

Evaporation of large black holes in AdS

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Abstract. The AdS/CFT correspondence offers a new perspective on the long-standing black hole information paradox. However, to be able to use the available gauge/gravity machinery one is forced to consider so-called ‘large’ black holes in AdS, and these objects are thermodynamically stable – they do not evaporate. We describe a simple toy model that allows large AdS black holes to decay, by coupling the emitted radiation to an external scalar field propagating in an auxiliary space. This effectively changes the properties of the boundary of AdS, making it partly absorbing. We demonstrate that the evaporation process never ceases by explicitly presenting (a) the transmission coefficient for a wave scattering from the bulk into auxiliary space and (b) the greybody factor for a black 3-brane in an AdS background. Therefore, the model provides an interesting framework to address the information paradox using AdS/CFT techniques.

1. Introduction

Despite many efforts over more than thirty years, the black hole information paradox still remains to be elucidated. This long-standing issue is directly related to foundational concepts in physics, namely *unitarity* and *locality*, and its resolution is expected to hold the key to a formulation of quantum gravity.

The famous information loss problem was posed by Hawking in 1976 [1] shortly after his revolutionary discovery that black holes radiate [2]. Furthermore, according to Hawking’s original computation the nature of the emitted radiation is that of a featureless blackbody spectrum [2]. It was soon realized that information seemed to be lost in the course of formation and consequent evaporation of black holes. This becomes evident when one considers the formation of a black hole by gravitational collapse of an initial matter distribution in a pure quantum state. Assuming a black hole does form and that it decays completely according to Hawking’s law, the final state would be a thermal (i.e., maximally mixed) state, leading to an obvious violation of unitarity. It should be noted that one does not need a black hole to evaporate completely to arrive at a paradox: information must start being recovered in the Hawking particles as soon as the black hole radiates half its initial entropy if the evolution is unitary [3]. So one just needs to start with a large enough black hole.

The AdS/CFT correspondence [4, 5, 6], holographically relating gravitational dynamics in an AdS background with a gauge theory living on the boundary of AdS, offers a new perspective on this old problem. Shortly after the discovery of this duality it was pointed out that the manifest unitarity of the dual field theory should prevent any information loss of occurring [7]. Of course, this argument by itself does not resolve the paradox; one would still need to explain what went

wrong in Hawking's argument. But the power of the gauge/gravity duality may well shed some light and provide indications for its resolution.

In order to take advantage of the AdS/CFT correspondence we should consider a black hole in AdS. These objects come in two classes: *small* and *large*. As in the asymptotically flat background case, small AdS black holes are thermodynamically unstable (they have negative specific heat) and do not have a clear interpretation in terms of the dual CFT [8, 9]. On the other hand, large AdS black holes have positive specific heat and are classically stable [10]. These objects are dual to a high temperature thermal state in the CFT [8]. To make progress exploiting the AdS/CFT correspondence we should then consider a large black hole in AdS. However, as it stands there is a clear obstacle to the formulation of the information paradox using large AdS black holes: these objects are stable and they do not evaporate!

The main purpose of the evaporon model [11, 12] is to provide a *simple* setting in which large black holes in AdS can decay, allowing to monitor the evaporation process employing the AdS/CFT duality and ultimately addressing the information paradox.

2. The evaporon model

The thermodynamical stability of large AdS black holes can be traced back to the total reflectivity of the time-like boundary of AdS. All the radiation emitted by the black hole is reflected at the boundary and is reabsorbed before the black hole evaporates. The system approaches an equilibrium state consisting of a black hole with temperature T_H and the surrounding Hawking radiation at the same temperature. The rate at which the black hole emits is compensated by the rate of absorption.

In order to allow large AdS black holes to evaporate, a simple model in the context of the AdS₅/CFT₄ correspondence was constructed in [11]. Recall that the gauge/gravity correspondence relates the CFT generating functional for a gauge invariant single trace operator \mathcal{O} to extrema of the type IIB supergravity action with prescribed boundary conditions for the bulk field Φ dual to the operator \mathcal{O} . The non-normalizable solution for the supergravity field determines the source term for the dual operator on the CFT side. In the proposed model, we work at the level of the supergravity action and in the large N limit in which all the bulk fields decouple so we can consider Φ independently.

