

Quantum induced superradiance

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Superradiance, the phenomenon enabling energy extraction through radiation amplification, is not universal to all black holes. We show that semiclassical backreaction can induce superradiance, even when absent at the classical level. Specifically, we compute the quasinormal modes of a massless scalar field probing a family of rotating “quantum” black holes in three-dimensional anti-de Sitter space, accounting for all orders of backreaction due to quantum conformal matter. A subset of these modes is identified as superradiant, leading to the formation of quantum black hole “bombs.” Further, we find all thermally unstable black holes exhibit modal linear instabilities, though the converse does not hold generally. Our findings thus suggest a semiclassical version of the Hollands-Wald conjecture on black hole stability. This motivates us to propose a stability criterion for quantum black holes.

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Introduction. Like the Penrose process for extracting energy from a rotating black hole [1], superradiance occurs when a reflected wave has a larger amplitude than the incoming wave [2–9]. To wit, a wave of frequency ω and azimuthal number m incident on a Kerr black hole of mass M and spin J yields the change $dM/dJ = \omega/m$. The first law of black hole mechanics then reads [10]

$$dM = \frac{\omega \kappa_g}{8\pi G} \frac{dA}{(\omega - m\Omega)}, \quad (1)$$

for black hole surface gravity κ_g , horizon area A , and angular velocity Ω . Assuming the wave obeys the weak energy condition, such that the horizon area never decreases [11], amplification occurs when $\omega < m\Omega$, yielding black hole energy extraction, $dM < 0$ [12].

Black hole superradiance predates evaporation via thermal radiation [16] and serves as the foundation for the Blandford-Znajek mechanism [17], thought to drive jet formation in astrophysical black holes. Moreover, since repeated amplification can trigger a dynamical instability [6,18–20], superradiance is central to black hole stability and the ultimate fate of gravitational collapse.

Not all black holes exhibit superradiance, however, nor do all types of classical waves. Firstly, superradiance occurs

when scattering bosonic fields off of Kerr-Newman (KN) black holes, but not fermionic fields [21–24]. Further, large four-dimensional KN black holes with anti-de Sitter (AdS) asymptotics have no superradiance if the bosonic fields obey reflective boundary conditions at the timelike boundary [25]; superradiance can be restored if the fields obey transparent boundary conditions [26]. Another notable exception is the AdS₃ Bañados-Teitelboim-Zanelli (BTZ) black hole [27,28], which does not exhibit superradiance for neutral scalar or fermionic fields obeying Dirichlet boundary conditions at spatial infinity [29,30]. Real scalars obeying generalized boundary conditions can lead to superradiance of the BTZ [31–33], although not all superradiant modes are unstable [34].

The aforementioned examples of black hole superradiance are wholly *classical*. Both the black hole spacetime and incident wave are treated classically, adhering to general relativity. It is natural to wonder how semiclassical quantum effects modify the situation, specifically, if there are scenarios where, classically, there is no superradiance, but quantum effects instigate its onset. Particularly challenging is the influence of semiclassical backreaction due to quantum fields in an otherwise classical spacetime. A complete treatment requires solving the semiclassical Einstein equations, a difficult open problem.

Here we show quantum backreaction can induce black hole superradiance. Our study centers on an *exact* black hole in semiclassical gravity, the “quantum” BTZ (qBTZ) black hole [36,37]. The qBTZ solution arises from brane-world holography [38], a framework where a $(D-1)$ -dimensional end-of-the-world brane is coupled to general relativity in a D -dimensional asymptotically AdS space,

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which has a dual description as a conformal field theory (CFT) living on the AdS boundary. A higher curvature gravity theory is induced on the brane coupled to a CFT with large central charge and an ultraviolet cutoff. In this formalism, AdS black holes that localize on branes may be understood as quantum-corrected black holes from the brane perspective and can be exactly constructed to all orders of backreaction in three dimensions [36,39].

Quantum BTZ black holes. We focus on the quantum BTZ (qBTZ) black hole [37]. Its construction relies on an AdS₃ brane [40,41] intersecting the AdS₄ C-metric [42] black hole horizon [43,44]. Aided by holography, the geometry and thermodynamics of the qBTZ are known analytically and nonperturbatively in backreaction. Here, we summarize the essentials of the solution. Further details can be found in the Supplemental Material [45].

