Tianjin University

4D de Sitter from 6D gauged supergravity with Green-Schwarz counterterm

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Introduction

Nobel Prize in Physics 2011

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Saul Perlmutter

Prize share: 1/2

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Brian P. Schmidt

Prize share: 1/4

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Adam G. Riess

Prize share: 1/4

The Nobel Prize in Physics 2011 was divided, one half awarded to Saul Perlmutter, the other half jointly to Brian P. Schmidt and Adam G. Riess "for the discovery of the accelerating expansion of the Universe through observations of distant supernovae"

The Astronomical Journal, 116:1009–1038, 1998 September

OBSERVATIONAL EVIDENCE FROM SUPERNOVAE FOR AN ACCELERATING UNIVERSE AND A COSMOLOGICAL CONSTANT

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Received 1998 March 13; revised 1998 May 6

The Astrophysical Journal, 517:565–586, 1999 June 1

MEASUREMENTS OF Ω AND Λ FROM 42 HIGH-REDSHIFT SUPERNOVAE

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Introduction

Our universe is likely to be asymptotic to de Sitter space

$$S = \int d^4x\sqrt{-g}(R - 2\Lambda + \text{matter})$$

which is the standard model of cosmology (ΛCDM). However, as an effective field theory, this is extremaly unnatural, $E_\Lambda/E_{EW} \sim 10^{-12}$, $E_{EW}/M_P \sim 10^{-17}$.

What is the origin of Λ, why it is so small?

- Quantum vacuum energy: this requires fine-tuning. Quantum field theory predicts a vacuum energy which is 10^{60} larger than the observed value.
- Classical: How quantum effects are suppressed, is there any hidden symmetry such as supersymmetry?
- Dynamical rather than being constant.

String theory is the leading candidate of quantum gravity

Construction of classical de Sitter vacuum in string/M theory is difficult

- Maldacena-Nunez no-go theorem: D-dimensional gravity ($D > 2$) compactified to d dimensions has no warped de Sitter solutions,

$$ds_D^2 = \Omega^2(y) ds_d^2 + g_{mn} dy^m dy^n, \quad ds_d^2 = L_{dS}^2 \frac{-dt^2 + dx_i dx_i}{t^2}$$

IF:

- The gravity theory has no higher derivative corrections.
- The potential is non-positive, $V \leq 0$ (massive IIA is treated separately).
- The massless fields have positive kinetic terms.
- $G_{N,d}$ is non-zero.

The Refined de Sitter Conjecture [Ooguri, Palti, Shiu and Vafa, '18]

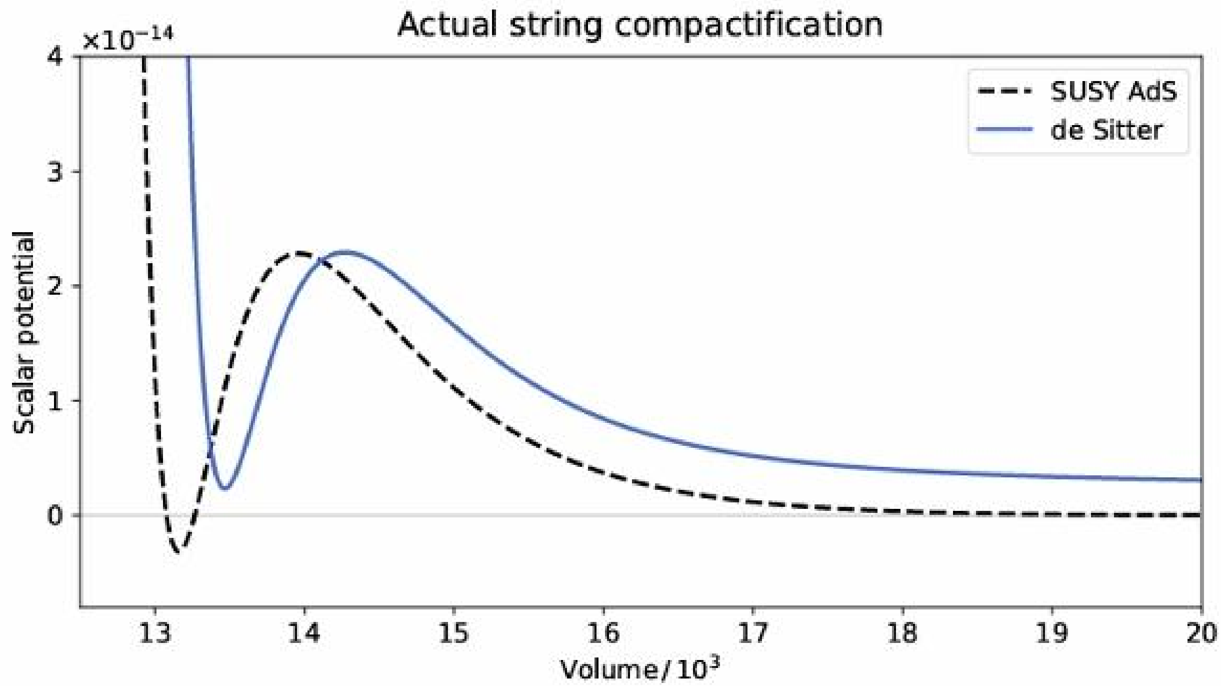
- $$|\nabla V| \geq \frac{c}{M_p} V \ ,$$

- $$\min(\nabla_i \nabla_j V \leq -\frac{c'}{M_p^2} V) \quad (*)$$

where c, c' are order one parameters. The left-hand side of $(*)$ is the minimum eigenvalue of the Hessian $\nabla_i \nabla_j V$.

By adding non-perturbative /or higher derivative terms, we come to KKLT and LVS models, neither of which is fully satisfactory. See [F. Quevedo et. al, '23] for a recent review.

