Black holes from anomalies

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- Large Schwarzschild & Kerr black holes in AdS_5
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Beyond high T: instabilities, “weak gravity” behaviors & hairy BH’s

High temperature limit of CFT on $S^{D-1} \times R$

- CFT at high T: Euclidean CFT on $S_r^{D-1} \times S_\beta^1$, at $\beta \equiv T^{-1} \ll r$.
- Two contributions to the "free energy" $\log Z\,(\beta, \ldots)$ at $\beta \ll r$.

\- 0-mode path integral on S^{D-1}:

  Absent in standard partition function, but may exist in SUSY indices.

  Cannot contribute to the divergent leading part at $\beta \to 0$.

\- Kaluza-Klein (KK) modes on S^1: divergent contributions to $\log Z$ at $\beta \to 0$.

- KK modes → effective action of background fields (chemical potentials) on S^{D-1}.

\- Universal background fields: metric, gravi-photon & dilaton:

$$ds_D^2 = ds_{D-1}^2 + e^{-2\Phi}(d\tau + a)^2 \qquad \tau \sim \tau + \beta \ \ , \ \ \beta e^{-\Phi} \sim \text{radius of } S^1$$

\- E.g. on S^3 from temperature β^{-1} and angular velocities $\omega_{1,2}/\beta$:

$$ds^2 = r^2\left[d\theta^2 + \sum_{i=1}^2 n_i^2\left(d\phi_i - \frac{i\omega_i}{\beta}d\tau\right)^2\right] + d\tau^2 \ = \ r^2\left[d\theta^2 + \sum_i n_i^2 d\phi_i^2 + \frac{r^2(\sum_i \omega_i n_i^2 d\phi_i)^2}{\beta^2(1 - r^2\sum_i \frac{n_i^2\omega_i^2}{\beta^2})}\right] + e^{-2\Phi}\,(d\tau + a)^2$$

$$e^{-2\Phi} \ = \ 1 - r^2\sum_i \frac{n_i^2\omega_i^2}{\beta^2} \ \ , \quad a = -i\frac{r^2\sum_i \omega_i n_i^2 d\phi_i}{\beta(1 - r^2\sum_i \frac{n_i^2\omega_i^2}{\beta^2})} \qquad\qquad (n_1, n_2) = (\cos\theta, \sin\theta)$$

Effective action & derivative expansion

- In CFT, small β = small r^{-1} = derivative expansion on S^{D-1}.

- ∞-tower of derivative expansion. Most coefficients depend on coupling constants. On S^3,

$$\underbrace{\frac{A_1}{\beta^3}\int d^3x e^{3\Phi}\sqrt{g}}_{\text{cosmological constant}} + \underbrace{\frac{A_2}{\beta}\int d^3x\sqrt{g}e^{\Phi}R}_{\text{Einstein-Hilbert}} + \underbrace{A_3\int d^3x\sqrt{g}(\beta e^{-\Phi})(\beta^{-1}da)^2}_{\text{gravi-photon kinetic}} + \cdots$$

- However, Chern-Simons terms on S^{D-1} (for even D) are coupling independent.
- 't Hooft anomalies determine them. Sub-leading in the normal free energy. On S^3,

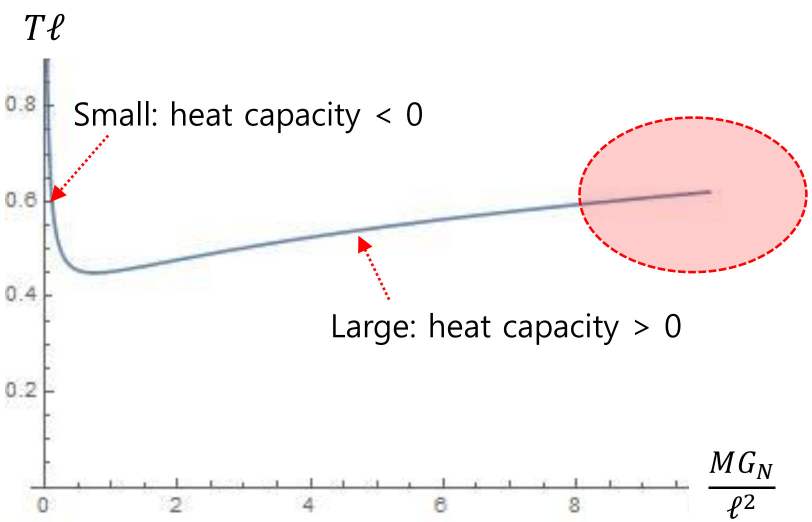
$$\kappa\beta^{-2}\int a\wedge da + \kappa_I\beta^{-1}\int \boxed{\mathcal{A}^I}\wedge da + \kappa_{IJ}\int \mathcal{A}^I\wedge d\mathcal{A}^J + \cdots$$