In Boyer-Lindquist coordinates $(\bar{t}, \bar{r}, \bar{\phi})$, the neutral qBTZ black hole has line element [37]

$$ds^2 = -\left(\frac{\bar{r}^2}{\ell_3^2} - 8\mathcal{G}_3 M - \frac{\ell \mu \eta^2}{r}\right) d\bar{t}^2 + \left(\frac{\bar{r}^2}{\ell_3^2} - 8\mathcal{G}_3 M + \frac{(4\mathcal{G}_3 J)^2}{\bar{r}^2} - \ell \mu (1 - \tilde{a}^2)^2 \eta^4 \frac{r}{\bar{r}^2}\right)^{-1} d\bar{r}^2 + \left(\bar{r}^2 + \frac{\mu \ell \tilde{a}^2 \ell_3^2 \eta^2}{r}\right) d\bar{\phi}^2 - 8\mathcal{G}_3 J \left(1 + \frac{\ell}{x_1 r}\right) d\bar{\phi} d\bar{t}. \quad (2)$$

Here, r is a function of coordinate $\bar{r}$, which is kept for convenience, and μ is a mass parameter for the AdS₄ black hole. Quantities M and J denote the mass and angular momentum of the black hole

$$8\mathcal{G}_3 M = 4 \frac{-\kappa x_1^2 + \tilde{a}^2(4 - \kappa x_1^2)}{(3 - \kappa x_1^2 - \tilde{a}^2)^2}, \quad 4\mathcal{G}_3 J = \frac{4\ell_3 \tilde{a}(1 - \kappa x_1^2 + \tilde{a}^2)}{(3 - \kappa x_1^2 - \tilde{a}^2)^2}, \quad (3)$$

with $\eta \equiv 2x_1/(3 - \kappa x_1^2 - \tilde{a}^2)$. Here, ℓ is the UV cutoff length scale, $\mathcal{G}_3 = G_3/\sqrt{1 + (\ell/\ell_3)^2}$ is the “renormalized” Newton’s constant, and ℓ_3 is the AdS₃ length. Further, $\tilde{a} \equiv ax_1^2/\ell_3$, for rotation parameter a , $\kappa = \pm 1$, 0 corresponds to different brane slicings, and x_1 is a real parameter whose precise meaning is not relevant for us here.

The inner/outer event horizons $\bar{r}_{\pm}$ are generated by orbits of the Killing vector $\chi = \partial_{\bar{t}} + \Omega_{\pm} \partial_{\bar{\phi}}$, where $\Omega_{\pm}$ denotes the horizon angular velocity relative to a *nonrotating* frame at spatial infinity

$$\Omega_{\pm} \equiv \frac{a}{r_{\pm}^2 + a^2 x_1^2} \left(1 + \frac{r_{\pm}^2 x_1^2}{\ell_3^2}\right) = \frac{d\bar{\phi}}{d\bar{t}}. \quad (4)$$

The ergoregion corresponds to orbits of the Killing vector $\partial_{\bar{t}}$, where the $\bar{t}\bar{t}$ -metric component of (2) vanishes.

The metric (2) is understood as a *quantum* black hole as it is a solution to the induced semiclassical theory on the brane at all orders in quantum backreaction. The dimensionless parameter $\nu \equiv \ell/\ell_3$ controls the strength of backreaction, while the geometry (2) represents the quantum-corrected black hole due to the backreacting CFT₃. This is most easily seen in the static limit ($J = 0$), where the metric takes the same form as the quantum-corrected BTZ solution due to perturbative backreaction of a conformally coupled scalar [48,49]. Notably, however, the metric (2) is valid for any strength of backreaction, i.e., any $\nu \geq 0$. As $\nu \rightarrow 0$ with fixed c_3 , backreaction vanishes, gravity weakens, and the metric resembles the classical BTZ black hole. Conversely, large ν corresponds to strong backreaction; in this limit, higher-derivative terms in the induced action become important.

Combined, (a, κ, x_1) parametrize a family of quantum black holes at fixed backreaction [37]. It is worth distinguishing between solutions with $\kappa = +1$ and those with $\kappa = -1$. The former can have $M < 0$ and are considered quantum-dressed conical defects, while the latter have $M \geq 0$ and represent quantum-corrected BTZ black holes. Depending on the range of parameters, qBTZ can have pathological features, e.g., naked closed timelike curves. Dismissing such traits requires parameter restrictions, particularly $1 - \tilde{a}^2 \geq 0$ and $\eta > 0$ [37], though the full parameter space has not yet been fully explored.

