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Dilaton black holes near the horizon

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Generic $U(1)^2$ four-dimensional (4D) black holes with unbroken $N = 1$ supersymmetry are shown to tend to a Robinson-Bertotti-type geometry with a linear dilaton and doubling of unbroken supersymmetries near the horizon. Purely magnetic dilatonic black holes, which have unbroken $N = 2$ supersymmetry, behave near the horizon as a 2D linear dilaton vacuum $\otimes S^2$. This geometry is invariant under 8 supersymmetries; i.e., half of the original $N = 4$ supersymmetries are unbroken. The supersymmetric positivity bound, which requires the mass of the 4D dilaton black holes to be greater than or equal to the central charge, corresponds to positivity of mass for a class of stringy 2D black holes.

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The evaporation of stringy $U(1)^2$ charged black holes [1, 2] may be understood as the process of restoration of supersymmetry [3]. It is likely that the end points of the process of evaporation for charged dilatonic black holes are stable remnants which are zero-temperature extreme black holes with some unbroken supersymmetry. Those black holes have the minimum possible mass $M = \frac{1}{\sqrt{2}}(|Q| + |P|)$. It was shown in [3] that the extreme solutions saturate the supersymmetry bound of $N = 4, d = 4$ supergravity, or dimensionally reduced superstring theory. When both electric and magnetic charges are present, only one bound of $N = 4$ supersymmetry is saturated and the corresponding solution has unbroken $N = 1$ supersymmetry. In the case of only electric or only magnetic charges, i.e., just U(1) black holes, both supersymmetry bounds of $N = 4$ supersymmetry are saturated and these black holes have unbroken $N = 2$ supersymmetry [3].

Several properties of extreme dilatonic black holes were investigated in [3]. The purpose of this paper is to investigate the properties of extreme supersymmetric stringy black holes in the neighborhood of the horizon, to find what kind of a geometry they tend to and what happens with unbroken supersymmetries in those geometries. We would also like to understand the relation to two-dimensional (2D) dilatonic black holes and whether the investigations of the supersymmetry and of the geometry near the horizon of the 4D black hole may be an important factor for understanding the late stages of evaporation of 2D black holes.

We will start by describing the behavior of the familiar extreme Reissner-Nordström black hole near the horizon [4], where the geometry becomes that of the Robinson-Bertotti solution and has twice as many supersymmetries as the extreme black hole.

Then we will study properties of electric-magnetic dilaton black holes near the horizon. The difference with the Reissner-Nordström black hole near the horizon will show up in the existence of a nonconstant dilaton, the linear dependence on the coordinate being proportional to the dilaton charge. We will find, however, that the essential properties of the geometry at the horizon are very close

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to those of the Robinson-Bertotti solution and also that doubling of supersymmetries takes place.

The geometry near the horizon and the supersymmetry properties of purely magnetic extreme black holes is the next topic to be studied. The fact that near the horizon these solutions form a geometry which is the direct product of a 2D linear dilaton vacuum and a two-sphere of constant curvature has been established before [5]. It was argued that, at scales where the two-sphere radius, which is proportional to the magnetic charge P, can be neglected, the black-hole physics may be described by a 2D effective theory. This is one of the reasons for the recent interest in two-dimensional black holes. It has also been established before [3] that the extreme magnetic dilaton black hole has 8 unbroken parameters of $N = 2$ supersymmetry, i.e., one-half of the original 16 parameters of $N = 4$. In this paper we analyze what happens with unbroken supersymmetries of the extreme black hole at the horizon. We find that the solution (a direct product of the 2D linear dilaton vacuum and a two-sphere of constant curvature) has 8 unbroken parameters of $N = 2$ supersymmetry, exactly as does the total geometry of the extreme black hole; there is no doubling of the supersymmetries to the maximum possible $N = 4$. Finally, the relation between the supersymmetric positivity bound $M \geq \frac{1}{\sqrt{2}}|P|$ in 4D and positivity of mass in 2D is exhibited.

