Infinite Lorentz boost along the M-theory circle and non-asymptotically flat solutions in supergravities

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Abstract

Certain non-asymptotically flat but supersymmetric classical solution of the type IIA supergravity can be interpreted as the infinitely-boosted version of the D-particle solution along the M-theory circle. By a chain of $T$-dual transformations, this analysis also applies to yield non-asymptotically flat solutions from the asymptotically flat and (non)-extremal solutions with intersecting D-strings and D five-branes of the type IIB supergravity compactified on a five-torus. Under $S$-duality, the non-asymptotically flat solutions in this context can in particular be used to describe the $(1+1)$-dimensional CGHS type black holes via spontaneous compactifications.
I. INTRODUCTION

D-branes and string dualities revolutionized our understanding of the non-perturbative aspects of the string theory and black hole physics [1]. Black holes in various types of supergravities, which are in turn the low energy limit of the string theories, are connected to each other via string dualities, and in some cases microscopic descriptions in terms of D-branes are possible [2]. The emergence of M-theory, whose low energy limit is believed to be described by the eleven dimensional supergravity, considerably enriches these recent developments. The most notable current approach to realize the quantum version of M-theory is the Matrix theory [3]. In its formulation, choosing the infinite momentum frame along the extra M-theory circle is a key technical tool.

In the long distance limit, D-branes are described by the black hole type brane solutions of supergravities. In obtaining these solutions, it is conventionally required that the space-time is asymptotically flat. In [4], however, it has been shown that by taking the light-cone compactification for each of these D-brane solutions, it is possible to find non-asymptotically flat and supersymmetric solutions which result also from appropriately setting a number of constants of integration [1]. Some of these non-asymptotically flat solutions, via $U$-duality, give the geometry of lower dimensional, in particular two and three dimensional, black holes [1]. The appearance of the lower dimensional black holes tensored with a torus and/or sphere has been observed in a different context. For the NS five-branes, it was noted in Ref. [6] that the metric near the horizon approaches the two dimensional black hole of Callan-Giddings-Harvey-Strominger (CGHS) [7].

The relationship between the D-brane solutions and the non-asymptotically flat solutions is the main focus of this paper; in the context of the eleven dimensional supergravities, we find that they are related to each other by an infinite Lorentz boost along the M-theory circle. We illustrate this property in section II for the simplest case of D-particles. Our re-

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1This kind of phenomena has been observed in a slightly different context in [3] as well.
results in section II are not new in a sense that the infinite Lorentz boost along the M-theory
circle has been recently investigated in [8], and our presentation adopts the results of that
reference. The novelty here is our explanation that the difficulties usually associated with
non-asymptotically flat solutions can be resolved in the context of the eleven dimensional
supergravity and under the infinite Lorentz boost. Via a chain of $T$-dualities, the consider-
atation in section II can be extended to arbitrary D $p$-branes. In section III, we consider the
application of the infinite boost for the D 5-branes in the case where we have intersecting
D-strings and D 5-branes with a Kaluza-Klein momentum along the circle where D-strings
are wrapped. For each asymptotically flat and (non)-extremal solution, we show that we
can obtain a non-asymptotically flat solution which has the same space-time structure near
the horizon as the asymptotically flat (non)-extremal solutions. Furthermore, we find that
these non-asymptotically flat, (non)-extremal solutions of the type IIB supergravity turn
into lower dimensional black holes, such as the CGHS black holes, under $S$-duality via a
spontaneous compactification. When combined with the Lorentz transformation along the
M-theory circle, the $U$-dual multiplets of the black holes in supergravities get enlarged to
include black holes in differing space-time dimensions.

II. PRELIMINARY: INFINITELY-BOOSTED D-PARTICLE SOLUTIONS

The D-particle solution of the ten-dimensional type IIA supergravity can be written as
follows. We have the ten-dimensional metric

$$ds_{10}^2 = -\frac{1}{\sqrt{f}} dt^2 + \sqrt{f} (dx_1^2 + \cdots + dx_9^2),$$

(1)

and the dilaton $\phi$ and the R-R one-form gauge field $A$ are given by

$$e^{2\phi} = f^{3/2}, \quad A_t = 1 - \frac{1}{f},$$

where

$$f = 1 + \frac{Q}{R_s^2 r^7}. $$
Here $f$ represents a nine-dimensional harmonic function and $r^2 = x_1^2 + \cdots + x_9^2$. The integer $Q$ counts the number of D-particles, and we are using a unit where the eleven dimensional Einstein action has the unit coefficient in front of the curvature tensor. The ten dimensional gravitational constant is therefore inversely proportional to the eleventh circle radius $R_s$, and the energy of one D-particle is given by $E = 1/(gl_s) = 1/R_s$, where $g$ is the string coupling constant and $l_s$ is the ten dimensional string scale.

