Minimal supergravity inflation

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Inflatino-less Cosmology

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We construct inflationary models in the context of supergravity with orthogonal nilpotent superfields [1]. When local supersymmetry is gauge-fixed in the unitary gauge, these models describe theories with only a single real scalar (the inflaton), a graviton and a gravitino. Critically, there is no inflatino, no sgoldstino, and no sinflaton in these models. This dramatically simplifies cosmological models which can simultaneously describe inflation, dark energy and SUSY breaking.

Introduction. In this letter we show that the recently constructed models of supergravity with orthogonal nilpotent superfields [1] significantly simplify construction of cosmological models which simultaneously describe inflation, dark energy and SUSY breaking. Now one can achieve this goal using the absolutely minimal number of ingredients: graviton, gravitino, and a single real scalar, the inflaton. Our methods apply to a broad class of inflationary theories, but they are especially suitable for describing α-attractors [2], which provide a very good fit to the latest cosmological data [3].

Orthogonal nilpotent superfields. Global supersymmetry models with orthogonal nilpotent superfields were studied in [4, 5]. Their generalization to local supergravity interacting with orthogonal nilpotent multiplets was presented in [1]. The models depend on two constrained chiral superfields, $\mathbf{S}$ and $\mathbf{\Phi}$ [6]. A stabilizer $\mathbf{S}$ has a nilpotency of degree two, $\mathbf{S}^n = 0$ for $n \geq 2$. This constraint removes the complex scalar $S(x)$, sgoldstino, from the bosonic spectrum. The chiral inflaton multiplet $\mathbf{\Phi}$ has as a first component a complex scalar $\Phi(x) = \phi(x) + ib(x)$. One can form a real superfield $\mathbf{B} \equiv \frac{1}{2i}\Big(\mathbf{\Phi} - \bar{\mathbf{\Phi}}\Big)$ with the first component $b(x)$, the sinflaton, and impose the orthogonality constraint $\mathbf{SB} = 0$. As a result, fields in the inflaton multiplet are no longer independent: the sinflaton $b(x)$, inflatino χ^ϕ and the auxiliary F^ϕ become functions of the spin 1/2 field χ^s in the S-multiplet. All of these fields vanish in the unitary gauge $\chi^s = 0$, which fixes the local supersymmetry of the action.

It follows from $\mathbf{S}^2 = 0$ and $\mathbf{SB} = 0$ that $\mathbf{B}$ is nilpotent of degree 3, $\mathbf{B}^m = 0$ for $m \geq 3$, [4, 5, 7]. The second real superfield which can be formed from the chiral inflaton superfield is $\mathbf{A} = \frac{1}{2}\Big(\mathbf{\Phi} + \bar{\mathbf{\Phi}}\Big)$. It starts with one real inflaton scalar field, $\phi(x)$.

The unusual property of these models is that in the unitary gauge, fixing local supersymmetry, there is only one real scalar $\phi(x)$, a massless graviton and a massive gravitino. There is no sgoldstino, no sinflaton and no inflatino. An essential property of these models is that the form of the potential is different from the standard supergravity potentials, as shown in [1]:

$$V = e^K(|D_S W|^2 - 3W^2)\Big|_{S=\bar{S}=\Phi-\bar{\Phi}=0} . \qquad (1)$$

First, the 3 scalars S, $\bar{S}$ and $\Phi - \bar{\Phi}$ vanish, there is no need to stabilize them. Secondly, the terms which would normally be present in the potential, quadratic and linear in $D_\Phi W$, are absent, despite the fact that $D_\Phi W$ can be arbitrary. This is because the auxiliary field F^ϕ from the inflaton multiplet is fermionic as a consequence of the orthogonality constraint $\mathbf{SB} = 0$.

Here we will study supergravity models with constrained superfields

$$K(\mathbf{S}\,,\bar{\mathbf{S}};\mathbf{\Phi},\bar{\mathbf{\Phi}}), \;\; W = f(\mathbf{\Phi})\mathbf{S} + g(\mathbf{\Phi}), \;\; \mathbf{S}^2 = \mathbf{SB} = 0. \qquad (2)$$

The consistency of these models with the constraints was studied in [1] where it was shown that the superfield Kähler potential can be also brought to the form

$$K(\mathbf{S}\,,\bar{\mathbf{S}};\mathbf{\Phi},\bar{\mathbf{\Phi}}) = \mathbf{S}\,\bar{\mathbf{S}} + h(\mathbf{A})\mathbf{B}^2 . \qquad (3)$$

This means, in particular that the Kähler potential in supergravity vanishes when the bosonic constraints are imposed:

$$K(S\bar{S};\Phi,\bar{\Phi})\Big|_{S=\bar{S}=\Phi-\bar{\Phi}=0} = 0 . \qquad (4)$$

Note, however, that the bosonic constraints $S = \bar{S} = \Phi - \bar{\Phi} = 0$, when deriving the supergravity action, have to be applied only after the relevant derivatives over S and $\bar{S}$ and Φ and $\bar{\Phi}$ are taken.

