Weak Field Black Hole Formation in Asymptotically AdS Spacetimes

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Based on arXiv [0904.0464] with Shiraz Minwalla

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References

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- Shu Lin, Edward Shuryak, "Toward the AdS/CFT Gravity Dual for High Energy Collisions. 3. Gravitationally Collapsing Shell and Quasiequilibrium" arXiv:0808.0910 [hep-th]
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Introduction

- The AdS/CFT correspondence can be used to study the dynamical passage of a system from a pure state to an approximately thermalized state .
- This process is dual to the process of black hole formation via gravitational collapse.
- In this talk we study asymptotically *AdS* collapse processes, in a weak field limit, that display rich dynamics while allowing for analytic control.

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The Basic Setup

An *AdS* collapse process can be set up as follows.

- Start with vacuum AdS.
- At time $= 0$ weakly perturb the system at the boundary by turning on the non normalizable component of a massless field for a finite duration.
- The perturbation creates an ingoing shell of the massless field which sometimes collapses to form black holes.

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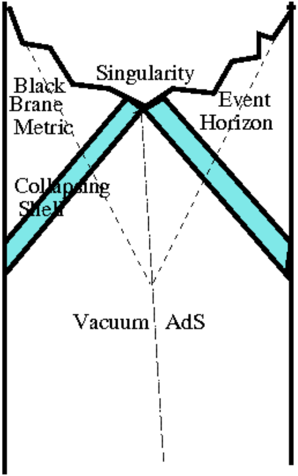
Penrose diagram



Translationally invariant collapse

- Our system: Negative cosmological constant Einstein gravity with a minimally coupled massless scalar field

$$S = \int d^{d+1}x\sqrt{g}\left(R + d(d-1) - \frac{1}{2}(\partial\phi)^2\right)$$

- We study spacetimes of the form (Eddington-Finkelstein gauge)

$$ds^2 = 2dr\, dv - g(r,v)dv^2 + f^2(r,v)dx_i^2$$
$$\phi = \phi(r,v).$$

- f, g and ϕ are independent of (x_i) : Translational invariance.

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$$ds^2 = 2dr\,dv - g(r,v)dv^2 + f^2(r,v)dx_i^2$$

The lines of constant v are null ingoing geodesics.

- When
$$g(r,v) = r^2, \quad f(r,v) = r$$
it is the metric of Poincare patch AdS space in Eddington Finkelstein coordinate.
- When
$$g(r,v) = r^2\left(1 - \frac{M}{r^3}\right), \quad f(r,v) = r$$
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Set up

- Initial condition:
  For $v < 0$ the metric is pure AdS and dilaton field is zero.
- Boundary condition:
  Metric is asymptotically Poincare AdS. Non normalizable component $\phi_0(v)$ of the dilaton field turned on.

$$\lim_{r \to \infty} \phi(r, v) = \phi_0(v)$$

$$\phi_0(v) \neq 0 \quad 0 \leq v \leq \delta t$$

$$|\phi_0(v)| \sim \epsilon \ll 1$$

- Goal: Given above data,
  determine space-time metric and dilaton for all r and v.

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Naive perturbation expansion

- Note that δt can be scaled away. Qualitatively the amplitude ϵ is the only parameter. In this talk we work in the limit of small ϵ
- A weak boundary perturbation creates a small amplitude wave that propagates from the boundary into the bulk of AdS.
- As the amplitude of the wave is small it is natural to attempt to construct our spacetime in a perturbation expansion (in ϵ) about empty AdS space. We refer to this as the naive perturbative expansion.

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Naive perturbation expansion

- We now implement the naive perturbation expansion in ϵ around the empty AdS.

$$f(r,v) = \sum_{n=0}^{\infty} \epsilon^n f_n(r,v), \quad g(r,v) = \sum_{n=0}^{\infty} \epsilon^n g_n(r,v)$$

$$\phi(r,v) = \sum_{n=0}^{\infty} \epsilon^n \phi_n(r,v) \quad \text{with}$$

$$f_0(r,v) = r, \quad g_0(r,v) = r^2, \quad \phi_0(r,v) = 0.$$

First correction

- We have determined f, g and ϕ upto 5th order in ϵ by solving the bulk equations. At first and second order