The inclusion of the first term, 1, in the harmonic function $f$ is necessary if we require that the space-time is asymptotically flat. On the other hand, under a different choice of the constants of integration, we have solutions of the form

$$ds_{10}^2 = -\frac{1}{\sqrt{h_0}}dt^2 + \sqrt{h_0}(dx_1^2 + \cdots + dx_9^2)$$

along with

$$e^{2\phi} = h_0^{3/2}, \quad A_t = -\frac{1}{h_0}$$

where $h_0$ is another harmonic function given by

$$h_0 = \frac{k}{r^7}.$$ 

This solution, although it satisfies the ten-dimensional equations of motion and supersymmetric, is clearly non-asymptotically flat. There are at least two serious problems that prevent us from regarding the solution (2) as a physically sensible one. First, the limit $k \to 0$ is singular. In contrast, in the case of the D-particle solution (1), the limit $Q \to 0$ simply represents the flat vacuum solution. Secondly, and more seriously, the ADM type mass of the solution (2) is difficult to define due to the non-asymptotic flatness in the long distance limit, unlike the D-particle solution. In fact, the second problem is related to the first problem, since ideally we hope to compute the energy of non-zero $k$ solution relative to the $k = 0$ solution.

Even with these difficulties, the solutions (2) are interesting; for the dynamical processes in the background geometry of (2), the ADM type mass might be defined to give a finite
answer for the mass, as is familiar from the lower dimensional gravity theories. Furthermore
they are sensible locally; the near horizon behavior \( r \to 0 \) limit is the same as the usual
D-particle solutions. The more detailed explanation of the relationship between these two
solutions is the theme of this subsection. In what follows, we will find that the satisfactory
answers to the two issues raised in the above can naturally be given in the context of the type
IIA/M-theory. A clear hint is provided by the inspection of the form of the functions \( f \) and
\( h_0 \). In the large charge and fixed distance limit, the form of the function \( f \) asymptotically
becomes that of the function \( h_0 \). Therefore, in view of the fact that the D-particle number \( Q \)
is the quantized Kaluzu-Klein momentum along the M-theory circle, the natural expectation
is that the solution (2) is the infinitely boosted version of the solution (1) along the M-theory
circle. To see this more explicitly, we lift the the D-particle solution (1) into the eleven-
dimensional supergravity and rewrite it as

\[
\begin{align*}
\text{ds}_{11}^2 &= e^{-2\phi/3} \text{ds}_{10}^2 + e^{4\phi/3} (dx_{11} - A_t dt)^2 \\
&= -\frac{1}{f} dt^2 + f(dx_{11} - (1 - \frac{1}{f})dt)^2 + dx_1^2 + \cdots + dx_9^2 ,
\end{align*}
\]

where we use the standard dual description between the M-theory on a circle and the IIA
supergravity. In our convention, the M-theory circle, parameterized by \( x_{11} \), has the period
of \( R_s \) and we identify

\[
x_{11} \simeq x_{11} + R_s , \quad t \simeq t
\]

to produce the two dimensional cylinder parameterized by \((x_{11}, t)\). \(^2\) For our subsequent
consideration, it is convenient to introduce a set of asymptotic light-cone coordinates \( x^\pm =
x_{11} \pm t \) and rewrite the metric Eq.(3) in terms of \( x^\pm \) as

\[
\text{ds}_{11}^2 = dx^+ dx^- + h dx^- dx^- + dx_1^2 + \cdots + dx_9^2
\]

\(^2\)We closely follow the notation and the idea of Ref. \( ^8 \) in our paper with a slight modification of
the boost parameter due to our coordinate choice.
where \( h = f - 1 \).

