

Trapping of waves and null geodesics for rotating black holes

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We present dynamical properties of linear waves and null geodesics valid for Kerr and Kerr–de Sitter black holes and their stationary perturbations. The two are intimately linked by the geometric optics approximation. For the null geodesic flow the key property is the *r-normal hyperbolicity* of the trapped set and for linear waves it is the distribution of quasinormal modes: the exact quantization conditions do not hold for perturbations but the bounds on decay rates and the statistics of frequencies are still valid.

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I. INTRODUCTION

The Kerr solutions [1] to Einstein equations are considered as physically relevant models of rotating black holes. The Kerr metrics depend on two parameters: mass M and rotational parameter a ; the special case $a = 0$ is the Schwarzschild metric. The *Kerr–de Sitter* solutions describe rotating black holes in the case of positive cosmological constant, $\Lambda > 0$ —see (1) below for the formula for the metric and Fig. 1 for the plot of admissible values of the parameters. Because of the observed *cosmic acceleration* [2], the current cosmological Λ CDM model assumes $\Lambda > 0$. As explained below, $\Lambda > 0$ makes the study of the topic of this article, *quasinormal modes* for black holes (QNMs), mathematically more tractable while not affecting the description of the physical phenomenon of ringdown [3].

The classical dynamics of Kerr black holes is concerned with the behavior of null geodesics of the corresponding metric, that is, the trajectories of photons in the gravitation field of the black hole. The key dynamical object is the *trapped set*, consisting of all null geodesics in phase space (position and momentum space) which never cross the event horizon of the black hole or escape to infinity. In other words, this is the set where the strength of gravitational fields forces photons to travel on bounded orbits.

In the case of a Schwarzschild black hole ($a = 0$) the time slice of the trapped set is just the phase space of a sphere (mathematically, the cotangent bundle of a sphere) called the *photon sphere*: along the photon sphere, all photons travel on closed orbits. A traveler who crosses the photon sphere, although still visible to outsiders, is forced to cross the black hole horizon eventually. In the case of nonzero angular momentum ($a \neq 0$) the trapped set is no longer the phase space or cotangent bundle of a smooth spatial set; instead it becomes a nontrivial object in the phase space. The photons are *trapped* because of the strength of the gravitational field but most of them (that is, a set of full measure) no longer travel along closed orbits—see Fig. 2 for a visualization of the trapped set and (2) for the analytic description. Although the trapped set is no longer

the phase space of a spatial object, it remains a smooth five-dimensional manifold. The symplectic form on the phase space of a time slice [see (3)] restricts to a nondegenerate form on the trapped set. That means that the time slice of the trapped set is a smooth symplectic manifold.

A remarkable feature of the geodesic flow on Kerr(–de Sitter) metrics is its complete integrability [4] in the sense of Liouville–Arnold [5]: there exist action variables which define invariant tori on which the motion is linear.

In this article we describe another important feature of the dynamics: *r-normal hyperbolicity*. It means that the flow is hyperbolic in directions normal to the trapped set in ways r -fold stronger than the flow on the trapped set—see (4) for a mathematical definition. This property, unlike complete integrability, is known to be stable under perturbations [6]: a small C^r (r times differentiable) stationary perturbation of the metric will destroy complete integrability but will preserve C^r structure of the trapped set and *r-normal hyperbolicity*. For Kerr black holes the condition holds for *each* r and hence regular perturbations will maintain the regular structure of the trapped set of light trajectories [7,8].

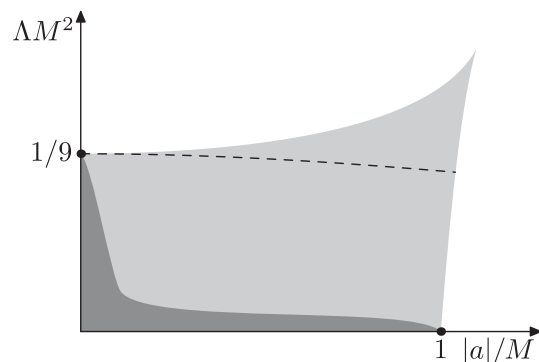


FIG. 1. Numerically computed admissible range of parameters for the subextremal Kerr–de Sitter black hole (light shaded) [30] and the range to which our results apply (dark shaded). QNMs are defined and discrete for parameters below the dashed line, $(1 - \alpha)^3 = 9\Lambda M^2$; see [8], Sec. 3.2.

