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The Volkov–Akulov–Starobinsky supergravity

I. Antoniadis [a,1], E. Dudas [b], S. Ferrara [a,c,d], A. Sagnotti [e,*]

[a] Department of Physics, CERN Theory Division, CH-1211 Geneva 23, Switzerland
[b] Centre de Physique Théorique, École Polytéchnique, F-91128 Palaiseau, France
[c] INFN – Laboratori Nazionali di Frascati, Via Enrico Fermi 40, I-00044 Frascati, Italy
[d] Department of Physics and Astronomy, University of California, Los Angeles, CA 90095-1547, USA
[e] Scuola Normale Superiore and INFN, Piazza dei Cavalieri 7, I-56126 Pisa, Italy

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ABSTRACT

We construct a supergravity model whose scalar degrees of freedom arise from a chiral superfield and are solely a scalaron and an axion that is very heavy during the inflationary phase. The model includes a second chiral superfield X, which is subject however to the constraint $X^2 = 0$ so that it describes only a Volkov–Akulov goldstino and an auxiliary field. We also construct the dual higher-derivative model, which rests on a chiral scalar curvature superfield $\mathcal{R}$ subject to the constraint $\mathcal{R}^2 = 0$, where the goldstino dual arises from the gauge-invariant gravitino field strength as $\gamma^{mn}\mathcal{D}_m\psi_n$. The final bosonic action is an $R + R^2$ theory involving an axial vector A_m that only propagates a physical pseudoscalar mode.

1. Starobinsky models of inflation in supergravity

It was recently shown how to embed in supergravity [1] a class of models [2] including the Starobinsky potential [3], which affords an excellent agreement with recent PLANCK data [4] for an inflationary epoch of about 60 e-folds. However, the models based on the "old minimal" supergravity rest on a pair of chiral superfields (see also [5–7] for closely related work), and thus involve three scalar fields in addition to the inflaton.

The construction reflects the Starobinsky duality [3,8] between an $R + R^2$ action and a special scalar-gravity system, encompassed by the master action

$$\mathcal{S} = \int d^4x\sqrt{-g}\left[\frac{1}{2}R + \chi(R-\varphi) + \alpha\varphi^2\right], \tag{1.1}$$

which reduces to

$$\mathcal{S}_1 = \int d^4x\sqrt{-g}\left[\frac{1}{2}(1+2\chi)R - \frac{1}{4\alpha}\chi^2\right] \tag{1.2}$$

upon integration over φ. On the other hand, χ enters Eq. (1.1) as a Lagrange multiplier imposing the constraint $\varphi = R$, and enforcing it leads to the dual form

$$\mathcal{S}_2 = \int d^4x\sqrt{-g}\left[\frac{1}{2}R + \alpha R^2\right]. \tag{1.3}$$

The embedding of this model in (higher-derivative) supergravity takes the form [9,2]

$$\mathcal{L} = \left[-S_0\bar{S}_0 + h\left(\frac{\mathcal{R}}{S_0}, \frac{\overline{\mathcal{R}}}{\bar{S}_0}\right)S_0\bar{S}_0\right]_D + \left[W\left(\frac{\mathcal{R}}{S_0}\right)S_0^3\right]_F, \tag{1.4}$$

where here and in the following subscripts identify D and F superspace densities, S_0 is the chiral compensator field and $\mathcal{R}$ is the chiral scalar curvature superfield, defined via the curved chiral projector Σ as

$$\mathcal{R} = \frac{\Sigma(\bar{S}_0)}{S_0}. \tag{1.5}$$

In Eq. (1.4) h is a real function of the chiral superfield $\mathcal{R}/S_0$ and its conjugate $\overline{\mathcal{R}}/\bar{S}_0$ that contains an $\mathcal{R}\overline{\mathcal{R}}$ term and W is a chiral function. The presence of the $\mathcal{R}\overline{\mathcal{R}}$ term in h brings about an R^2 term in components.