Under a Lorentz boost along the M-theory circle given by the boost parameter

\[
\beta = \frac{R}{\sqrt{R^2 + 4R_s^2}}
\]

the original spatial circle with the period \( R_s \) approaches the light-cone circle with the period \( R \), as we take the limit \( \beta \to 1 \) and \( R_s \to 0 \) while keeping \( R \) fixed. To see this, we note that the unboosted identification under the lattice translation

\[
x^+ \simeq x^+ + R_s, \quad x^- \simeq x^- + R_s
\]

changes into

\[
x^+ \simeq x^+ + \frac{2}{1 + \sqrt{1 + 4R_s^2/R^2}} \frac{R_s^2}{R}, \quad x^- \simeq x^- + \frac{1 + \sqrt{1 + 4R_s^2/R^2}}{2} \frac{1}{R}
\]

under the boost. We also find that the effect of the Lorentz boost on the metric corresponds to the substitution

\[
h \to h_0 = \frac{4}{(1 + \sqrt{1 + 4R_s^2/R^2})^2} \frac{R_s^2}{R^2} h
\]

in Eq. (4). Under the infinite-boost, the spatial circle becomes a light-cone circle with the radius \( R \) and we have a discrete light cone where we identify \( x^- \simeq x^- + R \). Thus, to obtain a ten-dimensional, i.e., the type IIA side form of the infinitely boosted solutions, we have to compactify Eq.(3) along the asymptotic light-cone circle. Upon the dimensional reduction along the light-cone circle \( x^- \) with the radius \( R \), the equations (3) with \( h \) given in Eq. (3) reduces to Eq.(2) with

\[
k = Q/R^2.
\]

The ten dimensional time coordinate \( t \) now should be identified with the light-cone time \( t = x^+/2 \) as is clear from the construction. As expected, this consideration shows that the non-asymptotically flat solution (2) is the infinitely-boosted version of the D-particle solution along the M-theory circle.
The original D-particle solutions have the light-cone energy \( p_+ = 0 \), due to the zero transversal momentum and the zero rest mass of the graviton, and the momentum \( p_- = Q/R_s \). After the infinite boost, the momentum \( p_- \) transforms into \( p_- = Q/R \). This can be mostly clearly seen if we write the eleven dimensional form of the (infininitely-boosted) solution Eq.(2) as

\[
\begin{align*}
\Delta s^2_{11} &= dx^+dx^- + h_0dx^2_+ + dx^2_1 + \cdots + dx^2_9 \\
&\simeq dx^+dx^- + \frac{P}{r^7}\delta(x^-)dx^2_- + dx^2_+ + \cdots + dx^2_9.
\end{align*}
\]

(7)

Here \( p_- = Q/R \) represents the momentum of the graviton moving along the asymptotic light-cone circle, \( h_0 = Q/(R^2r^7) \), and we replaced \( 1/R \) factor with the delta function defined on the finite range \( 0 \leq x^- \leq R \). Eq. (8) is the eleven dimensional Aichelberg-Sexl metric \[11\] representing the gravitational shock-wave travelling along the light-cone circle with the momentum \( p_- \), as was pointed out in Ref. \[9\]. In the Kaluza-Klein dimensional reduction from eleven dimensions to ten dimensions, zero mode part of the solution Eq.(7) is the same as that of Eq.(8). The difference comes only from the massive higher modes, justifying our approximate identification of these two equations.

Our consideration so far provides us with the answers to the questions we posed earlier. First, when we take the limit \( k \to 0 \), the eleven dimensional version Eq.(7) of the solution Eq.(2) indeed becomes a eleven dimensional flat solution, even if Eq.(2) has a singular limit. We note that we write Eq.(7) as

\[
\Delta s^2_{11} = h_0^{-1/2}( -\frac{1}{\sqrt{h_0}}d(\frac{x^+}{2})^2 + \sqrt{h_0}(dx^2_1 + \cdots + dx^2_9)) + h_0(dx^- + \frac{1}{h_0}d(\frac{x^+}{2}))^2.
\]

