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The holographic dual of a SUSY vector model and tensionless open strings

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Abstract: We describe a supersymmetric example of the holographic duality between 3 dimensional vector models and higher spin gauge theories in AdS_4, first proposed in a bosonic context by Klebanov and Polyakov. We argue that a particular Fradkin-Vasiliev type supersymmetric higher spin gauge theory in AdS_4 with 16 supersymmetries is dual to the singlet sector of bilinear operators of a free $\mathcal{N} = 4$, SU(N) vector model defined on the boundary. Starting from the duality between type-IIB on $AdS_5 \times S^5$ with a D5–brane as an $AdS_4 \times S^2$ subspace and the 4 dimensional SU(N) SYM with a defect, we recover the duality between our vector model and the higher spin gauge theory. In this case, we propose that the higher spin gauge theory is a truncation of the open string theory on the world volume of the D5-brane in its tensionless string limit. We also comment on a possible Higgs mechanism in our model.

Keywords: Superstrings and Heterotic Strings, AdS-CFT and dS-CFT Correspondence, Space-Time Symmetries, Supersymmetry and Duality.

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Contents

1. Introduction

Recently Klebanov and Polyakov [1] have proposed that the holographic dual of the singlet sector of bilinear operators of the bosonic vector models should be certain bosonic Fradkin-Vasiliev [2, 3] type higher spin gauge theories in AdS spaces. In particular, they argued that the 3-dimensional critical $O(N)$ vector model should be dual to the minimal higher spin gauge theory in AdS_4 which contains one massless physical field for each even spin $s = 0, 2, 4, \ldots$ This duality has been motivated by the fact that the free $O(N)$ vector model for the real fields $\phi^a(\vec{x})$ in the vector representation of $O(N)$:

$$S = \frac{1}{2} \int d^3x \, \partial^\mu \phi^a \partial_\mu \phi^a \tag{1.1}$$

has singlet bilinear conserved currents of the form

$$J_{(\mu_1 \cdots \mu_s)} = \phi^a \overset{\leftrightarrow}{\partial}_{(\mu_1} \cdots \overset{\leftrightarrow}{\partial}_{\mu_s)} \phi^a + \cdots \tag{1.2}$$

which are primaries for $s = 0, 2, 4, \ldots$ A prescription, similar to that in the usual AdS/CFT correspondence, to find the correlation functions of these currents from the higher spin theory has also been proposed:

$$\left\langle \exp \int d^3x \, h_0^{(\mu_1 \cdots \mu_s)} J_{(\mu_1 \cdots \mu_s)} \right\rangle = e^{S[h_0]} \tag{1.3}$$

where, as usual, h_0 is the boundary value of the massless spin s field in the bulk and J is the corresponding dual CFT operator. For example, the massless scalar in AdS_4 is proposed to be dual to the spin zero 'current' $\phi^a\phi^a$. The well known infrared conformal fixed point of the 3-dimensional vector model, in this language, is nicely interpreted as the one obtained by deforming (1.1) by a 'double trace' operator $(\phi^a\phi^a)^2$.

It is of interest to understand this duality in more detail. Some of the recent developments in this direction include [4]–[7].

In this paper, we consider supersymmetric versions of this duality. As a particular example, we propose that the non-minimal $\mathcal{N} = 4$ higher spin gauge theory [8] in AdS_4 based on the higher spin superalgebra $hs_0(4|4;1)$ should be holographically dual to the singlet sector of bilinear operators of the free supersymmetric 3-dimensional $\mathcal{N} = 4$ theory with fields in the vector representation of $\mathrm{SU}(N)$.

It has been suggested in the literature that the Klebanov-Polyakov (KP) duality should be realised as a truncation of the duality between M-theory on $AdS_4 \times X^7$ and three dimensional CFT on the boundary in its free field theory limit. However, it is not yet clear if and how the KP type duality can be realised in a string theory context. Therefore, it is interesting to study any possible realisations of such dualities in string theory. Here, we propose an embedding of our supersymmetric version into type-IIB string theory.

We start with the well known correspondence between a 4-d 'defect' conformal field theory and the type-IIB string theory on $AdS_5 \times S^5$ plus a D5–brane with its world-volume embedded as an $AdS_4 \times S^2$ subspace [9, 10]. The D5–brane acts as a co-dimension one defect in the boundary $\mathcal{N} = 4$, $d = 4$, $\mathrm{SU}(N)$ SYM. On the 3 dimensional defect, there is matter in the fundamental (vector) representation of the gauge group $\mathrm{SU}(N)$. We take a decoupling limit, the free field theory limit $\lambda = 4\pi g_{\mathrm{YM}}^2 N \to 0$ of the defect CFT, that separates the 4-d SYM dynamics from that of the 3-d vector model. In this limit the bulk theory becomes that of the tensionless closed and open strings. By taking recourse to a strong form of the AdS/dCFT duality, we argue that the holographic dual of the singlet sector of bilinear operators of the free vector model should be a consistent truncation of the open string theory living on the D5–brane in the tensionless string limit.

The paper is organised as follows. In section 2 we describe a supersymmetric version of the KP duality. In section 3 we deal with the subject of the string theory embedding of the duality of section 2. We also comment on how the higher string modes could become massive through a generalised Higgs mechanism when λ is turned on in the defect CFT. In section 4 we give conclusions and open questions.

2. A supersymmetric version of Klebanov-Polyakov duality

In this section we describe the holographic dual of a supersymmetric Fradkin-Vasiliev type higher spin gauge theory in AdS_4. We start with a brief review of the higher spin gauge theory that we are interested in. For a general introduction to higher spin gauge theories of Fradkin-Vasiliev type, see [11] and for more details of the particular theory being considered here, see [8].

