PROCEEDINGS OF SCIENCE

$\mathbf{R}+\mathbf{R}^2$ supergravity, auxiliary fields and cosmology

Ioannis Dalianis
Physics Division, National Technical University of Athens,
15780 Zografou Campus, Athens, Greece
E-mail: dalianis@mail.ntua.gr

Fotis Farakos*
Institute for Theoretical Physics, Masaryk University,
611 37 Brno, Czech Republic
E-mail: fotisf@mail.muni.cz

In this article we deal with the old-minimal higher curvature supergravity and we examine the auxiliary fields sector. These fields become propagating and can endow the theory with new properties. Regarding the inflationary observables, the models yield similar predictions to the standard Starobinsky model, however, the initial conditions problem is ameliorated. Moreover, in an R-symmetry violating setup, higher curvature supergravity can both drive inflation and lead to de Sitter vacua at the end of inflation.

Proceedings of the Corfu Summer Institute 2014 "School and Workshops on Elementary Particle Physics and Gravity",
3-21 September 2014
Corfu, Greece

*Speaker.

PoS(CORFU2014)147

1. Old-minimal higher curvature supergravity

Supersymmetric theories contain an equal number of fermionic and bosonic degrees of freedom. On-shell this leads to an equal number of propagating fermionic modes and propagating bosonic modes. Off-shell, the number of propagating modes do not always match, and auxiliary fields are essentially introduced. For the old-minimal supergravity [1] we have the multiplet

$$e_m^a\,,\ \ \psi_m^\alpha\,,\ \ M\,,\ \ b_m. \tag{1.1}$$

The graviton e_m^a and the gravitino ψ_m^α contain the propagating modes whereas M is a complex auxiliary scalar and b_m a real auxiliary vector. This assignment of auxiliary and propagating modes holds only for the supergravity theories with no higher curvature terms. When higher curvature terms are introduced some of the auxiliary fields become propagating [2]. For the old-minimal one finds a pair of massive chiral multiplets [3] and for the new-minimal one finds a massive vector multiplet [4]. The 20/20 $R+R^2$ has been constructed in the linearized level [5], using a formalism which can also describe higher superspin theories [6–8].

Here we will work with the old-minimal supergravity. As we will see, these new multiplets which arise due to the dynamical auxiliary fields can have novel applications for supersymmetry breaking, inflation, initial conditions problem and late time cosmology. Therefore, higher curvature supergravity is equipped with the tools to address various aspects of modern cosmology.

We will work with a generic form of old-minimal higher curvature supergravity [3, 9–11]

$$\mathscr{L} = -3\int d^4\theta\, E\, f(\mathscr{R},\bar{\mathscr{R}}) \tag{1.2}$$

which can be also written as

$$\mathscr{L} = \frac{3}{8}\int d^2\Theta\, 2\mathscr{E}\,(\bar{\mathscr{D}}^2-8\mathscr{R})f(\mathscr{R},\bar{\mathscr{R}})+c.c. \tag{1.3}$$

up to total derivatives. Here $\mathscr{R}$ is the supergravity chiral superfield and $2\mathscr{E}$ is the chiral density, while E is the full superspace density. For the moment we set $M_P = 1$; we will restore dimensions later. In component form, the bosonic sector of (1.3) is [10, 11]

$$\begin{aligned} e^{-1}\mathscr{L} = &-\frac{1}{2}\,(f+Mf_M+\bar{M}f_{\bar{M}}-4M\bar{M}f_{M\bar{M}}-2b^m b_m f_{M\bar{M}})\,R-\frac{3}{4}f_{M\bar{M}}R^2 \\ &+3f_{M\bar{M}}\,\partial M\partial\bar{M}-3f_{M\bar{M}}\,(\nabla^m b_m)^2+i\,(f_M\partial^m M-f_{\bar{M}}\partial^m\bar{M})\,b_m \\ &-i\,(f_M M-f_{\bar{M}}\bar{M})\,\nabla^m b_m-\frac{1}{3}M\bar{M}\,\{f-2\,(f_M M+f_{\bar{M}}\bar{M})+4M\bar{M}f_{M\bar{M}}\} \\ &+\frac{1}{3}b_m b^m\,\{f+f_M M+f_{\bar{M}}\bar{M}-4M\bar{M}f_{M\bar{M}}-b_n b^n f_{M\bar{M}}\} \end{aligned} \tag{1.4}$$

