



A Scattering Theory for Linear Waves on the Interior of Reissner–Nordström Black Holes

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Abstract. We develop a scattering theory for the linear wave equation $\square_g \psi = 0$ on the interior of Reissner–Nordström black holes, connecting the fixed frequency picture to the physical space picture. Our main result gives the existence, uniqueness and asymptotic completeness of finite energy scattering states. The past and future scattering states are represented as suitable traces of the solution ψ on the bifurcate event and Cauchy horizons. The heart of the proof is to show that after separation of variables one has *uniform boundedness* of the reflection and transmission coefficients of the resulting radial o.d.e. over all frequencies ω and ℓ . This is non-trivial because the natural T conservation law is sign-indefinite in the black hole interior. In the physical space picture, our results imply that the Cauchy evolution from the event horizon to the Cauchy horizon is a Hilbert space isomorphism, where the past (resp. future) Hilbert space is defined by the finiteness of the degenerate T energy fluxes on both components of the event (resp. Cauchy) horizon. Finally, we prove that, in contrast to the above, for a generic set of cosmological constants Λ , there is no analogous finite T energy scattering theory for either the linear wave equation or the Klein–Gordon equation with conformal mass on the (anti) de Sitter–Reissner–Nordström interior.

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1. Introduction

One of the most stunning predictions of general relativity is the formation of **black holes**, defined by the property that information cannot propagate from their interior region to outside far-away observers. Fortunately, we can count ourselves among the latter; nevertheless, if a group of physicists were so courageous as to cross the *event horizon* and enter a black hole, they could still

very well perform experiments and compare the outcomes among themselves. Indeed, the problem of determining the fate of these black hole explorers (and their laboratories) has led to some of the most central conceptual puzzles in gravitational physics.

In view of the above, there has been a lot of recent activity analyzing the Cauchy problem on black hole interiors, e.g., [15–17, 31, 47]. However, for certain physical processes it is more natural to consider the **scattering problem** (see [18] for scattering on the exterior of black holes). With this paper, we initiate the mathematical study of the finite energy scattering problem on black hole interiors. Specifically, we will consider solutions of the wave equation on what can be viewed as the most elementary interior, that of Reissner–Nordström. The Reissner–Nordström metrics constitute a family of space–times, parameterized by mass M and charge Q , which satisfy the Einstein–Maxwell system in spherical symmetry [41, 45] and admit an additional Killing vector field T . For vanishing charge $Q = 0$, the family reduces to Schwarzschild. We shall moreover restrict in the following to the subextremal case where $0 < |Q| < M$. In addition to the bifurcate event horizon, these black hole interiors then admit an additional bifurcate inner horizon, the so-called *Cauchy horizon*. Our past and future scattering states will be defined as suitable traces of the solution on the bifurcate event horizon and bifurcate Cauchy horizon, respectively, restricted to have finite L^2 energy ³ and on each component of the horizons.

In the rest of the introduction we will state our main results for the scattering problem on the interior of Reissner–Nordström (Theorems 1–5), relate them to existing literature on fixed frequency scattering, and draw links to various recent results in the interior and exterior of black holes. Finally, we will see that the existence of a bounded scattering map for the wave equation on Reissner–Nordström turns out to be a very fragile property; we shall show that there does not exist an analogous scattering theory in the presence of a cosmological constant (Theorem 6) or Klein–Gordon mass (Theorem 7).

The scattering problem on Reissner–Nordström interior. In this paper, we will establish a scattering theory for finite energy solutions of the linear wave equation

$$\square_g \psi = 0, \quad (1.1)$$

on the interior of a Reissner–Nordström black hole, from the bifurcate event horizon $\mathcal{H} = \mathcal{H}_A \cup \mathcal{H}_B \cup \mathcal{B}_-$ to the bifurcate Cauchy horizon $\mathcal{CH} = \mathcal{CH}_A \cup \mathcal{CH}_B \cup \mathcal{B}_+$, as depicted in Fig. 1. The first main result of our paper is Theorem 1 (see Sect. 3.1) in which we will show *existence*, *uniqueness* and *asymptotic completeness* of finite energy scattering states. In this context, *existence* and *uniqueness* mean that for given finite energy data ψ_0 on the event horizon \mathcal{H} , there exist unique finite energy data on the Cauchy horizon \mathcal{CH} arising from ψ_0 as the evolution of (1.1). With *asymptotic completeness* we denote the property that all finite energy data on the Cauchy horizon \mathcal{CH} can indeed be achieved from finite energy data on the event horizon \mathcal{H} . This provides a way to construct solutions with desired asymptotic properties which is a

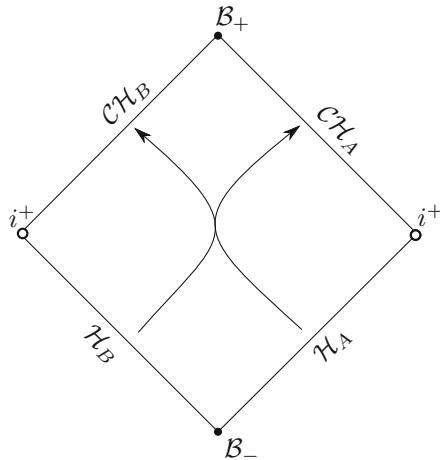


FIGURE 1. Penrose diagram of the interior of the Reissner–Nordström black hole and visualization of the scattering map

necessary first step to properly understand quantum theories in the interior of a Reissner–Nordström black hole (cf. [44, 23, 51]). The energy spaces on the event and Cauchy horizon are associated with the Killing field and generator of the time translation T . Indeed, T is null on the horizons and, in particular, is the generator of the event and Cauchy horizon \mathcal{H} and \mathcal{CH} . Because of the sign-indefiniteness of the energy flux of the vector field T on the bifurcate event (resp. Cauchy) horizon [see already (1.4)], we define our energy space by requiring the finiteness of the T energy on both components separately of the event (resp. Cauchy) horizon. These define Hilbert spaces with respect to which the scattering map is proven to be bounded.

Finally, it is instructive to draw a comparison between the interior of Reissner–Nordström and the interior of Schwarzschild ($Q = 0$). As opposed to Reissner–Nordström discussed above, the Schwarzschild interior terminates at a singular boundary at which solutions to (1.1) generically blow up (see [11]). In contrast, the non-singular and, moreover, Killing, Cauchy horizons (see Fig. 1) of Reissner–Nordström immediately yield natural Hilbert spaces of finite energy data to consider. In view of this, Reissner–Nordström with $Q \neq 0$ can be considered the most elementary interior on which to study the scattering problem. Furthermore, in view of the recent work [7], we have that the causal structure of Reissner–Nordström is stable in a weak sense (see the discussion below about related works in the interior).