The extreme Reissner-Nordström black hole near the horizon has been investigated before [4]. There are several properties of the extreme Reissner-Nordström black hole which make it interesting. One is that it has zero temperature, and so is stable to emission of Hawking radiation. The extreme Reissner-Nordström configuration also possesses unbroken $N = 1$ supersymmetry [6, 4] when viewed as a bosonic configuration of $N = 2, d = 4$ supergravity. A further property is that near its horizon the extreme Reissner-Nordström geometry asymptotes to a Robinson-Bertotti geometry, which is a maximally supersymmetric [4] (and homogeneous) configuration of $N = 2, d = 4$ supergravity. To see this, consider the extreme Reissner-Nordström (RN) metric in isotropic coordinates $\{x^i\}$:

$$ds^2_{\rm RN} = V^2_{\rm RN} dt^2 - V^{-2}_{\rm RN} d\mathbf{x}^2 \ , \tag{1}$$

where $|\mathbf{x}| = \rho$ and

$$V^{-1}_{\rm RN}(\rho) = 1 + M/\rho \ . \tag{2}$$

The Maxwell field is

$$F_{\rm RN} = \pm dV_{\rm RN} \wedge dt \ , \tag{3}$$

and $M = |Q|$. Near the horizon $\rho \to 0$, the function $V^{-1}_{\rm RN} \to \frac{M}{\rho} = V^{-1}_{\rm RB}$, i.e., the extreme Reissner-Nordström metric tends to the Robinson-Bertotti (RB) metric there, and the Maxwell field tends to that of the RB configuration:

$$\begin{aligned} ds^2_{\rm RB} &= \frac{\rho^2}{M^2}\, dt^2 - \frac{M^2}{\rho^2}\, d\rho^2 - M^2\, d\Omega^2 \ , \\ F_{\rm RB} &= \pm \frac{1}{M} d\rho \wedge dt \ . \end{aligned} \tag{4}$$

Since the extreme Reissner-Nordström geometry admits $N = 1$ supersymmetry and the Robinson-Bertotti geometry admits $N = 2$ supersymmetry, one may call this phenomenon "doubling of supersymmetries near the black-hole horizon." To explain this doubling of supersymmetries, let us consider the supersymmetric transformation of the *gravitino field strength* [7]:

$$\delta \Psi_{ABCI} = C^+_{ABCD}\epsilon^{D+}_I + C^-_{ABCD}\epsilon^{D-}_I \ . \tag{5}$$

In Eq. (5), two-dimensional spinor notation is used. The supersymmetry parameters $\epsilon^\pm_I, I = 1, 2$, are defined in Eqs. (37) of Ref. [7], and $C^\pm$ are the following combinations of the Weyl (C_{ABCD}) and Maxwell (F_{CD}) spinors:

$$C^\pm_{ABCD} \equiv C_{ABCD} \pm \nabla_{AB'}\, F_{CD}\, V^{-1}\, K^{B'}_B \ , \tag{6}$$

where $K^{B'}_B$ is the Killing vector. For the extreme Reissner-Nordström black hole, we have either $C^+_{ABCD} = 0$ or $C^-_{ABCD} = 0$ (depending on the sign of the charge Q), which is a relation between the Weyl spinor and the derivative of the Maxwell spinor. In the first case the unbroken $N = 1$ supersymmetry parameter of the extreme Reissner-Nordström black hole is ϵ^+_I, in the second case it is ϵ^-_I.

Near the horizon, the Reissner-Nordström geometry becomes that of the Robinson-Bertotti solution, which is *conformally flat and has a covariantly constant Maxwell field*:

$$C_{ABCD} = 0, \qquad \nabla_{AB'}\, F_{CD} = 0 \ . \tag{7}$$

Thus, both combinations of the Weyl and Maxwell spinors which enter the supersymmetry variation of the gravitino field strength (5) are vanishing in this geometry. This property ensures that all 8 parameters (ϵ^+_I and ϵ^-_I) of the original $N = 2$ supersymmetry are unbroken in the Robinson-Bertotti background.

We see that the sign of the charge controls which supersymmetries are unbroken in the Reissner-Nordström solution, but that no such phenomenon occurs for the Robinson-Bertotti solution where *all* supersymmetries are unbroken due to the separate vanishing of both terms making up the generalized Weyl curvatures $C^\pm$ in Eq. (6).

Notice also that the extreme Reissner-Nordström metric is asymptotically Minkowskian, since at infinity only the constant term in $V^{-1}_{\rm RN}$ survives. Minkowski space is also a maximally supersymmetric geometry. Therefore, in some sense, the extreme RN solution may be thought of as a soliton; it interpolates between two candidate vacua for $N = 2, d = 4$ supergravity [6].

