the different attractors, we introduce a special class of ξ attractors which is fully equivalent to α attractors with the identification $\alpha = 1 + (1/6\xi)$. While r is generically predicted to be of the order $1/N^2$, there is no theoretical lower bound on r in this class of models.

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Introduction.—The data releases by Wilkinson Microwave Anisotropy Probe and Planck brought attention to a mysterious fact: two different models, the Starobinsky model [1] and the Higgs inflation model [2], make the same prediction, matching well with observational data—both of Planck2013 [3] as well as Planck2014. In the leading approximation in $1/N$, where N is the number of e-folds, the spectral index n_s and tensor-to-scalar ratio r are given by

$$n_s = 1 - \frac{2}{N}, \qquad r = \frac{12}{N^2}. \tag{1}$$

This could be a coincidence, but further investigation revealed the existence of several broad classes of different models having the same predictions in the leading approximation in $1/N$, practically independent of the details of the model.

The first class of these theories was conformal attractors [4], which described a broad variety of different models including the Starobinsky model. Further investigation revealed the existence of α attractors [5,6], which generalized the models of conformal attractors, but predicted, for not-too-large values of the parameter α, that

$$n_s = 1 - \frac{2}{N}, \qquad r = \frac{12\alpha}{N^2}. \tag{2}$$

The Lagrangian of the α-attractor models of a real scalar field ϕ looks as follows in the Einstein frame,

$$\mathcal{L}_E = \sqrt{-g}\left[\frac{1}{2}R - \frac{\alpha}{(1-\phi^2/6)^2}\frac{(\partial\phi)^2}{2} - \alpha f^2(\phi/\sqrt{6})\right]. \tag{3}$$

It was shown in Refs. [4–6] that the predictions (2) of this class of models are stable with respect to major changes of the inflaton potential, which has a functional freedom in terms of an arbitrary f. In this context, the Starobinsky model [1] corresponds to a special choice for this function with $\alpha = 1$.

Note that both the kinetic and potential energies have an overall coefficient α. While the former appears in all versions of α-attractor models, the latter is a matter of choice since the functions f are nearly arbitrary. However, by placing α in front of it, one reaches an important goal: while the parameter r is proportional to α, both the parameter n_s and the amplitude of the scalar perturbations of the metric are independent of it for this class of theories.

Another class of models [2,7,8] described cosmological attractors with a nonminimal coupling to gravity,

$$\mathcal{L}_J = \sqrt{-g}\left[\frac{1}{2}\Omega(\phi)R - \frac{1}{2}K_J(\phi)(\partial\phi)^2 - V_J(\phi)\right], \tag{4}$$

which we refer to as the Jordan frame. For $\Omega = 1 + \xi\phi^2$, $V_J = \lambda\phi^4$, and $K_J = 1$, it is described as the Higgs inflation [2]. In a more general class of models, one retains the same functional relation between the nonminimal coupling and scalar potential,

$$V_J(\phi) = \frac{\lambda}{\xi^2}(\Omega(\phi) - 1)^2 \qquad (5)$$

but allows for a different form of these functions. For instance, the universal attractor models are based on $\Omega = 1 + \xi f(\phi)$ with an arbitrary function f, and $K_J = 1$ [7].

In the class of induced inflation models [8], one has $\Omega = \xi f_{\mathrm{ind}}(\phi) > 0$ and $K_J = 1$. This class of theories is equivalent to the class of universal attractors up to the redefinition $f_{\mathrm{ind}}(\phi) = f(\phi) + \xi^{-1}$ [9]. However, it is convenient to consider these two classes of models separately by defining universal attractors as the theories where $\Omega = 1$ in the limit $\phi \to 0$ and induced inflation as the theories where $\Omega = 0$ in the limit $\phi \to 0$. The inflationary predictions of all of these models depend on ξ but coincide with Eq. (1) in the large-ξ limit and are stable with respect to certain further modifications of $V_J(\phi)$ to be discussed in this Letter. Other choices of Ω and K_J have been also discussed in the literature. In this Letter, we will call all models of this type ξ attractors.

In addition to models with one attractor point, there were double attractors [10]; their predictions interpolated between the predictions of α attractors with small α or induced inflation at large ξ and the predictions $r = 4(1 - n_s) = 8/N$ of the simplest chaotic inflation model $\frac{1}{2}m^2\phi^2$ in the opposite parameter limit.

