Kaluza-Klein Black Holes
in String Theory

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Abstract

Exact solutions of heterotic string theory corresponding to four-dimensional magnetic black holes in $N = 4$ supergravity are described. The solutions describe the black holes in the throat limit, and consist of a tensor product of an $SU(2)$ WZW orbifold with the linear dilaton vacuum, supersymmetrized to $(1,0)$ world sheet SUSY. One dimension of the $SU(2)$ model is interpreted as a compactified fifth dimension, leading to a four-dimensional solution with a Kaluza-Klein gauge field having a magnetic monopole background; this corresponds to a solution in $N = 4$ supergravity, since that theory is obtained by dimensional reduction of string theory.

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1. Introduction

The study of black holes in string theory has been pursued quite vigorously in the last few years, with the hope that long-standing puzzles may be resolved in the context of an apparently consistent (but incomplete) theory of quantum gravity.

The first place to search for black hole solutions is in the low energy effective theory of string theory in four dimensions, and many have been found. They fall into two categories: those that involve only the fundamental fields of string theory, namely the metric, antisymmetric tensor, and $E_8 \times E_8$ gauge fields; and those that involve other fields resulting from compactification, such as Kaluza-Klein gauge fields and scalars, etc. Solutions in the first category include the four dimensional solutions of refs. [1,2], as well as the five-dimensional solutions of [3,4]. Into the latter category fall the black holes of $N = 4$ supergravity [5], since all the gauge fields of that theory arise from reduction of the ten dimensional metric and antisymmetric tensor [6]. These are the solutions of interest in this paper.

After studying solutions of the low energy effective action, it is desirable to continue and find the corresponding exact vacua for string theory, i.e. backgrounds whose sigma model is conformally invariant. The first such exact black hole was constructed in [7]; this was a two dimensional solution. Five dimensional solutions were studied in [4,8]. Here two types of exact solution were found: asymptotically flat solutions having sufficient supersymmetry to be unrenormalized, and solutions corresponding to the infinite throat of an extremal black hole, without the asymptotically flat region.

The first four dimensional exact solutions, corresponding to the magnetic black holes of [1,2], were constructed in [9]; here only the infinite throat of the extremal black hole was exactly constructable. In this paper we show that a similar construction gives the throat limit of some black holes in the second category above, namely extreme magnetic black holes in $N = 4$ supergravity.

The outline of the paper is as follows. In section two we discuss the low energy solution whose exact analog we will find. We give it as a solution of five dimensional string theory, with compact fifth dimension, and show the correspondence with $N = 4$ theory. In section three we give the corresponding exact solution in bosonic string theory, and in section four we discuss supersymmetrization of the model to a heterotic model. Finally, in section five we explore the low-lying levels of the theory, and construct operators giving throat-widening deformations of the model.
2. Low Energy Solution

$N = 4$ supergravity in four dimensions is a truncation of dimensionally reduced ten dimensional $N = 1$ supergravity, which in turn is the low energy theory of the heterotic string, with the $E_8 \times E_8$ degrees of freedom truncated. (Note that by truncation we mean removal of fields from a theory in a way consistent with the equations of motion.) Therefore to find the exact string solution corresponding to an $N = 4$ solution, one must work in dimensions greater than four and find the appropriate correspondence between fields. The solution we are interested in here is the extreme magnetic black hole; in particular, only one of the six $N = 4$ gauge fields is nonvanishing, suggesting that we need only one extra dimension to find the string analog. Thus we begin with the bosonic action for string theory in five dimensions, with $E_8 \times E_8$ gauge fields truncated, and assuming an unspecified internal configuration for the remaining five dimensions:

$$S_5 = \int d^5x \sqrt{-\hat{g}} e^{-2\hat{\phi}} \left[ \hat{R} + 4(\nabla\hat{\phi})^2 - \frac{1}{12} \hat{H}^2 \right]$$

where hats refer to five-dimensional objects, and $\hat{H}_{\mu\nu\rho} \equiv (\hat{B}_{\mu\nu,\rho} + \hat{B}_{\nu\rho,\mu} + \hat{B}_{\rho\mu,\nu})$. We will assume the fifth co-ordinate $x^5$ has a range $0 \to 4\pi$; this is its natural range in later sections, and it enters into the dimensional reduction as a contribution to the dilaton shift (see below). Now we reduce to four dimensions, defining the four dimensional fields $g_{\mu\nu}, \ell, A^1, A^2, \phi, B_{\mu\nu}$ by

