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Thermal de Sitter Solutions from String Effective Theories

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[In effective supergravities from Strings]

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[In Suprerstrings with broken SUSY]

1. Introduction

• **Superstrings and M-theory compactifications can give 4d vacua with exact or spontaneously broken supersymmetries.**

• **The most interesting phenomenologically are those where**

$$N = 8,4 \longrightarrow N = 1 \longrightarrow N = 0$$

• **The underlying D = 10 theories encode** $N \geq 4$ **constrained structure which can be used to obtain useful information on the effective** $N = 1$ **supergravity.**

• **The 4d** $N = 1$ **theories, typically include moduli fields whose vacuum expectation values are undetermined.**

Some of these moduli come from the : dilaton field Φ**, internal metric** G_{IJ}**, and** p**-form fields** F^p**.**

Generating a potential for the moduli is essential in order

(i) to reduce the number of massless scalars

(ii) to induce supersymmetry breaking

(iii) to determine the space-time background geometry

In $N \geq 4$ supergravities, the only known available tool for generating a potential is the "gauging" which is induced by:

i) Non trivial Fluxes and Branes

ii) Non-perturbative corrections

iii) Perturbative Radiative Corrections

iv) Thermal Corrections

2. $N = 4$ Gauging $\leftrightarrow$ $N = 1$ Superpotential

"Gauging" $\longrightarrow$ **We introduce in the theory a gauge group G acting on the vector fields in the gravitational and/or vector super-multiplets.**

The important fact is: The kinetic terms of the fields in the gauged theory, remain the same as in the ungauged theory.

When translated into the language of the "daughter $N = 1$" obtained from the "mother $N \geq 4$" by a consistent orbifold (CY) truncation, the gauging gives non-trivial modifications to the structure of the superpotential W, whereas the Kähler potential K remains the same.

To be more precise consider the case of $N = 4$ theory constructed either on:

- **Heterotic on T^6**
- **Type IIA or IIB on $K3 \times T^2$**
- **Type IIA, IIB on orientifolds**
- **. . .**

Independently of our starting point, the $N = 4$ scalar manifold is identical for all cases.

$$M = \left(\frac{SU(1,1)}{U(1)}\right)_S \times \left(\frac{SO(6,6+n)}{SO(6) \times SO(6+n)}\right)_{T_A,U_A,Z_I}$$

After $Z^2 \times Z^2$ orbifold (CY) projections

$N = 4 \rightarrow N = 1$ **and** $M \rightarrow K$

$$K = \left(\frac{SU(1,1)}{U(1)}\right)_{S} \times \left(\frac{SO(2,2+n_1)}{SO(2)\times SO(2+n_1)}\right)_{T_1,U_1,Z_1^I} \times$$

$$\left(\frac{SO(2,2+n_2)}{SO(2)\times SO(2+n_2)}\right) \times \left(\frac{SO(2,2+n_3)}{SO(2)\times SO(2+n_3)}\right)$$

$$K = -\log\left(S+\bar{S}\right)$$

$$-\log\left((T_1+\bar{T}_1)(U_1+\bar{U}_1)-(Z_1+\bar{Z}_1)^2\right)$$

$$-\log\left((T_2+\bar{T}_2)(U_2+\bar{U}_2)-(Z_2+\bar{Z}_2)^2\right)$$

$$-\log\left((T_3+\bar{T}_3)(U_3+\bar{U}_3)-(Z_3+\bar{Z}_3)^2\right).$$

The above choice of parametrization is a solution to the $N = 4$ constraints after $Z^2 \times Z^2$ orbifold projections $N = 4 \rightarrow N = 1$

S-manifold

$$|\phi_0|^2 - |\phi_1|^2 = \frac{1}{2} \qquad \longrightarrow$$

$$\phi_0 - \phi_1 = \frac{1}{(S+\bar{S})^{1/2}}, \qquad \phi_0 + \phi_1 = \frac{S}{(S+\bar{S})^{1/2}}$$