background gauge fields (~ chemical potential) for global symmetries

- CFT w/ AdS dual: The expansion constrains AdS black hole thermodynamics.

Kerr black holes in AdS_5

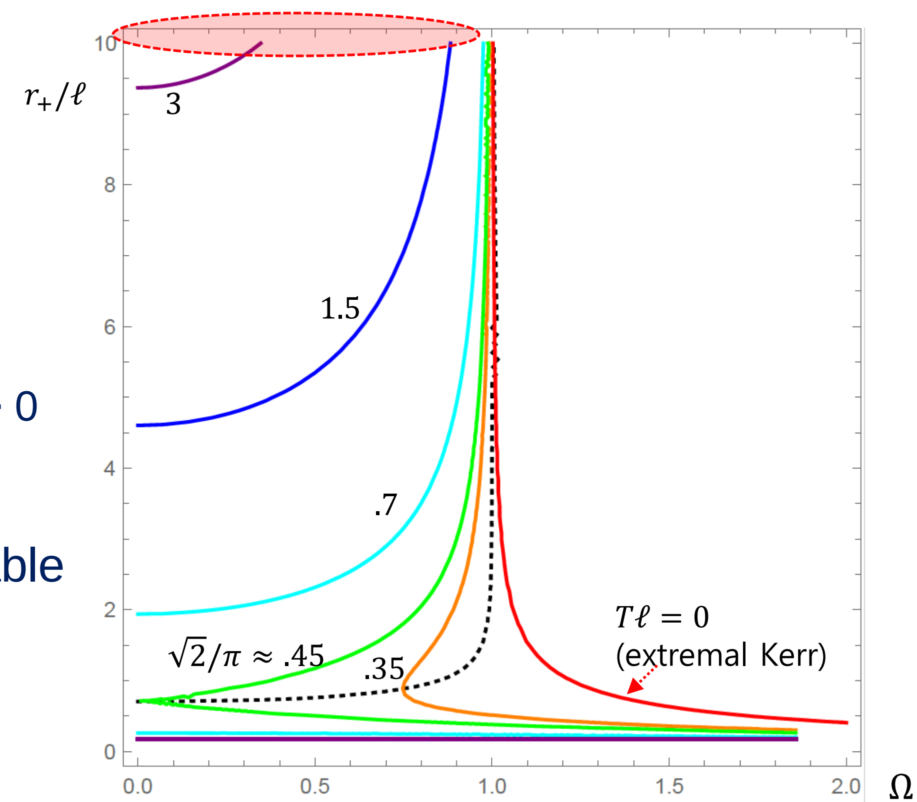
- Schwarzschild BH's: Small & large BH's (ℓ: AdS radius)



$$T = \frac{r_+}{\pi\ell^2} + \frac{1}{2\pi r_+} \qquad r_+^2 = -\frac{\ell^2}{2} + \ell\sqrt{\frac{\ell^2}{4} + \omega M}$$

$$\omega \equiv \frac{16\pi G_N}{3\mathrm{vol}(S^3)}$$

- Kerr BH's at $J_1 = J_2$: isothermal lines →
- Grand canonical ensemble is stable above the dotted line: heat capacity, susceptibility > 0



- Our studies are expected to capture stable BH's in the regions at high T.