Fixed frequency scattering. It is well known that the wave equation (1.1) on Reissner–Nordström spacetime allows separation of variables which reduces it to the radial o.d.e.

$$u'' - V_\ell u + \omega^2 u = 0, \quad (1.2)$$

with potential V_ℓ [see already (2.37)], where $\omega \in \mathbb{R}$ is the time frequency and $\ell \in \mathbb{N}_0$ is the angular parameter. Indeed, most of the existing literature concerning scattering of waves in the interior of Reissner–Nordström mainly considers fixed frequency solutions, e.g., [5, 21, 22, 33–35, 52]. For a purely *incoming* (i.e., supported only on \mathcal{H}_A) fixed frequency solution with parameters (ω, ℓ) , we can associate transmission and reflection coefficients $\mathfrak{T}(\omega, \ell)$ and $\mathfrak{R}(\omega, \ell)$. The transmission coefficient $\mathfrak{T}(\omega, \ell)$ measures what proportion of the incoming wave is transmitted to \mathcal{CH}_B , whereas the reflection coefficient specifies the proportion reflected to \mathcal{CH}_A . An essential step to go from fixed frequency scattering to physical space scattering is to prove **uniform boundedness** of $\mathfrak{T}(\omega, \ell)$ and $\mathfrak{R}(\omega, \ell)$. This is non-trivial in view of the discussion of the energy identity (1.4). In this paper, we indeed obtain this uniform bound in Theorem 1 (see Sect. 3.2). In particular, the regime $\omega \rightarrow 0, \ell \rightarrow \infty$ is the most involved frequency range to prove uniform boundedness. As we shall see, the proof relies on an explicit computation at $\omega = 0$ (see [21]) together with a careful analysis of special functions and perturbations thereof.

The uniform boundedness of the scattering coefficients is the main ingredient to prove the boundedness of the scattering map in Theorem 1. Moreover, it allows us to connect the separated picture to the physical space picture by means of a Fourier representation formula. This is stated as Theorem 3 (see Sect. 3.3). A somewhat surprising, direct consequence of the Fourier representation of the scattered data on the Cauchy horizon is that purely incoming compactly supported data lead to a solution which vanishes at the future bifurcation sphere \mathcal{B}_+ . This is moreover shown to be a necessary condition for the existence of a bounded scattering map (Corollary 3.1).

Comparison to scattering on the exterior of black holes. On the exterior of black holes, the scattering problem has been studied more extensively; see the pioneering works [2, 3, 11, 13] the book [18] and related results on conformal scattering in [32, 37, 4, 49]. Note that for the exterior of a Schwarzschild or Reissner–Nordström black hole, the uniform boundedness of the scattering coefficients or, equivalently, the boundedness of the scattering map, can be viewed a posteriori¹ as a consequence of the global T energy identity

$$\int_{\mathcal{H}^-} |T\psi|^2 + \int_{\mathcal{I}^-} |T\psi|^2 = \int_{\mathcal{H}^+} |T\psi|^2 + \int_{\mathcal{I}^+} |T\psi|^2. \quad (1.3)$$

Considering only incoming radiation from \mathcal{I}^- , this statement translates into $|\mathfrak{R}|^2 + |\mathfrak{T}|^2 = 1$ for the reflection coefficient \mathfrak{R} and transmission coefficients \mathfrak{T} . In the interior, however, due to the different orientations of the T vector field on the horizons (cf. Fig. 2), boundedness of the scattering map is not at all manifest. In particular, the global T energy identity on the interior of a Reissner–Nordström black hole reads

$$\int_{\mathcal{H}_A} |T\psi|^2 - \int_{\mathcal{H}_B} |T\psi|^2 = \int_{\mathcal{CH}_B} |T\psi|^2 - \int_{\mathcal{CH}_A} |T\psi|^2 \quad (1.4)$$

¹Note that proving (1.3) requires first establishing some form of qualitative decay toward i^+ and i^- .

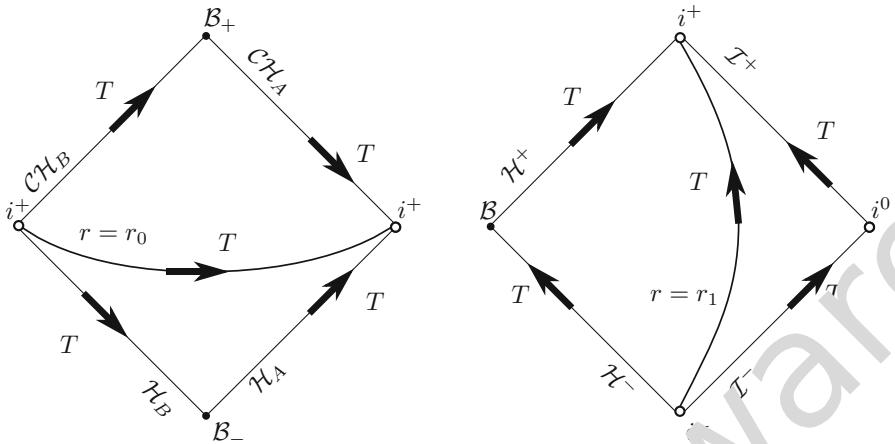


FIGURE 2. Interior of Reissner–Nordström (left) and exterior of Schwarzschild or Reissner–Nordström (right). In both diagrams, the arrows denote the direction of the T Killing vector field. Note that we have the identifications $\mathcal{H}_A = \mathcal{H}^+$ and $\mathcal{B}_- = \mathcal{B}$.

from which we cannot deduce boundedness of the scattering map even a posteriori. (Indeed, identity (1.4) corresponds only to the “pseudo-unitarity” statement of Theorem 1.) Again, considering only ingoing radiation from \mathcal{H}_A , this translates to

$$|\Sigma(\omega, \ell)|^2 - |\Re(\omega, \ell)|^2 = 1 \quad (1.5)$$

for the reflection coefficient \Re and the transmission coefficient \mathfrak{T} . Hence, while for fixed $|\omega| > 0$ and ℓ it is straightforward to show that \mathfrak{T} and \Re are finite, there is no a priori obvious obstruction from (1.5) for these scattering coefficients to blow up in the limits $\omega \rightarrow 0, \pm\infty$ and $\ell \rightarrow \infty$.

Moreover, note that in the exterior, the Killing field T is timelike, so the radial o.d.e. (1.2) should be considered as an equation for a fixed time frequency wave on a constant time slice. In the interior, however, the Killing field T is spacelike, so the radial o.d.e. (1.2) is rather an evolution equation for a constant spatial frequency.

The Schwarzschild family can be viewed as a special case ($a = 0$) of the two-parameter Kerr family, describing rotating black holes with mass parameter M and rotation parameter a [26].² On the exterior of Kerr many other difficulties arise: superradiance, intricate trapping, presence of ergoregion, etc. [8]. Nevertheless, using the decay results in [8], a definitive physical space scattering theory for Kerr black holes has been established in [9] (see also the earlier [19]). The proof on the exterior of Kerr involved first establishing a scattering

²Both Kerr and Reissner–Nordström can be viewed as special cases of the Kerr–Newman spacetime. For decay results on Kerr–Newman, see [6].

map from past null infinity \mathcal{I}^- to a constant time slice Σ and then concatenating that map with a second scattering map from Σ to the future event horizon \mathcal{H}^+ and future null infinity \mathcal{I}^+ . In the interior, however, we will directly show the existence of a “global” scattering map from the event horizon \mathcal{H} to the Cauchy horizon \mathcal{CH} . Indeed, due to blueshift instabilities (see [10]), we do not expect that the analogous concatenation of scattering maps (event horizon \mathcal{H} to spacelike hypersurface Σ and then from Σ to the Cauchy horizon \mathcal{CH}) as in the Kerr exterior, to be bounded in the interior of Reissner–Nordström.

