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Twisting S-branes

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Abstract: Smooth time dependent supergravity solutions corresponding to analytic continuations of Kerr black holes are constructed and limits with a local de Sitter phase are found. These solutions are non-singular due to a helical twist in space and a fine tuning of the energy flow in the spacetime. For the extremal limit in which the mass and twist parameters are equal the S-brane undergoes de Sitter expansion. Subextremal limits show the formation and decay of a twisted circle and closed string tachyon condensation backreaction effects can be followed. For small values of the twist deformation, a short lived ergosphere envelopes the S-brane and leads to the production of closed timelike curves.

Keywords: Superstrings and Heterotic Strings, p-branes, Tachyon Condensation.

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Contents

1. Introduction

Time dependent tachyon condensation [1]–[4] has received considerable attention recently due to exciting experimental results regarding the accelerating nature of our universe. Attempts to find a source for de Sitter space in the same way D-branes have near horizon anti de Sitter space geometries, have included the introduction of Spacelike branes [1] which describe the formation and decay of an unstable brane. While the formation and decay of an unstable brane might be expected to be a smooth process, the supergravity S-brane solutions of refs. [1, 3] contain singularities. It was not clear if these singularities were a deficiency of the supergravity approximation or arose from other factors, see ref. [4] for further details and discussions on related issues.

Gutperle and Strominger proposed several resolutions of the singularity problem including deforming the S-brane R-symmetry to obtain smooth spacetimes. The R-symmetry

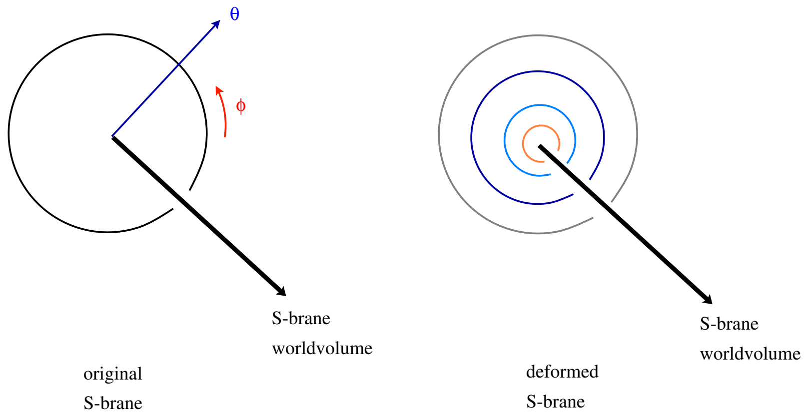

Figure 1: Due to the original R-symmetry, the directions transverse to S-brane worldvolume uniformly evolve in time. By deforming the R-symmetry it is possible to tune the energy flow in the transverse directions.

corresponds to the symmetry transverse to the S-brane worldvolume so reducing the R-symmetry could localize the S-brane in space as well as time. These deformed solutions would represent finely tuned incoming energy forming a localized and unstable object which then decays into outgoing radiation. Non-singular S-brane configurations have also recently been discussed in ref. [5].

Inspired by the discussion of how to resolve S-brane singularities in ref. [1] in this paper we begin an analysis of a class of solutions with reduced R-symmetry. These solutions are analytic continuations of rotating Kerr black holes and are free from curvature singularities. To distinguish these new solutions we note that the original S0-brane in four dimensions was a spacetime of the form

$$R^{1,1} \times H^2 \tag{1.1}$$

and so had $\mathrm{SO}(1,2)$ R-symmetry. The hyperbolic space H^2 was transverse to the one dimensional worldvolume of the S0-brane. For these new twisted S-branes the spacetime manifold $\mathcal{M}$ is a fiber bundle over $R^{1,1}$ with fiber H^2 and so the hyperbolic space is non-trivially fibered over the S-brane worldvolume

$$R^{1,1} \ltimes H^2 \tag{1.2}$$

which reduces the R-symmetry to $\mathrm{SO}(2)$. At a fixed radial coordinate r the Kerr solution depends on θ, while at a fixed moment in time τ the S-brane has non-trivial dependence in the radial θ direction as shown in figure 1. Energy flows towards the S-brane worldvolume, but the singularity is avoided because the regions farther from the origin move more slowly.

Although the time and space coordinates of the metric are effectively interchanged under Wick rotation, the spacetime structure is not described by a ninety degree rotation of the usual Kerr black hole Penrose diagram which is conventionally drawn for the equatorial plane $\theta = \pi/2$. This S-brane is smooth and free from curvature singularities due to the hyperbolic nature of the spacetime. For the Kerr solution θ is periodic so there is a special value at which to put the ring singularity. Under analytic continuation, there is no equivalent location to put the ring singularity and it therefore vanishes.

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These solutions are classified by a twist parameter, a, which characterizes the R-symmetry breaking, and what we shall call a mass parameter, M. The extremal $a = M$, subextremal $a > M$ and superextremal $a < M$ cases have distinct features and a general discussion of their properties is given in section 2. Although the Kerr solution has closed timelike curves near the ring singularity, these twisted S-branes are free from closed timelike curves for large twists $a \geq M$. In fact for $a > M$ there are no horizons, no ergospheres and no curvature singularities. Section 3 contains a discussion of the extremal limit which is distinguished by having a localized two dimensional de Sitter space limit, $dS_2 \ltimes H^2$. In section 4 the subextremal case is shown to be useful in the study of closed string tachyon condensation backreaction on spacetime. The superextermal case is discussed in section 5 and is found to produce closed timelike curves and an ergosphere.

2. S-branes with reduced R-symmetry

2.1 Review of Sp-brane with SO$(1, D - p - 2)$ R-symmetry

In four dimensions the S0-brane metric is

$$ds^2_{S0} = \left(1 - \frac{2M}{\tau} - \frac{e^2}{\tau^2}\right) dz^2 - \frac{d\tau^2}{1 - \frac{2M}{\tau} - \frac{e^2}{\tau^2}} + \tau^2(d\theta^2 + \sinh^2\theta d\phi^2) \qquad (2.1)$$

where the S-brane worldvolume lies along the z direction and θ, ϕ are transverse directions. Time symmetric S-branes are obtained by setting $M = 0$. This solution is the analytic continuation of Reissner Nordstrom black holes

$$\tau \to ir\,, \qquad z \to it\,, \qquad \theta \to i\theta\,, \qquad M \to iM\,. \qquad (2.2)$$

The Penrose diagrams for these solutions are ninety degree rotations of the Reissner Nordstrom Penrose diagrams, with the difference being that every point on the Penrose diagram represents a two dimensional hyperbola instead of a two sphere; the original SO(3) rotational symmetry becomes the SO(1,2) R-symmetry after the analytic continuation. In general starting from a p-brane in D dimensions, the SO$(D - p - 1)$ rotational symmetry transverse to the brane worldvolume becomes the SO$(1, D - p - 2)$ Sp-brane R-symmetry.

