Intersecting Non-extreme $p$-Branes and Linear Dilaton Background

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We construct the general static solution to the supergravity action containing gravity, the dilaton and a set of antisymmetric forms describing the intersecting branes delocalized in the relative transverse dimensions. The solution is obtained by reducing the system to a set of separate Liouville equations (the intersection rules implying the separability); it contains the maximal number of free parameters corresponding to the rank of the differential equations. Imposing the requirement of the absence of naked singularities, we show that the general configurations are restricted to two and only two classes: the usual asymptotically flat intersecting branes, and the intersecting branes some of which are asymptotically flat and some approach the linear dilaton background at infinity. In both cases the configurations are black. These are supposed to be relevant for the description of the thermal phase of the QFT’s in the corresponding Domain-Wall/QFT duality.

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I. INTRODUCTION

Recently it has been shown that the general supergravity solution describing a charged $p$-brane without naked singularities can be either the standard black asymptotically flat (AF) $p$-brane, or the black brane which asymptotically approaches the dilaton background (LDB). The LDB is known to be the near-horizon limit of the extremal dilatonic branes; it is relevant for the description of (non-conformal) quantum field theory (QFT) in the Domain-Wall/QFT correspondence and (in the case of NS5 branes) little string theories (LST). According to the standard reasoning, the same configuration endowed with an event horizon should describe the thermal version of the same QFT (LST). The asymptotically LDB black branes obtained must thus describe the class of quantum field theories in the thermal phase, or, in the particular case of NS5-branes, the thermodynamics of little string theory. This was extensively studied recently in the case of one-brane solutions, so it is interesting to present a more general intersecting brane framework.

The simplest example of the LDB endowed with an event horizon is the four-dimensional charged dilatonic black hole with the LDB asymptotics. This configuration was also identified as “the horizon plus throat” geometry arising in the near-horizon limit of the near extremal dilatonic black hole. The near-horizon limit of the BPS dilatonic black hole is the LDB itself, and in the theories admitting the 1/2 BPS branes the limiting configuration preserves the half of SUSY of the initial theory. The near horizon limit of black dilatonic black hole keeps the memory about the event horizon and turns out to be a non-supersymmetric configuration whose BPS limit is the LDB. The brane generalization of this construction has interesting particular cases in ten dimensions.

On the other hand, it has been also known that a more general class of solutions consisting of various kinds of black branes can be constructed within the same framework (for the time-dependent branes see ). It is then natural to ask if the above results can be extended to such general solutions. The purpose of this paper is to generalize the construction to the case of intersecting branes (see for review ). We show that the solutions are restricted to either the asymptotically flat black branes or asymptotically LDB ones and their mixed system if we impose the condition that there are no naked singularities. This generalization is interesting in that it opens the possibility of extending the Domain-Wall/QFT correspondence to more general configurations.

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This paper is organized as follows. In the next section, we start with the action for the $D$-dimensional gravity coupled to the dilaton and an arbitrary number of form fields of various ranks. We summarize the field equations, metric ansatz and the background for forms. We then derive the most general static solutions in the theory requiring the usual intersection rules for the branes. The obtained solutions have a large number of parameters. In Sec. III, we fix some of them by the requirement that the solutions do not have naked singularities, and show that the resulting solutions consist of either AF black branes or the black branes which asymptotically approach the LDB. In Sec. IV, we transform the solutions to the more familiar coordinates, and show that these reduce to the known solutions of intersecting AF branes and/or branes in LDB. Of course, the latter can be obtained from the known intersecting AF branes by taking the near-horizon limit, but a naive application of this rule produces the solutions consisting only of intersecting branes with the LDB asymptotics; our additional solutions of mixed type of AF and asymptotically LDB black branes can not be simply obtained in such a limit. Not only this, our results show that these are the only solutions which can be obtained by imposing the requirement that the space-time is free of naked singularities.

II. INTERSECTING NON-EXTREME $p$-BRANES

We consider the action describing gravity coupled to a dilaton $\phi$ and $m$ different $n_A$-form fields in $D$ dimensions

$$S = \int d^Dx \sqrt{-g} \left( R - \frac{1}{2} \partial_{\mu} \phi \partial^{\mu} \phi - \sum_{A=1}^{m} \frac{1}{2n_A} e^{a_A \phi} F^2_{[n_A]} \right).$$

All form fields are coupled to the unique dilaton with individual coupling constants $a_A$. This action is the simplest one to describe intersecting supergravity $p$-branes, it incorporates various (truncated) supergravity models for different choice of the parameters $D$, $n_A$, and $a_A$. The corresponding equations of motion read

$$R_{\mu\nu} - \frac{1}{2} \partial_{\mu} \phi \partial^{\nu} \phi = \frac{m}{2(n_A - 1)!} \left[ (F^2_{n_A})_{\mu\nu} - \frac{n_A - 1}{n_A(D - 2)} F^2_{[n_A]} g_{\mu\nu} \right] = 0,$$  \hspace{1cm} (2)

$$\partial_\mu \left( \sqrt{-g} e^{a_A \phi} F^{\nu_1 \nu_2 \cdots \nu_{n_A}} \right) = 0,$$  \hspace{1cm} (3)

$$\frac{1}{\sqrt{-g}} \partial_\mu \left( F_{\nu_1 \nu_2 \cdots \nu_{n_A}} \right) - \sum_{A=1}^{m} \frac{a_A}{2n_A} e^{a_A \phi} F^2_{[n_A]} = 0.$$  \hspace{1cm} (4)

In addition, each form field satisfies the Bianchi identity

$$\partial_{[\mu} F_{\nu_1 \nu_2 \cdots \nu_{n_A}]} = 0.$$  \hspace{1cm} (5)

