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Non-perturbative $SUSY$ and $\mathcal{R}$-symmetry breaking ; Inflationary behavior of the no-scale effective $SUGRAS$

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C. Kounnas, D. Lüst and N. Toumbas, Fortsch.Phys. 63 (2015) 12 ; arXiv:1409.7076

1. Introduction

This work is motivated by the observation that many of the superstring theories with spontaneously broken susy define at low energies effective theories in a flat Minkowski space-time similar to the no-scale supergravity models characterized by :

$$V \geq 0 \quad \text{with} \quad \text{classically untetermined} \quad m_{3/2}(T, U, S, \ldots \text{moduli})$$

Cremmer, Ferrara, Kounnas, Nanopoulos[82], Ellis, Kounnas, Nanopoulos[83,84], Ellis, Lahanas, Nanopoulos, Tamvakis[84]...

The minimal $\mathcal{M}_n$=$SU(1,1+n)/U(1) \times SU(1+n)$ no-scale model is extended by other moduli fields of the string theory like:
The dilaton field S, the complex structure fields U_I, ... :

Witten[85], Ferrara, Kounnas, Porrati[86], Ferrara, Kounnas, Zwirner[94], ...

$$K = -3\log Y(\psi^I) - \log(S+\bar{S}) - \log(U_1+\bar{U}_1) - \log(U_2+\bar{U}_2) - \log(U_3+\bar{U}_3) + \ldots$$

$\psi^I \equiv \{T, \Phi^i\}$ denote the fields of the $\mathcal{M}_n$ manifold

$$Y(\psi^I) = T + \bar{T} - |\Phi^i|^2$$

When the superpotential is independent of T the F-part of the potential is positive semi-definite $V_F \geq 0$ thanks to the no-scale structure of the $\mathcal{M}_n$ sub-manifold.

Ellis, Kounnas, Nanopoulos[84], Ferrara Kounnas Zwirner[94]...

$$V_F = \frac{3Y\sum_i |W_{\Phi^i}|^2 + |W - (S+\bar{S})W_S|^2 + \sum_I |W - (U_I + \bar{U}_I)W_{U_I}|^2}{Y^3\,(S+\bar{S})(U_1+\bar{U}_1)(U_2+\bar{U}_2)(U_3+\bar{U}_3)} \geq 0$$

The positive semi-definite structure of V_F follows from the properties of the $SU(1,n)$ Kälher manifold and the particular $SU(1,1)_{S,U_I}/U(1)$ manifolds of the (S, U_I)-moduli.

The negative contribution, $-3e^K|W|^2$, of the $N=1$ sugra potential is cancelled by the contribution of $\psi \equiv \{T, z^i\}$ fields in V_F provided W is T independent.

$$G = K + \log|W|^2\,, \qquad W_T = 0$$

$$e^K\left(G_\psi K^{\psi\bar\psi}G_{\bar\psi}-3\right)|W|^2=\sum_i\left(\frac{3|W_{\Phi^i}|^2}{Y^2}\right)\frac{1}{(S+\bar S)(U_1+\bar U_1)(U_2+\bar U_2)(U_3+\bar U_3)}$$

The D-part of the potential is always positive semi-definite $V_D \geq 0$

$$V_D=\sum_\alpha g_\alpha^2\left[\,G_A\,(T^\alpha)_B^A\;z^B\,\right]^2\;\geq 0\,,\qquad z^B=\{T,\Phi^i,S,U^I,\dots\}\,.$$

The only way to obtain non-trivial susy-breaking contributions from the D-terms is when they are associated to the gauging of $U(1)_R$ symmetries under which the superpotential W transforms non-trivially under $U(1)_R$:

$$U(1)_R:\quad W\longrightarrow e^{i\xi_W\theta_R}\,W,\qquad \Phi\longrightarrow e^{i\xi_\Phi\theta_R}\,\Phi$$

In these specific cases a Fayet-Iliopoulos term is generated giving rise to a non trivial (metastable) de Sitter vacua Kounnas, Lüst, Toumbas[14], Ferrara, Kehagias, Porrati[14] ...

$$V_D = \frac{1}{2} g_R^2 \left(\xi_W + 3\xi_\Phi \frac{|\Phi|^2}{Y} + \ldots \right)^2$$

The " $\ldots$ " stand for the contribution of the other fields $T, S, U_I, \ldots$