In addition to the $M > e$ charged black holes there are the $M < e$ solutions with a naked timelike singularity. The ninety degree rotation of the naked singularity Penrose diagram has a cosmological nature and is equivalent to the FRW Penrose diagram. Cosmological implications of such S-brane solutions were discussed in refs. [6] with the two cases, $M > e$ and $M < e$, corresponding to two universes due to the curvature singularity. Depending on the mass to charge ratio, there is either a big bang/crunch or a timelike singularity.

2.2 Twisting S-branes

S-branes originally were homogenous and purely time dependent constructions. Their supergravity solutions were analytic continuations of supergravity solutions of p-branes which were isotropic in the transverse directions. Deforming S-branes and localizing them in space is interesting and would require a solution dependent on two variables, time and

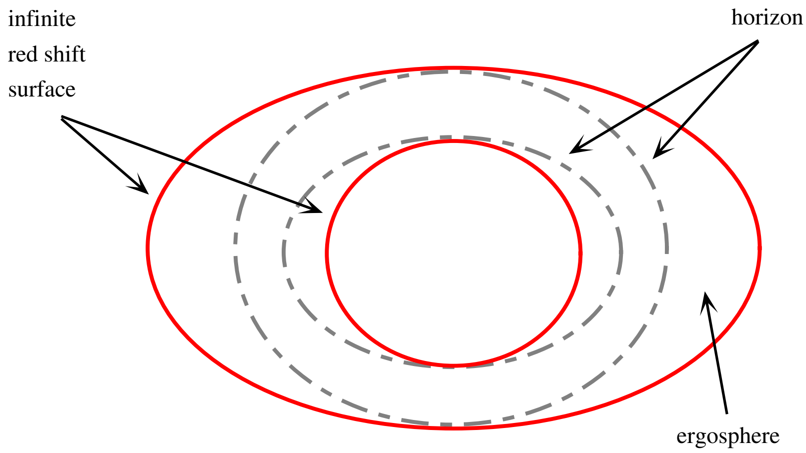

Figure 2: For $a < M$ the Kerr black hole has two horizons and an ergosphere.

a radial distance from the S-brane worldvolume. Solutions which depend on two variables are unfortunately generally difficult to obtain. Given the fact though that the S0-brane is the analytic continuation of a black hole solution it is possible to ask if there are any known non-isotropic black holes. The most general solution with mass, charge and angular momentum in four dimensions is the rotating charged black hole which is the focus of this paper.

The electrically charged Kerr-Newman metric in Boyer-Lindquist coordinates is

$$
\begin{aligned}
ds^2_{Kerr} &= -\left(\frac{\Delta - a^2\sin^2\theta}{\rho^2}\right) dt^2 - \frac{2a\sin^2\theta(r^2+a^2-\Delta)}{\rho^2} dt\, d\phi + \\
&\quad + \left[\frac{(r^2+a^2)^2 - \Delta a^2\sin^2\theta}{\rho^2}\right]\sin^2\theta d\phi^2 + \frac{\rho^2}{\Delta}dr^2 + \rho^2 d\theta^2 \qquad (2.3)\\
A &= -\frac{er}{\rho^2}(dt - a\sin^2\theta\, d\phi) \qquad (2.4)\\
\rho^2 &= r^2 + a^2\cos^2\theta \qquad \Delta = r^2 + a^2 + e^2 - 2Mr\,. \qquad (2.5)
\end{aligned}
$$

Here M, e, a are the mass, charge and angular momentum parameters while A is the vector potential. The interesting features of this solution persist in the absence of charges so we will set the electric (and magnetic) charges to zero. One may consider the angular momentum parameter to be a different charge carried by the black hole. The neutral Kerr metric is

$$
ds^2_{Kerr} = -\frac{\Delta}{\rho^2}(dt - a\sin^2\theta d\phi)^2 + \frac{\sin^2\theta}{\rho^2}[(r^2+a^2)d\phi - adt]^2 + \frac{\rho^2}{\Delta}dr^2 + \rho^2 d\theta^2 \qquad (2.6)
$$

which has been written in a form convenient for checking various properties. When $a < M$, the Kerr black hole has two horizons and a ring singularity. For the extremal case $a = M$, the horizons become degenerate and when $a > M$ the singularity is naked.

To obtain the twisted S-brane metric we analytic continue three coordinates along with the angular momentum and mass parameters

$$
t \mapsto iz\,, \qquad \theta \mapsto i\theta\,, \qquad r \mapsto i\tau\,; \qquad a \to ia\,, \qquad M \to iM \qquad (2.7)
$$

where our convention is not to introduce minus signs in the analytic continuation. The Wick rotation has the following effect

$$\rho^2 \;\to\; -\tau^2 - a^2\cosh^2\theta = -\rho^2\,, \qquad \Delta \;\to\; -\tau^2 + 2M\tau - a^2 = -\Delta\,. \tag{2.8}$$

The twisted S0-brane metric we will focus on is

$$\begin{aligned} ds^2_{\text{twisted S0}} &= \frac{\Delta + a^2\sinh^2\theta}{\rho^2}dz^2 + \frac{2a\sinh^2\theta(\tau^2+a^2-\Delta)}{\rho^2}dzd\phi + \\ &\quad + \left[\frac{(\tau^2+a^2)^2 + \Delta a^2\sinh^2\theta}{\rho^2}\right]\sinh^2\theta d\phi^2 + \rho^2\left(-\frac{d\tau^2}{\Delta} + d\theta^2\right) \qquad (2.9)\\ &= -\frac{\rho^2}{\Delta}d\tau^2 + \rho^2 d\theta^2 + \frac{\Delta}{\rho^2}(dz - a\sinh^2\theta d\phi)^2 + \frac{\sinh^2\theta}{\rho^2}[(\tau^2+a^2)d\phi + adz]^2 \\ \Delta &= \tau^2 + a^2 - 2M\tau \qquad (2.10) \end{aligned}$$

where we grouped the terms to hi-light the dependence on the factor Δ.

