An exact AdS3/CFT2 duality

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Based mainly on work with

Lorenz Eberhardt and Rajesh Gopakumar.

Motivation

It is generally believed that the CFT dual of string theory on

$$\mathrm{AdS}_3 \times \mathrm{S}^3 \times \mathbb{T}^4$$

is on the same moduli space of CFTs that also contains the symmetric orbifold theory

$$\mathrm{Sym}_N(\mathbb{T}^4) \equiv (\mathbb{T}^4)^N / S_N$$

[Maldacena '97], …. see e.g. [David et.al. '02]

Motivation

However, it is not known what precise string background is being described by the symmetric orbifold theory itself.

see however [Larsen, Martinec '99]

On the other hand, there is an explicitly solvable world-sheet theory for strings on this background in terms of an sl(2,R) WZW model.

[Maldacena, (Son), Ooguri '00 & '01]

However, it is not known what precise dual CFT (on the above moduli space) this corresponds to.

Motivation

In fact, the only consensus was that the actual symmetric orbifold theory **cannot** be dual to the WZW model…

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In fact, the only consensus was that the actual symmetric orbifold theory **cannot** be dual to the WZW model…

The basic reason for this is that the WZW model describes the background with pure NS-NS flux, which is known to have long string solutions.

[Seiberg, Witten '99], [Maldacena, Ooguri '00]

These long strings live near the boundary of AdS, and they give rise to a continuum of excitations that are not present in the actual symmetric orbifold theory.

Symmetric orbifold

In a separate development, the higher spin version of the AdS/CFT duality was studied, and it was found that the symmetric orbifold theory should correspond to the tensionless limit of string theory on AdS3.

[MRG, Gopakumar '10 - '14]

The tensionless limit arises when the spacetime geometry is of string size, i.e. in the deep stringy regime.

In the context of the WZW description, this should be the situation where the level of the sl(2,R) affine theory takes the smallest possible value.

WZW model

This led us to study the spacetime spectrum of the k=1 sl(2,R) WZW model systematically.

As will be explained in more detail below, we found that the k=1 theory indeed has massless higher spin fields, and that its spectrum resembles that of the symmetric orbifold theory in the large N limit.

[MRG, Gopakumar, Hull ’17],
[Ferreira, MRG, Jottar ’17],
[MRG, Gopakumar ’18]
[Giribet, Hull, Kleban, Porrati, Rabinovici ’18]

WZW model

However, the k=1 theory in the NS-R formalism is not really well-defined.

In particular, the full WZW model is in this case

$$\mathfrak{sl}(2)_k^{(1)} \oplus \mathfrak{su}(2)_k^{(1)} \oplus [\mathfrak{u}(1)^{(1)}]^{\oplus 4}$$

and at k=1

$$\mathfrak{su}(2)_1^{(1)} \cong \mathfrak{su}(2)_{-1} \oplus 3 \text{ free fermions}$$

non-unitary

WZW model

Furthermore, the WZW model still seems to contain a continuum of states (that are not present in the symmetric orbifold theory).

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As it turns out, both of these problems can be overcome by considering the alternative description of string theory on AdS3 x S3 in terms of the so-called hybrid formalism.

[Berkovits, Vafa, Witten ’99]

WZW model

In this formulation, the AdS3 x S3 part is described (for pure NS-NS flux) by a supergroup WZW model, namely

$$\mathfrak{psu}(1,1|2)_k$$

and this description continues to make sense also for k=1. However, something special happens for this value: as will be explained below, the representation theory is much more constrained for k=1, and in particular, the continuum of representations is not allowed any longer.

[Eberhardt, MRG, Gopakumar ’18]

WZW model

Taking this into account, we have shown that the resulting spacetime spectrum agrees precisely with that of the symmetric orbifold theory (in the large N limit)!

[Eberhardt, MRG, Gopakumar ’18]

Furthermore, we have shown that the structure of the symmetric orbifold correlators is precisely reproduced from the world-sheet analysis, and the same is true for the operator algebra.