for the dimensional reduction along the light-cone circle. Thus, the singular limits of the ten dimensional dilaton, graviton and graviphoton fields combine to give a well-behaved eleven dimensional limit, as we take \( k \to 0 \) (thereby \( h_0 \to 0 \)). Just as the \( Q = 0 \) D-particle solution gives the eleven dimensional flat space-time with a compact spatial circle, \( k = 0 \) non-asymptotically flat solution Eq.(2) gives the eleven dimensional flat space-time with
a compact light-cone circle. Secondly, we now have a better understanding of the energy of the non-asymptotically flat solutions Eq.(2). In the case of the D-particles, the energy computed from the eleven dimensional perspective is just the energy \( E = Q/R_s \) of the graviton travelling along the spatial circle, which is conjugate to the time coordinate \( t \) in Eq.(3). This time coordinate \( t \) is identical to the ten dimensional time and we thus recover the energy of D-particles, \( E = Q/(gl_s) \). In the similar spirit, since the solutions Eq.(7) are also asymptotically flat from the eleven dimensional point of view, we can unambiguously compute the light-cone energy \( p_+ = 0 \), both before and after the boost. Since the ten dimensional time in the case of non-asymptotically flat solutions corresponds to \( x^+/2 \), its conjugate energy is again \( 2p_+ = 0 \). This argument shows that the assignment of zero energy to the non-asymptotically flat solutions Eq.(3) in ten dimensions is a natural one.

**III. TYPE IIB SUPERGRAVITY ON FIVE-TORUS WITH INTERSECTING D-STRINGS AND D 5-BRANES**

Following the procedure explained in Section II, we can turn an asymptotically flat solution to a non-asymptotically flat one via the infinite boost along the M-theory circle. Using a chain of \( T \)-dualities, we can obtain corresponding non-asymptotically flat solutions for any D \( p \)-brane solutions, given our analysis of the D-particle solutions. The case of interest in this section, in particular, is the non-asymptotically flat solutions from the (non)-extremal solutions of the type IIB supergravity on a five-torus with intersecting D-strings and D 5-branes and with the Kaluza-Klein momentum along the circle where D-strings are wrapped. These configurations of D-branes were used in the computation of the black hole entropy via the D-brane technology, and they provide one of the simplest setting to discuss the quantum black hole physics. \[12\] By performing \( U \)-dual transformations to these well-understood configurations, we can better understand other configurations of D-branes and NS-branes. The question is then the action of the infinite boost along the M-theory circle on these configurations.
We consider D-strings wrapped along the circle parameterized by $x_9$, D 5-branes wrapped on the five-torus parameterized by $x_5, \ldots, x_9$, and the Kaluza-Klein momentum moving along the circle $x_9$. The charges produced by D-branes in the non-compact five dimensions ($t, x_1, \ldots, x_4$) are related to two real numbers $r_1$ and $r_5$, respectively. The metric, dilaton, and R-R two-form gauge field $A_{\mu\nu}$ are given by [2]

$$ds_{10}^2 = \frac{1}{\sqrt{f_1 f_5}}(-dt^2 + dx_9^2 + K(\cosh \sigma dt - \sinh \sigma dx_9)^2)$$

$$+ \sqrt{f_1 f_5}(\frac{1}{1-K}dr^2 + r^2d\Omega_3^2) + \sqrt{\frac{f_1}{f_5}}(dx_5^2 + \cdots + dx_8^2),$$

$$e^{-2\phi} = \frac{f_5}{f_1}, \quad A_{tx_9} = 1 - \frac{1}{f_1}, \quad (dA)_{ijk} = \frac{1}{2}\epsilon_{ijkl}\partial_l f_5$$

where we introduce three functions

$$f_1 = 1 + \frac{r_1^2}{r^2}, \quad f_5 = 1 + \frac{r_5^2}{r^2}, \quad K = \frac{r_0^2}{r^2}.$$ Here greek indices and latin indices are ten-dimensional and non-compact spatial four-dimensional, respectively. For non-compact spatial dimensions, we use the radial coordinate $r$ which becomes $r^2 = x_1^2 + x_2^2 + x_3^2 + x_4^2$ in the extremal limit, and $d\Omega_3^2$ denotes the unit three-sphere. For the extremal solutions, the real numbers $\sigma$ and $r_0$ take the limit $\sigma \to \infty$ and $r_0 \to 0$ while keeping $r_0^2 \sinh 2\sigma$ fixed.