2.3 Absence of curvature singularities and closed timelike curves

Several properties of this spacetime may be inferred from our knowledge of Kerr black holes. When the parameter $a = 0$ this gives the original S0-brane of eq. 2.1, so this is a new twisted deformation of known S-branes. For large values of τ, τ acts as the time coordinate, although for various parameter ranges there are horizons where Δ changes sign and so τ is not necessarily always timelike. The θ dependence means that the metric can represent a spatially localized source and due to the metric cross term $g_{z\phi}$ there is a coupling of two directions z and ϕ.

For large values of angular momentum

$$a^2 > m^2 \;\to\; \Delta > 0 \tag{2.11}$$

the Kerr solution does not have a horizon, where $\Delta = 0$. It is easy to check that after Wick rotation eq. 2.11 still holds, so the twisted S-brane does not have horizons for large values of the parameter a. The Kerr solution is also known to have a ring curvature singularity at $\rho^2 = 0$

$$\rho^2 = 0 \;\to\; r = 0\,, \qquad \theta = \frac{\pi}{2}\,. \tag{2.12}$$

The curvature for the S-brane is similarly dependent on

$$\rho^2 = \tau^2 + a^2\cosh^2\theta\,. \tag{2.13}$$

The difference however is that while τ can be zero, there are no zeros for hyperbolic cosine so $\rho \neq 0$ for real values of θ. Therefore the hyperbolic nature of the transverse space guarantees that this S-brane has no curvature singularities. Intuitively the reason that the singularity disappears is due to the fact that there is a special and unique value of θ at which to put the ring singularity in the case of the rotating black hole, but there is no equivalent value of θ in hyperbolic space to put a singularity. The conformal structure of this new spacetime is very different from the original S-branes and Kerr spacetimes. As shown by this example, however, reducing the R-symmetry is one way to cure the S-brane singularity problem.

Although for real values of θ the spacetime is smooth, it might be interesting to examine the fact that there is a singularity for $\theta = -i\pi/2$. Analogous observations have been made about Sen's rolling tachyon solution which can be thought of as a periodic array of D-branes in imaginary time [7]. For this S-brane however the singularities lie on a circle at an imaginary θ value. It would be interesting to clarify this point and understand if the singularity in complex directions corresponds to a specific tachyonic mode instability.

A known problem with the Kerr solution is the existence of closed timelike curves. The problem is that the periodic angular coordinate ϕ changes from being timelike to spacelike since the coefficient

$$g_{\phi\phi} = \left(\frac{(r^2+a^2)^2 - a^2\Delta\sin^2\theta}{r^2+a^2\cos^2\theta}\right)\sin^2\theta \tag{2.14}$$

changes sign from positive to negative values for

$$(r^2+a^2)^2 < a^2\sin^2\theta(r^2+a^2-2Mr)\,. \tag{2.15}$$

Near the ring singularity at $\theta = \pi/2$, and for r small and negative, the above condition is satisfied and we have closed timelike curves.

It is then imperative to check if this twisted S-brane also has closed timelike curves. One difference is that the time direction is now along τ and

$$g_{\tau\tau} = -\frac{\rho^2}{\Delta} \tag{2.16}$$

which is governed by the sign of Δ, which we have said is always positive for $a > M$. When the twist parameter a is large relative to M, then z and ϕ coordinates are always spacelike coordinates. It is also possible to explicitly check that the condition $g_{\phi\phi} < 0$ implies

$$(\tau^2+a^2)^2 < -a^2\sinh^2\theta(\tau^2+a^2-2M\tau) \tag{2.17}$$

which is satisfied if and only if

$$\tau^2+a^2-2M\tau < 0\,. \tag{2.18}$$

For $a \geq M$ there are no closed timelike curves and the spacetime does not rotate. When the angular momentum parameter is small, $a < M$, the S-brane does have an ergosphere and closed timelike curves as discussed in section 5.

Adding charge does not significantly change the properties of this S-brane. For electrically charged S-branes we find that the solution is still smooth and free from closed timelike curves as long as $a^2 - e^2 \geq M^2$.

2.4 General properties

The Kerr solution has two null vectors

$$l^\mu = \frac{r^2+a^2}{\Delta}\left(\frac{\partial}{\partial t}\right)^\mu + \frac{a}{\Delta}\left(\frac{\partial}{\partial\phi}\right)^\mu + \left(\frac{\partial}{\partial r}\right)^\mu \tag{2.19}$$

$$n^\mu = \frac{r^2+a^2}{2\rho^2}\left(\frac{\partial}{\partial t}\right)^\mu + \frac{a}{2\rho^2}\left(\frac{\partial}{\partial\phi}\right)^\mu - \frac{\Delta}{2\rho^2}\left(\frac{\partial}{\partial r}\right)^\mu \tag{2.20}$$

and likewise the twisted S-brane has two null vectors

$$\tilde{l}^{\mu} = \frac{\tau^2+a^2}{\Delta}\left(\frac{\partial}{\partial z}\right)^{\mu} + \frac{a}{\Delta}\left(\frac{\partial}{\partial\phi}\right)^{\mu} + \left(\frac{\partial}{\partial\tau}\right)^{\mu} \tag{2.21}$$

$$\tilde{n}^{\mu} = -\frac{\tau^2+a^2}{2\rho^2}\left(\frac{\partial}{\partial z}\right)^{\mu} - \frac{a}{2\rho^2}\left(\frac{\partial}{\partial\phi}\right)^{\mu} + \frac{\Delta}{2\rho^2}\left(\frac{\partial}{\partial\tau}\right)^{\mu} \tag{2.22}$$

$$\tilde{l}^2 = 0 = \tilde{n}^2\,, \qquad \tilde{l}\cdot\tilde{n} = -1\,. \tag{2.23}$$

From these two null vectors it is possible to construct a Killing two tensor which is useful for simplifying calculations and checking properties of this S-brane. The S-brane metric is manifestly independent of z and ϕ so there are also two Killing vectors, $\partial/\partial z$ and $\partial/\partial\phi$. An interesting point is that near $\Delta = 0$, the null vectors are also Killing vectors.