$$\phi_1(r,v) = \phi_0(v) + \frac{\dot{\phi}_0}{r}$$

$$f_2(r,v) = -\frac{\dot{\phi}_0^2}{8r}$$

$$g_2(r,v) = -\frac{C_2(v)}{r} - \frac{3}{4}\dot{\phi}_0^2$$

$$C_2(v) = -\int_{-\infty}^{v} dt\, \frac{\dot{\phi}_0 \dddot{\phi}_0}{2}$$

▸ Higher order results

Validity of perturbative expansion

$$\phi_1(r, v) = \phi_0(v) + \frac{\dot{\phi}_0}{r}$$

- Note that $\frac{\dot{\phi}_0}{r} \sim \frac{\epsilon}{r\delta t}$. Consequently $\phi_1 \gg 1$ for $r\delta t \ll \epsilon$. Therefore this perturbation theory, which is an expansion in the amplitude of ϕ, breaks down for $r\delta t \ll \epsilon$. Singular at $r = 0$.
- More generally it can be shown that naive expansion is valid whenever

$$r\delta t \gg Max\left\{\epsilon\sqrt{\frac{v}{\delta t}}, \epsilon\right\} \text{ and } \frac{v}{\delta t} \ll \epsilon^{-\frac{2}{3}}$$

- So naive perturbation breaks down at small r and large v.

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Validity of perturbative expansion

- We will show below that the small r region where perturbation fails is shielded from the boundary by an event horizon.
- Breakdown at large v (IR divergence) will be important. It is a consequence of an aspect of our solution that is visible already at small v and so is reliably displayed in naive perturbation theory.
- For $r \gg \frac{\epsilon}{\delta t}$ and $\frac{v}{\delta t} \ll \epsilon^{-\frac{2}{3}}$,where naive perturbation is valid, the small ϵ limit of the metric is given by

$$\begin{aligned} \lim_{\epsilon \to 0} ds^2 &= 2dr\,dv - \left(r^2 - \frac{C_2(v)}{r}\right) dv^2 + r^2 dx_i^2 \\ &\neq 2dr\,dv - r^2 dv^2 + r^2 dx_i^2 \quad \text{for } r \lesssim \frac{\epsilon^{\frac{2}{3}}}{\delta t} \end{aligned} \quad (1)$$

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- In other words our spacetime is not uniformly well approximated by empty AdS even in the limit $\epsilon \to 0$.
- Consequently the naive perturbation expansion misidentifies the starting point for perturbation theory. This results in infrared divergences in this expansion.
- These divergences may be cured by perturbing around the correct leading order spacetime. As we have seen this spacetime is given at early times by

$$ds^2 = 2dr\,dv - \left(r^2 - \frac{C_2(v)}{r}\right) dv^2 + r^2 dx_i^2$$

  This follows from naive perturbation theory within its regime of validity.
- In fact the above metric is correct to leading order at all times as we now explain.

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Resummed perturbation Theory

- For $v > \delta t$ the solution is an unforced normalizable solution to the equations of motion. It turns out that the solution is completely determined by two pieces of initial data: mass density $M(\delta t) \approx C_2(\delta t)$ and dilaton function $\phi(r, \delta t)$.
- Naive expansion (valid at $v = \delta t$) determines both of these perturbatively in ϵ. It turns out while $C_2(\delta t) \sim \mathcal{O}(\epsilon^2)$, $\phi(r, \delta t) \sim \mathcal{O}(\epsilon^3)$.
- This turns out to imply that the subsequent solution ($v > \delta t$) is a small perturbation around static black brane of energy density $C_2(\delta t)$
- Consequently

$$\lim_{\epsilon \to 0} ds^2 = 2dr\, dv - \left(r^2 - \frac{C_2(v)}{r} \right) dv^2 + r^2 dx_i^2$$

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Summary

- To the leading order in ϵ the spacetime takes the Vaidya form for all v but for $r > \frac{\epsilon}{\delta t}$

$$ds^2 = 2dr\,dv - \left(r^2 - \frac{M(v)}{r}\right) dv^2 + r^2 dx_i^2$$

  where the mass function $M(v) = C_2(v) = -\int_{-\infty}^{v} dt\, \frac{\dot{\phi}_0 \dddot{\phi}_0}{2}$.