Injectivity of the reflection map and blueshift instabilities. We can also conclude from our work that there is always non-vanishing reflection to the Cauchy horizon \mathcal{CH}_A arising from non-vanishing purely ingoing radiation at \mathcal{I}^-_A . This follows from the fact that in the separated picture and for fixed ℓ , the reflection coefficient $\Re(\omega, \ell)$ can be analytically continued to the strip $|\text{Im}(\omega)| < \kappa_+$ and hence, only vanishes on a discrete set of points on the real axis. This is shown in Theorem 4 (see Sect. 3.4).

We will also deduce from the Fourier representation of the scattered data on the Cauchy horizon \mathcal{CH} , and a suitable meromorphic continuation of the transmission coefficient, that there exist purely incoming compactly supported data on the event horizon \mathcal{H} leading to solutions which fail to be C^1 on the Cauchy horizon \mathcal{CH} . This C^1 -blowup of linear waves puts on a completely rigorous footing the pioneering work of Chandrasekhar and Hartle [5]. We state this as Theorem 5 (see Sect. 3.5).

For generic solutions arising from localized data on an asymptotically flat hypersurface, one expects a stronger instability, namely non-degenerate energy blowup at the Cauchy horizon \mathcal{CH} . Such non-degenerate energy blowup was proven in [27] for generic compactly supported data on an asymptotically flat Cauchy hypersurface. Later, for the more difficult Kerr interior, non-degenerate energy blowup was proven in [31] assuming certain polynomial lower bounds on the tail of incoming data on the event horizon \mathcal{H} and in [10] for solutions arising from generic initial data along past null infinity \mathcal{I}^- with polynomial tails.

Finally, we mention the forthcoming work [30] which studies the problem of non-degenerate energy blowup from a scattering theory perspective and also uses the non-triviality of reflection to establish results related to mass inflation for the spherically symmetric Einstein–Maxwell–scalar field system (cf. [28, 29]).

Related results on the interior. There has been a lot of recent progress studying the interior of black holes. In particular, new insights were gained concerning the stability of the Cauchy horizon and the strong cosmic censorship conjecture.

For the Cauchy problem for (1.1) on the interior of both a fixed Kerr and a Reissner–Nordström black hole, the works [16, 17, 24] establish uniform boundedness (in L^∞) and continuity up to and including the Cauchy horizon for solutions arising from smooth and compactly supported

data on an asymptotically flat spacelike hypersurface. Such data in particular give rise to solutions with polynomial decay along the event horizon.

In contrast, for the scattering problem considered in the present paper, we are required to work with spaces which are symmetric with respect to the event and Cauchy horizons. This naturally leads to the rougher class of finite T energy data in the statement of Theorem 1. Note that for such data on the Cauchy horizon, continuity or boundedness (in L^∞) does not necessarily hold true.

Turning finally to the full nonlinear dynamics of the Einstein equations, it is shown in [7] that the Kerr Cauchy horizon is C^0 -stable. Thus, the existence of a Cauchy horizon, a very natural setting parameterizing scattering data in the interior, is not a pure artifact of symmetry but rather a stable property at least in a weak sense. On the other hand, in [28, 29, 50] it is proven that for a suitable Einstein-matter system under spherical symmetry, the Cauchy horizon, while C^0 -stable, is generically C^2 -unstable. Finally, we mention that for the Schwarzschild black hole ($Q = 0$), which does not admit a Cauchy horizon, it is shown in [15] that solutions to (1.1) generically blow up at the spacelike singularity $\{r = 0\}$.

Breakdown of T energy scattering for $\Lambda \neq 0$ or $\mu \neq 0$. If a cosmological constant $\Lambda \in \mathbb{R}$ is added to the Einstein–Maxwell system, we can consider the analogous (anti-) de Sitter–Reissner–Nordström family of solutions whose interiors have the same Penrose diagram as depicted in Fig. 1. For very slowly rotating Kerr–de Sitter and Reissner–Nordström–de Sitter spacetimes, boundedness, continuity, and regularity up to and including the Cauchy horizon has been shown for solutions to (1.1) arising from smooth and compactly supported data on a Cauchy hypersurface [25]. However, remarkably, there is no analogous T energy scattering theory for either the linear wave equation (1.1) or the Klein–Gordon equation with conformal mass. This is the statement of Theorem 6 (see Sect. 3.6). The reason for this failure is the unboundedness of the transmission coefficient \mathfrak{T} and reflection coefficients \mathfrak{R} in the vanishing frequency limit $\omega \rightarrow 0$. To be more precise, we will prove that there does not exist a T energy scattering theory of the wave or Klein–Gordon equation in the interior of (anti-) de Sitter–Reissner–Nordström black hole for generic subextremal black hole parameters (M, Q, Λ) . In particular, for fixed $0 < |Q| < M$, there is an $\epsilon > 0$ such that there does not exist a T energy scattering theory for all $0 \neq |\Lambda| < \epsilon$.

Similarly, we prove in Theorem 7 (see Sect. 3.7) that there does not exist a T energy scattering theory for the Klein–Gordon equation $\square_g \psi - \mu \psi = 0$ on the Reissner–Nordström interior for a generic set of masses μ . This is in contrast to the exterior, where T energy scattering theories were established for massive fields in [3, 36].

It remains an open problem to formulate an appropriate scattering theory in the cosmological setting and for the Klein–Gordon equation as well as for the interior of Kerr.

Outline. This paper is organized as follows. In Sect. 2, we shall set up the spacetime, introduce the relevant energy spaces, and define the scattering coefficients in the separated picture. In Sect. 3, we state the main results of this paper, Theorems 1–7. Section 4 is devoted to the proof of Theorem 2. Then, the statement of Theorem 2 allows us to prove Theorem 1 in Sect. 5. Finally, in the last two sections are show our non-existence results: In Sect. 6, we prove Theorem 6 and in Sect. 7, we give the proof of Theorem 7.

2. Preliminaries

In this section, we will define the background differentiable structure and metric for the Reissner–Nordström spacetime and introduce some convenient coordinate systems.

2.1. Interior of the Subextremal Reissner–Nordström Black Hole

The global geometry of Reissner–Nordström was first described in [20]. For completeness, we will explicitly construct in this section the coordinates for the region considered. We start, in Sect. 2.1.1, by defining a Lorentzian manifold corresponding to the interior of the Reissner–Nordström black hole without the horizons. Then, in Sect. 2.1.2, we will attach the boundaries which will correspond to the event horizon and Cauchy horizon.