The notation $(F^2_{n_A})_{\mu\nu}$ is defined as

$$(F^2_{n_A})_{\mu\nu} := F_{\mu \nu_{1} \cdots \nu_{n_A}} F^{\alpha_1 \alpha_2 \cdots \alpha_{n_A}}.$$  \hspace{1cm} (6)

The most general ansatz for the metric describing the set of intersecting non-extremal branes is

$$ds^2 = -e^{2B} dt^2 + \sum_{\rho=1}^{P} e^{2C_{\rho}} dx^2_{\rho} + e^{2A} dr^2 + e^{2D} d\Sigma_{k,\sigma}^2,$$  \hspace{1cm} (7)

where the overall transverse space is described by

$$d\Sigma_{k,\sigma}^2 = \tilde{g}_{ab} dz^a dz^b = \begin{cases} d\psi^2 + \sin^2 \psi d\Omega_{-1-k}^2, & \sigma = +1, \\ d\psi^2 + \psi^2 d\Omega_{-1-k}^2, & \sigma = 0, \\ d\psi^2 + \sinh^2 \psi d\Omega_{-1-k}^2, & \sigma = -1, \end{cases}$$  \hspace{1cm} (8)

satisfying

$$\tilde{R}_{ab} = \sigma(k - 1) \tilde{g}_{ab}.$$  \hspace{1cm} (9)

The choice of $\sigma$ corresponds to different symmetries of the overall transverse space, namely, $SO(k)$ for $\sigma = 1$, $E(k)$ for $\sigma = 0$, and $SO(1,k-1)$ for $\sigma = -1$. Two latter cases correspond to topologically non-trivial solutions which are mostly known in the (multidimensional) black hole case. In what follows we will derive the general solution valid...
for all \( \sigma \), but later on we restrict to the case \( \sigma = 1 \). For a discussion of the topological solutions (in the single brane case) see, e.g., [32, 33]. The total number of dimensions occupied by the brane world-volumes is \( p \). Each brane of the set is specified by the particular form field labeled by \( n_A \).

The Ricci-tensor for this metric has the following non-vanishing components:

\[
R_{tt} = e^{2B-2A}(B'' + B' \Lambda'),
\]

\[
R_{x\rho x\rho} = -e^{2C_{\rho'}-2A}(C_{\rho''} + C_{\rho}' \Lambda'),
\]

\[
R_{rr} = -\Lambda'' - A'' + A' \Lambda' + A'^2 - B'^2 - \sum_{\rho=1}^{p} C_{\rho}^2 - kD'^2,
\]

\[
R_{ab} = [-e^{2D-2A}(D'' + D' \Lambda') + \sigma(k - 1)] \bar{g}_{ab},
\]

where

\[
\Lambda = -A + B + \sum_{\rho=1}^{p} C_{\rho} + kD.
\]

We consider both the electrically and magnetically charged branes. In the electric case each form field is given by

\[
F_{[n_A]} = f_{A}^{\text{elec}} dt \wedge dx_{\rho_1} \wedge \cdots \wedge dx_{\rho_{n_A}} \wedge dr,
\]

which satisfies the Bianchi identity automatically. It supports the \( p_A(= n_A - 2) \)-brane in the set of intersecting branes. The equation of motion for the form fields is solved as

\[
F_{[n_A]} = h_{A} \exp \left( A + B - \sum_{\rho=1}^{p} C_{\rho} - kD + 2 \sum_{\rho=1}^{p} C_{\rho} \delta^\rho_A - a_{A} \phi \right),
\]

where \( \delta^\rho_A = 1 \) for \( x^\rho \) belonging to the world-volume of the \( p_A \)-brane.

The form fields for the magnetic branes read

\[
F_{[n_A]} = h_{A} dx_{\rho_1} \wedge \cdots \wedge dx_{\rho_{n_A}-k} \wedge \text{vol}(\Sigma_{k, \sigma}),
\]

in which case the set \( x^{\rho_i} \) does not belong to the world-volume of the \( p_A(= D - n_A - 2) \)-brane. It is easy to see that the field equations for these form fields are satisfied indeed.

To solve the Einstein equations we change the independent variable as

\[
d\tau = e^{-\Lambda} dr,
\]

and present the radial part of the metric as

\[
e^{2A} dr^2 = e^{2A} \tau^2, \quad A = \Lambda + A.
\]

The Einstein and dilaton equations to be solved are (with derivatives with respect to \( \tau \))

\[
B'' = \sum_{A=1}^{m} \frac{b_A^2(D - p_A - 3)}{2(D - 2)} e^{A g_A},
\]

\[
C_{\rho''} = \sum_{A=1}^{m} \frac{b_A^2 \delta^{(\rho)}_A}{2(D - 2)} e^{A g_A},
\]

\[
D'' = -\sum_{A=1}^{m} \frac{b_A^2(p_A + 1)}{2(D - 2)} e^{A g_A} + \sigma(k - 1)e^{2A - 2D},
\]

\[
\phi'' = -\sum_{A=1}^{m} \frac{\epsilon_{AA} b_A^2}{2} e^{A g_A},
\]

where

\[
G_A = -\epsilon_{AA} a_A \phi + 2B + 2 \sum_{\rho=1}^{p} C_{\rho} \delta^\rho_A,
\]
and we have defined

$$\delta^{(p)} = \begin{cases} D - p_A - 3, & \text{for } x^\rho \text{ belonging to } p_A\text{-brane} \\ -(p_A + 1), & \text{otherwise} \end{cases}$$ (25)

Note that this can be also written as

$$\delta^{(p)} = (D - 2)\delta^p - (p_A + 1).$$ (26)

The $rr$ component of Einstein equations gives

$$A'' - A'^2 + B'^2 + \sum_{\rho=1}^p C_{\rho}^2 + kD'^2 = -\frac{1}{2}\delta'^2 + \sum_{A=1}^m \frac{b_A^2 (D - p_A - 3)}{2(D - 2)} e^{G_A}.$$ (27)

