Title: Following the state of an evaporating charged black hole into the quantum gravity regime

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Abstract:

In this talk I will discuss the energy probability density function of an evaporating near-extremal charged black hole. At sufficiently low energies, such black holes experience large quantum metric fluctuations in the AdS_{2} throat which are governed by a Schwarzian action. These fluctuations modify Hawking evaporation rates, and therefore also affect how the black hole state evolves over time. In previous work on Schwarzian-corrected Hawking radiation, the black hole was taken to be in the microcanonical or canonical ensemble [arXiv:2411.03447]. However, we find that an initially fixed-energy or fixed-temperature state does not remain so in the regime where Schwarzian corrections are important. We consider three decay channels: the emission of massless scalars, photons, and entangled pairs of photons in angular momentum singlet states. In each of the three cases, we find that in the very low energy, quantum dominated regime, the probability distribution of the black hole energy level occupation tends toward a particular attractor function that effectively depends on only one combination of time and energy. This function is independent of the initial state and gives new predictions for the energy fluxes and Hawking emission spectra of near-extremal charged black holes.

Following the state of an evaporating charged black hole into the quantum gravity regime

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March 13, 2025

Based on arXiv:2503.02051

The semiclassical treatment of black hole thermodynamics breaks down in the extremal limit.[1]

This is because the near-horizon metric has soft modes whose action becomes unsuppressed at low temperatures.

For a RN black hole, quantum metric fluctuations begin to modify the semiclassical description at the scale ($E \equiv M - Q$)

$$E_{\text{breakdown}} \equiv \frac{\pi}{r_+ S_0}$$

Today: how these quantum gravity effects correct observables of near-extremal charged black holes in 4d flat space

[1] Preskill et al. 1991; Maldacena, Michelson, and Strominger 1999.

Outline

Review: Quantum gravity corrected (microcanonical) emission rates

1. Breakdown of semiclassical treatment for near-extremal charged BH $\rightarrow$ the role of JT gravity
2. Low-energy effective theory for BH scattering
3. Example: the quantum gravity corrected (microcanonical) emission rate of neutral massless scalar particles

These quantum metric fluctuations modify how the black hole evolves over time.

Quantum gravity corrected time evolution of the black hole state

1. What is the probability density function of black hole energy, under...
   1.1 scalar emission?
   1.2 photon emission?
   1.3 di-photon emission?
2. How does this change the Hawking emission spectra?

A puzzle from the 90s

This story was reviewed by A.W. in their talk this morning.

The semiclassical analysis of RN black holes breaks down at sufficiently low energies. We can see this by thinking about Hawking radiation.

Letting $E \equiv M - Q$, the energy temperature relation for $E \ll 1$ is

$$E \sim Q^3 T_H^2 \qquad E \ll 1$$

The energy of a typical Hawking quanta is $\sim T_H$. When $T_H \sim E$, the emission of one quanta changes the temperature by a substantial amount. This occurs at $E_{\text{brk}} \sim 1/Q^3 \sim 1/(r_+ S_0)$.

This breakdown is due to modes of the metric exhibiting large quantum variance at low energies. They must be treated quantum mechanically.

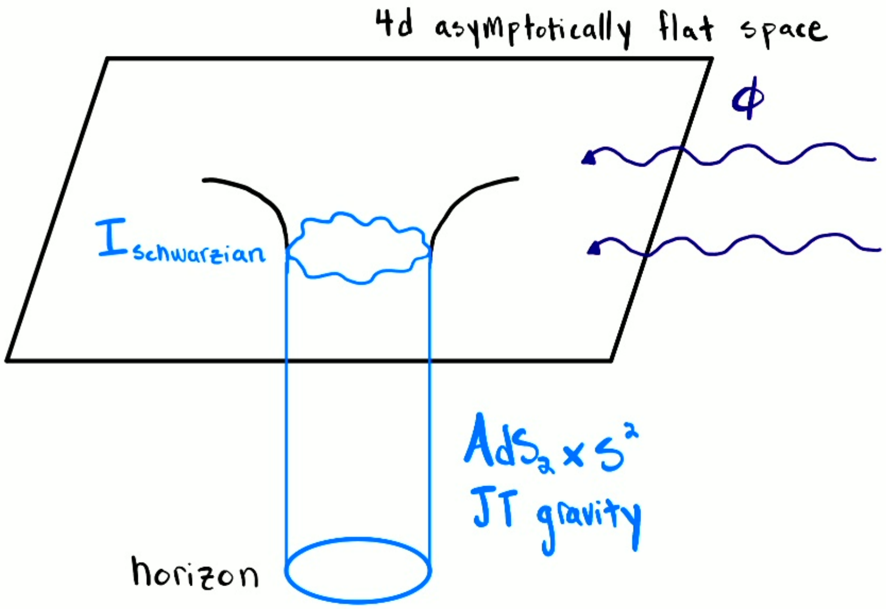
Within the last decade, it was understood how to do so in JT gravity.

Resolution in JT gravity

As the black hole approaches extremality, the near-horizon region develops a throat-like spatial geometry with the approximate metric $AdS_2 \times S^2$.

Performing a dimensional reduction to AdS_2 in the throat produces a theory of 2d gravity coupled to matter (JT gravity).

The only dynamics occur on the boundary. The boundary action evaluates to a Schwarzian time derivative.



Resolution in JT gravity

The gravitational path integral over this Schwarzian action can be performed exactly $\rightarrow$ used to calculate the modified BH thermodynamics & density of states below E_{brk}[2].

For JT + matter, correlation functions of matter operators dual to AdS_2 fields have also been computed exactly.[3]

As we will now explain, from these correlators we can extract Schwarzian corrections to observables such as the 4d absorption cross section and Hawking emission rates.[4]

[2] Iliesiu and Turiaci 2021; Ghosh, Maxfield, and Turiaci 2020; Iliesiu, Murthy, and Turiaci 2022; Kapec et al. 2024; Rakic, Rangamani, and Turiaci 2024; Kolanowski et al. 2024; Moitra et al. 2019; Modak, Singh, and Panda 2025; Boruch et al. 2022; Heydeman et al. 2022.

[3] Mertens, Turiaci, and Verlinde 2017; Mertens 2018; Lam et al. 2018; Blommaert, Mertens, and Verschelde 2018; Iliesiu, Pufu, et al. 2019; Kitaev and Suh 2019; Yang 2019; Suh 2020.

[4] Brown et al. 2024; Maulik, Meng, and Pando Zayas 2025; Emparan 2025; Bai and Korwar 2023.

Low energy effective theory

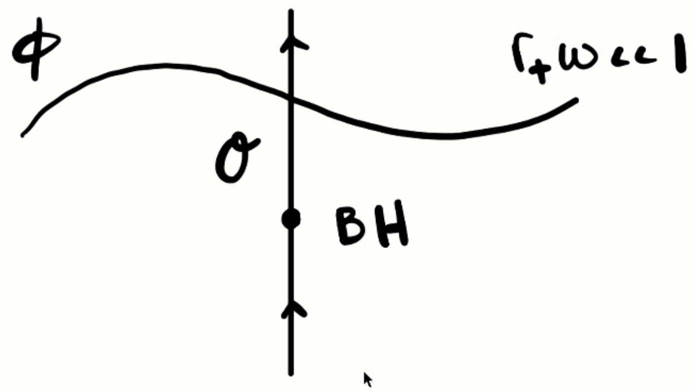
For illustration: neutral, massless scalar field coupled to a near-extremal RN black hole.

The idea: low-energy ($r_+\omega \ll 1$) interactions of the BH with the probe scalar are captured by an effective theory which replaces the BH by a quantum system living at a point in Minkowski spacetime.