In detail, the components of the chiral superfield $\mathcal{R}$ are [1,10,2]

$$\mathcal{R} = \left(\bar{u} \equiv S + iP, \gamma^{mn}\mathcal{D}_m\psi_n, -\frac{1}{2}R - \frac{1}{3}A_m^2 + i\mathcal{D}^m A_m - \frac{1}{3}u\bar{u}\right), \tag{1.6}$$

where u and A_m are the "old minimal" auxiliary fields of $N = 1$ supergravity and ψ_n is the gravitino field. The action (1.4) can be recast in a two-derivative dual form proceeding along the lines of

* Corresponding author.
E-mail addresses: Ignatios.Antoniadis@cern.ch (I. Antoniadis), dudas@cpht.polytechnique.fr (E. Dudas), Sergio.Ferrara@cern.ch (S. Ferrara), sagnotti@sns.it (A. Sagnotti).
[1] On leave of absence from CPhT, École Polytéchnique, F-91128 Palaiseau Cedex, France.

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Eqs. (1.1) and (1.2) and making use of a pair of chiral multiplets Λ and C as in [9,2], so that

$$\mathcal{L} = \left[-S_0\bar{S}_0 + h(C,\bar{C})S_0\bar{S}_0\right]_D + \left[\left\{\Lambda\left(C - \frac{\mathcal{R}}{S_0}\right) + W(C)\right\}S_0^3\right]_F. \tag{1.7}$$

Eliminating Λ in (1.7) yields the constraint $\mathcal{R} = S_0 C$, and one recovers in this fashion the original Lagrangian (1.4). Notice, to this end, that letting

$$W(C) = Cg(C) + W_0, \tag{1.8}$$

on account of the identity [1,2]

$$\left[f(\Lambda)\mathcal{R}S_0^2\right]_F + \text{h.c.} = \left[\left(f(\Lambda) + \overline{f}(\overline{\Lambda})\right)S_0\bar{S}_0\right]_D + \text{tot. deriv.}, \tag{1.9}$$

which holds for any chiral superfield Λ and for any function f, $g(C)$ can be shifted away redefining Λ into another chiral superfield Λ' [2], so that the Lagrangian can be cast in the form

$$\mathcal{L}_{\text{dual}} = \left[\left[-1 - \Lambda' - \overline{\Lambda}' + h(C,\bar{C})\right]S_0\bar{S}_0\right]_D + \left[\left(W_0 + C\Lambda'\right) + \text{h.c.}\right]_F. \tag{1.10}$$

Letting

$$\Lambda' = T - \frac{1}{2}, \tag{1.11}$$

one is then finally led to a standard $N = 1$ supergravity with a Kähler potential K and a superpotential W given by

$$K = -3\ln\left[T + \overline{T} - h(C,\bar{C})\right], \qquad W = C\left(T - \frac{1}{2}\right) + W_0. \tag{1.12}$$

The case in which $h(C,\bar{C})$ is a pure Kähler transformation of the form $h(C) + \bar{h}(\bar{C})$, so that C is not dynamical, was considered in [11], where it was referred to as $f(R)$ supergravity, but this class of models does not reproduce the $R + R^2$ bosonic terms and the Starobinsky potential [2,12,13].

The Lagrangian (1.12) contains Starobinsky's inflaton ϕ, which is related to T according to $\text{Re}(T) = \exp(\sqrt{2/3}\phi)$, and setting to zero the other three fields one can recover exactly, for $W_0 = 0$, the scalar potential of the original Starobinsky model. However, it was shown in [5] that for a minimal choice $h(C,\bar{C}) = C\bar{C}$ the complex scalar direction C is *unstable* during the inflationary phase. Non-minimal (and therefore non-universal) Kähler functions are thus needed to arrive at a satisfactory model, and in particular Kallosh and Linde [5] showed that

$$h(C,\bar{C}) = C\bar{C} - \zeta(C\bar{C})^2 \tag{1.13}$$

can stabilize the C direction for sufficiently large positive values of ζ.

