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PP-waves from nonlocal theories

Mohsen Alishahiha

Institute for Studies in Theoretical Physics and Mathematics (IPM)
P.O. Box 19395-5531, Tehran, Iran
E-mail: alishah@theory.ipm.ac.ir

Alok Kumar

Institute of Physics
Bhubaneswar 751 005, India
E-mail: kumar@iopb.res.in

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Abstract: We study the Penrose limit of ODp theory. There are two different PP-wave limits of the theory. One of them is a ten dimensional PP-wave and the other a four dimensional one. We observe the later one leads to an exactly solvable background for type-II string theories where we have both NS and RR fields. The Penrose limit of different branes of string (M-theory) in a nonzero B/E field (C field) is also studied. These backgrounds are conjectured to provide dual description of NCSYM, NCOS and OM theory. We see that under S-duality the subsector of NCSYM_4 and NCOS_4 which are dual to the corresponding string theory on PP-wave coming from NCYM_4 and NCOS_4 map to each other for given null geodesic.

Keywords: Superstring Vacua, Penrose limit and pp-wave background.

Contents

1. Introduction

Recently, PP-wave backgrounds in string theory have attracted much interest, since they possess several interesting features, such as exact solvability in Green-Schwarz formalism etc.. In fact an interesting observation in [1] is that type-IIB string theory on the 10-dimensional PP-wave background with constant five-form RR field strength is exactly solvable. Moreover it has also been shown [2] that this background is a maximally supersymmetric solution of type-IIB supergravity.

Another interesting fact is that such a $D = 10$ maximally supersymmetric PP-wave background can be obtained as a Penrose limit of another one with maximal supersymmetry, namely $AdS_5 \times S^5$ [3, 4]. On the other hand type-IIB string theory on $AdS_5 \times S^5$ is believed to be dual to the $\mathcal{N} = 4$ SYM theory in four dimensions. Indeed it has been shown [4] that taking Penrose limit of $AdS_5 \times S^5$ has a corresponding limit for the gauge theory as well. By making use of this fact the authors of [4] have been able to identify the excited string states with a class of gauge invariant operators in the large-N limit of $\mathcal{N} = 4$ SYM SU(N) theory in four dimensions. Such an identification is very difficult for $AdS_5 \times S^5$ itself mainly because the string theory in this maximally supersymmetric background has not proved to be exactly solvable yet.

Considering the fact that the string theory on a PP-wave background, obtained by taking the Penrose limit of $AdS_5 \times S^5$, is exactly solvable one might wonder if this is the case for other gravity backgrounds of string theory. Therefore it would be important to study the Penrose limit of different known supergravity solution in string theory. Actually soon after [4], several papers appeared discussing Penrose limit of different gravity solutions in string theory [5] (see also [6]).

Recently the Penrose limit of supergravity solutions which are supposed to be dual to some non-local theories have been studied in [7], where the authors considered little string

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theory, (0,2) theory and four dimensional noncommutative gauge theory.[1] These theories are conjectured to be dual to the theories on the worldvolume of NS5-brane in type-II string theories, M5-branes and D3-brane with B field in their decoupling limit, respectively.

The aim of this paper is to further study the PP-waves from non-local theories. The theories we shall consider is ODp-theory, NCOS and OM theory [13, 14, 15]. The corresponding supergravity solutions have also been studied in [15]–[21].

Our work is also motivated by a recent analysis of "time dependent" PP-wave background [22] (see also [23, 24]), where particle motion has been shown to be exactly solvable, with fixed points of the hamiltonian being related to the time-independent harmonic motion of massive particles. In addition, non-constant 'masses' appear in the Green-Schwarz worldsheet theory and turn out to be of importance in undestanding the RG flow in the corresponding gauge theory. More precisely, different fixed points of the renormalization group flow in guage theory side, give rise to different expressions for masses and thereby determine the validity of the supergravity description. In this paper we will also present several examples of non-constant 'time dependent', masses and discuss their implications for our solutions.

The organization of the paper is as following. In section 2 we will study the Penrose limit of ODp theory. We shall see that there are two different limits, one of which leads to a background in which the string theory can be exactly solved. In section 2 the Penrose limit of $NCOS_4$ is considered where we observe that the resulting PP-wave is S-dual to one coming from $NCSYM_4$. In section 4 we shall consider the Penrose limit of OM-theory. The last section is devoted to the conclusions and comments.

2. Penrose limit of type-II NS5-branes in the presence of RR field

There are decoupled theories (ODp) on the worldvolume of type-II NS5-branes in the presence of nonzero p-form RR filed [13, 15].[2] The excitations of these theories include light open Dp-branes. The gravity description of these theories have been studied in [19, 20].

In this section we will consider the Penrose limit of these theories. We shall consider two different null geodesics for these backgrounds leading to two PP-wave solutions in type-II string theories.

The gravity description of ODp theory is given by

$$\begin{aligned}
l_s^{-2}ds^2 &= (1+a^2r^2)^{1/2}\left[-dt^2+\sum_{i=1}^{p}dx_i^2+\frac{\sum_{j=p+1}^{5}dx_j^2}{1+a^2r^2}+\frac{N}{r^2}(dr^2+r^2d\Omega_3^2)\right],\\
A_{0\cdots p} &= \frac{l_s^{(p+1)}}{\tilde g}a^2r^2\,,\qquad A_{(p+1)\cdots 5}=\frac{l_s^{(5-p)}}{\tilde g}\,\frac{a^2r^2}{1+a^2r^2}\,,\\
e^{2\phi} &= \tilde g^2\frac{(1+a^2r^2)^{(p-1)/2}}{a^2r^2},\qquad l_s^2dB=2N\epsilon_3\,,\quad a^2=\frac{l_{\rm eff}^2}{N}\,,
\end{aligned}\tag{2.1}$$

[1] For definition of these theories see [8]–[12].