We now consider the change of the function $f_5$ of D 5-branes under the analysis of the section II. For this purpose, we first apply T-dual transformations to all of the compact coordinates from $x_5$ to $x_9$. These transformations turn the D 5-branes into D-particles, the D-strings into D 4-branes wrapped on a four torus parameterized by $x_5, \ldots, x_8$, and the Kaluza-Klein momentum into fundamental strings winding the $x_9$ circle. Secondly, we lift the resulting ten dimensional type IIA solution into the eleven dimensional solution. From the eleven dimensional perspective, we have longitudinal 5-branes (originally D-strings), longitudinal membranes (originally Kaluza-Klein momentum), and the momentum along the
M-theory circle (originally D 5-branes). Straightforward calculation of the transformations from Eq. (9) shows that we have

\[ ds_{11}^2 = \frac{1}{f_1^{1/3}(1 + K \sinh^2 \sigma)^{2/3}} \left\{ (1 + \frac{K}{2f_5})dx^+dx^- + (h + \frac{K}{4f_5})dx^-dx^- + \frac{K}{4f_5}dx^+dx^+ \right\} \]  

\[ + \frac{f_1^{2/3}}{(1 + K \sinh^2 \sigma)^{2/3}}dx_5^2 + f_1^{2/3}(1 + K \sinh^2 \sigma)^{1/3}(\frac{1}{1 - K}dr^2 + r^2d\Omega_3^2) \]

\[ + (1 + K \sinh^2 \sigma)^{1/3} \left( dx_5^2 + \ldots dx_8^2 \right) \]

for the eleven dimensional metric where we introduce \( h = f_5 - 1 \), quite similar to Eq. (7).

Now we perform the infinite-boost along the M-theory circle while taking \( R_s \to 0 \). As in section II, \( h \) represents D-particles and thus \( h \) is of the order of \( O(R_s^{-2}) \). Since \( K \) is the contribution from the Kaluza-Klein momentum along the \( x_9 \) direction, we can assume that the value of its charge is smaller than the order of \( O(R_s^{-1}) \). Under the infinite boost, the expression in the curly bracket of Eq. (10) transforms into

\[ (1 + \frac{K}{2f_5})dx^+dx^- + \frac{R_s^2}{R^2}(h + \frac{K}{4f_5})dx^-dx^- + \frac{R_s^2}{R^2} \frac{K}{4f_5}dx^+dx^+ \]  

\[ \simeq dx^+dx^- + h_5 dx^-dx^- + \frac{K}{4h_5}dx^+dx^+. \]

In obtaining the second line from the first line in the above equation, we retain only the leading order term in \( R_s \), which turns out to be the zeroth order in \( R_s \). We note that \( h_5 = (R_s^2/R^2)h \) is the zeroth order in \( R_s \). Since the M-branes in our consideration are all longitudinal, the functions \( f_1 \) and \( K \) do not change. The additional effect of this boost is to change the spatial circle into the light-cone circle along the \( x^- \) direction. Thus, we plug Eq. (11) into Eq. (10) and compactify the boosted eleven dimensional solution along the light-cone circle to get the corresponding type IIA solution. As a final step, after the T-dual transformations to each of the circles of the five-torus, we go back to the original type IIB theory. The fields now look
\[
 ds_{10}^2 = \frac{1}{\sqrt{f_1 h_5}} (-dt^2 + dx_5^2 + K(\cosh \sigma dt - \sinh \sigma dx_9)^2) 
\]

\[
 + \sqrt{f_1 h_5} \left( \frac{1}{1-K} dr^2 + r^2 d\Omega_3^2 \right) + \sqrt{\frac{f_1}{h_5}} (dx_5^2 + \cdots + dx_8^2), 
\]

\[
e^{-2\phi} = \frac{h_5}{f_1}, \quad A_{tx_9} = 1 - \frac{1}{f_1}, \quad (dA)_{ijk} = \frac{1}{2} \epsilon_{ijkl} \partial_l h_5
\]

where we use

\[
h_5 = \frac{R^2}{R^2(f_5 - 1)} = \frac{\bar{r}_5^2}{r^2}.
\]

Eq. (12) is identical to Eq.(9), except for the fact that the function \(f_5\) is changed into the function \(h_5\). For a generic non-extremal solutions, for which we can not use the BPS equations, it is not immediately apparent that Eq. (12) satisfies the field equations of the type IIB supergravity. However, we can verify this fact by straightforward calculations. Since \(h_5\) approaches zero as \(r\) goes infinity, the solution Eq. (12) is clearly not asymptotically flat.

