

Ramp and plateau in bulk correlators within the disk topology in Jackiw-Teitelboim gravity

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We study bulk two-point correlation functions of a massless scalar field in Jackiw-Teitelboim gravity around the eternal black hole saddle. While same-side correlators exhibit exponential decay, two-sided correlators, at the next-to-leading order in steepest-descent approximation, exhibit a ramp followed by a plateau after the initial semiclassical exponential decay. Our results indicate that the late-time saturation of two-sided correlators is already visible within the perturbative saddle expansion of the bulk path integral, without invoking nonperturbative e^{-S} effects. Finally, we show that the dip time, defined as the minimum of the correlator, grows inversely with the black hole temperature, consistent with expectations from the holographic dual.

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I. INTRODUCTION

Black holes are known to radiate. Hawking's seminal computation [1,2] shows that an evaporating black hole emits an almost thermal spectrum characterized only by its macroscopic parameters (mass, charge, angular momentum). If this conclusion is extrapolated to the end point of evaporation, one is led to the conclusion that pure states evolve into mixed states, apparently violating unitarity and threatening the foundations of quantum mechanics. This conflict is commonly referred to as the black hole information paradox [3,4].

One might hope that the missing information is merely hidden behind the event horizon, inaccessible to an exterior observer until the very final stages of evaporation. Page famously argued, however, that if black hole evaporation is unitary, the entanglement entropy of the Hawking radiation must begin to decrease long before the black hole shrinks to Planckian size [4,5]. The corresponding Page time typically occurs when roughly half of the initial Bekenstein-Hawking entropy has been radiated away, at which point the black hole is still semiclassical and large. The unavoidable conclusion is that the standard Hawking calculation—firmly rooted in semiclassical gravity—must already miss

an essential ingredient well before curvatures become extreme.

Confronted with this puzzle, it is natural to seek gravitational models where quantum effects are tractable and information flow can be analyzed exactly. Jackiw-Teitelboim (JT) [6,7] gravity—two-dimensional dilaton gravity with a negative cosmological constant—has emerged as a particularly fruitful setting. Classically, it is simple: the metric has constant negative curvature, and the dilaton carries the nontrivial gravitational degree of freedom. Quantum-mechanically, it is nontrivial: the theory admits black holes with finite entropy and a well-defined Hamiltonian description of boundary dynamics governed by the Schwarzian action [8–10]. Moreover, semiclassical JT gravity is conjectured to be the gravitational dual of the universal low-energy limit of Sachdev-Ye-Kitaev (SYK)-like models (from now on, simply SYK theory) [11,12]. This theory corresponds to a large number of interacting Majorana fermions at low energies (with respect to the coupling constants)—see, e.g., Ref. [13].

In SYK, disorder-averaged correlators exhibit the characteristic dip-ramp-plateau structure associated with chaotic quantum systems [12,14]. These features are most prominently discussed in the context of spectral observables and ensemble-averaged quantities, where nonperturbative effects related to level discreteness play a central role. By contrast, the two-sided correlators in the thermofield-double (TFD) state of two entangled SYK systems, which is the boundary dual of bulk observables connecting the two asymptotic regions of the eternal black hole, have received far less attention. This raises the question of whether the late-time behavior of such TFD correlators necessarily relies on nonperturbative spectral effects.

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In the bulk JT-gravity description, recent efforts to reproduce the ramp in the boundary correlators incorporate connected wormhole topologies into the gravitational path integral [14–19]. The leading connected contribution arises from the so-called double-trumpet Euclidean wormhole [18], which yields the linear ramp, while the plateau is commonly attributed to genuinely nonperturbative effects in Newton’s constant [20,21]. In these approaches, the gravitational path integral is supplemented by a topological term that governs the relative weighting of geometries and enables the inclusion of connected wormhole contributions. While such a sum over topologies can be implemented at the boundary, its precise interpretation for bulk correlators remains subtle.¹ From the viewpoint of the steepest-descent expansion, one may instead ask whether the relevant saddle structure is already encoded in the original gravitational action, without the need to introduce additional topological weighting.

Following the philosophy of Ref. [22], in contrast to previous approaches, we work entirely within the disk topology for bulk correlators and show that while the same-side (LL) correlator exhibits the typical semiclassical decay, a dip-ramp-plateau structure arises already from the perturbative expansion of the correlation function evaluated across two maximally entangled black hole exteriors. At leading order, the two-sided (LR) correlator exhibits pure exponential decay. However, at the next perturbative order in the gravity coupling constant, subleading corrections generate a linear ramp that ultimately saturates at a constant, power-suppressed plateau. This shows that nondecaying late-time correlations are already encoded within perturbative fluctuations around the principal Lorentzian saddle of the gravitational path integral, without invoking additional topologies or genuinely non-perturbative e^{-S} effects.

II. JT GRAVITY

Jackiw-Teitelboim gravity is a two-dimensional model of dilaton gravity [6,7,18] defined by a bulk action involving the product of the dilaton scalar field Φ and the Ricci scalar R , supplemented by the Gibbons-Hawking-York (GHY) boundary term. To focus on 2D anti-de Sitter space (AdS_2), we set $\Lambda = -\frac{1}{L^2}$ and work in units where the AdS_2 length is set to $L = 1$:

$$S_{\text{JT}}[g, \Phi] = \frac{1}{16\pi G_N} \left(\int_{\mathcal{M}} \Phi(R+2)\sqrt{-g}d^2x + 2 \oint_{\partial\mathcal{M}} \Phi|_{\partial}(K-1)\sqrt{-h}dx \right). \quad (1)$$

¹C. G. thanks Vijay Balasubramanian for pointing this out.

A. Metric solutions

In two dimensions, the Ricci scalar R fully determines the local geometry, since the Riemann tensor has only one component. Thus, varying Eq. (1) with respect to Φ fixes the background manifold to be locally AdS_2 through $R = -2$. This determines the metric up to coordinate transformations.

In global coordinates, the AdS_2 metric takes the form

$$ds^2 = \frac{1}{\cos^2(z_G)} (-dt_G^2 + dz_G^2) = -\frac{1}{\cos^2(\frac{1}{2}(u_G - v_G))} du_G dv_G, \quad (2)$$

where $z_G \in [-\frac{\pi}{2}, \frac{\pi}{2}]$, $t_G \in \mathbb{R}$, and the light-cone coordinates are defined as $u_G = t_G + z_G$ and $v_G = t_G - z_G$. These coordinates cover the maximal extension of the AdS_2 manifold.

One can get to the Poincaré patch through the transformations $u_G = 2\tan^{-1}(u_P) - \frac{\pi}{2}$ and $v_G = 2\tan^{-1}(v_P) + \frac{\pi}{2}$:

$$ds^2 = \frac{1}{z_P^2} (-dt_P^2 + dz_P^2) = -\frac{4}{(u_P - v_P)^2} du_P dv_P, \quad (3)$$

where $z_P \in (0, \infty)$, and $t_P \in \mathbb{R}$. This patch covers only a region inside the global manifold. The Poincaré time t_P does not correspond to the global time t_G , and therefore, one expects the Poincaré vacuum to be different from the global vacuum. However, as was shown in Refs. [23,24], this is not the case, and both vacua are in fact equivalent.

