Logarithmic Operators Fold D branes into AdS$_3$

John Ellis[a], **N.E. Mavromatos**[a,b] and **Elizabeth Winstanley**[b]

[a] *Theory Division, CERN, CH–1211 Geneva 23, Switzerland,*

[b] *Theoretical Physics (University of Oxford), 1 Keble Road, Oxford, OX1 3NP, U.K.*

Abstract

We use logarithmic conformal field theory techniques to describe recoil effects in the scattering of two Dirichlet branes in D dimensions. In the particular case that a $D1$ brane strikes a $D3$ brane perpendicularly, thereby folding it, we find that the recoil space-time is maximally symmetric, with AdS$_3 \otimes$ E$_{D-3}$ geometry. We comment on the possible applications of this result to the study of transitions between different background metrics.

Non-perturbative techniques have provided many new insights into the theory formerly known as strings [1]. In particular, a rich variety of classical soliton solutions have been found using the technology of Dirichlet (D) branes [2]. These must be taken into account in any study of possible ground states and the spectra of excitations. A crucial ingredient in such a study is understanding the interactions between D branes and 'elementary' closed-string states, and with other D branes. In general, since recoil perturbs the classical D-brane solutions, one might expect that such interactions, would involve departures from conformal symmetry and criticality.

One example of this recoil phenomenon was given in [3], where the perturbation of a D particle by interaction with a closed-string state in flat space was considered. In the low-energy limit, where the closed-string energy $E \ll M_D$, the mass of the D particle, it was shown that the background metric was perturbed to an anti-de-Sitter (AdS) form [4]. Moreover, the perturbation was described by a logarithmic pair of world-sheet operators corresponding to the position and velocity of the struck D brane, and it was shown that momentum was conserved during the interaction and recoil process.

The purpose of this paper is to extend and generalize this previous discussion by considering a generic interaction between two extended D branes that do not have particle interpretations. Such constructions may be relevant in the context of viewing the observable world as a $D3$ brane embedded in a higher-dimensional space-time, with the extra 'bulk' dimensions being transverse to the $D3$-brane coordinates [5–7].

Remarkably, we find just one case in which the recoil space has an AdS geometry, namely AdS$_3$, which arises when a macroscopic string ($D1$ brane) moving along some bulk direction in space-time hits a Dp brane perpendicular to the direction of motion, as shown in Fig. 1. This physical problem is analogous to a straight stick hitting a flat sheet of paper, which we would expect to fold after the impact. As in the closed-string/D-particle scattering case considered previously, the D-brane/D-brane interaction is described by a logarithmic pair of world-sheet operators. In this case, they do indeed describe folding of the struck D brane. We are also able to demonstrate that momentum is conserved during the interaction, as in the previous closed-string/D-particle scattering case.