Although for $W_0 = 0$ the model admits a supersymmetric ground state, supersymmetry is broken during the inflationary phase, and C plays the role of a goldstino superfield driving the breaking of supersymmetry [2,5,7]. In what follows we shall explain how a minimal and universal embedding of the Starobinsky potential in supergravity emerges once the ordinary chiral superfield C is replaced with a chiral superfield X satisfying, as in [14–17], the constraint $X^2 = 0$.

2. The Volkov–Akulov Lagrangian

It has been known for some time that the Volkov–Akulov Lagrangian [19] can be recast in a manifestly supersymmetric form introducing a chiral superfield X that satisfies identically the constraint [14–17]

$$X^2 = 0. \tag{2.1}$$

This eliminates the scalar component of X in favor of a goldstino bilinear, so that in two-component notation

$$X = \frac{GG}{2F_X} + \sqrt{2}\theta G + \theta^2 F_X, \tag{2.2}$$

and the complete Volkov–Akulov Lagrangian is then

$$\mathcal{L}_{VA} = [X\overline{X}]_D + [fX + \text{h.c.}]_F, \tag{2.3}$$

where the subscripts denote again D and F superspace densities. Notice that supersymmetry can be realized *off-shell*, as emphasized in [17,18], insofar as F_X is not replaced by the solution of its algebraic equation of motion. In supergravity, the off-shell couplings of the goldstino to the gravity multiplet can be found in a similar way, but taking into account the constraint (2.1) the most general couplings of X to supergravity rest on a Kähler potential K and a superpotential W of the form

$$K = -3\log(1 - X\overline{X}) \equiv 3X\overline{X}, \qquad W = fX + W_0, \tag{2.4}$$

since terms linear in X or $\overline{X}$ can be reabsorbed in W, and as a result [20]

$$V = \frac{1}{3}|f|^2 - 3|W_0|^2, \qquad m_{3/2}^2 = |W_0|^2. \tag{2.5}$$

The supergravity Lagrangian resulting from Eq. (2.4) does encode the proper goldstino couplings, and in particular in two-component notation the fermionic mass terms read

$$\mathcal{L}_{mass} = -m_{3/2}\left(\psi_m + \frac{i}{\sqrt{6}}\sigma_m\bar{G}\right)\sigma^{mn}\left(\psi_n + \frac{i}{\sqrt{6}}\sigma_n\bar{G}\right) + \text{h.c.} \tag{2.6}$$

3. The minimal Starobinsky Lagrangian

The usual embedding of the Starobinsky Lagrangian in supergravity rests, in the "old minimal" two-derivative formulation, on the gravitational supermultiplet coupled to a pair of additional chiral multiplets. As we have anticipated, the corresponding action is not unique, and non-minimal terms are actually needed [5] to stabilize the scalar fields during the inflationary phase.

A minimal universal model obtains if supergravity is coupled to the *constrained* goldstino multiplet X described in the preceding section and to a chiral multiplet T containing the inflaton.[2] Taking into account the constraint of Eq. (2.1), in this *off-shell* formulation the Lagrangian is determined by

$$K = -3\ln[T + \overline{T} - X\overline{X}], \qquad W = MXT + fX + W_0, \tag{3.1}$$

where $M^2 = \frac{3}{4\alpha}$ from the comparison with Eq. (1.1), while the corresponding scalar potential [20] is simply

$$V = \frac{|MT + f|^2}{3(T + \overline{T})^2}, \tag{3.2}$$

since the scalar component of X is not a dynamical field but a goldstino bilinear. Let us stress, however, that X contributes to the scalar potential via its derivatives, since

$$F_{\overline{X}} = e^{\frac{K}{2}}(K_{X\overline{X}})^{-1}\overline{W}_{\overline{X}}, \tag{3.3}$$

includes a bosonic contribution although it contains no elementary scalar field.