Kerr black holes in AdS_5

- Plug in our S^3 background fields to:

$$\frac{A_1}{\beta^3}\int d^3x e^{3\Phi}\sqrt{g} + \frac{A_2}{\beta}\int d^3x\sqrt{g}e^{\Phi}R + A_3\int d^3x\sqrt{g}(\beta e^{-\Phi})(\beta^{-1}da)^2 + \cdots$$

- Leading terms in derivative expansion. (Take $\Omega \equiv \omega_1 = \omega_2$ for simplicity.)

$$S = \log Z + \beta E + 2\Omega J = A_1\frac{\beta}{(\beta^2-\Omega^2)^2} + A_2\frac{3\beta^3 - 4\beta\Omega^2}{(\beta^2-\Omega^2)^2} + A_3\frac{\beta\Omega^2}{(\beta^2-\Omega^2)^2} + \beta E + 2\Omega J + \mathcal{O}(\beta^0)$$

- Extremize in β, Ω & solve for E, J, S. 2 more eqns.
- Expressions for S, E, J in terms of β, Ω, and the unknowns A_i.

- Fitting three A_i's, one reproduces all divergent parts of E, J, S carried by Kerr BH's. (Here, $\Omega = -a\beta + \cdots$.) Although 3 parameters are fitted by hand, still quite nontrivial constraints.

$$E = \frac{\pi l^2}{G}\left(\frac{3+a^2}{8(1-a^2)^3}\frac{l^4\pi^4}{\beta^4} - \frac{3+3a^2-2a^4}{8(1-a^2)^3}\frac{l^2\pi^2}{\beta^2} - \frac{3-11a^2+6a^4-2a^6}{16(1-a^2)^3} + \mathcal{O}(\beta^2)\right)$$

$$J = \frac{\pi l^2}{G}\left(\frac{a}{4(1-a^2)^3}\frac{l^4\pi^4}{\beta^4} - \frac{2a-a^3}{4(1-a^2)^3}\frac{l^2\pi^2}{\beta^2} + \frac{a}{8(1-a^2)^3} + \mathcal{O}(\beta^2)\right)$$

$$S = \frac{\pi^2 l^3}{G}\left(\frac{1}{2(1-a^2)}\frac{l^3\pi^3}{\beta^3} - \frac{3-2a^2}{4(1-a^2)^2}\frac{l\pi}{\beta} + \mathcal{O}(\beta)\right)$$

analysis by Nahmgoong (2018)

Comments

- This calculus is related to various well-known studies on black holes in AdS.

- Black D3-brane factor: [Klebanov, Tseytlin] (1996)

- In $AdS_5 \times S^5$, the first coefficient A_1 is simply the famous "$3/4$" factor for black D3-branes, since large Schwarzschild BH is locally a black brane.
- Once this strong coupling factor is fixed, the remaining physics of large Kerr BH is determined from kinematics. (And similarly, NLO terms all constrained up to A_2, A_3.)

- Hydrodynamic calculus: [Bhattacharyya, Lahiri, Loganayagam, Minwalla] (2007)

- Solve Navier-Stokes equation for quark-gluon plasma: Hydrodynamics valid at $\ell_{mfp} \ll r$.
- Equivalent to $T \gg r^{-1}$ for rotating plasma (Kerr BH), but in general different.
- E.g. hydrodynamics applicable to certain extremal BH's (at $T = 0$), but not to BPS BH's.

- On the other hand, our approach doesn't apply to general extremal BH's, but applies to indices for BPS black holes at large charges.

High T expansion of indices: expectation

- So far, we constrained high T partition functions (with unknowns $A_1, A_2, A_3 \dots$).