The thermodynamics of the quantum BTZ black hole is inherited from the AdS₄ bulk black hole, with M , J , and temperature T [37,44,50]. Explicit expressions are known analytically and presented in the Supplemental Material [45]. Notably, the Bekenstein-Hawking area entropy [16,51,52] of the classical AdS₄ black hole is identified to be the semiclassical generalized entropy [53], S_{gen} , accounting for both (higher-curvature) gravitational and von Neuman entropy of matter, $S_{\text{gen}} = S_{\text{grav}} + S_{\text{vN}}^{\text{mat}}$. Thus, for fixed ℓ and ℓ_3 , the thermodynamic first law is

$$dM = T dS_{\text{gen}} + \Omega dJ, \quad (5)$$

and is valid for any strength of backreaction, i.e., any ν .

Quantum black hole superradiance. As with classical black holes, superradiance of a rotating quantum black hole can plausibly occur when $\omega < m\Omega$. For quantum black holes, this will only happen, however, when the generalized second law [53] holds, $dS_{\text{gen}} \geq 0$. Indeed, a test wave of frequency ω impinging a neutral rotating quantum black hole will also obey $dM/dJ = \omega/m$. The quantum black hole, e.g., qBTZ, will then change according to the semiclassical first law (5), such that

$$dM = \frac{\omega T}{(\omega - m\Omega)} dS_{\text{gen}}. \quad (6)$$

Thus, if $\omega < m\Omega$, energy may be extracted from the black hole, suggesting it is unstable.

The above argument is a semiclassical generalization of the one proposed by Bekenstein [10]. Notably, the scattering wave need not be treated quantum mechanically. Replacement of the classical horizon area for generalized entropy is a result of the quantum matter backreacting on the geometry, independent of the test wave.

Admittedly, the reasoning leading to (6) is only heuristic. Indeed, when backreaction is switched off, the qBTZ becomes the classical BTZ black hole, which recall does not always exhibit superradiance. A reason to expect the qBTZ black hole features superradiance opposed to its classical sibling, is that, like higher-dimensional Kerr-AdS black holes [25], there are regions which rotate faster than the speed of light. Thence, energy escaping across the horizon will be negative, allowing for superradiance to occur. We present this in the Supplemental Material [45].

Perturbing quantum BTZ: To directly test superradiance, we perturb the rotating qBTZ with a massless probe scalar field Φ bound to the brane and analyze the associated quasinormal modes (QNMs), resonant modes of linear perturbations satisfying ingoing (outgoing) boundary conditions at the horizon (infinity) [54]. Our approach follows [55,56]. Let the scalar field have action

$$I_\Phi = \int dr d^2x \sqrt{-g} \partial_a \Phi \partial^a \Phi. \quad (7)$$

Decompose the field into its Fourier modes

$$\Phi(r, \bar{t}, \bar{\phi}) = \sum_{m=-\infty}^{\infty} \int d\omega e^{-i(\omega \bar{t} - m \bar{\phi})} \bar{\Phi}_m(\omega, r), \quad (8)$$

for frequency ω and azimuthal number m . [Here, for convenience, we work in coordinates $(\bar{t}, r, \bar{\phi})$, such that the coordinate system is not identical to those given in metric (2).] Solutions to the scalar equation of motion, $\square\Phi = 0$, near the asymptotic AdS₃ boundary behave as

$$\Phi \sim \Phi_{(0)} r^{-\Delta_-} + \Phi_{(+)} r^{-\Delta_+}, \quad \Delta_{\pm} = 1 \pm 1, \quad (9)$$

with scalar operator dimension $\Delta_{\pm}$, and coefficients $\Phi_{(0)}$ and $\Phi_{(+)}$ are functions of $(\bar{t}, \bar{\phi})$. Physically, $\Phi_{(0)}$ represents the source and $\Phi_{(+)}$ the vacuum expectation value of the scalar. Meanwhile, near the horizon,

$$\Phi \sim \Xi_1 e^{-ir_*(\omega - m\Omega)} + \Xi_2 e^{+ir_*(\omega - m\Omega)}, \quad (10)$$

for tortoise coordinate r_* (near the horizon)

$$r_* = \frac{\log(r - r_+)}{2\kappa_g}, \quad \kappa_g = \sqrt{-\frac{1}{2} \nabla_a \chi_b \nabla^a \chi^b}, \quad (11)$$

with surface gravity κ_g . To obtain QNMs, we use boundary conditions that disallow outgoing flux across the horizon and force the field to vanish at infinity, imposing $\Xi_2 = 0$ and $\Phi_{(0)} = 0$, respectively.