For models with orthogonal nilpotent superfields, the inflaton action is very simple. Taking into account (4) and with a proper normalization of S, we find

$$e^{-1}\mathcal{L}_{\rm infl} = -K_{\Phi,\bar{\Phi}}\partial\Phi\partial\bar{\Phi} + [3g^2(\Phi) - f^2(\Phi)]\Big|_{\Phi=\bar{\Phi}} . \qquad (5)$$

Kähler potentials in models with $\mathbf{S}^2 = \mathbf{SB} = 0$. Many successful inflationary models in supergravity are based on theories where the Kähler potential either vanishes along the inflaton direction, or can be represented in such form after some Kähler transformations, see for example [8–14, 16]. In models with $\mathbf{S}^2 = \mathbf{SB} = 0$, where $B = (\mathbf{\Phi} - \bar{\mathbf{\Phi}})/(2i)$, this requirement is naturally satisfied (3), (4).

Here we study the cosmological models with orthogonal nilpotent superfields (2) over several different Kähler potentials. The simplest Kähler potential with a flat direction describing a canonically normalized inflaton field $\phi = \mathrm{Re}\,\Phi$ is given by [9, 10]

$$K = -\frac{1}{4}(\Phi - \bar{\Phi})^2 + S\bar{S} \ . \tag{6}$$

Here the geometry of the moduli space is flat.

We will be especially interested in the Kähler potentials for a broad class of cosmological attractors describing Escher-type hyperbolic geometry [11, 12] of the inflaton moduli space. Compatibility of the constraints $\mathbf{S}^2 = \mathbf{SB} = 0$ with the hyperbolic geometry is demonstrated in the Appendix. Examples of such Kähler potentials include

$$K = -\frac{3}{2}\alpha \log\left[\frac{(1-\Phi\bar{\Phi})^2}{(1-\Phi^2)(1-\bar{\Phi}^2)}\right] + S\bar{S} \ . \tag{7}$$

It describes hyperbolic geometry in disk variables. The same geometry can be described in half-plane variables by the Kähler potential

$$K = -\frac{3}{2}\,\alpha \log\left[\frac{(\Phi+\bar{\Phi})^2}{4\Phi\bar{\Phi}}\right] + S\bar{S} \, . \tag{8}$$

These two versions correspond to equivalent ways of describing the Kähler geometry of α-attractors. See refs. [12–14] and the Appendix of our paper for a discussion of this issue.

One may also consider these Kähler potentials with the term $S\bar{S}$ under the logarithm. In all of these cases, the Kähler potential vanishes for $\Phi = \bar{\Phi}$, and $K_{S,\bar{S}} = 1$ or can be brought to $K_{S,\bar{S}} = 1$ by a holomorphic transformation defined in [1]. The inflaton action is given by (5), and the inflaton potential is given by a simple expression

$$V = f^2(\phi) - 3g^2(\phi) \ . \tag{9}$$

This result is similar to the expression $V = f^2(\phi)$ for the family of models with $W = Sf(\Phi)$ developed in [10]. The new generation of models is different in two respects. First of all, it describes a non-vanishing gravitino mass

$$m_{3/2}(\phi) = g(\phi) \ . \tag{10}$$

Additionally, it may also describe non-vanishing vacuum energy (cosmological constant) at the minimum of the potential. Without any loss of generality one may assume that the minimum of the potential corresponding to our vacuum state is at $\phi = 0$. The cosmological constant is equal to

$$\Lambda = f^2(0) - 3g^2(0) \ . \tag{11}$$

The condition that $\phi = 0$ is a minimum implies that $f'(0) = \sqrt{3}g'(0)$, up to small corrections vanishing in the limit $\Lambda \to 0$.

These conditions, plus the requirement that the functions $f(\phi)$ and $g(\phi)$ are holomorphic, leave lots of freedom to describe observational data. Indeed there are many ways to do so, depending on the choice of the Kähler potential.