2 . D-breaking: The emergence of the $\mathcal{R}^2 + \tilde{\xi}\mathcal{R}$ gravity

Dictated by string theory a common form of the superpotenial for large T

$$W = A_{ijk}\, \Phi^i \Phi^j \Phi^k + w(S, U_I) \equiv \left[\Phi^3\right] + w\,.$$

it contains two typical terms:

i) the R-symmetric susy terms $\left[\Phi^3\right]$

ii) the w -terms which can be generated by fluxes and/or non-perturbative terms:

$$w = \hat{w}\, e^{-\beta S}$$

The initial $U(1)_R$ symmetry is generically broken by the w-terms.
It can be restored if $U(1)_R$ acts nontrivially on w or when w is trivial.

$$U(1)_R: \quad \Phi \longrightarrow e^{i\xi_\Phi \theta_R}\, \Phi \qquad w \longrightarrow e^{i\xi_W \theta_R}\, w\,; \quad \xi_W = 3\xi_\Phi\,, \qquad \text{OR} \qquad w \equiv 0.$$

Although the classical $U(1)_R$ is restored in both cases, at the quantum level $U(1)_R$ is anomalous due to its action on the chiral fermions (the susy partners of Φ_i).

It is therefore necessary for the anomaly cancellation to extend the action of $U(1)_R$ to T field as axion translation : Kounnas, Lüst, Toumbas[14]

$$U(1)_R: \quad T \longrightarrow T + i\xi_T\, \theta_R \qquad \text{with} \qquad \xi_T = \mathrm{Tr}\, Q$$

For $w \equiv 0$, V_D becomes the only relevant part and has the well known form of the Starobinsky inflationary potential for $W_{\Phi^i} = 0$ (and $W = 0$ since $\xi_\Phi\, \Phi^i\, W_{\Phi^i} = 3W$)

$$V = (V_F + V_D)\,|_{W_{\Phi^i}=0} = \mu^2 \left(1 - \tilde{\xi}\, e^{-\alpha\phi} + \hat{\Phi}_i^2\right)^2 ; \qquad \tilde{\xi} = \left(\frac{\xi_T}{\xi_\Phi}\right), \quad \mu^2 = \frac{1}{2} g^2 (3\xi_\Phi)^2$$

In terms of canonically normalized fields,

$$\hat{\Phi}_i = \Phi_i\, e^{-\frac{\alpha}{2}\phi}, \qquad \alpha\phi = \log Y, \quad Y = (T + \bar{T} - |\Phi|)^2 ; \qquad \alpha = \sqrt{\frac{2}{3}},$$

the action in the Einstein frame in terms of the no-scale modulus ϕ and Φ_i or $\hat{\Phi}_i$:

$$S_E = \int d^4x \sqrt{|g|} \times$$

$$\frac{1}{2}\left[\mathcal{R} - \partial_\mu\phi\, \partial^\mu\phi - e^{-\alpha\phi}\, \partial_\mu\Phi_i \partial^\mu\Phi_i - 2\mu^2 \left(1 - \tilde{\xi}\, e^{-\alpha\phi} + \hat{\Phi}_i^2\right)^2 - 2\, V_F(\hat{\Phi}_i; S, U_I, \ldots) + \ldots\right]$$

- When $\tilde{\xi} = 0$ the full classical action is invariant under a scale symmetry:

$$\alpha\phi \to \alpha\phi + 2\sigma\,, \qquad \Phi_i \to e^{\sigma}\Phi_i\,, \qquad g_{\mu\nu} \to g_{\mu\nu}\,,$$

(V_F is function of $\hat{\Phi}_i$; there is no dependence on ϕ in the classical action !)

- When $\tilde{\xi} \neq 0$ The vacuum structure changes drastically. The scale violating term $\{-\tilde{\xi} e^{-\alpha\phi}\}$ induces a slow inflationary transition

  de Sitter phase $< \mathcal{R} >= 4\mu^2 \quad \longrightarrow \quad$ flat susy phase $< \mathcal{R} >= 0$.

  The off-shell potential has three characteristic phases: Kounnas, Lüst, Toumbas[14]

  (i) Approximate de Sitter phase with negligible scale breaking term $\tilde{\xi} e^{-\alpha\phi} \ll 1$.
  In this phase $\hat{\Phi}_i = 0$ thanks to the induced de Sitter mass $m_i^2 = \frac{2}{3}\mu^2$.
  The supersymmetry is broken by the $U(1)_R$ Fayet-Iliopoullos D-term; ($V_F = 0$).