$$\hat{g}_{\mu\nu} = \left( g_{\mu\nu} + \ell^2 A^1_{\mu} A^1_{\nu} \right) \left( \ell^2 A^1_{\mu} \right)$$

$$A^2_{\mu} = \hat{B}_{\mu5}$$

$$\phi = \hat{\phi} - \frac{1}{2} \ln 4\pi \ell$$

$$B_{\mu\nu} = \hat{B}_{\mu\nu}.$$

With these definitions, standard formulas give (see e.g. [10])

$$S_4 = \int d^4x \sqrt{-g} e^{-2\phi} \left[ R + 4(\nabla\phi)^2 - \frac{\nabla\ell^2}{\ell^2} - \frac{\ell^2}{4} (F^1)^2 + \frac{1}{4\ell^2} (F^2)^2 - \frac{1}{12} (H + A^1 \wedge F^2)^2 \right]$$

The term involving $H$ may be dualized into a scalar axion, leading to its replacement by $\frac{1}{2} (\nabla\phi)^2 + \frac{1}{12} (F^1)^2 (F^2)^2$. Now we truncate the theory, following [6], by noticing that $\ell = \ell_0 = \text{const.}$ is consistent with the equations of motion if $\frac{1}{\ell_0} F^2 = \pm \ell_0 F^1$. With this motivation we define new gauge fields

$$A^\pm = \frac{1}{\sqrt{2}} \left( \frac{\ell}{2} A^1 \pm \frac{1}{2\ell} A^2 \right)$$
and one finds that $\ell = \ell_0$, $A^- = 0$ is a consistent truncation. Defining finally $F \equiv F^+$, the action becomes

$$I = \int d^4x \sqrt{-g} e^{-2\phi} \left[ R + 4(\nabla \phi)^2 - F^2 + \frac{1}{2}(\nabla a)^2 + \frac{1}{2}a F \tilde{F} \right].$$

This is the bosonic action of $N = 4$ supergravity, with 5 of 6 vector fields removed. We will be considering only purely magnetic black hole solutions, which satisfy $F \tilde{F} = 0$; thus we will take $a = \text{const}$. These black holes are a subset of those described in [5].

Now we write down the low energy extremal magnetic black hole solution, using the five dimensional variables. With co-ordinates $t, r, \theta, \phi, \zeta$ the metric is

$$ds^2 = -dt^2 + dr^2 + R^2 d\theta^2 + (R^2 \sin^2 \theta + Q^2 \ell_0^2 \cos^2 \theta) d\phi^2 + \ell_0^2 Q \cos \theta d\phi d\zeta$$

(2.2)

where $Q$ is a parameter measuring charge, and $R = R(r)$ satisfies

$$R' = 1 - \frac{Q \ell_0}{R}.$$

Note $r \in (-\infty, \infty)$. The antisymmetric tensor and dilaton are

$$\hat{B}_{\phi \zeta} = \ell_0^2 Q \cos \theta$$

$$\hat{\phi} - \hat{\phi}_0 = -\frac{1}{2} \ln \left( 1 - \frac{Q \ell_0}{R} \right).$$

(2.3)

The dimensionally reduced quantities are determined from the above using (2.1). In particular one finds $A^- = 0$, so this is a solution of the truncated theory as well; therefore it corresponds to a solution of the $N = 4$ theory. The solution describes a spherically symmetric asymptotically flat black hole, with no singularity but rather an infinitely long throat having asymptotic radius $Q \ell_0$.

Before going on to consider the corresponding exact solution, we pause to discuss the supersymmetry properties of this background. In [4], it is shown that this solution keeps two supersymmetries of the $N = 4$ theory unbroken; thus the five dimensional solution, with the other five dimensions flat, should preserve half the supersymmetries of the $D = 10$, $N = 1$ theory; and in fact this may be checked. One finds that unbroken supersymmetries are generated by four spinors $\epsilon_a$ having definite four-dimensional chirality (in the four dimensions spanned by $r, \theta, \phi, \zeta$). Furthermore, the spinors $\epsilon_a$ may be used to construct three complex structures, covariantly constant with respect to the generalized connection $\Omega^- \equiv \hat{\omega} - \frac{1}{2} \hat{H}$; this leads to enhanced $(4,0)$ worldsheet supersymmetry for the
heterotic sigma model with this background (see also [4], and references therein, for similar constructions and studies of extended supersymmetry in sigma models.)