- At early times ($v \ll \frac{\delta t}{\epsilon^{\frac{2}{3}}}$) corrections to the Vaidya form are systematically captured by naive perturbation.
- Late time perturbation can be handled by a resummed perturbation expansion about the Vaidya spacetime. Corrections expressed in terms of universal functions that can be computed numerically. Resummed perturbation at leading order

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Time scales

- Note that the energy density M of the black brane is of $\mathcal{O}\left(\frac{\epsilon^2}{\delta t^3}\right)$ and so its temperature, $M^{\frac{1}{3}}$ is of $\mathcal{O}\left(\frac{\epsilon^{\frac{2}{3}}}{\delta t}\right)$. It follows that the black brane is formed over a time scale much smaller than the inverse temperature, the natural time scale associated with the brane.
- In particular in the limit $\epsilon \to 0$, $\delta t \to 0$, $\frac{\delta t}{\epsilon^{\frac{2}{3}}}$ held fixed describes the instantaneous formation of a black brane finite temperature. Spacetime for this formation process is simply given by AdS space for $v < 0$ and the blackbrane for $v > 0$. All corrections to this metric are suppressed by powers of $\epsilon^{\frac{2}{3}}$ and so vanish in this limit.

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Event horizon

- Recall that perturbation breaks down at $r \lesssim \frac{\epsilon}{\delta t}$. We will now show that this region is shielded from the boundary by an event horizon.
- The event horizon is the unique null manifold that asymptotes to $r = M^{\frac{1}{3}}$ at late times. Can be determined order by order in ϵ.
- It turns out that to the leading order the position of the event horizon

$$r_H(v) = M^{\frac{1}{3}} \quad v > 0$$

$$r_H(v) = \frac{M^{\frac{1}{3}}}{1 - M^{\frac{1}{3}} v} \quad v < 0$$

- Note $M^{\frac{1}{3}} \gg \frac{\epsilon}{\delta t}$. Consequently an arbitrarily small ϵ perturbation produces a black brane. singularity formation is always shielded behind an event horizon. Demonstration of cosmic censorship at small ϵ in translationally invariant AdS collapse.

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CFT dual

- The gravity solution describes a CFT in $R^{(d-1,1)}$
  - Initially in a vacuum state
  - During $(0, \delta t)$ perturbed by a translationally invariant source of amplitude ϵ .
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Rapid, scale dependent thermalization

To the leading order the spacetime reduces to that of a uniform blackbrane immediately after $v > \delta t$. It follows at leading order

- Our system responds to additional boundary perturbations at $v > \delta t$ as if it was precisely thermal.
- One point functions of local operators, which probe the spacetime only near the boundary, attain their thermal values almost instantaneously (for $v > \delta t$).
- Multipoint correlators of local operators and one point functions of non local operators like Wilson loops, which probe the spacetime away from the boundary, attain their thermal values over longer scale dependent time scales.

Subleading corrections in ϵ display slower quasinormal mode type decay over the linear response time scale $M^{-\frac{1}{3}} \gg \delta t$. Consequently all of the above is valid only at leading order in ϵ.

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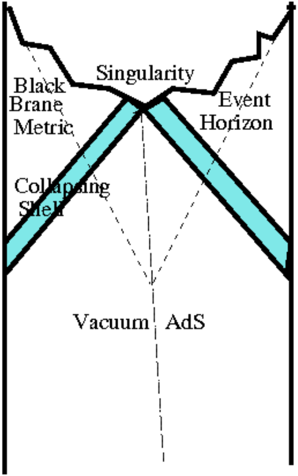

Forcing with slow spatial variation and fluid dynamics

- Let us generalize our discussion to allow for a forcing function that breaks the translational invariance but on a length scale $L \gg$ the thermal scale $M^{-\frac{1}{3}}$.
- At leading order we expect the spacetime to be tubewise well approximated by the Vaidya form with a spatially varying temperature.

$$ds^2 = 2dr\,dv - \left(r^2 - \frac{M(v,\vec{x})}{r}\right) dv^2 + r^2 dx_i^2$$

where

$$M(v,\vec{x}) = C_2(v,\vec{x}) + \mathcal{O}(\epsilon^4)$$

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- Subsequent evolution governed by the equations of boundary fluid dynamics.
- Note that the fluid dynamics is valid already for $v > \delta t \ll M^{-\frac{1}{3}}$. This is true even though correlation functions over length scales of the mean free path ($M^{-\frac{1}{3}}$) are far from thermal at these times. Consequently fluid dynamics is the precise dynamical description of our system well before it has locally equilibriated on the scale of mean free path.

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Spherically symmetric collapse in flat space

- We study the collapse of a spherically symmetric null shell, propagating inwards from $\mathcal{I}^-$ in an asymptotically flat space.
- Near $\mathcal{I}^-$ the shell takes the form

$$\lim_{r\to\infty}\phi(r,v)=\frac{\psi(v)}{r}$$

$$\psi(v)=0,\quad (v<0)$$

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- The shell is qualitatively characterized by thickness δt and Schwarzschild radius $r_H\sim\epsilon_f^2\delta t$.