Even though the expression of the potential $V = f^2(\phi) - 3g^2(\phi)$ is valid for all choices of the Kähler potentials described above, the field ϕ in the theories with the Kähler potentials (7) and (8) is not canonically normalized. In the theory (7) the field ϕ is related to the canonically normalized inflaton field φ as follows:

$$\phi = \tanh\frac{\varphi}{\sqrt{6\alpha}} \ . \tag{12}$$

Meanwhile for the theory (8) one has

$$\phi = e^{-\sqrt{\frac{2}{3\alpha}}\varphi} \ . \tag{13}$$

Thus, the potential $V = f^2(\phi) - 3g^2(\phi)$, expressed in terms of a canonically normalized field φ, depends on the choice of the Kähler potential. In the next section we will describe several realistic inflationary models in this context.

Inflationary models.

Model 1: $f(\phi) = M\phi^2 + a, \quad g(\phi) = b.$

The potential in this model is

$$V = M^2\phi^4 + 2aM\phi^2 + a^2 - 3b^2 \ . \tag{14}$$

The cosmological constant in this model, and all other models we present here, is equal to

$$\Lambda = a^2 - 3b^2. \tag{15}$$

In realistic models we should have $\Lambda \sim 10^{-120}$ due to an almost precise cancellation between a^2 and $3b^2$ in accordance with a string landscape scenario. The gravitino mass is $m_{3/2} = b$, which nearly coincides with $a/\sqrt{3}$. For $b \sim 10^{-15}$ one can have the gravitino mass in the often discussed TeV range. To have a proper amplitude of scalar perturbations one should have $M \sim 10^{-5} \gg a,\ b$.

If we consider a model with the simplest canonical Kähler potential (6), the potential (14) is quartic with respect to the canonically normalized inflaton field, which rules out this simple model.

The situation instantly improves in the theory with the logarithmic Kähler potential (7), which yields the following potential in terms of the canonically normalized field φ:

$$V = M^2\tanh^4\frac{\varphi}{\sqrt{6\alpha}} + 2aM\tanh^2\frac{\varphi}{\sqrt{6\alpha}} + a^2 - 3b^2 \ . \tag{16}$$

This is the typical T-model α attractor potential [2]. Inflation occurs at the plateau where $\tanh\frac{\varphi}{\sqrt{6\alpha}} \approx 1$. In this regime the second term in (16) is much smaller than the first term, and both terms are much greater than Λ, so inflation is described by the quartic T-model potential

$$V = M^2 \tanh^4 \frac{\varphi}{\sqrt{6\alpha}} \ . \tag{17}$$

The observational predictions of this model for $\alpha \lesssim 10$ practically coincide with the predictions of the simpler model $V = M^2 \tanh^2 \frac{\varphi}{\sqrt{6\alpha}}$, for the same number of e-foldings N [2]. However, at the end of inflation in the model (17) the inflaton field begins to oscillate in the approximately quartic potential $\sim \varphi^4$. The average equation of state during this stage is the same as of the hot plasma, $p = \rho/3$, as if reheating finishes immediately after inflation. This increases the required number of e-foldings by $\Delta N \sim 3$ [15]. In its turn, this leads to a slight increase of the spectral index n_s, which may provide even better fit to the recent Planck data.

Model 2: $f(\phi) = M\phi^2 + a, \quad g(\phi) = m\phi^2 + b$
The potential is

$$V = (M^2 - 3m^2)\phi^4 + 2(Ma - 3m\, b)\phi^2 + a^2 - 3b^2 \ . \tag{18}$$

This model is very similar to the previous one, but there is one potentially interesting difference: The gravitino mass depends on the inflaton, $m_{3/2} = m\phi^2 + b$.

Model 3: $f(\phi) = \sqrt{\frac{M^2}{2}\phi^2 + a^2}, \qquad g(\phi) = b$
The potential is

$$V = \frac{M^2}{2}\phi^2 + a^2 - 3b^2 \ . \tag{19}$$

The potential for ϕ is exactly quadratic, plus a cosmological constant.

In the theory with the logarithmic Kähler potential (7) this potential becomes a potential for the simplest α-attractor model of the canonically normalized field φ:

$$V = \frac{M^2}{2} \tanh^2 \frac{\varphi}{\sqrt{6\alpha}} + a^2 - 3b^2 \ . \tag{20}$$

The gravitino mass is $m_{3/2} = b$.