[Eberhardt, MRG, Gopakumar, to appear]
[Eberhardt, MRG ’19]

Plan of talk

1. **Introduction and Motivation**
2. The NS-R construction
3. The supergroup hybrid formulation
4. Correlators and operator algebra
5. Conclusions and Outlook

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NS-R WZW model

Let us begin by reviewing some basic facts about the WZW model based on sl(2,R). [Maldacena, Ooguri '00]

In the susy case, the relevant chiral algebra is

$$\mathfrak{sl}(2,\mathbb{R})^{(1)}_k \cong \mathfrak{sl}(2,\mathbb{R})_{k+2} \oplus 3 \text{ free fermions}$$

bosonic: $J^3_n, J^\pm_n$

decoupled

The free fermions sit in the usual NS/R representations.

NS-R WZW model

The representations of the bosonic sl(2,R) affine algebra are characterised by the sl(2,R) reps of the highest weights. There are 2 classes of sl(2,R) reps that appear:

[Maldacena, Ooguri '00]

Discrete lowest weight reps:

$$\mathcal{D}_j^+ : \qquad C = -j(j-1)\ , \quad J_0^- |j,j\rangle = 0$$

Continuous reps:

$$C_\alpha^j : \quad C = -j(j-1) = \tfrac{1}{4} + p^2\ , \quad |j,m\rangle \text{ with } m \in \alpha + \mathbb{Z}$$

$$(j = \tfrac{1}{2} + ip)$$

No-ghost theorem

Because of the Maldacena-Ooguri (unitarity) bound,

$$\text{MO-bound :} \qquad \frac{1}{2} < j < \frac{k+1}{2}$$

[Petropoulos '90]
[Hwang '91]
[Evans, MRG, Perry '98]
[Maldacena, Ooguri '00]

the (discrete) spectrum is bounded from above. Additional states are spectrally flowed images of these two classes of representations

They are not Virasoro highest weight, and are therefore best described in terms of the spectral flow w.

[Maldacena, Ooguri '00]
see also [Henningson et.al. '91]

Physical states

This description is covariant, i.e. we need to impose the physical state condition, e.g. in NS sector

$$G_r^{\text{tot}}\Phi = 0 \quad (r > 0)$$

$$(L_0^{\text{tot}} - \tfrac{1}{2})\Phi = 0 \ .$$

In particular, the second condition (mass-shell condition) implies that

$$\frac{C}{k} + h_0 + N = \frac{1}{2} \qquad \text{(NS-sector)}$$

Casimir of sl(2,R)

World-sheet conformal
dim. of internal CFT

Dual CFT

The dual (`spacetime') CFT lives on the boundary of AdS3, and we have the identifications

$$L_0^{\mathrm{CFT}} = J_0^3 \ , \qquad L_1^{\mathrm{CFT}} = J_0^- \ , \qquad L_{-1}^{\mathrm{CFT}} = J_0^+ \ ,$$

with a similar relation for the right-movers.

With these preparations at hand, we can now study the physical spectrum of the (spacetime) theory for k=1.

As we shall see, the interesting part of the spectrum comes from the spectrally flowed continuous reps.

Continuous reps

For the spectrally flowed continuous reps, the mass-shell condition (in the NS sector) is at k=1

$[\frac{C}{k} + h_0 + N = \frac{1}{2}]$

$[\sigma^w(L_n) = L_n - wJ_n^3 - \frac{k}{4}w^2\delta_{n,0}]$

$$C - wm - \frac{1}{4}w^2 + N = \frac{1}{2} \qquad \text{where } C = \frac{1}{4} + p^2$$

Here m is the J_0^3 eigenvalue before spectral flow, and we have set $h_0 = 0$ (for simplicity).

For the continuous representations we can simply solve this equation for m. For the case of p=0 we then get

Continuous reps

$$C - wm - \frac{1}{4}w^2 + N = \frac{1}{2} \qquad \text{with } C = \tfrac{1}{4}$$

$$m = \frac{1}{w}\left[N - \frac{w^2+1}{4}\right]$$

Continuous reps

$$C - wm - \frac{1}{4}w^2 + N = \frac{1}{2} \qquad \text{with } C = \tfrac{1}{4}$$

$$m = \frac{1}{w}\left[N - \frac{w^2+1}{4}\right]$$

Then observing that the actual J_0^3 eigenvalue is

$[\sigma^w(J_n^3) = J_n^3 + \frac{wk}{2}\,\delta_{n,0}]$

$$h = m + \frac{w}{2} = \frac{N}{w} + \frac{w^2-1}{4w} \; .$$

Full spectrum

$$h = m + \frac{w}{2} = \underbrace{\frac{N}{w}}_{\text{w-twisted modes}} + \underbrace{\frac{w^2-1}{4w}}_{\text{ground state energy in w-twisted sector}} \ .$$

Symmetric orbifold formula for cycle length w!