[2] The constrained superfield X was previously considered, in a different context also related to inflationary models, in [21].

The form of (3.2) reflects the no-scale structure [22] of the T kinetic term and its coupling to the goldstino superfield X, and in particular the constant superpotential W_0 does not enter V while it determines the gravitino mass. The complete bosonic Lagrangian is

$$\mathcal{L} = \frac{R}{2} - \frac{3}{(T+\overline{T})^2}|\partial T|^2 - \frac{|MT+f|^2}{3(T+\overline{T})^2}, \tag{3.4}$$

and letting

$$T = e^{\phi\sqrt{\frac{2}{3}}} + ia\sqrt{\frac{2}{3}}, \tag{3.5}$$

it finally becomes

$$\mathcal{L} = \frac{R}{2} - \frac{1}{2}(\partial\phi)^2 - \frac{1}{2}e^{-2\phi\sqrt{\frac{2}{3}}}(\partial a)^2 - \frac{1}{12}\big(M + fe^{-\sqrt{\frac{2}{3}}\phi}\big)^2 - \frac{M^2}{18}e^{-2\phi\sqrt{\frac{2}{3}}}a^2. \tag{3.6}$$

If $Mf < 0$, after a shift of ϕ and a rescaling of the axion a, one can bring the Lagrangian to the form

$$\mathcal{L} = \frac{R}{2} - \frac{1}{2}(\partial\phi)^2 - \frac{1}{2}e^{-2\phi\sqrt{\frac{2}{3}}}(\partial a)^2 - \frac{M^2}{12}\big(1 - e^{-\sqrt{\frac{2}{3}}\phi}\big)^2 - \frac{M^2}{18}e^{-2\phi\sqrt{\frac{2}{3}}}a^2. \tag{3.7}$$

This is precisely a minimal Starobinsky Lagrangian where ϕ is accompanied by an axion field a that is however much heavier during the inflationary phase where the vacuum values ϕ_0 are large and positive, so that

$$m_\phi^2 \simeq \frac{M^2}{9}e^{-2\phi_0\sqrt{\frac{2}{3}}} \ll m_a^2 \equiv \frac{M^2}{9}. \tag{3.8}$$

4. Dual gravitational formulation

In the conformal compensator formalism [1], the Lagrangian of the previous section reads

$$\mathcal{L} = -\big[\big(T+\overline{T}-|X|^2\big)S_0\bar{S}_0\big]_D + \big[(MXT + fX + W_0)S_0^3 + \text{h.c.}\big]_F, \tag{4.1}$$

and can be recast in the form

$$\mathcal{L} = \big[|X|^2 S_0\bar{S}_0\big]_D + \left[\left(T\left(-\frac{\mathcal{R}}{S_0} + MX\right) + fX + W_0\right)S_0^3 + \text{h.c.}\right]_F \tag{4.2}$$

resorting to the identity

$$\big[(T+\overline{T})S_0\bar{S}_0\big]_D = \big[T\mathcal{R}S_0^2\big]_F + \text{h.c.}, \tag{4.3}$$

where $\mathcal{R}$ is the chiral supergravity multiplet. Notice that T enters Eq. (4.2) as a Lagrange multiplier, whose equation of motion is the constraint

$$X = \frac{1}{M}\frac{\mathcal{R}}{S_0}, \tag{4.4}$$

and as a result the constraint $X^2 = 0$ of Eq. (2.1) translates into a similar constraint on $\mathcal{R}$:

$$\mathcal{R}^2 = 0. \tag{4.5}$$

Conversely, one can start from the dual gravitational theory

$$\mathcal{L} = -\left[S_0\bar{S}_0 - \frac{\mathcal{R}\overline{\mathcal{R}}}{M^2}\right]_D + \left[W_0 + \xi\frac{\mathcal{R}}{S_0}S_0^3 + \sigma\mathcal{R}^2 S_0\right]_F, \tag{4.6}$$