There have been many previous studies of string theory on an AdS$_3$ background [8,9], which can be described using a critical conformal field theory. Our analysis opens the way to studies of transitions between this and other string backgrounds. It confirms that such transitions are described by non-conformal logarithmic field theories. We also recall that the boundary of the AdS$_3$ background is equivalent to an appropriate coset of the $SL(2, R)$ manifold that appears in the analysis of two-dimensional black holes in string theory [10]. Perhaps the mathematical analysis presented here could be used as the basis for studies of the formation of such black-hole backgrounds, much as transitions between different two-dimensional string black holes have already been described in terms of logarithmic pairs of world-sheet operators [11]. However, the study of such a possibility lies beyond the scope of this article.
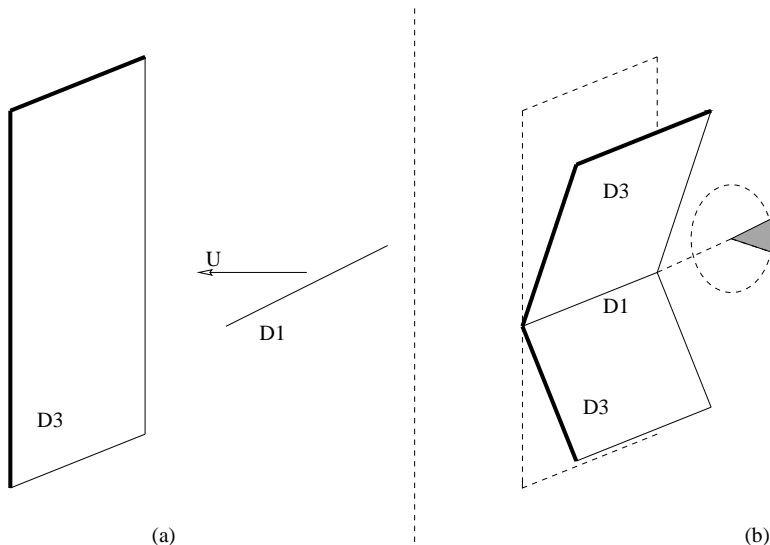
FIG. 1. *Schematic representation of the folding effect in D-brane/D-brane collisions: (a) a D1 brane moving with velocity U along a 'bulk' direction perpendicular to a D3 brane embedded in a D-dimensional Euclidean space time E_D strikes the D3 brane (b), which is then folded, and the space-time around it is distorted into $AdS_3 \otimes E_{D-3}$. The dashed circle around the D1 direction in (b) indicates the angular deficit that appears when the bulk direction along which the D1 brane was moving is compactified to a circle.*

We first review the world-sheet formalism based on logarithmic operators that was developed in a series of papers [12, 11, 13], for the mathematical description of the recoil of a D brane when struck by a closed-string state or by another D brane. Although closed-string/D-brane scattering has been studied using this formalism, we are unaware of any treatment of D-brane/D-brane scattering in the relevant literature so far. Logarithmic conformal field theory [14] lies on the border between finite conformal field theories and general renormalizable two–dimensional quantum field theories. It is the relevant tool [12,11,13] for this problem, because the recoil process involves a change of state (transition) in the string background, and as such is not described by a conformal field theory. This change of state induced by the recoil process can be described a change in the σ–model background, and as such is a non–equilibrium process. This is reflected [11, 13] in the logarithmic operator algebra itself.

As discussed in references [12, 11, 13] in the case of D–brane string solitons, their recoil after interaction with a closed-string (graviton) state is characterized by a σ model deformed by a pair of logarithmic operators [14]:

$$C_\epsilon^I = \epsilon \Theta_\epsilon(X^I), \qquad D_\epsilon^I = X^I \Theta_\epsilon(X^I), \qquad I \in \{0, \ldots, 3\} \tag{1}$$

defined on the boundary $\partial\Sigma$ of the string world sheet. Here $X^I, I \in \{0, \ldots, 3\}$ obey Neumann boundary conditions on the string world sheet, and denote the brane coordinates. The remaining $y^i, i \in \{4, \ldots, 9\}$ denote the transverse bulk directions. In the case of D particles, which were examined in [12, 11, 13], the index I takes the value 0 only, in which case the

3

operators (1) act as deformations of the conformal field theory on the world sheet. The operator $U_i \int_{\partial\Sigma} \partial_n X^i D_\epsilon$ describes the movement of the D brane induced by the scattering, where U_i is its recoil velocity, and $Y_i \int_{\partial\Sigma} \partial_n X^i C_\epsilon$ describes quantum fluctuations in the initial position Y_i of the D particle. It has been shown rigorously [13] that the logarithmic conformal algebra ensures energy–momentum conservation during the recoil process: $U_i = \ell_s g_s(k_i^1 + k_i^2)$, where $k^1(k^2)$ is the momentum of the propagating closed string state before (after) the recoil, and g_s is the string coupling, which is assumed here to be weak enough to ensure that D branes are very massive, with mass $M_D = 1/(\ell_s g_s)$, where ℓ_s is the string length.