The coupling between ϕ and the BH is described by

$$H_{\text{int}} = gO(t)\phi(t,\vec{0}) \qquad\qquad r_+\omega \ll 1$$

O is an operator living on the point particle worldline which acts on the BH Hilbert space. In the language of AdS/CFT, it is the $\Delta = 1$ primary operator dual to the massless scalar in AdS_2.



Low energy effective theory

H_{int} is weakly coupled. One way to see this is $[O] = 1$ and $[\phi] = 1$, so $[g] = -1$.

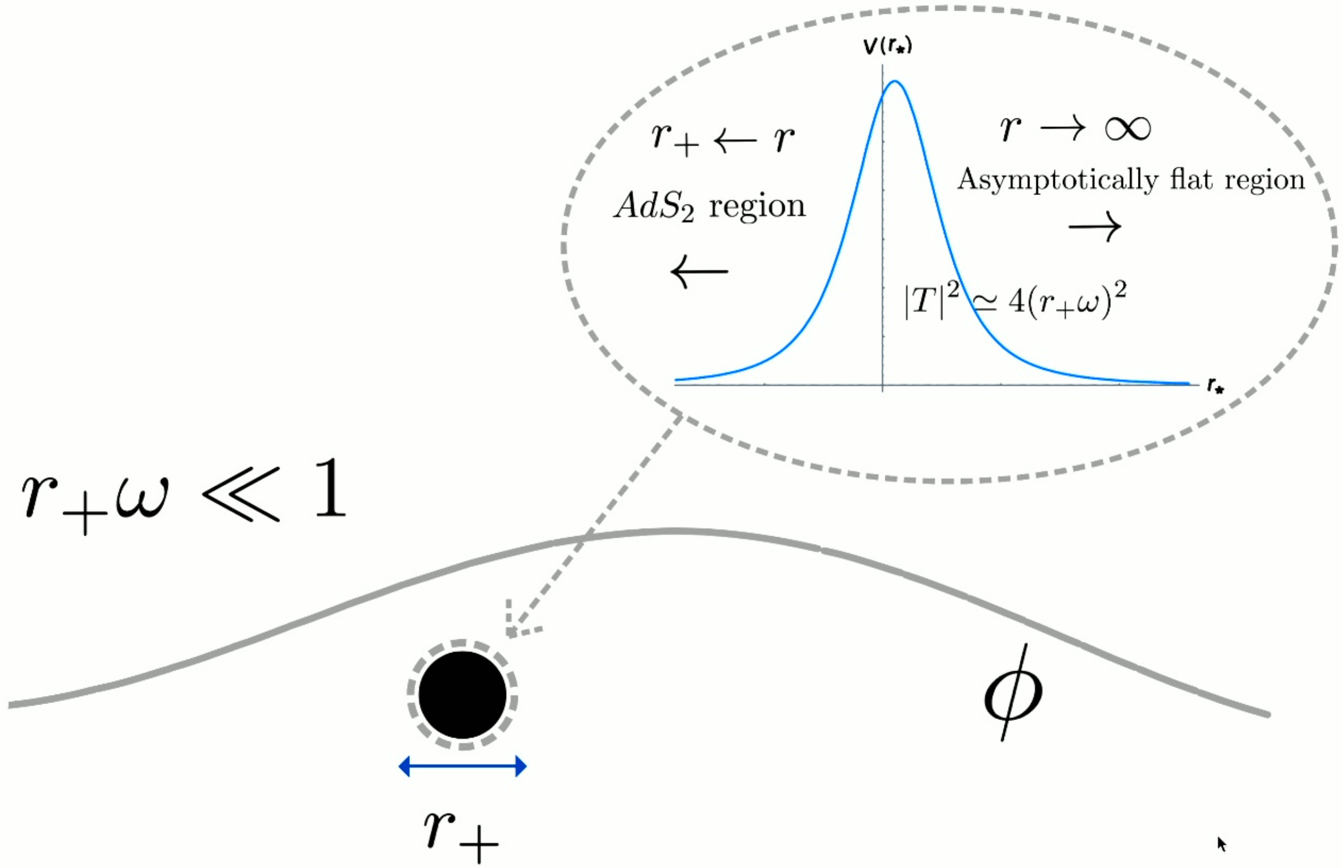
Another way: ϕ feels a potential barrier $V(r)$ separating the AdS_2 region from the asymptotically flat region. Transmission probability $|T|^2 \propto (r_+\omega)^2$ is small at low frequencies. (See next slide)

$\langle OO \rangle$ has been computed in JT gravity including the gravitational path integral over the Schwarzian mode. Once we have expressed an observable in terms of $\langle OO \rangle$, those previous results immediately tell us how it is modified by the Schwarzian, simply by expanding $\langle OO \rangle$ at low energies $E \ll E_{\text{brk}}$.

We determine g by a matching computation (to, e.g. the emission rate or absorption cross section). In this case, $g = 2r_+$.

Low energy effective theory

$$r_+\omega \ll 1$$



Example: spontaneous emission rate of scalars

$\mathcal{H} = \mathcal{H}_{\rm BH} \otimes \mathcal{H}_{\rm matter}$. $|i\rangle = |\psi_i, 0\rangle$ and $|f\rangle = |\psi_f, \vec{q}\rangle$, where $\vec{q}$ is the 3-momentum of the emitted particle and $|\psi_{i,f}\rangle$ are the initial and final states of the black hole. The amplitude for transition from $|i\rangle$ to $|f\rangle$ is

$$A_{i\to f} = -ig\frac{1}{\sqrt{2|\vec{q}|}}\int_0^T dt \langle\psi_f|O(t)|\psi_i\rangle e^{i|\vec{q}|t}$$

The total emission probability is

$$\Gamma(T) \equiv \sum_{E_i} p(E_i) \sum_{|\psi_f\rangle} \int \frac{d^3q}{(2\pi)^3} |A_{i\to f}|^2 = g^2 T \int_0^\infty \frac{d\omega}{2\pi}\frac{\omega^2}{\pi}\frac{1}{2\omega}\int_{-T}^{T} dt\ e^{-i\omega t}\langle O(t)O(0)\rangle$$

where we relabelled $|\vec{q}| \to \omega$. The number of particles emitted per unit time is

$$\frac{dN}{dt} = \lim_{T\to\infty}\frac{1}{T}\Gamma(T) = g^2\int_0^\infty \frac{d\omega}{2\pi}\frac{\omega}{2\pi}\int_{-\infty}^{\infty} dt e^{-i\omega t}\langle O(t)O(0)\rangle$$

The meaning of $\langle\rangle$ depends on the initial state of the black hole.

Schwarzian corrections

We expressed the emission rate in terms of $\langle OO \rangle$. In the semiclassical regime, this is a conformal two point function.