  (ii) Flat-Minkowski supersymmetric vacuum with $W_{\Phi^i} = 0$, $3W = \xi_\Phi \, \Phi^i W_{\Phi^i} = 0$
  which are satisfied with $\hat{\Phi}_i = 0 \longrightarrow SU(n)$-symmetric vacuum with $\tilde{\xi} e^{-\alpha\phi} \sim 1$.

  (iii) Scale-breaking dominant phase with $\xi e^{-\alpha\phi} \gg 1$; $|\hat{\Phi}_i|^2 \neq 0$, with broken $SU(n)$
  Here also the vacuum is supersymmetric since $W_{\Phi^i} = 0$, $3W = \xi_\Phi \, \Phi^i W_{\Phi^i} = 0$.

Therefore the $w = 0$ case provides an interesting slow inflationary transition very similar to the $SU(1,n)$-supersymmetric extension of the minimal Starobinsky model:

$$\tilde{\xi}\mathcal{R} + \mathcal{R}^2 + (\ \Phi\text{-matter}\ ) + (S, U_I + \ldots)\text{-moduli terms.}$$

Utilizing the equations of motion for Φ^i, S, U_I

$$\longrightarrow \qquad W_{\Phi^i} = 0, \qquad 3W = \xi_\Phi\, \Phi^i W_{\Phi^i} = 0 \qquad \text{and} \qquad T\text{-}\bar{T}\,, S, U_I\ = \text{const.}$$

we obtain in the Jordan frame the conformally equivalent action setting:

$$g_{\mu\nu} \rightarrow g_{\mu\nu}\, e^{\log\left(2t + \tilde{\xi} - |\Phi^i|^2\right)}\,, \qquad 2t \equiv T + \bar{T} - \tilde{\xi}\ :$$

$$S_J = \int d^4x \sqrt{|g|} \left\{ \frac{1}{2}\left(2t + \tilde{\xi} - |\Phi^i|^2\right)\mathcal{R} - 4\mu^2\, t^2 - 3\, g^{\mu\nu}\partial_\mu \Phi^i \partial_\nu \bar{\Phi}_i \right\}$$

In this frame the t field appears as a Lagrange multiplier and so,

t can be eliminated by its algebraic equation of motion to obtain:

$$S_J = \int d^4x \sqrt{|g|} \left\{ \frac{1}{2} \left(\frac{\mathcal{R}^2}{8\mu^2} + (\tilde{\xi} - |\Phi^i|^2)\mathcal{R} \right) - 3\, g^{\mu\nu} \partial_\mu \Phi^i \partial_\nu \bar{\Phi}_i \right\}$$

Notice that the $\tilde{\xi}$-term originated by the gauging of the T-translation, and which is necessary for the $U(1)_R$ anomaly cancelation, violates the scale invariance.

• Although the $w = 0$ it gives interesting slow inflationary transition, the over all output it is not satisfactory !

• In the final Flat-Minkowski vacuum the susy turns out to be unbroken !!!

• Many experts in our field think that this is a general statement !

• The main obstruction however comes from the property that if in the final Flat-Minkowski vacuum $W_{\Phi^i} = 0$ implies that $3W = \xi_\Phi\, \Phi^i W_{\Phi^i} = 0$ as well !

This imposes that $m_{3/2} = 0$; susy is always unbroken in the final Flat-Minkowski vacuum.

Can we bypass this obstruction in the cases where $w \neq 0$? (see later)

3. F-susy breaking via fluxes: Absence of R-symmetry

Flat no-scale realization

When $w \neq 0$ the $W_{\Phi^i}{=}0$ minimization does not implies $W = 0$.

Only the susy part $\left[\Phi^3\right] = 0 \quad \longrightarrow \quad < \ W = 0 + w \ >$,

thanks again to the "magic" properties of the $\mathcal{M}_n$ ψ sub-manifold.

Assuming non-trivial fluxes, (similar to those appearing in type IIB), a non-trivial w can by generated, (like for instance)

Derendinger, Kounnas, Zwirner[04], Derendinger, Kounnas, Petropoulos, Zwirner[04]

$$w = A(S + \epsilon U_1)(U_2 + \epsilon U_3), \qquad \epsilon = \pm 1\,.$$

The F-part of the potential is always positive semi-definite $V_F \geq 0$.