A natural question now arises: how much of this behavior carries over to extreme dilaton black holes?

Doubling of supersymmetries near the horizon for extreme electric-magnetic dilaton black holes will be explained in what follows.

The action we will use is the part of the SO(4) version of the $N = 4, d = 4$ supergravity action without an axion,

$$I = \frac{1}{16\pi}\int d^4x\,\sqrt{-g}[\,-R + 2\,\partial^\mu\phi\cdot\partial_\mu\phi - (e^{-2\phi}F^{\mu\nu}F_{\mu\nu} + e^{2\phi}\tilde{G}^{\mu\nu}\tilde{G}_{\mu\nu})]\ , \tag{8}$$

where $\tilde{G}_{\mu\nu}$ is related to the nondually rotated field $G_{\mu\nu}$ as

$$\tilde{G}^{\mu\nu} = \frac{i}{2}\frac{1}{\sqrt{-g}}\,e^{-2\phi}\,\epsilon^{\mu\nu\lambda\delta}\,G_{\lambda\delta}\ . \tag{9}$$

All notation is that of [3]. For extreme supersymmetric dilatonic black holes, the fields are built out of two functions H_1 and H_2 [3] :

$$\begin{aligned} &ds^2 = e^{2U}dt^2 - e^{-2U}d\mathbf{x}^2\ ,\\ &A = \psi dt, \qquad \tilde{B} = \chi dt\ ,\\ &F = d\psi\wedge dt, \qquad \tilde{G} = d\chi\wedge dt\ ,\\ &e^{-2U} = H_1H_2, \qquad e^{2\phi} = H_2/H_1\\ &\sqrt{2}\,\psi = \pm H_1^{-1}, \qquad \sqrt{2}\,\chi = \pm H_2^{-1}\ , \end{aligned} \tag{10}$$

where the condition on the functions H_1, H_2 is that they be harmonic,

$$\partial_i\partial_i H_1 = 0\ , \qquad \partial_i\partial_i H_2 = 0\,. \tag{11}$$

We have used isotropic coordinates $\{x^i\}$, where $\rho^2 = x^ix^i$ and

$$H_1 = e^{-\phi_0}(1+\sqrt{2}|Q|/\rho)\ , \qquad H_2 = e^{+\phi_0}(1+\sqrt{2}|P|/\rho)\ . \tag{12}$$

The mass M and dilaton charge Σ are related to the U(1) electric Q and magnetic P charges as

$$M = \frac{|P|+|Q|}{\sqrt{2}}\ , \qquad \Sigma = \frac{|P|-|Q|}{\sqrt{2}}\ . \tag{13}$$

It was shown in [3] that the bosonic background (10), (11) admits supercovariant Killing spinors of $N=4, d=4$ supergravity, i.e., that there exist nontrivial solutions of the equations

$$\delta\Psi_{\mu I}(\epsilon) = \delta\Lambda_I(\epsilon) = 0\ , \qquad I = 1,2,3,4, \tag{14}$$

describing the supersymmetry variation of 4 gravitinos and 4 dilatinos in the background (10), (11).

Specifically, there is always some unbroken $N=1$ supersymmetry for $PQ\neq 0$ extreme black holes (one-quarter of $N=4$ supersymmetry). For example, for $P>0, Q>0$ the unbroken $N=1$ supersymmetry for the solution (10), (11) is one combination of third and fourth supersymmetry, $\epsilon_+^{34}, \epsilon_{34}^+$, in the notation of Ref. [3], the first and the second being broken. The space-time dependence of the Killing spinor in the canonical geometry (10) is given by $\epsilon = e^{\frac{1}{2}U}\epsilon_0$ where ϵ_0 is a constant spinor.