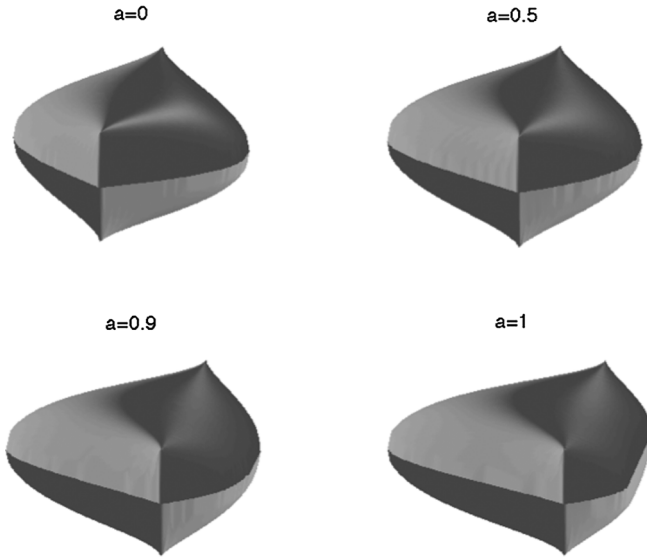


FIG. 2. Visualization of the trapped set for different values of a , with $\Lambda = 0$. The figures show the four-dimensional set $K \cap \{\xi_t = 1\}$ (K is the five-dimensional trapped set) projected to the coordinates $(x, y, z) = (\xi_\varphi, \theta, \xi_\theta)$. For $a = 0$ this corresponds to the visualization of the phase space of the 2-sphere: the sphere is parametrized by the coordinates $0 \leq \theta \leq \pi$, $0 \leq \varphi < 2\pi$, $(\sin \theta \cos \varphi, \sin \theta \sin \varphi, \cos \theta)$. The conjugate coordinates are denoted ξ_θ and ξ_φ and the restriction to $\xi_t = 1$ means that $\xi_\theta^2 + \sin^{-2}\theta \xi_\varphi^2 = 27M^2$. The vertical singular interval in front corresponds to $\theta = 0$, with the symmetrical interval in the back corresponding to $\theta = \pi$: the coordinates (θ, φ) on the sphere are singular at that point. The structure of K becomes more interesting when $a > 0$ as shown in the three examples. The additional coordinate, not shown in the figures, r is a function of ξ_φ and ξ_t only. When $a = 0$, we have $r = 3M$, but r gets larger to the left ($\xi_\varphi > 0$) and smaller to the right ($\xi_\varphi < 0$) when $a > 0$. When $a = 1$ we see the flattening in the (θ, ξ_θ) plane at extremal values of ξ_φ : the trapped set touches the event horizon $r = 1$ which results in lack of decay, and some QNMs have null imaginary parts [22,23]. Dynamically and invariantly this corresponds to the vanishing expansion rates—see Fig. 4.

The classical dynamical features are crucial for the behavior of gravitational waves emitted by black holes, especially during the *ringdown* phase, when a black hole spacetime settles down after a large cosmic event such as a binary black hole merger. Gravitational waves are expected to be observable by the existing detectors, once they are running at full capacity, and to provide information about the parameters of astrophysical black holes. During the ringdown phase, the behavior of gravitational waves is driven by the linearized system [3] and is much simpler to simulate numerically than the merger phase [9–11]. At ringdown, gravitational waves have a fixed set of complex frequencies, known as *quasinormal modes* (QNMs) [3] and depending only on the parameters of the black hole, rather than the specifics of the event. The simplest model of ringdown is obtained by solving the linear scalar wave equation for the Kerr(–de Sitter)

black hole spacetimes, and in that case quasinormal modes can be rigorously defined. More complicated linearizations have also been studied [3,12] but we concentrate on the simplest setting here. On the relevant time and space length scales, the value of the cosmological constant Λ does not have a physical effect on the ringdown since gravitational waves are generated in a neighborhood of the black hole but $\Lambda > 0$ makes the mathematical definition of QNMs much easier by eliminating the polynomial falloff for waves [3], Sec. 5.1.

In a more general physical or geometric context of scattering theory, quasinormal modes, also known as resonances, replace bound states (eigenvalues), for systems which allow escape of energy [13] and simultaneously describe oscillations (real parts of the mode) and decay (imaginary parts). They appear in expansions of waves—see (6) below, just as waves in bounded regions are expanded using eigenvalues. This dynamical interpretation immediately suggests that the distribution of QNMs is related to the trapping on the classical level—see [14,15] for a discussion and recent experimental results in the setting of microwave billiards.

The relation between dynamics and distribution of QNMs and resonances has been particularly well studied in problems where a reduction to one dimension is possible. More generally, complete integrability allows quantization rules which can be used to describe resonances in the semiclassical or high energy limit. In the setting of black holes this goes back to [16].

For Schwarzschild black holes the Regge–Wheeler reduction [see (7) below] produces a one-dimensional potential similar to the Eckart barrier potential $\cosh^{-2}x$ for which resonances are given by $\pm\sqrt{3}/2 - i(n + 1/2)$ —see [17] for a review in the context of chemistry, [18] for a mathematical discussion in the Schwarzschild case, and [19] for a general study of Pöschl–Teller potentials. Putting together different angular momenta produces an (approximate) lattice of resonances.

When $a \neq 0$, that is, in the genuine Kerr case, the degenerate QNMs split in a way similar to the Zeeman effect. They have been recently studied using WKB methods based on the completely integrable structure [20–23] and the Zeeman-like splitting has been rigorously confirmed.

The point of this article is to describe recent mathematical results [7,8,18,20,21,24,25] which apply to *stationary perturbations of Kerr metrics* and do not depend on the completely integrable structure. They are based on the use of the r -normal hyperbolicity of the trapped set and show that many features of QNMs studied using WKB methods available in the completely integrable case persist for perturbations. The r -normal hyperbolicity of black hole dynamics [7] has not been discussed in the physics literature but the importance of normal hyperbolicity in molecular dynamics has been explored [26]. It would be interesting to consider the stability of r -normally

hyperbolic dynamics under more general, nonstationary, perturbations.

