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Rarita-Schwinger field in the AdS/CFT correspondence

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Abstract: A free Rarita-Schwinger field in the Anti-de Sitter space is considered. We show that the usual action can be supplemented by a boundary term that can be interpreted as the generating functional of the correlation functions in a conformal field theory on the boundary of the Anti-de Sitter space.

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Contents

1. Introduction

Recently it has been proposed by Maldacena [1] that there is an exact correspondence between string theory on Anti-de Sitter space (AdS) and a certain superconformal field theory (CFT) on its boundary. For instance, the large N and the large t'Hooft coupling limit of $\mathcal{N}=4$ four dimensional $SU(N)$ Super Yang–Mills can be described by IIB supergravity compactified on $AdS_5 \times S^5$. In [2, 3] the recipe for computing field theory observables via AdS was suggested, namely, the action for a field theory on AdS considered as the functional of the asymptotic value of the field on the boundary of the AdS space is interpreted as the generating functional for the correlation functions in a conformal field theory on the boundary. For example, for the scalar field one has

$$Z_{AdS}[\phi_0] = \int_{\phi_0} \mathcal{D}\phi \, \exp(-S[\phi]) = Z_{CFT}[\phi_0] = \left\langle \exp\left(\int_{\partial AdS_d} d^d x \mathcal{O}\phi_0\right)\right\rangle , \tag{1.1}$$

where ϕ_0 is the boundary value of ϕ and it couples to the scalar $\mathcal{O}$ on the boundary. Hence calculating the action on the AdS_{d+1} side allows one to obtain the correlation functions on the CFT_d side.

Recently, scalar, spinor, vector and graviton fields have been considered on various products of AdS spaces; two, three and four-point functions were calculated and various non-renormalization theorems suggested [2]–[21]. One can deduce the scaling dimensions of operators in the conformal field theory from the masses of particles in string theory. Spinor fields of spin 1/2 were discussed in [7] (see also [6]) where it was pointed out that the ordinary Dirac action vanishes on-shell; so to get the generating functional in conformal field theory one has to supplement the action by a certain boundary term, which produces the correlation function in conformal field theory.

In this note we will discuss the Rarita-Schwinger field on AdS_{d+1}. It can couple to the supersymmetry current, thus it is natural to reproduce the two-point correlation function of the supersymmetry currents from the partition function on the AdS side by finding the appropriate boundary term to add to the action, since the bulk action vanishes onshell. In section 2 we consider the Rarita-Schwinger field on AdS_{d+1} and find the solution of the corresponding Dirichlet problem. In section 3 we review the two-point function of operators coupling to gravitini, which is restricted by conformal invariance. Finally, in section 4 we find the boundary term on the AdS_{d+1} side which reproduces the correlator on CFT_d side, and the relation between the scaling dimension of the operator in the CFT_d and the mass of the particle in the AdS_{d+1} space.

2. The Rarita-Schwinger field on AdS_{d+1}

Here we will consider the Rarita-Schwinger field on AdS_{d+1}. The action for the Rarita-Schwinger field in curved space reads

$$I_0 = \int_{AdS} d^{d+1}x\sqrt{G}\bar{\Psi}_\mu(\Gamma^{\mu\nu\lambda}D_\nu - m\Gamma^{\mu\lambda})\Psi_\lambda\,, \tag{2.1}$$

where

$$\Gamma^{\mu\nu\lambda} = \Gamma^{[\mu}\Gamma^\nu\Gamma^{\lambda]},\quad \Gamma^{\mu\nu} = \Gamma^{[\mu}\Gamma^{\nu]}\,, \tag{2.2}$$

$\Psi_\mu = (\Psi_{\mu\alpha})$ is the Rarita-Schwinger field, $\Gamma^\mu = \gamma^a e^\mu_a$, $\{\gamma^a,\gamma^b\} = 2\delta^{ab}$, e^μ_a is a vielbein such that the metric $G_{\mu\nu} = e^a_\mu e^b_\nu \delta_{ab}$; and by [...] we denoted antisymmetrization of the indices; $\mu,\nu,a,b = 0,1,...,d$.