In general, QNM frequencies are complex, $\omega = \text{Re}(\omega) + i \text{Im}(\omega)$, signaling the decay or growth of the field. In the Supplemental Material [45], we display an example of the QNM spectrum of a scalar fluctuation with vanishing momentum ($m = 0$). At $\nu = 0$, the QNMs agree with the classical BTZ spectrum [57–59]. At vanishing momentum m , these QNMs are purely imaginary and move along the imaginary axis as the black hole angular momentum is increased. Quantum backreaction strongly affects the QNM spectrum and, in particular, causes the modes of the BTZ black hole to acquire a finite real part, hence causing a transition from purely dissipative to propagating modes. This same behavior was seen for static qBTZ [55].

Superradiant modes: Superradiance manifests itself as unstable perturbations that grow in time. According to the decomposition (8), if the imaginary part of ω is positive, $\text{Im}(\omega) > 0$, the field will grow exponentially, indicating an instability; a stable mode has $\text{Im}(\omega) < 0$. Hence, the onset of the instability occurs when $\text{Im}(\omega) = 0$, i.e., $\omega \in \mathbb{R}$. This can be made more precise by the following argument. Since the geometry is stationary, the field equations are invariant under $\bar{t} \rightarrow -\bar{t}$ and $\omega \rightarrow -\omega$; hence, the complex conjugate $\bar{\Phi}^*$ is also a solution to the wave equation. Indeed, if one takes the tortoise coordinate

$$\frac{dr_*}{dr} = \frac{H'(r_+)r_+}{2r\kappa_g H(r)}, \quad (12)$$

then one can show that the differential equation satisfied by Φ can be put in the form

$$\bar{\Phi}''(r_*) + V_{\text{eff}}(r(r_*)) \bar{\Phi}(r_*) = 0, \quad (13)$$

where

$$V_{\text{eff}}(r) = V(r) [(\tilde{a}^2 \ell_3^3 \omega - \tilde{a} n (\ell_3^2 + r^2 x_1^2) + \ell_3 r^2 x_1^2 \omega)^2 - \ell_3^2 r^2 x_1^4 H(r) ((\tilde{a}^2 - 1)^2 \Delta^2 m^2 r^2 + (n - \tilde{a} \ell_3 \omega)^2)], \quad (14)$$

$$V(r) = \frac{4\kappa_g^2}{(\tilde{a}^2 - 1)^2 \Delta^2 H'(r_+)^2 \ell_3^2 r_+^2 x_1^4}. \quad (15)$$

By Abel's identity, the Wronskian is independent of r_* , $W(\bar{\Phi}, \Phi^*)(r_*) = W(\bar{\Phi}, \Phi^*)$, and hence conserved along the radial direction. The Wronskians evaluated at the