T_A, U_A, Z^I_A-manifolds

$$|\sigma^1_A|^2 + |\sigma^2_A|^2 - |\rho^1_A|^2 - |\rho^2_A|^2 - |\Phi^I_A|^2 = \frac{1}{2}$$

$$(\sigma^1_A)^2 + (\sigma^2_A)^2 - (\rho^1_A)^2 - (\rho^2_A)^2 - (\Phi^I_A)^2 = 0$$

$$\sigma^1_A = \frac{1 + T_A U_A - (Z^I_A)^2}{2Y_A^{1/2}}, \qquad \sigma^2_A = i\frac{T_A + U_A}{2Y_A^{1/2}}$$

$$\rho^1_A = \frac{1 - T_A U_A - (Z^I_A)^2}{2Y_A^{1/2}}, \qquad \rho^2_A = i\frac{T_A - U_A}{2Y_A^{1/2}}$$

$$\Phi^I_A = \frac{iZ^I_A}{2Y_A^{1/2}}, \qquad K_A = -log\ Y_A$$

The superpotential of the $N=1$ supergravity is determined by the gravitino mass terms in $N=4$ after the $Z^2 \times Z^2$ orbifold projections.

$$e^{K/2}W = (\phi_0 - \phi_1)\ f_{IJK}\ \Phi_1^I \Phi_2^J \Phi_3^K$$
$$+(\phi_0 + \phi_1)\ \bar{f}_{IJK}\ \Phi_1^I \Phi_2^J \Phi_3^K$$

$$\Phi_A^I = \left\{\ \sigma_A^1,\ \sigma_A^2;\ \rho_A^1,\ \rho_A^2,\ \Phi_A^I\right\}$$

f_{IJK} **and** $\bar{f}_{IJK}$ **are the gauge structure constants of the $N=4$ "mother" theory.**

In the heterotic, the contributions from f_{IJK} give rise to a perturbative "electric gauging". The contributions from $\bar{f}_{IJK}$ provide the non-perturbative "magnetic gauging".

In general, the breaking of SUSY requires a gauging with non-zero f_{IJK} involving the fields

$\sigma^1_A,\ \sigma^2_A;\ \rho^1_A,\ \rho^2_A \longrightarrow$ **gauging involving the** $N = 4$ **graviphotons** $\longrightarrow$ **gauging of the R-symmetry.**

In string and M-theory, f_{IJK} **and** $\bar{f}_{IJK}$ **are generated by non-zero electric and magnetic fluxes; RR and fundamental** p**-form fields:**

- **3-form fluxes** H^3**, in the NS-sector of heterotic, type IIA and type IIB**

- F^p**, p-form fluxes, in M-theory and in the RR sector of type IIA and type IIB**

- F^2 **2-form fluxes, in heterotic (**$E_8 \times E_8$ **or** $SO(32)$**) as well as in type I**

- ω^3 spin connection, 3-form geometrical fluxes, in all strings and M-theory.

Many special cases have been studied in the literature .

- H^3 in heterotic

Derendinger, Ibanez, Nilles, 85, 86;
Dine, Rohm, Seiberg, Witten, 85;
Strominger, 86; Rohm, Witten, 86;
Derendinger, Kounnas, Petropoulos, 06;
. . .

- ω^3, H^3, F^2, exact string solution via freely acting orbifold.

Generalization of the Scherk–Schwarz gauging to superstring theory.

Rohm, 84; Kounnas, Porrati, 88
Ferrara, Kounnas, Porrati, Zwirner, 89
Kounnas, Rostand, 90
Kiritsis, Kounnas, Petropoulos, Rizos, 99

Antoniadis, Dudas, Sagnotti, 99
Antoniadis, Derendinger, Kounnas, 99
Derendinger, Kounnas, Petropoulos, Zwirner, 04
Derendinger, Kounnas, Petropoulos, 05,06
...