The correct specification of the logarithmic pair in equation (3) entails a regulating parameter $\epsilon \to 0^+$, which appears inside the $\Theta_\epsilon(t)$ operator:

$$\Theta_\epsilon(X^I) = \int \frac{d\omega}{2\pi} \frac{1}{\omega - i\epsilon} e^{i\omega X^I}.$$

In order to realize the logarithmic algebra between the operators C and D, one takes [12]:

$$\epsilon^{-2} \sim \ln[L/a] \equiv \Lambda, \tag{2}$$

where L (a) are infrared (ultraviolet) world–sheet cutoffs. The recoil operators (3) are slightly relevant, in the sense of the renormalization group for the world–sheet field theory, having small conformal dimensions $\Delta_\epsilon = -\epsilon^2/2$. Thus the σ model perturbed by these operators is not conformal for $\epsilon \neq 0$, and the theory requires Liouville dressing [15, 16, 11]. Momentum conservation is assured when the Liouville field is identified with the time variable.

In the case of Dp branes, the pertinent deformations are slightly more complicated. As discussed in [12], the deformations are given by

$$\sum_I g_{Ii}^D \int_{\partial\Sigma} \partial_n X^i D_\epsilon^I \qquad \text{and} \qquad \sum_I g_{Ii}^C \int_{\partial\Sigma} \partial_n X^i C_\epsilon^I . \tag{3}$$

The $0i$ components of the two-index couplings g_{Ii}^α, $\alpha \in \{C, D\}$ include the collective momenta and coordinates of the D brane as in the D–particle case above, but now there are additional couplings g_{Ii}^α, $I \neq 0$, which describe the folding of the D brane. Such a folding may be caused by scattering with another macroscopic object, namely another D brane, propagating in a transverse direction, as shown schematically in Fig. 1 for the case of a $D1$ brane hitting a $D3$ brane. This situation is the most interesting to us, since it generates an AdS$_3$ space, as we show below. For symmetry reasons, in the situation depicted in Fig. 1, the folding of the $D3$ brane occurs symmetrically around the axis of the $D1$ brane. In this case, the precise logarithmic operator deformations shown in (3), which pertain only to the spatial region $y_i > 0$ for the Dirichlet coordinates, should be supplemented with their counterparts for the $y_i < 0$ region as well. This would, in principle, require additional $\Theta(\pm y_i)$ factors, which would complicate the analysis without introducing any new points of principle. Therefore, for simplicity, we restrict ourselves here to the $y_i > 0$ patch of space-time, away from the hypersurface $y_i = 0$. This will be implicit in what follows.

The folding couplings $g_{Ii}^D \equiv g_{Ii}$, $I \in \{0, \ldots, p\}$, $i \in \{p+1, \ldots, 9\}$, are relevant couplings with world–sheet renormalization–group β functions of the form

4

$$\beta_{g_{Ii}} = \frac{d}{dt} g_{Ii} = -\frac{1}{2t} g_{Ii}, \qquad t \sim \epsilon^{-2} \; . \tag{4}$$

This implies that one may construct an exactly marginal set of couplings $\overline{g}_{Ii}$ by redefining

$$\overline{g}_{Ii} \equiv \frac{g_{Ii}}{\epsilon} \; . \tag{5}$$

The renormalized couplings $\overline{g}_{0i}$ were shown in [13] to play the rôle of the physical recoil velocity of the D brane, while the remaining $\overline{g}_{Ii}$, $I \neq 0$, describe the folding of the Dp brane for $p \neq 0$.

As discussed in [11, 3], the deformations (1) create a local distortion of the space-time surrounding the recoiling folded D brane, which may be determined using the method of Liouville dressing. In [11, 3] we concentrated on describing the resulting space-time in the case when a D particle, embedded in a D-dimensional space time, recoils after the scattering of a closed string off the D–particle defect. To leading order in the recoil velocity u_i of the D particle, the resulting space-time was found, for times $t \gg 0$ long after the scattering event at $t = 0$, to be equivalent to a Rindler wedge, with apparent 'acceleration' ϵu_i [3], where ϵ is defined above (2). For times $t < 0$, the space-time is flat Minkowski [1].