2.1 A non-minimal $\mathcal{N}=4$ higher spin gauge theory in AdS_4.

Let us consider the non-minimal $\mathrm{U}(1)_f$ extended $\mathcal{N}=4$ higher spin gauge theory constructed in [8].[1] This theory is defined in AdS_4 background and is based on the higher spin algebra $hs_0(4|4;1)$. This infinite dimensional superalgebra contains

$$\mathrm{OSp}(4|4) \times \mathrm{U}(1)_f \subset hs_0(4|4;1) \tag{2.1}$$

as its maximal finite dimensional subalgebra. The spectrum of massless[2] physical fields of this theory, arranged into levels of $\mathcal{N}=4$ supermultiplets, is given in the table 1 [8].

The spectrum contains eight fields of each spin $s = 0, 1/2, 1, 3/2, 2, \ldots$ At $l=-1$, we have an $\mathcal{N}=4$ SYM multiplet and at the next level a supergravity multiplet and so on.

This spectrum is identical to the massless spectrum of higher spins obtained as a tensor product [14, 15] of two Dirac singletons [16] of OSp(4|4) superalgebra.

$$\mathrm{Spectrum}[hs_0(4|4;1)] \approx \Xi \otimes \Xi \tag{2.2}$$

where

$$\Xi = D\left(\frac{1}{2}, 0; \frac{1}{2}, 0\right) \oplus D\left(1, \frac{1}{2}; 0, \frac{1}{2}\right) \tag{2.3}$$

is the Dirac singleton. The infinite dimensional UIRs of the higher spin algebra $hs_0(4|4;1)$ are decomposed into those of OSp(4|4) and denoted by $D(\Delta, s; j_1, j_2)$. The quantum numbers are those of the maximal compact subalgebra

$$\mathrm{SO}(2) \times \mathrm{SO}(3) \times \mathrm{SU}(2)_H \times \mathrm{SU}(2)_V \subset \mathrm{OSp}(4|4)\,. \tag{2.4}$$

One can set up the higher spin field equations [8] which we will not reproduce here. At the level $l=-1$, the vector multiplet has six scalars which are the lowest weight components of the UIRs $D(1,0;1/2,0)$ and $D(2,0;0,1/2)$. Similarly, the scalars at the next level are the lowest weight components of $D(1,0;0,0)$ and $D(2,0;0,0)$. The rest of the fields have $\Delta = s+1$ and all the 6s in the table above are made out of $(1,0)\oplus(0,1)$ and all the 4s are made out of $(1/2,0)\oplus(0,1/2)$ of $\mathrm{SU}(2)_H \times \mathrm{SU}(2)_V$. The whole spectrum is neutral under $\mathrm{U}(1)_f$ by construction. Next, we construct a dual vector model.

$(l,j) \backslash s$	0	$\frac{1}{2}$	1	$\frac{3}{2}$	2	$\frac{5}{2}$	3	$\frac{7}{2}$	4	$\cdots$
$(-1,1)$	6	4	1							
$(0,0)$	$1+\bar{1}$	4	6	4	1					
$(0,1)$			1	4	6	4	1			
$(1,0)$					1	4	6	4	1	
$\vdots$						$\vdots$				

Table 1: The higher spin spectrum based on the algebra $hs_0(4|4;1)$.

[1] It is an extension of the minimal $\mathcal{N}=4$ theory of [12, 13, 11].

[2] Recall that in AdS_4 the massless fields have $\Delta = s+1$ for $s \geq \frac{1}{2}$ and $\Delta = 1, 2$ for $s=0$.

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2.2 The dual vector model

We look for a $d = 3$, SU(N) vector model with $\mathcal{N} = 4$ supersymmetry with the following properties. It should admit conserved currents $\mathcal{J}_s^{(i)}(\vec{x})$ (with $i = 1, 2, \ldots 8$ and $s = 0, 1/2, 1, \ldots$) belonging to the singlet sector of SU(N) and are bilinear in the fundamental fields. It should admit $\mathrm{OSp}(4|4) \times \mathrm{U}(1)_f$ symmetry group. The conserved currents $\mathcal{J}_s^{(i)}(\vec{x})$ should have the right quantum numbers of OSp(4|4) and neutral under $\mathrm{U}(1)_f$ with multiplicities as the corresponding higher spin fields listed in the previous subsection.

We have seen that the bulk field content can be obtained as a tensor product of two Dirac singletons (2.3). Therefore, it is natural to choose the field content of the boundary theory to have the same quantum numbers as the lowest weight states of the singleton. That is, a doublet of complex scalar fields $q^m(\vec{x})$, $m = 1, 2$, and a doublet of spinors $\psi^i(\vec{x})$, $i = 1, 2$, both in the fundamental representation of SU(N) with their conjugates transforming in the anti-fundamental. The field content is summarised below:

Field	Δ	Spin (s)	$\mathrm{SU}(2)_H$	$\mathrm{SU}(2)_V$	SU(N)	$\mathrm{U}(1)_f$
q^m	$\frac{1}{2}$	0	$\frac{1}{2}$	0	N	1
ψ^i	1	$\frac{1}{2}$	0	$\frac{1}{2}$	N	1

(2.5)

where m and i are the double indices of $\mathrm{SU}(2)_H$ and $\mathrm{SU}(2)_V$ respectively. A free supersymmetric action with this field content can be written down:

$$S_{\text{boundary}} = \int d^3x \left[(\partial^k q^m)^\dagger \partial_k q^m - i\overline{\psi}^i \rho^k \partial_k \psi^i \right] \tag{2.6}$$

where ρ_k are the Dirac matrices in $d = 3$. This free theory is superconformal and admits $\mathrm{OSp}(4|4) \times \mathrm{U}(1)_f$ as its symmetry group. Eight of its supercharges correspond to the ordinary supersymmetries and the remaining eight generate superconformal supersymmetries.