- **Simultaneous presence of NS, RR H^3 and F^3.**

Frey, Polchinski, 02
Giddings, Kachru, Polchinski, 02
Kachru, Schulz, Trivedi, 03
Kachru, Schulz, Tripathy, Trivedi, 03
Derendinger, Kounnas, Petropoulos, Zwirner, 2004
...

3. Combined Fluxes, Gauging and Moduli Stabilization.

• **Flat gaugings, no-scale models;** **stabilization of the four moduli out of the seven main moduli.**

(i) Scherk–Schwarz, perturbative, ω^3-fluxes.

$$W = a \; (\; T_2 U_1 + T_1 U_2 \;)$$

$V \geq 0$**, flat in** S, T_3, U_3 **directions**

$$m_{3/2}^2 = \frac{|a|^2}{s t_3 u_3}$$

(ii) Scherk–Schwarz non-perturbative, ω^3, F^2**,** H^3, F^6 **-fluxes in type IIA**

$$W = a(\; S T_1 + T_2 T_3 \;) + ib(\; S + T_1 T_2 T_3 \;)$$

$$m_{3/2}^2 = \frac{|2a|^2 + |2b|^2}{u_1 u_2 u_3}$$

(iii) F^3, H^3-fluxes in type IIB

$$W = a(\ SU_1 + U_2U_3\ ) + ib(\ S + U_1U_2U_3\ )$$

$$m_{3/2}^2 = \frac{|2a|^2 + |2b|^2}{t_1t_2t_3}$$

The stabilization of U_I moduli is a generic situation in type IIB $\longrightarrow$ No-scale models.

The vanishing of the potential in the flat direction of the T_I-moduli can be modified only by (non-) perturbative and/or thermal corrections.

(iv) $SO(3) \times SO(1,2)$, $E_3^c \times E_3^{nc}$ gaugings
ω^3, F^2, H^3, F^6 -fluxes

$$W = a(\ ST_1+ST_2+ST_3\ )+a(\ T_1T_2+T_2T_3+T_3T_1\ )$$
$$+i3b(\ S + T_1T_2T_3\ )$$

$$m_{3/2}^2 = \frac{|6a|^2 + |6b|^2}{u_1u_2u_3}$$

- **Non-compact gaugings,** $SO(1,2)$**,** E_3^{nc} $V > 0$ **, runaway solutions.**

(i) No-modulus stabilization

$$W = F_0$$

$$m_{3/2}^2 = \frac{|F_0|^2}{st_1t_2t_3u_1u_2u_3}, \qquad V = 4m_{3/2}^2$$

(ii) Two-moduli stabilization

$$W = F_0 + F_2 \ T_1T_2,$$

$$m_{3/2}^2 = \frac{|2F_0|^2}{st_3u_1u_2u_3}, \qquad V = 2m_{3/2}^2$$

(iii) Three-moduli stabilization E_3^{nc} **gauging** F_0, F_2, F_4, F_6 **-fluxes**

$$W = a(\ 1 + T_1T_2 + T_2T_3 + T_3T_1\ ) \\ +ib(\ T_1 + T_2 + T_3 + T_1T_2T_3\ )$$

$$m_{3/2}^2 = \frac{|4a|^2 + |4b|^2}{s u_1 u_2 u_3}, \qquad V = m_{3/2}^2$$

- **Compact gaugings $SU(2)$, E_3^c**

(i) Stabilization of six-moduli,
NS_5 brane solution + linear dilaton

$$W = \omega_3(\ T_1U_1 + T_2U_2 + T_3U_3\ ) - F_0$$

$$V = -2m_{3/2}^2, \qquad m_{3/2}^2 = \frac{|2F_0|^2}{s}$$

(ii) Stabilization of all moduli, AdS$_4$-solution

$$W = iB[\ 2S + 5T_1T_2T_3 + 2(U_1 + U_2 + U_3) \\ -3(T_1 + T_2 + T_3)\ ]+ \\ A[\ 2S(T_1 + T_2 + T_3) - (T_1T_2 + T_2T_3 + T_3T_1) \\ +6(T_1U_1 + T_2U_2 + T_3U_3) - 9\ ]$$

$$S = T_A = U_A = 1, \quad m_{3/2}^2 = \frac{25|B|^2}{128}, \quad V = -3m_{3/2}^2$$