Lastly we discuss the topology of the five dimensional solution. For the solution (2.2), one might expect the topology $R^2 \times S^2 \times S^1$; however, the monopole nature of the fields makes this impossible, and requires $S^1$ to be fibred over $S^2$ in a non-trivial way. What will appear in the sequel (and what appeared in [4]) is the space $R^2 \times S^3$, where $R^2$ is parametrized by $t, r$ and $S^3$ by $\theta, \phi, \zeta$ (and the $S^3$ may have further discrete identifications due to orbifolding.)

3. The Exact Solution

We are looking for a conformally invariant sigma model to describe the black hole throat. In this limit, the function $R$ is constant, and the radial direction decouples from the angular directions and becomes flat (but the dilaton will still vary with $r$). Thus to begin with we look for a conformally invariant sigma model for the angular directions alone. This is provided by an $SU(2)$ WZW model, as follows. For $g \in SU(2)$, we use the Euler angle parametrization $g = e^{i\phi S_3} e^{i\theta S_2} e^{i(\zeta - \phi) S_3}$

where $S_i = \frac{1}{2}\sigma_i$, and $\theta \in [0, \pi]$, $\phi \in [0, 2\pi]$, $\zeta \in [0, 4\pi]$. Then for the two components of the WZW action at level $k$ (bosonic for now) we compute

\[
S_1 = \frac{k}{16\pi} \int d^2 x \Tr \partial_\alpha g^{-1} \partial^\alpha g = \frac{k}{16\pi} \int d^2 x \left[ \frac{1}{2}(\nabla \theta)^2 + (1 - \cos \theta)(\nabla \phi)^2 + \frac{1}{2}(\nabla \zeta)^2 - (1 - \cos \theta)\nabla_\alpha \phi \nabla^\alpha \zeta \right]
\]

\[
S_2 = \frac{k}{24\pi} \int d^3 x \epsilon^{\alpha\beta\gamma} \Tr [g^{-1}\partial_\alpha g g^{-1}\partial_\beta g g^{-1}\partial_\gamma g]
\]

Comparing this to the sigma model action (with conveniently chosen $\alpha'$)

\[
S = \frac{1}{32\pi} \int d^2 x \partial_\alpha X^\mu \partial_\beta X^\nu (\eta^{\alpha\beta} G_{\mu\nu} + \epsilon^{\alpha\beta} B_{\mu\nu})
\]

gives the metric and antisymmetric tensor. The metric comes out to be

\[
G_{\mu\nu} = \begin{pmatrix}
k & 0 & 0 \\
0 & k \sin^2 \theta + k(\cos \theta - 1)^2 & k(\cos \theta - 1) \\
0 & k(\cos \theta - 1) & k
\end{pmatrix}
\]
with the co-ordinates ordered $\theta, \phi, \zeta$, and the torsion is (up to gauge)

$$\hat{B}_{\phi\zeta} = k(\cos \theta - 1).$$

We see this takes the form of the solution (2.2); note that $A^1_{\phi} = \cos \theta - 1$ is a gauge transform of $A^1_{\phi} = \cos \theta$, with gauge parameter $\Lambda = \phi$. The parameters have the values

$$Q = 1$$
$$\ell_0^2 = k.$$

Thus we have an interpretation of the $SU(2)$ WZW model as an exact bosonic string vacuum corresponding to a black hole in $N = 4$ supergravity (in the throat limit.) To get the complete black hole throat model in heterotic string theory, we must add the radial theory, which is the linear dilaton vacuum, and we must supersymmetrize both to $(1,0)$ supersymmetry. Also we must add an internal conformal field theory to fill out a central charge of $(15,26)$.