[2] For $p=1,2$ see also [18].

where $l_{\rm eff}$ and $\tilde{g}$ are effective string tension and effective string coupling of the theory which are the parameters of the theory after taking the decoupling limit [13]. Here ϵ_3 is the volume of S^3 part of the metric.

To study the Penrose limit of the above gravity solution we will first rescale $t \to \sqrt{N}t$ and consider a null geodesic in the (t, r, ψ) plane at the fixed point with respect to other coordinates.[3] This geodesic is generated by tangent vector $\dot{t}\partial_t + \dot{r}\partial_r + \dot{\psi}\partial_\psi$, where dot denotes derivative with respect to the affine parameter. Since ∂_t and ∂_ψ are killing vectors, they define constant of motion along geodesic. In fact, defining $h = 1 + a^2r^2$,

$$E = h^{1/2}\dot{t}\,, \qquad J = h^{1/2}\dot{\psi} \tag{2.2}$$

are conserved quantities corresponding to the energy and angular momentum. For a null geodesic we get

$$\dot{r}^2 = \frac{r^2}{h}(1 - l^2)\,, \tag{2.3}$$

where $l = J/E$ and we have also rescaled the affine parameter by E. One could solve this differential equation to find u as a function of r. This could be inverted to find $r(u)$ and then we can plug this into the other expressions in the metric to write the metric in terms of u.

Consider a coordinate transformation from $(t, r, \psi) \to (u, v, x)$ which is more suitable coordinate to write the metric. The coordinate transformation is given by

$$\begin{aligned} dr &= \frac{r}{\sqrt{h}}(1 - l^2)^{1/2}du\,, \\ dt &= \frac{1}{\sqrt{h}}du - dv + l dx\,, \\ d\psi &= \frac{l}{\sqrt{h}}du + dx\,. \end{aligned} \tag{2.4}$$

Substituting this change of coordinate in the metric and rescaling the coordinate as following

$$u \to u\,, \qquad v \to \frac{v}{N}\,, \qquad \theta \to \frac{z}{\sqrt{N}}\,, \qquad x \to \frac{x}{\sqrt{N}}\,, \tag{2.5}$$

with all other coordinates remaining fixed, and then taking the Penrose limit ($N \to \infty$) we find the PP-wave limit metric of ODp-theory as following

$$\begin{aligned} ds^2 &= 2dudv - \frac{l^2}{\sqrt{h}}z^2du^2 + h^{1/2}\Big((1 - l^2)dx^2 + d\vec{z}^2 + dy^2 + y^2 d\Omega_{p-1}^2\Big) + \\ &\quad + h^{-1/2}(dw^2 + w^2 d\Omega_{4-p}^2) \end{aligned} \tag{2.6}$$

in this metric h is a function of u. This is the PP-wave in Rosen coordinate.

To find the PP-wave in Brinkman we use the following transformation

$$u \to u\,, \qquad x \to \frac{x}{h^{1/4}\sqrt{1 - l^2}}\,, \qquad z \to \frac{z}{h^{1/4}}\,, \qquad y \to \frac{y}{h^{1/4}}\,, \qquad w \to h^{1/4}w \tag{2.7}$$

[3]Here we parameterize the 3-sphere in (2.1) as $d\Omega_3^2 = d\theta^2 + \cos^2\theta\, d\psi^2 + \sin^2\theta\, d\beta^2$.

and

$$v \to v - \frac{1}{2}\frac{\partial_u h^{1/4}}{h^{1/4}}(x^2+y^2+z^2-w^2)\,. \tag{2.8}$$

Using this transformation the metric reads

$$ds^2 = 2dudv + \left(F_1\,(x^2+y^2+z^2) + F_2 w^2 - \frac{l^2}{h}z^2\right)du^2 + dx^2 + d\vec{y}^2 + d\vec{z}^2 + d\vec{w}^2\,, \tag{2.9}$$

where

$$F_1 = \frac{\partial_u^2 h^{1/4}}{h^{1/4}}\,, \qquad F_2 = 2\left(\frac{\partial_u h^{1/4}}{h^{1/4}}\right)^2 - \frac{\partial_u^2 h^{1/4}}{h^{1/4}}\,. \tag{2.10}$$

We note, however, that this is not the whole story. We need to know which fields survive in this limit. In fact there is a free parameter in the theory whose scaling behavior is to be taken into account. Namely, we need to know the behavior of effective string length in the large-N limit. Indeed keeping $l_{\rm eff}$ fixed in large-N limit means that $a \to 0$ and therefore $h \to 1$. In this case, for the fixed scale where we have $l = 1$, we just get a PP-wave which is the same as if we had considered the Penrose limit of NS5-branes, so we get

$$\begin{aligned} ds^2 &= -2dudv - \mu^2 z^2 du^2 + d\vec{z}^2 + d\vec{Y}_6^2\,, \\ H_{uz_1z_2} &= 2\mu\,, \end{aligned} \tag{2.11}$$

with constant dilaton. Here we have properly rescaled du and dv by μ. This is exactly the NS PP-wave which was derived in [25]. This background represents an exact string background to all orders in the worldsheet theory [26, 27]. This background has also been studied in [28] (see also [7]). It is also easy to see that the solution (2.11) does solve supergravity equations of motion. To see this we note that the μ dependence appears nontrivially only in the Ricci tensor R_{++}, where one obtains a constant contribution of $2\mu^2$, which is canceled by $2\mu^2$ contribution coming from B field.