***** Fluxes/Branes $\leftrightarrow$ Gaugings*****

$\longrightarrow$

Explicit Superstring Constructions with different interesting features:

- **Stabilization of four moduli, $V \geq 0$: No-scale models.**

- **Stabilization of less than four moduli, $V > 0$: de Sitter like, runaway solutions with possible cosmological interest.**

- **Models based on compact "gaugings", $V < 0$: Domain-Wall Solutions, Five-brane solutions with non trivial Dilaton or else.**

- **Models with all seven moduli stabilized and which admit a supersymmetric AdS_4 vacuum.**

4. Cosmological-Inflationary Solutions

• I would like to show here, that several
String Effective No-scale Supergravities
admit
de Sitter-like Backgrounds
once the
i) Radiative Corrections,
and
ii) Temperature Corrections
are both included in the
effective potential of the theory.

the cosmological term Λ
and
the curvature term $\frac{k}{a^2}$

will be both generated, in an elegant way, in the effective no-scale supergravities.

• **Due to the presence of non trivial fluxes the supersymmetry is generally broken.**

In the case of no-scale models however, the potential is semi-positive definite
$V \geq 0$
with at least one flat direction Φ **in the scalar field space.**

The field Φ **is the supersymmetric partner of the Goldstino and defines the gravitino mass term** $m_{3/2}$**.**

$$m_{3/2}^2 \; = \; |W|^2 e^K = \; e^{2\alpha\Phi}, \qquad \text{with} \;\; V \equiv 0.$$

We are working in unites of the gravitation scale $M = 2.4 \times 10^{18} \; GeV$**.**

At the quantum level, the potential gets corrections because of the supersymmetry breaking.

The masses of the bosonic and fermionc degrees of freedom differ by an amount proportional to $m_{3/2}$.

What we know about the quantum structure of the corrected potential ?

1. Assuming an ultraviolet cutoff Λ, the terms proportional to Λ^4 are proportional to the difference of the bosonic and fermionic degrees of freedom:
$n_B - n_F \equiv 0$

2. The term proportional to Λ^2 is always proportional to $m_{3/2}^2$.
The numerical coefficient depends only on the geometry of the Kälher potential and the gauge kinetic function f of the effective supergravity theory.
At the string level the exact form of this term is calculable!

Here, for our purpose, it is sufficient to absorbed this coefficient to the definition of Λ^2

$$\Delta V = \Lambda^2 \, m_{3/2}^2$$

Λ^2 can be positive or negative depending of the model under consideration.

3. The remaining corrections scale logarithmically with respect to the supersymmetry breaking scale, $m_{3/2}$.

$$\Delta V = \alpha(m_{3/2}^4 + A log \, \frac{m_{3/2}}{\mu} \,)$$

$$+ \quad Q_2^2 \, m_{3/2}^2 (q_2 + log \, \frac{m_{3/2}}{\mu})$$

$$+ Q_0^4 (q_0 + log \, \frac{m_{3/2}}{\mu} \,) + ...$$

Where $Q_{2,4}$ denotes supersymmetric mass terms and μ the renormalization scale.

The dots denote threshold terms that their role will be specified later.

• The other corrections we need are those which are due to the thermal fluctuation of the fields.

Thus, it is necessary to evaluate the effective potential at finite temperature T.