where $f = f\left(-\frac{1}{6}M,-\frac{1}{6}\bar{M}\right)$, $f_M = \partial f/\partial M$ and $f_{\bar{M}} = \partial f/\partial\bar{M}$. A very interesting discussion on (1.4) can be found in [10]. Note that this is a theory of curvature and curvature square terms only, as far as gravitation is concerned. Nevertheless it does not propagate the same degrees of freedom as standard gravity. The $R+R^2$ theory on top of the dynamical degrees of freedom of the metric, also gives rise to an additional real scalar propagating degree of freedom known as the scalaron. For a gravitational theory the counting of the degrees of freedom ends here. For the supergravitational embedding, as can be seen from (1.4), the scalaron comes with supersymmetric scalar partners (and of course fermion superpartners which we have not included (1.4)). The counting of the total scalar supersymmetric degrees of freedom is [2, 10–12]

- M: 2 real scalar degrees of freedom,
- $\nabla^m b_m$: 1 real scalar degree of freedom,
- $R^2 \to scalaron$: 1 real scalar degree of freedom.

The above degrees of freedom reside inside appropriate supersymmetric multiplets as shown in [2], and are the reason why one needs exactly two chiral superfields ($\mathscr{S}$ and $\mathscr{T}$) to perform the duality of the theory (1.4) to standard supergravity; the degrees of freedom should match.

These generic pure supergravity self-couplings can be brought in first order form by introducing appropriate Lagrange multipliers, giving rise to an equivalent theory which contains two chiral superfields coupled to standard supergravity, denoted by $\mathscr{T}$ and $\mathscr{S}$. We start from the Lagrangian

$$\mathscr{L} = \frac{3}{8}\int d^2\Theta 2\mathscr{E}\,(\bar{\mathscr{D}}^2 - 8\mathscr{R})f(\mathscr{S},\bar{\mathscr{S}}) + c.c. + 6\int d^2\Theta 2\mathscr{E}\,\mathscr{T}(\mathscr{S}-\mathscr{R}) + c.c. \tag{1.5}$$

Indeed, from the superspace equations of motion of $\mathscr{T}$ we have $\mathscr{S} = \mathscr{R}$ which leads to Lagrangian (1.3). On the other hand we may rewrite Lagrangian (1.5) as

$$\begin{aligned}\mathscr{L} =& \frac{3}{8}\int d^2\Theta 2\mathscr{E}\,(\bar{\mathscr{D}}^2 - 8\mathscr{R})\left[\mathscr{T} + \bar{\mathscr{T}} + f(\mathscr{S},\bar{\mathscr{S}})\right] + c.c.\\ &+ 6\int d^2\Theta 2\mathscr{E}\,\mathscr{T}\mathscr{S} + c.c.\end{aligned} \tag{1.6}$$

which is nothing but standard supergravity

$$\mathscr{L} = \frac{3}{8}\int d^2\Theta 2\mathscr{E}\,(\bar{\mathscr{D}}^2 - 8\mathscr{R})e^{-\frac{1}{3}\mathscr{K}} + c.c. + \int d^2\Theta 2\mathscr{E}\,\mathscr{W} + c.c. \tag{1.7}$$

with Kähler potential

$$\mathscr{K} = -3\ln\left\{\mathscr{T} + \bar{\mathscr{T}} + f(\mathscr{S},\bar{\mathscr{S}})\right\} \tag{1.8}$$

and superpotential

$$\mathscr{W} = 6\,\mathscr{T}\mathscr{S}. \tag{1.9}$$

Thus, the bottom line of the above discussion is that the standard supergravity theory coupled to the chiral superfields $\mathscr{T}$ and $\mathscr{S}$ with a Kähler potential given by (2.4) and a superpotential given by (2.5) is equivalent to a supergravity theory containing only pure supergravity sector and its higher derivatives. This procedure was initially established in the language of superconformal supergravity in [3], and here we have followed it closely. The function $f(\mathscr{R},\bar{\mathscr{R}})$ in principle may contain terms which violate the R-symmetry [11] .