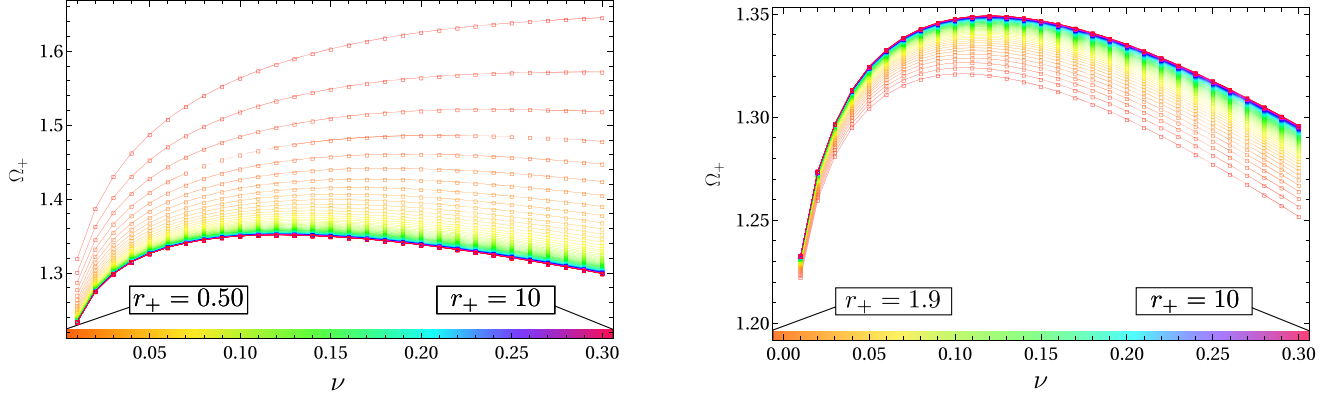


FIG. 1. The onset of superradiance for the rotating qBTZ black hole. Left: dressed cones ($\kappa = +1$). Right: corrected black holes ($\kappa = -1$). The color of the points represents increasing horizon radius.

horizon and asymptotic boundary thus agree, $W(\bar{\Phi}, \bar{\Phi}^*)|_{r_+}^\infty = 0$. Using the expansions in (9) and (10),

$$\bar{\Phi}_{(0)}^* \bar{\Phi}_{(+)} - \bar{\Phi}_{(0)} \bar{\Phi}_{(+)}^* = i(\omega - m\Omega) |\Xi_1|^2, \quad (16)$$

where we absorbed a constant prefactor into the definition of the source and vacuum expectation value. The Ward-Takahashi identities then imply (cf. [60]),

$$\begin{aligned} \partial_t \mathcal{M} &= -2\pi(\Phi_{(0)}^* \partial_t \Phi_{(+)} - \Phi_{(0)} \partial_t \Phi_{(+)}^*), \\ &= 4\pi\omega(\omega - m\Omega) |\Xi_1|^2, \end{aligned} \quad (17)$$

where $\mathcal{M} = \int_\Sigma T_{ij} k^i n^j$ is the energy of the spacetime on a constant time, one-dimensional spacelike hypersurface Σ , with n^i the unit normal to Σ and k^i the timelike Killing vector. This implies the energy decreases, $\partial_t \mathcal{M} < 0$, when the condition $\omega < m\Omega$ is fulfilled.

To precisely obtain superradiant modes, we resort to numerics. At the onset of the instability, the QNM frequency must take the form $\omega = m\Omega_+$. We solve the equation of motion for the scalar field subject to the boundary condition of no outgoing modes at the horizon with $\omega = m\Omega_+$. The solution near the asymptotic boundary takes the form Eq. (9). We use a shooting method from the outer horizon and search for solutions with $\Phi_{(0)} = 0$. [61] We display the search in Fig. 1, where we plot Ω_+ as a function of ν for modes at the cusp of instability for rotating quantum dressed cones and corrected black holes.

Notably, at least for the lightest modes ($m = 1$), we find no parameter combinations for which superradiant modes exist for fully nonperturbative quantum corrected BTZ black holes ($\nu > 1$), and thus, such black holes appear to not exhibit superradiant instabilities for the massless scalar. Such instabilities only exist for weak backreaction, $\nu < 1$; see Fig. 2. Alternatively, light superradiant modes exist for all ν for the quantum-dressed conical defects (Supplemental Material [45]). Thus, while quantum-dressed cones are always dynamically unstable for the lightest modes,

quantum-corrected black holes are stable against the lightest modes for large backreaction.

Quantum black hole bombs: Repeated superradiant amplification leads to superradiant instabilities as superradiant modes become locally trapped near the black hole [5]. For AdS black holes, this trapping naturally arises due to the timelike boundary, forming a black hole “bomb” for small black holes ($r_+/\ell_4 \ll 1$) [62–66], while large ones ($r_+/\ell_4 > 1$) remain stable. Likewise, the onset of superradiant instability for the qBTZ, bolstered by the existence of superradiant QNMs (Fig. 2), implies the formation of quantum black hole bombs. Notably, quantum-dressed cones and quantum-corrected black holes need not be extremely small (e.g., $r_+/\ell_3 \sim 0.75$ in Fig. 5, see supplemental material [45]).

Superradiant instabilities are inherently dynamical; they do not necessarily coincide with thermal instability.