We now discuss the generalization of this phenomenon to the case of single D-brane/D-brane scattering. For definiteness, we first consider the case where there are y_i, $i = 1, \ldots m$ transverse (bulk) dimensions (with Dirichlet boundary conditions), and X^I, $I = m + 1, \ldots D - 1$ are longitudinal coordinates on the brane, which obey Neumann boundary conditions. In this generic case, we consider the folding of the recoiling brane along several directions, determined by couplings g_{Ii} of appropriate logarithmic operators of the form (3), and the recoil-velocity vector of the D brane along the i'th bulk direction is denoted by u_i. As mentioned previously, we restrict ourselves for simplicity to the region of space-time in which $y_i > 0$.

The folding/recoil deformations of the Dp brane (3) are relevant deformations, with anomalous dimension $-\epsilon^2/2$, which disturbs the conformal invariance of the σ model, and restoration of conformal invariance requires Liouville dressing [16]. To determine the effect of such dressing on the space-time geometry, it is essential to write [11] the boundary recoil deformations as a bulk world-sheet deformations

$$\int_{\partial\Sigma} \overline{g}_{Iz} x \Theta_\epsilon(x) \partial_n z = \int_\Sigma \partial_\alpha \left(\overline{g}_{Iz}^1 x \Theta_\epsilon(x) \partial^\alpha z \right) \tag{6}$$

where the $\overline{g}_{Iz}$ denote the renormalized folding/recoil couplings (5), in the sense discussed in [13]. As we have already mentioned, such couplings are marginal on a flat world sheet. From now on, for notational simplicity, we rename $\overline{g}_{Ii} \to g_{Ii}$. The operators (6) are marginal

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[1] There is hence a discontinuity at $t = 0$, which leads to particle production and decoherence for a low-energy spectator field theory observer who performs local scattering experiments long after the scattering, and far away from the location of the collision of the closed string with the D particle [3].

also on a curved world sheet, provided [16] one Liouville-dresses the (bulk) integrand by multiplying it by a factor $e^{\alpha_{Ii}\phi}$, where ϕ is the Liouville field and α_{Ii} is the gravitational conformal dimension, which is related to the flat-world-sheet anomalous dimension $-\epsilon^2/2$ of the recoil operator, viewed as a bulk world-sheet deformation, as follows [16]:

$$\alpha_{Ii} = -\frac{Q_b}{2} + \sqrt{\frac{Q_b^2}{4} + \frac{\epsilon^2}{2}} \tag{7}$$

where Q_b is the central-charge deficit of the bulk world-sheet theory. In the recoil prblem at hand, as discussed in [3], $Q_b^2 \sim \epsilon^4/g_s^2$ for weak folding deformations g_{Ii}^1. This yields $\alpha_{Ii} \sim -\epsilon$ to leading order in perturbation theory in ϵ, to which we restrict ourselves here.

We next remark that, as the analysis of [11] indicates, the X^I-dependent field operators $\Theta_\epsilon(X^I)$ scale as follows with ϵ: $\Theta_\epsilon(X^I) \sim e^{-\epsilon X^I}\Theta(X^I)$, where $\Theta(X^I)$ is a Heavyside step function without any field content, evaluated in the limit $\epsilon \to 0^+$. The bulk deformations, therefore, yield the following σ-model terms:

$$\frac{1}{4\pi\ell_s^2}\epsilon \sum_{I=m+1}^{D-1} g_{Ii}X^I e^{\epsilon(\phi_{(0)}-X_{(0)}^I)}\Theta(X_{(0)}^I) \int_\Sigma \partial^\alpha X^I \partial^\alpha y_i \tag{8}$$

where the subscripts (0) denote world-sheet zero modes. Upon the interpretation of the Liouville zero mode $\phi_{(0)}$ as target time, the deformations (8) yield space-time metric deformations in a σ-model sense, which were interpreted in [11] as expressing the distortion of the space-time surrounding the recoiling D-brane soliton.