2. R-symmetric models: Starobinsky inflation

The embedding of the Starobinsky model of inflation in old-minimal supergravity in superspace consists of reproducing the Lagrangian

$$e^{-1}\mathscr{L} = -\frac{M_P^2}{2}R + \frac{M_P^2}{12m^2}R^2. \tag{2.1}$$

This is achieved by [3, 9, 11, 13]

$$\mathscr{L} = -3M_P^2 \int d^4\theta\, E \left[1 - \frac{4}{m^2}\mathscr{R}\bar{\mathscr{R}} + \frac{\zeta}{3m^4}\mathscr{R}^2\bar{\mathscr{R}}^2\right]. \tag{2.2}$$

Modifications and further properties can be found in [12, 14–25]. Lagrangian (2.2) when expanded to components yields R^2 terms and kinematic terms for M and b_m. One may work directly with (2.2) but it is more convenient to turn to the dual description in terms of two chiral superfields: $\mathscr{T}$ and $\mathscr{S}$. The equivalent description of the above higher curvature supergravity reads

$$\mathscr{L} = \int d^2\Theta 2\mathscr{E}\left[\frac{3M_P^2}{8}(\bar{\mathscr{D}}^2 - 8\mathscr{R})e^{-\frac{\mathscr{K}}{3M_P^2}}\right] + c.c. + \int d^2\Theta 2\mathscr{E}\mathscr{W} + c.c., \tag{2.3}$$

with Kähler potential

$$\mathscr{K} = -3M_P^2 \ln\left\{1 + \frac{\mathscr{T}+\bar{\mathscr{T}}}{M_P} - 4\frac{\mathscr{S}\bar{\mathscr{S}}}{M_P^2} + \frac{1}{3}\zeta\frac{\mathscr{S}^2\bar{\mathscr{S}}^2}{M_P^4}\right\}, \tag{2.4}$$

and superpotential

$$\mathscr{W} = 6m\,\mathscr{T}\mathscr{S}. \tag{2.5}$$

For the lowest components we have $\mathscr{S}| = S$ and $\mathscr{T}| = T$. For appropriate ζ-parameter values one can see that the fields ImT and S are strongly stabilized to $S = 0$ and Im$T = 0$ [9, 11, 13]. Therefore one finds the effective model

$$e^{-1}\mathscr{L} = -\frac{M_P^2}{2}R - \frac{1}{2}\partial\varphi\partial\varphi - \frac{3}{4}m^2M_P^2\left(1 - e^{-\sqrt{\frac{2}{3}}\varphi/M_P}\right)^2. \tag{2.6}$$

From the analysis of the perturbations during inflation [26] and Planck data [27, 28] we get

$$m \simeq 1.3 \times 10^{-5} M_P. \tag{2.7}$$

For a discussion on the new-minimal supergravity embedding and the massive vector multiplet see [9, 29–35].

3. Initial conditions problem

In this section we work with the R-symmetric models and we also assume large ζ-parameter in (2.4), such that S is strongly stabilized. We want to study the impact of the dynamical auxiliary fields of supergravity on the initial conditions problem [35].

PoS(CORFU2014)147

Let us first remind the reader that for $R+R^2$ gravity, the minimum initially *homogeneous* region required for inflation to start has radius

$$\begin{aligned} D_{\text{homog}}(t_P,\, w_{R^2}) &= d_{\text{event}}(t_P, t_{\max}) + H^{-1}(t_{\text{INF}}) \frac{a(t_P)}{a(t_{INF})} \\ &\sim 4884\, l_P + H^{-1}(t_{\text{INF}})\, \frac{1}{47} \,\sim\, 8.7 \times 10^3\, l_P \end{aligned} \tag{3.1}$$

which is a rather large number. This is a result of the kinetic energy domination for $\rho \gg m^2 M_P^2$, since the Starobinsky model inflationary potential is bounded from above.

Supergravity on the other hand, allows to have access to M_P^4 energy densities, if $\text{Im}T$ takes large values. First, for large $\text{Im}T$ values this component also becomes dynamical and after the redefinition

$$\begin{aligned} \varphi &= \sqrt{\frac{3}{2}} M_P \ln\left[1 + 2\,(\text{Re}T/M_P)\right], \\ b &= \sqrt{6}\,\text{Im}T, \end{aligned} \tag{3.2}$$

we have

$$e^{-1}\mathscr{L} = -\tfrac{M_P^2}{2}R - \tfrac{1}{2}\partial\varphi\partial\varphi - \tfrac{1}{2}e^{-2\sqrt{\frac{2}{3}}\varphi/M_P}\partial b \partial b - V_{\text{sugra}R^2}(\varphi, b), \tag{3.3}$$

where

$$V_{\text{sugra}R^2}(\varphi, b) = \frac{3}{4}m^2 M_P^2 \left(1 - e^{-\sqrt{\frac{2}{3}}\varphi/M_P}\right)^2 + \frac{1}{2}m^2 e^{-2\sqrt{\frac{2}{3}}\varphi/M_P} b^2. \tag{3.4}$$

Second, the field S remains strongly stabilized and will not affect the evolution.

