Giant Gravitons on $AdS_4 \times \mathbb{C}\mathrm{P}^3$ and their Holographic Three-point Functions

S. Hirano[a], C. Kristjansen[b], D. Young[b]

[a] *Department of Physics, Nagoya University, Nagoya 464-8602, Japan*
[b] *Niels Bohr Institute, Blegdamsvej 17, DK-2100 Copenhagen, Denmark*

Abstract

We find a simple parametrization of the anti-symmetric giant graviton in $AdS_4 \times \mathbb{C}\mathrm{P}^3$, first constructed in [1], dual to the anti-symmetric Schur polynomial involving two bi-fundamental complex scalar fields of ABJM theory. Using this parametrization we evaluate in a semi-classical approach the three-point function of two such giant gravitons and one point-like graviton considering both extremal and non-extremal configurations. We likewise discuss the case of the symmetric giant graviton in $AdS_4 \times \mathbb{C}\mathrm{P}^3$. Finally, we provide an expression for the planar three-point function of chiral primary operators in ABJM at strong coupling and find that the results for the giant graviton three-point functions reduce to this expression in the point-like limit.

1 Introduction

Giant gravitons constitute an important entry in the AdS/CFT dictionary. In the gravity language giant gravitons represent extended higher dimensional objects, D- or M-branes, while in the field theory language they correspond to operators carrying higher representations of the gauge group. In particular, the latter characterization imply that giant gravitons encode information about finite-N gauge theory.

In the $AdS_5 \times S^5$ case [2] one has a simple and beautiful description of 1/2 BPS giant gravitons. On the string theory side the giant gravitons are D3-branes which wrap an S^3 inside either S^5 or AdS_5 while moving on a circle of S^5 with a fixed angular momentum [3–6]. The gauge theory dual of the S^5 giant graviton is the completely anti-symmetric Schur polynomial of a single complex scalar while the dual of the AdS_5 giant is the completely symmetric Schur polynomial [6, 7].

For the $AdS_4 \times \mathbb{CP}^3$ case [8] the situation is slightly more complicated. The simplest possible Schur polynomials are constructed out of two complex bi-fundamental scalars, see [9, 10] for a discussion of these. The gravity dual of the completely symmetric Schur polynomial is a D2-brane which wraps an S^2 inside AdS_4 and (after uplift to M-theory) rotates along a great circle of S^7 orthogonal to the compactification circle [11–13]. The gravity dual of the anti-symmetric Schur polynomial can be described in M-theory language as an M5-brane which wraps two S^5's intersecting at an S^3, all inside S^7, and which like its symmetric cousin rotates along a circle orthogonal to the compactification circle. Its maximal version was discussed in [11, 14], see also [15–17], but the general solution was first constructed in [1].

In the present letter we find an improved parametrization of the anti-symmetric giant graviton of $AdS_4 \times \mathbb{CP}^3$. This greatly simplified parametrization allows us to calculate analytically the three-point function involving two such giant gravitons and one point-like graviton in the holographic approach suggested for strings in [18–22] and generalized to branes in [23, 24]. Unlike what is the case for $\mathcal{N} = 4$ SYM, in ABJM theory three-point functions of 1/2 BPS chiral primary operators are not protected. The chiral primary operators are built from pairs of the four bi-fundamental complex scalar fields W_I of the ABJM theory and their conjugates $\bar{W}^I$ and are given by

$$\mathcal{O}_A = \frac{(4\pi)^{J/2}}{\sqrt{J/2}\,\lambda^{J/2}} (\mathcal{C}_A)^{I_1 \ldots I_{J/2}}_{K_1 \ldots K_{J/2}} \operatorname{Tr}\left(W_{I_1} \bar{W}^{K_1} \cdots W_{I_{J/2}} \bar{W}^{K_{J/2}} \right), \tag{1}$$

where $\lambda = N/k$ is the 't Hooft coupling of ABJM, and $\mathcal{C}_A$ is completely symmetric in upper and (independently) in lower indices, while the trace taken on any pair consisting of one upper and one lower index vanishes. The tensors are orthonormal, so that

$$(\mathcal{C}_A)^{I_1 \ldots I_{J/2}}_{K_1 \ldots K_{J/2}} (\bar{\mathcal{C}}_B)^{K_1 \ldots K_{J/2}}_{I_1 \ldots I_{J/2}} = \delta_{AB}, \tag{2}$$

and the two-point function is protected and given by $\langle \mathcal{O}_A(x) \bar{\mathcal{O}}_B(0) \rangle = \delta_{AB}/|x|^J$. The three-point function structure constants C_{123} are then defined as

$$\langle \mathcal{O}_1(x_1) \mathcal{O}_2(x_2) \mathcal{O}_3(x_3) \rangle = \frac{C_{123}}{|x_1 - x_2|^{\gamma_3} |x_2 - x_3|^{\gamma_1} |x_3 - x_1|^{\gamma_2}}, \tag{3}$$

where $\gamma_i = (\sum_j J_j - 2J_i)/2$. Using the fact that the chiral primary operators are dual to point-like gravitons, the structure constants were calculated at strong coupling and large-N long ago [25] using M-theory on $AdS_4 \times S^7$ and so corresponding to the case of ABJM with Chern-Simons level $k = 1$. This expression may then be generalized for arbitrary k, see appendix A. The result is (we take $J_3 \geq J_2 \geq J_1$)

$$C_{123}^{\lambda \gg 1} = \frac{1}{N} \left(\frac{\lambda}{2\pi^2} \right)^{1/4} \frac{\prod_{i=1}^{3} \sqrt{J_i + 1}\, (J_i/2)!\, \Gamma(\gamma_i/2 + 1)}{\Gamma(\gamma/2 + 1)} \sum_{p=0}^{\gamma_3} \frac{(\mathcal{C}_1)^{I_1 \ldots I_p I_{p+1} \ldots I_{J_1/2}}_{K_1 \ldots K_{\gamma_3 - p} K_{\gamma_3 - p+1} \ldots K_{J_1/2}} (\mathcal{C}_2)^{K_1 \ldots K_{\gamma_3 - p} L_1 \ldots L_{\gamma_1 - J_2/2 + p}}_{I_1 \ldots I_p M_1 \ldots M_{J_2/2 - p}} (\mathcal{C}_3)^{K_{\gamma_3 - p+1} \ldots K_{J_1/2} M_1 \ldots M_{J_2/2 - p}}_{I_{p+1} \ldots I_{J_1/2} L_1 \ldots L_{\gamma_1 - J_2/2 + p}}}{p!(\gamma_3 - p)!(\gamma_1 - J_2/2 + p)!(J_2/2 - p)!(\gamma_2 - J_1/2 + p)!(J_1/2 - p)!}. \tag{4}$$

where $\gamma = \gamma_1 + \gamma_2 + \gamma_3$. In contrast to the $\mathcal{N} = 4$ SYM case, we see that not only is there a λ-dependence, but also that there is a range of contractions of the $\mathcal{C}$ tensors. This freedom amounts to the number p of upper indices in $\mathcal{C}_1$ which are contracted with lower indices in $\mathcal{C}_2$. It is instructive to compare this result to the tree-level perturbative result, where to the leading order in $1/N$, only one such contraction can appear[1], which we denote by $\langle \mathcal{C}_1 \mathcal{C}_2 \mathcal{C}_3 \rangle_{\text{planar}}$. One then has

$$C_{123}^{\lambda \ll 1} = \frac{1}{N} \sqrt{(J_1/2)(J_2/2)(J_3/2)}\, \langle \mathcal{C}_1 \mathcal{C}_2 \mathcal{C}_3 \rangle_{\text{planar}} + \mathcal{O}(\lambda^2/N). \tag{5}$$