- KKLT scenario



No scale separation between the external de Sitter and internal Calabi-Yau manifold, i.e. $L_{dS} \gg L_{\rm int}$. [Lust, Vafa, Wiesner and Xu, '22; Bedroya and Steinhardt,'25]. **KKLT is very much a prescription, rather than an explicit ten-dimensional solution.**

- In Large Volume Scenario (LVS), scale separation is naturally achieved with the potential

$$V_{\rm LVS} \propto \frac{a\sqrt{\tau_4}e^{-4\pi\tau_4}}{\mathcal{V}} - \frac{b\tau_4 e^{-2\pi\tau_4}}{\mathcal{V}^2} + \boxed{\frac{c}{\mathcal{V}^3}}$$

where a, b, $c > 0$ and the last term comes from compactifying the partial leading α' correction [K. Becker, M. Becker, M. Haack and J. Louis, '02]

$$S = -\frac{1}{2\kappa_{10}^2}\int d^{10}x\sqrt{-g^{(10)}}e^{-2\phi}(R + 4(\partial\phi)^2 + \alpha'^3 c_1 J_0)\ ,$$

where $c_1 = \frac{\zeta(3)}{3\cdot 2^{11}}$. The higher order interaction is defined as

$$J_0 = t^{M_1N_1\ldots M_4N_4}t_{M_1'N_1'\ldots M_4'N_4'}R^{M_1'N_1'}{}_{M_1N_1}\ldots R^{M_4'N_4'}{}_{M_4N_4} + \frac{1}{4}E_8\ ,$$

However, J_0 is merely the leading term in the tree level 8-derivative superinvariant. There are many other terms depending on fluxes.

The effective action at order α'^3 in the NS-NS sector takes the schematic form ($\phi = 0$) [Garousi, '20]

$$S^{IIB}_{\mathrm{NS-NS,tree}} = c\alpha'^3 \int dx^{10} e \Big[[R^4]_2 + [R^3 H^2]_{22} + [R^2 (\nabla H)^2]_{22} + [R^2 H^4]_7 + [R(\nabla H)^2 H^2]_{106} + [RH^6]_1 + [(\nabla H)^4]_{12} + [(\nabla H)^2 H^4]_{77} + [H^8]_2 \Big] ,$$

There are a lot more terms including RR fluxes.

- Most of the work focus on stability of lowest-lying Kaluza-Klein modes. However, in non-supersymmetric setup, even if the lowest-lying Kaluza-Klein modes are stable, higher Kaluza-Klein modes can be tachyonic [Malek, Nicolai, Samtleben, '20] .

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Tachyonic Kaluza-Klein modes and the AdS swampland conjecture

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ABSTRACT: We compute the Kaluza-Klein spectrum of the non-supersymmetric SO(3) × SO(3)-invariant AdS_4 vacuum of 11-dimensional supergravity, whose lowest-lying Kaluza-Klein modes belong to a consistent truncation to 4-dimensional $\mathcal{N} = 8$ supergravity anc are stable. We show that, nonetheless, the higher Kaluza-Klein modes become tachyonic so that this non-supersymmetric AdS_4 vacuum is perturbatively unstable within 11-dimensional supergravity. This represents the first example of unstable higher Kaluza-Klein

Level 2	
m^2L^2	Irrep
-3.117	$(\mathbf{1}, \mathbf{1})$
-2.821	$(\mathbf{1}, \mathbf{1})$
-2.532	$(\mathbf{3}, \mathbf{3})$
-2.448	$(\mathbf{3}, \mathbf{3})$
-2.361	$(\mathbf{5}, \mathbf{5})$

Level 3	
m^2L^2	Irrep
-3.146	$(\mathbf{2}, \mathbf{2})$
-2.892	$(\mathbf{2}, \mathbf{2})$
-2.741	$(\mathbf{4}, \mathbf{4})$
-2.446	$(\mathbf{4}, \mathbf{4})$
-2.627	$(\mathbf{6}, \mathbf{6})$

Level 4	
m^2L^2	Irrep
-2.950	$(\mathbf{1}, \mathbf{1})$
-2.922	$(\mathbf{1}, \mathbf{1})$
-3.114	$(\mathbf{3}, \mathbf{3})$
-2.801	$(\mathbf{3}, \mathbf{3})$
-2.876	$(\mathbf{5}, \mathbf{5})$
-2.752	$(\mathbf{7}, \mathbf{7})$

Level 5	
m^2L^2	Irrep
-2.721	$(\mathbf{2}, \mathbf{2})$
-2.657	$(\mathbf{2}, \mathbf{2})$
-3.056	$(\mathbf{4}, \mathbf{4})$
-2.567	$(\mathbf{4}, \mathbf{4})$
-2.930	$(\mathbf{6}, \mathbf{6})$
-2.736	$(\mathbf{8}, \mathbf{8})$

Level 6	
m^2L^2	Irrep
-2.400	$(\mathbf{3}, \mathbf{3})$
-2.266	$(\mathbf{3}, \mathbf{3})$
-2.910	$(\mathbf{5}, \mathbf{5})$
-2.895	$(\mathbf{7}, \mathbf{7})$
-2.577	$(\mathbf{9}, \mathbf{9})$

Table 1: Tachyonic Kaluza-Klein modes, their masses, m, in terms of the AdS length L and their representations under SO(3) × SO(3).

Our Strategy

- We work directly in lower dimensions, proposing a bottom-up construction of classical de Sitter vacuum, without invoking non-perturbative effects.
- The model will be subject to certain Infra-red consistency conditions, so that it might be embeddable in quantum gravity.
- We will focus on de Sitter vacua. Standard model-like particle physics are beyond the scope of this talk.