2.1.1. The Interior Without Boundary. We will now give an explicit description of the differential structure and metric. The Reissner–Nordström solutions [41, 45] are a two-parameter family of spherically symmetric spacetimes with mass parameter $M \in \mathbb{R}$ and electric charge parameter $Q \in \mathbb{R}$ solving the Einstein–Maxwell system

$$Ric_{\mu\nu} - \frac{1}{2}g_{\mu\nu}F_{\mu\nu} = 8\pi \left(\frac{1}{4\pi} \left(F_\mu^\lambda F_{\lambda\nu} - \frac{1}{4}g_{\mu\nu}F_{\lambda\rho}F^{\lambda\rho} \right) \right),$$

$$\nabla^\mu F_{\mu\nu} = 0, \quad \nabla_{[\mu}F_{\nu]\lambda} = 0. \quad (2.1)$$

In this paper, we consider the subextremal family of black holes with parameter range $0 < |Q| < M$. Define the manifold \mathcal{M} by

$$\mathcal{M} = \mathbb{R} \times (r_-, r_+) \times \mathbb{S}^2, \quad (2.2)$$

where $r_\pm = M \pm \sqrt{M^2 - Q^2} > 0$. The manifold is parameterized by the standard coordinates $t \in \mathbb{R}$, $r \in (r_-, r_+)$, and a choice of spherical coordinates (θ, ϕ) on the sphere \mathbb{S}^2 . We denote the global coordinate vector field ∂_t by T :

$$T := \frac{\partial}{\partial t}. \quad (2.3)$$

Using the above coordinates, we equip \mathcal{M} with the Lorentzian metric

$$g_{Q,M} = - \left(1 - \frac{2M}{r} + \frac{Q^2}{r^2} \right) dt \otimes dt + \left(1 - \frac{2M}{r} + \frac{Q^2}{r^2} \right)^{-1} dr \otimes dr + r^2 g_{\mathbb{S}^2}, \quad (2.4)$$

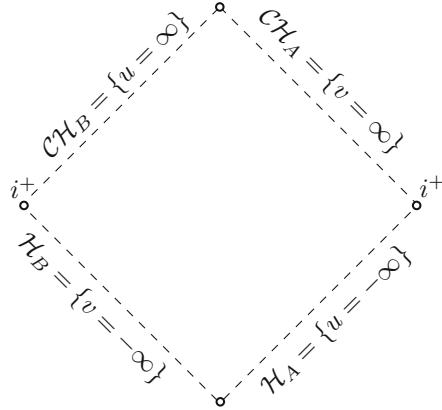


FIGURE 3. Penrose diagram of \mathcal{M} ; formally we have defined the boundary (not part of the manifold) by $\mathcal{C}\mathcal{H} = \mathcal{H}_A \cup \mathcal{H}_B$ and $\mathcal{CH} = \mathcal{CH}_A \cup \mathcal{CH}_B$

where g_{S^2} is the round metric on the 2-sphere. We also define the quantities

$$\Delta := r^2 - 2Mr + Q^2 = (r - r_+)(r - r_-), \quad \text{and} \quad h := \frac{\Delta}{r^2}. \quad (2.5)$$

Furthermore, define r_* by

$$dr_* := \frac{r^2}{\Delta} dr, \quad (2.6)$$

where we choose $r_*(\frac{r_+ + r_-}{2}) = 0$ for definiteness. Thus,

$$r_*(r) = r_+ - \frac{1}{r_+} \log |r - r_+| + \frac{1}{2\kappa_+} \log |r - r_-| + C \quad (2.7)$$

for a constant C only depending on the black hole parameters and

$$\kappa_\pm = \frac{r_\pm - r_\mp}{2r_\pm^2}. \quad (2.8)$$

We also introduce null coordinates

$$v = r_* + t \quad \text{and} \quad u = r_* - t \quad (2.9)$$

on \mathcal{M} . The Penrose diagram of \mathcal{M} is depicted in Fig. 3.

2.1.2. Attaching the Event and Cauchy Horizon. Now, we will attach the Cauchy and event horizon to the manifold. The Cauchy horizon characterizes the future boundary up to which the spacetime is uniquely determined as a solution to the Einstein–Maxwell system arising from data on the event horizon. Although the metric is smoothly extendible beyond the Cauchy horizon, any such extension fails to be uniquely determined from initial data, leading to a severe failure of determinism.

Attaching the event and Cauchy horizon gives rise to a manifold with corners. To do so, we first define the following two pairs of null coordinates.

Let $U_{\mathcal{H}}: \mathbb{R} \rightarrow (0, \infty)$ and $V_{\mathcal{H}}: \mathbb{R} \rightarrow (0, \infty)$ be smooth and strictly increasing functions such that

- $U_{\mathcal{H}}(u) = u$ for $u \geq 1$, $V_{\mathcal{H}}(v) = v$ for $v \geq 1$,
- $U_{\mathcal{H}}(u) \rightarrow 0$ as $u \rightarrow -\infty$, $V_{\mathcal{H}}(v) \rightarrow 0$ as $v \rightarrow -\infty$,
- there exists a $u_+ \leq 1$ such that $\frac{dU_{\mathcal{H}}}{du} = \exp(\kappa_+ u)$ for $u \leq u_+$,
- there exists a $v_+ \leq 1$ such that $\frac{dV_{\mathcal{H}}}{dv} = \exp(\kappa_+ v)$ for $v \leq v_+$.

This defines—in mild abuse of notation—the null coordinates $U_{\mathcal{H}}: \mathcal{M} - (0, \infty)$ via $U_{\mathcal{H}}(u)$ and $V_{\mathcal{H}}: \mathcal{M} \rightarrow (0, \infty)$ via $V_{\mathcal{H}}(v)$, where u, v are the null coordinates defined in (2.9).

Similarly, let $U_{\mathcal{CH}}: \mathbb{R} \rightarrow (-\infty, 0)$ and $V_{\mathcal{CH}}: \mathbb{R} \rightarrow (-\infty, 0)$ be smooth and strictly increasing functions such that

- $U_{\mathcal{CH}}(u) = u$ for $u \leq -1$, $V_{\mathcal{CH}}(v) = v$ for $v \leq -1$,
- $U_{\mathcal{CH}}(u) \rightarrow 0$ as $u \rightarrow \infty$, $V_{\mathcal{CH}}(v) \rightarrow 0$ as $v \rightarrow \infty$,
- there exists a $u_+ \geq -1$ such that $\frac{dU_{\mathcal{CH}}}{du} = \exp(\kappa_- u)$ for $u \geq u_+$,
- there exists a $v_+ \geq -1$ such that $\frac{dV_{\mathcal{CH}}}{dv} = \exp(\kappa_- v)$ for $v \geq v_+$.

As above, this defines null coordinates $U_{\mathcal{CH}}: \mathcal{M} \rightarrow (-\infty)$ via $U_{\mathcal{CH}}(u)$ and $V_{\mathcal{CH}}: \mathcal{M} \rightarrow (0, \infty)$ via $V_{\mathcal{CH}}(v)$, where u, v are the null coordinates defined in (2.9).