The calculation of $\langle OO \rangle$ at energies $E \ll E_{\text{brk}}$ involves integrating over the Schwarzian mode in the JT gravity partition function. For a black hole with $J = 0$ and fixed charge Q, the microcanonical result is

$$\int dt e^{-i\omega t} \langle E|O(t)O(0)|E\rangle = 2\pi\rho(E-\omega)|O_{E,E-\omega}|^2$$

$$\rho(E) = \frac{1}{2\pi^2 E_{\text{brk}}} e^{S_0} \sinh(2\pi\sqrt{2E_{\text{brk}}^{-1}E})\Theta(E)$$

$$|O_{E_1,E_2}|^2 = \frac{2e^{-S_0}\Gamma\left(\Delta \pm i\sqrt{2E_{\text{brk}}^{-1}E_1} \pm i\sqrt{2E_{\text{brk}}^{-1}E_2}\right)}{(2E_{\text{brk}}^{-1})^{2\Delta}\Gamma(2\Delta)}$$

Schwarzian corrections

The Schwarzian-corrected scalar emission rate is then

$$\left.\frac{dN}{dt}\right|_E = \frac{2r_+^2}{\pi}\int_0^\infty d\omega\,\omega\,\rho(E-\omega)|O_{E,E-\omega}|^2$$

In particular, stripping off the integral over final states, the microcanonical transition rate from a state $|E\rangle$ to $|E-\omega\rangle$ is

$$\gamma(E, E-\omega) = \frac{2r_+^2}{\pi}\omega|O_{E,E-\omega}|^2$$

Semiclassical result:

$$\left.\frac{dN}{dt}\right|_E \propto r_+^2\,(EE_{\text{brk}})^{3/2} \qquad E \gg E_{\text{brk}}$$

Quantum result:

$$\left.\frac{dN}{dt}\right|_E \propto r_+^2\sqrt{E_{\text{brk}}}E^{5/2} \propto r_+^2(EE_{\text{brk}})^{3/2}(E/E_{\text{brk}}) \qquad E \ll E_{\text{brk}}$$

Energy probability distribution

We have just calculated the scalar emission rate from a fixed energy state, i.e.

$$P(E, t=0) = \delta(E - E_0)$$

where $P(E, t)$ is the differential probability of energy level occupation at energy E and time t. These and other microcanonical rates were computed previously.[5]

Q: How does $P(E, t)$ evolve over time, as the BH undergoes Schwarzian-corrected Hawking radiation?

To answer this question, we need to solve

$$\frac{dP(E,t)}{dt} = -\int_0^E dE' \rho(E') \gamma(E, E') P(E, t) + \int_E^\infty dE'' \rho(E) \gamma(E'', E) P(E'', t)$$

[5] Brown et al. 2024.

Energy probability distribution

In the semiclassical limit $E \gg E_{\text{brk}}$ and $\omega \ll E$, the equation is solved by a time-dependent thermal state,

$$P_{\text{th}}(E,t) = Z(\beta(t))^{-1}\rho(E)e^{-\beta(t)E}$$

where $\dot{\beta} = \frac{\partial\beta}{\partial E}\frac{d\langle E\rangle}{dt}$ and $\frac{d\langle E\rangle}{dt}$ is the usual semiclassical expression for the total energy flux,

$$\frac{d\langle E\rangle}{dt} = -\mathcal{N}_{\ell,m}\int_0^\infty \frac{d\omega}{2\pi}\omega\frac{\mathcal{P}(\omega,\ell)}{e^{\beta\omega}-1}$$

$P_{\text{th}}(E,t)$ stops being a solution at energies $E \lesssim E_{\text{brk}}$. To study the solution below E_{brk}, let us continue with the example of scalar emission.

In the low energy limit $E \ll E_{\text{brk}}$ the operator matrix elements become constant, and the transition rate can be approximated as

scalar emission $\qquad \gamma(E,E') = \frac{1}{\pi}E_{\text{brk}}^2 r_+^2 e^{-S_0}(E-E') \qquad E \ll E_{\text{brk}}$

$P(E,t)$ solution - scalar emission

Now we have

$$\frac{1}{c_1}\frac{dP(E,t)}{dt} = -\int_0^E dE'\sqrt{E}(E-E')P(E,t) + \int_E^\infty dE''\sqrt{E''}(E''-E)P(E'',t)$$

where $\quad c_1 \equiv \frac{\sqrt{2E_{\rm brk}}r_+^2}{\pi^2}$

This can be solved analytically. Subject to the initial condition

$$P(E,t=0) = \delta(E-E_0)$$

for some $E_0 \ll E_{\rm brk}$, the solution for $t>0$ and $E<E_0$ is

$$P(E,t) = e^{-\mathcal{E}_0\tau}\delta(E-E_0) + \frac{3}{2}\frac{1}{E}\left(e^{-\mathcal{E}\tau} - e^{-\mathcal{E}_0\tau}\left(\frac{\mathcal{E}}{\mathcal{E}_0}\right)^{\frac{3}{5}} - \frac{3}{5}(\mathcal{E}\tau)^{\frac{3}{5}}\left[\Gamma\left(-\frac{3}{5},\mathcal{E}\tau\right) - \Gamma\left(-\frac{3}{5},\mathcal{E}_0\tau\right)\right]\right)$$

for $E<E_0$, $\quad$ where $\quad \mathcal{E}\equiv E^{5/2}, \quad \mathcal{E}_0 \equiv E_0^{5/2}, \quad \tau \equiv \frac{4}{15}c_1 t$

Attractor solution

Long time behavior? Consider the limit $\mathcal{E}_0\tau \to \infty$ with $\mathcal{E}\tau$ fixed. All the terms which depend on E_0 are exponentially suppressed, and $P(E,t)$ reduces to the simple function

$$P(E,t) \to \bar{P}(E,t) \equiv \frac{3}{2}\frac{1}{E}(\mathcal{E}\tau)^{3/5}\,\Gamma\left(\frac{2}{5},\mathcal{E}\tau\right) \qquad \text{as} \qquad \mathcal{E}_0\tau \to \infty$$

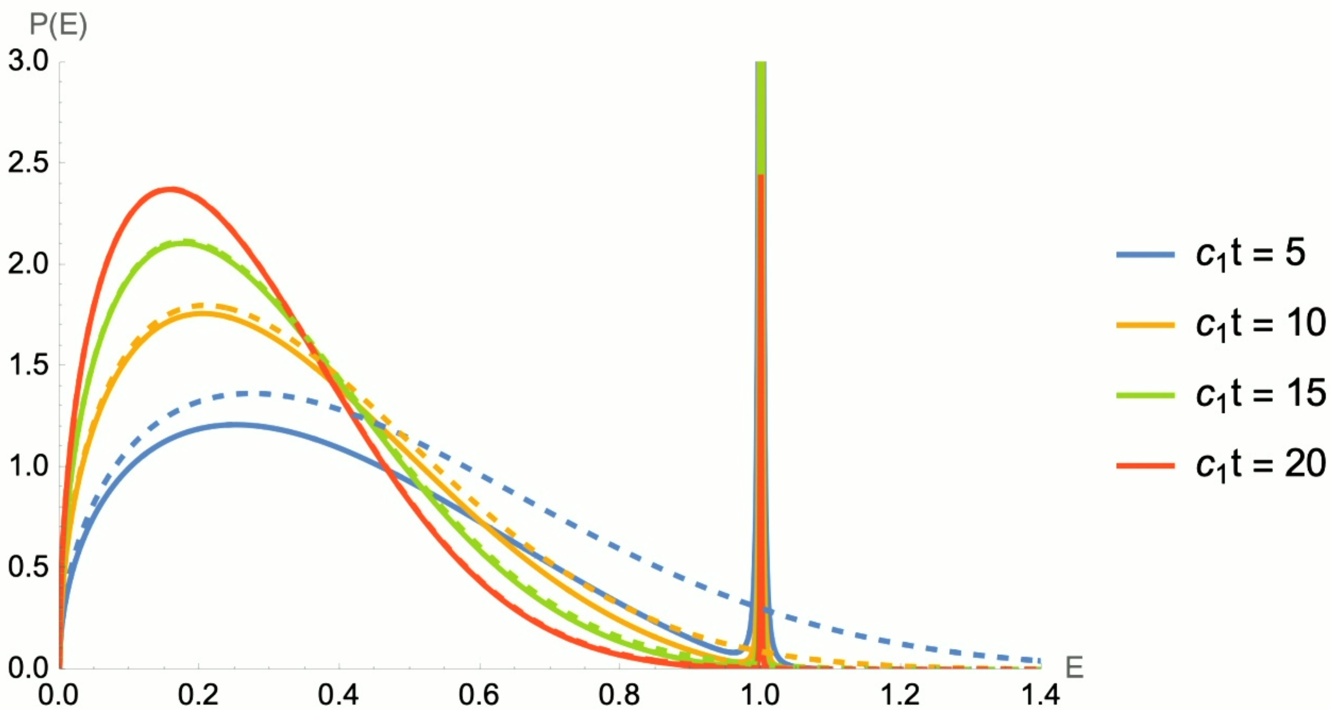
The time evolution of any other initial distribution $Q(E,t=0) = P_0(E)$ localized around some E_0 is given by integrating the delta function solution.

$$Q(E,t=0) = \int_0^\infty dE'\delta(E-E')P_0(E')$$

$$\Rightarrow \quad Q(E,t) \to \int_0^\infty dE' P_0(E')\bar{P}(E,t) = \bar{P}(E,t) \quad \text{as} \quad \mathcal{E}_0\tau \to \infty$$

This suggests we have an attractor solution given by $\bar{P}(E,t)$. We verify this by solving the equation numerically.