and in this form the nonlinear constraint (4.5) originates from the field equation of σ. One can also use the identity

$$\big[\sigma\mathcal{R}^2 S_0 + \text{h.c.}\big]_F = \left[\left(\sigma\frac{\mathcal{R}}{S_0} + \overline{\sigma}\frac{\overline{\mathcal{R}}}{\bar{S}_0}\right)S_0\bar{S}_0\right]_D + \text{tot. deriv.}, \tag{4.7}$$

which is a particular case of (1.9), and introducing two Lagrangian chiral superfields multipliers T and C according to

$$\begin{aligned}\mathcal{L}_{\text{dual}} &= -\left[\left(1 - \sigma C - \overline{\sigma}\overline{C} - \frac{C\overline{C}}{M^2}\right)S_0\bar{S}_0\right]_D \\ &\quad + \left[\left(-T\left(\frac{\mathcal{R}}{S_0} - C\right) + W_0 + \xi C\right)S_0^3 + \text{h.c.}\right]_F \\ &= -\left[\left(1 + T + \overline{T} - \sigma C - \overline{\sigma}\overline{C} - \frac{C\overline{C}}{M^2}\right)S_0\bar{S}_0\right]_D \\ &\quad + \big[\widetilde{W}(\sigma, T, C)S_0^3 + \text{h.c.}\big]_F,\end{aligned} \tag{4.8}$$

where

$$\widetilde{W}(T, C) = W_0 + (T+\xi)C, \tag{4.9}$$

a final shift $T \to T + \sigma C - \frac{1}{2}$ and the replacement of $\xi - \frac{1}{2}$ with f turn this expression into

$$\widetilde{W}(\sigma, T, C) = TC + \sigma C^2 + W_0 + fC. \tag{4.10}$$

Notice that the σ field equation enforces the constraint $X^2 = 0$, and finally letting $C/M = X$ and rescaling f to f/M yields the Lagrangian

$$\mathcal{L} = -\big[(T + \overline{T} - X\overline{X})S_0\bar{S}_0\big]_D + \big[W(T, X)S_0^3 + \text{h.c.}\big]_F, \tag{4.11}$$

where

$$W(T, X) = W_0 + (MT + f)X. \tag{4.12}$$

This is precisely the dual action coupled to a goldstino multiplet (3.1), which completes our proof of the duality with the higher-derivative supergravity form of the Starobinsky model (4.6).

In order to write explicitly the bosonic gravitational action, one introduces the Jordan scalar

$$e^{\phi\sqrt{\frac{2}{3}}} = 1 + 2\chi, \tag{4.13}$$

in terms of which the Lagrangian (3.7) becomes

$$\mathcal{L} = \frac{R}{2} - \frac{3}{(1+2\chi)^2}\left[(\partial\chi)^2 + \frac{1}{6}(\partial a)^2\right] - \frac{M^2}{3}\frac{\chi^2 + \frac{a^2}{6}}{(1+2\chi)^2}. \tag{4.14}$$

The transition to the Jordan frame is effected by the Weyl rescaling

$$g \to (1+2\chi)g, \tag{4.15}$$

and the resulting Lagrangian is

$$\mathcal{L} = \frac{1}{2}(1+2\chi)R - \frac{1}{2}\frac{(\partial a)^2}{1+2\chi} - \frac{M^2}{3}\left(\chi^2 + \frac{a^2}{6}\right). \tag{4.16}$$

In the "old minimal" supergravity formulation the axion should be traded for the longitudinal mode $\mathcal{D}\cdot A$ of the pseudovector auxiliary field A_m, and this last step brings about an interesting general link that can be deduced from the master Lagrangian

$$\mathcal{L}_a = -\frac{M^2}{18}a^2 + A^m\partial_m a + \frac{1}{2}(1+2\chi)A_m^2. \tag{4.17}$$