We now turn to a flat FLRW background and study the evolution of the fields and of the spacetime. The equations of motion for the fields φ and b read

$$\begin{aligned} &\ddot{\varphi} + 3H\dot{\varphi} + \sqrt{\frac{3}{2}} m^2 M_P e^{-\sqrt{\frac{2}{3}}\varphi/M_P} \left(1 - e^{-\sqrt{\frac{2}{3}}\varphi/M_P}\right) \\ &\quad - \sqrt{\frac{2}{3}} e^{-2\sqrt{\frac{2}{3}}\varphi/M_P} \left(m^2 b^2 - \dot{b}^2\right) = 0, \end{aligned} \tag{3.5}$$

and

$$\ddot{b} + 3H\dot{b} - 2\sqrt{\frac{2}{3}}\dot{\varphi}\dot{b} + m^2 b = 0. \tag{3.6}$$

The Friedmann equation reads

$$3H^2 M_P^2 = \frac{1}{2}\dot{\varphi}^2 + \frac{1}{2}e^{-2\sqrt{\frac{2}{3}}\varphi/M_P}\dot{b}^2 + V_{\text{sugra}R^2}(\varphi, b). \tag{3.7}$$

We want to study the evolution of this system which has been also discussed in [36–38] and in a different context in [39]. We choose initial conditions such that the initial energy density, $\rho_{\text{init}} = M_P^4$, is equally partitioned between the kinetic and the potential terms,

$$V(\varphi_{\text{init}},\, b_{\text{init}}) = \rho_{\text{kin, init}} = \frac{1}{2}M_P^4. \tag{3.8}$$

An *example* of such a set of values that realize these initial conditions is,

$$(\varphi_{\rm init}, b_{\rm init}) = (0, \frac{M_P^2}{m}) \quad \text{and} \quad \dot{\varphi}_{\rm init} = \dot{b}_{\rm init} = \frac{1}{\sqrt{2}} M_P^2 \,. \tag{3.9}$$

The system of the fields (φ, b) starts from nonzero values such that $V(\varphi_{\rm init}, b_{\rm init}) \sim M_P^4$. Due to the small mass of the b field, $m_b \ll H$, the φ will roll down the potential (3.4) which for constant b has the form

$$V_{{\rm sugra}R^2}(\varphi) \sim V_0 e^{-2\sqrt{\frac{2}{3}}\varphi/M_P} \,. \tag{3.10}$$

Initially, the approximation $\dot{b} \sim 0$ is a good one according to the numerical results. For the potential (3.10) the Friedmann equation have an exact solution of power law form given by the expressions

$$\begin{aligned} a \propto a_{\rm init} t^n \,, \quad n = 3/4 \\ \varphi = \sqrt{\frac{3}{2}} \ln\left(\sqrt{\frac{16}{15}\frac{V_0}{M_P^2}}\, t\right) \,. \end{aligned} \tag{3.11}$$

We see that $n = 2/[3(1+w)] = 3/4$. This corresponds to a barotropic fluid with equation of state $w = -1/9$, that is a negative pressure. Numerically we find that in the pre-inflation period the energy density decreases with a slower rate,

$$w_{{\rm sugra}R^2} \lesssim -1/9 \,, \quad a_{{\rm sugra}R^2}(t) \gtrsim t^{3/4} \,, \quad \rho_{{\rm sugra}R^2} \gtrsim \rho_{\rm init} a^{-8/3} \,, \tag{3.12}$$

because the actual system is a two-field one.

The approximation (3.11), to consider a constant equation of state $w = -1/9$, yields an event horizon radius $d_{\rm event} \propto (t_{\rm max}/t_P)^{1/4} \sim 87\, l_P$ for $t_{\rm max} = 10^5\, t_P$. The exact result can be obtained numerically for the varying equation of state $w_{{\rm sugra}R^2}$ (3.12). When we integrate from the Planck time until the beginning of inflation, which is found to be $t_{\rm INF} = 0.74 \times 10^5\, t_P$, we numerically take

$$d_{\rm event}\left(t_{\rm init} = t_P, t_{\rm INF}, w_{{\rm sugra}R^2}\right) \sim 29\, H_P^{-1} \,. \tag{3.13}$$