It is straightforward to check that the following combination separates:

$$(A - D)'' = \sum_{A=1}^m \frac{b_A^2}{2(D - 2)} \left[D - p_A - 3 + \sum_{\rho=1}^p \delta^{(p)} - (k - 1)(p_A + 1)\right] e^{G_A} + \sigma(k - 1)^2 e^{2(A - D)}$$ (28)

where we have taken into account the relation

$$D - p_A - 3 + \sum_{\rho=1}^p \delta^{(p)} = (k - 1)(p_A + 1),$$ (29)

which follows from (26). The general solution of Eq. (28) is

$$H := 2(A - D) = \begin{cases} \ln \left(\frac{\beta^2}{\sqrt{\tau - \tau_0}}\right) - \ln \left[\sinh^2 \left(\frac{\beta}{2}(\tau - \tau_0)\right)\right] & \sigma = 1, \\ \beta(\tau - \tau_0), & \sigma = 0, \\ \ln \left(\frac{\beta^2}{\sqrt{\tau - \tau_0}}\right) - \ln \left[\cosh^2 \left(\frac{\beta}{2}(\tau - \tau_0)\right)\right] & \sigma = -1. \end{cases}$$ (30)

Differentiating Eq. (21) twice and substituting (20) - (28) into this, we arrive at the following coupled system of equations for the set of $G_A$:

$$G''_A = \sum_B \left(\frac{\epsilon_{ABC} B}{2} + \frac{D - p_A - 3}{D - 2} + \sum_{\rho=1}^p \delta^{(p)} \delta^B \right) b^B e^{G_B},$$ (31)

where the relation (26) is understood. The matrix on the right had side of this equation is non-diagonal generally. However if we impose the standard *intersection rule* for $A \neq B$, 19:

$$\sum_{\rho=1}^p \delta^A \delta^B = \bar{p} = \frac{(p_A + 1)(p_B + 1)}{D - 2} - 1 - \frac{\epsilon_{ABC} B}{2},$$ (32)

where $\bar{p}$ relates to the brane on which $p_A$- and $p_B$-branes intersect, the non-diagonal part vanishes. In this case the system reduces to the set of separate Liouville equations, and it is integrable. Therefore, the intersection rule is a sufficient condition for integrability. It is an interesting question whether the intersection rules is a necessary condition for integrability, see the corresponding discussion in [34].

For the diagonal terms $A = B$, the expression in the parenthesis reduces to

$$\{\cdots\} = \frac{\Delta_A}{2}, \text{ where } \Delta_A = a^2_A + \frac{2(p_A + 1)(D - p_A - 3)}{D - 2}.$$ (33)

With this assumption, one obtains the set of the decoupled Liouville equations for $G_A$:

$$G''_A = \frac{\Delta_A b^2}{2} e^{G_A},$$ (34)
which leads to the solution

\[ G_A = \ln \left( \frac{a_A}{\Delta_A b_A^2} \right) - \ln \left[ \sinh^2 \left( \frac{\alpha_A (\tau - \tau_A)}{2} \right) \right]. \]

(35)

Using this result, we can immediately integrate Eqs. (20-23) to obtain

\[ B = \sum_{A=1}^{m} \frac{D - p_A - 3}{(D - 2)\Delta_A} \left( G_A + g_A^{(1)} \tau + g_A^{(0)} \right) + B^{(1)} \tau + B^{(0)}, \]

(36)

\[ C_\rho = \sum_{A=1}^{m} \frac{\delta^{(\rho)}_A}{(D - 2)\Delta_A} \left( G_A + g_A^{(1)} \tau + g_A^{(0)} \right) + C_\rho^{(1)} \tau + C_\rho^{(0)}, \]

(37)

\[ \phi = - \sum_{A=1}^{m} \frac{\epsilon A^{\alpha A}}{\Delta_A} \left( G_A + g_A^{(1)} \tau + g_A^{(0)} \right) + \phi^{(1)} \tau + \phi^{(0)}. \]

(38)

It follows from Eq. (24) that the constants are connected by the following 2m relations:

\[ g_A^{(0,1)} + 2B^{(0,1)} + 2 \sum_{\rho=1}^{p} C_\rho^{(0,1)} \delta^{(\rho)}_A - \epsilon^{(\rho)}_A a_A \phi^{(0,1)} = 0, \]

(39)

From the solution (30) for \( A - D \) and the gauge conditions (14) and (19), we can obtain the expressions for \( A \) and \( D \):

\[ A = \frac{k}{2(2k - 1)} H - \sum_{A=1}^{m} \frac{p_A + 1}{(D - 2)\Delta_A} \left( G_A + g_A^{(1)} \tau + g_A^{(0)} \right) - \frac{B^{(1)} + C^{(1)}}{k - 1} \tau - \frac{B^{(0)} + C^{(0)}}{k - 1}, \]

(40)

\[ D = \frac{1}{2(2k - 1)} H - \sum_{A=1}^{m} \frac{p_A + 1}{(D - 2)\Delta_A} \left( G_A + g_A^{(1)} \tau + g_A^{(0)} \right) - \frac{B^{(1)} + C^{(1)}}{k - 1} \tau - \frac{B^{(0)} + C^{(0)}}{k - 1}, \]

(41)

where the following constant parameters are introduced

\[ C^{(1,0)} = \sum_{\rho=1}^{p} C_\rho^{(1,0)}. \]

(42)