Although the twisted Sbrane metric is independent of the coordinate z, there is a metric cross term $g_{z\phi}$. Concretely the metric is invariant only under the combined discrete symmetry

$$z \to -z\,, \qquad \phi \to -\phi \tag{2.24}$$

demonstrating that there is a twist in space, which is a spiral or helix structure along these two directions. The cross term in the metric in these two spatial directions is reminiscent of twisted circles discussed in ref. [8]. Interestingly it is shown in section 4 that while supergravity S-branes were originally conceived to have implications for open string tachyon condensation, this twisted S-brane has implications for closed string tachyon condensation [9] in particular the time dependent formation and decay of a twisted circle.

The horizons of the Kerr black hole are defined to be where $g_{rr} = \infty$, so $\Delta = 0$, and they occur at

$$r_{\text{horizon}} = M \pm \sqrt{M^2 - a^2}\,. \tag{2.25}$$

There are also infinite red shift surfaces where $g_{tt} = 0$ at

$$r_{\text{ergosphere}} = M \pm \sqrt{M^2 - a^2\cos^2\theta} \tag{2.26}$$

which are the boundaries of the ergosphere and which begin farther out from the origin than the horizon

$$r_{\text{ergo}+} > r_{\text{horizon}+}\,. \tag{2.27}$$

For the S-branes the situation is similar but when $a > M$, $\Delta \neq 0$, the metric is smooth and there are no metric degeneracies. If $a \leq M$ then $g_{\tau\tau}$ does not go to zero but can go to infinity. Instead of an infinite red shifting surface there is a infinite blue shifting surface. There is also a surface on which $g_{zz} = 0$ which we regard as the boundary of an ergosphere. These points are further discussed in section 5.

2.5 Far from origin and late time behavior

It is interesting to examine properties of the twisted S-brane far from the origin. In the limit $\cosh^2\theta \gg \tau^2/a^2$ the metric is

$$ds^2 = e^{2\theta}\left[-\frac{a^2 d\tau^2}{\tau^2+a^2-2M\tau} + (\tau^2+a^2-2M\tau)d\phi^2 + a^2 d\theta^2\right] + dz^2 + \frac{\tau^2+a^2-\Delta}{a}dz\,d\phi\,. \tag{2.28}$$

At large values of θ, the effect of the cross term is small and can be neglected. If we also take the limit $\tau \gg M, a$ the metric is

$$ds^2 = e^{2\theta}\left[-\frac{a^2 d\tau^2}{\tau^2} + \tau^2 d\phi^2 + a^2 d\theta^2\right] + dz^2 \tag{2.29}$$

which we further rewrite so that θ acts as a radial coordinate

$$R = ae^{\theta}\,, \qquad T = \ln\frac{\tau}{\tau_0}$$

$$ds^2 = dR^2 - \frac{R^2}{\Delta}d\tau^2 + \frac{\Delta}{a^2}R^2 d\phi^2 + dz^2 \tag{2.30}$$

$$= e^{2\theta}[-dT^2 + \tau_0^2 e^{2T} d\phi^2 + a^2 d\theta^2] + dz^2$$

$$= dR^2 - R^2 dT^2 + \frac{\tau_0^2}{a^2}R^2 e^{2T} d\phi^2 + dz^2\,. \tag{2.31}$$

The metric in this limit is two dimensional Rindler space with exponential expansion in the angular direction which will have the interpretation of being the result of a passing pulse of energy.

Next examine eq. 2.28 for values of τ which minimize Δ. The physical significance of this time is that it will serve as an interesting initial condition for the time evolution of this solution. For the case $a > M$, the metric near $\tau = M$ is a Rindler wedge times a line of flat cones

$$ds^2_{\text{far cone}} = dR^2 - R^2 d\tau^2 + dz^2 + \frac{\Delta}{a^2}R^2 d\phi^2\,. \tag{2.32}$$

Returning to the usual Minkowski coordinates the above limit gives the metric

$$ds^2_{\text{far cone}} = -dt^2 + dx^2 + dz^2 + \frac{\Delta}{a^2}(x^2 - t^2)d\phi^2\,. \tag{2.33}$$

For $a = M$ there is exponential expansion in time for ϕ when $\tau \sim M$ so this is an exponentially expanding cone. For $a < M$ the expansion of the cone is linear in time for times close to $\tau \approx \tau_{+,-} = M \pm \sqrt{M^2 - a^2}$.

Examining the late time limit where $\tau^2/a^2 \gg \cosh^2\theta$ we have

$$ds^2 = -d\tau^2 + \tau^2 dH_2^2 + dz^2 - \frac{2M}{\tau}(a\sinh^2\theta d\phi + dz)^2 \tag{2.34}$$

a Milne wedge with a twist whose effect disappears at late times. For observers in the past and future light cones, the R-symmetry breaking is therefore spontaneous in time.

3. Extremal limit $a = M$

One of the original motivations for introducing Spacelike Branes was to find an object in string theory which could source de Sitter space. In fact there are E-branes [10] which do have a de Sitter space near horizon limit. These are obtained by starting from a BPS brane, and analytically continuing both the electric charge and the mass. Although this does lead to a de Sitter space limit, analytically continuing the charge leads to type-II* string theories with a wrong sign kinetic energy term for the field strength. Known charged

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S-branes in type-II string theories do not have such an extremal limit. To find an extremal limit related to de Sitter space and also stay in type-II string theories, one apparently needs a new type of charge.

The twisted S-branes in this paper are solutions to the vacuum Einstein equations but can be interpreted as having a charge arising from twisting and not from a field strength and so they should exist in the usual string theories. It is known that Kerr black holes do have an interesting extremal limit, $a = M$. We begin by examining this extremal S-brane limit and find that the S-brane worldvolume undergoes an expanding de Sitter phase.