Therefore we see that the chiral primary structure constant C_{123} is a highly non-trivial function of both the coupling λ and the charges defining the operators. In the extremal case, when $J_3 = J_1 + J_2$, the result at strong coupling simplifies dramatically, and only the planar contraction remains. One finds

$$C_{123}^{\lambda \gg 1}|_{J_3 = J_1 + J_2} = \frac{1}{N} \left(\frac{\lambda}{2\pi^2} \right)^{1/4} \sqrt{(J_1 + 1)(J_2 + 1)(J_3 + 1)}\, \langle \mathcal{C}_{J_1} \mathcal{C}_{J_2} \mathcal{C}_{J_3} \rangle_{\text{planar}}. \tag{6}$$

It is a very interesting direction of future research to determine the C_{123} at higher (or at all) orders of perturbation theory. Judging from the similarity between the strong coupling and tree-level results in the extremal case, it would appear that the extremal problem is far more tractable.

In this paper we provide a generalization of (4) to a specific case when two of the 1/2 BPS operators correspond to a specific giant graviton, and the remaining operator to a pointlike graviton. This implies taking two of the charges, J_2 and J_3, large and the remaining one, J_1, to be order one. We will find that taking the large $J_2 = J_3$ limit of (4) coincides with the small $J_2/N = J_3/N$ limit of the expressions we obtain for the two-giant, one point-like three-point structure constants.[2] This behaviour was

[1] The range of contractions in (4) includes all possible contractions, and so naturally includes the planar contraction.

[2] In the latter limit we first take $N, J_i \to \infty$ with J_i/N fixed and then J_i/N small (where $i = 2, 3$).

also observed in the context of $\mathcal{N} = 4$ SYM [24] where similar holographic three-point functions involving giant gravitons were studied. There, in addition, it was found that extremal correlators exhibited a structure very similar to the dual gauge theory correlators at tree level whereas a complete match was not observed. Later it was shown that one does obtain a complete match for non-extremal three-point functions [26]. In the present letter we present three-point functions of both types expecting that the non-extremal ones truly reflect the strong coupling behaviour of ABJM theory and hoping that the others could help shed light on the subtleties of the extremal case. We also calculate the three-point functions in ABJM perturbation theory at tree-level. Unsurprisingly, it is clear that like in the point-like case, a non-trivial function of the coupling and the charges defining the operators interpolates between weak and strong coupling.

We start by introducing the coordinate system which naturally leads to our improved parametrization in section 2 and move on to discussing in detail the anti-symmetric giant in section 3. In particular, we calculate in this section a number of extremal as well as non-extremal three-point functions involving two anti-symmetric giants and one point-like graviton. The same type of correlation functions are subsequently computed for symmetric giant gravitons in section 4 and in appendix B the dual ABJM three-point functions at tree-level are calculated. Finally, section 5 contains our conclusions.

2 The coordinate system

For the study of the anti-symmetric giant graviton in $AdS_4 \times \mathbb{C}\mathrm{P}^3$, it will prove particularly convenient to use the parametrization

$$Z_1 = re^{i(\chi/2+\phi)} Z \ , \qquad Z_2 = re^{-i(\chi/2-\phi)} \bar{Z}^{-1} \ , \qquad Z_3 = e^{\rho_3 + i(\theta_3+\phi)} \ , \qquad Z_4 = r_4 e^{i\phi} \ , \tag{7}$$

where $r_4^2 = 1 - 2r^2\cosh(2\rho) - e^{2\rho_3}$ and $Z = e^{\rho+i\theta}$. The Z_I cover the unit S^7

$$|Z_1|^2 + |Z_2|^2 + |Z_3|^2 + |Z_4|^2 = 1 \ , \tag{8}$$

once for

$$\begin{aligned} &0 \le e^{2\rho_3} \le 1 - 2r^2\cosh(2\rho) \equiv e^{2\rho_3^{\max}}, \\ -&\rho_{\max} \le \rho \le \rho_{\max} \quad \text{where} \quad \cosh(2\rho_{\max}) = 1/(2r^2), \\ &0 \le r \le 1/\sqrt{2}, \quad 0 \le \theta,\, \theta_3,\, \chi,\, \phi \le 2\pi. \end{aligned} \tag{9}$$

The Z_I are also in one-to-one correspondence with the four bi-fundamental complex scalars W_I of the ABJM theory. In terms of $z_i \equiv Z_i/Z_4$ $(i = 1, 2, 3)$, the S^7 metric is expressed as the $U(1)$ fibration over the Fubini-Study $\mathbb{C}\mathrm{P}^3$,

$$ds^2_{S^7} = \frac{dz_i d\bar{z}_j}{(1 + z_k\bar{z}_k)^2}\left[\delta_{ij}(1 + z_k\bar{z}_k) - \bar{z}_i z_j\right] + (d\phi + A)^2 \ , \tag{10}$$

with the standard 1-form $A = \frac{i}{2}(d - \bar{d})\ln(1 + z_k\bar{z}_k)$. The angle ϕ is the coordinate parametrizing the M-theory circle. The $S^7/\mathbb{Z}_k$ is obtained by restricting the range of the angle ϕ to $0 \le \phi \le 2\pi/k$.

3 The anti-symmetric giant graviton

The giant graviton dual to the Schur polynomial of the $U(N)$ adjoint field $W_1\bar{W}^2$ in antisymmetric representations was found in [1]. In M-theory, the giant graviton is an M5-brane in $S^7/\mathbb{Z}_k$ and described by the curve

$$Z_1\bar{Z}_2 = \alpha^2 e^{i\chi(t)} \ , \tag{11}$$

where α is the constant related to the size of the giant. The time t is that of the global AdS space, and the giant graviton rotates in the χ-direction. This curve reflects the property that the Schur polynomial of the maximal dimension N becomes a product of di-baryon operators, $\det W_1 \det \bar{W}^2$. Namely, when the giant is maximal ($\alpha = 0$), the curve becomes two S^5's ($Z_1 = 0$ and $Z_2 = 0$) intersecting at an S^3 ($Z_1 = Z_2 = 0$) dual to a product of two di-baryons.

In terms of the coordinates introduced in the previous section, the world volume of the M5 giant is spanned by $(t, \rho, \rho_3, \theta, \theta_3, \phi)$ where $-\rho_{\max} \leq \rho \leq \rho_{\max}$ where $\cosh(2\rho_{\max}) = 1/(2\alpha^2)$, $0 \leq e^{2\rho_3} \leq 1 - 2\alpha^2\cosh(2\rho) \equiv e^{2\rho_3^{\max}}$, $0 \leq \theta, \theta_3, k\phi \leq 2\pi$.