$P(E, t)$ for a black hole undergoing scalar emission below $E_{\rm brk}$



Solid lines: the numerical solution starting from a delta function distribution at $E_0 = E_{\rm brk}$. Dashed lines: the attractor solution $\bar{P}(E)$.

A scaling symmetry of $E\bar{P}$

$\bar{P}(E,t)$ exhibits a kind of scaling symmetry which fixes the time dependence of $\langle E(t)\rangle$, and likewise the dependence of $\langle dE/dt\rangle$ on $\langle E\rangle$.

$E\bar{P}$ depends only on the combination $E^{5/2}t$, so it is invariant under the rescaling

$$t \to \eta t \qquad\qquad E \to \eta^{-\frac{2}{5}} E$$

At long times, the expected energy is

$$\langle E(t)\rangle = \int_0^\infty dE E\bar{P}(E,t)$$
$$\propto \frac{1}{t^{\frac{2}{5}}}\int_0^\infty \frac{dz}{z^{\frac{3}{5}}} f(z) \qquad z \equiv E^{5/2}t$$

where $f(z)$ is some function of z. Likewise, this fixes

$$\langle \frac{dE}{dt}\rangle \propto \langle E\rangle^{7/2}$$

18 / 41

Evaporation history of a charged BH coupled to the standard model

Our universe contains no known massless scalars.

Recently Iliesiu, Penington, Usatyuk, and Brown studied the evaporation history of a charged BH coupled to the standard model.[6] See Penington's Strings 2025 talk.

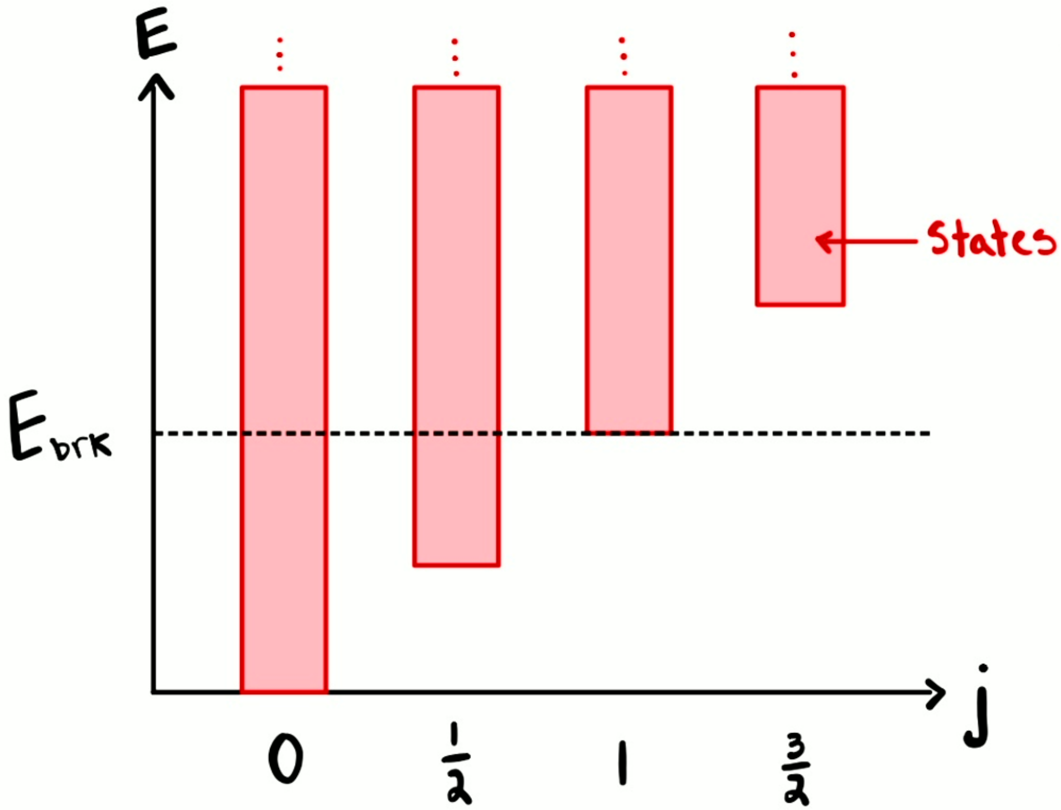
If the BH has $Q_{\text{initial}} > 1.8 \times 10^{44} q$, Schwinger pair production is exponentially suppressed, and the BH will lose energy faster than charge, driving it toward extremality.

It will spend the majority of its lifetime at energies $E \lesssim E_{\text{brk}}$.

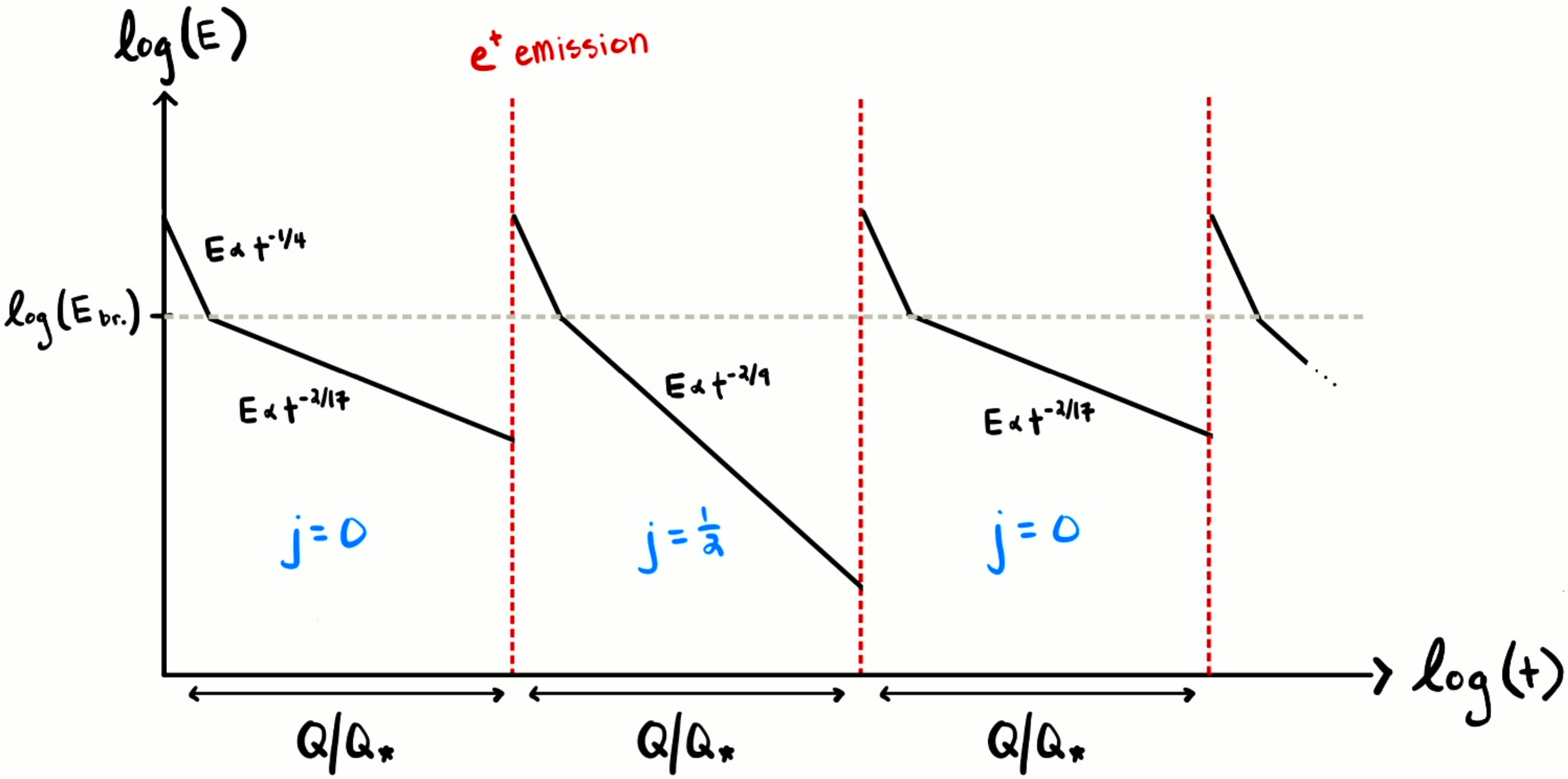
In this regime, the black hole alternates between two dominant radiation channels: emission of single photons from a black hole with angular momentum $j = 1/2$ or emission of entangled pairs of photons with zero net angular momentum ("di-photons") from a black hole with $j = 0$.