The physical scale T naturally sets the strength of the thermal fluctuations. Utilizing the appropriate renormalization group arguments, we must identified the scale

$$\mu \equiv T.$$

This identification resumes in an efficient way the logarithmic corrections.

The threshold corrections are organized in terms of O($\frac{m_i}{T}$) -expansion.
m_i **are the physical masses defined at the temperature scale** T**.**
Summarizing, the effective potential at finite temperature takes the form:

$$V_T = -p(m_i, T) + V(m_{3/2}, \Lambda, \mu = T)$$

where $p(m_i, T)$ **is the pressure density:**

$$p(m_i, T) = \sum_i p_i(m_i, T), \qquad \rho(m_i, T) = \sum_i \rho_i(m_i, T)$$

$$T \frac{\partial}{\partial T} p_i(m_i, T) = p_i(m_i, T) + \rho_i(m_i, T),$$

$$-m_i \frac{\partial}{\partial m_i} p_i(m_i, T) = \rho_i(m_i, T) - 3p_i(m_i, T)$$

$$p_i(m_i, T) = T^4 \ f_p^{B,F}(\frac{m_i}{T}),$$

$$\rho_i(m_i, T) = T^4 \ f_\rho^{B,F}(\frac{m_i}{T})$$

$$\rho_i(m_i, T) - 3p_i(m_i, T) = m_i^2 \ T^2 \ \Delta^{B,F}(\frac{m_i}{T})$$

5. Gravitational Equations

The energy density equation

$$3H^2 = \rho + \frac{1}{2}\dot{\Phi}^2 + V - \frac{3k}{a^2}$$

The conservations equation

$$\frac{d}{dt}\left(\rho + \frac{1}{2}\dot{\Phi}^2 + V\right) + 3H\left(\rho + p + \dot{\Phi}^2\right) = 0$$

The scalar-modulus equation

$$\ddot{\Phi} + 3H\dot{\Phi} + \frac{\partial}{\partial \Phi}(V - p) = 0$$

The above equations are the independent ones. Some other useful combinations are:

The trace equation

$$6\ \dot{H} + 12\ H^2 = (\rho - 3p + 4V - \dot{\Phi}^2) - \frac{6k}{a^2}$$

The entropy equation

$$\dot{H} = -\frac{1}{2}(\rho + p + \dot{\Phi}^2) + \frac{k}{a^2}$$

We find convenient, for the derivation of the solution, to organize the potential V and the thermal functions ρ, p terms according to their scaling properties with respect to $m_{3/2}$ and T.

$$V_4 = A \; m_{3/2}^4 \; log \; \left(\frac{m_{3/2}}{c_4 T}\right) \; ,$$

$$V_4' \equiv \frac{\partial}{\partial \Phi} V_4 \; = \; 4\alpha \; V_4 + \alpha \; A m_{3/2}^4$$

$$V_2 = \Lambda^2 m_{3/2}^2 + Q_2^2 \; m_{3/2}^2 \; log \; \left(\frac{m_{3/2}}{c_2 T}\right) \; ,$$

$$V_2' \; = \; 2\alpha \; V_2 + \alpha \; Q_2^2 \; m_{3/2}^2$$

$$V_0 = Q_0^4 \; log \; \left(\frac{m_{3/2}}{c_0 T}\right) \; ,$$

$$V_0' \; = \; \alpha \; Q_0^4$$

The thermal potential $-p(m_i, T)$ can be separated naturally in three classes:

i) The first class consists of all massless states: $\rho - 3p = 0$, $p' = 0$

ii) The second class consists of states with (supersymmetric masses) ***independent*** **of $m_{3/2}$.**

$\longrightarrow$

• If the temperature scale T is above their masses, their contribution to p and ρ is proportional to $\sim T^4$
their contribution to $\rho - 3p \sim m_i^2 \, T^2$

• If the temperature scale T is below their masses, then their contribution is exponentially suppressed $\sim e^{-m/T}$ and thus they are effectively decouple from the thermal system.