Lastly, we note the possibility of extending these models by orbifolding; this is accomplished by identifying points on the $SU(2)$ group manifold under $\zeta \rightarrow \zeta + \frac{4\pi}{m}$. Modular invariance requires $m$ to divide $k$, thus $k = mn$ (see [9]). To maintain the same co-ordinate range for the fifth co-ordinate, we define $\chi$ by $\zeta = \chi m$; then $\chi \in [0, 4\pi] \ (\text{recall that this range was assumed in (2.1)}).$ Then $G_{\mu\nu}$ may be rederived, and one finds the parameter values

$$Q = m$$
$$\ell_0^2 = \frac{n}{m}.$$

Thus one has a two-parameter family of models, with angular theory

$$\frac{SU(2)_{mn}}{Z(m)_{R}}.$$

However, in the remainder of the paper we focus on the case $m = 1$, since $m > 1$ complicates the expressions without any qualitative change.

4. The Heterotic Sigma Model

In this section we discuss the $(1,0)$ supersymmetric extension of the the WZW action (3.3). To supersymmetrize we start by coordinatizing the group manifold, so $g = g(x), \ldots$
$x \equiv (x^1, x^2, x^3)$. We will also need a basis for tangent space; for this a good choice is the left-invariant vierbein

$$L^a_\mu = \sqrt{2i} \Tr[S^a g^{-1} \partial_\mu g],$$

(4.1)

where the normalization is such that $L^a_\mu L^a_\nu$ is the group metric $g_{\mu\nu} = -\Tr[g^{-1} \partial_\mu g g^{-1} \partial_\nu g]$. The meaning of “left invariant” is the following: given a fixed element $h \in G$, we define a mapping $G \to G$ by left multiplication, $g \to hg$; or, in terms of coordinates, $x \to x'(x)$, where $x'$ has the appropriate form. A vierbein is pushed forward by this mapping:

$$e^a_\mu \to e'^a_\mu (x') \equiv \frac{\partial x^\nu}{\partial x'^\mu} e^a_\nu (x).$$

(4.2)

Then the vierbein is defined to be left-invariant if it is invariant under this mapping:

$$e'^a_\mu (x') = e^a_\mu (x').$$

(4.3)

One can verify that $L^a_\mu$ has this property. Left invariance is defined similarly for other types of tensor field.

Now we write the $(1,0)$ sigma model action, using the $(1,0)$ superfield

$$X^\mu = x^\mu + \theta E^\mu_a(x) \psi^a,$$

(4.4)

where $E^\mu_a$ is the inverse to $L^a_\mu$. The action is 

$$S_1 = \frac{1}{32\pi} \int d^2 x d\theta g_{\mu\nu}(X) D_+_\mu X_\nu \partial_- X^\nu$$

$$S_2 = \frac{1}{32\pi} \int d^2 x d\theta g_{\mu\nu}(X) D_+ X^\mu \partial_- X^\nu,$$

(4.5)

where $D_+ \equiv \frac{\partial}{\partial \theta} + i \theta \partial_+$ is the supercovariant derivative, and $G,B$ are the metric and torsion of the original bosonic WZW model. Working out this action in components, one finds

$$S_1 + S_2 = \text{WZW}_k + \frac{k}{8\pi} \int d^2 x \psi^a \partial_+ \psi^a.$$

(4.6)

It is just the bosonic WZW model at level $k$, plus free fermions.

This is simple enough, but confusion can arise when one considers different ways of introducing the fermions. In particular, choosing a different vierbein for the fermion part of equation (4.4) leads in general to a different theory; tangent space rotations are anomalous here. Furthermore, we now argue that in general one gets a theory in which the Kac-Moody symmetries $g \to f(z) g h(\bar{z})$ are anomalous; this is a disaster since without these
symmetries there is no reason to expect conformal invariance for the model. The argument
is simple: start with the action (4.6), and assume a different basis for the fermions, so the
fermion part becomes
\[ \int d^2 x \psi_a (M^{-1})^{a_b} \partial_+ (M^{b_c} \psi^c), \]  
where \( M^{a_b}(x) \) is some \( x \)-dependent rotation in tangent space; the inverse vierbein \( E^\mu_a M^{a_b} \)
will be left-invariant if and only if \( M^{a_b} \) is constant. Now, consider the transformation \( x \rightarrow x' \)
corresponding to \( g \rightarrow f(z)gh(\bar{z}) \). This leaves the WZW \( k \) part of the action invariant,
as usual, but on the fermions one requires the extra rotation \( \psi^a \rightarrow M^{-1}_{a_b}(x') M^{b_c}(x) \psi^c \).
But such a rotation is anomalous; therefore the Kac-Moody symmetry is as well. Thus, to
get the desired model, we must include the fermions with a left-invariant vierbein, as in
(4.4); in this case the fermions are free.