Model 4: $f(\phi) = \sqrt{\frac{M^2}{2}\phi^2 + a^2}, \quad g(\phi) = \sqrt{\frac{m^2}{2}\phi^2 + b^2}$
In this model one has

$$V = \frac{M^2 - 3m^2}{2}\phi^2 + a^2 - 3b^2 \ . \tag{21}$$

In the theory with the logarithmic Kähler potential (7) the potential of a canonically normalized inflaton field becomes

$$V = \frac{M^2 - 3m^2}{2} \tanh^2 \frac{\varphi}{\sqrt{6\alpha}} + a^2 - 3b^2 \ . \tag{22}$$

This model is very similar to Model 3, but the gravitino mass is ϕ-dependent, $m_{3/2} = \sqrt{\frac{m^2}{2}\phi^2 + b^2}$.

Model 5: $f(\phi) = \sqrt{F^2(\phi) + a^2}, \quad g(\phi) = \sqrt{G^2(\phi) + b^2}$
In this model

$$V = F^2(\phi) - 3G^2(\phi) + a^2 - 3b^2, \quad m_{3/2} = \sqrt{G^2(\phi) + b^2}. \tag{23}$$

Because of the freedom of choice of the holomorphic functions F and G, one can have a wide variety of potentials fitting all observational data even if the fields ϕ is canonically normalized, with the Kähler potential (6), see e.g. [16]. Meanwhile in the theories with the Kähler potential (7) one finds a family of T-model α-attractors with

$$V = F^2(\tanh \frac{\varphi}{\sqrt{6\alpha}}) - 3G^2(\tanh \frac{\varphi}{\sqrt{6\alpha}}) + a^2 - 3b^2 \ . \tag{24}$$

For a wide range of functions F and G, these theories have universal cosmological predictions for $\alpha \lesssim 10$ and any given number of e-foldings: $n_s = 1 - 2/N$, $r = 12\alpha/N^2$ [2]. However, by a proper choice of the function F one can modify the required number of e-foldings N, which can be useful for tuning the predictions for n_s.

Model 6: $f(\phi) = \sqrt{(1-\phi)^2 + a^2}, \qquad g(\phi) = b$
It is a particular version of Model 5 for $F(\phi) = M(1-\phi)$ and $G(\phi) = 0$. This yields

$$V = M^2(1-\phi)^2 + \Lambda, \quad \Lambda = a^2 - 3b^2, \quad m_{3/2} = b. \tag{25}$$

Using the half-plane Kähler potential (8) and the relation $\phi = e^{-\sqrt{\frac{2}{3\alpha}}\varphi}$ (13) one finds

$$V = M^2 \left(1 - e^{-\sqrt{\frac{2}{3\alpha}}\varphi}\right)^2 + \Lambda \ . \tag{26}$$

This represents the family of E-model α-attractors [2, 11], which reduces to the Starobinsky model for $\alpha = 1$, $\Lambda = 0$ and $m_{3/2} = 0$. Meanwhile our class of theories describes E-models with arbitrary α, Λ and $m_{3/2}$.

Conclusions. As one could see from the previous section, it is very easy to formulate and analyze models with orthogonal nilpotent fields. Previously, it was a much more complicated task because of certain constraints imposed on inflationary models with nilpotent fields, see e.g. [17–19] and the more advanced models presented in [14]. These constraints where required for simplification of the fermionic sector, but they are no longer required in the new class of models where the fermionic sector is trivial because the inflatino disappears in the unitary gauge. As a result, we have lots of flexibility in finding economical models containing only inflaton, graviton and gravitino, and yet capable of simultaneously describing inflation, dark energy and SUSY breaking. Note that SUSY breaking is achieved without introducing light Polonyi-type moduli, which plagued supergravity cosmology for more than three decades.

The absence of the inflatino also helps us argue that there there is no problem with the unitarity bound during inflation in these models. The effective cutoff in supergravity is the scale at which scattering amplitudes violate unitarity bound. In the theories with nilpotent fields in Minkowski space, this cutoff is expected at $\Lambda \simeq \sqrt{m_{3/2}M_p}$ [18]. It was argued in [5] that the cutoff during inflation is expected at $\Lambda \simeq \sqrt{HM_p}$. More precisely, it was shown in [20] that helicity 1/2 gravitino projectors in an expanding universe instead of $m_{3/2}$ involve the effective gravitino mass $\sqrt{H^2+m_{3/2}^2}$. Therefore the UV cutoff is expected at $\Lambda \simeq (H^2+m_{3/2}^2)^{1/4}M_p^{1/2} > \sqrt{HM_p}$. During inflation with $H \ll M_p$, this cutoff is much higher than the typical energy of inflationary quantum fluctuations $\sim H$.

In general, there could be some additional contributions to scattering due to gravitino-inflatino mixing, but in our models there is no inflatino. Therefore no violation of the unitarity bound is expected during inflation at sub-Planckian energy density.