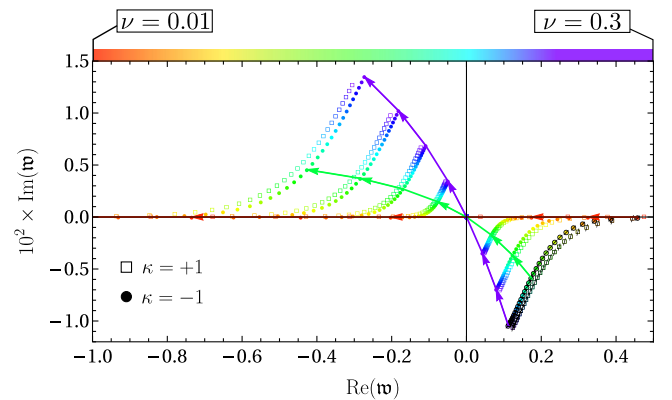


FIG. 2. Superradiant QNMs ($\mathfrak{w} = \omega/(2\pi T)$) in a corotating frame ($\mathfrak{w}_{\text{onset}} = 0$). Displayed are $m = 1$ modes at $x_1 = 0.21$. Solid circles denotes $\kappa = -1$, while the empty boxes correspond to $\kappa = +1$, and color represents increasing strength of backreaction. Lines with arrows indicating increasing values of a are drawn for $\nu = \{0.01, 0.15, 0.3\}$ to guide the eye. Note that all data for which $\text{Re}(\mathfrak{w}) > 0$ has $\text{Im}(\mathfrak{w}) < 0$.

For Kerr-AdS black holes, superradiant instability is consistent with thermal instability, as small black holes have negative heat capacity [67]. It is worth then to analyze the thermal stability of the rotating qBTZ black hole. Previous investigations include evaluating only the behavior of the heat capacity [68–72]. Here, we go beyond analyzing only the heat capacity, and instead study the sign of the eigenvalues of the Hessian matrix

$$\text{Hess}_S = \begin{pmatrix} \frac{\partial^2 S_{\text{gen}}}{\partial M^2} & \frac{\partial^2 S_{\text{gen}}}{\partial J \partial M} \\ \frac{\partial^2 S_{\text{gen}}}{\partial M \partial J} & \frac{\partial^2 S_{\text{gen}}}{\partial J^2} \end{pmatrix}, \quad (18)$$

for entropy S_{gen} (see [37] for explicit expressions). Analogous to ordinary thermal systems, we expect thermal instability if the Hessian has a positive eigenvalue. We find, at the onset of instability, the two eigenvalues of Hess_S for quantum-corrected black holes ($\kappa = -1$) are both negative, indicating a thermally stable system that exhibits dynamical instabilities. Meanwhile, for the quantum-dressed cones ($\kappa = 1$), Hess_S has only negative eigenvalues for $M > 0$, but for $M < 0$, it can have at least one positive eigenvalue over a specific range of parameters. Thus, for the dressed cones, there exist solutions that are both thermodynamically and dynamically unstable.

Discussion. We established that quantum backreaction can induce superradiance in black holes, which would otherwise not exhibit the phenomenon. Here, we focused on the rotating quantum BTZ black hole, finding superradiant QMNs giving rise to modal instabilities. Superradiance of course depends on the choice of positive frequency, a subtle issue for rotating black holes as they typically do not possess a globally timelike Killing vector. Here, we take $\omega > 0$. It is conceivable that other choices would imply no superradiance, as in Kerr-Newman-AdS₄ [26]. A more detailed mode analysis would shed more light on this subtlety. Further, a nonholographic perturbative analysis of backreaction due to a conformally coupled scalar in rotating BTZ [73] suggests no superradiance. Their corrected metric, however, is more complicated than the exact qBTZ [37], and evidently displays different physics. In fact, a reason why the qBTZ exhibits superradiance, in contrast with the classical BTZ or the perturbatively corrected solution, is due to the existence of a new macroscopic scale, $\ell \sim c_3 L_P$ [62]. Our work has many implications and avenues worth exploring.

Beyond rotating qBTZ: Our investigation here is nonexhaustive, being centered on classical massless probe scalars in rotating qBTZ. It would be straightforward to apply our methods to other types of quantum black holes. This includes the charged qBTZ [74,75], and their charged/rotating de Sitter and Minkowski counterparts [39,76–78]. Notably, in 3D vacuum general relativity, there are no de

Sitter or flat black holes to begin with. Should our results hold beyond AdS₃, not only would backreaction induce black hole horizons, but all 3D quantum black holes would exhibit superradiance.