3.1 The probe analysis

We shall work in the probe approximation. It is straightforward to find the low energy effective action, i.e., the DBI + WZ action, for the M5 giant:

$$S_{\rm DBI} = \\ -\frac{(2\pi)^3}{k} T_{\rm M5} R_{\rm S^7}^6 \alpha^2 \int_{-\infty}^{\infty} \frac{dt}{2} \int_{-\rho_{\max}}^{\rho_{\max}} d\rho \int_0^{e^{2\rho_3^{\max}}} de^{2\rho_3} \sqrt{\left(\cosh(2\rho) - 2\alpha^2\omega^2\right)\left(\cosh(2\rho) - 2\alpha^2\right)} \ , \tag{12}$$

where $T_{\rm M5} = \ell_P^{-6}/(2\pi)^5$, $R_{S^7}^6 = (2R_{\rm AdS})^6 = 2^3(2\pi)^2 kN\ell_P^6$, and $\omega \equiv \frac{d\chi}{dt}$. Meanwhile, the background 6-form potential is given by

$$C_6 = 2T_{\rm M5} R_{S^7}^6 e^{2\rho_3} r^2 \left(r^2 - \frac{1}{2}\cosh(2\rho)\right) d\rho \wedge d\rho_3 \wedge d\theta \wedge d\chi \wedge d\theta_3 \wedge d\phi + \cdots \ , \tag{13}$$

with an appropriate gauge choice.[3] We then find that

$$S_{\rm WZ} = +\frac{(2\pi)^3}{k} T_{\rm M5} R_{S^7}^6 \alpha^2 \int_{-\infty}^{\infty} \frac{dt}{2} \int_{-\rho_{\max}}^{\rho_{\max}} d\rho \int_0^{e^{2\rho_3^{\max}}} de^{2\rho_3} \left(\cosh(2\rho) - 2\alpha^2\right)\omega \ . \tag{14}$$

Introducing the new variable $x = \cosh(2\rho)$, the DBI + WZ action for the M5 giant yields

$$S_{\rm M5} = 8N\alpha^4 \int_{-\infty}^{\infty} dt \int_1^{\frac{1}{2\alpha^2}} dx \frac{\left(x - \frac{1}{2\alpha^2}\right)}{\sqrt{x^2 - 1}} \left[\sqrt{\left(x - 2\alpha^2\omega^2\right)\left(x - 2\alpha^2\right)} - \omega\left(x - 2\alpha^2\right)\right] \ . \tag{15}$$

[3] The 6-form potential is proportional to $(A + d\Lambda) \wedge dA \wedge dA \wedge d\phi$ where the 1-form $A = 2r^2\cosh(2\rho)d\theta + 2r^2\sinh(2\rho)d\chi + e^{2\rho_3}d\theta_3$. This gauge corresponds to the choice $\Lambda = -\theta$.

Note that the action vanishes when $\omega = 1$ which corresponds to the M5 giant moving at the speed of light, i.e., the giant graviton solution.

The R-charge/angular momentum of the M5 giant is fixed

$$L \equiv \frac{\partial L_{\rm M5}}{\partial \omega} = -8N\alpha^4 \int_1^{\frac{1}{2\alpha^2}} dx \frac{\left(x - \frac{1}{2\alpha^2}\right)}{\sqrt{x^2-1}} \left[2\alpha^2 \omega \sqrt{\frac{x - 2\alpha^2}{x - 2\alpha^2\omega^2}} + \left(x - 2\alpha^2\right) \right] , \quad (16)$$

where $L_{\rm M5}$ is the Lagrangian for the M5 giant. As in [1], we are unable to find ω or the Routhian $R(\alpha, L) \equiv L\omega - L_{\rm M5}(\alpha, \omega)$ as a function of L and α. However, we know that the Routhian is minimized when $\omega = 1$, as numerically checked in [1]. Hence the energy E of the giant graviton is equal to L, saturating the BPS bound. This agrees with the scaling dimension of the Schur polynomial, as L counts field-pairs, i.e. $W_1 \bar{W}^2$ which have conformal dimension 1.

It is easy to find the relation between the angular momentum L and the parameter α related to the size of the giant graviton ($\omega = 1$):

$$\frac{L}{N} = \sqrt{1 - 4\alpha^4} - 4\alpha^4 \log \frac{1 + \sqrt{1 - 4\alpha^4}}{2\alpha^2} \ . \quad (17)$$

The size is maximal when $\alpha = 0$ and zero when $\alpha = \frac{1}{\sqrt{2}}$. In the former case, the angular momentum is maximal $L = N$ (stringy exclusion principle), whereas it vanishes, $L = 0$, in the latter case.

The dimensional reduction to type IIA is straightforward. The M5 giant becomes a D4-brane. In particular, the maximal giant wraps two $\mathbb{CP}^2$'s intersecting at a $\mathbb{CP}^1$.

3.2 Holographic three-point functions

The three-point function between two of the giant gravitons described in the previous section and a chiral primary operator, corresponding to a point-like graviton, may be computed using the techniques described in [18, 22, 24]. The supergravity fluctuations corresponding to a chiral primary operator have been derived in [27–29], and used in a very similar context to the present one in [30], to which we refer the reader for details. The fluctuations are given by

$$\begin{aligned}
\delta g_{\mu\nu} &= h_{\mu\nu} = \frac{4}{J+2} \left[\nabla_\mu \nabla_\nu - \frac{1}{6} J(J-1) \right] s^J(X)\, Y_J(\Omega), \\
\delta g_{\alpha\beta} &= \frac{J}{3}\, g_{\alpha\beta}\, s^J(X)\, Y_J(\Omega), \\
\delta C_{\mu_1\mu_2\mu_3} &= 2\, \epsilon_{\mu_1\mu_2\mu_3\mu_4} \nabla^{\mu_4}\, s^J(X)\, Y_J(\Omega), \\
\delta C_{\alpha_1\ldots\alpha_6} &= -2\, \epsilon_{\alpha_1\ldots\alpha_7} \nabla^{\alpha_7}\, s^J(X)\, Y_J(\Omega),
\end{aligned} \quad (18)$$

where early greek indices refer to $S^7/\mathbb{Z}_k$ coordinates Ω while late greek indices refer to AdS_4 coordinates X, g is the metric and C_3 and C_6 are the three-form and six-form Ramond-Ramond potentials. $Y_J(\Omega)$ represents a scalar spherical harmonic[4] on $S^7/\mathbb{Z}_k$

[4] We take the radius of the $S^7/\mathbb{Z}_k$ to be 2, and the normalization of the spherical harmonics is given by $\int_{S^7/\mathbb{Z}_k} Y^J (Y^K)^* = \delta^{JK} 2^8 \pi^4 k^{-1} [(J/2)!]^2/(J+3)!$.