3.1 Review of extremal Reissner Nordstrom horizon

The extremal Reissner Nordstrom black hole metric

$$ds^2_{RN} = -\left(1 - \frac{2M}{r} + \frac{M^2}{r^2}\right) dt^2 + \frac{dr^2}{1 - \frac{2M}{r} + \frac{M^2}{r^2}} + r^2 dS_2^2 \tag{3.1}$$

has a near horizon limit

$$r = M + \delta \tag{3.2}$$

$$ds^2_{RN} = -\frac{\delta^2}{M^2} dt^2 + \frac{M^2 d\delta^2}{\delta^2} + M^2 dS_2^2 \tag{3.3}$$

which is $AdS_2 \times S_2$. Starting from this solution it is possible to obtain a near horizon limit which is de Sitter space instead of anti de Sitter space, if the mass and the charge are Wick rotated $e = M \to iM$. These are solutions however in type-II* theories with wrong sign kinetic energy terms since the field strength contribution to the action $F^2 \sim e^2$ becomes $-F^2$.

3.2 Near horizon limit

We begin by examining the terms in the S-brane metric eq. 2.9 along the z, ϕ directions

$$\frac{\Delta}{\rho^2} dz^2 + \frac{2\Delta a \sinh^2\theta}{\rho^2} dz d\phi + \frac{a^2 \Delta \sinh^4\theta}{\rho^2} d\phi^2 \tag{3.4}$$

and

$$\frac{a^2 \sinh^2\theta}{\rho^2} dz^2 + \frac{2a \sinh^2\theta (\tau^2 + a^2)}{\rho^2} dz d\phi + \frac{(\tau^2 + a^2)^2 \sinh^2\theta}{\rho^2} d\phi^2 \,. \tag{3.5}$$

Given that this S-brane is localized in space as well as time, we look at times which minimize Δ and also approach the S-brane worldvolume so θ is small

$$\tau = M + \delta \,, \qquad 0 \leq \Delta = (\tau - M)^2 = \delta^2 \ll |\tau^2 + a^2| \,, \qquad |\Delta a^2 \sinh^2\theta| \ll (\tau^2 + a^2)^2 \,. \tag{3.6}$$

In this limit the second and third terms in eq. 3.4 can be neglected and the metric is

$$ds^2_{\text{close extremal}} = -\frac{\rho^2}{\Delta} d\tau^2 + \rho^2 d\theta^2 + \frac{\Delta}{\rho^2} dz^2 + 2a^2 \sinh\theta^2 \left(d\phi + \frac{1}{2a} dz\right)^2 \tag{3.7}$$

$$= -\frac{2M^2 d\delta^2}{\delta^2} + \frac{\delta^2}{2M^2} dz^2 + 2M^2 \left[d\theta^2 + \sinh^2\theta \left(d\phi + \frac{dz}{2M}\right)^2\right] \tag{3.8}$$

a twisted hyperbolic space fibered non-trivially over two dimensional de Sitter space

$$dS_2 \ltimes H^2 \,. \tag{3.9}$$

Writing the metric in its exponential form

$$T = \sqrt{2}M \ln \frac{\delta}{\delta_0}\,, \qquad \delta = \delta_0 e^{\frac{T}{\sqrt{2}M}} \tag{3.10}$$

$$ds^2_{\text{close extremal}} = -dT^2 + \frac{\delta_0^2}{2M^2} e^{\frac{T}{\sqrt{2}M}} dz^2 + 2M^2 \left[d\theta^2 + \sinh^2\theta \left(d\phi - \frac{1}{2M} dz \right)^2 \right] \tag{3.11}$$

one finds that the Hubble parameter is $H^{-1} = \sqrt{2}M$. After a time $\tau \sim 2M$ the near S-brane limit of eq. 3.6 is no longer valid; due to the instability of this configuration, it eventually disperses. This causes the exponential expansion phase to slow down and the spacetime approaches flat space. If we add electric or magnetic charge to this S-brane, it is still possible to find an extremal limit similar to eq. 3.11, but the relative metric coefficients will be changed.

To remove the twist near the origin one may use the twisted coordinate with non-standard periodicity

$$\tilde{\phi} \to \phi - \frac{z}{2M} \tag{3.12}$$

to obtain a local region of $dS_2 \times H^2$. If we further periodically identify the z direction, this solution apparently demonstrates how flat spacetime decays to a Melvin universe in the far future. It would be interesting to embed this solution in string theory and see what role it plays in the proposed dualities and relations [11] between type 0 and type-II string theories.

While Wick rotating the electric charge leads to solutions in *-theories, Wick rotating the angular momentum parameter leads to a twist parameter. Instead of having S-branes charged under electro-magnetic fields, we deform the spacetime with a twist to obtain a region of de Sitter space near the S-brane worldvolume.

4. Subextremal case $a > M$

For the parameter range $a > M$ these subextremal S-branes have neither horizons nor ergospheres, and the near S-brane limit qualitatively differs from the extremal case. The time evolution of this solution is interpreted as showing the gravitational backreaction of local closed string tachyon condensation. In particular the solution shows the decay of a cone whose tip is a twisted circle. As the twist unwinds, this causes the cone to exponentially expand and open up into flat space.

4.1 Review of twisted circles

A twisted circle [8] is a non-standard identification of R^3

$$ds^2 = dr^2 + r^2 d\phi^2 + dz^2 \tag{4.1}$$

$$\phi \sim \phi + 2\pi n_1 + 2\pi n_2 RB\,, \qquad z \sim z + 2\pi n_2 R \tag{4.2}$$

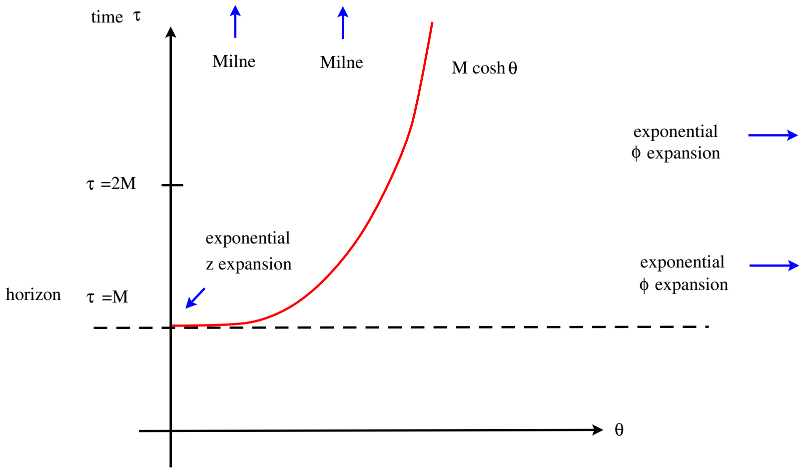

Figure 3: The neighborhood of the S-brane worldvolume undergoes a de Sitter expansion phase for times $\tau \sim M$. The expansion slows down into a Milne spacetime after $\tau \sim 2M$. Far from the origin, we have regions of exponential expansion of the ϕ direction. The curve $\tau = M\cosh\theta$ marks the boundary between the two types of behaviors.