We note, however, that since in the solution (2.1) a always appears in the combination of ar it is natural to consider a case in which the effective deformation parameter [16], ar, remains fixed in the Penrose limit. In this case, setting $ar = e^{\rho}$, the PP-wave obtained from the null geodesic (2.3) becomes

$$\begin{aligned} ds^2 &= 2dudv + \left(F_1\,(x^2+y^2+z^2) + F_2 w^2 - \frac{l^2}{h}z^2\right)du^2 + dx^2 + d\vec{y}^2 + d\vec{z}^2 + d\vec{w}^2\,, \\ e^{2\phi} &= \tilde{g}^2\frac{(1+e^{2\rho})^{(p-1)/2}}{e^{2\rho}}\,, \qquad A_{(p+1)\cdots 5} = \frac{l_s^{(5-p)}}{\tilde{g}}\,\frac{e^{2\rho}}{1+e^{2\rho}}\,, \\ dA^{(p+1)} &= \frac{2l(1-l^2)^{1/2}}{\tilde{g}}\,\frac{e^{2\rho}}{\sqrt{1+e^{2\rho}}}\,dx\wedge du\wedge dx_1\wedge\cdots\wedge dx_p\,, \\ dB &= \frac{2l}{\sqrt{1+e^{2\rho}}}\,du\wedge dz_1\wedge dz_2\,, \end{aligned} \tag{2.12}$$

where F_1 and F_2 are now given by

$$F_1 = \frac{(1-l^2)}{4}\,\frac{e^{2\rho}(4-e^{2\rho})}{(1+e^{2\rho})^3}\,, \qquad F_2 = \frac{(1-l^2)}{4}\,\frac{e^{2\rho}(3e^{2\rho}-4)}{(1+e^{2\rho})^3}\,. \tag{2.13}$$

Here ρ is a function of u given by

$$\sqrt{1+e^{2\rho}} - \tanh^{-1}\sqrt{1+e^{2\rho}} = (1-l^2)u\,. \tag{2.14}$$

It can be shown that this PP-wave solves the corresponding type-II supergravity equations of motion.

Since the PP-wave we found is time dependent (u-dependent) one can study the RG flow of the theory. For example, from point of view of the 2-dimensional quantum theory on string worldsheet the masses of fields are given by $-F_1$ and $-F_2$. But from the expression of F_1 and F_2, we observe that in UV or IR some directions get imaginary masses. This means the gravity description we started with is not a good one and we need to change our description which can be done either using S-duality (in type IIB) or lifting the theory to M-theory (in type IIA).

We now present another plane wave solution which is obtained by taking Penrose limit corresponding to a longitudinal null geodesic which also leads to an exactly solvable string theory independent of what a is. For this, we perform a rescaling of t and x_i, $(i=1,\ldots,5)$ by a factor $\sqrt{N}$ and make a coordinate transformation: $ar = e^{\rho}$. Then the metric as well as other fields in (2.1) read:

$$\begin{aligned}
ds^2 &= l_s^2 N h^{1/2}\Bigg[-dt^2 + \sum_{i=1}^{p} dx_i^2 + h^{-1}\sum_{j=p+1}^{5} dx_j^2 + d\rho^2 + \\
&\qquad\qquad + d\theta^2 + \sin^2\theta\, d\beta^2 + \cos^2\theta\, d\psi^2\Bigg]\,,\\
A_{0\cdots p} &= \frac{l_s^{(p+1)}}{\tilde g} N^{(p+1)/2} e^{2\rho}\,, \qquad A_{(p+1)\cdots 5} = \frac{l_s^{(5-p)}}{\tilde g}\,\frac{N^{(5-p)/2}e^{2\rho}}{1+e^{2\rho}}\,,\\
e^{2\phi} &= \tilde g^2\,\frac{(1+e^{2\rho})^{(p-1)/2}}{e^{2\rho}}\,, \qquad l_s^{-2}dB = 2N\epsilon_3\,.
\end{aligned} \tag{2.15}$$

Now consider a null geodesic around $\rho=\rho_0$=constant. In this case we perform the following coordinate transformation:

$$\begin{aligned}
\rho &= \rho_0 + (1+e^{2\rho_0})^{-1/4}\frac{r}{\sqrt{N}}\,, & x'_j &= (1+e^{2\rho_0})^{1/4}\frac{w_j}{\sqrt{N}}\,, \quad (j=p+1,\ldots,5)\,,\\
\theta &= (1+e^{2\rho_0})^{-1/4}\frac{z}{\sqrt{N}}\,, & x_i &= (1+e^{2\rho_0})^{-1/4}\frac{y_i}{\sqrt{N}} \quad (i=1,\ldots,p)\,,\\
t &= (1+e^{2\rho_0})^{-1/4}\left(x^+ + \frac{x^-}{N}\right), & \psi &= (1+e^{2\rho_0})^{-1/4}\left(x^+ - \frac{x^-}{N}\right),
\end{aligned} \tag{2.16}$$

and β is fixed. One then obtains in the Penrose limit ($N\to\infty$) a PP-wave background as following

$$ds^2 = -4dx^+dx^- - \mu^2 z^2 dx^{+2} + d\vec{z}^2 + d\vec{Y}_6^2\,, \tag{2.17}$$

where $d\vec{Y}_6^2 = dr^2 + d\vec{w}^2 + d\vec{y}^2$. Here we have also rescaled the longitudinal coordinates by $x^\pm \to x^\pm(1+e^{2\rho_0})^{\pm 1/4}$. In this limit, one gets also a nonvanishing $p+1$ form field and nonzero B field:

$$\begin{aligned} dA^{(p+1)} &= \frac{2\mu}{\tilde{g}}\, e^{2\rho_0}(1+e^{2\rho_0})^{-(p+1)/4}\, dx^+ \wedge dr \wedge dy_1 \wedge \cdots \wedge dy_p\,, \\ A_{(P+1)\cdots 5} &= \frac{1}{\tilde{g}}\frac{e^{2\rho_0}}{1+e^{2\rho_0}}\, dw_{(P+1)} \wedge \cdots \wedge dw_5\,, \\ dB &= 2\mu(1+e^{2\rho_0})^{-1/2}\, dx^+ \wedge dz_1 \wedge dz_2\,, \end{aligned} \tag{2.18}$$

and the constant dilaton after the Penrose limit is given by:

$$e^{2\phi} = \tilde{g}^2\, \frac{(1+e^{2\rho_0})^{(p-1)/2}}{e^{2\rho_0}}\,. \tag{2.19}$$

Here we have properly rescaled dx^+ and dx^- by μ. One can again check that this solution solves the type-II supergravities equations of motion. To see this we note that the μ dependence appears nontrivially only in the Ricci tensor R_{++}, where one obtains a constant contribution of $2\mu^2$, which is canceled by $2\mu^2(1+e^{2\rho_0})^{-1}$ contribution coming from B field and $2\mu^2 e^{2\rho_0}(1+e^{2\rho_0})^{-1}$ contribution from RR (p+1)-form.

We have therefore obtained a four dimensional PP-wave background from ODp theories. The form of the background is independent of the value of p one started with. The metric obtained for this background is very similar to that of NS PP-wave (2.11), though here we also have nonzero RR field. One can show that these backgrounds also correspond to exactly solvable string theories.

To see this one can write down the GS action for the superstring in this background. In the light-cone gauge the bosonic action leads to two massive fields (z_1, z_2) and six free scalar fields $\vec{Y}_6$. The fermion mass comes from the RR field and is dependent on the matrix $\Pi \equiv \gamma^1 \cdots \gamma^p$.

To be precise the bosonic 2-dimensional worldsheet action in the presence of non-zero B-field is given by

$$S_{\rm bos} = -\frac{1}{4\pi\alpha'}\int d\sigma^2 \left[\eta^{ab} G_{\mu\nu}\partial_a x^\mu \partial_b x^\nu + \epsilon^{ab} B_{\mu\nu}\partial_a x^\mu \partial_b x^\nu\right], \tag{2.20}$$

where $\eta = {\rm diag}(-1,1)$ is the worldsheet metric and we use a notation in which $\epsilon^{\tau\sigma} = 1$.

Plugging the R-NS PP-wave (2.17) and (2.18) in this action, one finds

$$\begin{aligned} S_{\rm bos} = -\frac{1}{4\pi\alpha'}\int d\sigma^2 \Big[&\eta^{ab}\left(2\partial_a x^+ \partial_b x^- + \partial_a z_i \partial_b z_i - \mu^2 z_i^2 \partial_a x^+ \partial_b x^+ + \partial_a \vec{Y}\partial_b \vec{Y}\right) + \\ &+ 2\mu(1+e^{2\rho_0})^{-1/2} z_2 \epsilon^{ab}\left(\partial_a x^+ \partial_b z_1 - \partial_a z_1 \partial_b x^+\right)\Big], \end{aligned} \tag{2.21}$$

where we have rescaled x^+ by μ. The action can be simplified using the light-cone gauge. In this gauge, setting, $x^+ = p^+\tau$, the action (2.21) reads

$$S_{\rm bos} = -\frac{1}{4\pi\alpha'}\int d\sigma^2 \left(\eta^{ab}\partial_a z_i \partial_b z_i + m^2 z_i^2 + \partial_a \vec{Y}\partial_b \vec{Y} + \frac{4m}{\sqrt{1+e^{2\rho_0}}} z_2 \partial_\sigma z_1\right), \tag{2.22}$$

where $m = \mu p^+$. The equations of motion are given by

$$\begin{aligned}
\eta^{ab}\partial_a\partial_b\vec{Y} &= 0\,,\\
\eta^{ab}\partial_a\partial_b z_2 - m^2 z_2 - \frac{2m}{\sqrt{1+e^{2\rho_0}}}\,\partial_\sigma z_1 &= 0\,,\\
\eta^{ab}\partial_a\partial_b z_1 - m^2 z_1 + \frac{2m}{\sqrt{1+e^{2\rho_0}}}\,\partial_\sigma z_2 &= 0
\end{aligned} \tag{2.23}$$

subject to the following boundary conditions

$$\partial_\sigma\vec{Y}\delta\vec{Y}|_{\rm boundary} = 0\,, \qquad \partial_\sigma z_2\delta z_2|_{\rm boundary} = 0\,, \qquad \partial_\sigma z_1\delta z_1 - \frac{2m}{\sqrt{1+e^{2\rho_0}}}\, z_2\delta z_1|_{\rm boundary} = 0\,. \tag{2.24}$$

Now we can proceed to solve the equation of motion subject to the above boundary condition. To do this one considers an ansatz as $z_i = \alpha_i e^{i(\omega_n\tau+n\sigma)}$ and $\vec{Y} = \vec{\beta}e^{i(\omega_n\tau+n\sigma)}$. Then the bosonic frequencies are given by

$$\begin{aligned}
\omega_n &= \pm\sqrt{m^2+n^2\pm 2mn(1+e^{2\rho_0})^{-1/2}}\,, &&\text{for } z_1, z_2\\
\omega_n &= \pm n\,, &&\text{for } \vec{Y}\,.
\end{aligned} \tag{2.25}$$

The general form of the solution is once again similar to the that discussed in the literature [1] (see also [29]–[34]).