Consider the extreme $PQ\neq 0$ dilatonic black holes near the horizon, i.e., in the limit $\rho\to 0$. The metric in (10) becomes

$$ds^2 = \frac{\rho^2}{M^2-\Sigma^2}\,dt^2 - \frac{M^2-\Sigma^2}{\rho^2}\,d\rho^2 - (M^2-\Sigma^2)\,d\Omega^2\ . \tag{15}$$

This metric is precisely the Robinson-Bertotti metric (4) familiar from $N=2$ supergravity [6]. The mass parameter of the Robinson-Bertotti metric is, in this case,

$$M_{\rm RB} = \sqrt{M^2-\Sigma^2} = \sqrt{2|PQ|}\ . \tag{16}$$

The dilaton for these solutions behaves as

$$e^{2\phi} = e^{2\phi_0}\left|\frac{P}{Q}\right|\left(1 - \frac{\sqrt{2}\,\Sigma}{M_{\rm RB}^2}\,\rho + O(\rho^2)\right)\ , \tag{17}$$

so we see that the term linear in ρ is proportional to the dilaton charge Σ. The electric and magnetic fields are given by

$$F = e^{\phi_0}\,\frac{1}{2Q}\,d\rho\wedge dt\ , \qquad \tilde{G} = e^{-\phi_0}\,\frac{1}{2P}\,d\rho\wedge dt\,. \tag{18}$$

Since the dilaton field has a term linear in ρ, we will call this solution a "Robinson-Bertotti type" geometry.

Let us now see if, near the horizon, the extreme dilaton black holes with $PQ\neq 0$ possess any additional supersymmetry and, if so, how the charges control which ones are unbroken.

First consider the dilatino transformations rules in the notation of [3]:

$$\begin{aligned} \frac{1}{2}\delta\Lambda_I &= -\gamma^\mu\epsilon_I\partial_\mu\phi\\ &+\frac{1}{\sqrt{2}}\sigma^{\mu\nu}\left(e^{-\phi}F_{\mu\nu}\alpha_{IJ} - e^{\phi}\tilde{G}_{\mu\nu}\beta_{IJ}\right)^-\epsilon^J\\ &=0\ . \end{aligned} \tag{19}$$

The first term in (19), $(-\gamma^a e_a^\mu\partial_\mu\phi)$, involves the *flat space* derivative of the dilaton, which vanishes near the horizon $\rho=0$:

$$e_a^\mu\partial_\mu\phi \sim \rho\to 0\,, \tag{20}$$

where we used Eqs. (17), (15). The second term in Eq. (19) is proportional to

$$[\alpha_{IJ} - {\rm sgn}(PQ)\beta_{IJ}]\epsilon^J\ , \tag{21}$$

so we see that this time it is the sign of PQ that controls which combinations of supersymmetries are broken (or unbroken).

For positive PQ, the term in square brackets in front of ϵ^J in (21) vanishes under the condition that

$$\epsilon^1 = \epsilon^2 = \epsilon_1 = \epsilon_2 = 0\ ; \tag{22}$$

i.e., the first and second supersymmetries are broken, and there are no constraints on the third and fourth supersymmetries. For negative PQ, the third and fourth supersymmetries are broken, and the first and second supersymmetries are not:

$$\epsilon^3 = \epsilon^4 = \epsilon_3 = \epsilon_4 = 0\ . \tag{23}$$

The next step is to investigate whether there are additional constraints coming from the supersymmetry transformation of the gravitino field strength. Consider first positive PQ. The variation of the first and the second gravitino vanishes because of Eqs. (22). The transformation rules for the third and fourth gravitino field strength are similar to those of $N = 2$ supergravity as given in Eq. (5). Therefore there are indeed no additional constraints coming from $\delta\Psi_{I\mu}$. There are several remarkable properties of the dilatonic black hole near the horizon making this $N = 2$ supersymmetry possible. The first is that the scalar curvature vanishes near the horizon,

$$R = 2g^{\mu\nu}\partial_\mu\phi \cdot \partial_\nu\phi = 2g^{rr}(\partial_r\phi)^2 \sim \rho^2 \to 0 \ , \qquad (24)$$

so that the Weyl spinor is identical to the Ricci spinor. The second is that the geometry is of Robinson-Bertotti type, as noted above, so that the covariant derivatives of the vector fields vanish and the Weyl spinor is zero. Finally, as shown in Eq. (20), the flat space derivative of the dilaton vanishes.