As in the $R + R^2$ gravity case, the event horizon increases as long as $w > -1$. It remains constant when the field configuration lies in the plateau with vanishing kinetic energy, $w \cong -1$. The numerical value of the total event horizon reads

$$d_{\rm event}\left(t_{\rm init} = t_P, t_{\rm max}, w_{{\rm sugra}R^2}\right) \sim 46\, H_P^{-1} \sim 80\, l_P \,. \tag{3.14}$$

The minimum initially *homogeneous* region required for inflation to start in the supergravity case has radius

$$\begin{aligned} D_{\rm homog}(t_P, w_{{\rm sugra}R^2}) &= d_{\rm event}(t_P, t_{\rm max}) + H^{-1}(t_{\rm INF}) \frac{a(t_P)}{a(t_{INF})} \\ &\sim 80\, l_P + H^{-1}(t_{\rm INF}) \frac{1}{5407} \sim 98\, l_P \,. \end{aligned} \tag{3.15}$$

That is, right after the Planck time the initial homogeneous volume is required to have radius at least 68 times the Planck length. The minimum number of the CDR is here

$$\frac{V_{\rm flat}(D_{\rm homog}, w_{{\rm sugra}R^2})}{V_{\rm flat}(l_P)} = \frac{\frac{4}{3}\pi D_{\rm homog}^3}{\frac{4}{3}\pi l_P^3} \sim 10^6 \text{ CDR} \,. \tag{3.16}$$

Compared to the non-supersymmetric case, in the $R+R^2$ supergravity the required initial homogeneous volume is about half a million times smaller,

$$\frac{\#\mathrm{CDR}_{R^2}}{\#\mathrm{CDR}_{\mathrm{sugra}R^2}} = \frac{D^3_{\mathrm{homog}}(t_P,\, w_{R^2})}{D^3_{\mathrm{homog}}(t_P,\, w_{\mathrm{sugra}R^2})} \sim 7\times 10^5\,. \tag{3.17}$$

This is a significant relaxation of the initial conditions problem.

In recent work [11] there has been found a new class of *R-symmetry violating $R+R^2$ models* which can both provide an inflationary sector and a hidden supersymmetry breaking sector, without invoking any matter superfields. These models we review in the next section. The new properties of these models which distinguish them from the R-symmetric $R+R^2$ old-minimal supergravity is that at the end of inflation the S field contribution starts to become important and the field configuration is driven towards the supersymmetry breaking vacuum. For these models it is also expected that the initial conditions problem is similar to the R-symmetric case that we analyzed here.

4. R-violating models

In this section we will also introduce R-violating terms. These models can provide inflationary candidates which fit the Planck data. Moreover, due to the supergravity dynamical auxiliary fields there is a new class of vacua which break supersymmetry. Hence the issue of vacuum selection is raised (see for example [40, 41]). It is interesting that the inflationary trajectory terminates in the supersymmetry breaking vacua.

An extensive discussion on these models can be found in [11]. Here we discuss their properties using a characteristic example

$$f(\mathscr{R},\bar{\mathscr{R}}) = 1+\gamma\frac{\mathscr{R}}{m}+\gamma\frac{\bar{\mathscr{R}}}{m}-2\frac{\mathscr{R}\bar{\mathscr{R}}}{m^2}+\frac{1}{9}\zeta\frac{\mathscr{R}^2\bar{\mathscr{R}}^2}{m^4} \tag{4.1}$$

where the second and third term source the R-symmetry breaking, and the equivalent Kähler potential is

$$\mathscr{K} = -3M_P^2\ln\left\{1+\frac{\mathscr{T}+\bar{\mathscr{T}}}{M_P}+\gamma\frac{\mathscr{S}+\bar{\mathscr{S}}}{M_P}-2\frac{\mathscr{S}\bar{\mathscr{S}}}{M_P^2}+\frac{1}{9}\zeta\frac{\mathscr{S}^2\bar{\mathscr{S}}^2}{M_P^4}\right\} \tag{4.2}$$

while the superpotential is

$$\mathscr{W} = 6m\,\mathscr{T}\,\mathscr{S}\,. \tag{4.3}$$

The scalar potential of this theory (which can be found in [11]) will in general have two classes of vacuum solutions. First there is the trivial vacuum

$$\langle T\rangle = \langle S\rangle = 0 \tag{4.4}$$

with no supersymmetry breaking and vanishing vacuum energy. Then there is the new class of vacua