Furthermore, one can introduce Rindler (or, equivalently, Schwarzschild) coordinates to describe two black hole exterior patches in the global AdS_2 manifold:

$$\text{Patch } \textcircled{I} \begin{cases} u_G = 2\tan^{-1} \left[\frac{\beta}{\pi} \tanh \left(\frac{\pi}{\beta} u_R^I \right) \right] - \frac{\pi}{2}, \\ v_G = 2\tan^{-1} \left[\frac{\beta}{\pi} \tanh \left(\frac{\pi}{\beta} v_R^I \right) \right] + \frac{\pi}{2}, \end{cases} \quad (4)$$

$$\text{Patch } \textcircled{II} \begin{cases} u_G = -2\tan^{-1} \left[\frac{\beta}{\pi} \tanh \left(\frac{\pi}{\beta} u_R^{II} \right) \right] + \frac{\pi}{2}, \\ v_G = -2\tan^{-1} \left[\frac{\beta}{\pi} \tanh \left(\frac{\pi}{\beta} v_R^{II} \right) \right] - \frac{\pi}{2}. \end{cases} \quad (5)$$

The constant β is related to the acceleration of an observer at fixed z_R (or, equivalently, to the “black hole” mass) and quantum-mechanically corresponds to the inverse horizon temperature [25].

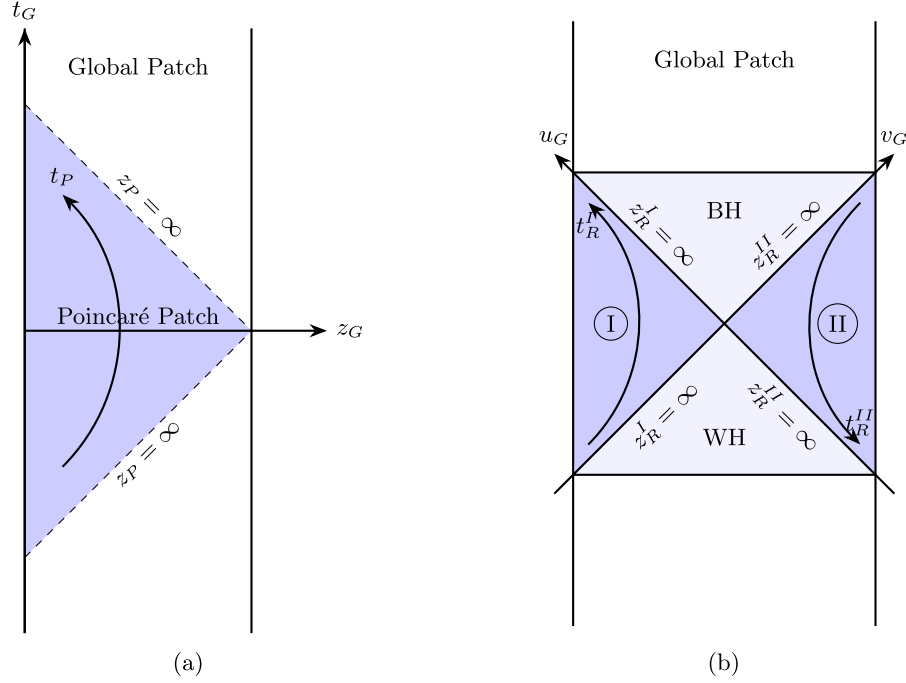


FIG. 1. Three coordinate systems on AdS_2 : (a) Global and Poincaré patches. (b) Global and Schwarzschild patches.

The resulting Rindler metric takes the form

$$ds^2 = \left(\frac{2\pi}{\beta}\right)^2 \frac{1}{\sinh^2\left(\frac{2\pi}{\beta} z_R\right)} (-dt_R^2 + dz_R^2) \\ = -\left(\frac{2\pi}{\beta}\right)^2 \frac{1}{\sinh^2\left(\frac{\pi}{\beta}(u_R - v_R)\right)} du_R dv_R, \quad (6)$$

where $z_R \in (0, \infty)$, and $t_R \in \mathbb{R}$ in both patches, each of which is causally disconnected from the other by an event horizon at $z_R^I, z_R^{II} = \infty$.

Using the radial coordinate $r = r_h \coth(r_h z_R)$, one can also rewrite this metric in the Schwarzschild form:

$$ds^2 = -f(r) dt_R^2 + \frac{dr^2}{f(r)}, \quad (7)$$

where $r \in (r_h, \infty)$, $f(r) = r^2 - r_h^2$, and the horizon is now located at $r_h = \frac{2\pi}{\beta}$.

Each of the patches in Fig. 1(b) and in Eqs. (4) and (5) possesses its own asymptotic boundary and its own time-like Killing vector, generating time evolution t_R within that patch. The two exterior regions are connected through the interior of the eternal black hole geometry, forming a smooth but nontraversable wormhole in the sense of a connected AdS_2 spacetime with two asymptotic boundaries. In the holographic dual description, these boundaries correspond to two decoupled boundary quantum systems [26]. Quantum-mechanically, the eternal geometry is dual to a thermofield-double (TFD) state of the two boundary theories [27], whose entanglement purifies each boundary

thermal density matrix and underlies the connected structure of the bulk spacetime [28].

B. Schwarzian dynamics

Putting the bulk term of Eq. (1) on shell selects the background manifold and reduces the JT action to the GHY boundary term [9,10,13]. Performing the gravitational path integral is then equivalent to summing over AdS_2 manifolds with different boundaries. To study the dynamics of these boundaries, we cut the manifold along a constant value of the dilaton in an arbitrary coordinate system.

Varying Eq. (1) with respect to the metric results in the vacuum equations of motion for the dilaton field:

$$\nabla_\mu \nabla_\nu \Phi - g_{\mu\nu} \nabla^2 \Phi + g_{\mu\nu} \Phi = 0. \quad (8)$$

Solving Eq. (8) in Poincaré coordinates [Eq. (3)], one obtains a divergent dilaton profile [13]:

$$\Phi(t_P, z_P) = \frac{A + B t_P + C(t_P^2 + z_P^2)}{z_P}, \quad (9)$$

where A , B , and C are arbitrary parameters of the theory. Clearly, the dilaton field diverges at the boundary $z_P = 0$. However, the theory can be regulated by introducing a near-boundary cutoff [10,13].