[MRG, Gopakumar ’18]
[Giribet, Hull, Kleban, Porrati, Rabinovici ’18]
see also [Giveon, Kutasov, Rabinovici, Sever ’05].

Note that for w=1 and N=0, this includes in particular chiral states (h=0) that correspond to massless higher spin fields!

[MRG, Gopakumar, Hull ’17]
[Ferreira, MRG, Jottar ’17]

Full symmetric orbifold

Thus we recover the full single-particle spectrum of the symmetric orbifold (in the large N limit).

[MRG, Gopakumar ’18]
[Giribet, Hull, Kleban, Porrati, Rabinovici ’18]

For $\mathrm{AdS}_3 \times \mathrm{S}^3 \times \mathbb{T}^4$ at k=1, criticality implies that the bosonic su(2) factor appears at level -1, and thus the analysis in the NS-R sector is a bit formal — in the hybrid formalism this will be cleaner (see below).

Which orbifold

In order to get a sense of what will happen, we can use that

$$\mathfrak{su}(2)_{-1} \oplus \mathfrak{u}(1) = 4 \text{ symplectic bosons}$$

[Goddard, Olive, Waterson '87]

The 4 symplectic bosons behave as ghosts (on the level of the partition function) and remove 4 of the 8 fermions.

This therefore suggests that we end up with 4+4 free bosons and fermions, i.e. with the spectrum of

$$\boxed{\text{symmetric orbifold of } \mathbb{T}^4}$$

[MRG, Gopakumar '18]

Continuum of states

However, the spectrum still seems to have a continuum (we earlier set p=0 by hand), which is not present in the symmetric orbifold theory.

There are also some discrete rep states that do not fit into the above.

Thus we have not quite managed yet to identify the world-sheet theory that corresponds to the symmetric orbifold.

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Hybrid formalism

In the hybrid formalism the world-sheet theory is described (for pure NS-NS flux) by the WZW model based on

$$\mathfrak{psu}(1,1|2)_k$$

together with the (topologically twisted) sigma model for T4. For generic k, this description agrees with the NS-R description a la MO.

[Troost '11], [MRG, Gerigk '11]
[Gerigk '12]

Hybrid formalism

For the following it will be important to understand the representation theory of

$$\mathfrak{psu}(1,1|2)_1$$

The bosonic subalgebra of this superaffine algebra is

$$\mathfrak{sl}(2)_1 \oplus \mathfrak{su}(2)_1$$

Thus only **n**=1 and **n**=2 are allowed for the highest weight states.

Hybrid formalism

One of the key differences to the NS-R formalism is that the fermions of

$$\mathfrak{psu}(1,1|2)_1$$

do not sit in the adjoint representation of the bosonic subalgebra, but rather in bispinor representations.

As a consequence, one cannot decouple the fermions as before and therefore obtain a negative level for su(2). In fact, the k=1 theory seems to be well-defined.

[Eberhardt, MRG, Gopakumar ’18]

Short representations

A generic representation of the zero mode algebra $\mathfrak{psu}(1,1|2)$ has the form

$$(C^{j}_{\alpha}, \mathbf{n})$$

$$(C^{j+\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n+1}) \quad (C^{j+\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n-1}) \quad (C^{j-\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n+1}) \quad (C^{j-\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n-1})$$

$$(C^{j+1}_{\alpha}, \mathbf{n}) \quad (C^{j}_{\alpha}, \mathbf{n+2}) \quad 2\cdot(C^{j}_{\alpha}, \mathbf{n}) \quad (C^{j}_{\alpha}, \mathbf{n-2}) \quad (C^{j+1}_{\alpha}, \mathbf{n})$$

$$(C^{j+\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n+1}) \quad (C^{j+\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n-1}) \quad (C^{j-\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n+1}) \quad (C^{j-\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n-1})$$

$$(C^{j}_{\alpha}, \mathbf{n})$$

continuous rep of sl(2,R)

rep of su(2) of dim = n+1.