Following [1] let us construct a set of conserved currents in this theory which are singlet bilinear operators and associate them with the corresponding bulk fields listed in the previous subsection.[3] First, there are eight scalar ($s = 0$) operators:

$$\overline{q}^m \sigma^I_{mn} q^n\,, \qquad \overline{\psi}^i \sigma^A_{ij} \psi^j\,, \qquad \overline{q}^m q^m \quad \text{and} \quad \overline{\psi}^i \psi^i \tag{2.7}$$

where I is a triplet index of $\mathrm{SU}(2)_H$ and A is a triplet index of $\mathrm{SU}(2)_V$. σ^I and σ^A denote the Pauli matrices. The three scalars $\overline{q}^m \sigma^I_{mn} q^n$ ($I = 1, 2, 3$) have $\Delta = 1$ and have the correct quantum numbers to be dual to three of the scalars at level $l = -1$ in table 2.5. Similarly, operators $\overline{\psi}^i \sigma^A_{ij} \psi^j$ correspond to the remaining three scalars in the $\mathcal{N} = 4$ SYM multiplet of the table 2.5. The remaining two operators $\overline{q}^m q^m$ and $\overline{\psi}^i \psi^i$ correspond to the scalars in the $(l, j) = (0, 0)$ multiplet of table 2.5.

At $s = 1/2$, we have

$$\psi^{\star i} q^m \quad \text{and} \quad \overline{q}^m \psi^i \tag{2.8}$$

[3] For details on the subject of higher spin conserved currents in conformal field theories, see refs. [17, 18, 11, 19, 20].

which are again eight in number and have the right quantum numbers to be associated with the eight spin half fields appearing in table 2.5. Let us now turn to the vectors. We find the following eight conserved currents with $s = 1$:

$$\overline{q}^m \overleftrightarrow{\partial}^k q^m \,, \qquad \overline{\psi}^i \rho^k \psi^i \,, \qquad \overline{q}^m \sigma^I_{mn} \overleftrightarrow{\partial}^k q^n \quad \text{and} \quad \overline{\psi}^i \sigma^A_{ij} \rho^k \psi^j \,. \tag{2.9}$$

We associate these currents with the eight vector fields in table 2.5 with the appropriate quantum numbers.

In general, at any given integer spin $s \geq 0$ we can construct eight conserved currents of the form

$$\mathcal{J}^\alpha_s(\vec{x}) = \overline{q}^m \sigma^\alpha_{mn} \overleftrightarrow{\partial}^{(k_1} \overleftrightarrow{\partial}^{k_2} \cdots \overleftrightarrow{\partial}^{k_s)} q^n + \cdots \,, \tag{2.10}$$

$$\mathcal{J}^\beta_s(\vec{x}) = \overline{\psi}^i \sigma^\beta_{ij} \rho^{(k_1} \overleftrightarrow{\partial}^{k_2} \cdots \overleftrightarrow{\partial}^{k_s)} \psi^j + \cdots \tag{2.11}$$

where $\alpha, \beta = 0, 1, 2, 3$ with $\sigma^0 = I_{2\times 2}$. Similarly, there are eight conserved currents for each half-integral spin:

$$\mathcal{J}^{im}_s = \psi^{\star i} \overleftrightarrow{\partial}^{(k_1} \overleftrightarrow{\partial}^{k_2} \cdots \overleftrightarrow{\partial}^{k_s)} q^m + \cdots \,, \tag{2.12}$$

$$\overline{\mathcal{J}}^{im}_s = \overline{q}^m \overleftrightarrow{\partial}^{(k_1} \overleftrightarrow{\partial}^{k_2} \cdots \overleftrightarrow{\partial}^{k_s)} \psi^i + \cdots \,. \tag{2.13}$$

Notice that these currents have the right quantum numbers of OSp(4|4) and are neutral under $\mathrm{U}(1)_f$ as required. For more complete expressions of these currents please see the appendix.[4]

Thus we conclude that our free $\mathcal{N} = 4$ SU(N) vector model (2.6) admits the right set of bilinear operators in its singlet sector that are in one to one correspondence with the fields of the higher spin gauge theory of the previous subsection. Of course, one needs to arrange these currents into supermultiplets which may involve taking various linear combinations of the basic currents listed above.

One can easily generalise the precription eq. (1.3) of [1] to evaluate the boundary correlators using the bulk higher spin theory to our supersymmetric context. The procedure followed here can clearly be generalised to construct the dual of any given higher spin gauge theory whose field content is contained in the tensor product of its Dirac singletons.

Let us now turn to finding an embedding of the model studied here into a string theory context.

3. Vector models and tensionless open strings

It has been conjectured in the literature that there exists an unbroken symmetric phase of M-theory on $AdS_4 \times X^7$ backgrounds dual to an SU(N) invariant free singleton theory on the boundary. To incorporate the Klebanov-Polyakov type dualities into this set up it has been suggested (see, for instance, [21] and references there in) that a massless sector

[4] Notice that all the higher spin fields can be packaged into the following set of bi-local collective fields in the sense of refs [7]: $\overline{q}^m(\vec{x})\sigma^\alpha_{mn}q^n(\vec{y})$, $\overline{\psi}^i(\vec{x})\sigma^\beta_{ij}\rho^\mu\psi^j(\vec{y})$, $\psi^{\star i}(\vec{x})q^m(\vec{y})$ and $\overline{q}^m(\vec{x})\psi^i(\vec{y})$ with $\rho^\mu = \{I_{2\times 2}, \rho^k : k = 0, 1, 2\}$.