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- Subleading corrections to the Vaidya metric back scatter a fraction of the incident wave to $\mathcal{I}^+$.
- As an application of our perturbative procedure we compute the fraction of the incident energy that is back scattered as a functional of the shape of the incident wave packet $\psi(v)$ to leading non trivial order in $\frac{1}{\epsilon_f^2}$.
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- Non normalizable component of dilaton is turned on uniformly over the boundary sphere of radius R and time interval δt.
- Qualitatively two parameters: Amplitude = ϵ and $x \equiv \frac{\delta t}{R}$.
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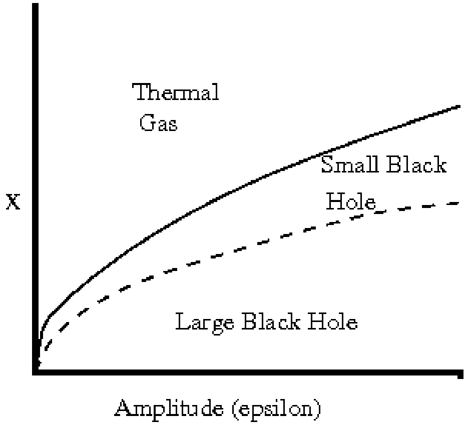

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- The gravitational solution is dual to a CFT on a sphere, initially in its vacuum state.
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- For odd $d > 3$ our results are qualitatively similar to those of $d = 3$. However collapse processes in even d are different, in some respects, from their odd d counterparts.

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- While in even bulk dimensions a massless field propagates along its light cone, in odd bulk dimension it spreads inside the light cone,
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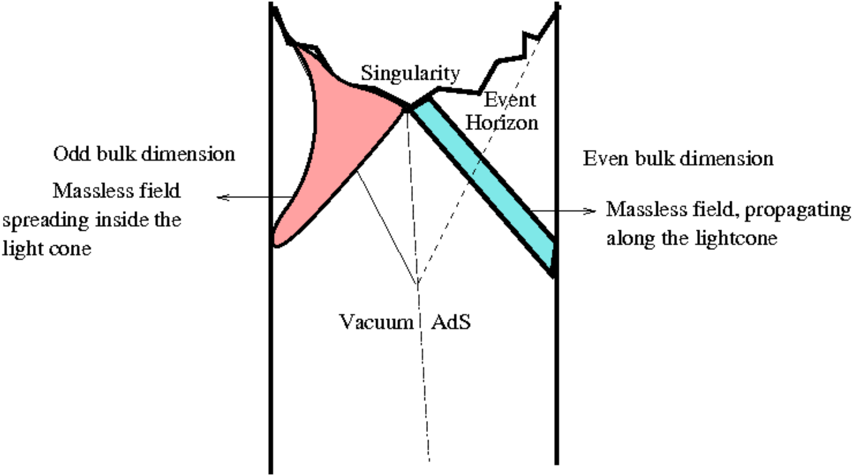

Explicit results upto fifth order

$$\phi_3(r,v) = \frac{1}{4r^3}\int_{-\infty}^{v} B(x)\,dx$$

$$f_4(r,v) = \frac{\dot{\phi}_0}{384r^3}\left\{\dot{\phi}_0^3 - 12\int_{-\infty}^{v} B(x)\,dx\right\}$$

$$g_4(r,v) = -\frac{C_4(v)}{r} + \frac{\dot{\phi}_0}{24r^2}\left\{-\dot{\phi}_0^3 + 3\int_{-\infty}^{v} B(x)\,dx\right\} + \frac{1}{48r^3}\left(3B(v)\dot{\phi}_0 - 4\dot{\phi}_0^3\ddot{\phi}_0 + 3\ddot{\phi}_0\int_{v}^{\infty} B(t)dt\right)$$

Where

$$C_4(v) = \int_{-\infty}^{v} dt\,\frac{3\dot{\phi}_0}{8}\left(\dot{\phi}_0^3 - \int_{-\infty}^{t} B(x)\,dx\right)$$