In the coordinates $x^\mu = (x^0,\mathbf{x}) = (x^0,x^i)$, $i = 1,\ldots,d$, the Euclidean AdS_{d+1} space is represented by the Lobachevsky upper half-space $x^0 > 0$ and the metric ds^2 is given by

$$ds^2 = G_{\mu\nu}dx^\mu dx^\nu = \frac{1}{x_0^2}(dx^0dx^0 + d\mathbf{x}\cdot d\mathbf{x})\,. \tag{2.3}$$

The boundary $M_d = \partial AdS_{d+1}$ is defined by the hypersurface $x^0 = 0$ plus a single point at $x^0 = \infty$. With the choice of the vielbein

$$e^a_\mu = (x^0)^{-1}\delta^a_\mu\,,\quad a = 0,\ldots,d \tag{2.4}$$

for which the corresponding nonvanishing components of the spin connection ω^{ab}_μ has the form [7]

$$\omega^{0j}_i = -\omega^{j0}_i = (x^0)^{-1}\delta^j_i\,, \tag{2.5}$$

the operators D_ν, $\not{D}$ are given by

$$D_\nu = \partial_\nu + \frac{1}{2}\omega^{bc}_\nu\Sigma_{bc} = \partial_\nu + \frac{1}{2x^0}\gamma_{0\nu},\quad \not{D} = x^0\gamma^0\partial_0 + x^0\mathbf{\Gamma}\cdot\nabla - \frac{d}{2}\gamma^0\,. \tag{2.6}$$

Using the condition

$$\Gamma^\mu\Psi_\mu = 0\,, \tag{2.7}$$

we can rewrite the Rarita-Schwinger equation

$$\Gamma^{\mu\nu\rho} D_\nu \Psi_\rho = m\Gamma^{\mu\nu}\Psi_\nu \tag{2.8}$$

in the following form for $\psi_a = e^\mu_a \Psi_\mu$:

$$x_0\gamma^\nu\partial_\nu\psi_a - \frac{d}{2}\gamma_0\psi_a = \gamma_a\psi_0 - m\psi_a\,. \tag{2.9}$$

The solution of the above is given by

$$\psi_0(x_0,\vec{x}) = \int d^d\vec{p}e^{i\vec{p}\vec{x}}(px_0)^{\frac{d+3}{2}}(i\frac{\hat{p}}{p}K_{m+1/2}(px_0) + K_{m-1/2}(px_0))c_0^-(\vec{p})\,, \tag{2.10}$$

$$\begin{aligned}\psi_i(x_0,\vec{x}) = &\int d^d\vec{p}e^{i\vec{p}\vec{x}}(x_0p)^{\frac{d+1}{2}}(i\frac{\hat{p}}{p}K_{m+1/2}(px_0) + K_{m-1/2}(px_0))c_i^-(\vec{p}) + \\ &+\int d^d\vec{p}e^{i\vec{p}\vec{x}}(x_0p)^{\frac{d+3}{2}}[\frac{p_i\hat{p}}{p^2}K_{m+3/2}(px_0) - \frac{ip_i}{p}K_{m+1/2}(px_0) + \\ &+\frac{\gamma_i}{x_0p}K_{m+1/2}(px_0)]c_0^-(\vec{p})\,.\end{aligned} \tag{2.11}$$

where $p = |\vec{p}|$, $\hat{p} = \gamma^i p_i$; K_ν is the modified Bessel function and c_0, c_i satisfy $\gamma_0 c_0^- = -c_0^-$ and $\gamma_0 c_i^- = -c_i^-$.