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of M-theory on $AdS_4 \times X^7$ should be dual to the singlet sector of $O(N^2-1)$ vector model. In what follows, we argue that a more natural set up for realising a Klebanov-Polyakov type duality is type-II string theories in the presence of D-branes.[5] We concentrate on the particular example discussed in the previous section.

As was outlined in the introduction, the set-up of interest is a Karch-Randall compactification [9] of type-IIB string theory. That is, we embed a D5–brane into the $AdS_5 \times S^5$ background as an $AdS_4 \times S^2$ subspace. This theory is expected to be dual to a 'defect' CFT on the boundary generalising the usual AdS/CFT correspondence [26]. The details of this duality along with a perturbative definition of the the dual field theory have been worked out in [10] (See also [27]). We are going to use this duality in an essential way. So let us first review briefly a few relevant details of this duality.

3.1 The bulk theory

The bulk side of this set up contains type-IIB closed strings propagating in the maximally supersymmetric $AdS_5 \times S^5$ background along with the D5–brane open strings propagating in the $AdS_4 \times S^2$ subspace. This system can be obtained by starting with N D3–branes along (x^1, x^2, x^6) directions plus one D5–brane along $(x^1, x^2, x^3, x^4, x^5)$ directions (with both D3 and D5–branes sitting at the same point in (x^7, x^8, x^9) directions) and then taking the near horizon limit of the D3–branes. Before taking the large-N limit, one has closed strings along with open strings in the Chan-Paton sectors 3-3, 3-5 and 5-5. The 3-3 sector fields are in the adjoint representation of $\mathrm{U}(N)$, 3-5 strings in the bifundamental representation of $\mathrm{U}(N) \times \mathrm{U}(1)$ and the 5-5 strings contain a U(1) gauge field. Taking the near horizon limit separates the 3-3 and 3-5 strings from the 5-5 and closed strings.

The isometry group of the resultant background is

$$\mathrm{SO}(2,3) \times \mathrm{SO}(3) \times \mathrm{SO}(3) \sim \mathrm{Sp}(4) \times \mathrm{SU}(2)_H \times \mathrm{SU}(2)_V \subset \mathrm{SO}(2,4) \times \mathrm{SO}(6)\,. \tag{3.1}$$

The latter is the isometry group of $AdS_5 \times S^5$ without the D5–brane. The presence of the D5–brane breaks half of the 32 supersymmetries of $AdS_5 \times S^5$ and we have

$$\mathrm{OSp}(4|4) \subset \mathrm{SU}(2,2|4) \tag{3.2}$$

as the unbroken supergroup. There is also a $\mathrm{U}(1)_f$ gauge field on the D5–brane. The low energy description of this system is studied in [10] and the modes are classified according to the representations of OSp(4|4). Since the supergravity and the DBI modes are expected to survive the tensionless string limit, they will be useful for our purposes later on. Let us review the relevant open string results found in [10] here. For more details see section 3.2 of ref. [10].

There are three sectors for the light 5-5 open string modes: (i) The angular fluctuation (ψ) of the D5–brane on S^5, (ii) The gauge field fluctuation (b_μ) of the D5–brane along the AdS_4 directions and (iii) two coupled sectors $(b+z)^{(\pm)}$ of gauge field components along S^2 and the remaining transverse fluctuations. Below, we summarise the spectrum of

[5] For discussions on higher spin theories from type-II closed strings, see [22, 23, 19, 24, 25].

these fluctuations in terms of their AdS_4 energies (the conformal dimensions of the dual operators):

$$
\begin{aligned}
&\text{(i)} \quad \Delta_+ = 2 + l\,, \qquad \Delta_- = 1 - l\,,\\
&\text{(ii)} \quad \Delta = 2 + l\,,\\
&\text{(iii)} \quad \Delta_\pm^{(+)} = \frac{3}{2} \pm \frac{1}{2}(2l+5)\,, \qquad \Delta_\pm^{(-)} = \frac{3}{2} \pm \frac{1}{2}|2l-3|\,.
\end{aligned} \tag{3.3}
$$

Here l labels the spherical harmonics over the S^2 that the D5–brane wraps, whose symmetry we denote by $\mathrm{SU}(2)_H$. The masses of these modes are related to their conformal dimensions by the usual relations. Notice that Δ_- is allowed only for $l = 0$, $\Delta_-^{(+)}$ is not possible at all and $\Delta_-^{(-)}$ is allowed only for $l = 1, 2$ by unitarity.[6] Under the $\mathrm{SU}(2)_V$ symmetry the vector b_μ is a singlet, the scalars ψ and $(b+z)^{(-)}$ are a triplet and a singlet respectively.

Further, for $l = 0$, the sectors (i) and (ii) contain three massless scalars (a triplet of $\mathrm{SU}(2)_V$) and a massless vector respectively. Similarly, the $\Delta_\pm^{(-)}$ sector contains another three massless scalars at $l = 1$ (a triplet of $\mathrm{SU}(2)_H$). There are also closed string supergravity modes which we do not describe here. Beyond the low energy limit, there will be higher stringy modes as well.