Following [1] one can write the quadratic fermionic action as following

$$\frac{i}{\pi}\int d^2\sigma\left(\eta^{ab}\delta_{pq} - \epsilon^{ab}(\sigma_3)_{pq}\right)\partial_a x^\mu\bar{\theta}^p\Gamma_\mu\mathcal{D}_b\theta^q\,, \tag{2.26}$$

where θ^p, $p = 1,2$ are two 10 dimensional spinors with same/different chiralities in type IIB/A, and $\sigma_3 = \mathrm{diag}(1,-1)$. The generalized covariant derivative, $\mathcal{D}$ is given by [35, 36]

$$\mathcal{D}_a = \partial_a + \frac{1}{4}\partial_a x^\rho\left[\left(\omega_{\mu\nu\rho} - \frac{1}{2}H_{\mu\nu\rho}\sigma_3\right)\Gamma^{\mu\nu} + \left(\frac{1}{2!}F_{\mu\nu}\Gamma^{\mu\nu}\sigma_0 + \frac{1}{2\times 4!}F_{\mu\nu\lambda\delta}\Gamma^{\mu\nu\lambda\delta}\sigma_1\right)\Gamma_\rho\right], \tag{2.27}$$

for type-IIA string theory, and

$$\mathcal{D}_a = \partial_a + \frac{1}{4}\partial_a x^\rho\left[\left(\omega_{\mu\nu\rho} - \frac{1}{2}H_{\mu\nu\rho}\sigma_3\right)\Gamma^{\mu\nu} + \left(\frac{1}{48}F_{\mu\nu\eta}\Gamma^{\mu\nu\eta}\sigma_1 + \frac{1}{480}F_{\mu\nu\lambda\delta\eta}\Gamma^{\mu\nu\lambda\delta\eta}\sigma_0\right)\Gamma_\rho\right], \tag{2.28}$$

for type-IIB string theory.[4]

In the light-cone gauge we set $x^+ = p^+\tau$, $\Gamma^+\theta^p = 0$, then the non-zero contribution to the (2.26) comes only from the term where both the "external" and "internal" $\partial_a x^\mu$ factors become $p^+\delta^\mu_+\delta^0_a$. Therefore the action in the light-cone gauge reads

$$-\frac{ip^+}{\pi}\int d\sigma^2\,\bar{\theta}^p\Gamma_+\left(\delta_{pq}\mathcal{D}_\tau + (\sigma_3)_{pq}\mathcal{D}_\sigma\right)\theta^q\,, \tag{2.29}$$

where the supercovariant derivatives are given by

$$\begin{aligned}
\mathcal{D}_\tau &= \partial_\tau + \frac{1}{4}p^+\left[\left(\omega_{\mu\nu+} - \frac{1}{2}H_{\mu\nu+}\sigma_3\right)\Gamma^{\mu\nu} + \left(\frac{1}{2}F_{\mu\nu}\Gamma^{\mu\nu}\sigma_0 + \frac{1}{48}F_{\mu\nu\lambda\delta}\Gamma^{\mu\nu\lambda\delta}\sigma_1\right)\Gamma_+\right],\\
\mathcal{D}_\sigma &= \partial_\sigma\,,
\end{aligned} \tag{2.30}$$

[4]Here we have fixed the numerical coefficient of the RR fields as those in [30].

for type-IIA string theory. Here σ_i's are Pauli matrices. Similarly one can write the supercovariant derivatives for type IIB in the light-cone gauge. We note that in our case the gravity solutions (2.15) have just one RR field which gives a mass terms for fermions.

As we see the fermion mass comes from the term corresponding to the RR $(p+1)$-form field. One can simply write down the equations of motion and solve them, though we will not study the fermionic solution in this paper.

As a particular example one can consider the case of $p = 5$. The background (2.1) will then provide the supergravity description of "New" six-dimensional gauge theory [37] (see also [38]). The supergravity description of new gauge theory has been studied in [39]. While we were typing our paper we received paper [33] where this model has been studied in detail.

3. Penrose limit of NCOS in four dimensions

The worldvolume theory of D-brane in the presence of B field with one leg along the time direction (E field) is an interacting theory which includes open string. These theories are known as noncommutative open string theory or in short NCOS [13]. The supergravity description of these theories have been studied in [17]. The corresponding gravity solution for NCOS in four dimensions is given by [14]

$$\begin{aligned} ds^2 &= R^2 h^{1/2}\left[r^2\left(-dt^2+dx_1^2+h^{-1}(dx_2^2+dx_3^2)\right)+\frac{dr^2}{r^2}+d\Omega_5^2\right] \\ e^{2\phi} &= g_s h\,, \qquad h = 1+a^4r^4\,, \\ B_{01} &= a^4r^4\,, \quad A_{23} = \frac{a^4r^4}{1+a^4r^4}\,, \end{aligned} \tag{3.1}$$

where $a^4 \sim b/\tilde{g}_s N$ with b is the noncommutative parameter and $\tilde{g}_s$ is effective string coupling appears in the theory after taking the decoupling limit.