$$\begin{aligned}\langle T\rangle &= M_P\langle t\rangle + iM_P\langle b\rangle = M_P\,t_0\\ \langle S\rangle &= M_P\langle s\rangle + iM_P\langle c\rangle = M_P\,s_0\end{aligned} \tag{4.5}$$

which will break supersymmetry with vanishing vacuum energy or small positive (or negative) vacuum energy depending on the parameters ζ and γ. For a Minkowski vacuum it is possible to relate s_0 to ζ and γ as follows

$$\zeta = \frac{1+2s_0^2}{3s_0^4} \tag{4.6}$$

$$\gamma = -2s_0 + \frac{2+4s_0^2}{3s_0}. \tag{4.7}$$

Therefore, the parameters ζ and γ for Minkowski vacua are parameterized by the vacuum expectation value of the s-field. An analysis of the inflationary properties of this model can be found in [11], where it is shown that it gives predictions consistent with the Planck data.

To further illustrate the features of the model we turn to a simple example where

$$\zeta = 8 \ , \ \ \gamma = 1. \tag{4.8}$$

The scalar potential will have a supersymmetry preserving vacuum (4.4) and on top of that there will be the vacuum

$$\begin{aligned} \langle T \rangle &= M_P t_0 = \frac{2}{3} M_P \\ \langle S \rangle &= M_P s_0 = \frac{1}{2} M_P \end{aligned} \tag{4.9}$$

which breaks supersymmetry with vanishing vacuum energy. A small deformation of the γ or ζ parameters leads to de Sitter, or anti-de Sitter vacua

$$\zeta = 8 \ , \ \gamma \gtrsim 1 \rightarrow \ \text{de Sitter vacuum with small vacuum energy.} \tag{4.10}$$

Let us also give a qualitative explanation of the inflationary predictions for this example. First, the fields b and c will be strongly stabilized to $b = 0$ and $c = 0$. Secondly, for large t values ($t > 30$) it is consistent to approximate

$$s_t \simeq \frac{5}{3} t^{-1} \ll 1. \tag{4.11}$$

This leads to the inflationary effective Lagrangian for t (s is strongly stabilized to $s \simeq 0$)

$$e^{-1}\mathscr{L} = -\frac{M_P^2}{2}R - \frac{3M_P^2}{(1+2s_t+2t)^2}\partial t \partial t - 6m^2 M_P^2 \frac{(t-s_t)^2}{(1+2s_t+2t)^2}. \tag{4.12}$$

After a redefinition of the t field

$$t = \frac{1}{2}\left(e^{\sqrt{\frac{2}{3}}\varphi/M_P} - 1 - 2s_t\right) \tag{4.13}$$

we have (ignoring derivatives on s_t because they will be highly suppressed)

$$e^{-1}\mathscr{L} = -\frac{M_P^2}{2}R - \frac{1}{2}\partial\varphi\partial\varphi - \frac{3}{2}m^2 M_P^2 \left(1 - (1+4s_t)e^{-\sqrt{\frac{2}{3}}\varphi/M_P}\right)^2. \tag{4.14}$$

From the number $\frac{2}{3}$ in the exponential we see that the model is similar to the original Starobinsky model and will give rise to the similar amount of gravitational waves which is favored by the Planck collaboration data [27,28]. Note that for the Starobinsky model in supergravity $s_t = 0$ always. For an approximate number of $N(t_*) \simeq 55$ e-folds, we find the value $t_* = 50$. Therefore the tensor-to-scalar ratio r [26] is

$$r = 16\varepsilon(t_*) \simeq 8 \times 10^{-3} \tag{4.15}$$

at the pivot scale t_*. Thus the model predicts very small amount of gravitational waves. Finally, for a ζ-parameter of order 10, s_0 is of order a half (as happens in our example), and the gravitino and inflaton mass turn out to be

$$m_{3/2}^2 \simeq m^2 \,, \quad m_{inf}^2 \simeq \frac{1}{4} m^2 \,. \tag{4.16}$$

We see that these models imply a relation of the inflaton mass to the gravitino mass $m_{3/2}$ via the new scale m, but notice that

$$m_{3/2}^2 > m_{inf}^2 . \tag{4.17}$$

Therefore higher curvature supergravity may account both for the inflationary phase and for the de Sitter vacuum at the end of inflation [11]. This property is related to the auxiliary fields which have become dynamical, due to the R^2 terms.