3.2 The boundary theory

Now, let us briefly review the relevant information about the boundary theory dual to the system described above. The intersection of the D5–brane on the boundary of $AdS_5 \times S^5$ is a co-dimension one defect in the 4-d CFT. The defect carries, on its world volume, a 3-dim vector model with one hypermultiplet in the fundamental representation of the gauge group.

Therefore, the boundary theory is a 'defect' CFT (or simply a dCFT). This is expected to be a superconformal field theory with 16 supersymmetries. The symmetry group is again the superconformal group $\mathrm{OSp}(4|4)$ of which the $\mathrm{SO}(2,3)$ is the conformal group and the $\mathrm{SO}(4) = \mathrm{SU}(2)_H \times \mathrm{SU}(2)_V$ is the global R-symmetry. There is also a $\mathrm{U}(1)_f$ global symmetry under which the defect fields are charged.[7] A lagrangian formulation of this theory is described in [10]. Here we will be content with describing the field content of the theory and refer the reader to the original work for more details. The 4-d ambient CFT (which we refer to as aFT) fields are the vector A_k, A_6 for $k = 0,1,2$, the scalars X_H^I and X_V^A where $I = 3,4,5$ and $A = 7,8,9$ and the Majorana spinors λ_{im} where i and m are the doublet indices of $\mathrm{SU}(2)_V$ and $\mathrm{SU}(2)_H$ respectively. All these fields are in the adjoint representation of the gauge group $\mathrm{SU}(N)$. Further, we have the fields on the defect: the 3–d scalars q^m and the fermions ψ^i which are doublets of $\mathrm{SU}(2)_H$ and $\mathrm{SU}(2)_V$ respectively. Both these fields are in the fundamental representation of the gauge group $\mathrm{SU}(N)$. Notice that the quantum numbers of these fields are the same as the ones given in eq. (2.5).

[6] Recall that the unitarity in AdS_4 requires the $\Delta \geq s + \frac{1}{2}$ for $s = 0, 1/2$ and $\Delta \geq s+1$ for $s = 1, 3/2, 2, \ldots$

[7] This is the remnant of the fact that the defect fields belonged to the sector of strings between D3 branes and a D5-brane. The $\mathrm{SU}(N)$ is a gauge symmetry but $\mathrm{U}(1)_f$ is only a global symmetry after taking the near horizon limit.

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The operators dual to the light open string modes described in the previous subsection are to be identified with gauge invariant single trace operators made of the fundamental fields of the field theory on the defect (which we refer to as dFT) and the adjoint fields of aFT restricted to the defect.

A dictionary of such identifications can be found in section 5 of ref. [10] which we reproduce below for later reference.

Mode	Δ	$\mathrm{SU}(2)_H$	$\mathrm{SU}(2)_V$	Operator
b_μ	$l+2$	$l \geq 0$	0	$i\overline{q}^m \overleftrightarrow{D}^k q^m + \overline{\psi}^i \rho^k \psi^i$
ψ	$l+2,\, 1-l$	$l \geq 0$	1	$\overline{\psi}_i \sigma^A_{ij} \psi_j + 2\overline{q}^m X^A_V q^m$
$(b+z)^{(-)}$	$l,\, 3-l$	$l \geq 1$	0	$\overline{q}^m \sigma^I_{mn} q^n$

(3.4)

where, in the last column, we indicated the dual operators corresponding to the lowest l modes which form a $\mathcal{N}=4$, $d=4$, U(1) SYM multiplet as expected from the supersymmetry counting.[8] The dual operators for higher l modes are also proposed in [10]. For example:

$$C^{I_0 \cdots I_l}_{l+1} = \overline{q}^m \sigma^{(I_0}_{mn} X^{I_1}_H \cdots X^{I_l)}_H q^n - \text{traces} \tag{3.5}$$

gives the operator dual to a mode with the $\mathrm{SU}(2)_H$ quantum number $l+1$ in the last row of (3.4).

The boundary theory is argued to be an exact CFT [10, 27]. Therefore, one expects that the $AdS_4 \times S^2$ embedding of the D5–brane is a solution to the string equations of motion. We are going to assume that there exists an exact world sheet boundary conformal field theory of this D5–brane in $AdS_5 \times S^5$ background. We will also assume that the duality is valid in its strongest form for all values of the 't Hooft coupling λ and N and treat the field theory as a definition of the dual string theory whenever necessary.

Having identified a type-IIB string theory context in which the 3-dimensional field content (2.5) of the supersymmetric vector model of section 2 is realised, we next aim to carefully extract the vector model from the boudary theory and obtain its dual bulk theory using the AdS/dCFT correspondence reviewed above.

3.3 Super KP from AdS/dCFT

To summarise, on the bulk side of the duality there is the theory of type-IIB closed strings in $AdS_5 \times S^5$ coupled to the open strings on a D5–brane embedded as an $AdS_4 \times S^2$ subspace. On the boundary, we have the defect CFT. Holography tells us that these two systems describe the same physics. Therefore, we write

$$\begin{aligned} &\mathcal{L}^{(\mathrm{Bulk})}_{IIB\ \mathrm{Closed\ Strings}} + \mathcal{L}^{(\mathrm{Bulk})}_{D5\text{-}\mathrm{Open\ Strings}} + \mathcal{L}^{(\mathrm{Bulk})}_{\mathrm{Int.}} \approx \\ &\qquad \approx \mathcal{L}^{(Bdy)}_{d=4,\mathcal{N}=4,\mathrm{SU}(N),SYM} + \mathcal{L}^{(Bdy)}_{d=3,\mathcal{N}=4,\mathrm{SU}(N)\ \mathrm{Vector\ Model}} + \mathcal{L}^{(Bdy)}_{\mathrm{Int.}} \end{aligned} \tag{3.6}$$

where we have denoted the theories by their lagrangian densities schematically. On the right hand side, the first two terms represent aFT and dFT respectively and the third denotes the interactions between them.