Therefore, we have the following situation near the horizon: for positive PQ the unbroken $N = 2$ supersymmetry consists of the third and the fourth ones, and for negative values of PQ it is reverse: the first and the second supersymmetries are unbroken whereas the third and the fourth are broken. Note that despite the doubling of unbroken supersymmetries near the horizon the unbroken supersymmetry never becomes equal to the maximal possible one: only half of the possible $N = 4$ supersymmetries are restored. This is different from the classical extreme Reissner-Nordström solution which, near the horizon, restores maximal supersymmetry of this theory, namely $N = 2$.

The differences with the classical extreme Reissner-Nordström solution near the horizon are also that a nontrivial dilaton dependence on ρ exists and is proportional to the dilaton charge Σ, and that the parameter $M^2_{\rm RB}$ in the Robinson-Bertotti type geometry is not the square of the mass of the extreme black hole but the difference $M^2 - \Sigma^2$. It is exactly this parameter in front of the angular part of the RB metric in Eqs. (4) and (15) which defines the area of the horizon of RN and dilaton black holes. It is this parameter which is never zero for the nontrivial classical extreme black hole, but becomes zero for dilatonic extreme black holes when $M^2 = \Sigma^2$. This happens in the limit when either the electric or magnetic charge vanishes. Such extreme black holes have unbroken $N = 2$ supersymmetry [3] and their geometry near the horizon is very different from the one described above where the limit to the horizon is taken with both electric and magnetic charges present.

Here we will consider the purely magnetic U(1) dilatonic black holes. The electric case in the canonical metric may be obtained trivially from it by the replacements $P \to Q$ and $\phi \to -\phi$. The magnetic black hole in the stringy metric turns out to have more interesting properties than the purely electric one [2, 3]. We use again the construction given in Eqs. (10) and (11). As explained above, this ansatz solves both the $N = 4$ supergravity equation of motion and the supersymmetry equations (14), leaving $N = 2$ supersymmetry unbroken. For purely magnetic extreme black holes we have, in isotropic coordinates,

$$H_1 = e^{-\phi_0} \ , \qquad H_2 = e^{+\phi_0}(1 + \sqrt{2}|P|/\rho) \ . \qquad (25)$$

The mass M and dilaton charge Σ are related to the magnetic charge P as follows:

$$M = \Sigma = |P|/\sqrt{2} \ . \qquad (26)$$

The metric for the extreme magnetic dilaton black hole already possesses $N = 2$ supersymmetry when viewed as a bosonic solution of $N = 4, d = 4$ supergravity [3]. However, as can be checked, going to the horizon does *not* result in additional supersymmetries. The geometry near the horizon just keeps all 8 unbroken supersymmetries of the black hole. To see this, let us consider the limit to the horizon of the extreme purely magnetic black hole. This limit has been studied before in [5]. However, we will investigate this limit by using our construction (10), (11), which automatically solves the Killing equations for the unbroken supersymmetry.

The metric has different behavior to that of the black holes of the previous section due to the fact that now $PQ = 0$. The most important difference is the fact that $g_{tt}^{-1} = g_{\rho\rho}$ is linear in ρ^{-1} rather then quadratic, as was the case for classical extreme Reissner-Nordström metric near the horizon and the $PQ \neq 0$ dilatonic extreme black holes considered above. In addition, the dilaton behaves differently. The metric, dilaton and vector fields near the horizon are

$$ds^2 = \frac{\rho}{\sqrt{2}|P|}dt^2 - \frac{\sqrt{2}|P|}{\rho}(d\rho^2 + \rho^2 d\Omega^2) \qquad (27)$$

$$= \frac{\rho}{\sqrt{2}|P|}\left[dt^2 - \frac{2P^2}{\rho^2}(d\rho^2 + \rho^2 d\Omega^2)\right] \ , \qquad (28)$$

$$e^{2\phi} = e^{2\phi_0}\sqrt{2}|P|/\rho \ , \qquad (29)$$

$$G = P\, e^{2\phi_0} \sin\theta\, d\theta \wedge d\phi \ . \qquad (30)$$

Now let us make a change of variables,

$$w = -\sqrt{2}\, P \ln\rho \ , \qquad (31)$$

so that the metric becomes

$$ds^2 = \frac{1}{\sqrt{2}|P|}\exp\left(-w/\sqrt{2}P\right)\left[dt^2 - dw^2 - 2P^2 d\Omega^2\right] \ , \qquad (32)$$