Despite a deepening understanding of the nature of these models [9], a direct link between the models with nonminimal coupling and the α attractors was missing, and their predictions coincided with each other only in certain limits. In this Letter, we aim to clarify both the relations and differences between these models and unravel the origin of the robust inflationary predictions (2).

Kinetic formulation.—Our starting point is a simple observation that can be phrased as "The inflationary predictions of models whose kinetic term is given by a Laurent series are determined by the order and the residue of the leading pole of the series." In the above, we have assumed minimal coupling to gravity, i.e., Einstein frame, as well as a smooth scalar potential at the location of the pole. Such a model can be summarized as

$$\mathcal{L} = \sqrt{-g}\left[\frac{1}{2}R - \frac{1}{2}K_E(\rho)(\partial\rho)^2 - V_E(\rho)\right]. \qquad (6)$$

The case where K_E is given by a Laurent series (where we have assumed the pole to be located at $\rho = 0$ without loss of generality)

$$K_E = \frac{a_p}{\rho^p} + \cdots, \qquad V_E = V_0(1 + c\rho + \cdots) \qquad (7)$$

is particularly interesting: it corresponds to a fixed point of the inflationary trajectory, which is characterized almost completely by the properties of this point. Indeed, in the limit of a large number of e-folds, one can assume that only the leading pole in K_E is relevant. This leads to the simple relation (where we will assume $p > 1$ for simplicity)

$$N = \int \frac{a_p}{c\rho^p}d\rho \sim \frac{a_p\rho^{1-p}}{c(p-1)}. \qquad (8)$$

Upon inverting this relation, one can calculate the spectral index and tensor-to-scalar ratio at leading order in $1/N$,

$$n_s = 1 - \frac{p}{p-1}\frac{1}{N}, \qquad r = \frac{8c^{(p-2/p-1)}a_p^{(1/p-1)}}{(p-1)^{(p/p-1)}}\frac{1}{N^{(p/p-1)}}. \qquad (9)$$

Indeed, the spectral index depends solely on the order of the pole, while the tensor-to-scalar ratio also involves the residue. Note that this yields the same relation between the $1/N$ coefficient of the spectral index and the $1/N$ power of the tensor-to-scalar ratio as stressed in Ref. [11]. Moreover, the kinetic formulation defines not only the power of $1/N$ but also the coefficient in the above formula for r.

The above holds for all values of $p > 1$, for example, hilltop inflation models [12] where $V_E = V_0[1 - (\varphi/\mu)^n]$ with $p = 2 - 2/n$, where n can be both negative and positive $n \geq 2$. However, in what follows, we will be mainly interested in the case $p = 2$: it is singled out as it allows for a superconformal and supergravity description and arises as a consequence of a nonminimal coupling to gravity. In particular, we will show that all cosmological attractors can be brought to the form (6) with a kinetic term that has a pole or order 2 at a location where the scalar potential is perfectly smooth. In other words, all attractors have a common denominator in the Laurent expansion (7). In this case, the general pole predictions (9), indeed, lead to Eq. (2) with the identification $a_p = \frac{3}{2}\alpha$. This provides a unified approach to their cosmological predictions, independent of the structure of the inflationary potentials, provided these are smooth at the point $\rho = 0$.

α attractors.—To demonstrate the equivalence of the above to α attractors, we start from the original formulation of the theory of conformal attractors and α attractors [4–6] given in a noncanonical field ϕ as Eq. (3). Its kinetic term has two poles of order 2, related by symmetry $\phi \to -\phi$. Without loss of generality, we will focus on the pole located at $\phi = \sqrt{6}$. Expanding around this pole, we find a Laurent expansion

$$K_E = \frac{3\alpha}{2}\frac{1}{(\phi - \sqrt{6})^2} - \frac{\sqrt{6}\alpha}{4}\frac{1}{\phi - \sqrt{6}} + \cdots. \qquad (10)$$

Indeed, we find the same leading pole of order 2 with residue $\frac{3}{2}\alpha$ in addition to subleading terms. Similarly, for a generic choice of the function f, the scalar potential is a Taylor series around the point $\phi = \sqrt{6}$.