iii) The third class consists of states where their masses are *mainly proportional to* $m_{3/2}$.

$$\longrightarrow \qquad p' = -\alpha(\rho - 3p)_\Phi .$$

The Φ-modulus equation:

$$\ddot{\Phi} + 3H\dot{\Phi} + \alpha\left(4V_4 + (\rho - 3p)_\Phi\right)$$

$$+\alpha\left(Am_{3/2}^4 + 2V_2 + Q_2^2 m_{3/2}^2 + Q_0^4\right) = 0$$

Using *the trace equation:*

$$6\ \dot{H} + 12\ H^2 = (\rho - 3p + 4V - \dot{\Phi}^2) - \frac{6k}{a^2}$$

the modulus equation becomes:

$$0 = \ddot{\Phi} + 3H\dot{\Phi} + \alpha\dot{\Phi}^2 + \alpha\left(6\ \dot{H} + 12\ H^2 + \frac{6k}{a^2}\right) +$$

$$\alpha\left(Am_{3/2}^4 - 2V_2 + Q_2^2 m_{3/2}^2 - 4V_0 + Q_0^4 - (\rho - 3p)_r\right)$$

Critical Solution

$$\alpha\dot{\Phi} = -H, \quad T = \frac{x}{a} \quad \text{and} \quad e^{\alpha\Phi} = m_{3/2} = \eta T$$

Assuming for the moment the absence of the logarithmic dependence in the V_4-term, (equivalent to the absence of $Am^4_{3/2}$) and using

$$(\rho - 3p)_r = M_2^2 \; T^2$$

M_2^2 is independent of Φ, SUSY mass term.

- **The existence of the critical solution:**

$$0 = (\; 6\dot{H} + 12 \;\; H^2) + \frac{6\alpha^2}{6\alpha^2 - 1}\left[Q_0^4(1 - 4log\frac{\eta}{c_0})\right]$$

$$+\frac{6\alpha^2}{6\alpha^2 - 1}\left[\frac{6k}{a^2} + \frac{1}{a^2}\left(\; x^2\eta^2 Q_2^2(1 - 2log\frac{\eta}{c_2}) - x^2 M_2^2\right)\right]$$

- **the consistency of the energy equation,**
- **the conservations equation, and**
- **the modulus equation**

imply that:

$$\Lambda_{cosm} = \frac{3}{2}\ \alpha^2\ Q_0^4,$$

$$k = \alpha^2\ x^2 \left[-\eta^2 Q_2^2 (1 - 2log\frac{\eta}{c_2}) + M_2^2 \right]$$

The lesson of this exercise is that :

- **the cosmological constant scale,** Λ_{cosm}
- **the curvature scale,** k

both are generated by :

"the scaling violating terms"

of the thermal effective potential.

- **In the absence of the scaling violating terms (or when these terms are negligible) the modulus Φ couples to the total trace of the energy momentum tensor.**

• **This special case was studied in 1986 by I. Antoniadis and C. Kounnas. They found that the the critical trajectory is the only stable solution under any field fluctuation and that this trajectory is an attractor at late times to a radiation evolving universe with :**

$$a^2 \sim t, \quad T^2 \sim m_{3/2}^2 \sim 1/t, \quad V \sim 1/t^2.$$

6. The time trajectory of a, T and Φ

The energy density equation

$$3H^2 = \rho + \frac{1}{2}\dot{\Phi}^2 - V - \frac{3k}{a^2}$$

in the background of the critical solution becomes:

$$\left(\frac{6\alpha^2 - 1}{6\alpha^2}\right) 3H^2 = \rho + V - \frac{3k}{a^2}$$

OR

$$\left(\frac{6\alpha^2 - 1}{6\alpha^2}\right) \; 3H^2 = \frac{C_\rho}{a^4} + \Lambda_{cosm} - \frac{3k'}{a^2}$$