- First, one attaches a (1+1)-dimensional flat spacetime $H = \{(t, z) | z \geq 0\}$ to the boundary of AdS₅ and inserts there a scalar field $\sigma(t, z)$ with a standard kinetic term. This field was dubbed the *evaporon*.
- Next, we add an interaction term that couples the evaporon to the supergravity field Φ (hereafter referred to as the dilaton, for simplicity). The interaction is localized on the boundary of AdS₅ (which is 'identified' with the boundary of the additional space H). This is a four-dimensional space parametrized by the timelike coordinate t and the non-holographic spatial coordinates \mathbf{x} . Typically, to regulate the computations one cuts off the AdS space¹ at a holographic coordinate $y = \varepsilon$ and the limit $\varepsilon \rightarrow 0$ is taken in the end. The interaction is chosen to take place at $y = \varepsilon$ and in such a way as to enforce fall-off conditions for the dilaton in accordance with the non-normalizable solution for a massless scalar field in AdS₅.
- Finally, one is required to add specific counterterms in the evaporon action to renormalize the theory. These counterterms were judiciously computed in [11] and their inclusion is crucial to obtain finite results when one takes the cutoff $\varepsilon \rightarrow 0$.

¹ We shall adopt coordinates suited for a Poincaré patch, for which the AdS metric can be expressed as $ds^2 = R^2/y^2 [-dt^2 + d\mathbf{x}^2 + dy^2]$. The boundary of AdS lies at $y = 0$.

According to the above discussion, our model is described by the following action for the evaporon-dilaton system (see reference [11] for more details):

$$\begin{aligned}
 S[\Phi, \sigma] &= S_\Phi[\Phi] + S_\sigma[\sigma] + S_{int}[\Phi, \sigma] + S_{c.t.}[\sigma], \quad (1) \\
 S_\Phi[\Phi] &= -\frac{1}{2\kappa^2} \int_{\text{AdS}} dt d^3\mathbf{x} dy \sqrt{-g} \frac{1}{2} g^{ab} \partial_a \Phi \partial_b \Phi, \\
 S_\sigma[\sigma] &= -\int_H dt dz \frac{1}{2} [-(\partial_t \sigma)^2 + (\partial_z \sigma)^2], \\
 S_{int}[\Phi, \sigma] &= \lambda \int_{\partial\text{AdS}} dt d^3\mathbf{x} \sqrt{-h} \Phi(t, \mathbf{x}, \varepsilon) \sigma(t, 0),
 \end{aligned}$$

where g_{ab} denotes the bulk metric, g is its determinant and h represents the determinant of the metric induced on the interaction surface $\partial\text{AdS} \equiv \{(t, \mathbf{x}, y) | y = \varepsilon\}$. Here, λ plays the role of a coupling constant and κ^2 is related to the 5-dimensional Newton constant G_5 through $\kappa^2 = 8\pi G_5$. The necessary counterterms mentioned previously, which are also localized on the interaction surface $\partial H = \{(t, z) | z = 0\}$, turn out to be

$$S_{c.t.}[\sigma] = \kappa^2 \lambda^2 R^5 V_3 \int_{\partial H} dt \left[-\frac{1}{4\varepsilon^4} \sigma^2 - \frac{1}{24\varepsilon^2} \dot{\sigma}^2 + \frac{1}{128} \ln\left(\frac{\varepsilon^2}{\mu^2}\right) \ddot{\sigma}^2 \right], \quad (2)$$

where $V_3 = \int d^3\mathbf{x}$ is the volume of the transverse space and μ is an IR cutoff needed to regulate the expression for the counterterm. Had we been considering AdS in global coordinates, the transverse space would have been S^3 and presumably the cutoff would be related to the AdS scale R .

2.1. The transmission coefficient

The equations of motion for the evaporon-dilaton system can be easily derived from the action (1), supplemented with the counterterms (2):

$$y^2 \partial_y^2 \Phi - 3y \partial_y \Phi - y^2 \partial_t^2 \Phi = -2\kappa^2 \lambda R \varepsilon \delta(y - \varepsilon) \sigma(t, 0), \quad (3)$$

$$-\partial_t^2 \sigma + \partial_z^2 \sigma = -\delta(z) \lambda V_3 R^4 \varepsilon^{-4} \Phi(t, \varepsilon) + \delta(z) 2\kappa^2 \lambda^2 R^5 V_3 \left[\frac{\sigma}{4\varepsilon^4} - \frac{\partial_t^2 \sigma}{24\varepsilon^2} - \frac{\partial_t^4 \sigma}{128} \ln\left(\frac{\varepsilon^2}{\mu^2}\right) \right]. \quad (4)$$

By construction, the evaporon only couples to dilaton zero-modes in transverse space, i.e., Φ is independent of the coordinates \mathbf{x} in the above equations. This is one of the simple features of the model.