Acknowledgments

It is a pleasure to thank A. Kehagias, A. Riotto and R. von Unge for our collaborations and discussions. This work was supported by the Grant agency of the Czech republic under the grant P201/12/G028.

References

[1] J. Wess and J. Bagger, "Supersymmetry and supergravity," Princeton, USA: Univ. Pr. (1992) 259 p

[2] S. Ferrara, M. T. Grisaru and P. van Nieuwenhuizen, "Poincare and Conformal Supergravity Models With Closed Algebras," Nucl. Phys. B **138**, 430 (1978).

[3] S. Cecotti, "Higher Derivative Supergravity Is Equivalent To Standard Supergravity Coupled To Matter. 1.," Phys. Lett. B **190**, 86 (1987).

[4] S. Cecotti, S. Ferrara, M. Porrati and S. Sabharwal, "New Minimal Higher Derivative Supergravity Coupled To Matter," Nucl. Phys. B **306**, 160 (1988).

[5] F. Farakos, A. Kehagias and K. Koutrolikos, "Linearized Non-Minimal Higher Curvature Supergravity," arXiv:1501.07562 [hep-th].

[6] S. J. Gates and K. Koutrolikos, "On 4D, $\mathcal{N} = 1$ massless gauge superfields of arbitrary superhelicity," JHEP **1406**, 098 (2014).

[7] Gates, S. James., Jr. and K. Koutrolikos, "On 4D, N = 1 Massless Gauge Superfields of Higher Superspin: Half-Odd-Integer Case," arXiv:1310.7386 [hep-th].

[8] S. J. Gates, Jr. and K. Koutrolikos, “On 4D, N = 1 Massless Gauge Superfields of Higher Superspin: Integer Case,” arXiv:1310.7385 [hep-th].

[9] F. Farakos, A. Kehagias and A. Riotto, “On the Starobinsky Model of Inflation from Supergravity,” Nucl. Phys. B **876**, 187 (2013) [arXiv:1307.1137].

[10] A. Hindawi, B. A. Ovrut and D. Waldram, “Four-dimensional higher derivative supergravity and spontaneous supersymmetry breaking,” Nucl. Phys. B **476**, 175 (1996) [hep-th/9511223].

[11] I. Dalianis, F. Farakos, A. Kehagias, A. Riotto and R. von Unge, “Supersymmetry Breaking and Inflation from Higher Curvature Supergravity,” JHEP **1501**, 043 (2015) [arXiv:1409.8299 [hep-th]].

[12] I. Antoniadis, E. Dudas, S. Ferrara and A. Sagnotti, “The Volkov-Akulov-Starobinsky supergravity,” Phys. Lett. B **733**, 32 (2014) [arXiv:1403.3269 [hep-th]].

[13] R. Kallosh and A. Linde, “Superconformal generalizations of the Starobinsky model,” JCAP **1306**, 028 (2013) [arXiv:1306.3214 [hep-th]].

[14] S. Ferrara, R. Kallosh and A. Linde, “Cosmology with Nilpotent Superfields,” arXiv:1408.4096 [hep-th].

[15] G. Dall’Agata and F. Zwirner, “On sgoldstino-less supergravity models of inflation,” JHEP **1412**, 172 (2014) [arXiv:1411.2605 [hep-th]].

[16] J. Ellis, D. V. Nanopoulos and K. A. Olive, “No-Scale Supergravity Realization of the Starobinsky Model of Inflation,” Phys. Rev. Lett. **111**, 111301 (2013) [Erratum-ibid. **111**, no. 12, 129902 (2013)] [arXiv:1305.1247 [hep-th]].

[17] J. Ellis, D. V. Nanopoulos and K. A. Olive, “Starobinsky-like Inflationary Models as Avatars of No-Scale Supergravity,” JCAP **1310**, 009 (2013) [arXiv:1307.3537].

[18] S. Renaux-Petel and K. Turzynski, “On reaching the adiabatic limit in multi-field inflation,” arXiv:1405.6195 [astro-ph.CO].

[19] K. Kamada and J. Yokoyama, “Topological inflation from the Starobinsky model in supergravity,” Phys. Rev. D **90**, no. 10, 103520 (2014) [arXiv:1405.6732 [hep-th]].

[20] S. Ferrara and A. Kehagias, “Higher Curvature Supergravity, Supersymmetry Breaking and Inflation,” arXiv:1407.5187 [hep-th].

[21] T. Terada, Y. Watanabe, Y. Yamada and J. Yokoyama, “Reheating processes after Starobinsky inflation in old-minimal supergravity,” arXiv:1411.6746 [hep-ph].