- Similar calculus for an index: This method is much more powerful.
- Naturally expect CS terms may determine the divergent part (coupling-independent).
- Index: For 4d SCFTs on $S^3 \times S^1$, index is defined by

$$Z(\omega_1, \omega_2) \equiv \mathrm{Tr}\left[e^{-\beta(E-\frac{3}{2}R-J_1-J_2)}e^{-\frac{1}{2}\Delta R}e^{-\omega_1 J_1-\omega_2 J_2}\right]$$

$$\Delta = \omega_1 + \omega_2 + 2\pi i \qquad E - \frac{3}{2}R - J_1 - J_2 \sim \{\mathcal{Q}, \mathcal{Q}^\dagger\}$$

- β is merely a regulator of the index. Easy to compute at $\beta \to 0^+$.
- Since leading term is $O(\beta^0)$, $\beta \to 0$ is a fake thermal circle parameter.
- True derivative expansion parameters are $|\omega_i| \ll 1$.

$$e^{-2\Phi} = 1 + \boxed{r^2 \sum_i \frac{n_i^2 \omega_i^2}{\beta^2}} \quad , \quad a = -i\frac{r^2 \sum_i \omega_i n_i^2 d\phi_i}{\beta(1 - r^2 \sum_i \frac{n_i^2\omega_i^2}{\beta^2})}$$

- Some reordering of naïve derivative expansion is inevitable.

Background field set-up

- The background fields for the index:

- same metric, dilaton, gravi-photon as before (up to small shifts of ω_i by β)
- Background $U(1)_R$ & other flavor symmetries' gauge fields:

$$A^I = -\frac{i\Delta_I}{\beta}d\tau \xrightarrow{\text{rearrange to 3d fields}} A^I = A^I_4(d\tau + a) + \mathcal{A}^I \qquad A^I_4 = -\frac{i\Delta^I}{\beta} \qquad \mathcal{A}^I = -A^I_4 a$$

- Non-CS terms:

Indeed, all vanish at $\beta \to 0$.