[6] Brown et al. 2024.

Evaporation history of a charged BH coupled to the standard model



2411.03447 Brown Iliesiu Penington Usatyuk '24



This motivates us to study the evolution of the black hole state under these two radiation channels in the Schwarzian regime.

Transition rates for photon, di-photon emission

Iliesiu et al. computed the Schwarzian-corrected transition rates in these channels. In the low energy limit $E \ll E_{\text{brk}}$,

$$\ell = 1 \text{ photon emission, } j = 1/2 \text{ BH} \qquad \gamma(E, E') = \frac{1}{9\pi} E_{\text{brk}}^6 r_+^8 e^{-S_0} (E - E')^3$$

$$\text{di-photon emission, } j = 0 \text{ BH} \qquad \gamma(E, E') = (8.2 \times 10^{-4}) \frac{640}{189\pi^3} E_{\text{brk}}^{10} r_+^{16} e^{-S_0} (E - E')^7$$

The primary difference between the three cases is in how γ scales with the energy of the emitted mode. The powers are

$a = 1$	scalar emission
$a = 3$	photon emission
$a = 7$	di-photon emission

For $a = 3, 7$ we also find attractor solutions with scaling symmetries. $E\bar{P}$ **depends only on** $E^{\frac{3}{2}+a} t$**.**

Photon emission attractor solution

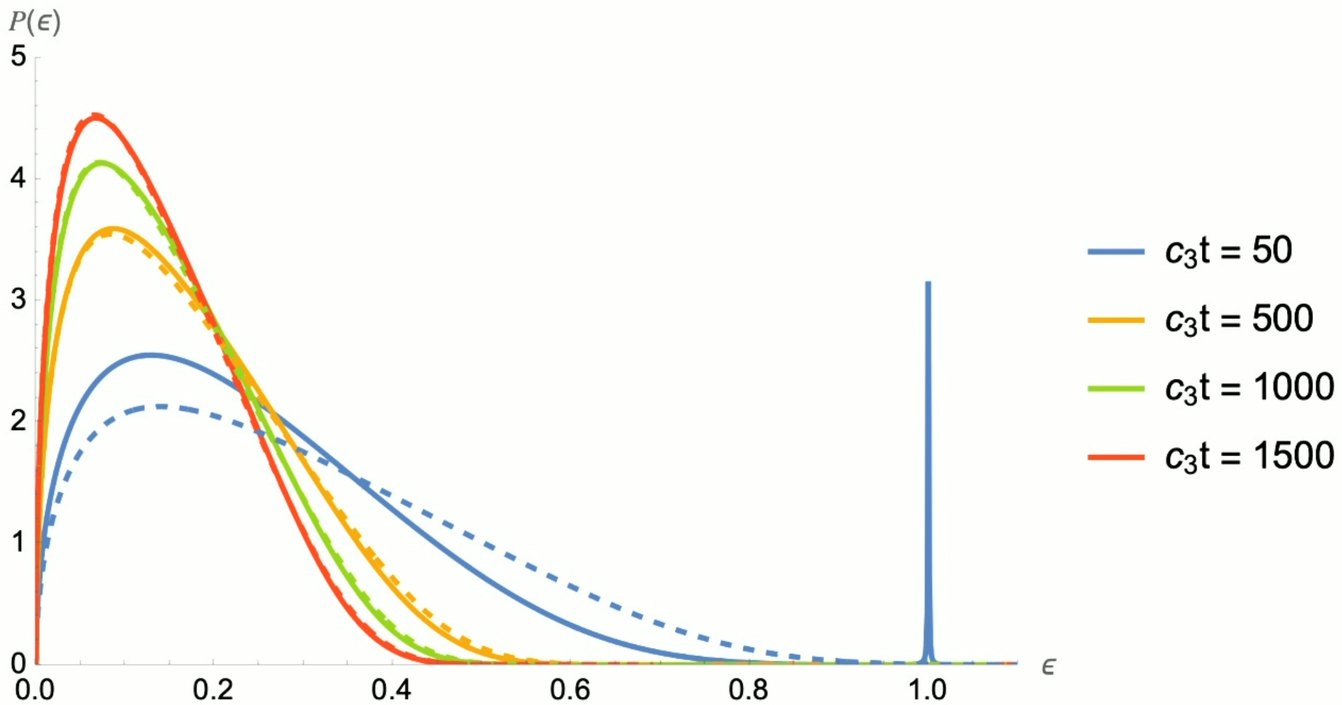
For BH with angular momentum $j = 1/2$, the energy spectrum starts at $M = Q + \frac{3}{8}E_{\text{brk}}$.
We denote $\varepsilon \equiv E - \frac{3}{8}E_{\text{brk}}$.

$$\begin{aligned}\tilde{P}(z) =& v_0 z^{1/3} + v_1 z \, {}_3F_3\left(\frac{2}{3}, \frac{4}{3} - \frac{i\sqrt{26}}{9}, \frac{4}{3} + \frac{i\sqrt{26}}{9}; \frac{11}{9}, \frac{13}{9}, \frac{5}{3}; -z\right) \\ &+ v_2 z^{5/9} \, {}_3F_3\left(\frac{2}{9}, \frac{8}{9} - \frac{i\sqrt{26}}{9}, \frac{8}{9} + \frac{i\sqrt{26}}{9}; \frac{5}{9}, \frac{7}{9}, \frac{11}{9}; -z\right) \\ &+ v_3 z^{7/9} \, {}_3F_3\left(\frac{4}{9}, \frac{10}{9} - \frac{i\sqrt{26}}{9}, \frac{10}{9} + \frac{i\sqrt{26}}{9}; \frac{7}{9}, \frac{11}{9}, \frac{13}{9}; -z\right)\end{aligned}$$

$$\tilde{\varepsilon} = \varepsilon^{\frac{9}{2}} \qquad \tau = \frac{32}{315} c_3 t \qquad z = \tilde{\varepsilon}\tau \qquad \tilde{P} \equiv \varepsilon P$$

The v_i are constant, order one coefficients. $c_3 = \frac{2\sqrt{2}}{9\pi^2} r_+^8 E_{\text{brk}}^{9/2}$.