For clarity, we now drop the subscripts (0) for the rest of this paper, and we work in a region of space-time on the $D3$ brane such that $\epsilon(\phi - X^I)$ is finite in the limit $\epsilon \to 0^+$. The resulting space-time distortion is therefore described by the metric elements

$$G_{0i} = (\epsilon^2 y_i + \epsilon u_i t)\Theta(t) + \epsilon g_{Ii}X^I\Theta(X^I) \ , \qquad i = 1, \ldots, m, \ \ I = m+1, \ldots D-1 \tag{9}$$

to leading order in ϵg_{Ii}. The presence of $\Theta(t), \Theta(X^I)$ functions and the fact that we are working in the region $y_i > 0$ indicate that the induced space-time is piecewise continuous [2]. In the general recoil/folding case considered in this article, the form of the resulting patch of the surrounding space-time can be determined fully if one computes the associated curvature tensors, along the lines of [3]. Working to leading non-trivial order in the parameter ϵ, we find, after some tedious but straightforward calculations, the following non-zero components of the Riemann tensor (where $i \neq j$ and $I \neq J$):

$$R_{titi} = \epsilon^2\delta(t) + \frac{1}{4}\sum_{I=m+1}^{D-1} \epsilon^2 \left[\Theta(X^I) + X^I\delta(X^I)\right]^2 g_{Ii}^2 + \mathcal{O}(\epsilon^4)$$

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[2] The important implications for non-thermal particle production and decoherence for a spectator low-energy field theory in such space-times were discussed in [3, 11], where only the D–particle recoil case was considered.

$$R_{titj} = \frac{\epsilon^2}{4} \sum_{I=m+1}^{D-1} g_{Ii}g_{Ij} \left[\Theta(X^I) + X^I \delta(X^I)\right]^2$$

$$R_{titI} = \mathcal{O}(\epsilon^3), \qquad R_{tiiI} = \mathcal{O}(\epsilon^4),$$

$$R_{tItI} = \frac{\epsilon^2}{4} \sum_{j=1}^{m} g_{Ij}^2 \left[\Theta(X^I) + X^I \delta(X^I)\right]^2 + \mathcal{O}(\epsilon^4)$$

$$R_{tItJ} = \frac{\epsilon^2}{4} \sum_{j=1}^{m} g_{Ij}g_{Jj} \left[\Theta(X^I) + X^I \delta(X^I)\right] \left[\Theta(X^J) + X^I \delta(X^J)\right] + \mathcal{O}(\epsilon^4)$$

$$R_{tiIJ} = \mathcal{O}(\epsilon^3), \qquad R_{ijij} = \mathcal{O}(\epsilon^4),$$

$$R_{tIiJ} = \mathcal{O}(\epsilon^3), \qquad R_{ijjI} = \mathcal{O}(\epsilon^3),$$

$$R_{tIiI} = -\frac{\epsilon}{2}g_{Ii} \left[\delta(X^I) + X^I \delta'(X^I)\right] + \mathcal{O}(\epsilon^3)$$

$$R_{ijIJ} = \frac{\epsilon^2}{4} \left[g_{Ii}g_{Jj} - g_{Ij}g_{Ji}\right] \left[\Theta(X^I) + X^I \delta(X^I)\right] \left[\Theta(X^J) + X^J \delta(X^J)\right] + \mathcal{O}(\epsilon^4)$$

$$R_{iIiI} = -\frac{\epsilon^2}{4}g_{Ii}^2 \left[\Theta(X^I) + X^I \delta(X^I)\right]^2 + \mathcal{O}(\epsilon^4)$$

$$R_{iIjI} = -\frac{\epsilon^2}{4}g_{Ii}g_{Ij} \left[\Theta(X^I) + X^I \delta(X^I)\right]^2 + \mathcal{O}(\epsilon^4)$$

$$R_{iIiJ} = -\frac{\epsilon^2}{4}g_{Ii}g_{Ji} \left[\Theta(X^I) + X^I \delta(X^I)\right] \left[\Theta(X^J) + X^J \delta(X^J)\right] + \mathcal{O}(\epsilon^4)$$

$$R_{iIjJ} = -\frac{\epsilon^2}{4}g_{Ij}g_{Ji} \left[\Theta(X^I) + X^I \delta(X^I)\right] \left[\Theta(X^J) + X^J \delta(X^J)\right] + \mathcal{O}(\epsilon^4) . \tag{10}$$