with angular momentum J, while $s^J(X)$ is a scalar field propagating on AdS_4 with mass-squared $J(J-6)/4$ and has a bulk-to-boundary propagator given by

$$\langle s^J(x,z)\, s^J(x_B,0)\rangle = \ell_p^{9/2}\frac{2^{J/2-1}\pi\sqrt{k}}{R_{\rm AdS}^{9/2}}\frac{J+2}{J}\sqrt{J+1}\frac{z^{J/2}}{((x-x_B)^2+z^2)^{J/2}}, \qquad (19)$$

where (x,z) are Poincaré coordinates on AdS_4, and x_B represents the boundary point. In terms of global coordinates $ds^2_{\rm AdS} = R^2_{\rm AdS}\left(-\cosh^2\mu dt^2 + d\mu^2 + \sinh^2\mu d\Omega_2^2\right)$

$$(x^0,x^1,x^2,z) = \frac{\mathcal{R}}{2(\cosh\mu\cos t - n_1\sinh\mu)}(\cosh\mu\sin t,\, n_2\sinh\mu,\, n_3\sinh\mu,\, 1), \quad (20)$$

where $\vec{n}\cdot\vec{n}=1$ and $\mathcal{R}$ denotes the separation along the boundary of the two giant gravitons. The fluctuation calculation is carried out in Euclidean space, so that $t\to -it$ and $x^0\to -ix^0$. The fluctuations (18) are pulled-back onto the Euclidean M5-brane while the field $s^J(X)$ is replaced by its bulk-to-boundary propagator connecting the insertion point on the brane to the boundary. The boundary point x_B, representing the position of the chiral primary field, is sent to infinity and the resulting expression is integrated over the on-shell world volume of the brane. The result is $-(\mathcal{R}/(2x_B^2))^{J/2}$ times the structure constant defining the three-point function. Note that for the antisymmetric giant of the previous section $\mu=0$.

We find the following fluctuations of the DBI and WZ parts of the Euclidean M5-brane Lagrangian density ($S=\int d^6\sigma\mathcal{L}$)

$$\begin{aligned}
\delta\mathcal{L}_{\rm DBI} &= \frac{R^6_{\rm AdS}}{(2\pi)^5\ell_p^6}\,Y_J(\Omega)\frac{1}{2}\sqrt{g}\left(2J+\frac{\cosh 2\rho}{\cosh 2\rho - 2\alpha^2\omega^2}\left[\frac{4}{J+2}\partial_t^2 - \frac{J^2}{J+2}\right]\right)s^J(X),\\
\delta\mathcal{L}_{\rm WZ} &= -\frac{R^6_{\rm AdS}}{(2\pi)^5\ell_p^6}\,\omega\,\sqrt{g_{S^7/\mathbb{Z}_k}}\,g^{r\beta}_{S^7/\mathbb{Z}_k}\partial_\beta Y_J(\Omega)\Big|_{r=\alpha}\; s^J(X),
\end{aligned} \qquad (21)$$

where $\sqrt{g} = 2^6\alpha^2 e^{2\rho_3}\sqrt{(\cosh 2\rho - 2\alpha^2)(\cosh 2\rho - 2\alpha^2\omega^2)}$, and $g_{S^7/\mathbb{Z}_k}$ is the metric on the $S^7/\mathbb{Z}_k$ of radius 2, parametrized as in section 2, so that on the classical solution, the coordinate r is set to α. Using the following spherical harmonic, dual to the ABJM theory operator ${\rm Tr}(W_1\bar{W}^2)^{J/2}$, we get the extremal correlator with a point-like graviton which is the degeneration of the antisymmetric giant graviton to a point

$$Y_J(\Omega) = \left(r^2 e^{i\chi}\right)^{J/2}. \qquad (22)$$

The resulting structure constant is

$$C^A_{L,L-\Delta,\Delta} = \frac{1}{N}\left(\frac{\lambda}{2\pi^2}\right)^{1/4} 2L\,\left(2\alpha^2\right)^{\Delta}\,\sqrt{2\Delta+1}, \qquad (23)$$

where $\lambda = N/k$, L is the number of field-pairs (i.e. $W_1\bar{W}^2$) in the giant graviton and $\Delta = J/2$ is the number of field-pairs in the chiral primary.

Given the subtleties associated with extremal correlators appearing in analogous computations in the AdS_5/CFT_4 correspondence [24], where agreement between gauge theory and holographic three point functions involving two giant gravitons is found only for the non-extremal case [26], we present here some calculations of non-extremal correlators. Specifically we consider the following operators [30]

$$
\begin{aligned}
\mathcal{O}_{1,0} &= \frac{2\pi}{\sqrt{3}\lambda}\mathrm{Tr}\Big[W_I\bar{W}^I - 4W_1\bar{W}^1\Big],\\
\mathcal{O}_{2,0} &= \frac{8\pi^2}{3\sqrt{5}\lambda^2}\mathrm{Tr}\Big[(W_I\bar{W}^I)^2 - 10W_I\bar{W}^I\,W_1\bar{W}^1 + 15(W_1\bar{W}^1)^2\Big],\\
\mathcal{O}_{3,0} &= \frac{16\pi^3}{3\sqrt{105}\lambda^3}\mathrm{Tr}\Big[(W_I\bar{W}^I)^3 - 18(W_I\bar{W}^I)^2\,(W_1\bar{W}^1)\\
&\qquad\qquad + 63(W_I\bar{W}^I)\,(W_1\bar{W}^1)^2 - 56(W_1\bar{W}^1)^3\Big],
\end{aligned}
\tag{24}
$$

which are normalized chiral primary operators. The corresponding spherical harmonics $Y_J(\Omega)$ are found by substituting (7) for the W_I and normalizing according to footnote 4

$$
\begin{aligned}
Y_2 &= \frac{1}{2\sqrt{3}}\left(1 - 4\,r^2e^{2\rho}\right),\\
Y_4 &= \frac{1}{3\sqrt{10}}\left(1 - 10\,r^2e^{2\rho} + 15\,r^4e^{4\rho}\right),\\
Y_6 &= \frac{1}{4\sqrt{35}}\left(1 - 18\,r^2e^{2\rho} + 63\,r^4e^{4\rho} - 56\,r^6e^{6\rho}\right).
\end{aligned}
\tag{25}
$$

The calculation then proceeds similarly to the extremal case. Using (21) one obtains

$$
C^A_{\mathcal{O}_{\Delta,0}} = \frac{1}{N}\left(\frac{\lambda}{2\pi^2}\right)^{1/4} \times \begin{cases} -\frac{\pi}{2}L, & \Delta = 1\\ -2\sqrt{2}L + \frac{11\sqrt{2}N}{9}(1-4\alpha^4)^{3/2}, & \Delta = 2\,,\\ -\frac{9\pi}{32\sqrt{5}}\left(2L(4+3\alpha^4) - 7N(1-4\alpha^4)^{3/2}\right), & \Delta = 3 \end{cases}
\tag{26}
$$

for the structure constant corresponding to the three-point functions involving two giant gravitons and one of $\mathcal{O}_{1,0}$, $\mathcal{O}_{2,0}$, or $\mathcal{O}_{3,0}$ respectively. As mentioned in the introduction, taking the small L/N (i.e. $\alpha \to 1/\sqrt{2}$) limit, in both the extremal and non-extremal cases, one obtains agreement with the large $J_2 = J_3 = 2L$ limit of the point-like result (4), as was the case in $\mathcal{N} = 4$ SYM [24]. Comparing with free-field contractions (46) and (51)-(54), we see that, as in the point-like case, a non-trivial function of λ and the charges interpolates between weak and strong coupling results.

4 The symmetric giant graviton

The giant graviton dual to the Schur polynomial of the $U(N)$ adjoint field $W_1\bar{W}^2$ in symmetric representations is an M2-brane wrapping the S^2 in global AdS_4 space and rotating along the great circle of $S^7/\mathbb{Z}_k$. This is the so-called AdS giant [4, 5].