The condition $\gamma_0\psi_0 + \gamma_i\psi_i = 0$ takes the form

$$c_0^-(\vec{p}) = -\frac{2ip_jc_j^-(\vec{p})}{(2m+d+1)p}, \qquad \gamma_i c_i^-(\vec{p}) = 0\,. \tag{2.12}$$

Similarly, the solution of the conjugate equation is

$$\begin{aligned}\bar{\psi}_0(x_0,\vec{x}) = &\int d^d\vec{p}e^{i\vec{p}\vec{x}}(px_0)^{\frac{d+3}{2}}\bar{c}_0^+(\vec{p})(-i\frac{\hat{p}}{p}K_{m+1/2}(px_0) + K_{m-1/2}(px_0))\,,\\ \bar{\psi}_i(x_0,\vec{x}) = &\int d^d\vec{p}\bar{c}_i^+(\vec{p})e^{i\vec{p}\vec{x}}(x_0p)^{\frac{d+1}{2}}(-i\frac{\hat{p}}{p}K_{m+1/2}(px_0) + K_{m-1/2}(px_0)) + \\ &+\int d^d\vec{p}\bar{c}_0^+(\vec{p})e^{i\vec{p}\vec{x}}(x_0p)^{\frac{d+3}{2}}[\frac{p_i\hat{p}}{p^2}K_{m+3/2}(px_0) - \\ &-\frac{ip_i}{p}K_{m+1/2}(px_0) - \frac{\gamma_i}{x_0p}K_{m+1/2}(px_0)]\,,\end{aligned} \tag{2.13}$$

with

$$\bar{c}_0^+(\vec{p}) = \frac{2ip_j\bar{c}_j^+(\vec{p})}{(2m+d+1)p}, \qquad \bar{c}_i^+(\vec{p})\gamma_i = 0\,. \tag{2.14}$$

Let us now consider the boundary conditions on the surface $x_0 = \epsilon$, $\epsilon \to 0$ [9, 15]. Without loss of generality we will consider the case of positive m (the other case can be treated similarly). Then fixing the positive components $\psi_i^{\epsilon,+}(\vec{x})$ of the boundary field we find that the negative components $\psi_i^{\epsilon,-}(\vec{x})$ vanish, where $\psi_i^{\epsilon,\pm}(\vec{x}) = \frac{1}{2}(1\pm\gamma_0)\psi_i(x_0 = \epsilon,\vec{x})$. To specify solutions we use the functions χ_i

$$\bar{\chi}_i^-(\vec{x}) = \epsilon^{m-d/2}\bar{\psi}_i^{\epsilon,-}(\vec{x}), \quad \chi_i^+(\vec{x}) = \epsilon^{m-d/2}\psi_i^{\epsilon,+}(\vec{x})\,. \tag{2.15}$$

The dependence of c_0 and c_i on them is given by (A.9) and (A.10) in appendix A.

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3. Correlator in CFT_d

Conformal invariance fixes the two-point correlation function for spin 3/2 primary fields of scaling dimension Δ up to normalization to be [26]

$$< O_{k\alpha}(x) O_{j\beta}(0) >= C \frac{\hat{x}_{\alpha\beta}}{x^{2\Delta+1}} (\delta_{kj} - 2\frac{x_k x_j}{x^2}) \,, \tag{3.1}$$

where i, j are vector indices, α, β are spinor indices and $\hat{x}_{\alpha\beta} = \gamma^i_{\alpha\beta} x_i$.

4. Boundary term

The Rarita-Schwinger action vanishes on-shell; thus in order to reproduce the two-point correlation function of spin 3/2 fields in the boundary conformal field theory one has to supplement the action by the boundary term. This is analogous to the spinor case [7] and so we similarly take

$$I \sim \lim_{\epsilon \to 0} \int_{M_d^\epsilon} d^d\vec{x} \sqrt{G_\epsilon} \bar{\Psi}_i G^{ij} \Psi_j \,, \tag{4.1}$$

where as in the spinor case M_d^ϵ is the surface of $x_0 = \epsilon$, G^{ij} is the induced metric on M_d^ϵ with determinant $G_\epsilon = \epsilon^{-2d}$. To compute the total action we following [9] consider the Dirichlet problem on M_d^ϵ.