Finally we have to fulfill the last equation (27). Using the intersection rules (32) and the constraints (39), this equation can be reduced to the following constant equations on the parameters:

\[ \frac{1}{2} \sum_{A=1}^{m} \frac{\alpha^2_A - (g_A^{(1)})^2}{\Delta_A} + (B^{(1)})^2 + \sum_{\rho=1}^{p} (C_\rho^{(1)})^2 + \frac{1}{k - 1} (B^{(0)} + C^{(1)})^2 + \frac{1}{2} (\phi^{(1)})^2 - \frac{k}{4(2k - 1)} \beta^2 = 0. \]

(43)

Let us count the number of free parameters. The total number of parameters appearing in the solutions is 5m + 2p + 6. It consists of 2 parameters of the function \( H (\beta, \tau_0) \), 4m parameters in \( G_A (\alpha_A, \tau_A, g_A^{(0,1)}) \), 2p + 4 parameters entering \( B, C_\rho, \phi (B^{(0,1)}, C_\rho^{(0,1)}, \phi^{(0,1)}) \) and m charge parameters \( b_A \). These have to satisfy 2m constraints (31) and one constraint (43), which can fix, for example, \( g_A^{(0,1)} \) and \( \phi^{(1)} \). Thus the remaining number of independent parameters is 3m + 2p + 5. However, not all of these 3m + 2p + 5 parameters are physical. The coordinate transformations allow us to fix some of them: by rescaling \( t, x_p \) one can absorb \( B^{(0)} \) and \( C_\rho^{(0)} \). Also, since the system of equations was autonomous, we are free to shift the coordinate \( \tau \) by a constant, so without loss of generality one can fix \( \tau_0 = 0 \). Therefore we are free to fix \( p + 2 \) and leave \( 3m + p + 3 \) physical parameters, basically \( 3m \) of \( b_A, \alpha_A \) and \( \tau_A \) and \( p + 3 \) of \( B^{(1)}, C_\rho^{(1)}, \phi^{(0)} \) and \( \beta \). Apparently, we also have a freedom of rescaling \( r \), but this is related to a change of the gauge function \( \Lambda \). So far we leave the constants in the expressions for the metric functions unfixed for later convenience.

### III. Fixing the Constants

From now on we will consider only the topologically simple case \( \sigma = 1 \). We are interested in solutions (possibly) possessing an event horizon and not plagued with naked singularities. The simple, though incomplete, way to reveal
the position of singularities and to impose the regularity condition on the horizon is to check the behavior of the Ricci scalar. Using the Einstein equations one can find the following expression for the Ricci scalar for the solution obtained:

\[ R = \frac{(k-1)^2}{2} e^{-2A} \left( \phi'^2 + \sum_{A=1}^{m} \frac{b_A^2(D-2p_A-4)}{(k-1)^2(D-2)} e^{G_A} \right). \]  

(44)

Using the explicit form of \( G_A \) in \([55]\), one can see that the points \( \tau = \tau_A \) are singular, unless some of them are zero, in which case the singularity can be avoided by imposing further conditions on the parameters. Other special points are \( \tau \to \pm \infty \) (we assume the \( \tau \)-coordinate to vary on the full real axis, and possibly to extend to the complex plane to ensure the change of signs of the exponential terms like \( e^2B \), see for details Ref. \([33]\)). These can correspond to the horizons. At the horizons the metric coefficient \( g_{AB} \) must vanish. Using Eq. \([50]\), one can see this can be the case in the limits \( \tau \to \pm \infty \) once suitable inequalities on the parameters are imposed. Combining this with the behavior of the Ricci scalar, we find, moreover, that one is free to choose for the regular event horizon any of these two limiting points, but then the other will be generically singular. We choose \( \tau \to -\infty \) as the horizon, then \( \tau \to \infty \) will be generically a null singularity, except for some special choice of parameters.

Let us investigate the behavior of the metric functions at the event horizon. Assuming without loss of generality \( \beta \geq 0, \alpha_A \geq 0 \), (we will not consider the possibility of imaginary values of these parameters which are also allowed by the overall reality of the solution), we find that, as \( \tau \to -\infty \), the functions \( G_A \) and \( H \) become linear in \( \tau \):

\[ G_A \sim \alpha_A \tau, \quad H \sim \beta \tau. \]  

(45)

This ensures vanishing of the metric component \( e^{2B} \) at the horizon.

An important further information can be extracted from the constraint equation \([21]\). It is convenient to rewrite it as follows:

\[-A'^2 + B'^2 + \sum_{\rho=1}^{p} C_{\rho}'^2 + kD'^2 + \frac{1}{2} \phi'^2 = \sum_{A=1}^{m} \frac{b_A^2}{2} e^{G_A} - \sigma k(k-1)e^H. \]  

(46)

From the radial geodesic equation, it follows that \( A' = B' \) at the horizon \([33]\), so the first two terms on the left hand side cancel. The right hand side vanishes at the horizon, so one is left with the sum of positive definite terms. Therefore, all the derivatives \( C_{\rho}', D', \phi' \) must separately vanish in the limit \( \tau \to -\infty \). More precisely, the condition \( A' = B' \) gives

\[ \frac{k}{2(k-1)} \beta - \sum_{A=1}^{m} \frac{\alpha_A + g^{(1)}_A}{\Delta_A} - \frac{B^{(1)} + C^{(1)}}{k-1} = 0. \]  

(47)

Moreover, using Eqs. \([37]\), \([38]\) and \([41]\), we find the following relations on the parameters involved, namely, from Eq. \([37]\) we obtain \( p \) relations

\[ C_{\rho}^{(1)} = \sum_{A=1}^{m} \frac{\phi_{\rho}^{(1)}}{D-2} \frac{\alpha_A + g^{(1)}_A}{\Delta_A}, \]  

(48)

from Eq. \([38]\)

\[ \phi^{(1)} = \sum_{A=1}^{m} \epsilon_A \alpha_A \frac{\alpha_A + g^{(1)}_A}{\Delta_A}, \]  