More specifically, the AdS giant of interest to us rotates along the great χ-circle; $r = \frac{1}{\sqrt{2}}, \rho = e^{\rho_3} = \theta = \phi = 0$, and $Z_1 \bar{Z}_2 = \frac{1}{2} e^{i\chi(t)}$.

The DBI + WZ action for the M2 giant is given by

$$S_{\text{M2}} = -4\pi T_{\text{M2}} R_{\text{AdS}}^3 \int dt \left[\sinh^2 \mu \sqrt{\cosh^2 \mu - \omega^2} - \sinh^3 \mu \right] , \tag{27}$$

where $ds_{\text{AdS}}^2 = R_{\text{AdS}}^2 \left(-\cosh^2 \mu dt^2 + d\mu^2 + \sinh^2 \mu d\Omega_2^2 \right)$ and $\omega \equiv \frac{d\chi}{dt}$. Note that $4\pi T_{\text{M2}} R_{\text{AdS}}^3 = N/\sqrt{2\lambda}$. The angular momentum yields

$$L \equiv \frac{\partial L_{\text{M2}}}{\partial \omega} = \frac{N}{\sqrt{2\lambda}} \frac{\omega \sinh^2 \mu}{\sqrt{\cosh^2 \mu - \omega^2}} . \tag{28}$$

The Routhian $R(\mu, L) = L\omega - L_{\text{M2}}(\mu, \omega)$ is minimized at $\omega = 1$ corresponding to the M2 giant moving at the speed of light. The energy E of the giant graviton is again L, saturating the BPS bound and the size of the giant is related to the angular momentum by

$$\sinh \mu = \sqrt{2\lambda} \frac{L}{N} . \tag{29}$$

There is no upper bound on the size of the giant in this case. The dimensional reduction to type IIA is trivial, and the M2 giant becomes a D2-brane.

A comment is in order: There is another type of M2 giant which rotates along the M-theory ϕ-circle. Without loss of generality, we can choose $Z_4 = e^{i\phi(t)}, Z_1 = Z_2 = Z_3 = 0$. The analysis is almost the same as the previous case, but there are slight differences. The energy is $E = kL/2$ and the size/angular momentum relation becomes

$$\sinh \mu = \sqrt{\frac{\lambda}{2}} \frac{kL}{N} , \tag{30}$$

where the angular momentum is $L \equiv \frac{\partial L_{\text{M2}}}{\partial \omega}$ with $\omega = \frac{d\phi}{dt}$. Upon dimensional reduction to type IIA, this M2 giant becomes a bound state of a D2-brane and L D0-branes. Since the D0-branes are monopoles in the dual field theory, the dual operator carries L units of monopole charge. The monopole operators can be labeled by the Cartan generators $H = \text{diag}(q_1, q_2 \cdots, q_N)$ and $H' = \text{diag}(q_1', q_2' \cdots, q_N')$ of two $U(N)$'s with $q_1 \geq q_2 \geq \cdots \geq q_N$ and $q_1' \geq q_2' \geq \cdots \geq q_N'$. In particular, the monopole operator $\mathcal{M}_k$ with the unit charges $q_1 = q_1' = 1$ is in the k-dimensional symmetric representations of two $U(N)$'s [31]. Note that, when the level $k = 1$, the operator $\mathcal{M}_1$ is bi-fundamental. Thus the dual operator is the Schur polynomial of the $U(N)$ adjoint field $W_4 \bar{\mathcal{M}}_1$ in the L-dimensional symmetric representation. For $k > 1$, the dual operator is the gauge invariant constructed from $(W_4)^{kL} (\bar{\mathcal{M}}_k)^L$. This operator has dimension $kL/2$ which agrees with the energy of the AdS giant.

4.1 Holographic three-point functions

The computation of the holographic three-point function between two symmetric giants and a chiral primary operator proceeds similarly to section 3.2. We parametrize the S^2 in AdS_4 using (see (20)) $\vec{n} = (\cos\vartheta, \sin\vartheta\sin\varphi, \sin\vartheta\cos\varphi)$. Using (18) we find that the variation of the Lagrangian density is

$$\begin{aligned}\delta\mathcal{L}_{\text{DBI+WZ}} =& \frac{T_{\text{M2}}R^3_{\text{AdS}}}{2}\sinh\mu\sin\vartheta\left[-\frac{J}{3}s + h_{tt} + h_{\vartheta\vartheta} + \frac{h_{\varphi\varphi}}{\sin^2\vartheta}\right] \\ & - 2T_{\text{M2}}R^3_{\text{AdS}}\cosh\mu\sinh^2\mu\sin\vartheta\,\partial_\mu s,\end{aligned} \tag{31}$$

where $s = s^J(X)Y_J(\Omega)$ and where μ is the global AdS_4 coordinate from (27).

Using the chiral primary corresponding to (22), i.e. the point-like degeneration of the giant itself, we find the following structure constant defining the extremal three-point function

$$\begin{aligned}C^S_{L,L-\Delta,\Delta} =& \frac{1}{N}\left(\frac{\lambda}{2\pi^2}\right)^{1/4} 2L\sqrt{2\Delta+1} \\ & \times\left(1+\frac{2L^2}{Nk}\right)^{-1-\Delta/2} {}_2F_1\left(1, 1+\Delta, 3/2, \frac{2L^2}{Nk+2L^2}\right).\end{aligned} \tag{32}$$

We also note the results for the structure constants corresponding to the three-point functions of two symmetric giants and one of the operators in (24), i.e. non-extremal correlators

$$\begin{aligned}C^S_{\mathcal{O}_{\Delta,0}} =& \frac{1}{N}\left(\frac{\lambda}{2\pi^2}\right)^{1/4} 2L\sqrt{2\Delta+1}\,\sqrt{\pi}\frac{\Gamma(1+\Delta/2)}{\Gamma(1/2+\Delta/2)}\,Y_{2\Delta} \\ & \times\left(1+\frac{2L^2}{Nk}\right)^{-1-\Delta/2} {}_2F_1\left(1+\Delta/2, 1+\Delta/2, 3/2, \frac{2L^2}{Nk+2L^2}\right),\end{aligned} \tag{33}$$

where $Y_{2\Delta} = -1/(2\sqrt{3}), -1/(12\sqrt{10}), 3/(16\sqrt{35})$ for $\Delta = 1, 2, 3$ respectively. We note that the expressions (32) and (23) (and similarly, (33) and (26)) agree in the point-like limit, when L/N is small (i.e. $\alpha \to 1/\sqrt{2}$). Thus the symmetric case also reduces to the point-like result (4), in the large $J_2 = J_3 = 2L$ limit. Comparing with free-field contractions (46) and (51)-(54), we see that, as in the point-like case, a non-trivial function of λ and the charges interpolates between weak and strong coupling results.

5 Conclusion

Our greatly simplified parametrization of the anti-symmetric giant graviton of $AdS_4 \times \mathbb{C}\text{P}^3$ made possible the calculation of holographic three-point functions. It is possible that further analytical results can now be obtained. One interesting example is the possibility of obtaining an instanton solution describing the tunneling of the anti-symmetric giant graviton to a point-like one. Such a solution is known to exist in the $AdS_5 \times S^5$ background [5].