The solutions of equations (3) and (4) can be immediately written. Imposing normalizability restricts the form of the dilaton wavefunction in the region near the boundary ($y < \varepsilon$), whereas the solutions for the evaporon are simply plane waves:

$$\Phi(t, y) = e^{-i\omega t} [\beta y^2 J_2(\omega y) + \gamma y^2 Y_2(\omega y)] + h.c. \quad \text{for } y > \varepsilon, \quad (5)$$

$$\Phi(t, y) = \alpha e^{-i\omega t} y^2 J_2(\omega y) + h.c. \quad \text{for } y < \varepsilon, \quad (6)$$

$$\sigma(t, z) = A e^{-i\omega(t+z)} + B e^{-i\omega(t-z)} + h.c. \quad (7)$$

The problem is now reduced to the scattering of waves in one dimension. Assuming an outgoing dilaton wave is incident on the interaction surface and $A = 0$, i.e., there is no evaporon wave coming from infinity and therefore it is purely outgoing, what fraction of energy is transmitted to the evaporon in the auxiliary space? The continuity of the dilaton at $y = \varepsilon$ and the discontinuity of the first derivatives of Φ and σ (determined by the right hand sides of equations (3) and (4))

can be used to relate the coefficients α, β, γ and B . The final result of this calculation yields the following formula for the transmission coefficient:

$$|\mathcal{T}|^2 = \frac{4 \operatorname{Im}(\beta\gamma^*)}{|\beta|^2 + |\gamma|^2 + 2 \operatorname{Im}(\beta\gamma^*)} = \frac{2}{1 + \frac{1}{4\pi} \left(\frac{\omega_4}{\omega}\right)^3 + \pi \left(\frac{\omega}{\omega_4}\right)^3 \left[1 + \frac{4}{\pi^2} (\ln(\mu\omega))^2\right]}, \quad (8)$$

where we have defined an effective coupling constant $\omega_4 \equiv 8(2\lambda^2\kappa^2R^5V_3)^{-1/3}$. The result (see Figure 1) depends logarithmically on the scale μ but all the dependence on the cutoff ε has canceled out, as necessary, after having taken the limit $\varepsilon \rightarrow 0$. Nevertheless, for this to be possible it was vital to include the counterterms in the action: the transmission coefficient would not be finite without their addition.

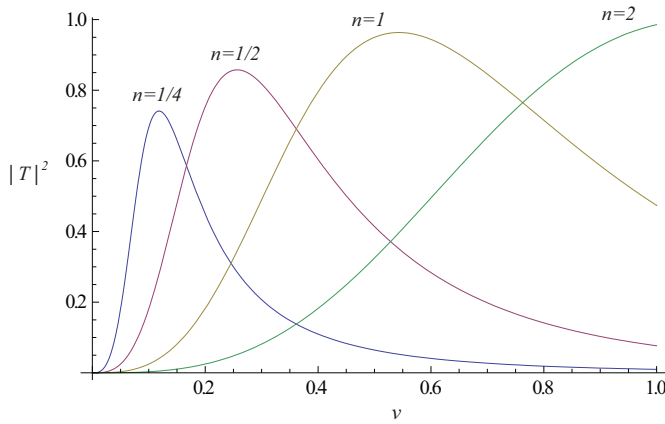


Figure 1. The transmission coefficient as a function of the frequency in units of the IR cutoff scale $\nu \equiv \mu\omega$. The result is shown for four different values of the effective coupling constant, equally in units of the IR cutoff, $n \equiv \mu\omega_4$.