(49)

and from Eq. \([41]\)

\[ \frac{1}{2(k-1)} \beta - \sum_{A=1}^{m} \frac{p_A + 1}{D-2} \frac{\alpha_A + g^{(1)}_A}{\Delta_A} - \frac{B^{(1)} + C^{(1)}}{k-1} = 0. \]  

(50)

Consequently

\[ C^{(1)} = \sum_{A=1}^{m} \frac{D-2 - k(p_A + 1)}{D-2} \frac{\alpha_A + g^{(1)}_A}{\Delta_A}. \]  

(51)
The above constraints (47)-(50) are consistent with (43). Combining (47) and (50), one can obtain

$$B^{(1)} = \frac{1}{2} \beta - \sum_{B=1}^{m} \frac{D - p_B - 3 \alpha_B + g_B^{(1)}}{D - 2} \Delta_B,$$

(52)

Furthermore, the constraints (39), after substituting (48) and (49), become

$$B^{(1)} = \frac{1}{2} \alpha_A - \sum_{B=1}^{m} \frac{D - p_B - 3 \alpha_B + g_B^{(1)}}{D - 2} \Delta_B.$$

(53)

Hence, all the parameters \( \alpha_A \) must be equal:

$$\alpha_A = \beta.$$  

(54)

Then the metric functions read

$$B = \sum_{A=1}^{m} \frac{D - p_A - 3}{(D - 2) \Delta_A} \left( G_A - \beta \tau + g_A^{(0)} \right) + \frac{1}{2} \beta \tau + B^{(0)},$$

(55)

$$C_\rho = \sum_{A=1}^{m} \frac{\delta_A^{(\rho)}}{(D - 2) \Delta_A} \left( G_A - \beta \tau + g_A^{(0)} \right) + C_\rho^{(0)},$$

(56)

$$\phi = - \sum_{A=1}^{m} \epsilon_{A A A} \Delta_A \left( G_A - \beta \tau + g_A^{(0)} \right) + \phi^{(0)},$$

(57)

$$A = \frac{k}{2(k-1)} H - \sum_{A=1}^{m} \frac{p_A + 1}{(D - 2) \Delta_A} \left( G_A - \beta \tau + g_A^{(0)} \right) - \frac{\beta}{2(k-1)} \tau - \frac{B^{(0)} + C^{(0)}}{k - 1},$$

(58)

$$D = \frac{1}{2(k-1)} H - \sum_{A=1}^{m} \frac{p_A + 1}{(D - 2) \Delta_A} \left( G_A - \beta \tau + g_A^{(0)} \right) - \frac{\beta}{2(k-1)} \tau - \frac{B^{(0)} + C^{(0)}}{k - 1}.$$

(59)

Though the parameters \( \alpha_A \) are the same for all \( A \), the functions \( G_A \) differ in position of singularities \( \tau_A \)

**A. Asymptotically flat solutions**

The asymptotic region is located at \( \tau \to \tau_0 \), and we have fixed the translational freedom by choosing \( \tau_0 = 0 \). The corresponding behavior of the function \( H \) is

$$H \simeq \frac{1}{\tau^2}.$$  

(60)

However, there are two different cases for \( G_A \) depending on the value of \( \tau_A \). For the case \( \tau_A \neq 0 \), the asymptotic value of \( G_A \), for \( \alpha_A = \beta \), is

$$G_A \simeq G_A^{(0)} := \ln \left[ \frac{\beta^2}{\delta_A^{(0)}} \right].$$

(61)

Therefore, asymptotically the solutions reduces to

$$B \simeq \sum_{A=1}^{m} \frac{D - p_A - 3}{(D - 2) \Delta_A} \left( G_A^{(0)} + g_A^{(0)} \right) + B^{(0)},$$

(62)

$$C_\rho \simeq \sum_{A=1}^{m} \frac{\delta_A^{(\rho)}}{(D - 2) \Delta_A} \left( G_A^{(0)} + g_A^{(0)} \right) + C_\rho^{(0)},$$

(63)

$$\phi \simeq - \sum_{A=1}^{m} \epsilon_{A A A} \Delta_A \left( G_A^{(0)} + g_A^{(0)} \right) + \phi^{(0)}.$$  

(64)
The quantities $B$ and $C_\rho$ are constants which can be set to zero by rescaling of time and world-volume coordinates
imposing the conditions

\[
B^{(0)} = -\sum_{A=1}^{m} \frac{D - p_A - 3}{(D - 2)\Delta_A} \left( G_A^{(0)} + g_A^{(0)} \right), \tag{65}
\]
\[
C_\rho^{(0)} = -\sum_{A=1}^{m} \frac{\delta^{(\rho)}_A}{(D - 2)\Delta_A} \left( G_A^{(0)} + g_A^{(0)} \right). \tag{66}
\]

For the dilaton, it is common to preserve the finite value $\phi_\infty$ at infinity, and this can be ensured by the redefinition
\[
\phi^{(0)} = \sum_{A=1}^{m} \frac{\epsilon_A a_A}{\Delta_A} \left( G_A^{(0)} + g_A^{(0)} \right) + \phi_\infty. \tag{67}
\]

Then the relation (69) gives
\[
\phi_\infty = -\frac{G^{(0)}_A}{\epsilon_A a_A}. \tag{68}
\]

This equation requires the quantities $G_A^{(0)}/\epsilon_A a_A$ to be identical for all $A$, which imposes a new set of relations between $b_A$ and $\tau_A$. The Ricci scalar diverges at $\tau_A$, so, to avoid naked singularities, one has to ensure that all $\tau_A$ lie inside the event horizon for black branes. If there is also an inner horizon, one can check that the real metrics correspond to location of singularities $\tau_A$ either outside the event horizon, or inside the inner horizon. The first possibility should be ruled out. The notable exception constitutes the case of $\tau_A = 0$. In this case there is no singularity, and this point can be interpreted as spatial infinity. The physical nature of this will be clearer later when we express the solution in the Schwarzschild-type coordinate.