The $NCOS_4$ is S-dual to $NCSYM_4$ [14]. Therefore the corresponding supergravity solutions also map to each other under S-duality. The Penrose limit of $NCSYM_4$ has been studied in [7], the resulting PP-wave metric in the Brinkman form of $NCSYM_4$ is given by

$$ds^2 = 2dudv - l^2(x^2+z^2+y_1^2+Fy_2^2)du^2+dx^2+d\bar{z}^2+dy_1^2+dy_2^2\,, \tag{3.2}$$

where F is a known function of u, for which we do not need an explicit form in our consideration.

According to the proposal of [4], there must be a subsector of $NCSYM_4$ theory which describes type-IIB string theory in this PP-wave background. Of course, neither this background may be a solvable one for type IIB, nor the corresponding gauge theory description may be tractable. Nevertheless one could ask if the relation between $NCOS_4$ and $NCSYM_4$ is still the same after Penrose limit. In other words, if there is a subsector of $NCOS_4$ describing type-IIB string theory in a PP-wave background coming from the Penrose limit of (3.1) which maps to the subsector of $NCSYM_4$ describing type-IIB string theory on the PP-wave background (3.2).

We note, however, that we can take several null geodesic each giving a PP-wave in both NCSYM and NCOS. But starting with a null geodesic in NCSYM theory, one can find only one null geodesic in NCOS which map together under S-duality. These two Penrose limits then give subsectors of NCSYM and NCOS theories which map togther under S-duality. To demonstrate this procedure let us consider the Penrose limit of (3.1).

Consider a null geodesic in the (t, r, ψ) plane at the fixed point with respect to other coordinates. This geodesic is generated by tangent vector $\dot{t}\partial_t + \dot{r}\partial_r + \dot{\psi}\partial_\psi$, where dot denotes derivative with respect to the affine parameter. Since ∂_t and ∂_ψ are killing vectors, they define constant of motion along geodesic, in fact

$$E = h^{1/2} r^2 \dot{t}\,, \qquad J = h^{1/2}\dot{\psi} \tag{3.3}$$

are conserved quantities corresponding to the energy and angular momentum. For a null geodesic we get

$$\dot{r}^2 = \frac{1}{h} - \frac{r^2 l^2}{h}\,, \tag{3.4}$$

where $l = J/E$ and we have also rescaled the affine parameter by E.

Consider a coordinate transformation from $(t, r, \psi) \to (u, v, x)$, which is more suitable coordinate to write the metric. The coordinate transformation is given by:

$$\begin{aligned} dr &= \left(\frac{1}{h} - \frac{r^2 l^2}{h}\right)^{1/2} du\,, \\ dt &= \frac{1}{r^2 h^{1/2}} du - dv + l dx\,, \\ d\psi &= \frac{l}{h^{1/2}} du + dx\,. \end{aligned} \tag{3.5}$$

Substituting this change of coordinate in the metric we get

$$\begin{aligned} ds^2 = R^2 \Big[& 2dudv - \frac{l^2}{h^{1/2}} \sin^2\theta\, du^2 - 2l \sin^2\theta\, dudx + h^{1/2}(\cos^2\theta - l^2 r^2) dx^2 + \\ & + 2l r^2 h^{1/2} dv dx - r^2 h^{1/2} dv^2 + r^2 h^{1/2} dw^2 + \\ & + r^2 h^{-1/2}(dy^2 + y^2\, d\alpha^2) + h^{1/2}(d\theta^2 + \sin^2\theta^2\, d\Omega_3^2) \Big]\,. \end{aligned} \tag{3.6}$$

Note that here we have also changed the coordinates as

$$\begin{aligned} dx_1^2 &\to dw^2\,, \\ dx_2^2 + dx_3^2 &\to dy^2 + y^2\, d\alpha^2\,. \end{aligned} \tag{3.7}$$

We now rescale the coordinate as following

$$u \to u\,, \qquad v \to \frac{v}{R^2}\,, \qquad \theta \to \frac{z}{R}\,, \qquad x \to \frac{x}{R}\,, \qquad w \to \frac{w}{R}\,, \qquad y \to \frac{y}{R}\,, \tag{3.8}$$

and all other coordinates remain fixed. Taking the Penrose limit in which $R \to \infty$ and keeping a fixed, we find the PP-wave metric of NCOS$_4$ as following

$$\begin{aligned} ds^2 = & \, 2dudv - \frac{l^2}{h^{1/2}} z^2 du^2 + h^{1/2}(1 - l^2 r^2) dx^2 + h^{1/2}(dz^2 + z^2 d\Omega_2^3) + \\ & + r^2 h^{1/2} dw^2 + r^2 h^{-1/2}(dy^2 + y^2 d\alpha^2) \end{aligned} \tag{3.9}$$

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in this metric h is a function of u. This is the PP-wave in Rosen coordinate. To find the PP-wave in Brinkman we use the following transformation

$$u \to u\,, \qquad x \to \frac{x}{h^{1/4}\sqrt{1-l^2r^2}}\,, \qquad z \to \frac{z}{h^{1/4}}\,, \qquad w \to \frac{w}{rh^{1/4}}\,, \qquad y \to \frac{h^{1/4}y}{r}\,, \tag{3.10}$$

and

$$v \to v + \frac{1}{2}\frac{\partial_u(h^{1/4}\sqrt{1-l^2r^2})}{h^{1/4}\sqrt{1-l^2r^2}}x^2 + \frac{1}{2}\frac{\partial_u h^{1/4}}{h^{1/4}}z^2 + \frac{1}{2}\frac{\partial_u(rh^{1/4})}{rh^{1/4}}w^2 + \frac{1}{2}\frac{\partial_u(r/h^{1/4})}{r/h^{1/4}}y^2\,. \tag{3.11}$$