$$\begin{aligned}
\frac{1}{(2\pi)^2}\int \beta^{-3}e^{3\Phi}\sqrt{g} &= \frac{\beta r^3}{2(\beta^2-r^2\omega^2)^2} = \frac{\beta}{2r\omega^4} + \mathcal{O}\left(\frac{\beta^3}{r^3\omega^6}\right) \qquad (2.62)\\
\frac{1}{(2\pi)^2}\int \beta^{-1}e^{\Phi}\sqrt{g}\,\mathcal{R}^{ab}{}_{ab} &= \frac{r(3\beta^3-4\beta r^2\omega^2)}{(\beta^2-r^2\omega^2)^2} = -\frac{4\beta}{r\omega^2} + \mathcal{O}\left(\frac{\beta^3}{r^3\omega^4}\right)\\
\frac{1}{(2\pi)^2}\int \beta e^{-\Phi}\sqrt{g}\,\mathcal{F}^I_{ab}\mathcal{F}^J_{ab} &= \frac{\beta\Delta^I\Delta^J r^3\omega^2}{(\beta^2-r^2\omega^2)^2} = \frac{\beta\Delta^I\Delta^J}{r\omega^2} + \mathcal{O}\left(\frac{\beta^3}{r^3\omega^4}\right)\\
\frac{1}{(2\pi)^2}\int \beta^3 e^{-3\Phi}\sqrt{g}\,(\nabla_c\mathcal{F}^I_{ab})(\nabla^c\mathcal{F}^{Jab}) &= \frac{2\beta^3 r\omega^2\Delta^I\Delta^J}{(\beta^2-r^2\omega^2)^2} = \frac{2\beta^3\Delta^I\Delta^J}{r^3\omega^2} + \mathcal{O}\left(\frac{\beta^5}{r^5\omega^4}\right)\\
\frac{1}{(2\pi)^2}\int \beta^3 e^{-3\Phi}\sqrt{g}\,(\nabla_c\mathcal{F}^I_{ab})(\nabla^a\mathcal{F}^{Jcb}) &= \frac{\beta^3 r\omega^2\Delta^I\Delta^J}{(\beta^2-r^2\omega^2)^2} = \frac{\beta^3\Delta^I\Delta^J}{r^3\omega^2} + \mathcal{O}\left(\frac{\beta^5}{r^5\omega^4}\right)\\
\frac{1}{(2\pi)^2}\int \beta e^{-\Phi}\sqrt{g}\,\mathcal{R}^{ab}{}_a{}^c\mathcal{R}_b{}^d{}_{cd} &= \frac{2\beta(8r^4\omega^4-8\beta^2r^2\omega^2+3\beta^4)}{r(\beta^2-r^2\omega^2)^2} = \frac{16\beta}{r} + \mathcal{O}\left(\frac{\beta^3}{r^3\omega^2}\right)\\
\frac{1}{(2\pi)^2}\int \beta e^{-\Phi}\sqrt{g}\,\mathcal{R}_{abcd}\mathcal{R}^{abcd} &= \frac{32\beta r^4\omega^4-16\beta^3r^2\omega^2+6\beta^5}{r(\beta^2-r^2\omega^2)^2} = \frac{32\beta}{r} + \mathcal{O}\left(\frac{\beta^3}{r^3\omega^2}\right)\\
\frac{1}{(2\pi)^2}\int \beta^3 e^{-3\Phi}\sqrt{g}\,\mathcal{F}^{Iab}\mathcal{F}^{Jc}_a\mathcal{R}_b{}^d{}_{cd} &= \frac{2\beta\Delta^I\Delta^J r\omega^2(\beta^2-2r^2\omega^2)}{(\beta^2-r^2\omega^2)^2} = -\frac{4\beta\Delta^I\Delta^J}{r} + \mathcal{O}\left(\frac{\beta^3}{r^3\omega^2}\right)\\
\frac{1}{(2\pi)^2}\int \beta^3 e^{-3\Phi}\sqrt{g}\,\mathcal{F}^I_{ab}\mathcal{F}^J_{cd}\mathcal{R}^{abcd} &= \frac{2\beta\Delta^I\Delta^J r\omega^2(\beta^2-4r^2\omega^2)}{(\beta^2-r^2\omega^2)^2} = -\frac{8\beta\Delta^I\Delta^J}{r} + \mathcal{O}\left(\frac{\beta^3}{r^3\omega^2}\right)\\
\frac{1}{(2\pi)^2}\int \beta^5 e^{-5\Phi}\sqrt{g}\,\mathcal{F}^{Iab}\mathcal{F}^{Jc}_a\mathcal{F}^{Kd}_b\mathcal{F}^L_{cd} &= \frac{\beta\Delta^I\Delta^J\Delta^K\Delta^L r^3\omega^4}{(\beta^2-r^2\omega^2)^2} = \frac{\beta\Delta^I\Delta^J\Delta^K\Delta^L}{r} + \mathcal{O}\left(\frac{\beta^3}{r^3\omega^2}\right)\\
\frac{1}{(2\pi)^2}\int \beta^5 e^{-5\Phi}\sqrt{g}\,\mathcal{F}^{Iab}\mathcal{F}^J_{ab}\mathcal{F}^{Kcd}\mathcal{F}^L_{cd} &= \frac{2\beta\Delta^I\Delta^J\Delta^K\Delta^L r^3\omega^4}{(\beta^2-r^2\omega^2)^2} = \frac{2\beta\Delta^I\Delta^J\Delta^K\Delta^L}{r} + \mathcal{O}\left(\frac{\beta^3}{r^3\omega^2}\right).
\end{aligned}$$

CS coefficients from anomalies

- So if one knows the CS coefficients, one can determine the Cardy free energy.

- Although a general presentation is possible, I will phrase the results for SCFTs where R-symmetry anomaly is related to the Weyl anomaly.
- There turn out to be two types of CS terms: gauge non-invariant & invariant

- Gauge non-invariant CS terms:

- Effective action $S_{\text{eff}} = -\log Z$ should respect 't Hooft anomaly: $\delta_\epsilon S_{\text{eff}} \sim \epsilon\, F \wedge F + \ldots$
- This should be reflected in the effective action for 3d background fields.
- Anomaly demands the existence of certain gauge non-invariant CS terms.