The gravitational path integral is defined for metrics which are asymptotically AdS_2 :

$$ds^2 = \frac{1}{z^2} (-dt^2 + dz^2) + \mathcal{O}(z^{-1}) \quad \text{as } z \rightarrow 0, \quad (10)$$

while imposing fixed divergent asymptotics for the dilaton field:

$$\Phi = \frac{a}{2z} + \mathcal{O}(z) \quad \text{as } z \rightarrow 0, \quad (11)$$

where a is the regularized value of the dilaton. Performing an arbitrary coordinate transformation $u_P(u, v)$, $v_P(u, v)$, the first boundary condition [Eq. (10)] forces $\partial_u v_P = \partial_v u_P = 0$ and $u_P(u, v) = v_P(u, v)$ at leading order in z as $z \rightarrow 0$:

$$\begin{aligned} u_P(u) &= F(u) + \mathcal{O}(z), \\ v_P(v) &= F(v) + \mathcal{O}(z). \end{aligned} \quad (12)$$

In the Poincaré patch, the boundary $z = \epsilon$ is therefore parametrized by

$$\begin{aligned} t_P(t) &= \frac{1}{2}(u_P + v_P) = F(t) + \mathcal{O}(\epsilon), \\ z_P(t) &= \frac{1}{2}(u_P - v_P) = \epsilon F'(t) + \mathcal{O}(\epsilon^2), \end{aligned} \quad (13)$$

with the induced metric $\sqrt{-h} = \frac{1}{\epsilon} + \mathcal{O}(\epsilon)$ and the extrinsic curvature trace

$$K = 1 + \epsilon^2 \{F(t), t\} + \mathcal{O}(\epsilon^4), \quad (14)$$

where the Schwarzian derivative is defined as

$$\{F, t\} := \frac{F'''(t)}{F'(t)} - \frac{3}{2} \left(\frac{F''(t)}{F'(t)} \right)^2. \quad (15)$$

The second boundary condition [Eq. (11)] forces the regulated value of the dilaton at this boundary to be $\Phi|_\partial = \frac{a}{2\epsilon} + \mathcal{O}(\epsilon)$.

Finally, the bulk gravitational action [Eq. (1)] simplifies to the Schwarzian action governing the dynamics of the boundary:

$$\begin{aligned} S_{\text{JT}}[g, \Phi] &= \frac{1}{8\pi G_N} \int_{\partial \mathcal{M}} \Phi|_\partial (K - 1) \sqrt{-h} dt \\ &= \frac{a}{16\pi G_N} \int \{F(t), t\} dt. \end{aligned} \quad (16)$$

III. SCALAR CORRELATORS IN JT GRAVITY

We are interested in calculating the vacuum correlation of a massless scalar field ϕ , including gravitational fluctuations. For simplicity, we assume a field decoupled from the dilaton, therefore ignoring backreaction.

A. Tree-level correlators

The AdS_2 vacuum is fully specified by the two-point function $G_{12} = \langle \phi(x_1) \phi(x_2) \rangle$. In Lorentzian spacetimes,

there exist various Green functions; we focus on the Hadamard function:

$$H_{12} = \langle \{ \phi(x_1), \phi(x_2) \} \rangle, \quad (17)$$

which is related to the Feynman propagator $G_{12}^F = i \langle T \phi(x_1) \phi(x_2) \rangle$ by $H_{12} = 2 \text{Im} G_{12}^F$.

The fixed background Hadamard function of the global AdS_2 vacuum is known to be [23]

$$H_{12} = -\frac{1}{2\pi} \ln \left| \frac{\cos(t_G - t'_G) - \cos(z_G - z'_G)}{\cos(t_G - t'_G) + \cos(z_G + z'_G)} \right|. \quad (18)$$

One can express this function in Rindler coordinates [Eqs. (4) and (5)] to obtain a thermal vacuum in either of the two BH exteriors. To evaluate the correlator in only one patch, we choose the coordinates corresponding to that patch. For example, in the left exterior,

$$H_{12}^{LL} = -\frac{1}{2\pi} \ln \left| \frac{\sinh \frac{\pi}{\beta} (u_{R,1}^I - u_{R,2}^I) \sinh \frac{\pi}{\beta} (v_{R,1}^I - v_{R,2}^I)}{\sinh \frac{\pi}{\beta} (v_{R,1}^I - u_{R,2}^I) \sinh \frac{\pi}{\beta} (u_{R,1}^I - v_{R,2}^I)} \right|. \quad (19)$$

Furthermore, one can also investigate the global vacuum correlation across both BH patches:

$$H_{12}^{LR} = -\frac{1}{2\pi} \ln \left| \frac{\cosh \frac{\pi}{\beta} (u_{R,1}^I + v_{R,2}^I) \cosh \frac{\pi}{\beta} (v_{R,1}^I + u_{R,2}^I)}{\cosh \frac{\pi}{\beta} (v_{R,1}^I + v_{R,2}^I) \cosh \frac{\pi}{\beta} (u_{R,1}^I + u_{R,2}^I)} \right|. \quad (20)$$

B. Gravity correlators

In the Rindler patch, the boundary reparametrization is given by $F(t_R) = \tanh \frac{\pi}{\beta} f(t_R)$ at leading order as $z_R \rightarrow 0$, where $f(t_R)$ is a function that satisfies

$$f(t_R + i\beta) = f(t_R) + i\beta, \quad f'(t_R) \geq 0. \quad (21)$$

The first relation enforces that f wraps around the Euclidean thermal circle once for each period β in $\tau = it_R$. The second requirement guarantees that f is monotone and thus orientation-preserving. These properties, after Wick rotation, identify f as an element of the diffeomorphism group of the circle, $\text{Diff}(\mathbb{S}^1)$, and justify interpreting it as a boundary time reparametrization.

Since the JT integral [Eq. (1)] always reduces to a boundary term after the background manifold is selected, integrating the Hadamard function over metrics is equivalent to an integration over the boundary function $f(t_R)$ [10,13,15,17–19]

$$\begin{aligned}\langle H_{12} \rangle &= \text{Re} \left(\frac{\int H_{12}[g] e^{iS_{JT}[g, \Phi]} \mathcal{D}g \mathcal{D}\Phi}{\int e^{iS_{JT}[g, \Phi]} \mathcal{D}g \mathcal{D}\Phi} \right) \\ &= \text{Re} \left(\frac{\int H_{12}[f] e^{iS_{JT}[f]} d\mu[f]}{Z(\lambda)} \right),\end{aligned}\quad (22)$$

where $Z(\lambda) = \int e^{iS_{JT}[f]} d\mu[f]$ is the vacuum normalization, $H_{12}[f]$ is the bulk Hadamard function evaluated on the metric parametrized by f , and $d\mu[f]$ is the measure of integration over the symplectic manifold $\text{Diff}(\mathbb{S}^1)/SL(2, \mathbb{R})$ [29].

After Wick-rotating the boundary action $iS_{JT} \rightarrow I$, rescaling the boundary circle to $\hat{\tau} = \frac{2\pi}{\beta}\tau$ and $\hat{f}(\hat{\tau}) = \frac{2\pi}{\beta}f(\tau)$, and evaluating the Pfaffian measure $d\mu[\hat{f}]$ in terms of a Grassmann field ψ [15,29], the bulk quantum gravity correlator evaluates to

$$\langle H_{12} \rangle = \text{Re} \left(\frac{\int H_{12}[\hat{f}] e^{-\frac{1}{2}I[\hat{f}, \psi]} \mathcal{D}\hat{f} \mathcal{D}\psi}{Z(\lambda)} \right), \quad (23)$$

where the Euclidean action is a functional of the boundary functions $\hat{f}(\hat{\tau})$ and $\psi(\hat{\tau})$:

$$I[\hat{f}, \psi] = \int_0^{2\pi} \frac{\dot{\hat{f}}^2}{\lambda^2 \hat{f}^2} - \frac{\dot{\hat{f}}^2}{\lambda^2} + \frac{\ddot{\psi} \psi}{\hat{f}^2} - \dot{\psi} \psi d\hat{\tau}, \quad (24)$$

with $\lambda^2 = \frac{8G_N \beta}{a}$.