5. The Spectrum

The previous section discussed the \((1,0)\) supersymmetric WZW model from the point
of view of the heterotic sigma model. The end result was that the model contains a bosonic
WZW model at some level \( k \) plus three free fermions. In this section we discuss the \((1,0)\)
model further, beginning from this point. We also add fields corresponding to the radial
direction, which make up the \((1,0)\) linear dilaton vacuum. We explore the low lying levels
of the theory, and construct operators corresponding to throat-widening deformations.

The content of the model is as follows: right and left moving level \( k \) \( SU(2) \) Kac-Moody
algebras, denoted \( J^i(z), \bar{J}^i(\bar{z}) \); a free boson (the radial coordinate), denoted \( r \); and finally
four free right moving fermions \( \psi^a \) for \((1,0)\) supersymmetry. The fermions are treated
exactly as in the usual \( D = 10 \) heterotic string, \( i.e. \) both R and NS sectors are included
with the standard GSO projection \([13]\). This, along with modular invariance of the bosonic
model \([14]\), guarantees modular invariance of the full theory. It also guarantees spacetime
supersymmetry and absence of tachyons.

As discussed in section 2, the black holes in question have \( N = 2 \) spacetime super-
symmetry; this implies that they must have \((4,0)\) supersymmetry on the world sheet \([15]\).
Furthermore, the \( SU(2) \) Kac-Moody algebra which is contained in the \((4,0)\) supersym-
metry algebra must have level 1, and therefore cannot be the same as the \( SU(2)_k \) of the
bosonic theory. Thus this model should contain quite a large symmetry algebra, a fact
which has been noted and studied before in [13, 14, 17]. The right-moving \((4, 0)\) supersymmetry algebra consists of the stress tensor \(T\), four supercharges \(G^i\), and an \(SU(2)\) current algebra at level one, given as follows:

\[
T = -\frac{1}{2} \partial r \partial r - \frac{1}{k + 2} J^i J^i - \partial \psi^a \psi^a - \frac{1}{\sqrt{2k + 4}} \partial^2 r
\]

\[
G^0 = \sqrt{2} \partial r \psi^0 + \frac{2}{\sqrt{k + 2}} J^i \psi^i + \frac{4}{\sqrt{k + 2}} \psi^1 \psi^2 \psi^3 - \frac{2}{\sqrt{k + 2}} \partial \psi^0
\]

\[
G^1 = \sqrt{2} \partial r \psi^1 + \frac{2}{\sqrt{k + 2}} (-J^1 \psi^0 + J^2 \psi^3 - J^3 \psi^2) - \frac{4}{\sqrt{k + 2}} \psi^0 \psi^2 \psi^3 - \frac{2}{\sqrt{k + 2}} \partial \psi^1
\]

\(G^2, G^3\) = cyclic perms.

\[
A^-_i = \psi^0 \psi^i + \frac{1}{2} \epsilon_{ijk} \psi^j \psi^k.
\]

The derivative terms at the end of the expressions for \(T\) and \(G^a\) are due to the linear dilaton in the \(r\) direction. This algebra closes on itself and has \(c = 6\) [17]. Also on the right moving side another level 1 \(SU(2)\) algebra may be constructed from the fermions, namely

\[
A^+_i = -\psi^0 \psi^i + \frac{1}{2} \epsilon_{ijk} \psi^j \psi^k.
\]

Meanwhile the left moving side has only the current algebra \(\tilde{J}^i\) at level \(k\), with stress tensor as above, minus fermion terms. Left moving central charge is 4.