[8]The mode $(b+z)^{(+)}$ is not listed above as it is always massive and drops out of our story.

Since we are interested in singling out the vector model, we first need to decouple both aFT and dFT. Therefore, we seek a limit of the boundary theory in which the term $\mathcal{L}^{(Bdy)}_{Int.}$ can be dropped. This indeed happens in the limit $\lambda \to 0$ where $\lambda = 4\pi g^2_{\rm YM} N$ is the 't Hooft coupling on the boundary[9] and $\lambda = R^4/\alpha'^2$ in the bulk. In this limit, both $\mathcal{L}^{(Bdy)}_{4-d,\mathcal{N}=4,\,{\rm SU}(N),\,SYM}$ and $\mathcal{L}^{(Bdy)}_{3-d,\mathcal{N}=4,\,{\rm SU}(N)\,Vector\ \ Model}$ become free field theories and the interaction lagrangian $\mathcal{L}^{(Bdy)}_{Int.}$ goes to zero. On the bulk side, the limit $\lambda \to 0$ is the tensionless string limit.[10] Therefore, in this limit, one expects that the whole string spectrum to become massless with infinitely many states at each spin.

Now that we have decoupled the aFT and dFT, we can in principle, extract the holographic dual of the vector model of dFT alone. But unfortunately we do not have much handle over the bulk side of the story. Therefore, we take recourse to indirect arguments based on the assumption that the strong form of the AdS/dCFT is true.

First we restrict ourselves to single trace operators of the boundary theory. The calculable quantities of interest in the CFT are the correlation functions of gauge invariant operators. In the limit $\lambda \to 0$, these operators are simply the singlets of SU(N). There are essentially three classes of these singlet operators. Denoting an arbitrary adjoint field of aFT by $\mathcal{A}$ and a fundamental field of dFT by $\mathcal{F}$ these operators, schematically, are of the type:

$$(a) \quad {\rm Tr}(\mathcal{A}\cdots\mathcal{A})\,, \qquad (b) \quad \overline{\mathcal{F}}\mathcal{A}\cdots\mathcal{A}\mathcal{F}\,, \qquad (c) \quad \overline{\mathcal{F}}\mathcal{F}\,. \tag{3.7}$$

The operators of type (a) correspond to the closed string states. The operators of type (b) and (c) (in general, combinations of those, like in eq. (3.4) correspond to the open string states.[11]

We are interested in the vector model and therefore in the correlation functions involving type (c) operators alone. Hence we need to 'subtract' the aFT from the boundary theory. At the level of correlation functions, this amounts to setting all the operators belonging to type (a) and (b) to zero which is achieved by simply setting the fields of the aFT to zero. This is consistent because we have decoupled aFT and dFT completely.

By 'subtracting' the aFT from the boundary theory, we are left with a much reduced set of operators (those of type (c)) on the boundary theory. Let us try to infer what this truncation of the set of operators of the boundary theory means for the bulk theory. Using the strong form of the AdS/dCFT correspondence, we conclude that setting type (a) operators to zero amounts to setting all the dual closed string states on $AdS_5 \times S^5$ to zero. Similarly, setting type (b) operators to zero corresponds to truncating the spectrum of tensionless open string states consistently.

More specifically, out of the spectrum of the light open string modes discussed in section 3.1 (all of which survive the tensionless string limit) we will be left with only the

[9] For this limit we keep N to be large but finite and take $g^2_{\rm YM} \to 0$.

[10] In a strict large-N limit one expects that $\mathcal{L}^{(Bulk)}_{Int.}$ also should go to zero. But we will be interested in keeping higher orders in $1/N$.

[11] There is, in general, operator mixing with multitrace operators which is crucial at higher orders in perturbation theory. Here, we mean the operators which reduce to the ones listed in (3.7) at the lowest order in the perturbation theory.

lowest l modes after the truncation. This is because the operators dual to any higher l mode contains powers of adjoint fields of aFT (like in (3.5)) which we set to zero as they belong to type (b) of (3.7). Thus, we get rid of all the higher KK modes of these light states in this truncation. About the higher string modes we again go back to the dual theory. For this we simply note that the free field theory that we obtain in the limit is exactly the one we considered in section 2. There, we constructed all the operators that could be dual to the string states that we are after. Therefore, we expect that the truncated spectrum of the tensionless open strings on $AdS_4 \times S^2$ should be identical to the spectrum of the higher spin gauge theory on AdS_4 discussed in section 2.

In other words, the higher spin gauge theory considered in section 2 is precisely the consistent truncation of the tensionless open string theory. The truncation is achieved by subtracting the $d = 4$ aFT from the AdS/dCFT correspondence that we started with. The symmetry group OSp(4|4) is now simply the superconformal group of the 3-d vector model and the $\mathrm{U}(1)_f$ is now associated with the U(1) gauge field on the D5–brane. The neutrality of the higher spin spetrum of section 2 under $\mathrm{U}(1)_f$ is simply the reflection of the fact that the open string states on a single D-brane are neutral under its U(1) gauge field. However, if we had more than one coincident D5-branes, the corresponding higher spin spectrum would have been charged under the the non-abelian gauge field on the branes.