To define the event horizon, we consider the coordinate chart $(U_{\mathcal{H}}, V_{\mathcal{H}}, \theta, \phi)$. We now define the event horizon without the bifurcation sphere as the union

$$\mathcal{H}_0 := \mathcal{H}_A \cup \mathcal{H}_B, \quad (2.10)$$

where

$$\mathcal{H}_A := \{U_{\mathcal{H}} = 0\} \times (0, \infty) \times \mathbb{S}^2 \quad \text{and} \quad \mathcal{H}_B := (0, \infty) \times \{V_{\mathcal{H}} = 0\} \times \mathbb{S}^2. \quad (2.11)$$

Analogously, we also define the Cauchy horizon without the bifurcation sphere in the coordinate chart $(U_{\mathcal{CH}}, V_{\mathcal{CH}}, \theta, \phi)$ as the union

$$\mathcal{CH}_0 := \mathcal{CH}_A \cup \mathcal{CH}_B, \quad (2.12)$$

where

$$\mathcal{CH}_A := (U_{\mathcal{CH}}, \infty) \times \{V_{\mathcal{CH}} = 0\} \times \mathbb{S}^2 \quad \text{and} \quad \mathcal{CH}_B := \{U_{\mathcal{CH}} = 0\} \times (0, \infty) \times \mathbb{S}^2. \quad (2.13)$$

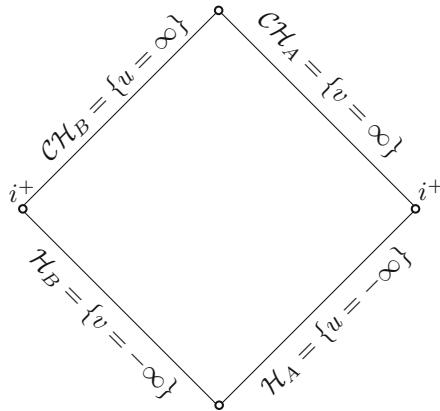
Therefore, we define the interior of the Reissner–Nordström spacetime without the bifurcation sphere as the manifold with boundary

$$\tilde{\mathcal{M}} := \mathcal{M} \cup \mathcal{H} \cup \mathcal{CH}. \quad (2.14)$$

The Lorentzian metric on \mathcal{M} extends smoothly to $\tilde{\mathcal{M}}$. In particular, the boundary of $\tilde{\mathcal{M}}$ consists of four disconnected null hypersurfaces. In Fig. 4, we have depicted its Penrose diagram. In mild abuse of notation, we shall also use the coordinate systems

$$(U_{\mathcal{H}}, v, \theta, \phi) \quad \text{on} \quad \mathcal{M} \cup \mathcal{H}_A, \quad (2.15)$$

$$(u, V_{\mathcal{H}}, \theta, \phi) \quad \text{on} \quad \mathcal{M} \cup \mathcal{H}_B, \quad (2.16)$$

FIGURE 4. Penrose diagram of $\tilde{\mathcal{M}}$

$$(u, V_{\mathcal{CH}}, \theta, \phi) \quad \text{on} \quad \mathcal{M} \cup \mathcal{CH}_A, \quad (2.17)$$

$$(U_{\mathcal{CH}}, v, \theta, \phi) \quad \text{on} \quad \mathcal{M} \cup \mathcal{CH}_B \quad (2.18)$$

In particular, we can write

$$\mathcal{H}_A = \{U_{\mathcal{H}} = 1\} \times \mathbb{S}^2 \times \mathbb{R}, \quad (2.19)$$

$$\mathcal{H}_B = \{u \in \mathbb{R}\} \times \{V_{\mathcal{H}} = 0\} \times \mathbb{S}^2, \quad (2.20)$$

$$\mathcal{CH}_A = \{v \in \mathbb{R}\} \times \{V_{\mathcal{CH}} = 0\} \times \mathbb{S}^2, \quad (2.21)$$

$$\mathcal{CH}_B = \{U_{\mathcal{CH}} = 0\} \times \{v \in \mathbb{R}\} \times \mathbb{S}^2. \quad (2.22)$$

Note also that the vector field T , initially defined on \mathcal{M} in (2.3), extends to a smooth vector field on $\tilde{\mathcal{M}}$ with

$$T|_{\mathcal{H}_A} = \frac{\partial}{\partial v}|_{\mathcal{H}_A}, \quad (2.23)$$

where $\frac{\partial}{\partial v}$ is the coordinate derivative with respect to local chart defined in (2.15). Similarly, we have

$$T|_{\mathcal{H}_B} = -\frac{\partial}{\partial u}|_{\mathcal{H}_B} \quad \text{w.r.t. to the local chart (2.16),} \quad (2.24)$$

$$T|_{\mathcal{CH}_A} = -\frac{\partial}{\partial u}|_{\mathcal{CH}_A} \quad \text{w.r.t. to the local chart (2.17),} \quad (2.25)$$

$$T|_{\mathcal{CH}_B} = \frac{\partial}{\partial v}|_{\mathcal{CH}_B} \quad \text{w.r.t. to the local chart (2.18).} \quad (2.26)$$

Note that T is a Killing null generator of the Killing horizons \mathcal{H}_A , \mathcal{H}_B , \mathcal{CH}_A , and \mathcal{CH}_B . Recall also that $\nabla_T T|_{\mathcal{CH}} = \kappa_- T|_{\mathcal{CH}}$ and $\nabla_T T|_{\mathcal{H}} = \kappa_+ T|_{\mathcal{H}}$, where κ_{\pm} is defined by (2.8).

At this point, we note that we can attach corners to \mathcal{H}_0 and \mathcal{CH}_0 to extend $\tilde{\mathcal{M}}$ smoothly to a Lorentzian manifold with corners. To be more precise, we attach the past bifurcation sphere \mathcal{B}_- to \mathcal{H}_0 as the point $(U_{\mathcal{H}}, V_{\mathcal{H}}) = (0, 0)$. Then, define $\mathcal{H} := \mathcal{H}_0 \cup \mathcal{B}_-$. Similarly, we can attach the future bifurcation

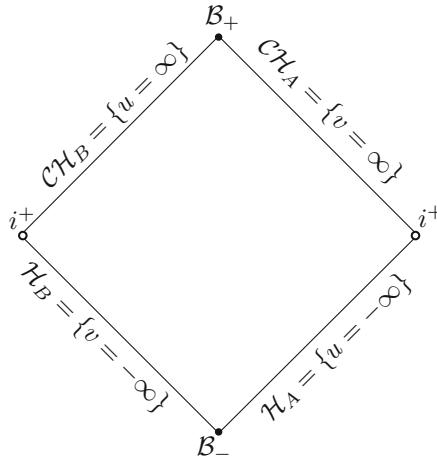


FIGURE 5. Penrose diagram of \mathcal{M}_{RN} which includes the bifurcate spheres \mathcal{B}_+ and \mathcal{B}_-

sphere \mathcal{B}_+ to the Cauchy horizon which will be denoted by \mathcal{CH} . We call the resulting manifold \mathcal{M}_{RN} . Further details are not given since the precise construction is straightforward and the metric extends smoothly. Moreover, the T vector field extends smoothly to \mathcal{B}_+ and \mathcal{B}_- and vanishes there. Its Penrose diagram is depicted in Fig. 5

Further details about the coordinate systems can be found in [42]. From a dynamical point of view, we can also consider the spacetimes $(\mathcal{M}_{\text{RN}}, g_{M,Q})$ as being contained in the Cauchy development of a spacelike hypersurface with two asymptotically flat ends, solving the Einstein–Maxwell system in spherical symmetry.

2.2. The Characteristic Initial Value Problem for the Wave Equation

In the context of scattering theory, we will be interested in solutions to the wave equation (1.1) arising from suitable characteristic initial data. Recall the following well-posedness result for (1.1) with characteristic initial data.

Proposition 2.1. *Let $\Psi \in C_c^\infty(\mathcal{H})$ be smooth compactly supported data on the event horizon \mathcal{H} . Then, there exists a unique smooth solution ψ to (1.1) on $\mathcal{M}_{\text{RN}} \setminus \mathcal{CH}$ such that $\psi|_{\mathcal{H}} = \Psi$.*

The previous proposition is well known, see [38, 46]. Analogously, we have the following for the backward evolution.