Short representations

A generic representation of the zero mode algebra $\mathfrak{psu}(1,1|2)$ has the form

$$(\mathcal{C}^j_\alpha, \mathbf{n})$$

$$(\mathcal{C}^{j+\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n+1}) \quad (\mathcal{C}^{j+\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n-1}) \quad (\mathcal{C}^{j-\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n+1}) \quad (\mathcal{C}^{j-\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n-1})$$

$$(\mathcal{C}^{j+1}_\alpha, \mathbf{n}) \quad (\mathcal{C}^j_\alpha, \mathbf{n+2}) \quad 2\cdot(\mathcal{C}^j_\alpha, \mathbf{n}) \quad (\mathcal{C}^j_\alpha, \mathbf{n-2}) \quad (\mathcal{C}^{j+1}_\alpha, \mathbf{n})$$

$$(\mathcal{C}^{j+\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n+1}) \quad (\mathcal{C}^{j+\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n-1}) \quad (\mathcal{C}^{j-\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n+1}) \quad (\mathcal{C}^{j-\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{n-1})$$

$$(\mathcal{C}^j_\alpha, \mathbf{n})$$

continuous rep
of sl(2,R)

Thus for k=1 need
a short rep!

Short representations

In fact, the only representations that are allowed are

$$(C^{j}_{\alpha}, \mathbf{2})$$

$$(C^{j+\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{1}) \qquad (C^{j-\frac{1}{2}}_{\alpha+\frac{1}{2}}, \mathbf{1})$$

and the shortening condition actually implies that this is only possible provided that

$$j = \frac{1}{2} \quad \longrightarrow$$

NO CONTINUUM!

[Eberhardt, MRG, Gopakumar ’18]

Short representations

The corresponding affine representation $\mathcal{F}_\lambda$ has in fact many null-vectors, and thus, after including the ghosts, the contribution from

$$\mathfrak{psu}(1,1|2)_1$$

just reduces to the zero-modes

$$Z_{\mathrm{AdS}_3\times \mathrm{S}^3} = \left| \sum_{m\in\mathbb{Z}+\lambda} x^m q^{-mw+\frac{w^2}{2}} \right|^2$$

[Morally similar to su(2) at level 1, which only describes the degrees of freedom of a single boson!]

Short representations

$$Z_{\mathrm{AdS}_3\times \mathrm{S}^3} = \left| \sum_{m\in\mathbb{Z}+\lambda} x^m q^{-mw+\frac{w^2}{2}} \right|^2$$

Finally, for every state from T4, exactly one term in this sum is picked out by the mass-shell condition:

topological sector!

One also finds that these are the only representations, i.e. no discrete representations appear.

Physical spectrum

The rest of the analysis works essentially as in the NS-R formulation. Because now the continuum has disappeared (and there are no discrete representations) we get exactly the (single-particle) spectrum of

[Eberhardt, MRG, Gopakumar '18]

$$\mathrm{Sym}_N(\mathbb{T}^4)$$

where, as before, the spectral flow parameter w is to be identified with the length of the single cycle twisted sector in the symmetric orbifold (in the large N limit).

Free field realisation

The affine algebra $\mathfrak{psu}(1,1|2)_1$ actually has a free field realisation in terms of

[Eberhardt, MRG, Gopakumar '18]

$$\begin{aligned}\mathfrak{psu}(1,1|2)_1 &\cong \frac{\mathfrak{u}(1,1|2)_1}{\mathfrak{u}(1)_U \oplus \mathfrak{u}(1)_V} \\ &\cong \frac{\text{2 pairs of symplectic bosons and 2 complex fermions}}{\mathfrak{u}(1)_U \oplus \mathfrak{u}(1)_V} \ .\end{aligned}$$

This allows us to calculate the characters, the fusion rules, etc., and show (with some effort) that the world-sheet theory is consistent.

see also [Gotz, Quella, Schomerus '06]
[Ridout '10]