3. Large black holes in AdS and their evaporation

The evaporon model, and in particular the previous computation of the transmission coefficient, were presented in the planar limit of AdS, in which the boundary becomes $\mathbb{R} \times \mathbb{R}^3$. To preserve the transverse translational symmetry it is natural to consider a black hole possessing a horizon with three flat directions, i.e., a black 3-brane in AdS₅. In the following subsection we shall outline the computation of the greybody factors for this geometry, originally performed in [12], to which the reader is referred for more details. An adaptation of the techniques used in [13] can be employed to derive an expression for the greybody factor in the low frequency regime.

The metric describing a black 3-brane with AdS₅ × S⁵ asymptotics was found in [14] as the decoupling limit of a stack of D3-branes:

$$ds^2 = H^{-1/2}(r) [-f(r)dt^2 + d\mathbf{x}^2] + H^{1/2}(r) [f^{-1}(r)dr^2 + r^2d\Omega_5^2], \quad (9)$$

where

$$H(r) = \left(\frac{R}{r}\right)^4, \quad f(r) = 1 - \left(\frac{r_H}{r}\right)^4. \quad (10)$$

This is a solution of type IIB string theory in 10 dimensions of the form $X \times S^5$, where the five-sphere has constant radius R . The five dimensional manifold X features a regular horizon at $r = r_H$ and in the limit $r \rightarrow \infty$ it approaches the metric of AdS₅ in Poincaré coordinates². Since the S^5 plays no role in the following we shall drop it from now on.

² The relation between the radial coordinate r and the previously used coordinate y is simply $y = R^2/r$.

The Hawking temperature T_H , as well as the energy and entropy per unit 3-volume of the black 3-brane are given by

$$T_H = \frac{r_H}{\pi R^2}, \quad \frac{E}{V_3} = \frac{3r_H^4}{2\kappa^2 R^5}, \quad \frac{S}{V_3} = \frac{2\pi r_H^3}{\kappa^2 R^3}. \quad (11)$$

These black objects have positive specific heat [10] and therefore do not evaporate if the boundary of AdS is totally reflective. In this sense, these black branes fall in the category of the *large* black holes in AdS.

3.1. The greybody factor

Recall that greybody factors correct the otherwise perfect blackbody emission spectrum of a black hole, by accounting for the fact that the Hawking radiation needs to traverse frequency-dependent potential barriers originated by the background geometry in order to reach asymptotic infinity.

Consider then a massless, minimally coupled scalar field Φ propagating in the geometry of the black brane, obtained from the metric (9) by dropping the five-sphere. It is convenient to Fourier transform and study modes with definite frequency ω and wavevector \mathbf{k} separately, so take $\Phi(t, \mathbf{x}, r) = e^{i\omega t - i\mathbf{k}\cdot\mathbf{x}}\phi_{\omega, \mathbf{k}}(r)$. The equation of motion for the dilaton wavefunction in the above geometry can be cast in Schrödinger-like form:

$$[\partial_x^2 + \omega^2 - V(r)](r^{3/2}\phi) = 0, \quad (12)$$

where the potential V is given by

$$V(r) = \frac{15}{4} \frac{r^2}{R^4} f(r)^2 + 6 \frac{r_H^4}{r^2 R^4} f(r) + m^2 \omega^2 f(r), \quad (13)$$

m is defined by $\mathbf{k}^2 = m^2 \omega^2$ (taking values between 0 and 1) and the tortoise coordinate x is determined (up to an additive constant) by $dx = H^{1/2}(r) f^{-1}(r) dr$.

Focusing on the near-horizon region ($r \simeq r_H$ and $V(r) \ll \omega^2$) the solution reduces to

$$\phi(r) = A e^{i\omega x}, \quad (14)$$

where we chose the linear combination that is purely ingoing at the horizon. In the asymptotic region the solutions may be expressed in terms of Hankel functions of degree two:

$$\phi(r) = C_1 (1 - m^2) \frac{\omega^2 R^4}{r^2} H_2^{(1)} \left(\sqrt{1 - m^2} \frac{\omega R^2}{r} \right) + C_2 (1 - m^2) \frac{\omega^2 R^4}{r^2} H_2^{(2)} \left(\sqrt{1 - m^2} \frac{\omega R^2}{r} \right). \quad (15)$$

The first and second terms are associated with the outgoing and ingoing parts of the wavefunction ϕ , respectively.