Framework

- 6D (1,0) supergravity with 8 chiral supercharges.

$$\underbrace{\left(e_\mu^m, \psi^+_{\mu A}, B^+_{\mu\nu}\right)}_{\text{graviton}}, \quad \underbrace{\left(B^-_{\mu\nu}, \chi^-_A, \varphi\right)}_{\text{tensor}}, \quad \underbrace{\left(A_\mu, \lambda^+_A\right)}_{\text{vector}}, \quad \underbrace{\left(4\phi, \psi^-\right)}_{\text{hyper}}$$

- Fermions are chiral (symplectic Majorana-Weyl) and transform under **2** of SU(2) R-symmetry groups (index *A*).
- The chiral fields induce local and global anomalies.
- n_T tensor multiplet: scalars parametrize $SO(1, n_T)/SO(n_T)$, n_H hypermultiplet: scalars parametrize $Usp(2, 2n_H)/(SU(2)_R \times Usp(2n_H))$.
- We consider gauging a subgroup of $SU(2)_R \times Usp(2n_H)$.

- There exist many consistent (1,0) models from string theory compactifications, which allows us to study similarity and difference of our model compared to stringy models.
- 6D models has diverse vacuum solutions: $\text{Mink}_4 \times S^2$, $dS_4 \times S^2$, $(A)dS_3 \times S^3$, $(A)dS_2 \times S^4$, $(A)dS_2 \times S^2 \times S^2$, reminiscent of 10D string theory, and allows us to study scale separation.
- Gauged R-symmetry is a crucial distinction between our model and string models, which leads to a non-trivial scalar potential (even in the absence of hyperscalars).
- Previous swampland bounds derived in [Kim, Shiu, Vafa, '19;Kim, Vafa, Xu,'24;Birkar, Lee,'25; Hamada, Loges,'25 $\cdots$] does not applies to models with R-symmetry gauging.

Local anomalies

- $G = \prod_i G_i \times U(1)_R$,

  Local anomalies from chiral spin-$\frac{3}{2}$, spin-$\frac{1}{2}$ and 2-form are [Alvarez-Gaume, Ginsparg, '85]

$$\mathrm{ind}\, iD_{3/2} = \frac{1}{(16\pi^2)^2} \int \left(\frac{245}{360} \mathrm{tr} R^4 - \frac{43}{288} (\mathrm{tr} R^2)^2 \right) \dim r - \frac{19}{6} \mathrm{tr} R^2\, \mathrm{tr}_r F^2 + \frac{10}{3} \mathrm{tr}_r F^4 \ ,$$

$$\mathrm{ind}\, iD_{1/2} = \frac{1}{(16\pi^2)^2} \int \left(\frac{1}{360} \mathrm{tr} R^4 + \frac{1}{288} (\mathrm{tr} R^2)^2 \right) \dim r + \frac{1}{6} \mathrm{tr} R^2\, \mathrm{tr}_r F^2 + \frac{2}{3} \mathrm{tr}_r F^4 \ ,$$

$$\mathrm{ind}\, iD_A = \frac{1}{(16\pi^2)^2} \int \left(\frac{28}{360} \mathrm{tr} R^4 - \frac{8}{288} (\mathrm{tr} R^2)^2 \right) \ ,$$

Cancellation of local anomalies

- $G = \prod_i G_i \times U(1)_R$,
  the local anomalies can be cancelled via Green-Schwarz mechanism provided that

$$\frac{I_8}{2\pi} = \tfrac{1}{2}\Omega_{\alpha\beta} Y^\alpha Y^\beta$$

$$Y^\alpha = \frac{1}{16\pi^2}\left(\frac{1}{2}a^\alpha \,\mathrm{tr}\, R^2 + b_i^\alpha\left(\frac{2}{\lambda_i}\mathrm{tr}\, F_i^2\right) + 2\,c^\alpha F_1^2\right)$$

  where λ_i is the index in the fundamental irreps and $\Omega_{\alpha\beta}$ is the SO(1,n_T) invariant metric.
- Deforming the 3-form field strength

$$H^\alpha_{(3)} = dB^\alpha_{(2)} \Rightarrow H^\alpha_{(3)} = dB^\alpha_{(2)} + v_z^\alpha \mathrm{CS}_{(3)z}, \quad dH^\alpha_{(3)} = Y^\alpha,$$

$$\delta B^\alpha_{(2)} = v_z^\alpha \Lambda_z F_{(2)z}, \quad \Delta S = \tfrac{1}{2}\int \Omega_{\alpha\beta} B^\alpha \wedge Y^\beta, \quad z = i, 1$$

Cancellation of global anomalies

The globally well-definedness of Green-Schwarz term implies [Monnier and Moore, '18]

$$e^{\frac{\mathrm{i}}{2}\int_{M_6}\Omega_{\alpha\beta}B^\alpha\wedge Y^\beta} \hookrightarrow \mathrm{Wu-Chern-Simons\ on\ } X_7, \quad M_6=\partial X_7$$

Independent of the choice of $X_7 \Rightarrow$ triviality of the partition of a TQFT on closed X_7 with spin structure which leads to

- $a^\alpha, b^\alpha_r, \frac{1}{2}c^\alpha \in \Lambda_S$ with Λ_S being a unimodular lattice ;
- $a \in \Lambda_S$ is a characteristic element $\quad a\cdot x = x\cdot x \mod 2$;
- $\Omega_7^{\mathrm{Spin}}(BG)=0$, which vanishes for U(1) or any simply-connected non-Abelian compact Lie group.

The concrete model

We consider $n_T = 1$, which enjoys a local diff-invariant Lagrangian description. Local supersymmetry determines [Riccioni, '01]

$$e^{-1}\mathscr{L} = \Big[\frac{1}{4}R - \frac{1}{4}\partial_\mu\phi\partial^\mu\phi - \frac{1}{12}e^{2\phi}H_{\mu\nu\rho}H^{\mu\nu\rho} - \frac{1}{4}f_z(\phi)F_{z\mu\nu}F_z^{\mu\nu}$$
$$\underline{+\frac{1}{8}e^{-1}v_{z-}\varepsilon^{\mu\nu\rho\sigma\tau\lambda}B_{\mu\nu}F_{z\rho\sigma}F_{z\tau\lambda}) + \frac{1}{2}e^{-1}\varepsilon^{\mu\nu\rho\sigma\lambda\tau}v_{z+}v_{z-}X_{z\mu\nu\rho}X_{z'\sigma\lambda\tau}} - V(\phi)\Big]$$

where the underlined terms are Green-Schwarz counterterms.