Our holographic three-point functions involving two giant and one point-like graviton reduce to the three-point function of three point-like gravitons calculated in the supergravity approach when the size of the giants approaches zero but remains larger than $\mathcal{O}(1)$. The supergravity three-point functions behave as $\lambda^{1/4}$ as $\lambda \to \infty$ (where $\lambda = \frac{N}{k}$) signaling that three-point functions of chiral primaries in ABJM theory are not protected. Hence, we do not expect to be able to recover our holographic three-point functions by a gauge theory computation. In the case of $\mathcal{N} = 4$ SYM, where three-point functions of chiral primaries are known to be protected, the method developed for calculating holographic three-point functions of giant and point-like gravitons [24] led to a complete match between gauge and string theory for non-extremal correlators [26], but extremal correlators did not match their gauge theory duals completely [24]. Based on these observations we expect that our non-extremal three-point correlators correctly encode the strong coupling behaviour of ABJM theory. It remains, however, of utmost importance to fully understand the subtleties of holographic three-point functions in the extremal case and we hope that our results for the $AdS_4 \times \mathbb{C}\mathrm{P}^3$ case will provide useful data for the future development of this topic.

An interesting outcome of our analysis is that holographic three-point functions are very different for anti-symmetric and symmetric giant gravitons in $AdS_4 \times \mathbb{C}\mathrm{P}^3$. This difference is not reflected by the dual correlation functions in ABJM theory when calculated at tree-level, cf. appendix B. As pointed out above, in ABJM theory three-point functions of 1/2 BPS operators are not protected and not even the lowest order loop correction to the three-point function of chiral primaries is known. Calculating such loop corrections constitutes another important future task.

Acknowledgments

We thank A. Bissi, P. Caputa, H. Shimada and K. Zoubos for useful discussions. SH would like to thank The Niels Bohr Institute for their great hospitality. CK and DY were supported by FNU through grant number 272-08-0329. SH was partially supported by the Grant-in-Aid for Nagoya University Global COE Program (G07).

A Chiral primary structure constants from supergravity

The strong coupling result for the three-point function structure constant was given for the $k = 1$ case in [25]. Restoring the k dependence is trivial, and one obtains the following expression

$$C_{123}^{\lambda\gg1} = \frac{1}{N}\left(\frac{\lambda}{2\pi^2}\right)^{1/4} \frac{1}{\Gamma(\gamma/2+1)} \prod_{i=1}^{3} \frac{\Gamma(\gamma_i/2+1)\sqrt{J_i+1}}{\sqrt{J_i!}} \times \frac{k\,2^{\gamma}\,(\gamma+3)!}{2^8\pi^4} \int_{S^7/\mathbb{Z}_k} \mathcal{Y}_{J_1}\mathcal{Y}_{J_2}\mathcal{Y}_{J_3}, \tag{34}$$

where the $S^7/\mathbb{Z}_k$ is taken to have radius 2, and the spherical harmonics $\mathcal{Y}_{J_i}$ are taken to be normalized as

$$\int_{S^7/\mathbb{Z}_k} \mathcal{Y}_J \bar{\mathcal{Y}}_K = \frac{2^8\pi^4}{k}\delta_{JK}\frac{J!}{2^J(J+3)!}. \tag{35}$$

We would like to evaluate the integral of three spherical harmonics, written using the $\mathcal{C}$ tensors appearing in the definition of the operators (1), i.e. using the harmonics

$$Y_J = (\mathcal{C}_A)^{I_1 \dots I_{J/2}}_{K_1 \dots K_{J/2}} Z_{I_1} \cdots Z_{I_{J/2}} \bar{Z}^{K_1} \cdots \bar{Z}^{K_{J/2}}. \tag{36}$$

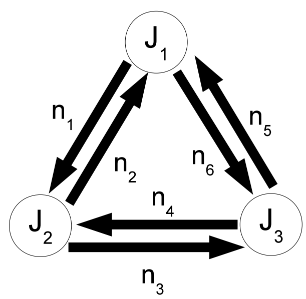
We use the following identity proven in [30]

$$\int_{S^7/\mathbb{Z}_k} Z_{I_1} \cdots Z_{I_m} \bar{Z}^{K_1} \cdots \bar{Z}^{K_m} = \frac{2^8\pi^4}{k(m+3)!} \sum_{\sigma \in S_m} \delta^{I_1}_{K_{\sigma(1)}} \cdots \delta^{I_m}_{K_{\sigma(m)}}, \tag{37}$$

to show that the Y_J are normalized according to footnote 4. The relation between the $\mathcal{Y}_J$ and the Y_J is then

$$\mathcal{Y}_J = \sqrt{\frac{J!}{2^J}} \frac{1}{(J/2)!} Y_J. \tag{38}$$

We will also require the integral over three Y_{J_i}. The identity (37) instructs us to count all possible contractions between the Z's and $\bar{Z}$'s. We use the following figure



to explain this counting. We have $J_i/2$ upper indices and $J_i/2$ lower indices in each of three $\mathcal{C}$ tensors defining the three spherical harmonics. We denote upper-to-lower (Z to $\bar{Z}$) contractions with an arrow pointing at the lower index. We thus have the following constraints

$$\begin{aligned} n_1 + n_6 &= J_1/2 = n_2 + n_5, \\ n_2 + n_3 &= J_2/2 = n_1 + n_4, \\ n_4 + n_5 &= J_3/2 = n_3 + n_6, \end{aligned} \tag{39}$$

which yields the solution

$$\begin{aligned} &n_1 = p, \qquad n_2 = \gamma_3 - p, \qquad n_3 = \gamma_1 - \frac{J_2}{2} + p, \\ &n_4 = \frac{J_2}{2} - p, \qquad n_5 = \gamma_2 - \frac{J_1}{2} + p, \qquad n_6 = \frac{J_1}{2} - p, \end{aligned} \tag{40}$$

where $\gamma_i = (\sum_j J_j - 2J_i)/2$ denotes the total number of contractions between the two operators other than the i^{th}. Assuming, w.l.o.g. that $J_3 \geq J_2 \geq J_1$, we see that

$$p \in [0, \gamma_3]. \tag{41}$$

We can now count possible contractions. We have $(J_1/2)!/[n_1! n_6!]$ ways of dividing the $J_1/2$ Z's into two groups of n_1 and n_6 respectively, while for the $J_1/2$ $\bar{Z}$'s we have $(J_1/2)!/[n_2! n_5!]$. Multiplying by similar factors for each of the three operators, we have

$$\left(\frac{\prod_{i=1}^{3} (J_i/2)!}{\prod_{j=1}^{6} n_j!} \right)^2 \tag{42}$$

many ways of splitting the various Z's and $\bar{Z}$'s into their requisite groups. We then have $\prod_{j=1}^{6} n_j!$ ways of contracting the groups together. We therefore find that

$$\begin{aligned} \int_{S^7/\mathbb{Z}_k} Y_{J_1} Y_{J_2} Y_{J_3} = \frac{2^8 \pi^4}{k\,(\gamma+3)!} \sum_{p=0}^{\gamma_3} \\ \frac{(J_1/2)!^2 (J_2/2)!^2 (J_3/2)!^2}{p!(\gamma_3 - p)!(\gamma_1 - J_2/2 + p)!(J_2/2 - p)!(\gamma_2 - J_1/2 + p)!(J_1/2 - p)!} \\ (\mathcal{C}_{J_1})^{I_1 \ldots I_p I_{p+1} \ldots I_{J_1/2}}_{K_1 \ldots K_{\gamma_3 - p} K_{\gamma_3 - p+1} \ldots K_{J_1/2}} \\ (\mathcal{C}_{J_2})^{K_1 \ldots K_{\gamma_3 - p} L_1 \ldots L_{\gamma_1 - J_2/2 + p}}_{I_1 \ldots I_p M_1 \ldots M_{J_2/2 - p}} \, (\mathcal{C}_{J_3})^{K_{\gamma_3 - p+1} \ldots K_{J_1/2} M_1 \ldots M_{J_2/2 - p}}_{I_{p+1} \ldots I_{J_1/2} L_1 \ldots L_{\gamma_1 - J_2/2 + p}} \, , \end{aligned} \tag{43}$$

and the expression (4) follows.