The flux per unit physical transverse area near the horizon can be shown to be

$$J_{hor} = \frac{r_H^3}{R^3} \omega |A|^2, \quad (16)$$

and the asymptotic flux per unit physical area may be decomposed as $J_{asy} = J_{in} - J_{out}$, where

$$J_{in} = \frac{2}{\pi} (1 - m^2)^2 \omega^4 R^3 |C_2|^2, \quad J_{out} = \frac{2}{\pi} (1 - m^2)^2 \omega^4 R^3 |C_1|^2. \quad (17)$$

The greybody factor is then defined [13] by the ratio of the (ingoing) flux at the horizon by the ingoing part of the asymptotic flux:

$$\Gamma_m(\omega) \equiv \frac{J_{hor}}{J_{in}} . \quad (18)$$

When the frequency is small compared to the temperature of the black brane $\omega \ll T_H$ an analytic expression can be obtained for the greybody factor by matching solutions valid in overlapping regions. The result obtained in [12] was

$$\Gamma_m(\omega) \simeq \frac{\pi}{2}(1 - m^2)^2 \tilde{\omega}^3 , \quad \text{for } \tilde{\omega} \ll 1/\pi . \quad (19)$$

Here we have used a dimensionless combination of the frequency, the AdS scale and the horizon radius:

$$\tilde{\omega} \equiv \omega \frac{R^2}{r_H} . \quad (20)$$

Besides m , this is the *only* other dimensionless parameter that enters the equation of motion.

A global analytic solution of equation (12) interpolating between the horizon and the boundary of AdS appears to be out of reach but one can proceed resorting to numerical methods. Essentially, one imposes the ingoing boundary condition at the horizon and then integrates the equation of motion all the way out to the boundary. We know that the solution has the behavior displayed in equations (14) and (15) when $r \rightarrow r_H$ and $r \rightarrow \infty$, respectively. To compute the greybody factor we just need to express the constants C_1 and C_2 characterizing the asymptotic solution in terms of the constant A which determines the solution near the horizon and employ expression (18).

The results are presented in figures 2 and 3 for three values of m . The greybody factor smoothly interpolates between 0 at low frequencies and its expected asymptotic value of 1 at high energies. In the low frequency regime the numerical results are in very good agreement with the analytic approximation (19). Reference [12] also analyzed the high frequency regime and, although not completely conclusive, there are clear indications that the greybody factor approaches the asymptotic value in a power-law fashion, specifically $1 - \Gamma_0(\tilde{\omega}) \sim \tilde{\omega}^{-8}$ for large frequencies (and $m = 0$). The dependence on m was also considered: the greybody factor flattens out for larger m and for $m = 1$ it vanishes identically, since this situation corresponds to propagation of the wave parallel to the brane.

3.2. The decay rate

Once the transmission coefficient and the greybody factor are known one can proceed to determine the rate of evaporation of the considered black brane in the context of the evaporon model. The asymptotic flux of evaporons per unit frequency interval is weighted by three terms: the blackbody spectrum, $[e^{\omega/T_H} - 1]^{-1}$, the greybody factor, $\Gamma(\omega)$, and the transmission coefficient at the dilaton-evaporon interface, $|\mathcal{T}|^2$. The energy emitted per unit time is given by

$$\frac{dE}{dt} = \int \frac{d\omega}{2\pi} \frac{\omega \Gamma_0(\omega)}{e^{\omega/T_H} - 1} |\mathcal{T}|^2(\omega) . \quad (21)$$

Note that, *by construction*, the evaporon only couples to the dilaton mode independent of the transverse directions \mathbf{x} . Thus, for this calculation only the greybody factor for $m = 0$ is needed.

The decay rate as written in (21) is the overlap of two functions: $|\mathcal{T}|^2$, which peaks roughly at $\omega \approx \omega_4$, and the remainder, peaking roughly at $\omega \approx T_H$. As the black hole decays the temperature will decrease and so one of the peaks will shift, whereas the other peak remains

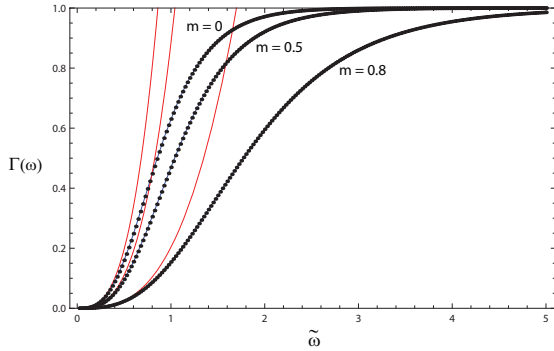


Figure 2. The greybody factor as a function of the dimensionless frequency $\tilde{\omega} = \omega R^2/r_H$ for three values of m . The dots correspond to the numerical results and interpolating lines have been added to guide the eye. The thin red lines represent to the analytic results valid in the limit of low frequencies.