and the dilaton becomes

$$\phi = \phi_0 + \tfrac{1}{2}\ln(\sqrt{2}|P|) + w/2\sqrt{2}P \ , \qquad (33)$$

or

$$e^{2\phi} = \sqrt{2}|P|\, e^{2\phi_0} \exp(w/\sqrt{2}P) \ . \qquad (34)$$

Going to the stringy metric via the transformation $ds^2_{\rm str} = e^{2\phi} ds^2$, we obtain

$$ds^2_{\rm str} = dt^2 - dw^2 - 2P^2 d\Omega^2 \ , \qquad (35)$$

which is the direct product of a flat 2D Minkowskian metric in the coordinates (t, w) and a transverse space of constant curvature $(1/\sqrt{2P^2})$. In these coordinates we see that the dilaton is linear. This is a well-known result in the context of 2D "dilaton gravity" [5, 8].

The fact that the direct product of the linear dilaton vacuum and a sphere of a constant curvature has 8 unbroken supersymmetries follows simply from the fact that this is a specific example of the construction (10), (11), which has been proven [3] to have unbroken $N = 2, d = 4$ supersymmetry. For positive P the unbroken supersymmetries are given by the 8 specific combinations [3] $\epsilon_{+}^{34}, \epsilon_{34}^{+}, \epsilon_{-}^{12}, \epsilon_{12}^{-}$ of the original 16 supersymmetries, the other 8 are broken. For negative P those which were broken become unbroken and vice versa.

Let us see why there is no doubling of supersymmetries near the horizon this time, as different from the previous cases. Consider again the dilatino transformation rules (19). Now that the metric behaves differently at the horizon, there is a contribution from the first term as well as from the vector field. The sum of those two contributions vanishes under the condition that half of the supersymmetries are broken. The unbroken supersymmetries are listed above: going near the horizon does not change the number of unbroken supersymmetries of the extreme purely magnetic (or electric) dilatonic black hole.

As argued in [5] the evaporation of or scattering by purely magnetic 4D black holes near extremality is closely related to the analogous processes for 2D black holes. It can be shown that the relation between the masses of the 2D and 4D black holes is

$$M_{2D} \sim \frac{M_{4D} - |z|}{|z|} , \tag{36}$$

where the two central charges of $N = 4$ supersymmetry are defined in [3] to be $|z_1| = |z_2| = |z| = |P|/\sqrt{2}$. The supersymmetry bound for 4D dilaton black holes derived in [3],

$$M_{4D} - |z| \geq 0 , \tag{37}$$

thus ensures that the mass of the 2D black holes is nonnegative:

$$M_{2D} \geq 0 . \tag{38}$$

The saturation of the bound in 4D takes place when the mass reaches the value of the central charge, $M_{4D} = |z|$, $N = 2$ supersymmetry becomes restored and the evaporation stops. In 2D this corresponds to total evaporation of the black hole $M_{2D} = 0$, i.e., to the linear dilaton vacuum.

Thus we conclude that stringy 2D black holes, which are related to 4D black holes as in Eqs. (36)–(38), may have better control over quantum corrections and stability of the extreme solution due to the existence of 8 unbroken supersymmetries and the corresponding nonrenormalization theorems [3]. However, the role of an additional two-sphere of a radius $2|z|$, which is required in the geometry (35), (33) to keep supersymmetry unbroken, is still to be understood. In particular, a supersymmetric embedding of 2D dilaton gravity may be looked for, which has the linear dilaton vacuum with flat two-dimensional metric as the solution with unbroken supersymmetry [9]. This supersymmetry will take place in a purely two-dimensional theory without an additional two-sphere. This kind of supersymmetry may provide constraints on the quantum theory in $d = 2$ and be useful in the context of a two-dimensional toy model of quantum gravity. However, if one considers investigation of 2D black holes as a way to get insight into a more complicated (but more realistic) theory of 4D black holes, one should remember that physics of the four-dimensional black holes near the horizon is effectively two-dimensional and described by 2D effective Lagrangian only in the low energy limit. It is not clear whether 2D black holes may have all 8 supersymmetries possessed by 4D black holes near the horizon. Therefore the problem of stable remnants may have quite different solutions, depending on whether the geometry is purely two dimensional or four dimensional.

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