By means of field redefinitions, one can change the form of the subleading terms and trade certain subleading corrections to others. For instance, in this case, one can redefine the field ϕ into a new variable ρ, such that the kinetic term becomes only a pole in ρ, without additional terms. This can be performed by

$$\frac{\phi}{\sqrt{6}} = \frac{1-\rho}{1+\rho}. \tag{11}$$

The Lagrangian of the α-attractor models (3) in the new variables ρ has

$$K_E = \frac{3\alpha}{2}\frac{1}{\rho^2}, \qquad V_E = \alpha f^2\left(\frac{1-\rho}{1+\rho}\right). \tag{12}$$

Finally, one can go to a canonical field φ with $K_E = 1$, where the scalar potential reads $V_E = \alpha f^2[\tanh(\varphi/\sqrt{6\alpha})]$. For $\alpha = 1$ and monomial functions f, they coincide with the T models from the theory of conformal attractors [4].

Note that the kinetic terms blow up at $\phi = \sqrt{6\alpha}$ or $\rho = 0$. While the subleading corrections are different, both cases have the same leading term; this corresponds to a pole of order 2 with residue $3\alpha/2$. It is this singularity that is responsible for the stability of predictions of these theories (2) with respect to strong deformations of the inflationary potential near the boundary of the moduli space at $\rho = 0$. Subleading corrections in either the Laurent expansion of the kinetic term or the Taylor expansion of the potential term are irrelevant in the large-N limit.

In terms of the canonical scalar field, this boundary is located at φ close to infinity. For generic functions f, the scalar potential will asymptote to a plateau at infinity and will have an exponentially suppressed falloff with leading term $e^{-\sqrt{2/3\alpha}\varphi}$. It is this leading term that determines all inflationary properties at large N.

Nonminimal coupling and special attractors.— Similarly, there is an interesting relation to ξ attractors based on a nonminimal coupling between the gravitational and inflationary sector. Therefore, we generalize our starting point to the Jordan frame (4). By means of a conformal transformation for $\Omega > 0$, it can be brought to the Einstein frame with

$$K_E = \left(\frac{3\Omega'^2}{2\Omega^2} + \frac{K_J}{\Omega}\right), \qquad V_E = \frac{V_J(\phi)}{\Omega^2}. \tag{13}$$

So far, only models with $K_J = 1$ have been considered, where the parameter ξ was a part of the choice of the function $\Omega(\phi)$ in Eq. (4). Now we will define a new class of theories, which we will call special attractors. They will be defined by the following choice of functions in Eq. (4),

$$K_J = \frac{1}{4\xi}\frac{(\Omega')^2}{\Omega}, \qquad V_J(\phi) = \Omega^2 U(\Omega). \tag{14}$$

Thus, we absorbed the ξ dependence into the factor K_J. Then the theory (4) in the Einstein frame becomes

$$\mathcal{L}_E = \sqrt{-g}\left[\frac{R}{2} - \frac{3\alpha}{4}\left(\frac{\partial\Omega}{\Omega}\right)^2 - U(\Omega)\right], \qquad \alpha \equiv 1 + \frac{1}{6\xi}. \tag{15}$$

In this theory, Ω becomes the field variable. Its kinetic term is exactly of the form (7) with a pole of order 2 and no subleading corrections. However, physically this does not correspond to the same limit; while the α attractors derive their attractor predictions from the region close to $\rho = 0$, inflation in the ξ attractors takes place at very large Ω very large. Therefore, it is natural to identify

$$\rho(\phi) = \Omega^{-1}(\phi). \tag{16}$$

Note that a pole of order 2 is exactly invariant under this redefinition and retains the same form.

In order for the kinetic energy to be well defined, one has to require that α is positive. There are three regions of the parameter ξ; the condition $\alpha > 0$ is satisfied in the first two of them: (i) $\xi > 0$, with $\alpha > 1$, or (ii) $-\infty < \xi < -\frac{1}{6}$ corresponding to $0 < \alpha < 1$, while (iii) intermediate regions with $-1/6 < \xi < 0$ lead to a wrong sign of the Einstein frame kinetic term. The limiting case with $\alpha = 1$ can be reached either in the limit $\xi \to \infty$ or $\xi \to -\infty$, while $\alpha = 0$ is accessible via $\xi \to -1/6$ from below.

It is important to take stock of the situation at this point. In particular, one can allow ξ to become negative (and α smaller than 1) at a very specific price: the Jordan frame kinetic term (14) has the wrong sign. While this could seem dangerous, for $-\infty < \xi < -\frac{1}{6}$ this danger is, in fact, fictitious, as it does not lead to negative kinetic terms and instability in the Einstein frame.