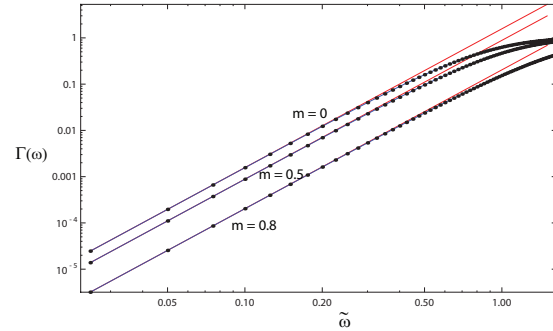


Figure 3. A close-up of the greybody factor around small $\tilde{\omega}$ is displayed in a log-log plot. The top curves correspond to $m = 0$, the middle ones to $m = 0.5$ and the lower ones to $m = 0.8$. The agreement between the numerical result and the analytic approximation valid in the regime low frequencies is manifest.

fixed at some value dependent on the choice of the coupling constant λ . Nevertheless, the overlap of the two functions is always *finite* and therefore the evaporation proceeds indefinitely.

Performing the integration (21) that would yield the *exact* rate of evaporation is not an easy task for two reasons: first, the greybody factor is only known numerically and for a limited range of frequencies and, secondly, due to our ignorance regarding the IR cutoff μ . However, we can make progress by obtaining a lower bound on the decay rate. This can be accomplished by truncating the domain of integration at a maximum frequency, ω_{max} . In particular, if $\tilde{\omega}_{max} \ll 1, \tilde{\omega}_4$ we can use the expression (19) valid at low frequencies and simultaneously the term in the transmission coefficient containing μ can be neglected³. One finds that

$$\frac{dE}{dt} > \frac{\pi r_H^5}{\omega_4^3 R^{10}} \int_0^{\tilde{\omega}_{max}} d\tilde{\omega} \frac{\tilde{\omega}^7}{e^{\pi\tilde{\omega}} - 1}. \quad (22)$$

This integral can be evaluated exactly and it is given in terms of polylogarithms. For $\tilde{\omega}_{max} \simeq 0.1$ its numerical value is 4×10^{-9} . In any event, the lower bound (22) allows us to obtain an upper limit for the evaporation time. For illustration purposes, using the relations (11) with $V_3 = \text{Vol}(S^3)R^3 = 2\pi^2 R^3$, we find that the time it takes for the horizon length to drop from $r_H = \infty$ down to $r_H = R$ is bounded by $\Delta t_{evap} \lesssim 10^{10} \omega_4^3 R^7 \kappa^{-2}$.

4. Conclusions

In conclusion, we have presented a toy model in which large black holes in AdS are allowed to evaporate in finite time. This was accomplished by effectively relaxing the boundary conditions on an AdS bulk field, coupling it to an external scalar field living in an auxiliary (asymptotically flat) space.

The main virtue of the evaporon model is that it provides a framework in which one may address the information paradox with AdS/CFT tools and the evaporation of the black hole can be followed with the CFT dual frame in a controlled way. Indeed, from the dual field theory point of view the interpretation is clear: we are adding an extra sector to the CFT and the

³ Here, we have changed variables from ω to $\tilde{\omega}$ as given by the definition (20).

evaporation of a large black hole in AdS corresponds to the transfer of energy to this infinite external reservoir.

Thus the expectation, as advocated since the early days of the AdS/CFT correspondence, is that there is no information loss in the process of formation and evaporation of black holes. Nevertheless, the issue of how the paradox is resolved is still unanswered. An interesting possibility for further exploration of the evaporon model would be to analyze also the formation of a black hole by collapsing a shell of matter in this context. It is our hope that this model may be useful to develop intuition and perhaps pinpoint the elusive ingredient necessary to settle the black hole information paradox.

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