In the following, we shall assume that the gravity coupling constant $\lambda \ll 1$ and fixed. This will justify a perturbative expansion in λ .

C. Perturbative expansion

The saddle of the action, Eq. (16), subject to the periodicity condition [Eq. (21)], is $\hat{f}_s(\hat{\tau}) = \hat{\tau}$ [10,13,18]. Evaluated at this saddle, the two-sided correlators $\langle H_{12} \rangle$ exhibit pure exponential decay at late times, reproducing the familiar semiclassical behavior that underlies the tension with the unitarity of the dual boundary theory [26]. From the point of view of the gravitational path integral, this reflects the fact that $H_{12}(\hat{f}_s) \rightarrow 0$ at large times. To determine whether residual correlations survive, one must therefore examine fluctuations around the saddle and include subleading terms in the steepest-descent expansion in the parameter λ , following the approach of Ref. [22]. If subleading corrections generate nondecaying contributions to $\langle H_{12} \rangle$, then late-time correlations do not exponentially vanish, despite the leading semiclassical decay.

For small λ , one can use the method of stationary phase to asymptotically expand the functional part of the integral [Eq. (23)] in powers of λ around the saddle point $\hat{f}_s$. In other words, we shall consider $\hat{f}(\hat{\tau}) = \hat{\tau} + \lambda\gamma(\hat{\tau}) + \dots$ at leading order in λ . By using the map of boundary correlators versus bulk correlators introduced in [19], in

Ref. [22] a linear ramp was found after an exponential decay for $\langle H_{12}^{LL} \rangle$ at quadratic order in λ . However, this behavior was in tension with the Schwarzian computation of Ref. [8]. Upon revisiting the calculation, we find that the discrepancy originates from a sign difference in the correlator $\langle \gamma(\hat{\tau}_1) \dot{\gamma}(\hat{\tau}_2) \rangle$ used in [22] (cf. Ref. [30]). Correcting this sign removes the linear growth in $\langle H_{12}^{LL} \rangle$ at order λ^2 , as the relevant contributions cancel. The sign structure reflects the time-reversal properties of the derivative of the two-point function. As we will show, once this correction is implemented, the linear ramp instead appears in the cross-boundary correlator H_{12}^{LR} connecting the two exteriors of the maximally extended geometry [Eq. (6)].

Expanding the action of Eq. (24) in powers of λ , we have

$$I[\hat{\tau} + \lambda\gamma(\hat{\tau}), \psi] = -\frac{2\pi}{\lambda^2} + I_0[\gamma, \psi] + \sum_{k=1}^{\infty} \lambda^k I_k[\gamma, \psi], \quad (25)$$

where

$$I_0[\gamma, \psi] = \int_0^{2\pi} \ddot{\gamma}^2 - \dot{\gamma}^2 + \ddot{\psi} \psi - \dot{\psi} \psi d\hat{\tau} \quad (26)$$

and

$$I_k[\gamma, \psi] = (-1)^k (k+1) \int_0^{2\pi} \dot{\gamma}^k (\ddot{\gamma}^2 + \ddot{\psi} \psi) d\hat{\tau}. \quad (27)$$

Similarly, one can expand the bulk correlator $H_{12}[\hat{\tau} + \lambda\gamma] = \sum_0^{\infty} \lambda^k H_k[\gamma]$ to obtain the full perturbative expansion:

$$\begin{aligned}\langle H_{12} \rangle &= \text{Re} \left(\frac{\int (\sum_{k=0}^{\infty} \lambda^k H_k) e^{\frac{\pi}{2\lambda^2} e^{-\frac{1}{2}I_0} e^{-\frac{1}{2} \sum_{n=1}^{\infty} \lambda^n I_n} \mathcal{D}\gamma}}{Z_{1\text{-loop}}(\lambda)} \right) \\ &= \sum_{k=0}^{\infty} \lambda^k \langle \Gamma_k \rangle,\end{aligned}\quad (28)$$

where

$$\begin{aligned}\Gamma_0 &= H_0, \\ \Gamma_1 &= \text{Re}(H_1 + H_0 I_1), \\ \Gamma_2 &= \text{Re} \left(H_2 + H_1 I_1 + \frac{1}{2} (I_1^2 - 3I_2) \right), \\ &\vdots\end{aligned}\quad (29)$$

and where we use the fact that the Schwarzian partition function is one-loop exact [29], so that the normalization $Z(\lambda)$ is given exactly by $Z_{1\text{-loop}}(\lambda)$. In particular, although the partition function receives no higher-loop corrections, the correlator itself is given by an infinite series in λ arising from the numerator expansion. For each order in λ , $\langle \Gamma_k[\gamma, \psi] \rangle$ is a Gaussian integral involving correlations of

the perturbation [30]

$$\begin{aligned} \langle T\gamma(\hat{\tau}_1)\gamma(\hat{\tau}_2) \rangle &= \int \gamma(\hat{\tau}_1)\gamma(\hat{\tau}_2)e^{-I_0[\gamma]}\mathcal{D}\gamma \\ &= \frac{1}{2\pi} \left(1 + \frac{\pi^2}{6} - \frac{1}{2}(\pi - \hat{\tau})^2 - (\pi - \hat{\tau})\sin(\hat{\tau}) \right. \\ &\quad \left. + \frac{5}{2}\cos(\hat{\tau}) \right), \end{aligned} \quad (30)$$

where $\hat{\tau} = |\hat{\tau}_1 - \hat{\tau}_2|$. All odd orders of the expansion vanish, and one is left with $\langle \Gamma_2 \rangle$ as the first quantum correction.

IV. RESULTS

To justify the perturbative expansion, we fix $\lambda^2 = 0.1$. Furthermore, we are interested in investigating the large-time behavior of the gravitational path integral of the correlators in Eqs. (19) and (20). For this, we fix a bulk point $z_{R,1} = z_{R,2} = z_R$, and set $t_{R,1} = 0$ and $t_{R,2} = t_R$. Finally, in order to compare the β dependence of the gravitationally dressed correlators in a geometrically well-defined way, we anchor the first bulk point at $t_P = 0$ and $z_P = 1$, and then find the corresponding z_R . For $\beta \leq \pi$, the chosen Rindler coordinate patch fails to cover the fixed Poincaré bulk point; we therefore restrict to $\beta > \pi$ for convenience. No physical singularity is associated with this $\beta = \pi$.

A. Bulk correlators

First, let us investigate the one-sided bulk correlator. As described above, we set $t_{R,1}^I = 0$, $t_{R,2}^I = t_R$, and $z_{R,1}^I = z_{R,2}^I = z_R = \frac{\beta}{\pi} \arctan(\frac{\pi}{\beta})$. One finds that this correlator decays in time at both orders λ^0 and λ^2 , as can be seen directly from Figs. 2 and 3. This is in agreement with Ref. [8].