It may be puzzling that the gravity theory expected for holography is now actually part of a tensionless open string theory living on a D5-brane. However, one expects a graviton (a massless spin two particle) in the open string spectrum in the tensionless string limit. One way of arguing for this is to look at the conserved energy-momentum tensor of the dual boundary CFT [10]:

$$T_{\mu\nu}(\vec{y}, x) = \mathcal{T}_{\mu\nu}(\vec{y}, x) + \delta(x)\delta^k_\mu \delta^l_\nu t_{kl}(\vec{y}) \tag{3.8}$$

where, $\vec{y}$ and x are the coordinates along and transverse to the defect. Neither $\mathcal{T}_{\mu\nu}$ nor t_{kl} is conserved by itself in the full theory. In the free CFT limit, however, they are both conserved individually. Since in this limit t_{kl} is quadratic in the fields of the dFT, it corresponds to an open string mode. Hence, we expect that one of the stringy modes of the 5-5 sector to become massless exactly in this limit and play the role of the graviton dual to this conserved energy momentum tensor.[12] Notice that there are eight spin 2 fields in the higher spin spectrum of table 2.5. But we expect that the actual graviton dual to the 3-d energy momentum tensor is a combination of the spin 2 states in the second and fourth rows of table 2.5 as the others do not have the same global quantum numbers (j_1, j_2) as t_{kl}.

Before concluding this section, we comment on the possibility of understanding how the higher open string states of the higher spin theory could attain masses once we go away from the decoupling limit discussed above.

3.4 The open string graviton and the higgs mechanism

For consistency we expect that, when λ is turned on, all the fields except for the massless $\mathcal{N} = 4$ SYM multiplet (at level $l = -1$ in table 2.5) acquire masses proportional to the string tension. From the boundary theory point of view, this means that the operators

[12] This was also noticed recently in [28].

dual to these states develop anomalous dimensions

$$\Delta - s - 1 \sim \mathcal{O}(\lambda)\,. \tag{3.9}$$

In the full interacting dCFT, one can see this happening in perturbation theory [10, 28]. From the bulk theory point of view, once we turn on the tension of the strings, we can no longer neglect interactions with closed string modes (and the rest of the open string modes which we set to zero by the consistent truncation explained above). But it may be possible to argue that the extra massless modes acquire mass for nonzero λ in the higher spin theory in AdS_4. A related issue has been considered in the literature [29] (see also [30]) for $s = 2$.[13] In ref. [29] it has been argued that depending on the boundary conditions set for the massless scalar, the graviton goes through a Higgs mechanism to acquire mass. By simply noting that the required mixing of boundary conditions occurs precisely when the defect interacts with the ambient CFT, we can take over the results of [29] into our context.

Massless UIRs of AdS_4 group with spin $s > 0$ have $\Delta = s + 1$ and massive ones have $\Delta > s + 1$. A massive UIR $D(\Delta,\, s;\, j_1,\, j_2)$ becomes reducible in the massless limit $\Delta \to s + 1$. For example, for a spin s massive UIR in AdS_4 denoted by $D(\Delta, s)$ one has the decomposition

$$D(\Delta, s) \to D(s + 1, s) \oplus D(s + 2, s - 1) \tag{3.10}$$

in the limit $\Delta \to s + 1$. The first piece on the r.h.s is massless and the second is a massive representation. This implies that in AdS_4 a spin s field must eat a massive particle of spin $s - 1$ to attain mass. Generalising this result to our case, the decomposition of a massive representation $D(\Delta, s, j_1, j_2)$ in the limit of $\Delta \to s + 1$ may be written as:

$$D(\Delta, s, j_1, j_2) \to D(s + 1, s, j_1, j_2) \oplus D(s + 2, s - 1, j_1, j_2)\,. \tag{3.11}$$

This gives the quantum numbers required for the Higgs fields. One should be able to find these representations in the tensor product of the scalar field representations of table 2.5 as explained in [5]. We postpone a more detailed analysis to the future.

4. Conclusion

We have provided a supersymmetric example of the holographic duality between vector models and higher spin gauge theories first considered by Klebanov and Polyakov [1] in a bosonic context. Our bulk theory is a particular non-minimal $\mathcal{N} = 4$ higher spin gauge theory in AdS_4 and based on the algebra $hs_0(4|4;1) \supset \mathrm{OSp}(4|4) \times \mathrm{U}(1)_f$. It contains eight physical fields at each spin $s = 0, 1/2, 1, 3/2, 2, \ldots$ The holographic boundary theory is a 3-dimensional $\mathcal{N} = 4$ $\mathrm{SU}(N)$ free vector model. Since our method is based on constructing

[13] There is, however, an essential difference between [29] and our case. In [29], Higgs mechanism for the the bulk (closed string) graviton was considered (*à la* Karch-Randall brane-world [9]). We are interested in the open string graviton.

the field theory using the singleton representations of the higher spin algebra, it should help constructing other examples where the higher spin field content is contained in the tensor product of the corresponding Dirac singleton.

We have suggested that the model studied here can be embedded into type-IIB string theory. We have argued that the bulk theory should be a truncation of the open string theory on a D5 brane, embedded as $AdS_4 \times S^2$ subspace of $AdS_5 \times S^5$ background of type-IIB, in its tensionless string limit. The truncation is understood from the boundary theory point of view simply as the 'subtraction' of the free aFT.