The action (4.1) can be rewritten in momentum space in the form

$$I \sim \int d^d\vec{p} \epsilon^{-d} (\bar{\psi}_i^+(\epsilon, \vec{p}) \psi_i^+(\epsilon, -\vec{p}) + \bar{\psi}_i^-(\epsilon, \vec{p}) \psi_i^-(\epsilon, -\vec{p})) \,. \tag{4.2}$$

Generally speaking the expression (4.2) contains divergences when $\epsilon \to 0$. Selecting the finite part, one gets the action (see appendix A for details of the derivation)

$$I \sim \int d^d p \bar{\chi}_i^+(\vec{p}) p^{2m} (\frac{\delta_{ij}\hat{p}}{p} - \frac{2(2m-1)}{d+2m-1} \cdot \frac{p_i p_j \hat{p}}{p^3}) \chi_j^-(-\vec{p}) \,, \tag{4.3}$$

where the boundary spinors χ_i and $\bar{\chi}_i$ are defined as in (2.15).

Taking the Fourier transform and integrating over $\vec{p}$ (we use integrals from appendix B) we get the expression for the boundary term on the AdS side

$$I \sim \int d^d x \int d^d y \bar{\chi}_i^+(x) \frac{\hat{x} - \hat{y}}{(x-y)^{2m+d+1}} \times (\delta_{ij} - 2\frac{(x-y)_i (x-y)_j}{(x-y)^2}) \chi_j^-(y) \,, \tag{4.4}$$

which matches the expression for the CFT correlator (3.1) with

$$\Delta = d/2 + m \,. \tag{4.5}$$

So we found the agreement between the two-point function derived from the AdS_{d+1} side and the CFT_d side once the appropriate boundary term was added.

The mass spectrum of IIB supergravity on $AdS_5 \times S^5$ has been worked out in [22]. In particular it was shown that there is a massless gravitino field which satisfies equation (2.8) with $m = 3/2$. The IIB supergravity on $AdS_5 \times S^5$ is dual to four dimensional $\mathcal{N} = 4$ Super Yang-Mills [1]. The supersymmetry current Σ^μ_α

$$\Sigma^\mu{}_{\alpha A} = -\sigma^{k\nu} F^-_{k\nu} \sigma^\mu{}_{\alpha\dot\alpha} \overline{\lambda}^{\dot\alpha}{}_A + 2i\overline{\varphi}_{AB} \overleftrightarrow{\partial^\mu} \lambda_\alpha{}^B + \frac{4}{3} i\sigma^{\mu\nu}{}_\alpha{}^\beta \partial_\nu (\overline{\varphi}_{AB} \lambda_\beta{}^B)\,, \tag{4.6}$$

(see [24, 25] for details) which couples to the Rarita-Schwinger field has the scaling dimension $\Delta = 7/2$. One gets this dimension from (4.5) if one takes $d = 4$ and $m = 3/2$. This provides a test of the AdS/CFT correspondence in this case. The 11d supergravity on $AdS_7 \times S^4$, is related to the theory on M5-branes [1] and thus to (2,0) theory in the large N limit. The spectrum of the supergravity theory was computed in [23] and that of the primary operators of the conformal algebra of arbitrary spin was discussed in [6, 5]. In particular the gravitino was discussed and the relation between the scaling dimension and the mass agrees with our analysis.

After this paper was completed and prepared the paper [27] appeared with similar results.

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A. Derivation of the finite action

Here we will present the details of the imposed boundary conditions and the derivation of the term (4.3). Let us set $x_0 = \epsilon$ in (2.11)

$$\psi_i^\epsilon(\vec{p}) = \int d^d\vec{x} e^{-i\vec{p}\vec{x}} \psi_i(\epsilon, \vec{x}) = (\epsilon p)^{\frac{d+1}{2}} (i\frac{\hat{p}}{p} K_{m+1/2}(p\epsilon) + K_{m-1/2}(p\epsilon)) c_i^-(\vec{p}) +$$
$$+ (\epsilon p)^{\frac{d+3}{2}} [\frac{p_i \hat{p}}{p^2} K_{m+3/2}(p\epsilon) - \frac{ip_i}{p} K_{m+1/2}(p\epsilon) + \frac{\gamma_i}{\epsilon p} K_{m+1/2}(p\epsilon)] c_0^-(\vec{p})\,, \tag{A.1}$$

and introduce

$$\psi_i^{\epsilon,\pm} = \frac{1}{2}(1 \pm \gamma_0)\psi_i^\epsilon\,. \tag{A.2}$$