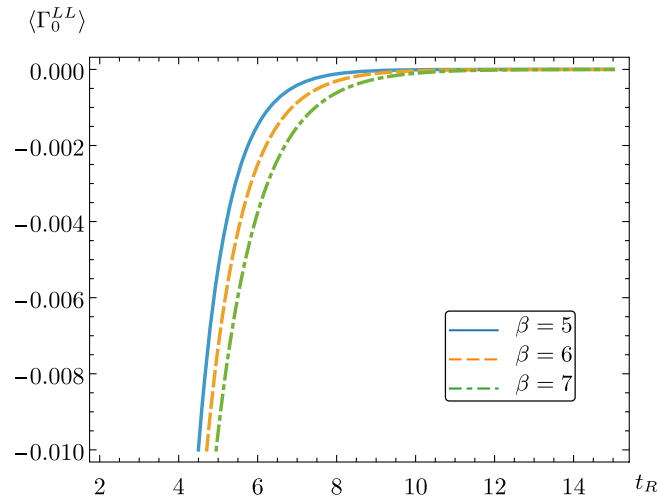


FIG. 2. One-sided bulk saddle correlator plotted for different values of β and at $z_P = 1$.

Specifically, taking the large-time limit yields

$$\begin{aligned} \langle \Gamma_0^{LL} \rangle &\sim -\frac{2}{\pi} e^{-\frac{2\pi}{\beta} t_R} \sinh^2\left(\frac{2\pi}{\beta} z_R\right), \quad t_R \rightarrow \infty, \\ \langle \Gamma_2^{LL} \rangle &\sim -\frac{8\pi}{\beta^2} t_R^2 e^{-\frac{2\pi}{\beta} t_R} \sinh^2\left(\frac{2\pi}{\beta} z_R\right), \quad t_R \rightarrow \infty. \end{aligned} \quad (31)$$

This behavior reflects the suppression of correlations between bulk points on the same side of the maximally extended geometry. The exponential decay is the standard semiclassical behavior of one-sided correlators in the eternal black hole geometry and is consistent with the analysis of Ref. [8]. In this approximation, the presence of the horizon gives rise to a thermal state, leading to late-time decay of one-sided correlators.

In contrast, the two-sided correlator develops a qualitatively different late-time structure. Setting $t_{R,1}^I = 0$, $t_{R,2}^I = t_R$, and $z_{R,1}^I = z_{R,2}^I = z_R = \frac{\beta}{\pi} \arctan(\frac{\pi}{\beta})$, the bulk two-sided correlator decays to zero for large time at tree level only, as shown in Fig. 4. At order λ^2 , the correlator asymptotes to a negative but constant value, as seen in Fig. 5.

Taking the large-time limit, one obtains

$$\begin{aligned} \langle \Gamma_0^{LR} \rangle &\sim \frac{2}{\pi} e^{-\frac{2\pi}{\beta} t_R} \sinh^2\left(\frac{2\pi}{\beta} z_R\right), \quad t_R \rightarrow \infty, \\ \langle \Gamma_2^{LR} \rangle &\sim -\frac{2z_R}{\pi\beta} \cosh\left(\frac{2\pi}{\beta} z_R\right) \sinh\left(\frac{2\pi}{\beta} z_R\right), \quad t_R \rightarrow \infty. \end{aligned} \quad (32)$$

Although the leading saddle contribution $\langle \Gamma_0^{LR} \rangle$ decays exponentially with time [Eq. (32) and Fig. 4], the next-to-leading correction $\langle \Gamma_2^{LR} \rangle$ alters the late-time behavior, generating a linear ramp that drives the correlator toward a constant, power-suppressed plateau (Fig. 5). Thus, non-vanishing correlations across the two exterior patches persist within the perturbative expansion.

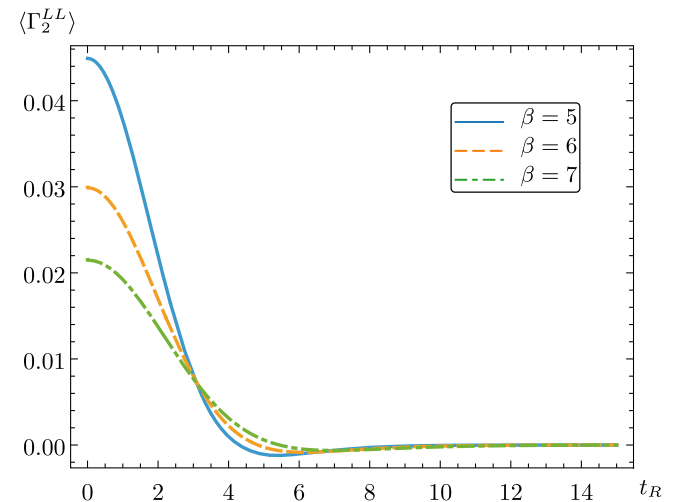


FIG. 3. First gravitational correction to the one-sided bulk correlator plotted for different values of β and at $z_P = 1$.

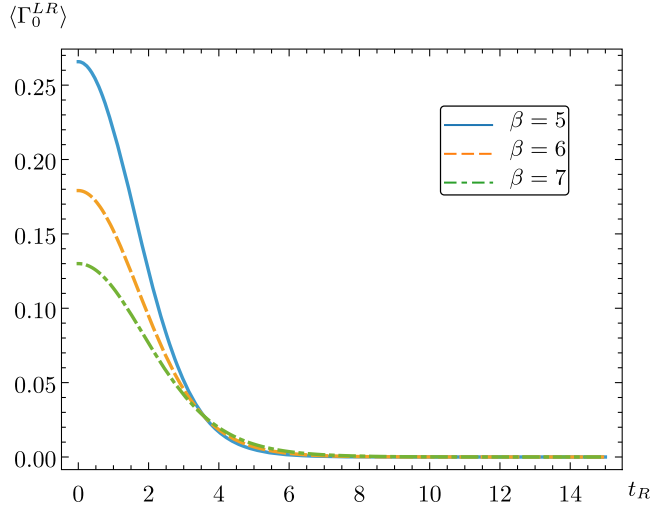


FIG. 4. Two-sided bulk saddle correlator plotted for different values of β and at $z_P = 1$.

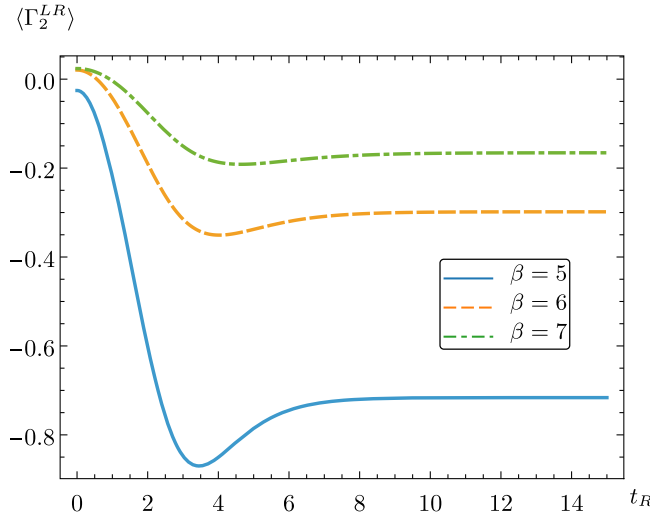


FIG. 5. First gravitational correction to the two-sided bulk correlator plotted for different values of β and at $z_P = 1$.