**B. Asymptotically LDB branes**

For the case $\tau_A = 0$, each $G_A$ has the same asymptotic behavior as $H$:
\[
G_A \simeq \ln \frac{1}{\tau^2}. \tag{69}
\]

In such case, we have
\[
B \simeq \sum_{A=1}^{m} \frac{D - p_A - 3}{(D - 2)\Delta_A} \ln \frac{1}{\tau^2}, \tag{70}
\]
\[
C_\rho \simeq \sum_{A=1}^{m} \frac{\delta^{(\rho)}_A}{(D - 2)\Delta_A} \ln \frac{1}{\tau^2}, \tag{71}
\]
\[
\phi \simeq -\sum_{A=1}^{m} \frac{\epsilon_A a_A}{\Delta_A} \ln \frac{1}{\tau^2}. \tag{72}
\]

This solution is not asymptotically flat, but asymptotically approaching the LDB.

**IV. SCHWARSCHILD-TYPE COORDINATES**

To present our solution in a more familiar form, we change the coordinates similarly to \[33\]. We map the horizon $\tau = -\infty$ to $r = r_H$,
\[
r_H = \mu^{\frac{1}{k-1}}, \quad \mu = \frac{\beta}{k-1}. \tag{73}
\]

and the internal horizon $\tau = \infty$ to $r = 0$ by choosing the gauge function
\[
\Lambda = \ln(r^{k} f), \quad \text{such that} \quad d\tau = \frac{dr}{r^{k} f}, \tag{74}
\]
where
\[ f = 1 - \frac{\mu}{r^{k-1}}. \] (75)

This corresponds to the coordinate transformation
\[ \tau = \frac{1}{(k-1)\mu} \ln f. \] (76)

Thus the region outside the event horizon \( r > r_H \) corresponds to the half axis \( (\infty, 0) \) of \( \tau \). Then we will have
\[ H = \ln \left( r^{2(k-1)} f \right). \] (77)

A. Asymptotically flat solutions

Introduce a new parameter \( q_A \) instead of \( \tau_A \) by
\[ q_A = \frac{\mu}{e^{\mu(k-1)\tau_A} - 1}. \] (78)

As we already noted, to ensure the absence of naked singularities, all \( \tau_A \) should be taken non-negative. Here we assume that \( \tau_A \) are strictly positive, so this definition leads to finite \( q_A \) (the case \( \tau_A = 0 \) will be considered below).

The function \( G_A \) then reads
\[ G_A = \ln \left( \frac{4(k-1)^2(\mu + q_A)q_A}{\Delta_A b^2_A} \frac{f}{h^2_A} \right), \] (79)

where
\[ h_A = 1 + \frac{q_A}{r^{k-1}}, \] (80)

are the harmonic functions.

In terms of the new coordinates, the metric components (7) read
\[ B = -2 \sum_{A=1}^{m} \frac{D - p_A - 3}{(D - 2)\Delta_A} \ln h_A + \frac{1}{2} \ln f, \] (81)
\[ C_\rho = -2 \sum_{A=1}^{m} \frac{\delta^{(\rho)}_A}{(D - 2)\Delta_A} \ln h_A, \] (82)
\[ \phi = 2 \sum_{A=1}^{m} \frac{\epsilon_{\rho A} q_A}{\Delta_A} \ln h_A + \phi^\infty, \] (83)
\[ A = -\frac{1}{2} \ln f + 2 \sum_{A=1}^{m} \frac{p_A + 1}{(D - 2)\Delta_A} \ln h_A, \] (84)
\[ D = \ln r + 2 \sum_{A=1}^{m} \frac{p_A + 1}{(D - 2)\Delta_A} \ln h_A, \] (85)

with a constraint following from (85):
\[ b^2_A = e^{\epsilon_{\rho A} q_A \phi^\infty} \frac{4(k-1)^2(\mu + q_A)q_A}{\Delta_A}. \] (86)

For the electric branes, the solution for the form field is
\[ f^\text{elec}_A = b_A e^{-a_A \phi^\infty} \frac{1}{r^2 h^2_A}. \] (87)

This solution corresponds to the one given in [19]. Free parameters of the solution are \( \mu, q_A \) and \( \phi^\infty \).
B. Asymptotically LDB branes

If \( \tau_A = 0 \), \( G_A \) takes the form

\[
G_A = \ln \left( \frac{4(k-1)^2}{\Delta_A b_A^2} \right)^{2(k-1)} f. \tag{88}
\]

Then the solution is, with the choice of parameters \( B^{(0)} = 0 \) and \( C_{\rho}^{(0)} = 0 \),

\[
B = 2 \sum_{A=1}^{m} \frac{D - p_A - 3}{(D - 2) \Delta_A} \ln \frac{r^{k-1}}{q_A} + \frac{1}{2} \ln f, \tag{89}
\]

\[
C_{\rho} = 2 \sum_{A=1}^{m} \frac{\delta^{(\rho)}_A}{(D - 2) \Delta_A} \ln \frac{r^{k-1}}{q_A}, \tag{90}
\]

\[
\phi = -2 \sum_{A=1}^{m} \epsilon_A q_A \ln \frac{r^{k-1}}{q_A} + \phi^\infty, \tag{91}
\]

\[
A = -\frac{1}{2} \ln f - 2 \sum_{A=1}^{m} \frac{p_A + 1}{(D - 2) \Delta_A} \ln \frac{r^{k-1}}{q_A}, \tag{92}
\]

\[
D = \ln r - 2 \sum_{A=1}^{m} \frac{p_A + 1}{(D - 2) \Delta_A} \ln \frac{r^{k-1}}{q_A}, \tag{93}
\]

where the parameters \( q_A \) are defined as

\[
q_A^{-2} = \frac{4(k-1)^2}{\Delta_A b_A^2} e^{q_A^{(0)}}. \tag{94}
\]