Then from the asymptotic behavior of the Bessel function for $m > 0$ we conclude that once we fix $\psi_i^{\epsilon,+}$, then $\psi_i^{\epsilon,-} \to 0$, so analagous to the spinor case $\psi_i^{\epsilon,+}$ can be used to determine the boundary data and

$$\psi_i^{\epsilon,+}(\vec{p}) = (\epsilon p)^{\frac{d+1}{2}} i\frac{\hat{p}}{p} K_{m+1/2}(p\epsilon)\cdot c_i^-(\vec{p}) + \qquad (A.3)$$

$$+(\epsilon p)^{\frac{d+3}{2}}[\frac{p_i\hat{p}}{p^2}K_{m+3/2}(p\epsilon) + \frac{\gamma_i}{\epsilon p}K_{m+1/2}(p\epsilon)]c_0^-(\vec{p})\,. \qquad (A.4)$$

Let us now find the dependence on c_0^- and c_i^- from $\psi_i^{\epsilon,+}$. Suppose for some a and f

$$c_i^- = b_i + \frac{p_i}{p}ac_0^- + f\gamma_i c_0^-\,.$$

Then in (A.3) if we require the first term in the right hand side to be equal to the left hand side we get

$$b_i = -\frac{i\hat{p}\psi_i^{\epsilon,+}}{pK_{m+1/2}(\epsilon p)}(\epsilon p)^{-\frac{d+1}{2}}, \qquad (A.5)$$

$$a = i\frac{(p\epsilon)K_{m+3/2}(p\epsilon)}{K_{m+1/2}(p\epsilon)}, \qquad (A.6)$$

$$f = \frac{i\hat{p}}{p} \qquad (A.7)$$

or

$$c_i^- = -\frac{i\hat{p}}{p}\frac{\psi_i^{\epsilon,+}}{K_{m+1/2}(p\epsilon)}(\epsilon p)^{-\frac{d+1}{2}} + [\frac{ip_i}{p}\frac{(p\epsilon)K_{m+3/2}(p\epsilon)}{pK_{m+1/2}(p\epsilon)} + \frac{i\hat{p}\gamma_i}{p}]c_0^-\,. \qquad (A.8)$$

The condition (2.12) gives the expresion for c_0 and c_i in the form

$$c_0^- = -\frac{2\hat{p}p_j}{p^2}\frac{(\epsilon p)^{-\frac{d+1}{2}}}{(2m+d-1)K_{m+1/2}(p\epsilon) - 2(\epsilon p)K_{m+3/2}(\epsilon p)}\psi_j^{\epsilon,+}\,, \qquad (A.9)$$

$$c_i^- = -\frac{i\hat{p}}{p}\frac{(\epsilon p)^{-\frac{d+1}{2}}}{K_{m+1/2}(p\epsilon)}\psi_i^{\epsilon,+} - (\frac{ip_i}{p}\frac{(p\epsilon)K_{m+3/2}(p\epsilon)}{pK_{m+1/2}(p\epsilon)} + \frac{i\hat{p}\gamma_i}{p})\times$$

$$\times(\frac{2\hat{p}p_j}{p^2}\frac{(\epsilon p)^{-\frac{d+1}{2}}}{(2m+d-1)K_{m+1/2}(p\epsilon) - 2(\epsilon p)K_{m+3/2}(\epsilon p)}\psi_j^{\epsilon,+})\,. \qquad (A.10)$$

Then from (2.11) we get

$$\psi_i(\epsilon,\vec{x}) = \int d^d\vec{p}e^{i\vec{p}\vec{x}}[\psi_i^{\epsilon,+} + (\epsilon p)^{\frac{d+1}{2}}K_{m-1/2}(\epsilon p)c_i^- - (\epsilon p)^{\frac{d+3}{2}}(\frac{ip_i}{p})K_{m+1/2}(\epsilon p)c_0^-]\,. \quad (A.11)$$