$$H_{\mu\nu\rho} = 3\partial_{[\mu}B_{\nu\rho]} - 6v_{z+}\mathrm{CS}_{z\mu\nu\rho}\ , \quad f_z(\phi) = v_{z+}e^{\phi} + v_{z-}e^{-\phi}.$$

In 10D heterotic or type I string theory, the GS counterms appears at 8-derivative level(not yet fully supersymmetrized).

$$\text{Scalar potential}: V = f_z^{-1}(\phi)g_z^2 C^I_{zAB}C_z^{IAB},$$

$$T_z^I \in H_z \in SU(2)_R \times Usp(2n_H), \quad C^I_{zAB} = (L^{-1}T_z^I L)_{AB}, \quad L(\varphi) \in \frac{Usp(2,2n_H)}{SU(2)_R \times Usp(2n_H)},$$

If $T_z \in SU(2)_R$, C_z is non-vanishing even for vanishing hyperscalars.

Variation of the action and anomaly matching

The one-loop effective action in Euclidean spacetime denote by Γ_M is the anomaly $\mathscr{A}$ given by

$$\delta \Gamma_M = \int I_6^1 \equiv \mathscr{A} \ ,$$

Under the gauge transformation

$$\delta B_{\mu\nu} = 2 v_{z+} \mathrm{Tr}_z \left(\Lambda \partial_{[\mu} A_{\nu]} \right) \ ,$$

the gauge variation of $\mathscr{L}_{GS}$ gives

$$\delta \mathscr{L}_{GS} = - v_{z+} v_{z'-} \mathrm{Tr}_z (\Lambda dA) \, \mathrm{Tr}_{z'} F^2 \ ,$$

which cancels precisely the one-loop gauge anomaly

$$I_6^1 + \delta \mathscr{L}_{GS} = 0 \ .$$

The gauge anomaly of the classical action implies supersymmetry anomaly, which must satisfy Wess-Zumino consistency condition

$$\delta_g \delta Q \mathscr{L} - \delta_Q \delta_g \mathscr{L} = 0 \ .$$

The simplest example with a single U(1) gauge group

We consider gauging just $U(1)_R \in SU(2)_R$ $(n_V = 1)$

$$n_T = 1 \Rightarrow n_H = 273 - 29n_T + n_V \Rightarrow n_H = 245.$$

After truncating out the hyperscalars, the bosonic Lagrangian is given by

$$\begin{aligned} e^{-1}\mathscr{L} = &\frac{1}{4}R - \frac{1}{4}\partial_\mu\phi\partial^\mu\phi - \frac{1}{12}e^{2\phi}H_{\mu\nu\rho}H^{\mu\nu\rho} - \frac{1}{4}f(\phi)F_{\mu\nu}F^{\mu\nu} \\ &+ \frac{1}{8}e^{-1}v_-\varepsilon^{\mu\nu\rho\sigma\tau\lambda}B_{\mu\nu}F_{\rho\sigma}F_{\tau\lambda} - \frac{1}{2}f^{-1}(\phi) + \text{higher order terms} , \end{aligned}$$

where $H_{\mu\nu\rho} = 3\partial_{[\mu}B_{\nu\rho]} + 6v_+A_{[\mu}\partial_\nu A_{\rho]}$ and

$$f(\phi) := v_+e^\phi + v_-e^{-\phi} \quad \text{(tree level+one loop)} .$$

Classically, the values of v_+, v_- are arbitrary. Anomaly freedom fixes them.

- $\langle e^{-2\phi} \rangle$ plays a role similar to the string loop expansion parameter in heterotic string. In string frame

$$V \propto \frac{\sqrt{-g}e^{2\phi}}{v_+ + v_- e^{-2\phi}} ,$$

- In the weak coupling limit, V contains infinite loop contributions characterized by power of $e^{-2\phi}$.
- It is consistent with S-duality exchanging $v_+ \to v_-$ and $\phi \to -\phi$.
- $v_- = 0$, it becomes Salam-Sezgin model [Salam-Sezgin, '84]. Properties of models with $v_- \neq 0$ have never been considered before.

$M_4 \times S^2$ type solutions

The ansatz

$$R_{\mu\nu\rho\sigma} = \frac{\varepsilon}{L^2}(g_{\mu\rho}g_{\nu\sigma} - g_{\mu\sigma}g_{\nu\rho}) \,, \quad R_{ijkl} = \frac{1}{a^2}(g_{ik}g_{jl} - g_{il}g_{jk}),$$
$$F_{ij} = k\,\varepsilon_{ij} \,, \quad \phi = \phi_0 \,, \quad f_0 \equiv v_+ e^{\phi_0} + v_- e^{-\phi_0} > 0 \,, \quad \text{rest} = 0 \,.$$

We find the following solutions:

$\frac{1}{2}-\mathrm{BPS}$ $\boxed{\mathrm{Mink}_4 \times S^2}$ $\varepsilon = 0 \,, \quad \phi = \phi_0 \,, \quad k = \pm\frac{1}{f_0} \,, \quad \frac{1}{a^2} = \frac{2}{f_0} > 0 \,,$

The Dirac quantization condition (fermions carrying unit $U(1)_R$ charge)

$$ka^2 = \frac{1}{2}n, \quad n = \pm 1 \,.$$

$M_4 \times S^2$ type solutions

If $v_+ v_- > 0$,

$$n = 0 : \quad L^2 = 12\sqrt{v_+ v_-} = 3a^2 \; , \quad k = 0 \; , \quad \phi_0 = \frac{1}{2} \ln \left(\frac{v_-}{v_+} \right) \; ,$$