In particular, we show how dynamical features (such as the maximal and minimal expansion rates) and statistical properties of the distribution of QNMs depend on $a = J/Mc$, the rescaled angular momentum. For the exact Kerr (or Kerr–de Sitter) black hole the counting law for QNMs and their maximal and minimal decay rates determine the mass and the angular momentum. For the perturbed case, they determine stable features such as the symplectic volume of the trapped set and classical decay rates.

The presentation is organized as follows: in Sec. II we discuss classical dynamics and define r -normal hyperbolicity in a precise way; in Sec. III we describe the challenges of rigorously defining of QNMs for Kerr and Kerr–de Sitter black holes. The difficulties come from the presence of the ergosphere which obstruct standard methods for defining resonances (lack of coercivity or ellipticity of the stationary wave equation) and from the “size” of infinity in the Kerr ($\Lambda = 0$) case. In Sec. IV we present the quantitative results about the distribution of QNMs valid for perturbations of Kerr black hole: bounds on imaginary parts of the modes (11), the counting law (12), and the consequences for solutions of the wave equation (13). The strongest results are subject to a pinching condition (10) which is valid for all but rapidly rotating rotating black holes. The results presented here and some of the figures have appeared in works aimed at the mathematical audience [8] and this is an attempt to relate them to an active field of research in physics.

II. THE TRAPPED SET OF NULL GEODESICS

We are interested in null geodesics and the wave equation $\square_g u = 0$, where g is the Kerr(–de Sitter) metric [1,4,7,8]. The Kerr–de Sitter metric is a generalization of Kerr to the case of a positive cosmological constant Λ . For a black hole of mass $M > 0$ (for $\Lambda > 0$ there is no unique definition of global mass; here we treat M simply as a parameter of the metric), rotating with speed a the space slice is

$$X = (r_+, r_C) \times \mathbb{S}^2,$$

where $r_C < \infty$ when $\Lambda > 0$, and $r_C = \infty$ when $\Lambda = 0$. The behavior of \square_g for r near r_C is dramatically different in the two cases: for $\Lambda > 0$ the metric is asymptotically hyperbolic in the sense of non-Euclidean geometry (infinity is *large* in the sense that the volume of balls grows exponentially in radius) and for $\Lambda = 0$ it is asymptotically Euclidean (infinity is *small* in the sense that the volume of balls grows polynomially). For solutions to the wave equation that produces differences in long time decay and in the behavior at low energies [27–29]. The surface $r = r_+$ is the event horizon of the black hole. When $\Lambda > 0$, $r = r_C$ is the cosmological horizon. While the two horizons have different physical interpretations, their mathematical roles in the study of wave decay and QNMs are remarkably similar.

The geodesic flow can be considered as a flow on the phase space (the position-momentum space) of $\mathbb{R} \times X$, that in mathematical terms on the cotangent bundle $T^*(\mathbb{R} \times X)$. We denote the coordinates on $\mathbb{R} \times X$ by (t, r, θ, φ) (see Fig. 2) and write $(\xi_t, \xi_r, \xi_\theta, \xi_\varphi)$ for the corresponding conjugate (momentum) variables. The flow is given by the classical Hamiltonian flow [5] for the Hamiltonian G ,

$$\begin{aligned} \dot{t} &= \partial_{\xi_t} G, & \dot{r} &= \partial_{\xi_r} G, & \dot{\theta} &= \partial_{\xi_\theta} G, \\ \dot{\varphi} &= \partial_{\xi_\varphi} G, & \dot{\xi}_t &= -\partial_t G, & \dot{\xi}_r &= -\partial_r G, \\ \dot{\xi}_\theta &= -\partial_\theta G, & \dot{\xi}_\varphi &= -\partial_\varphi G, \end{aligned}$$

where

$$\begin{aligned} G &= \rho^{-2}(G_r + G_\theta), & \rho^2 &= r^2 + a^2 \cos^2 \theta, \\ G_r &= \Delta_r \xi_r^2 - \frac{(1 + \alpha)^2}{\Delta_r} ((r^2 + a^2) \xi_t + a \xi_\varphi)^2, \\ G_\theta &= \Delta_\theta \xi_\theta^2 + \frac{(1 + \alpha)^2}{\Delta_\theta \sin^2 \theta} (a \sin^2 \theta \xi_t + \xi_\varphi)^2, \\ \alpha &= \frac{\Lambda a^2}{3}, & \Delta_r &= (r^2 + a^2) \left(1 - \frac{\Lambda r^2}{3}\right) - 2Mr, \\ \Delta_\theta &= 1 + \alpha \cos^2 \theta. \end{aligned} \quad (1)$$

The function G is the dual metric to the semi-Riemannian Kerr(–de Sitter) metric g (see for example [8], Sec. 3.1, for the formulas for g). It is also the principal symbol of \square_g in the sense that

$$\square_g = G(t, r, \theta, \varphi, \partial_t/i, \partial_r/i, \partial_\theta/i, \partial_\varphi/i), \quad i = \sqrt{-1},$$

modulo a first order differential operator. The limiting radii r_+, r_C solve $\Delta_r = 0$.