Substituting the expression (A.9) and (A.10) in (A.11) and defining the boundary spinors

$$\bar{\chi}_i^- = \lim_{\epsilon\to 0}\epsilon^{m-d/2}\bar{\psi}_i^{\epsilon,-}, \quad \chi_i^+ = \lim_{\epsilon\to 0}\epsilon^{m-d/2}\psi_j^{\epsilon,+}\,, \qquad (A.12)$$

we get

$$\psi_i = \epsilon^{-m+d/2} \int d^d\vec{p} e^{i\vec{p}\vec{x}} [\chi_i^+ - \frac{i\hat{p}}{p} \cdot \frac{K_{m-1/2}}{K_{m+1/2}} \chi_i^+ -$$
$$- \{ (\frac{ip_i}{p}) \epsilon p \frac{K_{m-1/2}K_{m+3/2} - K_{m+1/2}^2}{K_{m+1/2}} + (\frac{i\hat{p}\gamma_i}{p}) K_{m-1/2} \} \times$$
$$\times \{ \frac{2p_j\hat{p}}{p^2} \cdot \frac{1}{(2m+d-1)K_{m+1/2} - 2\epsilon p K_{m+3/2}} \chi_j^+ \}] . \tag{A.13}$$

We assumed here that the argument of the Bessel function is ϵp. The very similar expression can be easily written for the conjugate spinor. Now let us proceed with calculation of the boundary term:

$$I \sim \int d^d\vec{k} \epsilon^{-d} (\bar{\psi}_i^+(\vec{k}) \psi_i^+(-\vec{k}) + \bar{\psi}_i^-(\vec{k}) \psi_i^-(-\vec{k})) \tag{A.14}$$

as the functional of the boundary spinors χ_i^+ and $\bar{\chi}_i^-$:

$$I \sim 2 \int d^d\vec{k} \epsilon^{-2m} \bar{\chi}_i^-(\vec{k}) [- \frac{i\hat{p}}{p} \mathcal{D}(\epsilon p) - (\frac{2ip_ip_j\hat{p}}{p^3}) \mathcal{R}(\epsilon p)] \chi_j^+ , \tag{A.15}$$

where

$$\mathcal{D}(z) = \frac{K_{m-1/2}(z)}{K_{m+1/2}(z)}, \quad \mathcal{R}(z) = \frac{[\epsilon p \frac{K_{m-1/2}(z)K_{m+3/2}(z) - K_{m+1/2}^2(z)}{K_{m+1/2}(z)} + 2K_{m-1/2}(z)]}{(2m+d-1)K_{m+1/2}(z) - 2zK_{m+3/2}(z)} . \tag{A.16}$$

We notice that to get rid of singularities one has to take terms in the square brackets $\sim \epsilon^{2m}$. Using the asymptotics of the Bessel functions for small z:

$$K_\nu(z) = z^{-\nu} 2^{\nu-1} \Gamma(\nu) - z^\nu 2^{-1-\nu} \frac{\Gamma(\nu)\Gamma(1-\nu)}{\Gamma(\nu+1)} + \cdots \tag{A.17}$$

we get the answer (4.3).

B. Momentum integrals

We used the following integrals in the above:

$$\int d^dp p_i e^{i\vec{p}\vec{x}} p^\lambda = \frac{iA_\lambda (d+\lambda) x_i}{x^{d+\lambda+2}} \tag{B.1}$$

$$\int d^dp p_i p_j p_k e^{i\vec{p}\vec{x}} p^\lambda = \frac{iA_\lambda (d+\lambda)(d+\lambda+2)}{x^{d+\lambda+4}} \times$$
$$\times [(\delta_{ik}x_j + \delta_{ij}x_k + \delta_{jk}x_i) - (d+\lambda+4) \frac{x_ix_jx_k}{x^2}] , \tag{B.2}$$

where $A_\lambda = 2^{d+\lambda} \pi^{d/2} \frac{\Gamma(\frac{d+\lambda}{2})}{\Gamma(-\frac{\lambda}{2})}$, so that $A_{\lambda+2} = -(d+\lambda)(2+d)A_\lambda$.

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