The scalar curvature is then:

$$R = -2\epsilon^2 m\delta(t) - 2\epsilon^2 \sum_{i=1}^{m} \sum_{I=m+1}^{D-1} g_{Ii}[2\delta(X^I) + X^I \delta'(X^I)][(\epsilon y_i + u_i t)\Theta(t) + g_{Ji}X^J \Theta(X^J)]$$

$$-\frac{3\epsilon^2}{2} \sum_{i=1}^{m} \sum_{I=m+1}^{D-1} g_{Ii}^2 [\delta(X^I) + X^I \delta'(X^I)]^2 + \mathcal{O}(\epsilon^3) . \tag{11}$$

Away from the defect hypersurface $\{X^I = 0\}$, the resulting scalar curvature is negative and constant, with the form:

$$R \simeq -\frac{3}{2}\epsilon^2 \sum_{i=1}^{m} \sum_{I=m+1}^{D-1} g_{Ii}^2 + \mathcal{O}(\epsilon^4) \tag{12}$$

in the general D-brane/D-brane scattering case.

However, it is interesting to notice that the space is maximally symmetric, in the sense of the curvature being given in terms of the metric tensor as:

$$R_{abcd} = \mathcal{K} \left(g_{ac}g_{bd} - g_{ad}g_{bc}\right) \tag{13}$$

only in the case that only one of the g_{Ii} is non zero, e.g., g_{Xz}. This describes the situation depicted in Fig. 1: a $D1$ brane is moving along the z direction, perpendicular to a $D3$ brane embedded in a D-dimensional space-time. In this case, we obtain the form (13) for the Riemann tensor, with

$$\mathcal{K} = -\frac{\epsilon^2}{4} g_{Xz}^2. \tag{14}$$

The fact that the space is not in general maximally symmetric has been confirmed by a detailed analysis of the general case, whose details we do not discuss here.

In the special $D1/D3$ case, the resulting space-time away from the hypersurface $X = 0$ looks, for $t, y_i > 0$, like $AdS_3 \otimes E_{D-3}$, where E_{D-3} is a flat Euclidean $(D-3)$-dimensional space. String theory on AdS_3 has been the subject of many recent articles in the context of critical string or D–brane theory [9]. Much of the mathematical interest in string theories formulated on such spaces lies in the fact that such spaces have a correspondence with well-known group structures based on $SL(2, R)$: in particular, the boundary of the AdS_3 may be formulated as an appropriate coset of $SL(2, R)$. This is relevant to the study of many physical problems, such as the quantum Hall effect and high-energy scattering in QCD [17], as well as the two-dimensional black hole [10]. It is interesting and suggestive that a patch of the space-time found here resembles that structure. We also note that such a target space-time corresponds to the electrovac solution of gauged supergravity [8], which can be shown to leave space-time supersymmetry unbroken.

We conclude with some remarks about the metric (13), (14).