The trapped set consists of null geodesics that stay away from $r = r_+, r = r_C$ for all times. The variables (θ, ξ_θ) evolve according to the Hamiltonian flow of G_θ , and ξ_t, ξ_φ are conserved. The trapping depends on the evolution of (r, ξ_r) according to the flow of G_r , which is essentially the one-dimensional motion for a barrier top potential [18]. Under the assumptions that either $a = 0, 9\Lambda M^2 < 1$ or $\Lambda = 0, |a| < M$, and for nearby values of M, Λ, a [8], Prop. 3.2, [30] the trapped set K is given by

$$K = \{G = 0, \xi_r = 0, \partial_r G_r = 0, \xi \neq 0\}. \quad (2)$$

For $a = 0, \partial_r G_r = 0$ gives $r = 3M$, the radius of the photon sphere. For $a \neq 0$, a more careful analysis is required, but K is still a smooth submanifold of the characteristic set $\{G = 0\}$. Moreover, it is symplectic in the sense that the spatial symplectic form σ ,

$$\sigma = d\xi_r \wedge dr + d\xi_\theta \wedge d\theta + d\xi_\varphi \wedge d\varphi, \quad (3)$$

is nondegenerate on the surfaces $K \cap \{t = \text{const}\}$.

Let $\mathcal{C}_+ \subset \{G = 0\}$ be the positive light cone and

$$\varphi^t: \mathcal{C}_+ \rightarrow \mathcal{C}_+$$

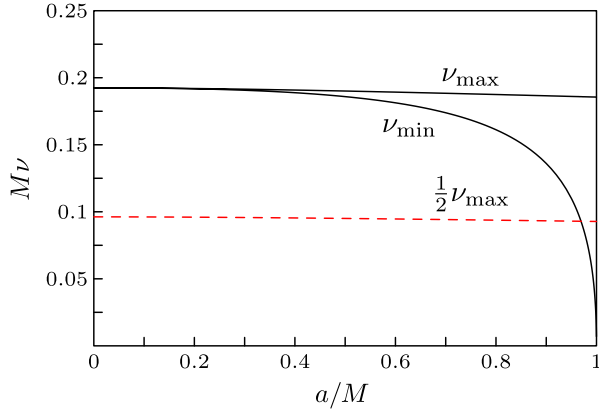


FIG. 3 (color online). The dependence of ν_{\max} and ν_{\min} on the parameters M and a in the case of $\Lambda = 0$. The dashed line indicates the range of validity of the pinching condition needed for the Weyl law (12).

the geodesic flow parametrized by t . The r -normal hyperbolicity condition asserts the existence of a C^r (r -times differentiable) splitting

$$T_K \mathcal{C}_+ = TK \oplus \mathcal{V}_+ \oplus \mathcal{V}_-,$$

invariant under the flow and such that for some constants $\nu > 0$, $C > 0$,

$$\begin{aligned} \sup_{(x,\xi) \in K} |d\varphi^{\mp t}|_{\mathcal{V}_\pm} &\leq C e^{-\nu t}, \\ \sup_{(x,\xi) \in K} |d\varphi^{\pm t}|_{TK} &\leq C e^{\nu|t|/r}, \quad t \geq 0. \end{aligned} \quad (4)$$

This means that the maximal expansion rates (Lyapunov exponents) on the trapped set are r -fold dominated by the expansion and contraction rates in the directions transversal to the trapped set. As shown in [6,24], Sec. 5.2, r -normal hyperbolicity is stable under perturbations: when G_ε is a time-independent (that is, stationary) Hamiltonian such that G_ε is close to G in C^r near K , the flow for G_ε is r -normally hyperbolic in the sense that the trapped set K_ε has C^r regularity and is symplectic and (4) holds. For Kerr (–de Sitter) metrics the flow is r -normally hyperbolic for all r as shown in [7,8,25], essentially because the flow on K is completely integrable.

Key dynamical quantities are the minimal and maximal expansion rates $0 < \nu_{\min} \leq \nu_{\max}$, characterized by inequalities true for all $\varepsilon > 0$, a constant C_ε depending on ε ,

$$C_\varepsilon^{-1} e^{-(\nu_{\max} + \varepsilon)t} \leq |d\varphi^t|_{\mathcal{V}_-} \leq C_\varepsilon e^{-(\nu_{\min} - \varepsilon)t},$$

$t > 0$. For Kerr(–de Sitter) metrics, the quantities ν_{\min} , ν_{\max} are obtained by taking the minimum and maximum of averages of the local expansion rate

$$\nu = \sqrt{-2\Delta_r \partial_r^2 G / |\partial_{\xi_t} G|},$$

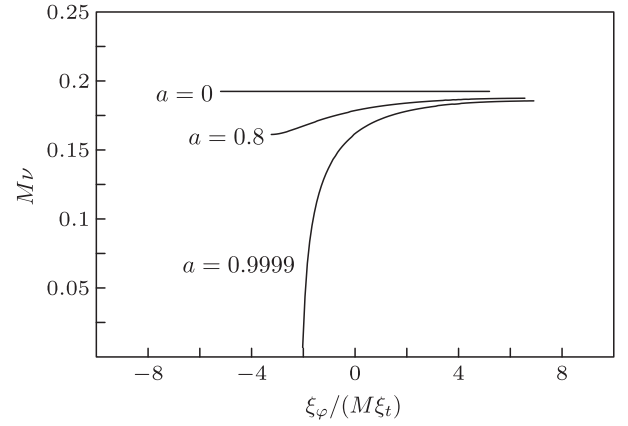


FIG. 4. The pointwise expansion rates ν on Liouville tori $\xi_\varphi / (M\xi_t) = \text{const}$, for $\theta = \pi/2$. When a approaches 1, $\nu = 0$ for some values of ξ_φ which shows that there is no gap and QNMs can be arbitrarily close the real axis [22,23].

on the Liouville tori [5] of the flow of G_θ on the trapped set.