Combining both contributions, the bulk two-sided gravity correlator $\langle H_{12}^{LR} \rangle \sim \langle \Gamma_0^{LR} \rangle + \lambda^2 \langle \Gamma_2^{LR} \rangle$ exhibits a dip-ramp-plateau structure, as shown in Fig. 6. In this framework, a ramp and plateau arise already at order λ^2 , providing a perturbative realization of the dip-ramp-plateau behavior entirely within the disk topology, without invoking additional topologies or genuinely nonperturbative e^{-S} effects.

Finally, defining t_{dip} as the time at which $\langle H_{12}^{LR} \rangle$ reaches its minimum, we find that t_{dip} scales linearly with β when the fixed bulk point lies well within the Rindler patch, as shown in Fig. 7. Linear dependence of the dip time on the inverse temperature is likewise encountered in nonperturbative analyses of spectral observables [18], although in that context additional entropy-dependent scales also enter. The same plot illustrates the increasing redshift of the

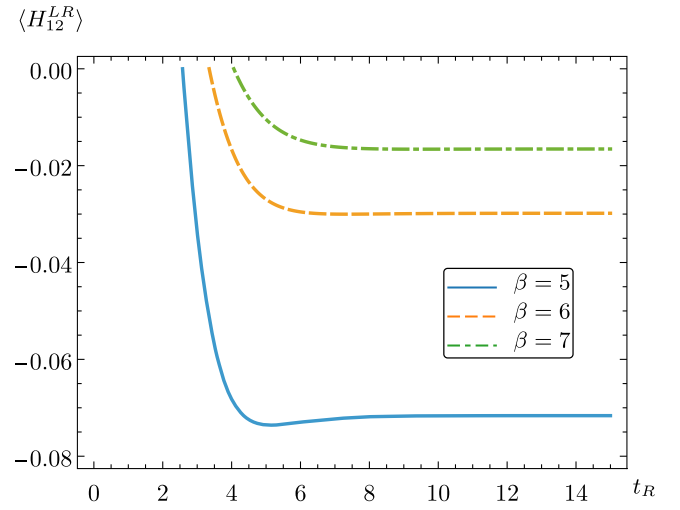


FIG. 6. Two-sided bulk gravity correlator plotted for different values of β , with $\lambda^2 = 0.1$, and at $z_P = 1$.

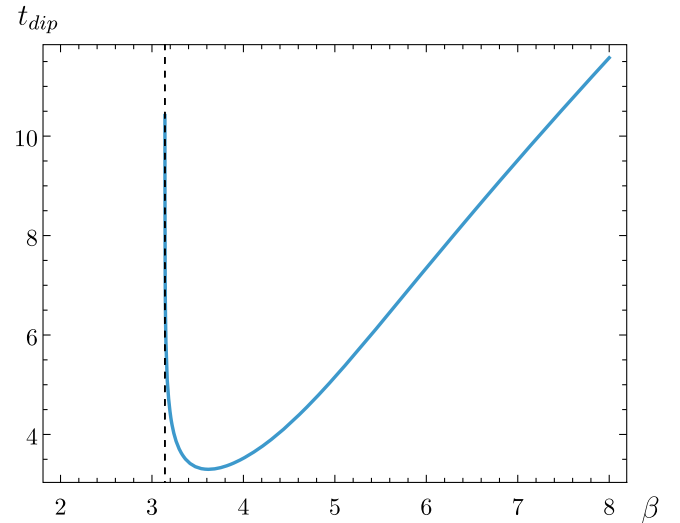


FIG. 7. Minimum of the two-sided bulk gravity correlator plotted as a function of β for $\lambda^2 = 0.1$, and at $z_P = 1$.

Rindler time coordinate t_R as the fixed bulk point approaches the boundary of the chosen Rindler patch near $\beta = \pi$.

B. Boundary correlators: Relation to AdS/CFT

To obtain the boundary CFT correlators, in AdS/CFT one takes the near-boundary limit of the bulk two-point functions. For a massless scalar field in AdS_2 , the scaling dimension is $\Delta = 1$; hence, the boundary correlators are extracted by taking the limit [31]

$$\langle \tilde{H}_{12} \rangle = \lim_{z \rightarrow 0} (z^{-2} \langle H_{12} \rangle). \quad (33)$$

Applying this to Eqs. (19) and (20) results in the boundary saddle CFT correlators:

$$\begin{aligned}\tilde{H}_{12}^{LL} &= -\frac{2\pi}{\beta^2} \frac{1}{\sinh^2\left(\frac{\pi}{\beta}(t_{R,1}^I - t_{R,2}^I)\right)}, \\ \tilde{H}_{12}^{LR} &= \frac{2\pi}{\beta^2} \frac{1}{\cosh^2\left(\frac{\pi}{\beta}(t_{R,1}^I + t_{R,2}^I)\right)}.\end{aligned}\quad (34)$$

One expects the boundary correlators to retain the features of their corresponding bulk functions. Setting $t_{R,1}^I = 0$ and $t_{R,2}^I = t_{R,2}^I = t_R$, and taking the large-time limit yields

$$\tilde{H}_{12}^{LL} \sim -\frac{8\pi}{\beta^2} e^{-\frac{2\pi}{\beta}t_R}, \quad t_R \rightarrow \infty, \quad (35)$$

$$\tilde{H}_{12}^{LR} \sim \frac{8\pi}{\beta^2} e^{-\frac{2\pi}{\beta}t_R}, \quad t_R \rightarrow \infty, \quad (36)$$

$$\langle \tilde{\Gamma}_2^{LL} \rangle \sim -\frac{32\pi^3}{\beta^4} t_R^2 e^{-\frac{2\pi}{\beta}t_R}, \quad t_R \rightarrow \infty, \quad (37)$$

$$\langle \tilde{\Gamma}_2^{LR} \rangle \sim -\frac{4}{\beta^2}, \quad t_R \rightarrow \infty. \quad (38)$$

Clearly, the one-sided correlator decays exponentially for large time at both orders λ^0 and λ^2 . However, the cross-boundary correlator decays exponentially at order λ^0 only. Furthermore, one observes that the height of the plateau goes to zero as $\beta \rightarrow \infty$, in accordance with Ref. [32].

Plotting the full expression for the next-to-leading order of both correlators, one obtains the same dip-ramp-plateau features, shown in Figs. 8 and 9. The crucial difference is that the height of the plateau of the one-sided correlator is zero for all values of β , whereas the height of the plateau of the cross-boundary correlator is β -dependent.