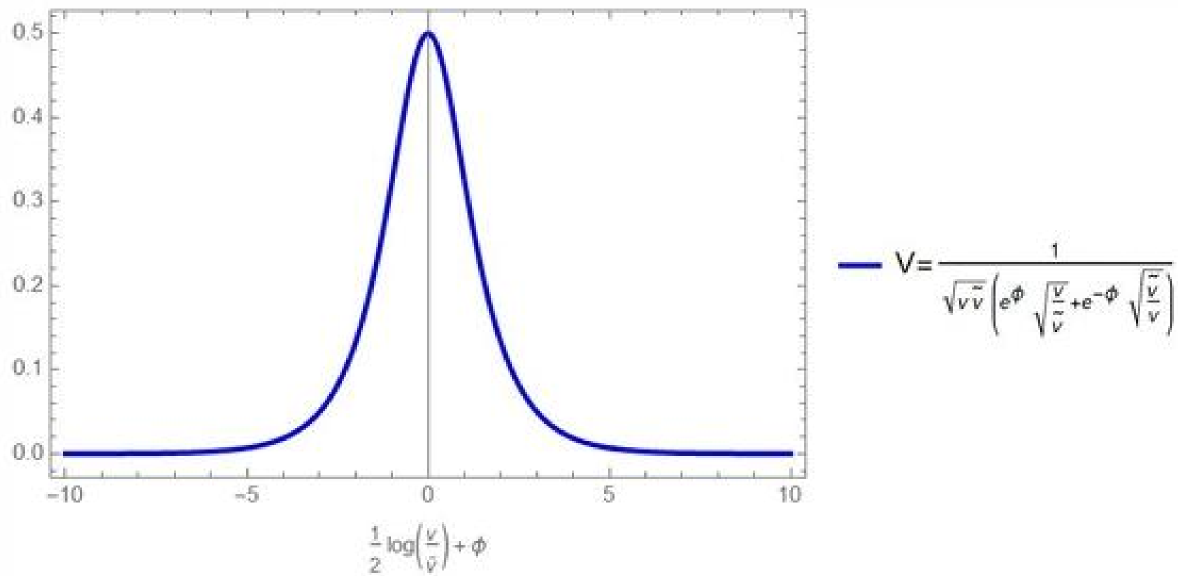
$$n = \pm 1 : \quad L^2 = \frac{27}{2}\sqrt{v_+ v_-} = \frac{9}{2}a^2 \quad k = \pm \frac{1}{6\sqrt{v_+ v_-}} \; , \quad \phi_0 = \frac{1}{2} \ln \left(\frac{v_-}{v_+} \right) \; .$$

There is also $\mathrm{AdS}_4 \times S^2$ solution to the classical equation of motion, which violates the Dirac quantization condition.

dS_6 solution

This dS_6 solution is given by

$$R_{\mu\nu} = \frac{5}{L^2} g_{\mu\nu} = \frac{1}{4\sqrt{v_+ v_-}} g_{\mu\nu} \,, \quad \phi = \frac{1}{2}\ln(v_-/v_+) \,, \quad H_{(3)} = 0 \,, \quad F_{(2)} = 0 \,.$$



About this vacuum, $g_{\mu\nu}$, A_μ and $B_{\mu\nu}$ and dilatino are massless, the gravitino, which eats the gaugino, is in a unitary massive representation, and the dilaton is tachyonic with $m^2 = -10/L^2$.

Absence of local anomaly implies

$$I_8 = \frac{2\pi}{(16\pi^2)^2}\left(\mathrm{tr}R^2 + 2c_+F^2\right)\left(\mathrm{tr}R^2 + 2c_-F^2\right).$$

The anomaly coefficients lie on the lattice of signature (1,1). From anomaly cancellation $I_6^1 + \delta\mathscr{L}_{GS} = 0$, we read off v_+v_-.

$$v_+v_- = \frac{c_+c_-}{128\pi^3}\,.$$

Thanks to the scaling symmetry, the results depend only on v_+v_-

$$g_{\mu\nu} \to g_{\mu\nu}, \quad B_{\mu\nu} \to \lambda B_{\mu\nu}, \quad A_\mu \to A_\mu, \quad e^\phi \to \lambda^{-1}e^\phi, \quad v_+ \to \lambda v_+, \quad v_- \to \lambda^{-1}v_-$$

Existence of de Sitter requires $c_+c_- > 0$, can we actually have such models?

Diagonal gauging of $U(1)_{R+}$

- Hyperfermions are neutral under $U(1)_R$, $v_+ v_- < 0$.

$$\frac{1}{2\pi} I_8 = \frac{1}{(16\pi^2)^2}\Bigg\{ (tr\,R^2)^2 - \frac{1}{6}\mathrm{tr}\,R^2\Big[(n_V - 20)F^2 + \sum_{i=1}^{n}\Big(\mathrm{Tr}_{ad}F_i^2 - \mathrm{tr}_r F_i^2\Big)\Big]$$
$$+\frac{2}{3}\Big[\boxed{-(n_V+4)F^4} + \sum_{i=1}^{n}\Big(-\mathrm{Tr}_{ad}F_i^4 + \underline{\mathrm{tr}_r F_i^4} - 6F^2\mathrm{Tr}_{ad}F_i^2\Big)$$
$$\underline{+6\sum_{i<j,r,s} n_{rs}^{ij}(\mathrm{tr}_r F_i^2)(\mathrm{tr}_s F_j^2)\Big]}\Bigg\},$$

- In order to have $v_+ v_- > 0$, we must have hyperfermions charged under the $U(1)_R$. Gauge the diagonal sum

$$U(1)_{R+} := T + T_3, \quad T_3 \in SU(2)_R$$

where

$$T = \sum_{i=1}^{n_H} q^i h_i\ , \quad h_i \in u(1)^{n_H} \in Usp(2n_H)$$

- Such a diagonal gauging was also noted in [Suzuki, Tachikawa, '05].