The dilatation factor in frond of $3H^2$,
can be absorbed by a redefinition of :
$C_\rho, \quad \Lambda_{cosm}$ **and** $k' \longrightarrow$

The time dependence of a is similar to
the radiation-deformed de Sitter Solutions!

$$a_+^2 = \frac{3}{\Lambda}\left(ch^2\left(\sqrt{\frac{\Lambda}{3}}\; t\right) + \epsilon_\rho^2\right)$$

$$a_-^2 = \frac{3}{\Lambda}\left(-sh^2\left(\sqrt{\frac{\Lambda}{3}}\; t\right) + \epsilon_\rho^2\right)$$

H.Firouzjahi, S.Sarangi, S.H.H.Tye, 04 ;
S.Sarangi, S.H.H.Tye, 05;
R. Brustein, S.P. de Alvis, 06;
C. Kounnas, H. Partouche, 07; . . .

It is interesting that the a_- and a_+ are connected by a Gravitational (Φ-Dilatonic) Instanton

$$a_E^2 = \frac{3}{\Lambda}\left(cos^2\left(\sqrt{\frac{\Lambda}{3}}\ \tau\right) + \epsilon_\rho^2 \right)$$

with a transition probability $(a_- \rightarrow a_+)$

$$< \Psi_- || \Psi_+ > = \ P \ \sim e^{\left(\frac{3}{\Lambda}\ -\ \frac{\chi\rho}{\Lambda^2}\right)}$$

- **Where $\Psi_\pm$ is the wave-function of the universe.**

J.B. Hartle, S.W. Hawking, 83;
A. Vilenkin, 82, 83;
A.D. Linde, 84;
H.Firouzjahi, S.Sarangi, S.H.H.Tye, 04;
S.Sarangi, S.H.H.Tye, 05;
R. Brustein, S.P. de Alvis, 06;
C. Kounnas, H. Partouche, 07; . . . . . . ;

• χ_ρ is proportional to the number of the effective marginal and/or thermal degrees of freedom at the temperature scale T_0, defined at the transition point.

C. Kounnas, H. Partouche, 07;

7. String Perspectives and Conclusion

At classical string level it seems that it is difficult to construct exact cosmological string solutions and is even more difficult to obtain de Sitter like inflationary solutions even in lower than four dimensions.

• Consider for instance the euclidian version of dS^3 which is an exact conformal field theory base on $SU(2)_k$ WZW model. However, it does not admit any real-time analytic continuation due to the existence of a non trivial torsion H_{ijk} which becomes imaginary!!

I.Antoniadis, C.Bachas, A.Sagnotti, 90;
P.K Townsend 01;
J. Sonner, P.K Townsend 06;
C.Bachas, C.Kounnas,
D.Orlando, M.Petropoulos 07;

. . .

- **Indeed the only known cosmological solution based on an exact conformal field theory is that of $SL(2,R)/U(1)_{-|k|} \times K$. Its euclidian version is also well defined by the gauged WZW parafermionic T-fold $SU(2)/U(1)_k$.**

C.Kounnas, D.Luest 92;
L.Coralba, M.S.Costa, C.Kounnas 02;
C.Kounnas, N.Toumbas, J.Troost O7;

. . .

- **In all String cosmological models with a well define euclidian version, (like for instance the $SL(2,R)/U(1)_{-|k|} \times K$),**

the super-string analog of:
"the Stringy wave function of the universe"
can be unambiguously defined.

Furthermore the transition probabilities can be easily calculated at the string level.
C.Kounnas, N.Toumbas, J.Troost O7.

Concluding Remark

- **Our proposal however goes even beyond the scope of the above statement; namely** ***towards to :***

The plausible existence of cosmological super-string solutions (Inflationary or not)

which are generated dynamically
at the quantum sting level
from a flat classical space-time and
spontaneously broken supersymmetry
(no-scale radiative-induced cosmology).