Varying $\mathcal{L}_a$ with respect to A_m yields indeed

$$A_m = -\frac{\partial_m a}{1+2\chi}, \qquad \mathcal{L}_a = -\frac{M^2}{18}a^2 - \frac{1}{2}\frac{(\partial a)^2}{1+2\chi}, \tag{4.18}$$

while varying $\mathcal{L}_a$ with respect to a yields

$$a = -\frac{9}{M^2}\mathcal{D}\cdot A, \qquad \mathcal{L}_a = \frac{9}{2M^2}(\mathcal{D}\cdot A)^2 + \frac{1}{2}(1+2\chi)A_m^2. \tag{4.19}$$

Using these results in (4.16) one finds

$$\mathcal{L} = \frac{1}{2}(1+2\chi)\big(R + A_m^2\big) - \frac{M^2}{3}\chi^2 + \frac{9}{2M^2}(\mathcal{D}\cdot A)^2, \tag{4.20}$$

and finally, eliminating χ via its algebraic field equation

$$\chi = \frac{3}{2M^2}\big(R + A_m^2\big), \tag{4.21}$$

one reaches the bosonic terms of the higher-derivative supergravity Lagrangian.

The final redefinition

$$A_m \to \sqrt{\frac{2}{3}}A_m \tag{4.22}$$

recasts this Lagrangian in the notation of [10],

$$\mathcal{L} = \frac{1}{2}\left(R + \frac{2}{3}A_m^2\right) + \frac{3}{4M^2}\left(R + \frac{2}{3}A_m^2\right)^2 + \frac{3}{M^2}(\mathcal{D}\cdot A)^2, \tag{4.23}$$

and the linearized higher-derivative terms then reproduce precisely the combination

$$\frac{3}{4M^2}\big[R^2 - 4A^\rho\partial_\rho\partial^\mu A_\mu\big], \tag{4.24}$$

in agreement with [23].

The bosonic Lagrangian (4.23) propagates one scalar and one pseudoscalar degree of freedom, in addition to gravity. The scalar degree of freedom draws its origin, as is well known, from the R^2 term, which is manifest in our construction, while the need for the pseudoscalar one was pointed out long ago in [23]. The field equation for A_m following from Eq. (4.23) can be turned into a Klein–Gordon equation for $\mathcal{D}\cdot A$, since

$$\partial_m(\mathcal{D}\cdot A) - \frac{M^2}{9}A_m = 0 \quad \to \quad \Box(\mathcal{D}\cdot A) - \frac{M^2}{9}(\mathcal{D}\cdot A) = 0. \tag{4.25}$$

As a result, there is indeed one (pseudo-)scalar degree of freedom whose mass, $\frac{M^2}{9}$, coincides with the mass of the dual axion a.

Let us conclude by stressing the dual gravitational interpretation of the constraint (4.4), which translates into the component relations

$$\bar{u} \equiv S + iP = M\frac{GG}{2F_X}, \tag{4.26}$$

$$\gamma^{mn}\mathcal{D}_m\psi_n = MG, \tag{4.27}$$

$$-\frac{1}{2}R - \frac{1}{3}A_m^2 + i\mathcal{D}^m A_m - \frac{1}{3}u\bar{u} = MF_X, \tag{4.28}$$

and thus links the (pseudo-)scalar auxiliary fields to the goldstino. Notice that in this dual formulation the goldstino is determined by the gauge-invariant expression in Eq. (4.27). All in all, the off-shell formulation is crucial for the consistency of the theory, and indeed the spacetime curvature R is not fixed in any way by the constraint (4.5), but is dynamically determined by the expectation value $\langle F_X\rangle$ of the auxiliary field F_X. Notice also that Eq. (4.26) implies that u is nilpotent, a fact that we used in deriving the Lagrangian of Eq. (4.23).

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