Our previous form strength parameter \( b_A \) is related to \( q_A \) via

\[
b_A^2 = e^{q_A a_A \phi^\infty} \frac{4(k-1)^2}{\Delta_A}. \tag{95}
\]

The solution for the electric form field is

\[
f_A^{elec} = b_A e^{-q_A \phi^\infty} \frac{r^{k-2}}{q_A^2}. \tag{96}
\]

This second possibility for \( \tau_A \) leads to an intersecting generalization of the solution of Ref. [1]. The BPS limit of these solutions corresponds to \( \mu = 0 \). It can be recognized that in this case one deals with a solution known earlier as the near-horizon limit of the BPS intersecting branes. This corresponds to the usual rule of omitting constants in the harmonic functions describing the metric. Here we reproduce these metrics as a particular case of the general supergravity solution. Our result is, however, not just to reproduce them by another technique, but to prove that no other solutions without naked singularities exist within the class of metrics considered, which is fixed by their isometries, up to the the mixed intersections below.

The black versions of these solutions with \( \mu \neq 0 \) should describe the thermal phase of QFT’s in the Domain-Wall/QFT correspondence associated to their BPS limit.

C. Mixed intersections

The mixed intersecting configurations, namely the black AF \( p \)-branes with the asymptotically LDB ones, can be obtained simply by picking up the corresponding term in the summation of \( A, B, C_{\rho}, D \) and \( \phi \) depending on whether \( p_A \) is AF brane or LDB. Again, the BPS limit of these solutions was found previously through the near-horizon considerations. It was noted that one can drop the constant term in the harmonic functions describing the system of intersecting branes not in all harmonic functions at the same time, but in some of them [35]. This corresponds to the BPS limit of our mixed intersection. To our knowledge, the black version of these solutions is new.
V. TWO INTERSECTING BRANES

We give the examples of two intersecting branes. According to the intersection rules, we can have solutions listed below:

IIA : 2 ⊥ 2(0), 4 ⊥ 4(2), 6 ⊥ 6(4), 0 ⊥ 4(0), 2 ⊥ 4(1), 2 ⊥ 6(2), 4 ⊥ 6(3), 4 ⊥ 8(4), 6 ⊥ 8(5),
IIB : 3 ⊥ 3(1), 5 ⊥ 5(3), 7 ⊥ 7(5), 1 ⊥ 3(0), 1 ⊥ 5(1), 3 ⊥ 5(2), 3 ⊥ 7(3), 5 ⊥ 7(4), 5 ⊥ 9(5),
M : M2 ⊥ M2(0), M2 ⊥ M5(1), M5 ⊥ M5(3),
NS : NS1 ⊥ NS5(1), NS5 ⊥ NS5(3), Dp ⊥ NS1(0), 0 ≤ p ≤ 8, Dp ⊥ NS5(p − 1), 1 ≤ p ≤ 6,

where those branes indicated only by numbers are D-branes, and the number in the parenthesis denotes the dimensions of the overlapping coordinates. We give explicit solutions for M and IIB theories in the following.

A. M2-M2 configurations

Let us consider the general solution for two M2-branes as an explicit example. In this case, the parameters have the values: \( D = 11, k = 5, n_A = 4, a_A = 0, p_A + 1 = 3, D - p_A - 3 = 6 \) and \( \Delta_A = 4 \). The AF solutions have the metric

\[
ds^2 = h_1^2 h_2^2 \left[ -f h_1^{-1} h_2^{-1} dt^2 + h_1^{-1} (dx_1^2 + dx_2^2) + h_2^{-1} (dx_3^2 + dx_4^2) + f^{-1} dr^2 + r^2 d\Omega_5^2 \right],
\]

where \( h_A = 1 + q_A/r^4 \), and the form fields

\[
f_{A}^{\text{elec}} = b_A/r^3 h_A^2, \quad b_A^2 = 16 (\mu + q_A)q_A.
\]

We recover the canonical intersecting two M2-branes.

For asymptotically LDB case, the solution we have derived corresponds to omitting the constant in the harmonic functions \( h_1 \) and \( h_2 \). It is

\[
ds^2 = \left( \frac{q_1}{r^4} \right)^2 \left( \frac{q_2}{r^4} \right)^2 \left[ -f \left( \frac{q_1}{r^4} \right)^{-1} \left( \frac{q_2}{r^4} \right)^{-1} dt^2 + \left( \frac{q_1}{r^4} \right)^{-1} (dx_1^2 + dx_2^2) + \left( \frac{q_2}{r^4} \right)^{-1} (dx_3^2 + dx_4^2) + f^{-1} dr^2 + r^2 d\Omega_5^2 \right],
\]

and

\[
f_{A}^{\text{elec}} = b_A/r^3 q_A^2, \quad b_A^2 = 16 q_A^2.
\]

The mixed intersecting configurations of two M2-branes is given by

\[
ds^2 = h_1^2 \left( \frac{q_2}{r^4} \right)^2 \left[ -f h_1^{-1} \left( \frac{q_2}{r^4} \right)^{-1} dt^2 + h_1^{-1} (dx_1^2 + dx_2^2) + \left( \frac{q_2}{r^4} \right)^{-1} (dx_3^2 + dx_4^2) + f^{-1} dr^2 + r^2 d\Omega_5^2 \right],
\]

where

\[
f_{A}^{\text{elec}} = b_1/r^3 h_1^2, \quad b_1^2 = 16 (\mu + q_1)q_1, \quad f_2^{\text{elec}} = b_2/r^3 q_2^2, \quad b_2^2 = 16 q_2^2.
\]