III. DEFINITION AND DISCRETENESS OF QUASINORMAL MODES

The *scattering resonances*, called QNMs in the context of black holes [12], replace eigenmodes when one switches from closed systems to open systems—see [14] for a recent experimental discussion. They are the frequencies ω of oscillating solutions to the wave equation

$$\square_g(e^{-i\omega t} v(r, \theta, \varphi)) = 0, \quad (5)$$

which continue smoothly across the event horizons.

Solutions to $\square_g u = 0$ are expected to have expansions

$$u(t, r, \theta, \varphi) \sim \sum_k e^{-it\omega_k} u_k(r, \theta, \varphi) \quad (6)$$

valid in a suitable sense [21,31]. The fact that QNMs ω_k form a discrete set in the lower half plane is nontrivial but it is now rigorously known in the case of Kerr–de Sitter and its perturbations [18,20,25,31,32].

In the simpler Schwarzschild–de Sitter case we indicate the reason for discreteness of the set of QNMs as follows. Equation (5) can be rewritten as $P(\omega)v = 0$, where $P(\omega)$ is obtained from $-\rho^2 \square_g$ by replacing ∂_t with $-i\omega$. The operator $P(\omega)$ is spherically symmetric; its restriction to the space of spherical harmonics with eigenvalue $\ell(\ell + 1)$, written in the Regge–Wheeler coordinate x [20], Sec. 4, is the Schrödinger operator

$$P_\ell(\omega) = -\partial_x^2 + \omega^2 V_1(x) + \ell(\ell + 1)V_2(x), \quad (7)$$

where the potentials V_1 and V_2 are real analytic (their Taylor series converge to their values) and satisfy

$$V_1(x) = -V_\pm^2 + \mathcal{O}(e^{-A_\pm|x|}), \quad V_2(x) = \mathcal{O}(e^{-A_\pm|x|}),$$

as $x \rightarrow \pm\infty$; here $A_{\pm} > 0$, $V_- = r_-^2$, and $V_+ = r_+^2$. (When $\Lambda = 0$ then $V_2 \sim x^{-2}$ as $x \rightarrow +\infty$ and that creates problems at low energies [27,28]. More precisely, it is expected that due to the slow decay of the potential, the resolvent is not holomorphic in a neighborhood of zero and even in simplest cases such as Schwarzschild, there is no mathematical argument excluding the possibility of accumulation of QNMs at zero.) A number $\omega \in \mathbb{C}$ is a QNM if there exists an angular momentum ℓ and a nonzero solution $v(x)$ to the equation $P_{\ell}(\omega)v = 0$ satisfying the outgoing condition: near $x = \pm\infty$, $e^{\mp iV_{\pm}\omega x}v(x)$ is a smooth function of $e^{-A_{\pm}|x|}$. The outgoing condition follows naturally from the requirement that $e^{-i\omega t}v(x)$ extends smoothly past the event horizon of the black hole. For fixed ℓ , it follows by standard one-dimensional methods that the set of all corresponding ω is discrete.

Showing that as $\ell \rightarrow \infty$, quasinormal modes corresponding to different values of ℓ do not accumulate is more delicate: we need to know that if $|\omega| \leq R$ and ℓ is large enough depending on R , then ω cannot be a QNM corresponding to ℓ . Assume the contrary and let $v(x)$ be the corresponding solution to the equation $P_{\ell}(\omega)v = 0$. We fix large $X > 0$ independently of ℓ , to be chosen later. The potential V_2 is everywhere positive; therefore for ℓ large enough depending on R , X , $\text{Re}(\omega^2 V_1(x) + \ell(\ell + 1)V_2(x)) > 0$ for $x \in [-X, X]$. If v satisfied a Dirichlet or Neumann boundary condition at $\pm X$, then integration by parts would give the impossible statement that $v = 0$ on $[-X, X]$, finishing the proof:

$$\begin{aligned} 0 &= \text{Re} \int_{-X}^X \bar{v} \cdot P_{\ell}(\omega)v dx \\ &= -\text{Re}(v'\bar{v})|_{-X}^X \\ &\quad + \int_{-X}^X |v'|^2 + \text{Re}(\omega^2 V_1 + \ell(\ell + 1)V_2)|v|^2 dx = 0, \end{aligned}$$

and the terms under the integral are all non-negative. This argument works also for v 's satisfying the defining properties of the QNMs, as described below. For Kerr–de Sitter black holes a separation procedure is still possible but it does not work for stationary perturbations. Nevertheless in both cases the discreteness of QNMs remains valid [24,25].

To indicate how this works for resonant states which not satisfy a boundary condition at $\pm X$ (after all, this ‘‘boundary’’ is completely artificial), we follow [20], Sec. 6. It suffices to prove the boundary inequalities

$$\pm \text{Re}(v'(\pm X)\overline{v(\pm X)}) < 0. \quad (8)$$

To prove (8), we cannot use integration by parts on the whole \mathbb{R} , since v does not lie in $L^2(\mathbb{R})$ and moreover the real part of our potential may become negative as $x \rightarrow \pm\infty$. We instead use the methods of complex analysis and real analyticity of V_1, V_2 .