We note that there are also other spherical harmonics with an unequal number Δ^+ of Z's and Δ^- of $\bar{Z}$'s such that $\Delta^+ - \Delta^- = mk$ where m is an integer. These correspond to states in ABJM with non-zero $U(1)_B$ charge, discussed for example in [30], and which require the presence of the monopole operators discussed beneath (30). Our analysis can also be carried out for this more general case using $n_1 + n_6 = \Delta_1^+$, $n_2 + n_5 = \Delta_1^-$, etc., and requiring that the total $U(1)_B$ charge, $\sum_i m_i = 0$. This then gives a generalization of (4) for these more general operators. It is not clear however, how to evaluate these more general structure constants in perturbation theory.

B The dual operators and their three-point functions

In ABJM theory one can construct operators which form Schur polynomials of a single $U(N)$ by combining two bi-fundamental scalar fields. Denoting the two complex bi-fundamental scalars as W_1 and $\bar{W}^2$ a $U(N)$ Schur polynomial can then be written as

$$\chi_{R_L}(W_1 \bar{W}^2) = \frac{1}{L!} \sum_{\sigma \in S_L} \chi_{R_L}(\sigma)\, (W_1 \bar{W}^2)^{i_{\sigma(1)}}_{i_1} \ldots (W_1 \bar{W}^2)^{i_{\sigma(L)}}_{i_L}, \tag{44}$$

where R_L denotes an irreducible representation of $U(N)$ described in terms of a Young tableau with L boxes. The sum is over elements of the symmetric group and

$\chi_{R_L}(\sigma)$ is the character of the element σ in the representation R_L. The calculation of two- and three-point functions of operators of the type (44) at tree-level is a purely combinatorial problem which can be solved in close analogy with the similar problem involving a single adjoint scalar appearing in $\mathcal{N}=4$ SYM, see [9].

The structure constant dual to the three-point function of two giants and one point-like graviton of $AdS_4 \times \mathbb{C}\mathrm{P}_3$ is

$$C_{L,L-\Delta,\Delta} \equiv \frac{\langle \bar{\chi}_L \, \chi_{L-\Delta} \mathrm{Tr}(W_1 \bar{W}^2)^\Delta \rangle}{\sqrt{\langle \bar{\chi}_L \, \chi_L \, \rangle \langle \bar{\chi}_{L-\Delta} \chi_{L-\Delta} \rangle \langle \mathrm{Tr}(\bar{W}^1 W_2)^\Delta \mathrm{Tr}(W_1 \bar{W}^2)^\Delta \rangle}}, \tag{45}$$

where here and in the following the expectation values are to be understood as expectation values in a zero-dimensional Gaussian complex matrix model with unit propagator. Furthermore, χ_L is the Schur polynomial corresponding to a Young tableau consisting either of a single column (anti-symmetric case) or a single row (symmetric case) with L boxes and we have suppressed the dependence of the W_I-fields. Expanding the single trace operator in the basis of Schur polynomials and making use of the known three-point functions of the Schurs from [9] one easily finds the following expression for the three-point functions in the limit $\Delta \ll L$, $L, N \to \infty$, $\frac{L}{N}$ fixed

$$C^A_{L,L-\Delta,\Delta} = \frac{1}{\sqrt{\Delta}} \left(1 - \frac{L}{N}\right)^\Delta, \qquad C^S_{L,L-\Delta,\Delta} = (-1)^{\Delta-1} \frac{1}{\sqrt{\Delta}} \left(1 + \frac{L}{N}\right)^\Delta, \tag{46}$$

where the superscript A refers to the anti-symmetric case and S to the symmetric one. For details we refer to [24].

To determine the tree-level contribution to the non-extremal ABJM three-point function, dual to the correlator of two giant gravitons and one point-like one of the type given in equation (24) one has to evaluate

$$C_{\mathcal{O}_{\Delta,0}} = \frac{\langle \bar{\chi}_L \, \chi_L \, \mathcal{O}_{\Delta,0} \rangle}{\langle \bar{\chi}_L \, \chi_L \, \rangle}. \tag{47}$$

Writing the operators out explicitly, stripping off the factors originating from the gauge theory propagators and furthermore exploiting the symmetry properties of the expectation values one finds that in the formula (47) one can replace $\mathcal{O}_{\Delta,0}$ by $\mathcal{O}^{eff}_{\Delta,0}$ given by

$$\mathcal{O}^{eff}_{1,0} = \frac{1}{2\sqrt{3}N} \mathrm{Tr}(-2 W_1 \bar{W}^1), \tag{48}$$

$$\mathcal{O}^{eff}_{2,0} = \frac{1}{6\sqrt{5}N^2} \mathrm{Tr}\left[7(W_1 \bar{W}^1)^2 - 4 W_1 \bar{W}^1 W_2 \bar{W}^2 - 4 W_1 \bar{W}^2 W_2 \bar{W}^1\right], \tag{49}$$

$$\mathcal{O}^{eff}_{3,0} = \frac{1}{12\sqrt{105}\,N^3} \mathrm{Tr}\left[-9(W_1 \bar{W}^1)^3 + 5(W_1 \bar{W}^1)^2 W_2 \bar{W}^2 + 5 W_1 \bar{W}^1 W_1 \bar{W}^2 W_2 \bar{W}^1 + 5 W_1 \bar{W}^1 W_2 \bar{W}^1 W_1 \bar{W}^2\right]. \tag{50}$$

A somewhat lengthy but in principle straightforward calculation along the lines of [26] gives

$$C^S_{\mathcal{O}_{1,0}} = C^A_{\mathcal{O}_{1,0}} = -\frac{1}{\sqrt{3}}\frac{L}{N}, \tag{51}$$

$$C^S_{\mathcal{O}_{2,0}} = -\frac{1}{6\sqrt{5}}\frac{L}{N}\left(8+\frac{L}{N}\right), \qquad C^A_{\mathcal{O}_{2,0}} = -\frac{1}{6\sqrt{5}}\frac{L}{N}\left(8-\frac{L}{N}\right), \tag{52}$$

$$C^S_{\mathcal{O}_{3,0}} = \frac{1}{12\sqrt{105}}\frac{L}{N}\left(6\left(\frac{L}{N}\right)^2+5+20\frac{L}{N}\right), \tag{53}$$

$$C^A_{\mathcal{O}_{3,0}} = \frac{1}{12\sqrt{105}}\frac{L}{N}\left(6\left(\frac{L}{N}\right)^2+5-20\frac{L}{N}\right). \tag{54}$$

References

[1] D. Giovannoni, J. Murugan, and A. Prinsloo, "The Giant graviton on AdS_4xCP^3 - another step towards the emergence of geometry," *JHEP* **1112** (2011) 003, `arXiv:1108.3084 [hep-th]`.