(i) If the folding is ignored: $g_{Xz} = 0$, the space-time resembles flat Minkowski space-time to $\mathcal{O}(\epsilon^2)$ [3], upon making the following transformation for $t > 0$:

$$\tilde{z} = z + \frac{1}{2}\epsilon u_z t^2, \qquad \tilde{t} = t. \tag{15}$$

This implies that the space-time induced by the recoiling D brane resembles, for $t \gg 0$ and to order $\mathcal{O}(\epsilon^2)$, a Rindler wedge space with 'acceleration' ϵu_z. This space is well known to produce a conical singularity when the time is compactified to a Euclidean-signature 'temperature', with deficit angle

$$\delta_{0z} \sim 2\pi \left(1 - 1/\epsilon u_z\right) . \tag{16}$$

Such conical singularities break target space-time supersymmetry. However, this particular conical singularity should be considered as a 'thermal' property of the space. Critical string theory in Rindler space is poorly understood at present, and the only case studied so far concerns the situation in which the deficit angle is quantized: $\delta_{0z} = 2\pi(1 - \frac{1}{N})$, where N a positive integer, in which case the situation is equivalent to a string moving on a Z_N orbifold [18] [3]. This is obviously not the case in (16), given that our recoil/folding analysis

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[3]Even in this case, however, the perturbative superstring vacuum is known to be plagued with tachyonic instabilities, indicating that a non-perturbative analysis is in order.

is valid for weakly-coupled string theory with $g_s \ll 1$, and hence very massive D branes that move slowly: $\epsilon|u_i| \ll 1$. The non-critical logarithmic conformal field theory string approach we advocate here may thus be of use in the generic study of strings in Rindler spaces.

(ii) A simpler (but also interesting) case in which a conical singularity is generated is that in which the cone is produced by the $D3$ brane folding alone, i.e., ignoring the recoil velocity $u_i \to 0^+$. We look at the asymptotic structure of the distorted space-time obtained by such a folding of the D brane along, say, the Y axis, for $X \gg 0$ and $t \gg 0$, suppressing the recoil contribution for simplicity. Upon performing the time transformation $t \to t - \frac{1}{2}\epsilon g_{Xz}Xz$, the line element of the above-mentioned asymptotic space-time becomes:

$$ds^2 = -dt^2 + dy_\perp^2 + \left(1 - \alpha^2\, z^2\right)\, dX^2 + \left(1 + \alpha^2\, X^2\right)\, dz^2 - 2\alpha\, z\, dX\, dt\ ,$$
$$\alpha \equiv \frac{1}{2}\epsilon g_{Xz}\ , \tag{17}$$

where $y_\perp$ denotes collectively the remaining $D - 3$ space-time coordinates. First of all, it is immediate to observe that there is no angular deficit for non-compact bulk dimension z. This is in agreement with the maximally-symmetric AdS$_3$ structure, and is in agreement with the above-mentioned fact that space-time supersymmetry is preserved by such space-times.

The situation changes drastically in case where the extra dimension z is *compact*, which is the case assumed in [5–7]. For simplicity, we assume that z lies on a circle S^1 of unit radius (in appropriate string units, in which all the other distances are measured). For compact z and at fixed $X \sim 1/\epsilon \gg 0$ and $t \gg 0$, such that $\alpha^2\, X^2 \simeq g_{Xz}^2/4$, we observe from the metric (17) that there exists a deficit angle in the circle around z:

$$\delta \simeq (\pi g_{Xz}^2/4) \tag{18}$$

implying the dynamical formation of a conical-like singularity due to the $D3$–brane folding in this asymptotic region of the bulk space-time.

Such singularities in general break bulk space-time supersymmetry [19]. However, in view of the fact that the folded D brane is an excited state of the string/D–brane system, the phenomenon should be viewed as a symmetry obstruction rather than a spontaneous breaking of symmetry, in the sense that, although the ground state of the string/D–brane system is supersymmetric, recoil produces a particular excited state that does not respect that symmetry [20].

(iii) As mentioned earlier, logarithmic field theory is the appropriate tool for discussing transitions between different classical string vacua described by different conformal field theories. It has been shown previously, in the context of closed-string/D-particle scattering, how one logarithmic pair of operators describes the transition between flat space and an AdS space [3, 4], and how another logarithmic pair describes transitions between different two-dimensional string black holes [11]. The logarithmic pair discussed here may open the way to a description of the transition between flat space and a two-dimensional black hole, but pursuing that possibility lies beyond the scope of this paper.

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