The characterization of v as a mode can be strengthened to say that $e^{\mp iV_{\pm}\omega x}v(x)$ is a *real analytic* function of $e^{-A_{\pm}|x|}$

(it has a convergent Taylor series in that variable), which means that for X large enough, we can extend $v(x)$ to a *holomorphic* function in $\{|\text{Re}z| \geq X\}$, and this extension is Floquet periodic: $v(z + 2\pi i/A_{\pm}) = e^{\mp 2\pi V_{\pm}\omega/A_{\pm}}v(z)$, $\pm \text{Re}z \geq X$. Now, consider the restriction of u to the vertical lines $\{\text{Re}z = \pm X\}$, $w_{\pm}(y) := v(\pm X + iy)$, $y \in \mathbb{R}$, and note that it solves the differential equation

$$(\partial_y^2 + \omega^2 V_1(\pm X + iy) + \ell(\ell + 1)V_2(\pm X + iy))w_{\pm} = 0. \quad (9)$$

The key difference between (9) and the equation $P_{\ell}(\omega)v = 0$ is that the potential $V_2(\pm X + iy)$ is no longer real valued. For instance, if V_2 were equal to $e^{-A_{\pm}|x|}$, then $V_2(\pm X + iy)$ would equal $e^{-XA_{\pm}}e^{\mp iA_{\pm}y}$, only taking real values when $y \in \pi A_{\pm}^{-1}\mathbb{Z}$. This means that Eq. (9) is *elliptic* (in the semiclassical sense, where we treat ∂_y as having same order as ℓ) except at a discrete set of points in the phase space $T^*\mathbb{R}$. Further analysis shows that $w_{\pm}(y)$ is concentrated in phase space near $y \in 2\pi A_{\pm}^{-1}\mathbb{Z}$, $\eta = \mp \ell\sqrt{V_2(\pm X)}$, in particular implying

$$|(\partial_y \pm i\ell\sqrt{V_2(\pm X)})w_{\pm}(0)| \leq C\ell^{1/4}|w_{\pm}(0)|,$$

and (8) follows from Cauchy-Riemann equations, since $v(\pm X) = w_{\pm}(0)$ and $v'(\pm X) = -iw'_{\pm}(0)$.

IV. DISTRIBUTION OF QUASINORMAL MODES

The distribution of QNMs ω_k can now be studied in the more general stable setting of r -normally hyperbolic trapped sets. Three fundamental issues are

- (a) distribution of decay rates, that is of the imaginary parts of QNMs;
- (b) asymptotics of the counting function;
- (c) expansion of waves in terms of QNMs.

(a).—We can bound the decay rates from below whenever the trapped set is normally hyperbolic, without requiring the stronger r -normal hyperbolicity assumption. The bound [33] is given by $\text{Im}\omega_k < -(\nu_{\min} - \varepsilon)/2$, for any $\varepsilon > 0$, once the frequency (the real part of ω_k) is large enough. In the case of r -normal hyperbolicity and under the pinching condition

$$\nu_{\max} < 2\nu_{\min}, \quad (10)$$

we get more detailed information [24]: there are additionally no QNMs with

$$-(\nu_{\min} - \varepsilon) < \text{Im}\omega_k < -(\nu_{\max} + \varepsilon)/2. \quad (11)$$

That means that the modes with least decay are confined to a band shown in Fig. 6. In the completely integrable case this follows from WKB constructions [20–23]—see Fig. 6—but this structure persists under perturbations. Figure 5 shows the accuracy of the estimate (11) for the numerically computed QNMs of exact Kerr black holes [9]. For a recent experimental investigation of the

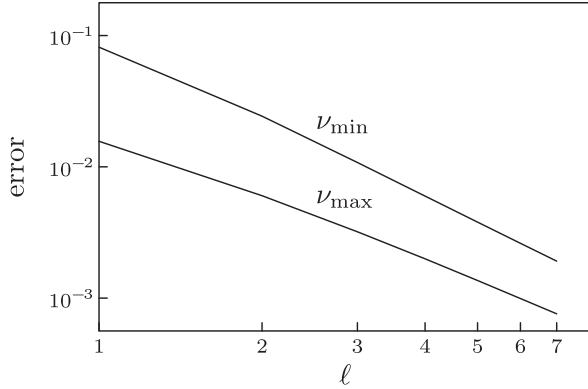


FIG. 5. A log-log plot of relative errors $\frac{\min |\operatorname{Im} \omega_k(\ell) - \nu_{\min}/2|}{\min |\operatorname{Im} \omega_k(\ell)|}$ and $\frac{\max |\operatorname{Im} \omega_k(\ell) - \nu_{\max}/2|}{\max |\operatorname{Im} \omega_k(\ell)|}$ where $\omega_k(\ell)$ are the numerically computed resonances in the first band corresponding to the angular momentum ℓ [9] and ν_{\min}, ν_{\max} are minimal and maximal expansion rates. The agreement is remarkable when ℓ increases.

distribution of decay rates and the relation to classical dynamics (more precisely, the topological pressure and classical escape rates) see [15].