[22] J. Alexandre, N. Houston and N. E. Mavromatos, “Starobinsky-type Inflation in Dynamical Supergravity Breaking Scenarios,” Phys. Rev. D **89**, no. 2, 027703 (2014) [arXiv:1312.5197 [gr-qc]].

[23] A. Kehagias, A. M. Dizgah and A. Riotto, “Comments on the Starobinsky Model of Inflation and its Descendants,” Phys. Rev. D **89**, 043527 (2014) [arXiv:1312.1155 [hep-th]].

[24] A. Kehagias, C. Kounnas, D. Lust and A. Riotto, “Black Hole Solutions in R^2 Gravity,” arXiv:1502.04192 [hep-th].

[25] Y. S. Myung, “Stability of black holes in f(R) gravity,” arXiv:1503.03559 [gr-qc].

[26] D. H. Lyth and A. Riotto, “Particle physics models of inflation and the cosmological density perturbation,” Phys. Rept. **314** (1999) 1 [hep-ph/9807278].

PoS(CORFU2014)147

[27] P. A. R. Ade *et al.* [Planck Collaboration], "Planck 2013 results. XXII. Constraints on inflation," arXiv:1303.5082 [astro-ph.CO].

[28] P. A. R. Ade *et al.* [Planck Collaboration], "Planck 2015 results. XX. Constraints on inflation," arXiv:1502.02114 [astro-ph.CO].

[29] S. Ferrara, R. Kallosh, A. Linde and M. Porrati, "Minimal Supergravity Models of Inflation," Phys. Rev. D **88**, no. 8, 085038 (2013) [arXiv:1307.7696 [hep-th]].

[30] S. Ferrara, R. Kallosh, A. Linde and M. Porrati, "Higher Order Corrections in Minimal Supergravity Models of Inflation," JCAP **1311**, 046 (2013) [arXiv:1309.1085 [hep-th], arXiv:1309.1085].

[31] S. Ferrara, P. FrÃľ and A. S. Sorin, "On the Topology of the Inflaton Field in Minimal Supergravity Models," JHEP **1404**, 095 (2014) [arXiv:1311.5059 [hep-th]].

[32] A. Ceresole, G. Dall'Agata, S. Ferrara, M. Trigiante and A. Van Proeyen, "A search for an $\mathcal{N} = 2$ inflaton potential," Fortsch. Phys. **62**, 584 (2014) [arXiv:1404.1745 [hep-th]].

[33] F. Farakos and R. von Unge, "Naturalness and Chaotic Inflation in Supergravity from Massive Vector Multiplets," JHEP **1408**, 168 (2014) [arXiv:1404.3739 [hep-th]].

[34] S. Ferrara and M. Porrati, "Minimal $R+R^2$ Supergravity Models of Inflation Coupled to Matter," Phys. Lett. B **737**, 135 (2014) [arXiv:1407.6164 [hep-th]].

[35] I. Dalianis and F. Farakos, "On the initial conditions for inflation with plateau potentials: the $R+R^2$ (super)gravity case," arXiv:1502.01246 [gr-qc].

[36] S. Ferrara, A. Kehagias and A. Riotto, "The Imaginary Starobinsky Model," Fortsch. Phys. **62**, 573 (2014) [arXiv:1403.5531 [hep-th]].

[37] R. Kallosh, A. Linde, B. Vercnocke and W. Chemissany, "Is Imaginary Starobinsky Model Real?," JCAP **1407**, 053 (2014) [arXiv:1403.7189 [hep-th]].

[38] K. Hamaguchi, T. Moroi and T. Terada, "Complexified Starobinsky Inflation in Supergravity in the Light of Recent BICEP2 Result," Phys. Lett. B **733**, 305 (2014) [arXiv:1403.7521 [hep-ph]].

[39] Z. Lalak, D. Langlois, S. Pokorski and K. Turzynski, "Curvature and isocurvature perturbations in two-field inflation," JCAP **0707** (2007) 014 [arXiv:0704.0212 [hep-th]].

[40] I. Dalianis and Z. Lalak, "Cosmological vacuum selection and metastable susy breaking," JHEP **1012**, 045 (2010) [arXiv:1001.4106 [hep-ph]].

[41] I. Dalianis and Z. Lalak, "Thermally Favourable Gauge Mediation," Phys. Lett. B **697** (2011) 385 [arXiv:1012.3157 [hep-ph]].