We assign unit charge to all hyperfermions,

$$q_i = 1,$$

the anomaly polynomial factorizes nicely

$$I_8 = \frac{2\pi}{(16\pi^2)^2}\left(\mathrm{tr}R^2 + 4F^2\right)\left(\mathrm{tr}R^2 + 40F^2\right)$$

The anomaly coefficients

$$a^\alpha = (2,2), \quad c^\alpha = (2,20)\ ,$$

lie on an even lattice with basis vectors

$$e_1^\alpha = (1,0), \quad e_2^\alpha = (0,1)\ ,$$

From anomaly cancellation $I_6^1 + \delta\mathscr{L}_{GS} = 0$, we read off v_+v_-.

$$v_+v_- = \frac{40}{32\pi^3}\ .$$

We parameterize the fluctuations around the background as follows

$$g_{MN} = \bar{g}_{MN} + h_{MN} \ , \quad \phi = \phi_0 + \varphi, \quad A_M = \bar{A}_M + a_M \ , \quad B_{MN} = \bar{B}_{MN} + \tilde{b}_{MN} \ .$$

We choose the de Donder-Lorentz gauge:

$$\nabla^i h_{ij} = \frac{1}{2}\nabla_j h^k_{\ k} \ , \quad \nabla^i h_{i\mu} = 0 \ ,$$
$$\nabla^i a_i = 0 \ , \qquad \nabla^i b_{iM} = 0.$$

Harmonic expansions under this gauge take the form

$$\varphi = \sum_{\ell \geq 0} \phi^{(\ell)} Y^{(\ell)} \ , h_{\mu\nu} = \sum_{\ell \geq 0} h^{(\ell)}_{\mu\nu} Y^{(\ell)} \ , h_{\mu i} = \sum_{\ell \geq 1} h^{(\ell)}_{\mu} Y^{(\ell)}_i \ , h_{ij} = g_{ij} \sum_{\ell \geq 0} N^{(\ell)} Y^{(\ell)}$$

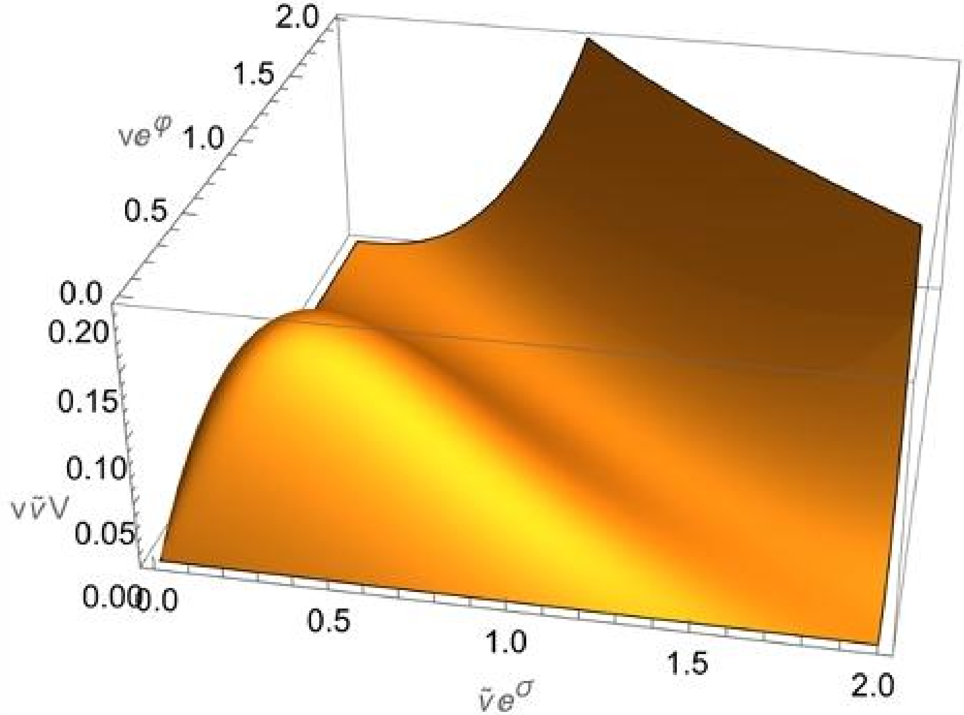
$$a_\mu = \sum_{\ell \geq 0} a^{(\ell)}_\mu Y^{(\ell)} \ , \ a_i = \sum_{\ell \geq 1} a^{(\ell)} Y^{(\ell)}_i \ , \ b_{\mu\nu} = \sum_{\ell \geq 0} b^{(\ell)}_{\mu\nu} Y^{(\ell)} \ , \ b_{\mu i} = \sum_{\ell \geq 1} b^{(\ell)}_\mu Y^{(\ell)}_i \ ,$$

$$b_{ij} = \varepsilon_{ij} b^{(0)} Y^{(0)} \ ,$$

Mass spectrum of fluctuations around $dS_4 \times S^2$

s	Δ	Comments
2	$\frac{3}{2}+\frac{3}{2}\sqrt{1-2\ell(\ell+1)}$	$\ell \geq 0$, massless for $\ell = 0$
$\frac{3}{2}$	$\frac{3}{2}+\frac{3}{2}\sqrt{-2\ell(\ell+1)}$	Dirac massive gravitino for $\ell \geq 1$ massive for $\ell = 0$ (discrete series)
1	$\frac{3}{2}+\frac{3}{2}\sqrt{\frac{1}{9}-2\ell(\ell+1)}$ $\frac{3}{2}+\frac{1}{2}\sqrt{1+4x_i(\ell)}$, $i=2,3,4$ $\frac{3}{2}+i\frac{\sqrt{5}}{2}$	$\ell \geq 1$, $\ell \geq 1$, $x_i(\ell)$ are roots of (5.16) $x_2 = 0$ massless for $\ell = 1$ $x_3 = 0$ for $\ell = 0$ (massless 2-form) $x_4 = -\frac{3}{2}$ for $\ell = 0$ (massive vector)
$\frac{1}{2}$	$\frac{3}{2}+\frac{3}{2}\sqrt{-2\ell(\ell+1)}$ $\frac{3}{2}+\frac{3}{2}\sqrt{-2\ell^2+2}$ $\frac{3}{2}+\frac{3}{2}\sqrt{-2\ell(\ell+2)}$	complex massive spinor for $\ell \geq 1$ massless for $\ell = 0$ complex massive spinor for $\ell \geq 2$ massless for $\ell = 1$ complex massive spinor for $\ell \geq 1$ massless for $\ell = 0$
0	$\frac{3}{2}+\frac{3}{2}\sqrt{1-2\ell(\ell+2)}$ $\frac{3}{2}+\frac{3}{2}\sqrt{3-2\ell^2}$ $\frac{3}{2}+\frac{3}{2}\sqrt{\frac{11}{3}-2\ell(\ell+1)}$	$\ell \geq 1$ $\ell \geq 0$, **tachyon** for $\ell = 0$ $\ell \geq 0$, **tachyon** for $\ell = 0$