For $\epsilon = -1$, the susy is unbroken with $V_F = 0$,

$$\epsilon = -1: \qquad S = U_1, \qquad U_2 = U_3,, \quad \text{with} \qquad m^2_{3/2} = 0.$$

For $\epsilon = 1$ the susy is broken with $V_F = 0$,

$$\epsilon = 1: \qquad S = \bar{U}_1, \qquad U_2 = \bar{U}_3, \quad \text{with} \qquad m^2_{3/2} = \frac{|\hat{A}|^2}{Y^3}$$

Here we recover the structure of the no-scale models where the susy breaking scale depends on the moduli T which is undetermined at the classical level.
The classical degeneracy of the vacuum is fixed at the quantum level.
Ellis, Kounnas, Nanopoulos[83,84], Ellis, Lahanas, Nanopoulos, Tamvakis[84]...

The above example is generic in superstrings when fluxes and (non-) perturbative effects are included. However, when $w \neq 0$ R-symmetry is generically broken.

This implies the absence of D-breaking via Fayet-Iliopoulos term which can induce a (metastable) de Sitter phase and slow inflationary transition of Starobinsky type.

3. D-and F-breaking: Non-perturbative R-symmetry

Including the typical non-perturbative term

$$w = \hat{w}\, e^{-\beta S}$$

the R-symmetry is restored provided that under $U(1)_R$
S is transforming non-trivially while $\hat{w}$ is neutral.

$$U(1)_R: \qquad S \longrightarrow S + i\xi_S\,\theta, \qquad \hat{w} \longrightarrow \hat{w}\,; \qquad \xi_S = -\frac{\xi_W}{\beta} = -3\frac{\xi_\Phi}{\beta}$$

The non-perturbative effects originating from gaugino condensation are of the form

$$w = \mu^3 \exp\left(-\frac{24\pi^2 S}{|b_0|} \right), \qquad \Re\, S = g^{-2}(\mu)\,; \qquad \beta = \frac{24\pi^2}{|b_0|}$$

where b_0 is a one-loop beta-function coefficient.

Dine, Rohm, Seiberg, Witten[85], Derendinger, Ibanez, Nilles[85,86], Derendinbger, Kounnas, Petropoulos[06], . . .

The expectation value of w defines the renormalization-group-invariant scale Λ of the confining gauge group in which gauginos condense, $\langle w \rangle = \Lambda^3$.

The renormalization scale μ is in general a modulus-dependent quantity.

Derendinger, Kounnas, Petropoulos[83,84], . . .

At the one-loop level (or in $N = 2$ theories to all orders) μ is independent of S.

At higher orders a non-trivial S-dependence appear.

In the case of a spectator gauge group like the E_8' in the heterotic
(only vector multiples in the E_8' gauge sector)
all the perturbative corrections can be included, since the exact β-function is known.

$$-\frac{\beta(g)}{g} = \frac{b_0\left(\frac{g^2}{16\pi^2}\right)}{1-\frac{b_0}{3}\left(\frac{g^2}{16\pi^2}\right)} \longrightarrow \Lambda^3 = \hat{\mu}^3(S+\delta)\,e^{-\beta S}$$

If in the gauge sector there is an extra charged matter then (up to two loop-level)

$$\Lambda^3 = \hat{\mu}^3(S+\delta)^{1+\frac{\gamma m}{b_0}}\ e^{-\beta S}$$

The explicit appearance of S is in general an obstruction to the R-symmetry
except iff the δ-term is linear in some of the U_I.
Restricting to the case of $\gamma_m = 0$ and choosing the simplest example of w :

$$w = A\,(U_2+U_3)\,(S+U_1-2\beta\gamma)\,e^{-\beta S} = \hat{w}\,e^{-\beta S}, \quad \hat{w} = A\,(U_2+U_3)\,(S+U_1-2\beta\gamma)\,.$$

The requirements concerning the R-symmetry are satisfied provided that
U_1 transforms non trivially under $U(1)_R$

$$U(1)_R: \qquad U_1 \longrightarrow U_1 + i\xi_1\,\theta \quad \text{with} \quad \xi_1 = -\xi_S \quad \text{so that} \;\; \hat{w} \longrightarrow \hat{w}$$

and also,

$$U(1)_R: \;\; \Phi^i \longrightarrow \Phi^i\, e^{i\xi_\phi \theta}\,, \qquad T \longrightarrow T + i\xi_T\,\theta\,, \qquad U_{2,3} \longrightarrow U_{2,3} + i\xi_{2,3}\,\theta$$

$$\text{with} \;: \qquad \xi_W = 3\xi_\Phi = -\beta\xi_S, \qquad \xi_1 = -\xi_S, \qquad \xi_2 = -\xi_3$$