This phenomenon is reminiscent of the Breitenlohner-Freedman bound in anti–de Sitter space. In that case, an apparent instability due to a negative mass can be cured by the nontrivial geometry, provided the mass satisfies the Breitenlohner-Freedman bound [13]. In our case, an apparent instability due to a negative kinetic energy can be cured by the nonminimal coupling in Jordan frame, provided the coefficient $1/(4\xi)$ of the negative term in Eq. (14) is sufficiently small such that α is positive.

One can represent the theory (15) in terms of a canonically normalized inflaton field φ defined by $\Omega = e^{\sqrt{2/3\alpha}\varphi}$ as

$$\mathcal{L}_E = \sqrt{-g}\left[\frac{R}{2} - \frac{1}{2}(\partial\varphi)^2 - U(e^{\sqrt{(2/3\alpha)}\varphi})\right]. \tag{17}$$

For the special choice $U(\Omega) = \alpha f^2(1 - \Omega/1 + \Omega)$, this theory coincides with the class of α attractors defined in Eq. (12), with $V_E = \alpha f^2[\tanh(\varphi/\sqrt{6\alpha})]$. In particular, for the simplest choice $f(x) = cx$, where c is some constant, one finds the α generalization of the simplest T-model potential [4,6]

$$V = \alpha c^2 \tanh^2 \frac{\varphi}{\sqrt{6\alpha}}. \tag{18}$$

For $f(x) = (cx/1 + x)$, which is equivalent to the choice $V_J = c^2(\Omega - 1)^2$, one finds the α-β model [5]

$$V = \alpha c^2 (1 - e^{-\sqrt{(2/3\alpha)}\varphi})^2, \tag{19}$$

which generalizes the Starobinsky potential. More general choices of potentials are possible; e.g., one can add to $U(\Omega)$ corrections

$$\Delta U(\Omega) = \sum_{i=2}^{\infty} c_i \Omega^{-i} = \sum_{i=2}^{\infty} c_i \rho^i. \tag{20}$$

This results in the subleading corrections in $e^{\sqrt{(2/3\alpha)}\varphi}$, which do not affect the inflationary predictions in the large-N limit.

Induced inflation.—Induced inflation is defined by Eq. (4) with $\Omega = \xi f(\phi)$ and the scalar potential given by the usual relation (5). This theory is well defined (i.e., it describes gravity instead of antigravity) only for $\Omega > 0$. Without any loss of generality, one can define this class of theories by requiring $\xi > 0$, $f(\phi) > 0$. Then, independent of the function $f(\phi)$, which, in principle, can be chosen arbitrary, the inflationary predictions of this model coincide with Eq. (1) in the limit of $\xi \to +\infty$ [8]. Moreover, in the opposite limit $\xi \to 0$, the predictions approximate those of quadratic inflation, again independent of the functional choice [10].

Remarkably, for the special case $\Omega = \xi\phi^2$ and $\xi > 0$, the induced inflation model in the Einstein frame is also represented by the special attractor action (15). In this model, $V_J = (\lambda/\xi^2)(\Omega - 1)^2$, and the Einstein frame potential for $\alpha = 1$ is given by Eq. (19) with $c^2 = \lambda/\xi^2$. This choice of c^2 here is not required; it was motivated by the desire to implement the Higgs inflation scenario [8]. But the potential (19) is different from the Higgs inflation potential anyway: it is not symmetric with respect to the change $\varphi \to -\varphi$, and it does not contain an important part of the potential at intermediate values of φ where the potential is quartic in φ. However, it is important that it belongs to the class of the special attractors. Moreover, it allows for the same generalization (20) of the scalar potential.

Universal attractors.—Finally, we wish to emphasize how the universal attractor models of Ref. [7] are related to α attractors and spell out how they fit in the present framework. The universal attractor models considered in Ref. [7] are defined by the choice $K_J = 1$ and $\Omega = 1 + \xi f(\phi)$ for an arbitrary function $f(\phi)$.