$P(\varepsilon, t)$ for a $j = 1/2$ black hole undergoing single photon emission below E_{brk}



Solid lines: numerical solution starting from $\delta(\varepsilon - E_{\text{brk}})$. Dashed lines: attractor solution.

Timescale for approach to the attractor

The relevant timescale for the approach to the attractor is about two orders of magnitude larger than in the case of scalar emission.

This is because the probability density is most naturally a function of $z \propto E^{3/2+a}t$, and has support when z is order one.

For the scalar, $z \propto E^{5/2}t$, so for E to decrease by a factor of 10, t must increase by a factor of $10^{5/2} \approx 3 \times 10^2$.

For photon emission, $z \propto E^{9/2}t$, so t must increase by a factor of $10^{9/2} \approx 3 \times 10^4$.

Because the distribution is evolving more slowly, it also takes longer for the numerical solution to converge to the attractor.

Di-photon emission attractor solution

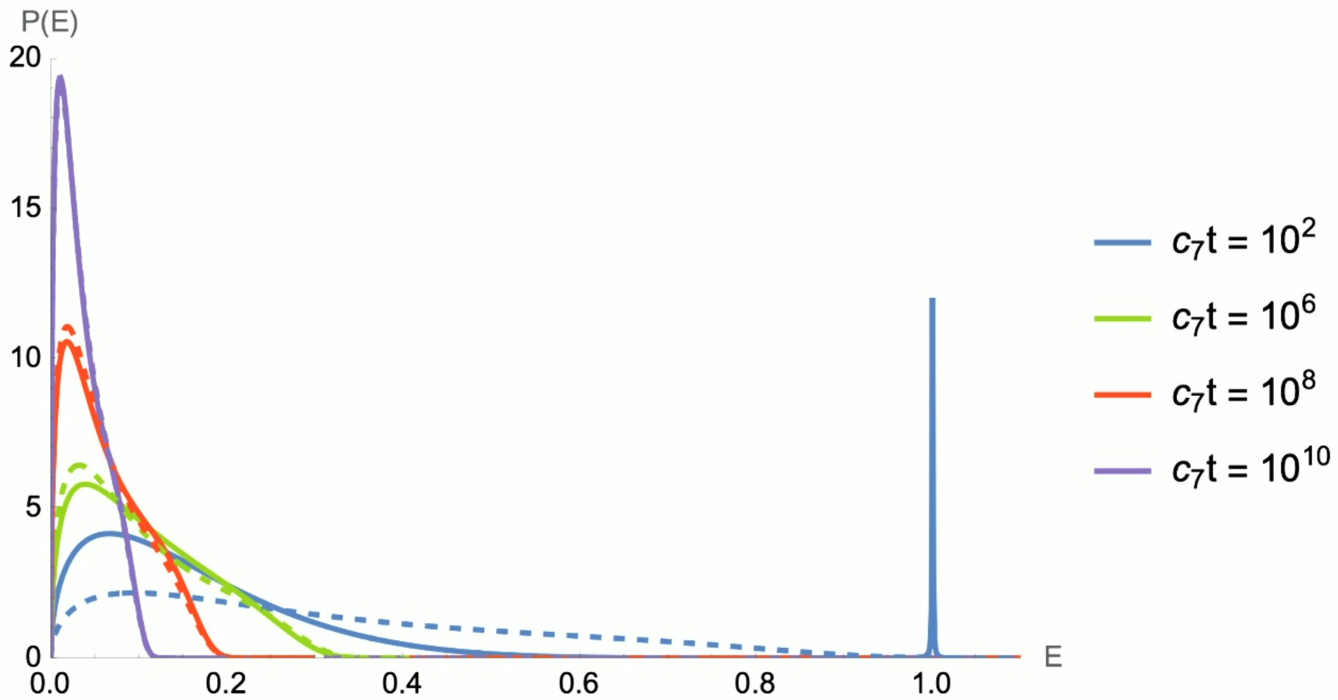
The solution is again a linear combination of generalized hypergeometric functions (parameters are suppressed)

$$\tilde{P}(z) = v_0 z^{3/17} + v_1 z^{5/17}\ {}_7F_7(z) + v_2 z^{7/17}\ {}_7F_7(z) + ... + v_7 z\ {}_7F_7(z)$$

$$\mathcal{E} \equiv E^{\frac{17}{2}} \qquad \tau = \frac{4096}{109395} c_7 t \qquad z = \mathcal{E}\tau \qquad \tilde{P} \equiv EP$$

The v_i are constant, order one coefficients. $c_7 \equiv (8.2 \times 10^{-4}) \times \frac{640\sqrt{2}}{189\pi^4} E_{\mathrm{brk}}^{17/2} r_+^{16}$.

$P(E,t)$ for a $j=0$ black hole undergoing di-photon emission below $E_{\rm brk}$



Solid lines: numerical solution starting from $\delta(E-E_{\rm brk})$. Dashed lines: attractor solution.

Corrected spectrum and emission rates

In all cases, the scaling symmetry fixes the time dependence of $\langle E(t) \rangle$ to be

$$\langle E(t) \rangle \propto \frac{1}{t^{\frac{2}{3+2a}}}$$

These are the same powers found by computing the energy flux dE/dt in a microcanonical state and integrating with respect to time.[7] However, the overall coefficients are different.

We will now discuss the results for $\langle E(t) \rangle$, $\langle dE/dt \rangle$, and the Hawking emission spectra $\frac{dE}{dtd\omega}$ in the various attractor states.

[7] Brown et al. 2024.

Scalar emission - corrected Hawking spectrum and energy flux

For $\langle E \rangle \ll E_{\text{brk}}, \quad E_{\text{brk}}^{5/2} c_1 t \gg 1$,

$$\langle E(t) \rangle = \frac{3}{5} \left(\frac{15}{4} \right)^{2/5} \Gamma \left(\frac{7}{5} \right) \frac{1}{(c_1 t)^{2/5}}$$