The total potential is a sum of positive semi-definite terms.
The minimization conditions with $V = 0$ simplifies enormously:

$$i) \;\; V_{\Phi^i} = 0 \;\; \rightarrow \qquad W_\Phi = [\,\Phi\,]^3_{\Phi^i} = 0 \qquad \longrightarrow \qquad W = \hat{w}\, e^{-\beta S}$$

$$ii) \;\; V_{U_2} = 0 \;\; \rightarrow \;\; W - (U_2 + \bar{U}_2) W_{U_2} = (-\bar{U}_2 + U_3)\hat{w} e^{-\beta S} = 0 \qquad \longrightarrow \qquad \bar{U}_2 = U_3$$

$$iii) \; V_{U_3} = 0 \;\; \rightarrow \;\; W - (U_3 + \bar{U}_3) W_{U_2} = (-\bar{U}_3 + U_2)\hat{w} e^{-\beta S} = 0 \qquad \longrightarrow \qquad U_2 = \bar{U}_3$$

$$iv)\ V_{U_1}=0 \ \rightarrow\ W-(U_1+\bar{U}_1)W_{U_1}=0 \quad \longrightarrow \quad (S+U_1-2\beta\gamma)=(U_1+\bar{U}_1)$$

$$v)\ V_S=0\ \rightarrow\ W-(S+\bar{S})W_S=0 \quad \longrightarrow$$
$$(\bar{S}+U_1-2\beta\gamma)=(S+\bar{S})-\beta(S+\bar{S})(S+U_1-2\beta\gamma)$$

The above equations have to be taken together with the constrain $V_D=0$ which can be seen as $V_{z^I}=0$ condition for all fields $z^I=(T,\Phi,U_I,S)$,

$$vi)\ V_D=0\ \rightarrow\ \frac{1}{2}g_R^2\left(\frac{3\xi_\Phi|\Phi^i|^2}{Y}-\frac{3\xi_T}{Y}-\frac{\xi_S}{(S+\bar{S})}-\sum_I\frac{\xi_I}{(U_I+\bar{U}_I)}-\beta\xi_S\right)^2=0\,.$$

Combining iv with v

$$(S+\bar{S})-(U_1+\bar{U}_1)-\beta\,(S+\bar{S})(U_1+\bar{U}_1)=0\ \rightarrow\ \left(\frac{-\xi_S}{(S+\bar{S})}+\frac{\xi_S}{(U_1+\bar{U}_1)}-\beta\xi_S\right)=0$$

Combining ii with iii

$$(U_2+\bar U_2)-(U_3+\bar U_3)=0 \qquad\qquad \rightarrow \quad \left(\frac{-\xi_2}{(U_2+\bar U_2)}+\frac{\xi_2}{(U_3+\bar U_3)}\right)=0$$

Inserting to vi we observe that $V_D=0$ implies that in the global vacuum $|\Phi^i|^2$ is fixed while T remains flat direction.

$$|\Phi^i|^2=\tilde\xi; \qquad \tilde\xi=\xi_T/\xi_\Phi \quad \text{and} \quad T \text{ arbitrary.}$$

• It is important to stress here that this property, although derived in a particular example, turns out to be generic for all models with R-symmetric w of the type $w=\hat w\, e^{-\beta S}$!

In the example under consideration
(setting $S=s+ia$, $U_1=u+ib$ and $2\beta\gamma$ real) :

$$u=s-2\beta\gamma, \quad a=-b \qquad \text{and} \qquad (s-u)=2\beta\, su, \quad \rightarrow \qquad su=\gamma$$

$$s=\sqrt{(\beta^2\gamma^2+\gamma)}+\beta\gamma\ , \qquad u=\sqrt{(\beta^2\gamma^2+\gamma)}-\beta\gamma.$$