where the unusual feature is that the z direction is coupled with the ϕ direction. To recover a more standard set of coordinates one can introduce a twisted coordinate

$$\tilde{\phi} = \phi - Bz \tag{4.3}$$

with standard periodicity $\tilde{\phi} \sim \tilde{\phi} + 2\pi n_1$. The metric then becomes

$$ds^2 = dr^2 + r^2(d\tilde{\phi} + Bdz)^2 + dz^2 \tag{4.4}$$

and under dimensional reduction of the z direction, the twist, which is characterized by the constant B, becomes a magnetic field.

4.2 Near extremal limit

Although when $a > M$ there is no horizon, the metric simplifies for $\tau = M$ since this minimizes $\Delta(\tau = M) = a^2 - M^2$. If we also take the near extremal limit, $a \approx M$, and near origin limits $\theta \ll 1$

$$\tau = M + \delta\,, \qquad \Delta = a^2 - M^2 + \delta^2\,, \qquad \rho^2 \sim 2a^2 \tag{4.5}$$

$$\delta^2 \ll a^2 - M^2 = \epsilon\,, \qquad \delta \ll a\,,\ M\,, \qquad \epsilon \ll M\delta \tag{4.6}$$

$$a\theta \ll \delta \ll \sqrt{\epsilon} \ll \sqrt{\delta M} \ll M, a \tag{4.7}$$

the metric is

$$ds^2_{NE-} = -\frac{\rho^2}{\epsilon}d\delta^2 + \rho^2 d\theta^2 + \frac{\epsilon}{\rho^2}dz^2 + \frac{\sinh^2\theta}{\rho^2}[\rho^2 d\phi + adz]^2\,. \tag{4.8}$$

Rescaling the time and z coordinates and taking the limit $\sinh\theta \sim \theta$ the metric becomes

$$z' = \frac{\sqrt{\epsilon}}{\rho}z = \sqrt{\frac{\epsilon}{2}}\frac{z}{M}\,, \qquad \delta = \sqrt{\frac{\epsilon}{2}}\frac{T}{M} \tag{4.9}$$

$$ds^2_{NE-} = -dT^2 + dR^2 + R^2\left[d\phi - \frac{1}{\sqrt{2\epsilon}}dz'\right]^2 + dz'^2\,. \tag{4.10}$$

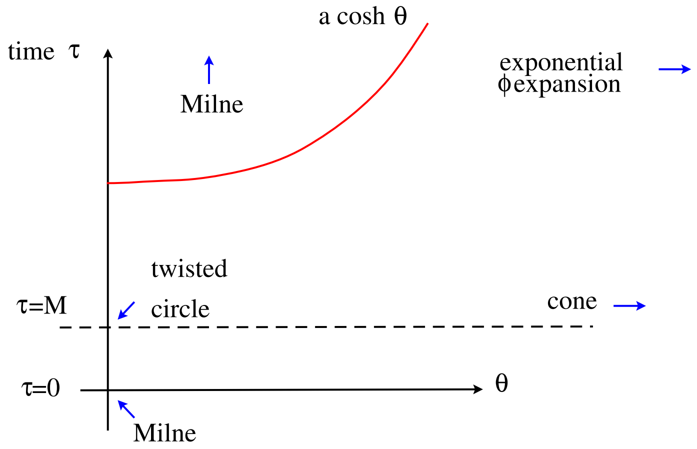

Figure 4: For the subextremal case $a > M$, near the origin we have a twisted circle at $\tau = M$ and which disperses at time $\tau = 2M$ and flattens out to Milne after $\tau = a$. Far from the origin, there are regions of Rindler near time $\tau = M$. After a time a pulse from the origin causes an exponential expansion in the angular directions and after the pulse passes the space becomes flat.

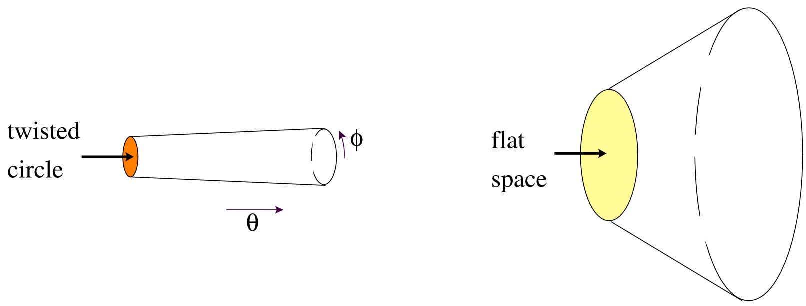

Figure 5: For the near extremal solution, the tip of the cone is a twisted circle at time $\tau = M$. Going forward in time the circle untwists and starting from the tip the cone opens up.

In this limit if we periodically identify the z direction, we obtain a twisted circle. As time passes, there will be an exponential expansion phase, and eventually the twist unwinds leaving a growing region of flat space as shown in figures 4, 5. Recalling that at $\tau = M$, the spacetime far from the origin is a cone (times a line see eq. 2.32), then we see that this solution represents the time evolution of a cone whose tip is replaced by a twisted circle. The importance of $\tau = M$ is that it serves as an interesting initial condition for time evolution for the system. Furthermore this solution is not symmetric in time so there are actually two ways in which the solution unwinds.

4.3 Compactifying the z direction

Although it does not seem necessary to periodically identify the z direction, it is natural so the metric has a twisted circle limit. In addition if we wish to have a sensible Kaluza-Klein compactification near $\tau = M$, it is necessary that the magnetic field be small relative to the compactification scale, $B \ll 1/R$. From eq. 4.10 we see that the magnetic field is small $B = (2\epsilon)^{-1/2}$ and so the compactification radius for the z direction must be smaller than $\epsilon^{1/2}$.