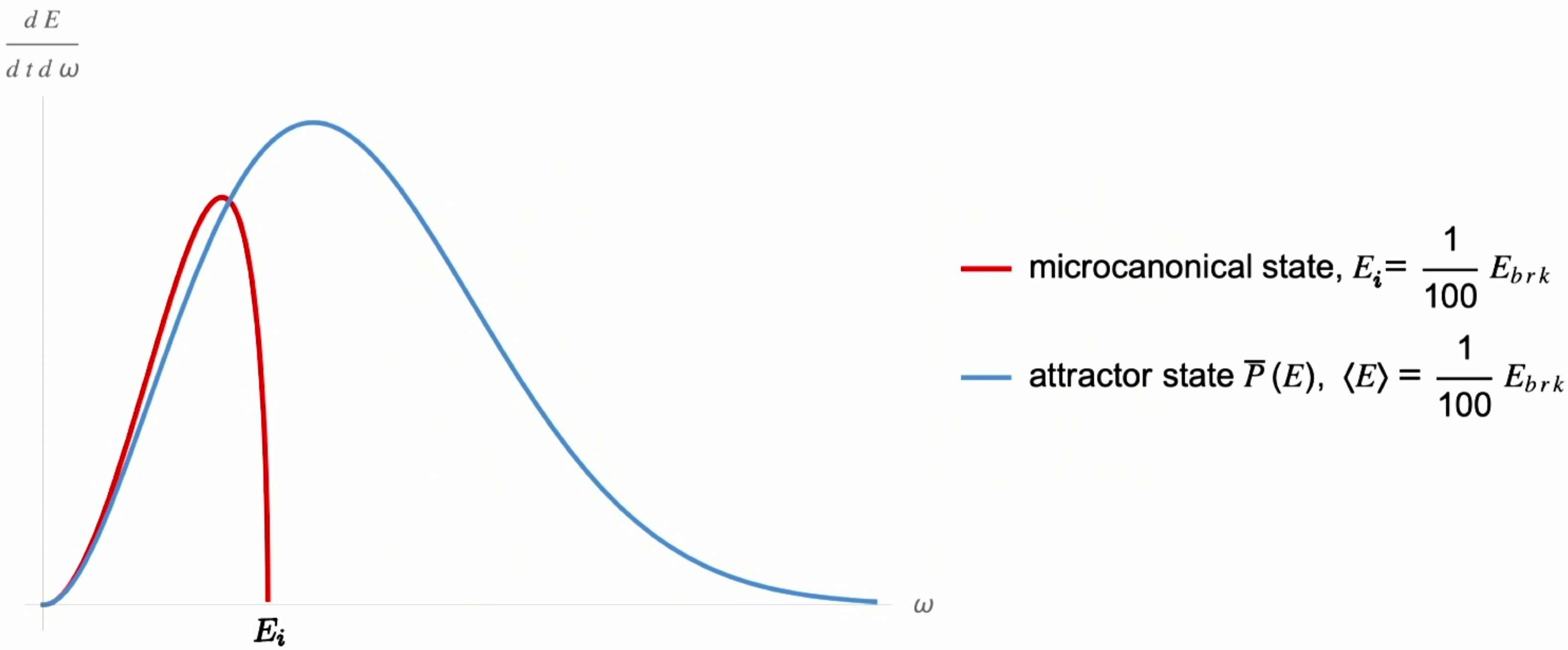
$$\langle \frac{dE}{dt} \rangle = -\frac{8}{27} \sqrt{\frac{5}{3}} \Gamma \left(\frac{7}{5} \right)^{-5/2} c_1 \langle E \rangle^{7/2}$$

We can compare this to the microcanonical result. The energy flux from a black hole in a state of fixed energy E_i is

$$\left. \frac{dE}{dt} \right|_{E_i} = -\frac{16}{105} c_1 E_i^{\frac{7}{2}} \qquad E_i \ll E_{\text{brk}}$$

Comparing at the expected energy $\langle E \rangle = E_i$, the energy flux in the attractor state is larger than the one in the microcanonical ensemble by a factor of ~ 3.4.

Scalar emission - corrected Hawking spectrum and energy flux



Hawking radiation into the $\ell = 0$ massless scalar mode at energy $\langle E \rangle = \frac{1}{100} E_{\text{brk}}$. Blue: the spectrum in the attractor distribution which the black hole would occupy at long times. Red: spectrum in a microcanonical state.

Photon emission - corrected Hawking spectrum and energy flux

For $\langle E\rangle \ll E_{\rm brk}, \quad \varepsilon \ll \langle E\rangle, \quad E_{\rm brk}^{9/2} c_3 t \gg 1$ the expected energy and energy flux of $\ell = 1$ photons from a fermionic black hole is

$$\langle\varepsilon\rangle \approx 0.706976 \times \frac{1}{(c_3 t)^{2/9}}$$

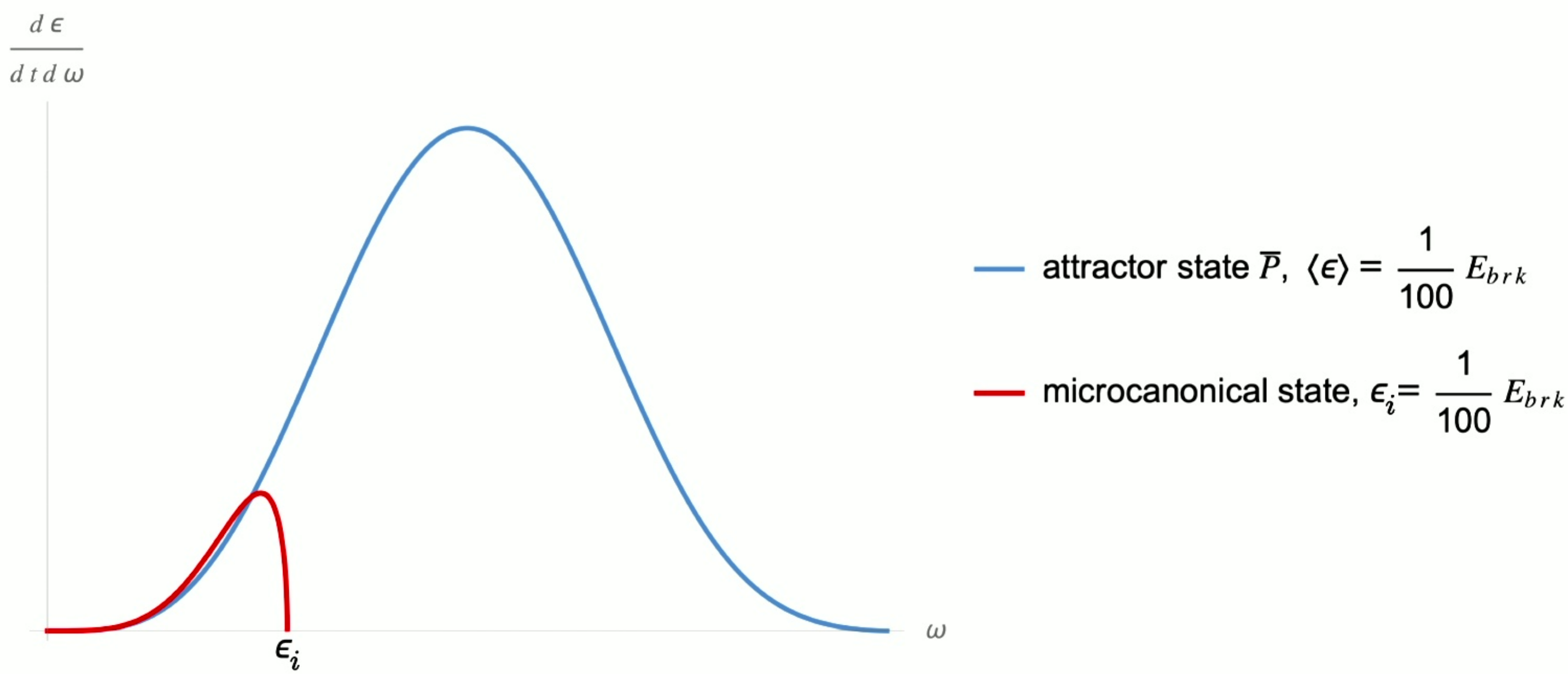
$$\langle\frac{d\varepsilon}{dt}\rangle \approx -182.659 \times c_3 \langle\varepsilon\rangle^{11/2}$$

We can again compare this to a black hole in a microcanonical state with the same expected energy, $E_i = \langle E\rangle$. In the microcanonical ensemble, the energy flux is

$$\left.\frac{d\varepsilon}{dt}\right|_{\varepsilon_i} = -\frac{256}{3465} c_3 \varepsilon_i^{11/2} \tag{1}$$

In this case, the energy flux in the attractor state is larger by a factor of ~ 14.3.