Proposition 2.2. *Let $\Psi \in C_c^\infty(\mathcal{CH})$ be smooth compactly supported data on the Cauchy horizon \mathcal{CH} . Then, there exists a unique smooth solution ψ to (1.1) on $\mathcal{M}_{\text{RN}} \setminus \mathcal{H}$ such that $\psi|_{\mathcal{CH}} = \Psi$.*

2.3. Hilbert Spaces of Finite T Energy on Both Horizon Components

Now, we are in the position to define the Hilbert spaces on the event $\mathcal{H} = \mathcal{H}_A \cup \mathcal{H}_B \cup \mathcal{B}_-$ and Cauchy horizon $\mathcal{CH} = \mathcal{CH}_A \cup \mathcal{CH}_B \cup \mathcal{B}_+$, respectively.

We will start with constructing the Hilbert space on the event horizon. Roughly speaking, it will be defined by requiring finiteness of the T energy flux on \mathcal{H}_A minus the T energy flux on \mathcal{H}_B . More precisely, let $C_c^\infty(\mathcal{H})$ be the vector space of smooth compactly supported functions on \mathcal{H} . Recall that the Killing vector field T is also a null generator of \mathcal{H} and vanishes at the past bifurcation sphere \mathcal{B}_- . This allows us to define the norm $\|\cdot\|_{\mathcal{E}_\mathcal{H}^T}^2$ on the vector space $C_c^\infty(\mathcal{H})$ as

$$\|\psi\|_{\mathcal{E}_\mathcal{H}^T}^2 := \int_{\mathcal{H}_A} J_\mu^T[\psi] n_{\mathcal{H}_A}^\mu \, d\text{vol}_{n_{\mathcal{H}_A}} - \int_{\mathcal{H}_B} J_\mu^T[\psi] n_{\mathcal{H}_B}^\mu \, d\text{vol}_{n_{\mathcal{H}_B}}, \quad (2.27)$$

where $\psi \in C_c^\infty(\mathcal{H})$, $\mathbf{T}[\psi]$ is the energy momentum tensor

$$\mathbf{T}[\psi]_{\mu\nu} := \text{Re}(\partial_\mu \psi \overline{\partial_\nu \psi}) - \frac{1}{2} g_{\mu\nu} \partial_\alpha \psi \overline{\partial^\alpha \psi}, \quad (2.28)$$

and $J^T[\psi] := \mathbf{T}[\psi](T, \cdot)$. In (2.27), the fluxes are defined with respect to future directed normal vector fields $n_{\mathcal{H}_A}$ and $n_{\mathcal{H}_B}$ on \mathcal{H}_A and \mathcal{H}_B , respectively.³ Moreover, recall from Fig. 2 that T is future (resp. past) directed on \mathcal{H}_A (resp. \mathcal{H}_B). Thus, the terms arising in (2.27) satisfy $\int_{\mathcal{H}_A} J_\mu^T[\psi] n_{\mathcal{H}_A}^\mu \, d\text{vol} \geq 0$ and $-\int_{\mathcal{H}_B} J_\mu^T[\psi] n_{\mathcal{H}_B}^\mu \, d\text{vol} \geq 0$. Again, in view of the fact that on the component \mathcal{H}_B the normal vector field T is past directed, we can choose $n_{\mathcal{H}_A} := T|_{\mathcal{H}_A}$ and $n_{\mathcal{H}_B} := -T|_{\mathcal{H}_B}$ as the future directed normal vector fields on \mathcal{H}_A and \mathcal{H}_B , respectively, such that we can rewrite the norm (2.27) as [using the coordinate charts (2.15) and (2.16)]

$$\|\psi\|_{\mathcal{E}_\mathcal{H}^T}^2 = \int_{\mathbb{R} \times \mathbb{S}^2} |\partial_v \psi|_{\mathcal{H}_A}^2 \, v \sin \theta \, d\theta \, d\varphi + \int_{\mathbb{R} \times \mathbb{S}^2} |\partial_u \psi|_{\mathcal{H}_B}^2 \, du \sin \theta \, d\theta \, d\varphi. \quad (2.29)$$

The norm (2.29) defines an inner product, hence its completion is a Hilbert space.

Definition 2.1. We define the Hilbert space of finite T energy $\mathcal{E}_\mathcal{H}^T$ on both components of the event horizon as the completion of smooth and compactly supported functions $C_c^\infty(\mathcal{H})$ on the event horizon $\mathcal{H} = \mathcal{H}_A \cup \mathcal{H}_B \cup \mathcal{B}_-$ with respect to the norm (2.27).

Analogously, we can consider the vector space $C_c^\infty(\mathcal{CH})$ and define the norm $\|\cdot\|_{\mathcal{E}_{\mathcal{CH}}^T}^2$ as the T energy flux on the component \mathcal{CH}_B minus the T energy flux on the component \mathcal{CH}_A :

$$\|\psi\|_{\mathcal{E}_{\mathcal{CH}}^T}^2 := \int_{\mathcal{CH}_B} J_\mu^T[\psi] n_{\mathcal{CH}_B}^\mu \, d\text{vol}_{n_{\mathcal{CH}_B}} - \int_{\mathcal{CH}_A} J_\mu^T[\psi] n_{\mathcal{CH}_A}^\mu \, d\text{vol}_{n_{\mathcal{CH}_A}}. \quad (2.30)$$

³A choice of such normal vectors fixes the volume form. Also note that this is the natural setup for energy estimates.

Again, in view of the orientation of the T vector field (cf. Fig. 2), this norm can be represented as [using the coordinate charts (2.17) and (2.18)]

$$\|\psi\|_{\mathcal{E}_{\mathcal{CH}}^T}^2 = \int_{\mathbb{R} \times \mathbb{S}^2} |\partial_v \psi|_{\mathcal{CH}_B}^2 dv \sin \theta d\theta d\varphi + \int_{\mathbb{R} \times \mathbb{S}^2} |\partial_u \psi|_{\mathcal{CH}_A}^2 du \sin \theta d\theta d\varphi. \quad (2.31)$$

Definition 2.2. We define the Hilbert space of finite T energy $\mathcal{E}_{\mathcal{CH}}^T$ on both components of the Cauchy horizon as the completion of smooth and compact! supported functions $C_c^\infty(\mathcal{CH})$ the Cauchy horizon $\mathcal{CH} = \mathcal{CH}_A \cup \mathcal{CH}_B \cup \mathcal{B}_+$ with respect to the norm (2.30).