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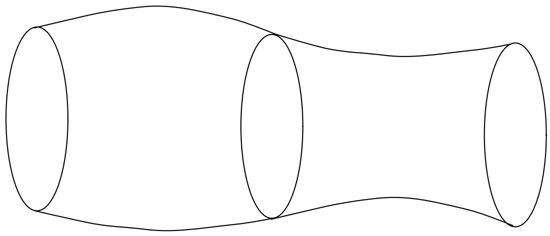

Figure 6: Size of the compactified z direction at fixed θ.

Under compactification this solution has the interpretation as the formation and decay of a fluxbrane. Fluxbranes have been shown to arise from a special scaling limit of a brane and anti-brane pair [12], so the parameter a might be interpreted as how much charge flows along the S-brane worldvolume; the D-brane and anti D-brane annihilate via the twist in the S-brane.

Under periodic identification of z we do not introduce closed timelike curves since we are in the limit $a > M$ where Δ is always positive. Compactifying in the extremal case $a = M$ we would obtain closed null curves and it would be interesting to discuss this in relation to work in ref. [13]. In addition, under compactification the extremal case would lead to a metric degeneracy at $\tau = M$ and $\theta = 0$. To observers bound to the S-brane worldvolume this would appear as a big bang type singularity which is smoothed out by opening a new spatial dimension.

The metric for the circle direction

$$g_{zz} = 1 - \frac{2M\tau}{\rho^2} \tag{4.11}$$

starts with value one at past infinity, increases, then decreases until we get to $\tau = 0$ where it is one again. Then the circle size is less than one, decreasing and going back to one at future infinity. At $\tau = M$ the circle reaches it minimum size which can go to zero in the extremal limit. As we go farther from origin, the time dependence is less pronounced and delayed relative to what occurs near the origin. See figure 6.

Under the usual Kaluza-Klein reduction this spacetime has the interpretation of a the time dependent formation and decay of a fluxbrane. If we use the twisted coordinates as in eq. 4.3 to remove the twist at $\tau = M$, an interesting interpretation of this is that we are seeing how flat space decays into a twisted circle at late times. As in the extremal case it will be interesting to embed this solution in string theory.

5. Superextremal case $a < M$

In this section we examine $a < M$ superextremal S-branes which are distinguished by their horizons, ergospheres and creation of closed timelike curves.

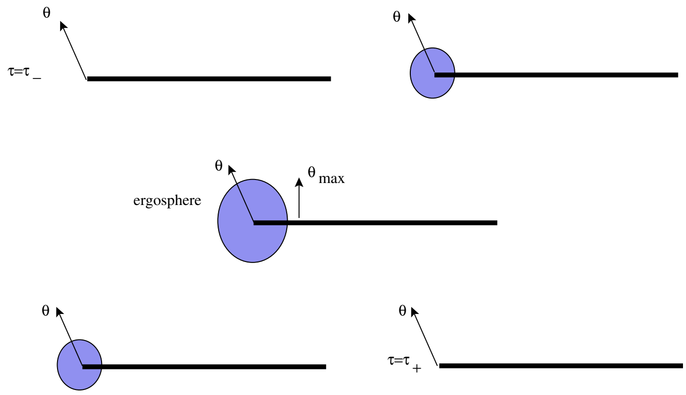

Figure 7: There is an ergosphere between the times τ_-, τ_+. The ergosphere only exists in a neighborhood of the S-brane and extends out to a maximum distance $\theta_{\max}$. The five figures are shown sequentially in time from left to right, top to bottom.

5.1 Ergosphere

For small values of the twist parameter, $a < M$, one might assume that these solutions become qualitatively similar to the original S-brane solutions but in fact a small twist parameter drastically distorts the spacetime structure. These solutions have an ergosphere which exists between two high blue shifting times.

For the Kerr black hole the ergosphere always lies outside the horizon. On the other hand S-branes have surfaces where $g_{\tau\tau} = \infty$

$$\tau_\pm = M \pm \sqrt{M^2 - a^2} \tag{5.1}$$

and the Killing horizons $g_{zz} = 0$ occur when

$$\tau_{\text{ergo}} = M \pm \sqrt{M^2 - a^2 \cosh^2 \theta} \tag{5.2}$$

with

$$\tau_- \leq \tau_{\text{ergo}} \leq \tau_+ \,. \tag{5.3}$$

We will call the region between the Killing horizon and the origin to be the S-brane's ergosphere. The ergosphere extends out to a maximum θ value for $\tau = M$ and then it disappears as illustrated in figure 7. In some sense the ergosphere is hidden.

Due to the change in spacetime signature this solution does rotate for a period of time between τ_- and τ_+.

5.2 Near extremal limit

For $a < M$ there are two non-degenerate roots of Δ and near $\tau = \tau_+ + \delta$ the time component of the metric is

$$-\frac{\rho^2}{\Delta} d\tau^2 = -\frac{\tau_+^2 + a^2}{\tau_{+-}\delta} d\delta^2 \,, \qquad \tau_{+-} = \tau_+ - \tau_- \,. \tag{5.4}$$

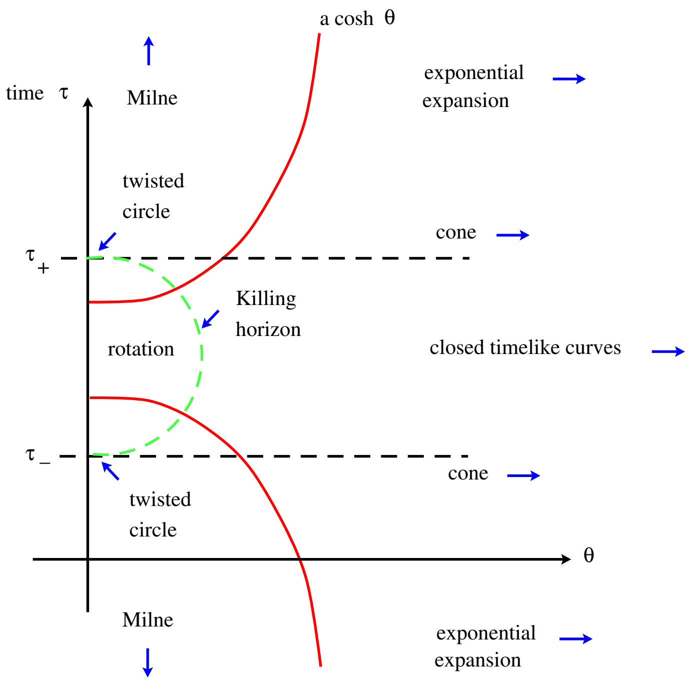

Figure 8: For the superextermal case $a < M$ there are two horizons and an ergosphere. Near the origin we have a twisted circle. Closed timelike curves exist for $\tau_- < \tau < \tau_+$.