Two scalars are tachyonic modes, one from the volume of the internal space and the other from the dilaton.



$$V = \frac{e^{\varphi+\sigma}}{2(v_- e^{\sigma} + v_+ e^{\varphi})}(2 - v_- e^{\sigma} - v_+ e^{\varphi})^2 \ .$$

The kinetic terms of σ and φ are canonically normalized. The effective scalar potential for $n = 0$ vacuum is unbounded from below.

- de Sitter vacuum breaks supersymmetry. The spin-1/2 goldstino is eaten by spin-3/2 gravitino, which belongs to the discrete series representation of SO(1,4) group.

  The unitary representation of fermions labeled by (Δ, s) of SO(1,4) are [Letsios , '23]

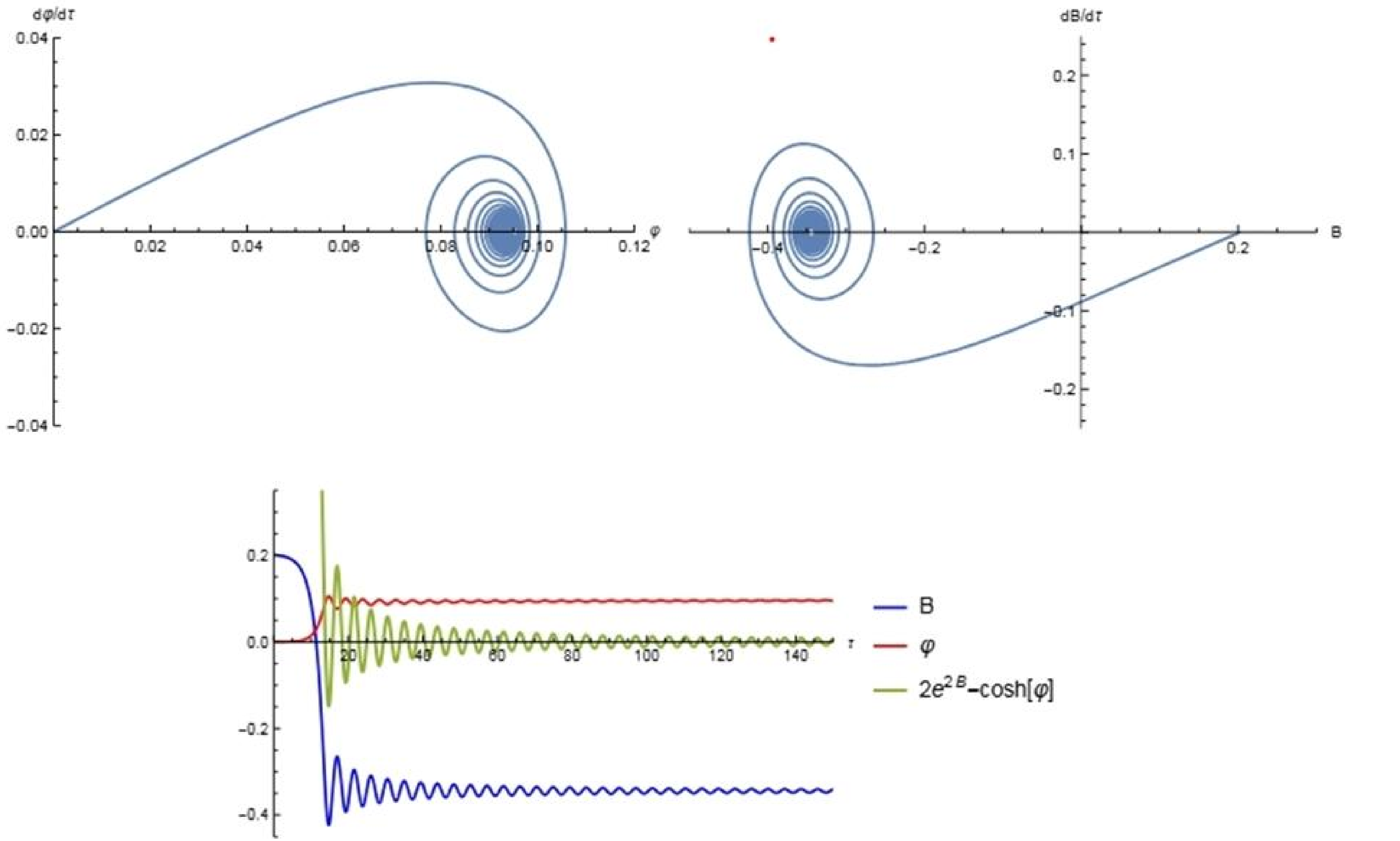
$$\mathscr{F}_{\Delta,s} \quad , \quad \Delta = \frac{3}{2} + i\mu \quad , \quad (\mu \in \mathbb{R}^+ \backslash \{0\}) \quad \text{Principal Series}$$

$$\mathscr{F}_{\Delta,s} \quad , \quad \Delta = 3/2 \quad , \quad \text{Discrete Series}$$

$$\mathscr{F}_{\Delta,s} \quad , \quad \Delta = 2 + t \quad , \quad \left(t = \frac{1}{2}, \ldots, s-1, s \in \mathbb{Z}_+ - \frac{1}{2}\right) \quad \text{Discrete Series}$$

- The fluctuations of tachyonic modes trigger a flow from $\mathrm{dS}_4 \times S^2$ towards $\mathrm{Mink}_4 \times S^2$.