Photon emission - corrected Hawking spectrum and energy flux



Hawking radiation into $\ell = 1$ photons from a black hole with $j = 1/2$ and energy $\langle E\rangle = \frac{1}{100} E_{\text{brk}}$. Blue: Spectrum in the attractor distribution which the black hole would occupy at long times. Red: microcanonical spectrum.

Di-photon emission - corrected Hawking spectrum and energy flux

For $\langle E \rangle \ll E_{\text{brk}}, \quad E_{\text{brk}}^{17/2} c_7 t \gg 1,$

$$\langle E(t) \rangle \approx 0.541649 \times \frac{1}{(c_7 t)^{2/17}}$$

$$\langle \frac{dE}{dt} \rangle \approx -21.5765 \times c_7 \langle E \rangle^{19/2}$$

For comparison, the microcanonical emission rate is

$$\left. \frac{dE}{dt} \right|_{E_i} = -\frac{65536}{2078505} c_7 E_i^{19/2} \quad E_i \ll E_{\text{brk}}$$

The energy flux in the attractor state is larger by a factor of ~ 684.

Why is the energy flux so much larger in the attractor state?

The microcanonical emission spectrum in the low energy limit is

$$\left.\frac{dN}{dtd\omega}\right|_{E_i} = c_a\, \omega^a \sqrt{E_i - \omega}\,\Theta(E_i - \omega) \qquad E_i \ll E_{\text{brk}}$$

The spectrum in the attractor state is given by integrating against the microcanonical one,

$$\langle \frac{dN}{dtd\omega} \rangle = \int_0^\infty dE\ P(E)\ \left.\frac{dN}{dtd\omega}\right|_E$$

The microcanonical emission rate has maximum

$$\left.\frac{dN}{dtd\omega}\right|_{E_i} (\omega_{\text{max}}) \propto E_i^{a+\frac{1}{2}}, \qquad \omega_{\text{max}} = \frac{2a}{2a+1} E_i$$

The contribution of each $|E_i\rangle$ to the total emission rate increases steeply with E_i. A significant proportion of states in $P(E)$ have energy $E_i > \langle E \rangle$, and contribute a significant amount of energy to the total flux.
This is also why the emission rates peak at a frequency larger than the expected energy of the black hole state.

Why is the energy flux so much larger in the attractor state?

The microcanonical emission spectrum in the low energy limit is

$$\left.\frac{dN}{dtd\omega}\right|_{E_i} = c_a\, \omega^a \sqrt{E_i - \omega}\,\Theta(E_i - \omega) \qquad E_i \ll E_{\text{brk}}$$

The spectrum in the attractor state is given by integrating against the microcanonical one,

$$\langle \frac{dN}{dtd\omega} \rangle = \int_0^\infty dE\ P(E)\ \left.\frac{dN}{dtd\omega}\right|_E$$

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$$\left.\frac{dN}{dtd\omega}\right|_{E_i}(\omega_{\text{max}}) \propto E_i^{a+\frac{1}{2}}, \qquad \omega_{\text{max}} = \frac{2a}{2a+1}E_i$$

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Conclusion

We discussed the energy probability density $P(E,t)$ of the charged black hole as it evolves towards extremality in the deep quantum gravity regime.

Below the breakdown scale E_{brk} where the Schwarzian becomes strongly coupled, the state of the black hole evolves toward a non-thermal, universal long-time distribution.

The attractor solutions effectively depend only on one combination of energy and time which fixes the powers of time and energy in $\langle E \rangle$ and $\langle dE/dt \rangle$.

The Hawking fluxes calculated in the attractor state can be much larger than those in a microcanonical state with the same expected energy.

Future directions

The evolving state of other types of near-extremal black holes such as Kerr-Newman[8] or near-BPS black holes in $\mathcal{N} = 2$ supergravity.

Possible phenomenological implications for the lifetime of near-extremal charged primordial black holes.

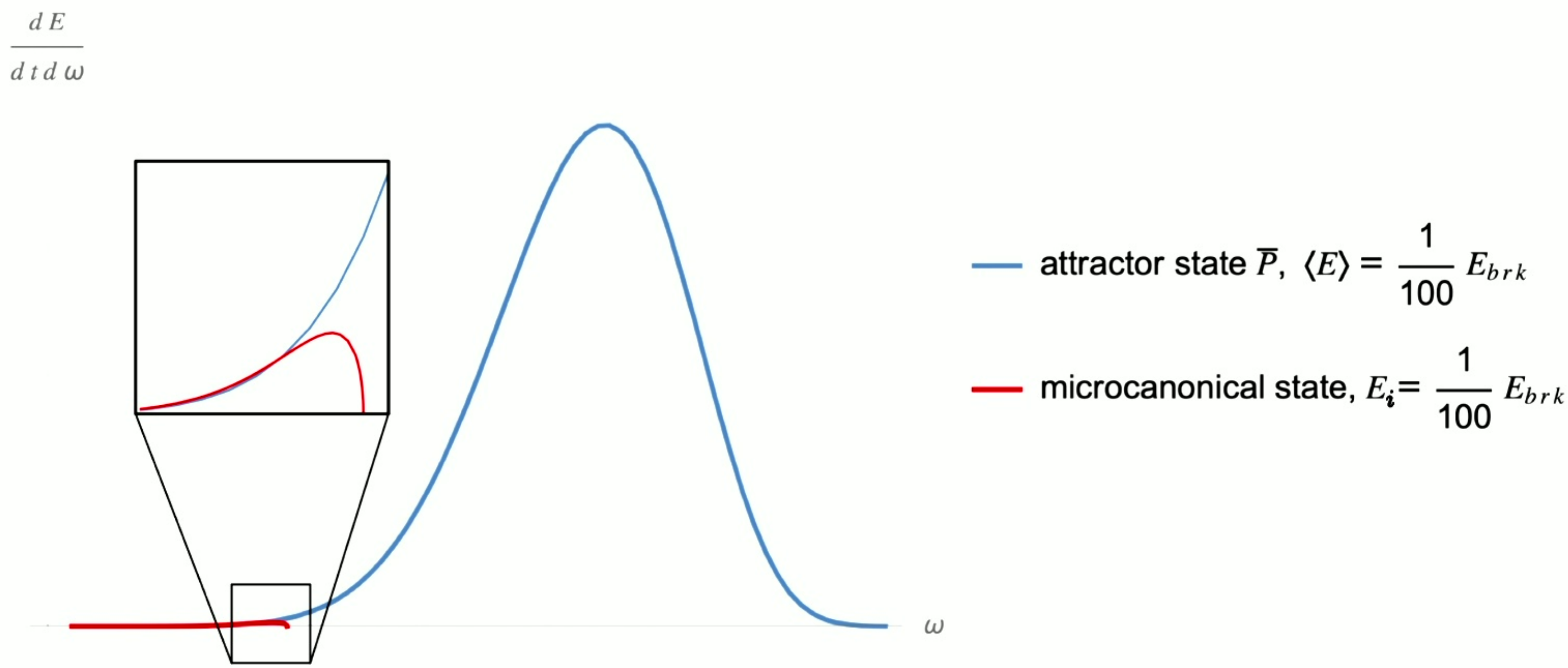
Schwarzian corrections to the decoherence rates of quantum systems in the exterior of near-extremal black holes, such as in the thought experiments proposed by Wald et al.[9]

[8] Maulik, Meng, and Pando Zayas 2025.

[9] Danielson, Satishchandran, and Wald 2022; Gralla and Wei 2024.

Thank you for your attention.

Di-photon emission - Hawking spectrum and energy flux



Hawking radiation into entangled singlet states from a black hole with $j = 0$ and energy $\langle E \rangle = \frac{1}{100} E_{\text{brk}}$. Blue: Spectrum in the attractor distribution which the black hole would occupy at long times. Red: microcanonical spectrum.