Explicit results upto fifth order

$$\phi_5(r,v) = \frac{1}{8r^5}\int_{-\infty}^{v} B_1(x)\,dx + \frac{1}{6r^4}\int_{-\infty}^{v} B_3(x)\,dx + \frac{5}{24r^4}\int_{-\infty}^{v} dy \int_{-\infty}^{y} B_1(x)\,dx + \frac{1}{4r^3}\int_{-\infty}^{v} B_2(x)\,dx + \frac{1}{6r^3}\int_{-\infty}^{v} dy \int_{-\infty}^{y} B_3(x)\,dx + \frac{5}{24r^3}\int_{-\infty}^{v} dz \int_{-\infty}^{z} dy \int_{-\infty}^{y} B_1(x)\,dx$$

where

$$B(v) = \dot{\phi}_0\left[-C_2(v) + \dot{\phi}_0\ddot{\phi}_0\right]$$

$$B_1(v) = \left(-\frac{9}{4}C_2(v) + \frac{7}{8}\dot{\phi}_0\ddot{\phi}_0\right)\int_{-\infty}^{v} B(x)\,dx + \frac{1}{2}C_2(v)\dot{\phi}_0^3 + \frac{3}{8}\dot{\phi}_0^2 B(v) - \frac{1}{6}\dot{\phi}_0^4\ddot{\phi}_0$$

$$B_2(v) = C_4(v)\dot{\phi}_0$$

$$B_3(v) = \frac{1}{24}\left(-30\dot{\phi}_0^2\int_{-\infty}^{v} B(x)\,dx + 7\dot{\phi}_0^5\right)$$

Resummed perturbation expansion

- Explicitly to the leading order for $v > \delta t$:

$$\phi = \frac{\phi_3^0(\delta t)}{M} \psi \left(\frac{r}{M^{\frac{1}{3}}}, (v - \delta t) M^{\frac{1}{3}} \right)$$

where

$$\phi_3^0(\delta t) = \frac{1}{4} \int_{-\infty}^{\delta t} \dot{\phi}_0 \left[-C_2(t) + \dot{\phi}_0 \ddot{\phi}_0 \right] \; dt$$

$$M = \int_{-\infty}^{\delta t} dt \; \frac{\ddot{\phi}_0^2}{2}$$

- $\psi(x, y)$. satisfies

$$\partial_x \left(x^4 \left(1 - \frac{1}{x^3} \right) \partial_x \psi \right) + 2x \partial_y \partial_x (x\psi) = 0$$

- Boundary condition: $\psi \sim \mathcal{O}(\frac{1}{x^3})$ at large x
- The initial condition: $\psi(x, 0) = \frac{1}{x^3}$.

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$$\phi_3^0(\delta t) = \frac{1}{4}\int_{-\infty}^{\delta t} \dot{\phi}_0 \left[-C_2(t) + \dot{\phi}_0\ddot{\phi}_0\right] \; dt$$

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- $\psi(x, y)$. satisfies

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Resummed perturbation expansion

- Explicitly to the leading order for $v > \delta t$:

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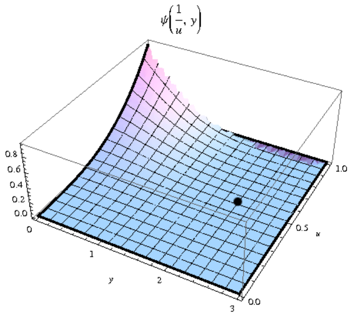

- This plot is obtained using Mathematica 6.
- Note exponential decay in time (quasinormal mode behavior).
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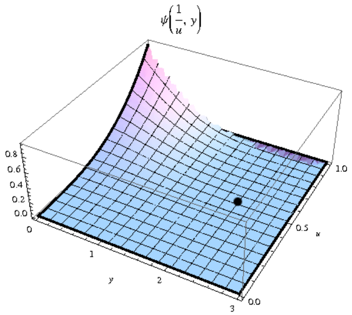

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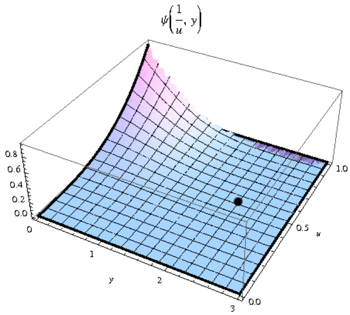

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Stability

- Large uncharged black holes: usually stable in AdS/CFT .
- Thermal gas phase is also stable.
  Its energy $\ll$ Critical energy density for Jean's instability.
- Small black hole phase: usually unstable in AdS/CFT.
- Here it is a two step thermalization:
  - Small black holes at $v \sim \delta t$.
  - Gregory-Lafflamme instability at $v \sim R^2 r_H \gg \delta t$.

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Choptuik phenomena

- Black hole and thermal gas phases are separated by critical surface $x^2 \sim \epsilon$
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