[Banerjee, Bhattacharya, Bhattacharyya, Jain, Minwalla, Sharma] (2012)

$$S_{\text{non-inv}} = -i\frac{\beta(5a-3c)}{8\pi^2}\int_{S^3} C_{IJK}\left(A_4^I \mathcal{A}^J \wedge d\mathcal{A}^K + A_4^I A_4^J \mathcal{A}^K \wedge da + \frac{1}{3}A_4^I A_4^J A_4^K a \wedge da\right)$$

[C_{IJK}: coefficients of cubic anomaly polynomial in the normalization $C_{111} = 6$ for $U(1)_R$ at $I = 1$]

CS coefficients from anomalies

- Gauge invariant CS terms:

- More elaborate arguments [Jensen, Loganayagam, Yarom] (2013) determine them anomalies. (See also [Di Pietro, Komargodski] (2014).)
- The basic idea is to study the effective action on $R^4 \sim R^2 \times R_r^+ \times S_\tau^1$, on which $\langle T^{r\tau}\rangle = 0$.
- Regarding the circle (near the origin…?) as the thermal circle, imposing $\langle T^{r\tau}\rangle = 0$ determines the gauge-invariant CS coefficients.

- CPT invariance forbids gravi-photon CS term and internal symmetries' CS terms:

$$\cancelto{\kappa = 0}{\kappa\beta^{-2}\int a\wedge da} + \kappa_I\beta^{-1}\int \mathcal{A}^I\wedge da + \cancelto{\kappa_{IJ} = 0}{\kappa_{IJ}\int \mathcal{A}^I\wedge d\mathcal{A}^J} + \cdots$$

- Mixed CS term: [Jensen, Loganayagam, Yarom] [Di Pietro, Komargodski]

$$S_{\text{inv}} \sim -i(a-c)\beta^{-1}\int_{S^3} k_I\mathcal{A}^I\wedge da$$

- Other CS terms, involving gravitational terms, determined as well.

[k_I: mixed anomaly coefficients. In the convention here, $k_1 = 1$ for $U(1)_R$ and other $k_I = 0$]

Universal Cardy free energy

- Plugging in our background fields, one obtains the Cardy free energy.

$$\log Z \sim (5a-3c)\frac{C_{IJK}\Delta^I\Delta^J\Delta^K}{6\omega_1\omega_2} + \frac{4\pi^2(a-c)k_I\Delta^I}{\omega_1\omega_2}$$

- For simplicity, just turn on chemical potential for superconformal R-symmetry.

$$\log Z \sim -\frac{16i}{27}\frac{3c-2a}{\omega_1\omega_2} \qquad C_{111}=6\ ,\quad k_1=1\ ,\quad \Delta^1=\frac{2\pi i}{3}$$

- $3c-2a>0$ is always met by all interacting SCFTs. [Hofman, Maldacena]

- Legendre transformation at macroscopic charges → entropy…? : **Subtle...!**

$$S(J_i+\tfrac{R}{2};\omega_i) = -\frac{16i}{27}\frac{3c-2a}{\omega_1\omega_2} + \omega_1(J_1+\tfrac{R}{2}) + \omega_2(J_2+\tfrac{R}{2})$$

- Naively doing it, one obtains 3 saddle points w/ “complex entropies”

$$S_* = 3\left[-i\frac{16(3c-2a)}{27}(J_1+\tfrac{R}{2})(J_1+\tfrac{R}{2})\right]^{\frac{1}{3}} = \left\{\begin{array}{c}\frac{\sqrt{3}-i}{2}\\ i\\ \frac{-\sqrt{3}-i}{2}\end{array}\right\}\times 3\left[\frac{16(3c-2a)}{27}(J_1+\tfrac{R}{2})(J_1+\tfrac{R}{2})\right]^{\frac{1}{3}}$$

Cardy entropy & black holes

- Naïve "complex entropies" at 3 saddles:

- "Degeneracy" $\Omega = e^{Re(S_*)} e^{iIm(S_*)}$ has an oscillating phase as a function of charges.

$$S_* = 3\left[-i\frac{16(3c-2a)}{27}(J_1+\tfrac{R}{2})(J_1+\tfrac{R}{2})\right]^{\frac{1}{3}} = \left\{\begin{array}{c}\frac{\sqrt{3}-i}{2}\\ i\\ \frac{-\sqrt{3}-i}{2}\end{array}\right\} \times 3\left[\frac{16(3c-2a)}{27}(J_1+\tfrac{R}{2})(J_1+\tfrac{R}{2})\right]^{\frac{1}{3}}$$

- Interpretation:

- In the index, $\Omega(J_i + R/2)$ can be macroscopic w/ either + or – overall sign.
- Macroscopic Legendre transformation doesn't know the precise quantized value of charges. So the calculation gets confused about the overall sign at a given charge sector.
- Proposal: Sign deprived $|\Omega(J_i + R/2)| \sim e^{Re(S_*)}$ is the degeneracy of the index.