Using this transformation the metric reads

$$ds^2 = 2dudv + (F_w w^2 + F_x x^2 + F_y y^2 + F_z z^2)du^2 + dw^2 + dx^2 + d\vec{y}^2 + d\vec{z}^2\,, \tag{3.12}$$

where

$$\begin{aligned} F_w &= \tfrac{\partial_u^2(rh^{1/4})}{rh^{1/4}}\,, & F_y &= \tfrac{\partial_u^2(r/h^{1/4})}{r/h^{1/4}}\,, \\ F_z &= -\tfrac{l^2}{h} + \tfrac{\partial_u^2 h^{1/4}}{h^{1/4}}\,, & F_x &= \tfrac{\partial_u^2(h^{1/4}\sqrt{1-l^2r^2})}{h^{1/4}\sqrt{1-l^2r^2}}\,, \end{aligned} \tag{3.13}$$

and one can show that $F_w = F_z = F_x$.

We note, however, that one could also consider the Penrose limit of the theory for case where the noncommutative parameters is kept fixed in the limit of $N \to \infty$. In this case $a \to 0$, thus $h \to 1$. Therefore the PP-wave solution we get is exactly the one coming from $AdS_5 \times S^5$.

Under S-duality, the metric in string frame changes as $ds^2 \to e^{-\phi}ds^2$, and therefore performing S-duality on the PP-wave solution (3.9) one gets

$$\begin{aligned} ds^2 &= 2h^{-1/2}dudv - \frac{l^2}{h}z^2du^2 + (1-l^2r^2)dx^2 + (dz^2 + z^2 d\Omega_2^3) + \\ &\quad + r^2dw^2 + r^2h^{-1}(dy^2 + y^2d\alpha^2)\,. \end{aligned} \tag{3.14}$$

On the other hand under S-duality the null geodesic is also changed such that the affine parameter of the null geodesic changes as $du \to h^{1/2}du$. Using this map the above PP-wave solution reads

$$\begin{aligned} ds^2 &= 2dudv - l^2z^2du^2 + (1-l^2r^2)dx^2 + (dz^2 + z^2 d\Omega_2^3) + \\ &\quad + r^2dw^2 + r^2h^{-1}(dy^2 + y^2d\alpha^2)\,, \end{aligned} \tag{3.15}$$

which is the PP-wave metric of NCSYM_4 in the Rosen coordinates, and moreover goes to the same form as (3.2) in the Brinkman coordinates. Note also that the null geodesic equation also maps to $\dot{r}^2 = 1 - l^2r^2$.

More generally for a PP-wave solution with string metric in the Rosen coordinates as

$$ds^2 = 2dudv + G(u)dx^2\,, \tag{3.16}$$

one can write the solution in the Brinkman form as following

$$ds^2 = 2dudv + \frac{1}{2}x^2F(G)du^2 + dx^2\,, \tag{3.17}$$

where F(G) is a function of G and its derivative. Under S-duality the metric changes to

$$ds^2 = 2dudv + \frac{1}{2}x^2\tilde{F}(G)du^2 + dx^2\,, \tag{3.18}$$

where

$$\tilde{F}(G) = F(G) + \frac{1}{2}(\partial_u\phi)^2 - \frac{\partial_u(G\partial_u\phi)}{G}\,, \tag{3.19}$$

here ϕ is the dilaton.

We therefore conclude that under S-duality the subsector of NCOS$_4$ describing type-IIB string theory in a PP-wave background, coming from the Penrose limit of (3.1), maps to a subsector of NCSYM$_4$ describing type-IIB string theory on the PP-wave background (3.2). Note that under S-duality the corresponding null geodesics are also mapped to each other. In fact only for these geodesics one can compare the corresponding subsectors. The Penrose limit of other NCOS can also be done in the same way as presented in this section.

4. Penrose limit of OM theory

The worldvolume theory of M5-brane in the presence of C-field is decoupled from the bulk gravity leading to a nontrivial interacting theory with light open membrane. This theory has been introduced in [13] and called OM-theory. Starting with N M5-branes on top of each other with a nonzero C-field parametrized by b, at the decoupling limit we find the gravity solution corresponding to OM-theory as following [16]

$$\begin{aligned}
l_p^{-2}ds^2 &= 4R^2h^{1/3}\Big[r^2\left(-dt^2 + dx_1^2 + dx_2^2 + h^{-1}(dx_3^3 + dx_4^2 + dx_5^2)\right) + \frac{dr^2}{r^2} + \\
&\qquad + \frac{1}{4}\left(d\theta^2 + \cos^2\theta\, d\psi^2 + \sin^2\theta\, d\Omega_2^2\right)\Big] \\
h &= 1 + \frac{b^3}{R^3}r^6\,, \qquad dC = l_p^3 N\epsilon_4\,, \\
C_{012} &= l_p^3\frac{b^{3/2}}{R^3}r^6\,, \qquad C_{345} = l_p^3\frac{b^{3/2}}{R^3}\frac{r^6}{1+\frac{b^3}{R^3}r^6}\,,
\end{aligned} \tag{4.1}$$

where ϵ_4 is the volume of S^4 part of the above metric.