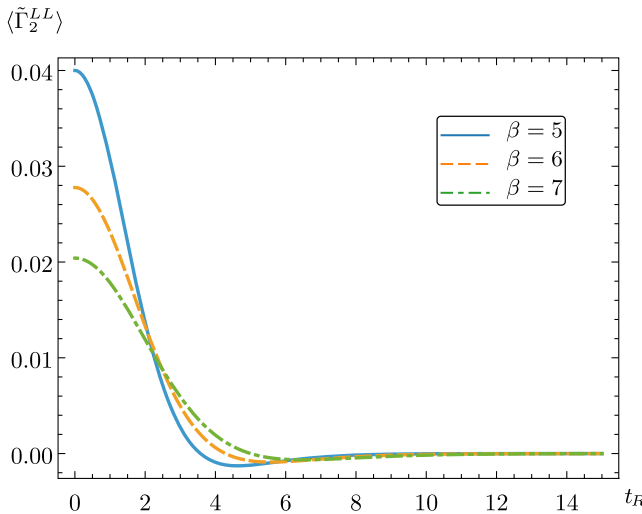


FIG. 8. First gravitational correction to the one-sided boundary correlator plotted for different values of β .

On the other hand, plotting both orders of the boundary gravity correlators, one clearly observes the complete dip-ramp-plateau structure emerging only for the cross-boundary correlator. As shown in Fig. 10, the one-sided boundary correlator decays exponentially at all temperatures, reflecting the loss of correlations across the horizon and reproducing the expected semiclassical behavior. In contrast, Fig. 11 shows the cross-boundary correlator, displaying a clear dip-ramp-plateau structure whose amplitude and timescale depend on β . The initial dip corresponds to the semiclassical decay, while the subsequent ramp and saturation indicate the reemergence of correlations between the two boundaries.

Within the perturbative framework, this late-time plateau provides a realization of the qualitative dip-ramp-plateau pattern familiar from studies of spectral observables in

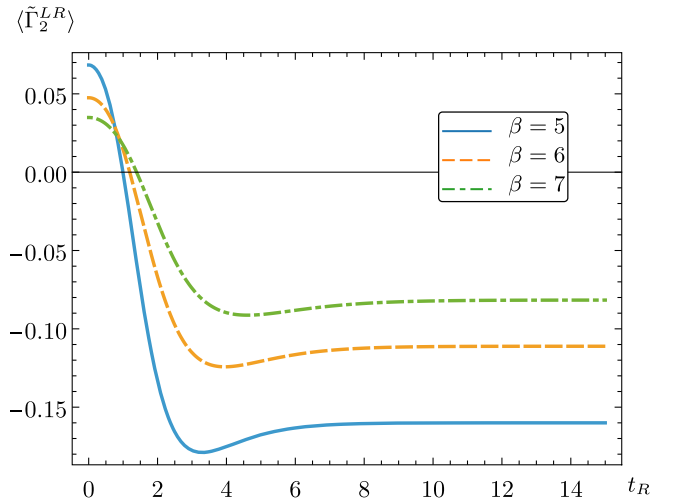


FIG. 9. First gravitational correction to the cross-boundary correlator plotted for different values of β .

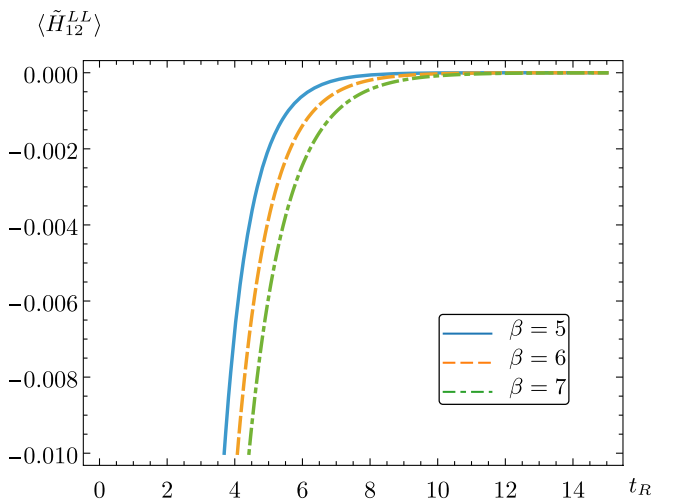


FIG. 10. One-sided boundary gravity correlator plotted for different values of β .

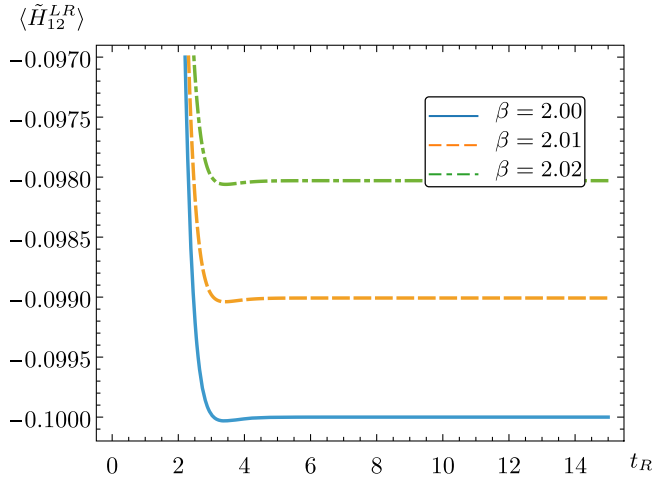


FIG. 11. Cross-boundary gravity correlator plotted for different values of β .

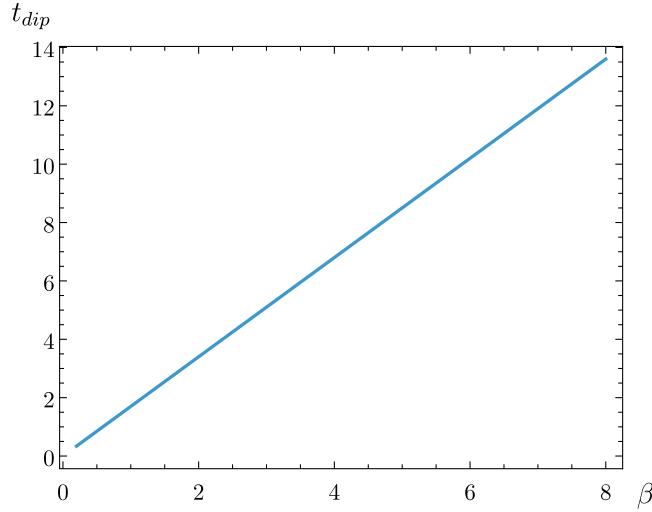


FIG. 12. Minimum of the cross-boundary gravity correlator plotted as a function of β for $\lambda^2 = 0.1$.

SYK [12,14]; however, here it arises entirely from fluctuations around the principal saddle rather than from genuinely nonperturbative effects.

Finally, defining t_{dip} as the time at which $\langle \tilde{H}_{12}^{LR} \rangle$ reaches its minimum, we find that t_{dip} scales linearly with β , as shown in Fig. 12. Linear dependence of the dip time on the inverse temperature is likewise encountered in nonperturbative analyses of spectral observables [18], although in that context, additional entropy-dependent scales also enter.