2.4. Separation of Variables

In this section, we show how the radial o.d.e. (1.2) arises from decomposing a general solution in spherical harmonics and Fourier modes. For concreteness, let ψ be a smooth solution to $\square_g \psi = 0$ such that on each $\{r = const.\}$ slice, ψ is compactly supported in the t variable.⁴ Then, we can define its Fourier transform in the t variable and also decompose ψ in spherical harmonics to end up with

$$\hat{\psi}_{m\ell}(r, \omega) := \int_{\mathbb{R} \times \mathbb{S}^2} e^{-i\omega t} Y_{m\ell}(\theta, \phi) \psi(t, r, \sigma, \psi, \sin \sigma) dt \frac{d\theta d\phi}{\sqrt{2\pi}}. \quad (2.32)$$

Due to the compact support on constant r slice, the wave equation $\square_g \psi = 0$ implies that

$$\hat{\psi}_{m\ell}(r, \omega) =: R_{m\ell}^{(\omega)}(r) =: R(r) \quad (2.33)$$

satisfies the radial o.d.e.

$$\Delta \frac{d}{dr} \left(\frac{d}{dr} R \right) - \Delta \ell(\ell+1) R + r^4 \omega^2 R = 0. \quad (2.34)$$

In Sect. 4, we will analyze solutions to (2.34) and denote a solution thereof with $R(r)$. Moreover, it is useful to introduce the function u defined as

$$u(r) := r R(r) \quad (2.35)$$

and consider $u = u(r_*)$ as a function of r_* , which is defined in (2.7). Using the r_* variable, the o.d.e. (2.34) finally reduces to

$$u'' + (\omega^2 - V_\ell) u = 0 \quad (2.36)$$

on the radial line with potential

$$V = V_\ell = \Delta \left(\frac{r(r_+ + r_-) - 2r_+ r_-}{r^3} + \frac{\ell(\ell+1)}{r^4} \right). \quad (2.37)$$

In Lemma A.3 in Appendix, it is proven that, as a function of r_* , the potential V_ℓ decays exponentially as $r_* \rightarrow \pm\infty$. In particular, this indicates that we have asymptotic free waves (asymptotic states) near the event and Cauchy horizon of the form $e^{\pm i\omega r_*}$ as $|r_*| \rightarrow \infty$. In order to construct these solutions, we use the following proposition for Volterra integral equations (see Lemma 2.4 of [48]).

⁴Note that we will prove later that such solutions arise from data which are dense in $\mathcal{E}_{\mathcal{H}}^T$.

Proposition 2.3. *Let $x_0 \in \mathbb{R} \cup \{+\infty\}$ and $g \in L^\infty(-\infty, x_0)$. Suppose the integral kernel K satisfies*

$$\alpha := \int_{-\infty}^{x_0} \sup_{\{x:y < x < x_0\}} |K(x, y)| dy < \infty. \quad (2.38)$$

Then, the Volterra integral equation

$$f(x) = g(x) + \int_{-\infty}^x K(x, y) f(y) dy \quad (2.39)$$

has a unique solution f satisfying

$$\|f\|_{L^\infty(-\infty, x_0)} \leq e^\alpha \|g\|_{L^\infty(-\infty, x_0)}. \quad (2.40)$$

If in addition K is smooth in both variables and

$$\int_{-\infty}^{x_0} \sup_{\{x:y < x < x_0\}} |\partial_x^k K(x, y)| dy < \infty \quad (2.41)$$

for all $k \in \mathbb{N}$, then the solution f is smooth on $(-\infty, x_0)$ and the derivatives can be computed by formal differentiation of (2.39).

Remark 2.1. Analogous results as in Proposition 2.3 also hold true for Volterra integral equations on intervals of the form $(-\infty, x_1)$ or $(x_0, +\infty)$.

This allows us to define the following fundamental pairs of solutions to the o.d.e. (2.36). In view of the exponential decay of the potential, it is straightforward to check that the assumptions of Proposition 2.3 are satisfied.

Definition 2.3. Let $\omega \in \mathbb{R}$ and $\ell \in \mathbb{N}_0$ be fixed. Define asymptotic state solutions u_1 and u_2 of the radial o.d.e. (2.36) as the unique solutions to the Volterra integral equations

$$u_1(\omega, r_*) = e^{i\omega r_*} - \int_{-\infty}^{r_*} \frac{\sin(\omega(r_* - y))}{\omega} V(y) u_1(\omega, y) dy, \quad (2.42)$$

$$u_2(\omega, r_*) = e^{-i\omega r_*} + \int_{-\infty}^{r_*} \frac{\sin(\omega(r_* - y))}{\omega} V(y) u_2(\omega, y) dy. \quad (2.43)$$

Analogously, define v_1 and v_2 as the unique solutions to the Volterra integral equations

$$v_1(\omega, r_*) = e^{i\omega r_*} - \int_{r_*}^{\infty} \frac{\sin(\omega(r_* - y))}{\omega} V(y) v_1(\omega, y) dy, \quad (2.44)$$

$$v_2(\omega, r_*) = e^{-i\omega r_*} - \int_{r_*}^{\infty} \frac{\sin(\omega(r_* - y))}{\omega} V(y) v_2(\omega, y) dy. \quad (2.45)$$

For $\omega = 0$, we set $\frac{\sin(\omega(r_* - y))}{\omega} \big|_{\omega=0} = r_* - y$ in the integral kernel in which case u_1 and u_2 coincide. We define

$$\tilde{u}_1(r_*) := u_1(0, r_*) = u_2(0, r_*) \quad (2.46)$$

and similarly,

$$\tilde{v}_1(r_*) := v_1(0, r_*) = v_2(0, r_*) \quad (2.47)$$

Since $u_1(0, r_*) = u_2(0, r_*)$ for $\omega = 0$, there exists another linearly independent fundamental solution \tilde{u}_2 solving the Volterra integral equation

$$\tilde{u}_2(r_*) = r_* + \int_{-\infty}^{r_*} (r_* - y)V(y)\tilde{u}_2(y)dy. \quad (2.48)$$

Similarly, we also have another fundamental solution, which is linearly independent from \tilde{v}_1 , solving

$$\tilde{v}_2(r_*) = r_* - \int_{r_*}^{\infty} (r_* - y)V(y)\tilde{v}_2(y)dy. \quad (2.49)$$

Since r_* is not uniformly bounded, we cannot apply Proposition 2.3 to construct \tilde{u}_2 and \tilde{v}_2 . Nevertheless, after switching to coordinates which are regular at \mathcal{H} or \mathcal{CH} , the existence of the desired solutions follows immediately from the usual local theory of regular singularities (see [44]).

Remark 2.2. Due to the exponential decay of the potential V_ℓ (see Lemma A.3 in Appendix), it follows from standard theory that the solutions $u_1(\omega, r_*)$, $u_2(\omega, r_*)$, $v_1(\omega, r_*)$ and $v_2(\omega, r_*)$ can be continued to holomorphic functions of ω in the strip $|\text{Im}(\omega)| < \kappa_+$ for fixed $r_* \in \mathbb{R}$. Indeed, in [5] it is shown that $u_1(\omega, r_*)$ is analytic in $\mathbb{C} \setminus \{im\kappa_+ : m \in \mathbb{N}\}$ with possible poles at $\{im\kappa_+ : m \in \mathbb{N}\}$ and similarly for u_2, v_1 , and v_2 . See also the proof of Proposition A.2 in Appendix.

This allows us now to define the reflection and transmission coefficients \mathfrak{R} and \mathfrak{T} .

Definition 2.4. Let $\omega \neq 0$. Then we define the transmission coefficient $\mathfrak{T}(\omega, \ell)$ and reflection coefficient $\mathfrak{R}(\omega, \ell)$ as the unique coefficients such that

$$u_1 = \mathfrak{T}v_1 + \mathfrak{R}v_2. \quad (2.50)$$

Using the fact that the Wronskian

$$\mathfrak{W}(f, g) := fg' - f'g \quad (2.51)$$

of two solutions f and g is independent of r_* , we can equivalently define the scattering coefficients as

$$\mathfrak{T} := \frac{\mathfrak{W}(u_1, v_2)}{\mathfrak{W}(v_1, v_2)} = \frac{\mathfrak{W}(u_1, v_2)}{-2i\omega} \quad (2.52)$$

and

$$\mathfrak{R} := \frac{\mathfrak{W}(u_1, v_1)}{\mathfrak{W}(v_2, v_1)} = \frac{\mathfrak{W}(u_1, v_1)}{2i\omega}. \quad (2.53)$$

The transmission and reflection coefficients satisfy a pseudo-unitarity property proven in the following.