We hope to return to this issue in the future, simultaneously with investigation of reheating in the new class of models. In particular, we expect that the absence of the inflatino should significantly simplify the theory of non-thermal gravitino production by an oscillating inflaton field [20–23].

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Appendix A. Hyperbolic geometry models with orthogonal nilpotent superfields

Three equivalent versions of α-attractor models with $S^2(x,\theta) = 0$ and hyperbolic geometry of the inflaton moduli space [11, 12] are given either by a disk geometry $Z\bar{Z} < 1$,

$$\begin{aligned} K &= -\frac{3}{2}\alpha \log\left[\frac{(1-Z\bar{Z})^2}{(1-Z^2)(1-\overline{Z}^2)}\right] + S\bar{S}\,, \\ W &= A(Z) + SB(Z)\,, \end{aligned} \tag{27}$$

or a half-plane geometry, $T+\bar{T} > 0$,

$$\begin{aligned} K &= -\frac{3}{2}\,\alpha \log\left[\frac{(T+\bar{T})^2}{4T\bar{T}}\right] + S\bar{S}\,, \\ W &= G(T) + SF(T)\,. \end{aligned} \tag{28}$$

or a Killing adapted geometry

$$\begin{aligned} K &= -3\alpha \log\left[\cosh\frac{\Phi-\bar{\Phi}}{\sqrt{6\alpha}}\right] + S\bar{S}\,, \\ W &= A(\tanh\frac{\Phi}{\sqrt{6\alpha}}) + SB(\tanh\frac{\Phi}{\sqrt{6\alpha}}) \\ &= G(e^{\sqrt{\frac{2}{3\alpha}}\Phi}) + SF(e^{\sqrt{\frac{2}{3\alpha}}\Phi})\,. \end{aligned} \tag{29}$$

In all of these models the Kähler potential and the superpotential, separately, are related by a change of variables

$$T = \frac{1+Z}{1-Z}, \quad Z = \tanh\frac{\Phi}{\sqrt{6\alpha}}, \quad T = e^{\sqrt{\frac{2}{3\alpha}}\Phi}. \tag{30}$$

Now we would like to impose the orthogonality constraint on our superfields. The Killing-adapted variable Φ in (29) is an unconstrained superfield whose scalar part is not restricted by the boundaries. Therefore we may start by imposing the orthogonality and the nilpotency constraint in the form

$$\mathbf{S}^2 = 0, \quad \mathbf{S}(\mathbf{\Phi}-\bar{\mathbf{\Phi}}) = 0, \quad (\mathbf{\Phi}-\bar{\mathbf{\Phi}})^n = 0,\ n \geq 3\,, \tag{31}$$

Using the relation between these variables, one can derive the related constraints for the disk and half-plane variables Z, and T, respectively.

$$\mathbf{S}(\mathbf{\Phi}-\bar{\mathbf{\Phi}}) = 0 \quad \Rightarrow \quad \mathbf{S}(\mathbf{Z}-\bar{\mathbf{Z}}) = 0 \quad \Rightarrow \quad \mathbf{S}(\mathbf{T}-\bar{\mathbf{T}}) = 0.$$

It can be shown also that the Kähler potentials in (27), (28), (29), take the form of eq. (3) due to orthogonality constraints.

For the Kähler potential (29), one finds that the corresponding $h(\mathbf{A}) = 1$ in eq. (3) and therefore the model with hyperbolic geometry in variables shown in (29) with constraints $\mathbf{B}^3 = 0$ is reduced to the simplest form of the shift-symmetric, inflaton-independent Kähler potential given in eq. (6). In disk and half-plane variables, the corresponding $h(\mathbf{A})$ are not trivial and lead to non-canonical kinetic terms for the inflaton.

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[6] We use bold face letters for superfields which in global case are functions of $(x, \theta, \bar{\theta})$. For example, the chiral superfield $\mathbf{S}$ means that $\mathbf{S} = S(x) + \sqrt{2}\theta\chi^s + \theta^2 F^s$, and $\mathbf{\Phi} = \Phi(x) + \sqrt{2}\theta\chi^\phi + \theta^2 F^\phi$.

[7] An easy way to see this is to notice that from $\mathbf{S}^2 = 0$ and $\mathbf{SB} = 0$ it follows that $b \sim \chi^s_\alpha \chi^s_{\dot\alpha} \partial^{\alpha\dot\alpha}\phi$, up to higher order in fermions, and since χ^s_α is a 2-component spinor, $b^m = 0$ for $m \geq 3$.

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