It is amusing that a large-N CFT is dual to an open string theory as opposed to a closed string theory. The role of the graviton is expected to be played by a massive open string mode which becomes massless in the tensionless string limit. One could consider variants of the model described in this paper. For example, taking an orientifold of the model considered here, where one replaces $AdS_5 \times S^5$ by $AdS_5 \times RP^5$, should result in an SO(N) version of our duality. One can use the notion of consistent truncation to get to the minimal bosonic higher spin gauge theory based on $hs(4)$ algebra as explained in [8]. Similarly, with an SO(N) analogue, one would end up with the proposal of Klebanov and Polyakov itself via consistent trucation.

One could have started with more than one (say M) coincident D5–branes in $AdS_5 \times S^5$. This should lead to a U(M) valued higher spin spectrum. Such higher spin theories have been considered by Vasiliev [11].

It would be interesting to study the tensionless limit of open string theory side explicitly based on the recent results [31]–[34]. Also, it may be possible to understand the higher spin spectrum considered here as that of a 'brane' in an appropriate higher spin gauge theory in AdS_5. The spectrum on AdS_4 subspace may perhaps be thought of as the goldstone and goldstino modes of broken higher spin symmetries of a corresponding AdS_5 theory. Related ideas have been considered recently in [35]. We note that we only tested the duality that we proposed here at the level of matching the spectra. However, the way we 'derived' it from a string theory context suggests that even the interactions of the higher spin theory considered should be reproducible from the dual modes. It will of course be interesting to check this explicitly. One may also be able to define the bulk theory in terms of the bilocal collective fields (see the footnote 4) as in [7]. Finally, we have not considered the possible 'double trace' deformations of our free theory. The allowed set of boundary conditions for the scalars in the bulk suggests the existence of other (interacting) CFTs [36]. It should be interesting to explore that aspect as well. We hope to return to some of the open questions in the future.

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A. The conserved currents

In this appendix we provide more detailed expressions of the conserved currents listed in section 2. We closely follow refs. [17, 18, 11, 19, 20].

We seek conserved tensor primary current of the free field theory of a complex scalar and a spinor in $d = 3$ given in (2.6). More complete expressions of the ones indicated in section 2 are as follows:

$$\left(\mathcal{J}_s^{(\overline{q}q)}\right)^{\alpha} = \sum_{k=0}^{s} \frac{(-1)^k(\partial_{i_1}\cdots\partial_{i_k}\overline{q}^m)\sigma^{\alpha}_{mn}(\partial_{i_{k+1}}\cdots\partial_{i_s}q^n)}{\Gamma(k+1)\Gamma(s-k+1)\Gamma(k+\frac{1}{2})\Gamma(s-k+\frac{1}{2})} - \text{traces}\,, \tag{A.1}$$

$$\left(\mathcal{J}_s^{(\overline{\psi}\psi)}\right)^{\beta} = \sum_{k=0}^{s-1} \frac{(-)^k(\partial_{i_1}\cdots\partial_{i_k}\overline{\psi}^i)\sigma^{\beta}_{ij}\rho_{i_{k+1}}(\partial_{i_{k+2}}\cdots\partial_{i_s}\psi^j)}{\Gamma(k+1)\Gamma(s-k)\Gamma(k+\frac{3}{2})\Gamma(s-k+\frac{1}{2})} - \text{traces}\,. \tag{A.2}$$

In the above expressions $s = 0, 1, 2, \ldots$ and all the indices $i_1, \ldots, i_s$ are understood to be completely symmetrised. Further, as stated earlier, the indices α, β take four values $\alpha, \beta = 0, 1, 2, 4$ and the corresponding $\sigma^{\alpha} = \{I, \sigma^i\,:\; i = 1, 2, 3\}$ with I being the 2×2 identity matrix and σ^i being the SU(2) Pauli matrices. For each spin s, we have a total of eight currents. Similarly, the half-integer spin currents are given as follows:

$$\begin{aligned}
\left(\mathcal{J}_s^{(\psi^{\star}q)}\right)^{m}_{i} = \sum_{l,n,p=0}^{\infty} & \frac{(-1)^l\Gamma(p+l+n+\frac{3}{2})}{\Gamma(l+1)\Gamma(n+1)\Gamma(p+1)2^p\Gamma(p+l+\frac{3}{2})\Gamma(p+n+\frac{3}{2})} \times \\
& \times \Bigg[\left(p+n+\frac{1}{2}\right)\delta_{s-l-n-2p-\frac{1}{2}}\eta_{i_1 i_2}\cdots\eta_{i_{2p-1}i_{2p}}\partial_{i_{2p+1}}\cdots\partial_{i_{2p+l}}\partial^{j_1}\cdots\partial^{j_p}\psi^{\star i} \times \\
& \quad \times \partial_{i_{2p+l+1}}\cdots\partial_{i_{p+l+n}}\partial_{j_1}\cdots\partial_{j_p}q^m - \\
& \quad - \frac{1}{2}\delta_{s-l-n-2p-\frac{3}{2}}\eta_{i_1 i_2}\cdots\eta_{i_{2p-1}i_{2p}}\rho_{i_{2p+1}}\partial_{i_{2p+2}}\cdots\partial_{i_{2p+l+1}}\partial^{j_1}\cdots\partial^{j_p}\rho_{i_{2p+l+2}}\psi^{\star i} \times \\
& \quad \times \partial_{i_{2p+l+3}}\cdots\partial_{i_{2p+l+n+2}}\partial_{j_1}\cdots\partial_{j_{p+1}}q^m\Bigg]
\end{aligned} \tag{A.3}$$

and finally there is another set of currents obtained by $\psi^{\star} \leftrightarrow \overline{q}$ and $q \leftrightarrow \psi$ in the expression (A.3). Again there are $4 + 4 = 8$ currents with half odd-integral spins as required for the model considered in section 2.

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