  Choose the maximal $Re(S_*)$ as the entropy of the index.

- This viewpoint is numerically justified: compare w/ fugacity expansion to high charge order.

- The entropy obtained this way precisely agrees with the Bekenstein-Hawking entropy of large BPS black holes in AdS_5.

Other even dimensions

- The method applies to non-Lagrangian SCFTs in even dimensions.

- 4d non-Lagrangian theories: Argyres-Douglas, Minahan-Nemeschansky, …
- This method reproduces the Cardy formula for 2d SCFTs [J. Kim, K. Lee, J. Park] (2018)

$$Z(\tau) \sim \mathrm{Tr}\left[e^{2\pi i \tau L_0}\right] \sim \exp\left[\frac{\pi i c}{12\tau}\right] \quad \text{at } \tau \to i0^+$$

- Or more generally, flavor-refined Cardy formula is also available.

- 6d SCFTs & AdS_7 black holes: [CKKN] [Nahmgoong]

- (2,0) theory on N M5-branes (A_{N-1} type + free (2,0) tensor):

$$\log Z \sim \boxed{-\frac{N^3}{24}\frac{\Delta_1^2\Delta_2^2}{\omega_1\omega_2\omega_3}} - \frac{N((\Delta_1+\Delta_2)^2+4\pi^2)((\Delta_1-\Delta_2)^2+4\pi^2)}{192\omega_1\omega_2\omega_3} + \frac{N(\Delta_1^2+\Delta_2^2-4\pi^2)(\omega_1^2+\omega_2^2+\omega_3^2)}{96\omega_1\omega_2\omega_3} + \mathcal{O}(\log\omega)$$

dominant term in the large N limit

- Similar analysis possible for 6d (1,0) SCFTs, e.g. E-string theory [Nahmgoong]

Instabilities at finite/zero T

- Physics is best illustrated w/ extremal BH's at zero T:
  - Most of them are non-BPS: $E_{ext} = 3\Delta Q + 2\Omega J > 3Q + 2J \equiv E_{BPS}$. That is, $\Delta, \Omega > 1$.
  - "Weak gravity conjecture": ∃ particles with large enough "charge/mass" ratio.

Unprotected extremal BH's should decay.

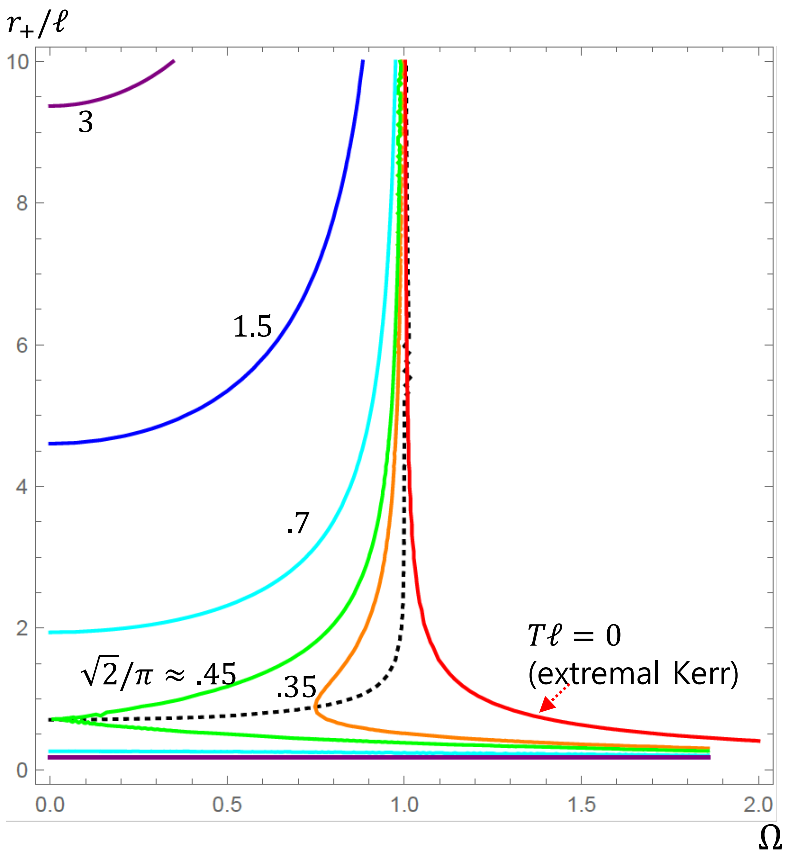
- Extremal Kerr BH's: $\Omega > 1$.
  - Large "charge/mass" particles are conformal descendants $\partial^n \hat{O}_{\Delta,j}$ w/ $n \gg \Delta, J$