[2] J. M. Maldacena, "The Large N limit of superconformal field theories and supergravity," *Adv.Theor.Math.Phys.* **2** (1998) 231–252, `arXiv:hep-th/9711200 [hep-th]`.

[3] J. McGreevy, L. Susskind, and N. Toumbas, "Invasion of the giant gravitons from Anti-de Sitter space," *JHEP* **0006** (2000) 008, `arXiv:hep-th/0003075 [hep-th]`.

[4] M. T. Grisaru, R. C. Myers, and O. Tafjord, "SUSY and goliath," *JHEP* **0008** (2000) 040, `arXiv:hep-th/0008015 [hep-th]`.

[5] A. Hashimoto, S. Hirano, and N. Itzhaki, "Large branes in AdS and their field theory dual," *JHEP* **0008** (2000) 051, `arXiv:hep-th/0008016 [hep-th]`.

[6] V. Balasubramanian, M. Berkooz, A. Naqvi, and M. J. Strassler, "Giant gravitons in conformal field theory," *JHEP* **0204** (2002) 034, `arXiv:hep-th/0107119 [hep-th]`.

[7] S. Corley, A. Jevicki, and S. Ramgoolam, "Exact correlators of giant gravitons from dual N=4 SYM theory," *Adv.Theor.Math.Phys.* **5** (2002) 809–839, `arXiv:hep-th/0111222 [hep-th]`.

[8] O. Aharony, O. Bergman, D. L. Jafferis, and J. Maldacena, "N=6 superconformal Chern-Simons-matter theories, M2-branes and their gravity duals," *JHEP* **0810** (2008) 091, `arXiv:0806.1218 [hep-th]`.

[9] T. K. Dey, "Exact Large R-charge Correlators in ABJM Theory," *JHEP* **1108** (2011) 066, `arXiv:1105.0218 [hep-th]`.

[10] S. Chakrabortty and T. K. Dey, "Correlators of Giant Gravitons from dual ABJ(M) Theory," *JHEP* **1203** (2012) 062, arXiv:1112.6299 [hep-th].

[11] D. Berenstein and D. Trancanelli, "Three-dimensional N=6 SCFT's and their membrane dynamics," *Phys.Rev.* **D78** (2008) 106009, arXiv:0808.2503 [hep-th].

[12] T. Nishioka and T. Takayanagi, "Fuzzy Ring from M2-brane Giant Torus," *JHEP* **0810** (2008) 082, arXiv:0808.2691 [hep-th].

[13] A. Hamilton, J. Murugan, A. Prinsloo, and M. Strydom, "A Note on dual giant gravitons in AdS(4) x CP**3," *JHEP* **0904** (2009) 132, arXiv:0901.0009 [hep-th].

[14] J. Murugan and A. Prinsloo, "ABJM Dibaryon Spectroscopy," *JHEP* **1105** (2011) 129, arXiv:1103.1163 [hep-th].

[15] N. Gutierrez, Y. Lozano, and D. Rodriguez-Gomez, "Charged particle-like branes in ABJM," *JHEP* **1009** (2010) 101, arXiv:1004.2826 [hep-th].

[16] Y. Lozano, M. Picos, K. Sfetsos, and K. Siampos, "ABJM Baryon Stability and Myers effect," *JHEP* **1107** (2011) 032, arXiv:1105.0939 [hep-th].

[17] M. Herrero, Y. Lozano, and M. Picos, "Dielectric 5-Branes and Giant Gravitons in ABJM," *JHEP* **1108** (2011) 132, arXiv:1107.5475 [hep-th].

[18] K. Zarembo, "Holographic three-point functions of semiclassical states," *JHEP* **1009** (2010) 030, arXiv:1008.1059 [hep-th].

[19] R. A. Janik, P. Surowka, and A. Wereszczynski, "On correlation functions of operators dual to classical spinning string states," *JHEP* **1005** (2010) 030, arXiv:1002.4613 [hep-th].

[20] E. Buchbinder and A. Tseytlin, "Semiclassical four-point functions in AdS_5xS^5," *JHEP* **1102** (2011) 072, arXiv:1012.3740 [hep-th].

[21] E. Buchbinder and A. Tseytlin, "On semiclassical approximation for correlators of closed string vertex operators in AdS/CFT," *JHEP* **1008** (2010) 057, arXiv:1005.4516 [hep-th].

[22] M. S. Costa, R. Monteiro, J. E. Santos, and D. Zoakos, "On three-point correlation functions in the gauge/gravity duality," *JHEP* **1011** (2010) 141, arXiv:1008.1070 [hep-th].

[23] D. Bak, B. Chen, and J.-B. Wu, "Holographic Correlation Functions for Open Strings and Branes," *JHEP* **1106** (2011) 014, arXiv:1103.2024 [hep-th].

[24] A. Bissi, C. Kristjansen, D. Young, and K. Zoubos, "Holographic three-point functions of giant gravitons," *JHEP* **1106** (2011) 085, arXiv:1103.4079 [hep-th].

[25] F. Bastianelli and R. Zucchini, "Three point functions of chiral primary operators in d = 3, N = 8 and d = 6, N = (2,0) SCFT at large N," *Phys. Lett.* **B467** (1999) 61–66, `arXiv:hep-th/9907047`.

[26] P. Caputa, R. d. M. Koch, and K. Zoubos, "Extremal vs. Non-Extremal Correlators with Giant Gravitons," `arXiv:1204.4172 [hep-th]`.

[27] B. Biran, A. Casher, F. Englert, M. Rooman, and P. Spindel, "THE FLUCTUATING SEVEN SPHERE IN ELEVEN-DIMENSIONAL SUPERGRAVITY," *Phys.Lett.* **B134** (1984) 179.

[28] L. Castellani, R. D'Auria, P. Fre, K. Pilch, and P. van Nieuwenhuizen, "THE BOSONIC MASS FORMULA FOR FREUND-RUBIN SOLUTIONS OF d = 11 SUPERGRAVITY ON GENERAL COSET MANIFOLDS," *Class.Quant.Grav.* **1** (1984) 339.

[29] F. Bastianelli and R. Zucchini, "Bosonic quadratic actions for 11-D supergravity on AdS(7/4) x S(4/7)," *Class.Quant.Grav.* **16** (1999) 3673–3684, `arXiv:hep-th/9903161 [hep-th]`.

[30] N. Drukker, J. Gomis, and D. Young, "Vortex Loop Operators, M2-branes and Holography," *JHEP* **0903** (2009) 004, `arXiv:0810.4344 [hep-th]`.

[31] I. R. Klebanov and G. Torri, "M2-branes and AdS/CFT," *Int. J. Mod. Phys.* **A25** (2010) 332–350, `arXiv:0909.1580 [hep-th]`.