Not all higher-dimensional Kerr-Newman black holes exhibit superradiance [25,26]. It would be interesting to see if backreaction in these backgrounds would likewise facilitate superradiance. This is technically challenging because exact quantum black hole solutions (nonperturbative in backreaction) in four or higher dimensions are not known. A first step would be to study backgrounds with perturbative backreaction; indeed, our findings show that superradiance is induced for quantum-corrected black holes at weak backreaction.

Finally, it would be worth considering other types of probes, e.g., massive scalars, or conformally coupled fields. It would also be interesting to extend to quantum matter probes, holographically dual to 3 + 1 dimensional bulk fields [79]. The quantum analog of classical superradiance is well known for Kerr and static, charged black holes [9,13,80–82], where the black holes spontaneously emit particles with classically superradiant modes. In principle, similar nonthermal radiation is expected for quantum test fields in a quantum black hole (see, e.g., [83]).

Testing cosmic censorship: General relativity predicts black holes formed under collapse develop singularities [84]. To maintain predictability, Penrose conjectured “cosmic censorship,” which in its weak form (WCC), states black hole singularities lie behind event horizons [1,85]. Tests of WCC, where one attempts to overspin or overcharge Kerr-Newman black holes via particle/wave scattering [86,87], have only shown WCC is robust. In fact, scalar matter that overspins Kerr, thus violating WCC, is repelled due to superradiance [88]. Similarly, the rotating qBTZ cannot be overspun [89]. It would be interesting to check if induced superradiance is the cause.

Stability of quantum black holes: Here, we uncovered unstable superradiant modes in the qBTZ geometry. This raises the question of the fate of such instabilities, suggesting the existence of “black resonators” [90] as exact solutions in semiclassical gravity. Further, at the onset of instability, quantum-corrected black holes were found to be thermally stable, while, for specific parameter ranges, quantum-dressed cones were thermally unstable. This is consistent with the Hollands-Wald stability criterion [91]—stationary black holes that locally extremize the entropy for fixed conserved charges are dynamically stable under linear perturbations that preserve said conserved charges—and suggests it extends semiclassically. For classical stationary and axisymmetric black holes, dynamical stability is equivalent to the following: non-negativity of the canonical energy $\mathcal{E}$, a bilinear form over linearized perturbations that fix appropriate conserved quantities. Conversely, for AdS

black holes, $\mathcal{E} < 0$ indicates a superradiant instability [92]. It would be interesting to develop such a criterion for quantum black holes. Naively, the criterion is to be replaced by

$$\mathcal{E}_{\text{sc}} = \delta^2 M - \Omega \delta^2 J - T \delta^2 S_{\text{gen}}, \quad \mathcal{E}_{\text{sc}} \geq 0, \quad (19)$$

where, as in the thermodynamic first law, the classical event horizon area has been supplanted with generalized entropy, and δ^2 indicates second-order variations. In fact, since the AdS C-metric belongs to the class of metrics the classical stability criterion [92] applies to, the criterion (19) automatically follows for the quantum BTZ black hole due to the identification of the bulk area-entropy and brane generalized entropy (as in the first law).

Weak cosmic censorship and black hole stability, moreover, are intimately linked. Indeed, $\mathcal{E} > 0$ implies the satisfaction of a local Penrose inequality (PI) [91], an inequality whose violation yields a counterexample to WCC [93]. The local PI is thus a necessary and sufficient condition for dynamical stability. Quantum matter, however, is known to violate the classical PI [94,95]. This motivates the need for (weak) quantum cosmic censorship [96,97] for which a quantum PI would be an input. To date, all quantum BTZ black holes obey such a proposed inequality [72] for all orders of backreaction. It would be worth seeing how these robust quantum inequalities relate to the proposed criterion (19).

Beyond black holes: Finally, superradiance is present in nongravitational systems, historically uncovered in studies in optics [99]. Likewise, we speculate the quantum induced

superradiance uncovered here is an instance of a more general phenomenon. It would be interesting to find other contexts where superradiance arises via semiclassical backreaction, e.g., Q -ball dynamics [100].

Note added. In remembrance of our extraordinary colleague and friend, Umut Gursoy, who passed away while this paper was under review.

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Data availability. The data that supports the conclusions contained in the manuscript is publicly available as an ancillary file [101].

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