$$\frac{J}{M} = \frac{\Delta + n/2}{j+n} \approx \frac{1}{2} > (2\Omega)^{-1} = \frac{J_{ext}}{M_{ext}}$$

  - BH emits spin & decay: "tachyonic" instability
  - General Kerr BH's are unstable at $\Omega > 1$. [Murata]

- Similar tachyonic decays for charged BH's:
  - Scalars emitted, forming "hairy BH's"

[Gubser] [Bhattacharyya, Minwalla, Papadodimas] …



Consequence of instabilities in the BPS sector

- A rather confusing property of the known BPS BH solutions in AdS is that, they carry one less parameters than the number of distinct charges.

- E.g. in $AdS_5 \times S^5$: BPS mass $M\ell = Q_1 + Q_2 + Q_3 + J_1 + J_2$. 5 distinct charges.
- 5 charges & 4 parameters: implies a "charge relation" [Gutowski, Reall] [Kunduri, Lucietti, Reall]

$$Q_1Q_2Q_3 + \frac{N^2}{2}J_1J_2 = \left(\frac{N^2}{2} + Q_1 + Q_2 + Q_3\right)\left(Q_1Q_2 + Q_2Q_3 + Q_3Q_1 - \frac{N^2}{2}(J_1 + J_2)\right)$$

- More general "hairy BH's" formed by decay of non-BPS extremal BH's…?
- Numerical studies [Bhattacharyya, Minwalla, Papadodimas] [Markeviciute, Santos] [Markeviciute]

- Probe picture of hairy BH's: [Choi, SK] in progress

- Slightly violate charge relation and construct hairy BH's from "known BH + probes"
- Hairy BH's formed only for over-rotating violation of charge relation: $J > J_{BH}(Q)$
- This analytically confirms numerical observations on BPS hairy BH's. [Markeviciute, Santos]
- Relations to the near-BPS extremal BH's are under investigation.

Conclusion & remarks

- One can make universal studies of radially quantized CFTs at high T.
- In general, this strongly constrains the physics of large AdS BH's
- "Anomalies + some SUSY" derive the physics of large BPS BH's in $AdS_{3,5,7}$.
- Rather abstract "degrees of freedom" measured by anomaly is related to the high T d.o.f.

- Today's talk, the published parts on the arXiv, is only a small subset of the recent advances in BPS AdS black holes.
- Large BPS AdS_{D+1} black holes from $SCFT_D$ Cardy indices for all $D = 3,4,5,6$.
- BPS black holes at general size in AdS_5 from $SCFT_4$ [Benini, Milan] [Cabo-Bizet, Murthy]

- Beyond large BH & high T limit, one expects a very rich spectrum of BH's.
- As a concrete example, we have already predicted the existence of over-rotating hairy BH's from CFT, in a 1/8-BPS sector of maximal SYM, by studying the so-called Macdonald index. [Gadde, Rastelli, Razamat,Yan]
- Can we construct the new solutions from gravity…?