Next we construct an operator corresponding to a throat widening deformation, as evidence that the model can be connected to an asymptotically flat region. Such a deformation should alter \(R(r)\) from constant to increasing; at the same time, it should keep \(\ell\) constant, to maintain the correspondence with the \(N = 4\) theory. For the purely bosonic theory such an operator was described in [13]; it is

\[
V_b = (\tilde{J}^i \tilde{\phi}^1)(J^1 \phi^1 + J^2 \phi^2)e^{\alpha r},
\]

where \(\phi^j_m, \tilde{\phi}^j_m\) are the WZW primary fields, and \(\alpha\) will be determined later. For the heterotic model we must supersymmetrize the right moving part, i.e. find an operator with the same bosonic content which is the highest component of a superfield. To do this we start with the superconformal primary field (lowest component of a superfield)

\[
W = -\sqrt{k + 2} (\psi^1 \phi^1_1 + \psi^2 \phi^1_2) e^{\alpha r}.
\]

Using the operator product

\[
J^i(z) \phi^1_j(0) \sim \frac{1}{z} \epsilon^{ijk} \phi^1_k
\]
one sees that $W$ has a $\frac{1}{z}$ singularity in its operator product with $G^0$, as it should. The coefficient of $\frac{1}{z}$ is the supersymmetric operator we are looking for. It is

$$V = (J^1\phi_1^1 + J^2\phi_1^2 - 4\phi_2^1\psi^1\psi^2) e^{\alpha r} - \alpha \sqrt{2k+4}\psi_0 (\psi^1\phi_1^1 + \psi^2\phi_2^1) e^{\alpha r} \tag{5.6}$$

where $\alpha = \sqrt{\frac{2}{k+2}}$ to give dimension $(1,1)$. Note that this deformation breaks the supersymmetries associated to $G^i$, leaving only $(1,0)$, and in particular eliminating spacetime supersymmetry. Thus the deformation is away from extremality.

Next we discuss the low-lying spectrum of the model, checking for massless or tachyonic states. The vertex operators (evaluated at $z = \bar{z} = 0$) take the general form

$$e^{-\beta r} e^{ipr} (\tilde{j}_K \cdot \tilde{\phi}_m^j) (j_K \cdot \phi_m^j) V_{\text{ferm}} V_{\text{int}} \tilde{V}_{\text{int}} \tag{5.7}$$

where $j_K$ and $\tilde{j}_K$ represent products of WZW raising operators, $V_{\text{ferm}}$ comes from the right moving fermions $\psi^a$, and the operators labelled “int” derive from the internal theory. Ghost factors have not been written in this expression. The term $e^{-\beta r} e^{ipr}$ is a primary field of the linear dilaton theory of the $r$ direction, where $\beta$ must be chosen such that the operator has real conformal weight for all real values of $p$. Noting the coefficient of the $\partial^2 r$ term in $T$, this fixes $\beta = \frac{1}{\sqrt{2k+4}}$, giving the operator a weight of

$$L_0 = \tilde{L}_0 = \frac{1}{2} p^2 + \frac{1}{2k+4} \tag{5.8}$$

Now we study the spectrum, beginning with the Ramond sector. Working in the $-\frac{1}{2}$ picture, a state must have weights $(0,1)$, thus

$$0 = \frac{1}{2} p^2 + \frac{1}{2k+4} + \frac{j(j+1)}{k+2} + |K| + \frac{1}{2} N_f + L_0^{\text{int}}$$

$$1 = \frac{1}{2} p^2 + \frac{1}{2k+4} + \frac{j(j+1)}{k+2} + |\tilde{K}| + \tilde{L}_0^{\text{int}}. \tag{5.9}$$

From the first equation above, we immediately have

$$p^2 < -\frac{1}{k+2},$$

thus ruling out massless and tachyonic states in the R sector. In the NS sector, we work in the $-1$ picture, in which states have weights $(\frac{3}{2},1)$, giving

$$\frac{1}{2} = \frac{1}{2} p^2 + \frac{1}{2k+4} + \frac{j(j+1)}{k+2} + |K| + \frac{1}{2} N_f + L_0^{\text{int}}$$

$$1 = \frac{1}{2} p^2 + \frac{1}{2k+4} + \frac{j(j+1)}{k+2} + |\tilde{K}| + \tilde{L}_0^{\text{int}}.$$
Evidently $p^2 \leq 0$ requires $|K| = |\tilde{K}| = N_f = 0$, and $\tilde{L}_0^{\text{int}} - L_0^{\text{int}} = \frac{1}{2}$. The simplest possibility is an NS fermion from the current algebra of the internal theory; however, such a state is eliminated by the GSO projection. Thus this model contains no tachyons or massless states propagating in the black hole throat.

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