In the limit when $\xi \to \infty$, the inflationary predictions of these models coincide with those of the induced models with $\Omega \approx \xi f(\phi)$, as well as those of special attractors and α attractors for $\alpha \approx 1$. In this limit, there is no need to make a choice $f(\phi) = \phi^2$ (as we did in the case of an exact relation between the α models and generalized induced inflation models above). In the limit $\xi \to \infty$, the second term in Eq. (13) can be neglected and we find

$$K_E = \frac{3}{2}\frac{1}{\rho^2}, \qquad V_E = \frac{\lambda}{\xi^2}(1 - \rho)^2, \tag{21}$$

where we have replaced the nonminimal coupling $\Omega(\phi)$ (which can be chosen arbitrarily) by its inverse ρ. Here we see again that the pole structure at $\rho = 0$ allows us to deform the potential and, instead of the function (5), consider any function with additional terms with higher powers of $\rho = e^{-\sqrt{(2/3)}\varphi}$.

Moreover, one can calculate the subleading corrections to the above kinetic term that arise for finite values of ξ. For instance, in the case of Higgs inflation with $f = \phi^2$, the full kinetic term for the field ρ is given by

$$K_E = \frac{3}{2}\frac{1}{\rho^2} + \frac{1}{4\xi}\frac{1}{(1 - \rho)\rho^2} = \frac{3\alpha}{2}\frac{1}{\rho^2} + \frac{1}{4\xi}\frac{1}{\rho} \cdots + . \tag{22}$$

While this has the same leading pole, subleading corrections will be different. A particularly acute difference with respect to the case of induced inflation discussed in the previous subsection is that the kinetic term is not necessarily positive definite. In particular, inflation takes place close to $\rho = 0$, while the Minkowski vacuum is located at $\rho = 1$. In the latter regime, the second term will always be dominant. Therefore, Higgs inflation does not allow one to take ξ negative even in the Einstein frame, in contrast to induced inflation; in addition to the condition $\alpha > 0$ from the inflationary regime, one also requires $\xi > 0$ from the cosmological era following inflation.

Discussion.—Provided the kinetic term of the inflaton is given by a Laurent series, its inflationary predictions are, to a large extent, determined by the properties of the leading pole and, therefore, robust to changes to the subleading terms, either in the kinetic or the potential energy. Such a pole of order 2 underlies the attractor properties of both α and ξ attractors and, therefore, yields the inflationary predictions (2).

Next, we have explicitly demonstrated the unity of these two types of attractors, either based on nontrivial kinetic terms or on nonminimal couplings; when transforming ξ attractors from the Jordan to the Einstein frame, one obtains α attractors and vice versa. Moreover, we have emphasized that there are special type of attractors whose kinetic term consists only of a single pole: both the original α attractors of [6] as well as induced inflation [8] are of this form. This is illustrated in Fig. 1.

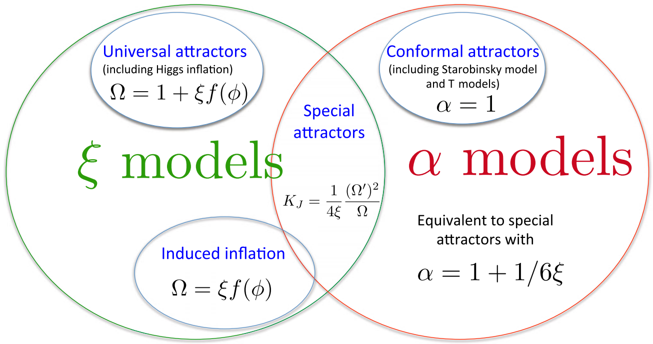

FIG. 1 (color online). Unification of cosmological attractors. The new class of special attractors is defined by Eq. (14) and is fully equivalent to α attractors with $\alpha = 1 + (1/6\xi)$.

The introduction of generalized ξ attractors including the special attractors (14) opens a simple way towards the unification of all presently known cosmological attractors. We have shown that the class of the special attractors is equivalent to α attractors with $\alpha = 1 + (1/6\xi) > 0$. This relation between both parameters, which is one of our main results, embodies the two viable ranges $\xi > 0$ and $\xi < -1/6$. In the Jordan frame, only the first of these has a positive kinetic term corresponding to $\alpha \geq 1$. However, similar to the Breitenlohner-Freedman bound, the theory is well defined for both cases: it has positive kinetic term in the Einstein frame, and it does not exhibit any instability. There is no theoretical lower bound on $r = 12\alpha/N^2$ in this class of models.

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