B. D1-D5 configurations

The other interesting case is the D1-D5 system. In this case, the parameters have the following values: \( D = 10, k = 3, n_A = 3, a_A = 1, \epsilon_1 = 1, \epsilon_2 = -1, p_1 + 1 = 2 = D - p_2 - 3, p_2 + 1 = 6 = D - p_1 - 3 \) and \( \Delta_A = 4 \). The AF solutions have the metric

\[
ds^2 = h_1^2 h_2^2 \left[ h_1^{-1} h_2^{-1} (-f dt^2 + dx_1^2) + h_1^{-1} (dx_2^2 + \cdots + dx_5^2) + f^{-1} dr^2 + r^2 d\Omega_5^2 \right],
\]

and the dilaton and the form fields

\[
e^{2\phi} = e^{2\phi_\infty} \frac{h_1}{h_2}, \quad f_1^{\text{elec}} = e^{-\phi_\infty} \frac{b_1}{r^3 h_1^2}, \quad b_A^2 = e^{\phi_\infty} 4 (\mu + q_A)q_A.
\]
where $h_A = 1 + q_A/r^2$. We recover the canonical intersecting D1-D5-branes in the Einstein frame.

Again, for LDB, our solution corresponds to omitting the constant in the harmonic functions $h_1$ and $h_2$. It is

$$ds^2 = \left(\frac{q_1}{r^2}\right)^{\frac{1}{2}} \left(\frac{q_2}{r^2}\right)^{\frac{1}{2}} \left[\left(\frac{q_1}{r^2}\right)^{-1} \left(\frac{q_2}{r^2}\right)^{-1} (-f dt^2 + dx_1^2) + \left(\frac{q_2}{r^2}\right)^{-1} (dx_2^2 + \cdots + dx_5^2) + f^{-1} dr^2 + r^2 d\Omega_3^2\right],$$

and

$$e^{2\phi} = e^{2\phi^\infty} \frac{q_1}{q_2}, \quad f_1^{\text{elec}} = e^{-\phi^\infty} \frac{b_1}{q_1^2}, \quad b_2^A = e^{\phi^\infty} 4 q_A^2.$$

There are also two mixed intersecting configurations of D1-D5-branes. One is given by

$$ds^2 = \left(\frac{q_1}{r^2}\right)^{\frac{1}{2}} \left(\frac{q_2}{r^2}\right)^{\frac{1}{2}} \left[\left(\frac{q_1}{r^2}\right)^{-1} (-f dt^2 + dx_1^2) + \left(\frac{q_2}{r^2}\right)^{-1} (dx_2^2 + \cdots + dx_5^2) + f^{-1} dr^2 + r^2 d\Omega_3^2\right],$$

where

$$e^{2\phi} = e^{2\phi^\infty} \frac{r^2 h_1}{q_2}, \quad f_1^{\text{elec}} = e^{-\phi^\infty} \frac{b_1}{r^3 h_1^2}, \quad b_1^2 = e^{\phi^\infty} 4 (\mu + q_1) q_1, \quad b_2^2 = e^{-\phi^\infty} 4 q_2^2,$$

and the other is

$$ds^2 = \left(\frac{q_1}{r^2}\right)^{\frac{1}{2}} \left(\frac{q_2}{r^2}\right)^{\frac{1}{2}} \left[\left(\frac{q_1}{r^2}\right)^{-1} h_2^{-1} (-f dt^2 + dx_1^2) + \left(\frac{q_2}{r^2}\right)^{-1} (dx_2^2 + \cdots + dx_5^2) + f^{-1} dr^2 + r^2 d\Omega_3^2\right],$$

where

$$e^{2\phi} = e^{2\phi^\infty} \frac{q_1}{r^2 h_2}, \quad f_1^{\text{elec}} = e^{-\phi^\infty} \frac{b_1 r}{q_1}, \quad b_1^2 = e^{\phi^\infty} 4 q_1^2, \quad b_2^2 = e^{-\phi^\infty} 4 (\mu + q_2) q_2.$$

### VI. CONCLUSIONS

In this paper we obtained the general solution for the non-BPS intersecting supergravity $p$-branes delocalized in relative transverse dimensions. This was done by fully integrating the Einstein equations with a suitable ansatz for the metric, the antisymmetric forms and the dilaton. Our solution differs from earlier ones by a larger number of free parameters, which is maximal in the present case. This, in particular, opens a way to construct individual brane solutions which are either asymptotically flat, or asymptotically approaching the LDB. For intersecting branes, imposing the conditions on parameters which ensure the absence of naked singularities, we found that the resulting configuration is an intersection of some AF branes with some asymptotically LDB branes. The intersection rules remain the same as previously known. In the case when all branes are AF, we recover the solution of [19]. When all individual branes are asymptotically LDB, we obtain the solution which corresponds to omitting constant terms in the harmonic functions involved (the overall solution being black). Apart from these two extremes, there are intersections of a number of AF branes with a number of LDB branes within the intersection rules.

The linear dilaton backgrounds were extensively used in the context of DW/QFT (NS/LST) dualities [2, 3, 5, 6]. The LDB solutions endowed with the event horizon describe the thermal phase of the QFT (LST) associated with the linear dilaton background [7, 8, 9]. Although this was not our intention here, our new intersecting solutions with LDB asymptotics may serve a basis for further studies along these directions.

The BPS versions of our solutions involving intersections with LDB branes correspond to partial omission of the constant terms in the harmonic functions involved. It has to be emphasized that our results are based on the complete integration of the Einstein equations for metric of chosen symmetry, so there are no other other solutions within this class which are free from naked singularities. Various suggestions for $p$-brane solutions with extra parameters which obey the same metric ansatz, but do not reduce to the standard BPS or black branes, generically suffer from naked singularities.

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