Flow from $dS_4 \times S^2$ to $Mink_4 \times S^2$



$U(1)_{R+} \times U(1)'$ model

In this case we have $n_H = 2 + 244 = 246$. The number of hyperini carring charges $(\pm p, \pm q)$, $(p, q = 1, 2, 3)$ under $U(1)_{R+} \times U(1)'$ is denoted by n_{pq} [Suzuki and Tachikawa, '05]

$$n_{pq} = \begin{pmatrix} 2 & 4 & 6 \\ 150 & 4 & 2 \\ 6 & 62 & 10 \end{pmatrix} .$$

With the above charge assignments, the resulting anomaly polynomial factorizes as

$$I_8 = \frac{2\pi}{(16\pi^2)^2} \left(\mathrm{tr} R^2 + 30 F^2 + 76 F^{'2} \right) \left(\mathrm{tr} R^2 + 196 F^2 + 24 F^{'2} \right) ,$$

where F and F' are the $U(1)_{R+}$ and $U(1)'$ field strengths. The corresponding anomaly coefficient vectors are

$$a^\alpha = (2, 2), \quad c^\alpha = (15, 98), \quad c^{'\alpha} = (38, 12) ,$$

where a, $\frac{1}{2}c$ and $\frac{1}{2}c'$ must lie on certain unimodular lattice. We find that such a lattice indeed exists with basis vectors

$$e_1^\alpha = (1, 0), \quad e_2^\alpha = (\tfrac{1}{2}, 1) .$$

Conclusion and discussions

- We observe that it is possible to obtain de Sitter vacuum in 6D gauged supergravity, without invoking orientifolds, α' corrections and so on.
- Higher order corrections take the schematic form

$$\Delta\mathscr{L} = \frac{1}{\kappa^2}\Big[\kappa(a^+ e^{\phi} + a^- e^{-\phi})\mathrm{Riem}^2 + \kappa^3(a^+ v_- + a^- v_+)(a^+ e^{\phi} + a^- e^{-\phi}) f F^4 + \cdots\Big] ,$$

The higher order corrections are suppressed by $\frac{1}{c_+} + \frac{1}{c_-}$. We need models with large anomaly coefficients associated with $U(1)_{R+}$.

Within the $U(1)_{R+}$ models. Among 245 hypers, n of them carry charge q and the rest carry charge p.

n	q	p	c_-	c_+	c_+/c_-
1	185	101	186542	23732	0.13
1	185	115	241502	30260	0.13
2	1	9999	1799640020	224955002	0.06
2	25	11	2232	324	0.15
2	53	157	444212	55400	0.12
2	241	577	6003570	747924	0.12

The effective coupling $e^{-2\phi}$ is controlled by c_+/c_-.

- Certain de Sitter vacuum in IIB string theory can have tachyonic perturbations [Andriot, '21]
- In the past, tachyon field has been proposed as candidate of dark energy.

Open questions

- The simple toy model does not exhibit scale separation between dS_4 and S^2

$$L^2_{\mathrm{dS}} = \frac{27}{2}\sqrt{v_+ v_-}, \quad L^2_{S^2} = 3\sqrt{v_+ v_-}, \quad L^2_{\mathrm{dS}}/L^2_{S^2} = \frac{9}{2}\,.$$

- Preliminary study shows the $U(1)_{R+} \times U(1)'$ seems to exhibit scale separation.
- If we turn on hyperscalars, will there be new vacuum?
- Can higher order corrections to scalar potential stablize the dilaton and breathing mode?
- Find string theory embedding or eventually rule out this class of models by other consistency criteria.

Thank you for listening!

Mass spectrum of fluctuations around $\text{Mink}_4 \times S^2$

s	$a^{-2}m^2$	Comments
2	$\ell(\ell+1)$	$\ell \geq 0$
$\frac{3}{2}$	$\ell(\ell+1)$	$\ell \geq 1$: complex massive gravitino $\ell = 0$: massless gravitino
1	$\ell(\ell+1)$ $\ell(\ell+1)$ $\ell(\ell+1)+1-\frac{v\tilde{v}}{2a^4} \pm \left[4\ell(\ell+1)\left(1-\frac{v\tilde{v}}{4a^4}\right)+\left(1-\frac{v\tilde{v}}{2a^4}\right)^2\right]^{\frac{1}{2}}$	$\ell \geq 1$ $\ell \geq 1$ m_+^2 for $\ell \geq 0$ $m_-^2 = 0$ for $\ell = 0$ (massless 2-form) $m_-^2 = 0$ for $\ell = 1$
$\frac{1}{2}$	$\ell(\ell+1)$ $\ell(\ell+1)+1-\frac{v\tilde{v}}{2a^4} \pm \left[4\ell(\ell+1)\left(1-\frac{v\tilde{v}}{4a^4}\right)+\left(1-\frac{v\tilde{v}}{2a^4}\right)^2\right]^{\frac{1}{2}}$	$\ell \geq 1$: complex massive spinor m_+^2 for $\ell \geq 0$: complex massive spinor m_-^2 for $\ell \geq 2$: complex massive spinor $m_-^2 = 0$ for $\ell = 0, 1$
0	$\ell(\ell+1)$ $\ell(\ell+1)+1-\frac{v\tilde{v}}{2a^4} \pm \left[4\ell(\ell+1)\left(1-\frac{v\tilde{v}}{4a^4}\right)+\left(1-\frac{v\tilde{v}}{2a^4}\right)^2\right]^{\frac{1}{2}}$	$\ell \geq 1$ m_+^2 for $\ell \geq 0$ $m_-^2 = 0$ for $\ell = 0$

The massless states for a 4D $N = 1$ supergravity multiplet, Yang-Mills multiplet and a scalar multiplet.

A lattice is said to be *even* if

$$v \cdot v \in 2\mathbb{Z} \quad \forall v \in \Gamma \, .$$