The condition (10) is called pinching because it pinches the ratio of the maximal and the minimal transversal expansion rates. In the absence of this condition, the gap (11) between the first band of QNMs and the faster decaying bands disappears, making it difficult to obtain a counting asymptotics (12). Physically, (10) could be interpreted as the requirement that there be no interaction between QNMs from different bands.

(b).—The relation between the density of high energy states and phase space volumes defined by the classical

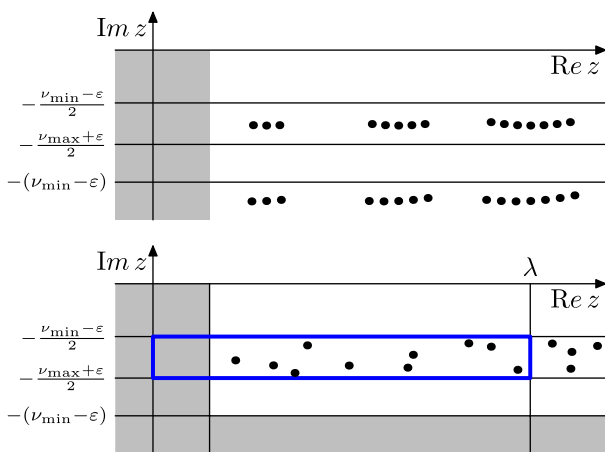


FIG. 6 (color online). A schematic comparison between QNMs in a completely integrable case and in the general r -normally hyperbolic case. The former lie on a fuzzy lattice and are well approximated by WKB construction based on quantization conditions [20–23]. When trapping is r -normally hyperbolic and the pinching condition $\nu_{\max} < 2\nu_{\min}$ holds, quasinormal modes are still localized to a strip with dynamically determined bounds, and their statistics are given by the Weyl law (12).

Hamiltonian is one of the basic principles of quantum mechanics or spectral theory. It states that for closed systems the number, $N_{\hat{H}}(\lambda)$, of energy levels of \hat{H} , a quantization of H (for instance the Dirichlet Laplacian on a bounded domain), below energy λ^2 (we think of λ as frequency which is natural when considering QNMs) satisfies the *Weyl law*

$$N_{\hat{H}}(\lambda) \sim (2\pi)^{-\dim X} \operatorname{vol}_{T^*X}(H \leq \lambda^2) \sim C_{\hat{H}} \lambda^{-\dim X}.$$

Here $\operatorname{vol}_{T^*X}$ denotes the phase space volume calculated using the volume $\sigma^{\dim X}/(\dim X)!$ obtained from the symplectic form σ [see (3)].

For open systems QNMs replace real energy levels and the counting becomes much more tricky [14]. In the case of exact Kerr(–de Sitter) black holes the WKB constructions can be used to show that the number, $N_{\text{QNM}}(\lambda)$, of QNMs with

$$|\omega_k| \leq \lambda, \quad \operatorname{Im} \omega_k \geq -(\nu_{\min} - \varepsilon)$$

satisfies the asymptotic law $N_{\text{QNM}}(\lambda) \sim c \lambda^2$. The constant c has a geometric interpretation: in a scattering problem the total phase space T^*X (the cotangent bundle of X) is replaced by the trapped set [34], and c corresponds to the symplectic volume of the trapped set.

The same law is proved [8] for perturbations of Kerr–de Sitter, using completely different ideas based on r -normal hyperbolicity rather than symmetries of the metric and separation of variables. Under the assumptions of r -normal hyperbolicity (4) and pinching (10), we have

$$N_{\text{QNM}}(\lambda) \sim \frac{\lambda^2}{(2\pi)^2} \operatorname{vol}(K \cap \{\xi_i^2 \leq 1\} \cap \{t = 0\}), \quad (12)$$

where the volume is taken using the symplectic form on $K \cap \{t = \text{const}\}$ [8], Theorem 3. We note that just as $\dim X = \frac{1}{2} \dim T^*X$ in the exponent of the Weyl law, here

$$2 = \frac{1}{2} \dim(K \cap \{t = \text{const}\});$$

that is, the effective phase space is now the trapped set. For exact Kerr(–de Sitter) metrics with several values of Λ the volume as function of a is shown in Fig. 7. The volume is finite provided that $(1 - \Lambda a^2/3)^3 > 9\Lambda M^2$; see Fig. 1.

We should stress that normally hyperbolic behavior (unlike r -normal hyperbolicity) is often unstable under perturbations as shown by examples of hyperbolic quotients where a small perturbation can change the dimension of the trapped set (Fig. 1 in [35]), leading to a fractal Weyl law, that is a law in which the exponent 2 changes to half of the fractal dimension of the trapped set—see [14] for recent experiments on that.