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Vertex operators

In order to calculate the correlation functions of the dual CFT, introduce the vertex operators on world-sheet

see also [Kutasov, Seiberg '99]

$$V(\psi, x; z) = e^{xJ_0^+}\, e^{zL_{-1}}\, V(\psi, 0; 0) e^{-zL_{-1}}\, e^{-xJ_0^+}$$

spacetime CFT coordinate

worldsheet coordinate

identify fields and states

Vertex operators

In the following concentrate on spectrally flowed image of affine primary — corresponds to ground state of twisted sector:

spectral flow sector

$$V_h^w(x;z)$$

labels rep of affine primary

Here we work in x-basis, in which sl(2,R) acts via

$$J_0^+ = -\frac{\partial}{\partial x}$$

$$J_0^3 = h - x\frac{\partial}{\partial x}$$

$$J_0^- = 2hx - x^2\frac{\partial}{\partial x} \ .$$

Correlation functions

The corresponding correlation functions

$$\Big\langle \prod_{i=1}^{n} V_{h_i}^{w_i}(x_i; z_i) \Big\rangle$$

have a remarkable localisation property on the world-sheet.

[Eberhardt, MRG, Gopakumar, to appear]

To see this, consider

$$\gamma(z) = \frac{\big\langle J^3(z) \prod_{i=1}^{n} V_{h_i}^{w_i}(x_i; z_i) \big\rangle}{\big\langle J^+(z) \prod_{i=1}^{n} V_{h_i}^{w_i}(x_i; z_i) \big\rangle}$$

Covering map

This function behaves as

$$\gamma(z) = \frac{\left\langle J^3(z) \prod_{i=1}^{n} V_{h_i}^{w_i}(x_i; z_i) \right\rangle}{\left\langle J^+(z) \prod_{i=1}^{n} V_{h_i}^{w_i}(x_i; z_i) \right\rangle}$$

$$\sim x_i + c_i (z - z_i)^{w_i} \qquad \text{as } z \to z_i$$

i.e. it defines the covering map of the spacetime sphere by the world-sheet, where

$\gamma(z_i) = x_i$ with ramification of order w_i

Covering surface

This is a very restrictive condition! For a fixed choice of ramification points x_i together with the orders w_i, there are only finitely many such covering maps, i.e. only finitely many (discrete choices) for z_i.

As a consequence, the world-sheet integral localises to a finite number of points.

More conceptually, it shows that the world-sheet is the covering surface of the spacetime CFT.

cf. [Lunin, Mathur '00]
[Pakman, Rastelli, Razamat '09]

Relation to sym orbifold

These covering maps control the calculation of twist field correlators of symmetric orbifold theories: in fact, for the ground states, the entire amplitude just comes from the conformal factor associated to the covering map.

[Lunin, Mathur '00]
[Pakman, Rastelli, Razamat '09]

Up to checking this conformal factor, this makes the matching to the symmetric orbifold manifest. Since our arguments work for arbitrary genus of the world-sheet surface, this reproduces, in particular, the relation between

higher genus ~ 1/N correction

[Eberhardt, MRG, Gopakumar, to appear]

Algebra structure

We have furthermore shown that also the algebraic structure of the dual CFT can be recovered from the world-sheet.

This can be read off from the spectrum generating operators of the spacetime CFT: DDF operators.

[Giveon, Kutasov, Seiberg ’98]
[Kutasov, Seiberg ’99]
[Eberhardt, MRG ’19]

Bosonic case

The basic idea can already be explained in the bosonic case, say for the spacetime Virasoro generators

[Giveon, Kutasov, Seiberg '98]

$$\mathcal{L}_m = \oint \mathrm{d}z \left((1-m^2)\gamma^m J^3 + \frac{m(m-1)}{2}\gamma^{m+1}J^+ + \frac{m(m+1)}{2}\gamma^{m-1}J^- \right)(z) \ .$$

sl(2,R) currents

$\beta\gamma$-system of Wakimoto rep of sl(2,R)

Spacetime Virasoro algebra

[Giveon, Kutasov, Seiberg '98]

These generators satisfy a spacetime Virasoro algebra

$$[\mathcal{L}_m, \mathcal{L}_n] = (m-n)\,\mathcal{L}_{m+n} + \tfrac{k}{2}\,\mathcal{I}\,m(m^2-1)\,\delta_{m+n,0}$$

whose central term equals

$$\mathcal{I} = \oint \mathrm{d}z\ \big(\gamma^{-1}\partial\gamma\big)(z)\ .$$

The value of this integral depends on the spectral flow sector, and one finds

$$\mathcal{I} = w \cdot \mathbf{1}$$

Fractional modes

Since the central term vanishes in the unflowed sector, it follows that $\log(\gamma)$ and hence γ^r is single-valued.