Further properties of the non-asymptotically flat solutions can be learned if we take the \(S\)-dual transformation of the type IIB theory solution. The original D-strings and D 5-branes described by Eq.(9) become the fundamental strings and NS 5-branes, respectively, under the \(S\)-dual transformations. In the case of the solution Eq. (12), it transforms into

\[
 ds_{10}^2 = \frac{1}{f_1} (-dt^2 + dx_5^2 + K(\cosh \sigma dt - \sinh \sigma dx_9)^2) + \frac{\bar{r}_5^2}{r^2(1-K)} dr^2
\]

\[
 + \bar{r}_5^2 d\Omega_3^2 + (dx_5^2 + \cdots + dx_8^2),
\]

As a verification of this, we can resort to Ref. \([5]\) where the general \(s\)-wave sector static solutions of the type IIB supergravity on \(T^5\) for the same brane and internal momentum configurations are given. In fact, by setting \(\sinh(\sqrt{c_1}(\bar{c} - \bar{c}_1)) = 0\) for the solutions reported in that paper, we can indeed confirm that the ten dimensional field equations are satisfied for Eq. (12).
\( e^{-2\phi} = \frac{f_1}{h_5}, \quad B_{tx_9} = 1 - \frac{1}{f_1}, \quad (dB)_{ijk} = \frac{1}{2} \epsilon_{ijkl} \partial_l h_5 \)

under the \( S\)-dual transformation, where \( B_{\mu\nu} \) is the NS-NS two-form gauge field. The solution Eq. (13) is again non-asymptotically flat in the five dimensional space-time. In fact, we observe that a spontaneous compactification of the would-be non-compact coordinates occurs. The metric Eq. (13) shows that the ten dimensional manifold \( \mathcal{M}_{10} \) becomes a tensor product \( \mathcal{M}_{10} = \mathcal{M}_3 \times S^3 \times T^4 \), where \( S^3 \) is a three-sphere \( (d\Omega_3^2) \) with a fixed radius and \( T^4 \) is a four-torus parameterized by \( x_5, \cdots, x_8 \) with fixed torus moduli.

We can further compactify the metric Eq. (13) upon the \( x_9 \) circle to get the two dimensional solutions. These solutions turn out to be the two-dimensional black holes studied in Ref. [14]. Furthermore, if we set \( r_1 = \sigma = 0 \), we end up exactly recovering the CGHS black hole, where \( \tilde{r}_5^2 \) plays the role of the cosmological constant. [7] We can easily see this by comparing Eq. (13) with the metric of the CGHS model in a coordinate system where the radial coordinate \( r \) is related to the dilaton \( \phi \) via \( r = \exp(-\phi) \). When we approach the extremality with the vanishing Kaluza-Klein momentum and the vanishing \( r_1 \), we have the linear dilaton vacuum, which is the vacuum of the CGHS model and clearly has zero ADM mass. This behavior is consistent with our previous mass assignment; the mass contribution from \( \tilde{r}_5^2 \), after the infinite boost along the M-theory circle, vanishes. On the contrary, for the NS five-branes, the same conditions \( r_1 = 0 \) and the vanishing Kaluza-Klein momentum at the extremality give the ADM mass proportional to \( \tilde{r}_5^2 \) according to the BPS mass formula.

**IV. CONCLUDING REMARKS**

One interesting result from our analysis is that certain black holes in differing space-time dimensions are connected to each other via \( U\)-dualities and the eleven dimensional diffeomorphism (the infinite Lorentz boost along the M-theory circle being part of it). This aspect is consistent with the idea of [13] where it is shown that the essential features of the low energy effective string theories are universal. If this universality and the relationship between higher and lower dimensional black holes are confirmed at the level of the dynamics,
we will be able to have better understanding of this universal behavior from the relatively easier investigation of the lower dimensional gravity theories (such as the CGHS model) tensored with appropriate internal theories. In this process, some non-asymptotically flat solutions we discussed in this paper are expected to play a role, and we have shown that some bothering features of them can be understood in the framework of the type IIA/M theory. Given our results, one urgent next step will be to study the (quantum) dynamics happening on the background of the non-asymptotically flat solutions. Related to this issue is the recent work \[3\], where the solution Eq. (13) with \(r_1 = 0\) and \(\sigma = 0\) describes the near horizon dynamics of the transversal five-branes and it is shown that the semiclassical analysis is reliable in some regimes of the parameters.

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