In this section we would like to consider the Penrose limit of OM theory. To do this, in the same way as in the previous sections, let us consider a null geodesic in the (t, r, ψ) plane at the fixed point with respect to other coordinates. This geodesic is generated by tangent vector $\dot{t}\partial_t + \dot{r}\partial_r + \dot{\psi}\partial_\psi$, where dot denotes derivative with respect to the affine parameter. Since ∂_t and ∂_ψ are killing vectors, they define constant of motion along geodesic, in fact

$$E = h^{1/3}r^2\dot{t}\,, \qquad J = \frac{h^{1/3}}{4}\dot{\psi} \tag{4.2}$$

are conserved quantities corresponding to the energy and angular momentum. For a null geodesic we get

$$\dot{r}^2 = \frac{1}{h^{2/3}} - \frac{4r^2l^2}{h^{2/3}}\,, \tag{4.3}$$

Now consider a coordinate transformation from $(t, r, \psi) \to (u, v, x)$ which is more suitable coordinate to write the metric. The coordinate transformation is given by

$$\begin{aligned} dr &= \left(\frac{1}{h^{2/3}} - \frac{4r^2l^2}{h^{2/3}}\right)^{1/2} du\,, \\ dt &= \frac{1}{r^2h^{1/3}}du - dv + l dx\,, \\ d\psi &= \frac{4l}{h^{1/3}}du + dx\,. \end{aligned} \tag{4.4}$$

One can use the first equation to find r as a function of u

$$u = \int \frac{h^{1/3}dr}{\sqrt{1-4r^2l^2}} \tag{4.5}$$

which could be inverted to find $r(u)$. Then we can plug this to the other expression in the metric to write the metric (4.1) in terms of the u. Substituting this change of coordinate in the metric (4.1), and rescaling the coordinate as following:

$$u \to u\,, \qquad v \to \frac{v}{R^2}\,, \qquad \theta \to \frac{z}{R}\,, \qquad x \to \frac{x}{R}\,, \qquad w \to \frac{w}{R}\,, \qquad y \to \frac{y}{R}\,, \tag{4.6}$$

with all other coordinates remaining fixed, we obtain by taking the Penrose limit ($R \to \infty$ while b/R fixed), the PP-wave metric of OM theory in Rosen coordinate as following:

$$\begin{aligned} ds^2 &= 2dudv - \frac{4l^2}{h^{1/3}}z^2du^2 + h^{1/3}\left(\frac{1}{4} - l^2r^2\right)dx^2 + \frac{h^{1/3}}{4}(dz^2 + z^2d\Omega_2^2) + \\ &\quad + r^2h^{1/3}(dw^2 + w^2d\alpha^2) + r^2h^{-2/3}(dy^2 + y^2d\tilde{\Omega}_2^2)\,. \end{aligned} \tag{4.7}$$

Note that in the above we have changed the coordinates as

$$\begin{aligned} dx_1^2 + dx_2^2 &\to dw^2 + w^2\,d\alpha^2\,, \\ dx_3^2 + dx_4^3 + dx_5^2 &\to dy^2 + y^2\,d\tilde{\Omega}_2^2\,. \end{aligned} \tag{4.8}$$

In the same way, as in previous sections, one can write the PP-wave in the Brinkman coordinates which is:

$$ds^2 = 2dudv + \left(F_w w^2 + F_x x^2 + F_y y^2 + F_z z^2\right)du^2 + d\vec{w}^2 + dx^2 + d\vec{y}^2 + d\vec{z}^2\,, \tag{4.9}$$

where

$$\begin{aligned} F_w &= \frac{\partial_u^2(rh^{1/6})}{rh^{1/6}}\,, & F_y &= \frac{\partial_u^2(r/h^{1/3})}{r/h^{1/3}}\,, \\ F_z &= -\frac{16l^2}{h^{2/3}} + \frac{\partial_u^2h^{1/6}}{h^{1/6}}\,, & F_x &= \frac{\partial_u^2\left(h^{1/6}\sqrt{\frac{1}{4} - l^2r^2}\right)}{h^{1/6}\sqrt{\frac{1}{4} - l^2r^2}}\,. \end{aligned} \tag{4.10}$$

One can show once again that $F_w = F_z = F_x$. We also note that in this limit we get only one C field corresponding to the N M5-branes as following

$$(dC)_{u,\vec{z}} = 32l_p^3h^{-5/6}\,, \tag{4.11}$$

all others go to zero in the Penrose limit.

Note that in getting the above PP-wave we have assumed that the noncommutative parameter b is also going to infinity in large-N limit. But we could consider the case in which the noncommutative parameter remains fixed in the Penrose limit. In this case $h \to 1$ and then we just get maximally supersymmetric 11-dimensional PP-wave.

5. Conclusions

In this paper we have studied the PP-wave limit of ODp, NCOS and OM theories. We have shown that there are two different Penrose limits one can take for ODp theories in which one of them leads to an exactly solvable theory. The metric of this PP-wave is very similar to that considered by Nappi-Witten [25]. But, besides the NS B field we also have an RR field in this background. This R-NS PP-wave also gives an exact solution of worldsheet string theory. We have studied the light-cone GS action for this model.

We have shown that under S-duality the subsector of $NCOS_4$ which is dual to type-IIB string theory on PP-wave of NCOS, maps to the subsector of $NCSYM_4$, dual to type-IIB string theory on PP-wave of NCSYM. In fact, this statement is true only for those null geodesics which are S-dual to each other.

An interesting fact about the Penrose limit is that it can wash away the noncommutativity effects. More precisely, if we assume that the noncommutativity parameter is kept fixed in taking the Penrose limit, the PP-wave solution we get is the same as if we had considered the Penrose limit of the corresponding commutative theory. According to BMN [4] proposal, this means, for example, that type-IIB string theory on maximally supersymmetric PP-wave background has a gauge theory description in terms of a subsector of non-commutative SYM theory in four dimensions with 16 supercharges. We however leave the detail of this correspondence for further study.

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