In the w'th spectrally flowed sector we have

$$\big(\sigma^w(\gamma)\big)^m(z) = z^{mw}\gamma^m(z) \ ,$$

and hence the spacetime Virasoro generators remain well-defined for

$$[\mathcal{L}_m = \oint \mathrm{d}z \left((1-m^2)\gamma^m J^3 + \frac{m(m-1)}{2}\gamma^{m+1}J^+ + \frac{m(m+1)}{2}\gamma^{m-1}J^- \right)(z)]$$

$$m \in \frac{1}{w}\mathbb{Z} \quad \longleftarrow \quad \text{modes of w-cycle twisted sector!}$$

[Eberhardt, MRG '19]

Spacetime CFT

The analysis works similarly for the other generators, and also the central terms and ground state energies work out exactly as expected.

In fact, the whole analysis can be done for general k, and we are led to deduce that the continuous world-sheet reps for

$$\mathrm{AdS}_3 \times X$$

give rise to the spacetime CFT

$$\mathrm{Sym}^N\Big(\underbrace{\Big[\text{Liouville with } c^{\mathrm{L}} = 1 + \frac{6(k-3)^2}{k-2}\Big] \times X}_{c=6k}\Big)$$

in the large N limit.

[Eberhardt, MRG '19]
see also [Argurio, Giveon, Shomer '00],

Discrete reps

We should stress that the entire symmetric orbifold comes from the continuous representations on the world-sheet, while the discrete representations on the world-sheet include states that lie below the Liouville gap (and do not seem to fit into this theory nicely).

Reminiscent of the difficulties Maldacena & Ooguri noted already in the calculation of the correlation functions of the discrete states….

Susy case

The supersymmetric generalisation works similarly, and the corresponding statement is that the continuous world-sheet reps of string theory on

$$\mathrm{AdS}_3 \times \mathrm{S}^3 \times \mathbb{T}^4$$

give rise to the spacetime CFT

$$\boxed{\mathrm{Sym}^N\left([\mathcal{N}=4\ \text{ Liouville with } c=6(k-1)] \times \mathbb{T}^4\right)}$$

in the large N limit.

[Eberhardt, MRG ‘19]

Susy case

For k=1, the Liouville part becomes trivial, and we thus recover again

$$\mathrm{Sym}^N\left(\mathbb{T}^4\right)$$

In that case, there are no discrete representations, and thus the statement is clean.

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Conclusions and Outlook

We have given strong evidence that the symmetric orbifold theory is exactly dual to string theory with one unit of NS-NS flux (k=1):

$$\mathrm{Sym}_N(\mathbb{T}^4) \quad = \quad \mathrm{AdS}_3 \times \mathrm{S}^3 \times \mathbb{T}^4$$

1 unit of NS-NS flux

This background describes a tensionless string theory, where massless higher spin fields are present.

Conclusions and Outlook

$$\mathrm{Sym}_N(\mathbb{T}^4) \quad = \quad \mathrm{AdS}_3 \times \mathrm{S}^3 \times \mathbb{T}^4$$

1 unit of NS-NS flux

Both sides are explicitly solvable and have free field realisations.

This opens the door for all sorts of quantitative tests of the (stringy) duality.

Conclusions and Outlook

$$\mathrm{Sym}_N(\mathbb{T}^4) \quad = \quad \mathrm{AdS}_3 \times \mathrm{S}^3 \times \mathbb{T}^4$$

1 unit of NS-NS flux

The world-sheet theory exhibits signs of a topological string theory:

- only short representations of psu(1,1|2) appear
- correlation functions localise to isolated points

cf [Aharony, David, Gopakumar, Komargodski, Razamat ’07] [Razamat ’08], [Gopakumar ’11], [Gopakumar, Pius ’12]

Future directions

Many directions for future work:

- match null vectors of dual theories

  [Dei, Eberhardt, MRG, in progress]

- understand topological structure directly
- prove by some sort of field redefinition
- Euclidean path integral
- study deformations…

Thank you!