(c).—The expansion (6) is rigorously established for slowly rotating black holes [21,31] and heuristically it is

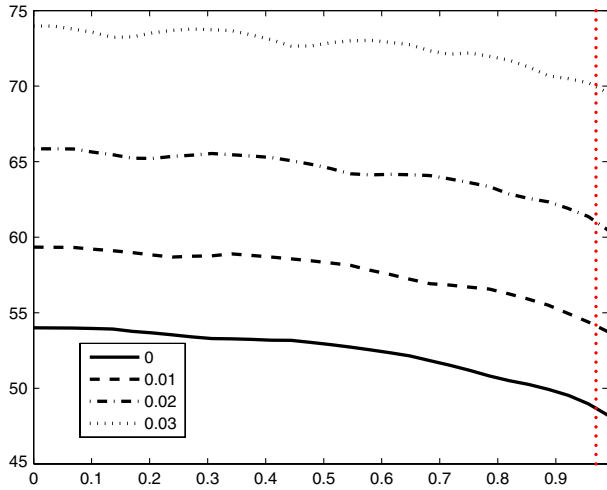


FIG. 7 (color online). Numerically computed constant in (12), $\text{vol}(K \cap \{\xi_t^2 \leq 1\} \cap \{t = 0\})/4\pi^2$, for $\Lambda = 0, 0.01, 0.02, 0.03$ and $M = 1$. The vertical line shows the value of a at which the pinching condition (10) fails for $\Lambda = 0$. For smaller values of a the Weyl law (12) holds. We note that as a increases the QNMs split (see Fig. 6) and hence we do expect the counting function to decrease, in agreement with the behavior of the volume. From the volume and the gap giving the decay rate one can read off a and M .

one of the motivations for studying quasinormal modes [12]. For rapidly rotating black holes, or for their perturbations satisfying (10) a more robust version can be formulated using projector onto the states associated to quasinormal modes in the first band shown in Fig. 6. The solution to the wave equation $\square_g u = 0$ with initial data localized near frequency $\lambda \gg 1$ can be decomposed as

$$u = u_{\text{QNM}} + u_{\text{DEC}},$$

where, for $0 \leq t \leq T \log \lambda$,

$$\square_g u_{\text{QNM}}(t), \quad \square_g u_{\text{DEC}}(t) = \mathcal{O}(\lambda^{-\infty});$$

that is, we have rapid decay (faster than any negative power) when the frequency λ is large. This means that both terms solve the wave equation approximately at high energies times bounded logarithmically in λ . We then have, again for $0 \leq t \leq T \log \lambda$,

$$\begin{aligned} \|u_{\text{QNM}}(t)\|_{\mathcal{E}} &\leq C e^{-(\nu_{\min} - \varepsilon)t/2} \|u_{\text{QNM}}(0)\|_{\mathcal{E}}, \\ \|u_{\text{QNM}}(t)\|_{\mathcal{E}} &\geq C^{-1} e^{-(\nu_{\max} + \varepsilon)t/2} \|u_{\text{QNM}}(0)\|_{\mathcal{E}}, \\ \|u_{\text{QNM}}(0)\|_{\mathcal{E}} &\leq C \sqrt{\lambda} \|u(0)\|_{\mathcal{E}}, \\ \|u_{\text{DEC}}(t)\|_{\mathcal{E}} &\leq C \lambda e^{-(\nu_{\min} - \varepsilon)t} \|u(0)\|_{\mathcal{E}}, \end{aligned} \quad (13)$$

where strictly speaking errors $\mathcal{O}(\lambda^{-\infty}) \|u(0)\|_{\mathcal{E}}$ should be added to the right-hand sides. The norm $\|\cdot\|_{\mathcal{E}}$ is the standard energy norm in any sufficiently large compact

subset of X [in the case of exact Kerr–de Sitter, we can take $(r_+ + \delta, r_C - \delta) \times \mathbb{S}^2$]—see [8], Theorem 2. The term $u_{\text{QNM}}(t)$ corresponds to the part of the solution dominated by the QNMs in the first band and it has the natural decay properties dictated by the imaginary parts of these QNMs. In fact, $u_{\text{QNM}}(t)$ can be physically interpreted as the radiation coming from light rays traveling along the trapped set. The directions in which such a light ray radiates towards infinity can be described in terms of the geometry of the flow, and the amplitude of the radiated waves can be calculated using the global dynamics of the flow near the trapped set [24], Sec. 8.5.

V. CONCLUSIONS

We have shown that for Kerr–de Sitter metrics and their perturbations quasinormal modes are rigorously defined and form a discrete set in the lower half plane, provided that the parameters of the black hole satisfy

$$(1 - \Lambda a^2/3)^3 > 9\Lambda M^2 > 0;$$

see Fig. 1. This is due to the size of infinity when $\Lambda > 0$ and the compactness of the trapped set at finite energies.

If one neglects the issues of long time decay and of behavior at low energies, then the results are also valid in the case of $\Lambda = 0$. On the length scales involved in the ringdown phenomenon, which in principle would lead to the detection of black hole parameters through QNMs, the (small) value of Λ is not relevant but $\Lambda > 0$ is a more convenient mathematical model.

The main dynamical feature of the set on which photons are trapped (the trapped set) is its r -normal hyperbolicity for any r —see (4). Because of the stability of this property the main features of the distribution of quasinormal modes are preserved for perturbations: the decay rates are bounded from below in terms of the minimal expansion rate [$\text{Im}\omega_k \leq -(\nu_{\min} - \varepsilon)/2$] and under the pinching conditions, the least decaying modes are confined to a strip where they satisfy a counting law (12)—see Fig. 6.

The r -normal hyperbolicity is valid for all rotating black holes but the pinching condition (10) needed for the finer results (12) and (13) fails in the case of very fast rotation—see Fig. 3.

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