Taking the same near S-brane limits as in eq. 3.6 and using a coordinate transformation

$$T = \frac{2\sqrt{2}\rho\delta^{1/2}}{\tau_{+-}^{1/2}}, \qquad \delta = \frac{\tau_{+-}}{4\rho^2}T^2, \qquad z' = \frac{\tau_{+-}z}{2\rho^2} = \frac{\tau_{+-}z}{4M^2} \tag{5.5}$$

$$\begin{aligned} ds^2_{NE+} &= -dT^2 + \rho^2 d\theta^2 + \frac{\tau_{+-}\delta}{\rho^2}dz^2 + \frac{\sinh^2\theta}{\rho^2}[(\tau_+^2 + a^2)d\phi + adz]^2 \\ &= -dT^2 + T^2 dz'^2 + 2M^2\left\{d\theta^2 + \sinh^2\theta\left[d\phi + 4M\frac{dz'}{\tau_{+-}}\right]^2\right\} \end{aligned} \tag{5.6}$$

the metric again describes a twisted circle for small θ which then untwists.

One can further probe the near S-brane limit in the spirit of ref. [14] by using a coordinate choice which introduces a small hyperboloid near the S-brane worldvolume

$$\tau - \tau_+ = T\frac{1+\cosh\tilde{\theta}}{2}, \qquad \sinh^2\theta = T\frac{\cosh\tilde{\theta} - 1}{\tau_+ - M} \tag{5.7}$$

$$d\tau = dT\frac{1+\cosh\tilde{\theta}}{2} + T\frac{\sinh\tilde{\theta}}{2}d\tilde{\theta}, \qquad d\theta = \frac{dT(\cosh\tilde{\theta}-1) + T\sinh\tilde{\theta}d\tilde{\theta}}{2\sqrt{(\tau_+ - M)T(\cosh\tilde{\theta}-1)}} \tag{5.8}$$

$$\Delta = (T - \tau_+)(T - \tau_-) = \tau_{+-}T\frac{1+\cosh\tilde{\theta}}{2}. \tag{5.9}$$

Using the twisted coordinate $\tilde{\phi} = \phi - Bz$, the metric becomes

$$ds^2_{twNE+} = \frac{M^2}{\tau_{+-}T/2}[-dT^2 + T^2(d\tilde{\theta} + \sinh^2\tilde{\theta}d\tilde{\phi}^2)] +$$

$$
\begin{aligned}
&+\frac{T\tau_{+-}}{2\rho^2}\left[1+\cosh\tilde{\theta}+\frac{2}{\tau_{+-}^2/2}(a+B(\tau_+^2+a^2))^2(\cosh\tilde{\theta}-1)\right]dz^2+\\
&+\frac{2T(\cosh\tilde{\theta}-1)}{2M^2\tau_{+-}/2}(\tau_+^2+a^2)(a+B(\tau_+^2+a^2))dzd\tilde{\phi}-\\
&-\frac{T(\tau_+^2+a^2)^2}{4M^2\tau_{+-}/2}(\cosh\tilde{\theta}-1)^2d\tilde{\phi}^2\\
=&\ \frac{M^2}{\tau_{+-}T/2}[-dT^2+T^2dH_2^2]-\frac{T}{\tau_{+-}/2}\left[(\cosh\tilde{\theta}-1)d\tilde{\phi}-\frac{1+2BM}{M}dz\right]^2+\\
&+\frac{T\tau_{+-}}{4M^2}\left[1+\frac{2M^2(1+2BM)}{\tau_{+-}^2/2}\right](1+\cosh\tilde{\theta})dz^2
\end{aligned}
\tag{5.10}
$$

which is similar to a near origin limit of a hyperbolic space version of Taub-NUT (with an extra term). It would be interesting to find the exact solution corresponding to this limit which would be a new type of S-brane. This follows from the fact that the isometry generated by the Killing vector along the twisted direction has an isolated fixed point which is at $\tau = \tau_+$ and $\theta = 0$. This new solution should be interesting in the same way that "nut" solutions are canonical examples of isolated fixed points for euclidean spacetimes. Apparently this new solution is not an analytic continuation of Taub-NUT solutions but an exact solution is expected to be found from a modified ansatz.

It would be interesting to further explore several properties of this solution. Although closed timelike curves do exist they apparently are created only for a finite duration. Given the special properties of the Kerr ergosphere, there could also be an interesting physical significance of these S-brane ergosphere regions.

6. Discussion and further directions

Spacelike branes play an interesting role in time dependent physics. In this paper we found smooth S-branes with reduced R-symmetry and which have a spacetime twist which may interpreted as a type of charge. A feature of these solutions is that they have extremal limits in which the S-brane worldvolume undergoes an expanding de Sitter phase. Although the lifetime and amount of inflation of this phase are infinite, the configuration is unstable and eventually disperses. Near (sub)extremal limits of this solutions are also interesting and were shown to describe the decay of a cone with a twisted circle at its tip. Superextremal solutions produce closed timelike curves.

It is tempting to consider the twisted S-brane to be an example of a closed string analogue of an inhomogeneous [2, 4, 15] open string tachyon solution some of which had interpretations as describing the formation of lower dimensional branes and strings. For the rolling tachyon, the solutions with energy equal to the the tachyon potential correspond to the creation of a brane while solutions with less energy are smooth. For the twisted S-brane when $a > M$ the solution is smooth while the extremal limit, $a = M$, should then correspond to the formation of an object. Reducing the R-symmetry is effectively analogous to introducing inhomogeneities in the rolling tachyon.

Several avenues are available to pursue such as generalizations of these solutions, their embedding into string theory and relationships to cosmology. Further details will appear in a following paper.

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