,

Both the real parts of S and U_1 are fixed while
their imaginary values are restricted so that $a+b=0$.
The gravitino mass at the minimun $V=0$ reads :

$$m_{3/2}^2 = \frac{|\hat{w}|^2\, e^{-2\beta\, s}}{4\gamma\, Y^3} \neq 0\,, \qquad Y = (T + \bar{T} - |\Phi^i|^2) = 2t - \tilde{\xi}$$

with $$\hat{w} = 2\, A\, (s - 2\beta\gamma) = 2\, A\, u$$

There is no-dependence on U_2, U_3 . The effective theory is just the $SU(1,1+n)$ no scale model with a constant term in W proportional to Λ^3 giving rise to naturally small susy breaking scale:

$$m_{3/2}^2 = \frac{|\Lambda|^6}{(T' + \bar{T}')^3}\,, \qquad T' + \bar{T}' = T + \bar{T} - \tilde{\xi}\,.$$

4. Inflationary transitions

The off-shell inflationary properties of the potential $V = V_F + V_D$ depends crucially on the initial values of the fields T, Φ^I, S and U_I.

Assuming $1/g_R^2 = S + \bar{S}$, V_F and V_D scale like:

$$V_F \sim \frac{[\Phi_i^4]}{Y^2\, s\, u\, u_2\, u_3} + \frac{e^{-2\beta s}}{Y^3\, su}, \qquad V_D \sim \frac{1}{s}$$

- For large $Y = e^{\alpha\phi}$ the potential is exponentially dominated by V_D.
- The V_D dominance is even more efficient in the large S-limit.

It follows that depending on the initial off-shell values of T, Φ^I, S and U_I the inflationary transitions can be either:

- Fast transition giving rise to "multi component chaotic inflationary models"

OR

- Exponentially slow transition giving rise to slow-roll inflation from
de Sitter $\longrightarrow$ flat space.

Setting for instance S and $U_{2,3}$ and to their minima with $s = g_R^{-2}$ and $U_2 = \bar{U}_3$:

$$V \sim \frac{1}{2} g_R^2 (3\xi_\Phi)^2 \left(|\hat{\Phi}^i|^2 - \tilde{\xi} e^{-\alpha\phi} - \frac{1}{\beta} e^{-\sqrt{2}\phi_u} + \hat{\xi} \right)^2 \longrightarrow \mu^2 = \frac{1}{2} g_R^2 (3\xi_\Phi)^2 \, (\hat{\xi})^2$$

$$\hat{\xi} \equiv \left(1 + \frac{1}{2\beta} g_R^2 \right) \qquad \text{and} \qquad \phi_u = \frac{1}{\sqrt{2}} \log(U_1 + \bar{U}_1)$$

5. Conclusions

– In this talk I first review the minimal no-scale model based on

$$\mathcal{M}_n = SU(1,1+n)/U(1) \times SU(1+n)$$

which is classically equivalent to the $\mathcal{R}^2$ coupled in a scale invariant matter.

– The stringy extension consists to include stringy moduli fields, namely, the dilaton and the complex stucture moduli, S and U_I.
The scalar manifold is extended by the factor:

$$\mathcal{K}_n = \frac{SU(1,1)_S}{U(1)} \times \frac{SU(1,1)_{U_1}}{U(1)} \times \frac{SU(1,1)_{U_2}}{U(1)} \times \frac{U(1,1)_{U_3}}{U(1)}$$

– The presence of the additional moduli turns out to be crucial since both susy and R-symmetry can be spontaneously broken in the global Minowski vacuum after the inflationary transition.

– The R-symmetry acts non trivially on the non perturbative superpotential which involves flux terms and the non-perturbative exponential factor: $\hat{w}(S, U_I)\, e^{-\beta S}$.

– The semi-classical R-symmetry becomes non-anomalous at the quantum level if the Peccei-Quinn axion shift symmetries associated to the T , S , U_I moduli-fields are gauged consistently.

– For large initial values of $T + \bar{T}$ and $U + \bar{U}$ the extended theory can describe successfully inflationary transition from de Sitter to the flat Minkowski space.

– The susy is broken in both phases:

- In the de Sitter by a Fayet-Iliopoulos D-term

- In the flat global vacuum by the non-perturbative susy breaking with $w \neq 0$ which breaks spontaneously both susy and R-symmetry

– The effective low energy theory (after fixing S, U_I and $|\Phi^i|^2$) is identical to M_n no scale model with naturally small gravitino mass,

$$m_{3/2}^2 = \frac{|\Lambda|^6}{(T' + \bar{T}')^3}\,, \qquad\qquad 2T' = 2T - \tilde{\xi}$$

T' and $m_{3/2}^2$ are undetermined at the (semi-)classical level.

Thank you for your attention

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