V. PERTURBATIVE VS NONPERTURBATIVE

It is useful to clarify the relation between the perturbative and nonperturbative regimes in the language of the gravitational path integral. Schematically, the partition function admits a saddle-point expansion of the type

$$Z \sim e^{-S_{\text{dom}}/\lambda^2} \sum_{n=0}^{\infty} a_n \lambda^n + \sum_i e^{-S_i/\lambda^2} \sum_{n=0}^{\infty} b_{n,i} \lambda^n,$$

where S_{dom} denotes the action of the dominant saddle (the disk geometry in our case), and the S_i 's correspond to subleading saddles exponentially suppressed in $1/\lambda^2$ capturing genuinely nonperturbative contributions—e.g., higher-topology configurations.

In our analysis, the dip-ramp-plateau structure arises entirely from the perturbative series around the dominant disk saddle. The entropy-suppressed effects associated with additional saddles are not required to generate the ramp and plateau at the order we considered. This does not exclude the eventual relevance of nonperturbative contributions at parametrically late times; rather, it shows that qualitatively nontrivial late-time structure is already encoded in the perturbative sector of the theory.

From the perspective of the dual description, this separation between power-suppressed and entropy-suppressed contributions maps directly onto the large- N expansion of the partition function. In the SYK/JT correspondence, the gravitational coupling in the bulk scales as $\lambda^2 \sim G_N \sim 1/N$ in the large- N limit. For the particular cross-boundary observable we study, the natural dual is the two-sided correlator in a pair of maximally entangled SYK systems prepared in the thermofield-double state. In this setting, the plateau we find corresponds to a contribution of order $1/N$ in the boundary theory, whereas entropy-controlled nonperturbative effects are expected to scale as $e^{-S} \sim e^{-N}$ [12]. The cross-boundary plateau obtained from the disk saddle is thus parametrically larger than genuinely nonperturbative contributions, and therefore dominates the late-time behavior within the regime of validity of the expansion.

More broadly, our results suggest that the purely semiclassical picture, characterized by exponential decay of all correlators in a fixed black hole background, is too naive. Once subleading corrections in the perturbative expansion are included, qualitatively new late-time structures can emerge. This indicates that essential ingredients relevant to the information paradox may already be encoded in controlled quantum fluctuations around a single saddle geometry [22]. It would be interesting to compute the cross-boundary two-point function in a pair of maximally entangled SYK systems prepared in the thermofield-double state, beyond leading order in $1/N$, to get a direct test of the perturbative mechanism identified here in the context of AdS/CFT. We leave this interesting question for future work.

VI. CONCLUSION

In this work, we have shown that a dip-ramp-plateau structure emerges perturbatively from bulk correlators in Jackiw-Teitelboim gravity within the disk topology, without invoking topology change or genuinely nonperturbative saddles. By evaluating the gravitationally dressed

Hadamard function through the next-to-leading term in the steepest-descent expansion of the Schwarzian path integral, we found that correlators evaluated in a single black hole exterior region decay exponentially at all orders considered, reproducing the semiclassical suppression. In contrast, correlators linking the two exteriors display a qualitatively different behavior: after the initial semiclassical decay, the $O(\lambda^2)$ correction generates a linear ramp followed by saturation to a constant value. At this perturbative order, both bulk and boundary observables exhibit the complete dip-ramp-plateau pattern. Higher orders in λ^2 eventually lead to exponential growth, indicating that the perturbative series is asymptotic. This is similar to the case of the same-side correlators studied in Ref. [33]. The implications of this asymptotic structure are left for future research.

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DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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- [1] S. W. Hawking, *Nature (London)* **248**, 30 (1974).
 - [2] S. W. Hawking, *Commun. Math. Phys.* **43**, 199 (1975); **46**, 206(E) (1976).
 - [3] S. W. Hawking, *Phys. Rev. D* **14**, 2460 (1976).
 - [4] D. N. Page, *Phys. Rev. Lett.* **71**, 3743 (1993).
 - [5] D. N. Page, *Phys. Rev. Lett.* **71**, 1291 (1993).
 - [6] R. Jackiw, *Nucl. Phys.* **B252**, 343 (1985).
 - [7] C. Teitelboim, *Phys. Lett.* **126B**, 41 (1983).
 - [8] J. Maldacena, D. Stanford, and Z. Yang, *Prog. Theor. Exp. Phys.* **(2016)** 12C104.
 - [9] A. Almheiri and J. Polchinski, *J. High Energy Phys.* **11** (2015) 014.
 - [10] J. Engelsöy, T. G. Mertens, and H. Verlinde, *J. High Energy Phys.* **07** (2016) 139.
 - [11] S. Sachdev, *Phys. Rev. Lett.* **105**, 151602 (2010).
 - [12] J. S. Cotler, G. Gur-Ari, M. Hanada, J. Polchinski, P. Saad, S. H. Shenker, D. Stanford, A. Streicher, and M. Tezuka, *J. High Energy Phys.* **05** (2017) 118; **09** (2018) 2.
 - [13] J. Maldacena and D. Stanford, *Phys. Rev. D* **94**, 106002 (2016).
 - [14] P. Saad, S. H. Shenker, and D. Stanford, *arXiv:1806.06840*.
 - [15] P. Saad, S. H. Shenker, and D. Stanford, *arXiv:1903.11115*.
 - [16] P. Saad, *arXiv:1910.10311*.
 - [17] D. Stanford and E. Witten, *Adv. Theor. Math. Phys.* **24**, 1475 (2020).
 - [18] T. G. Mertens and G. J. Turiaci, *Living Rev. Relativity* **26**, 4 (2023).
 - [19] A. Blommaert, T. G. Mertens, and H. Verschelde, *J. High Energy Phys.* **09** (2019) 060.
 - [20] D. Marolf and H. Maxfield, *J. High Energy Phys.* **08** (2020) 044.
 - [21] P. Saad, S. H. Shenker, D. Stanford, and S. Yao, *J. High Energy Phys.* **09** (2024) 133.
 - [22] C. Germani, *Phys. Rev. D* **106**, 066018 (2022).
 - [23] M. Spradlin and A. Strominger, *J. High Energy Phys.* **11** (1999) 021.
 - [24] U. H. Danielsson, E. Keski-Vakkuri, and M. Kruczenski, *J. High Energy Phys.* **01** (1999) 002.
 - [25] W. G. Unruh, *Phys. Rev. D* **14**, 870 (1976).
 - [26] J. M. Maldacena, *J. High Energy Phys.* **04** (2003) 021.
 - [27] J. Maldacena and L. Susskind, *Fortschr. Phys.* **61**, 781 (2013).
 - [28] S. Ryu and T. Takayanagi, *J. High Energy Phys.* **08** (2006) 045.
 - [29] D. Stanford and E. Witten, *J. High Energy Phys.* **10** (2017) 008.
 - [30] Y.-H. Qi, Y. Seo, S.-J. Sin, and G. Song, *Phys. Rev. D* **99**, 066001 (2019).
 - [31] E. Witten, *Adv. Theor. Math. Phys.* **2**, 253 (1998).
 - [32] T. G. Mertens, G. J. Turiaci, and H. L. Verlinde, *J. High Energy Phys.* **08** (2017) 136.
 - [33] L. Griguolo, J. Papalini, and D. Seminara, *J. High Energy Phys.* **05** (2021) 140.