Proposition 2.4 (Pseudo-unitarity in the separated picture). *The transmission and reflection coefficients satisfy*

$$1 = |\mathfrak{T}|^2 - |\mathfrak{R}|^2. \quad (2.54)$$

Proof. First, note that any solution to the o.d.e. (2.36) satisfies the identity

$$\operatorname{Im}(\bar{u}u') = \operatorname{const}. \quad (2.55)$$

Applying this to the solution $u_1 = \mathfrak{T}v_1 + \mathfrak{R}v_2$ shows the claim. \square

In the following, we shall see that the reflection and transmission coefficients are regular at $\omega = 0$.

Proposition 2.5. *Let $\ell \in \mathbb{N}_0$ be fixed. Then, the scattering coefficients $\mathfrak{R}(\omega, \ell)$ and $\mathfrak{T}(\omega, \ell)$ are analytic functions of ω in the strip $\{\omega \in \mathbb{C} : |\operatorname{Im}(\omega)| < \kappa_+\}$ with values for $\omega = 0$ given by*

$$\mathfrak{R}(0, \ell) = \frac{(-1)^\ell}{2} \left(\frac{r_-}{r_+} - \frac{r_+}{r_-} \right), \quad (2.56)$$

$$\mathfrak{T}(0, \ell) = \frac{(-1)^\ell}{2} \left(\frac{r_-}{r_+} + \frac{r_+}{r_-} \right) \quad (2.57)$$

In particular, the reflection coefficient $\mathfrak{R}(\omega, \ell)$ only vanishes on a discrete set of points ω .

Moreover, the reflection and transmission coefficients $\mathfrak{R}(\omega, \ell)$ and $\mathfrak{T}(\omega, \ell)$ are analytic functions on $\mathbb{C} \setminus \mathbb{P}$ with possible poles at $\mathbb{P} = \{im\kappa_+ : m \in \mathbb{N}\} \cup \{ik\kappa_- : k \in \mathbb{Z} \setminus \{0\}\}$.

Proof. From the analyticity of u_1, u_2, v_1 , and v_2 in the strip $|\operatorname{Im}(\omega)| < \kappa_+$ (cf. Remark 2.2), we conclude that \mathfrak{T} and \mathfrak{R} are holomorphic in $\{\omega \neq 0 \in \mathbb{C} : |\operatorname{Im}(\omega)| < \kappa_+\}$ with a possible pole at $\omega = 0$. In the following, we shall show that $\{\omega = 0\}$ is removable singularity. Indeed, we will compute the explicit value of the reflection and transmission coefficient at $\omega = 0$ and deduce that for fixed $\ell \in \mathbb{N}_0$, the transmission coefficient $\mathfrak{T}(\omega, \ell)$ and the reflection coefficient $\mathfrak{R}(\omega, \ell)$ are analytic functions on the strip $\{\omega \in \mathbb{C} : \operatorname{Im}(\omega) | < \kappa_+\}$ (cf. unpublished work of McNamara cited in [21]). To do so, note that from Proposition 4.2 in Sect. 4.1.3 we conclude the pointwise limits

$$u_1(\omega, r_*) \rightarrow \tilde{u}_1(r_*), \quad (2.58)$$

$$v_1(\omega, r_*) \rightarrow \tilde{v}_1(r_*) = (-1)^\ell \frac{r_+}{r_-} \tilde{u}_1(r_*), \quad (2.59)$$

$$v_2(\omega, r_*) \rightarrow \tilde{v}_2(r_*) = (-1)^\ell \frac{r_+}{r_-} \tilde{u}_1(r_*) \quad (2.60)$$

as $|\omega| \rightarrow 0$. Using the definition in (2.50) of $\mathfrak{T}(\omega, \ell)$, $\mathfrak{R}(\omega, \ell)$, and the condition $1 + |\mathfrak{R}|^2 = |\mathfrak{T}|^2$ (cf. Proposition 2.4), we deduce that the limits $\lim_{\omega \rightarrow 0} \mathfrak{R}(\omega, \ell)$ and $\lim_{\omega \rightarrow 0} \mathfrak{T}(\omega, \ell)$ exist and moreover can be computed to be (2.56) and (2.57). Note that (2.56) and (2.57) have been established in [22]. Also note that in view of the analyticity properties of u_1, v_1 , and v_2 , the $\mathfrak{R}(\omega, \ell)$ and $\mathfrak{T}(\omega, \ell)$ are analytic functions on $\mathbb{C} \setminus \mathbb{P}$ with possible poles at $\mathbb{P} = \{im\kappa_+ : m \in \mathbb{N}\} \cup \{ik\kappa_- : k \in \mathbb{Z} \setminus \{0\}\}$. \square

2.5. Conventions

Let X be a point set with a limit point c (e.g., $X = \mathbb{R}, [a, b], \mathbb{C}$). Throughout this paper, we will use the symbols \lesssim and \gtrsim , where the implicit constants might depend on the black hole parameters M and Q . In particular, for functions (or constants) $a(x), b(x) > 0$ the notation $a \lesssim b$ means that there is a constant $C = C(M, Q) > 0$ such that $a(x) \leq Cb(x)$ for all $x \in X$. We will also make use of the notation \lesssim_ℓ or \gtrsim_ℓ which means that the constant may additionally also depend on ℓ . We also write $a \sim b$ if there are constants $C(M, Q), \tilde{C}(M, Q) > 0$ such that $Ca(x) \leq b(x) \leq \tilde{C}a(x)$ for all $x \in X$.

We shall also make use of the standard Landau notation O and o [3, 44]. To be more precise, as $x \rightarrow c$ in X

$$f(x) = O(g(x)) \quad \text{means} \quad \left| \frac{f(x)}{g(x)} \right| \leq C(M, Q) \quad (2.61)$$

and

$$f(x) = o(g(x)) \quad \text{means} \quad \frac{f(x)}{g(x)} \rightarrow 0 \quad (2.62)$$

We will also use the notation O_ℓ if the constant C in (2.61) may additionally depend on ℓ .

3. Main Theorems

In this section, we will formulate our main theorems.

Theorem 1, which we state in Sect. 3.1, establishes the existence of a scattering map S^T of the form

$$S^T : \mathcal{E}_{\mathcal{H}}^T \rightarrow \mathcal{E}_{\mathcal{CH}}^T, \quad (3.1)$$

which is a Hilbert space isomorphism, i.e., a bounded and invertible map with bounded inverse. Theorem 1 will be proven in Sect. 5. In the separated picture, the boundedness of S^T corresponds to the uniform boundedness of the transmission and reflection coefficients which is stated as Theorem 2 in Sect. 3.2. Theorem 2 will be proven in Sect. 4 (and later used in the proof of Theorem 1).

Sect. 3.3 is then devoted to Theorem 3, which connects our physical space scattering theory to the fixed frequency scattering theory. (We will infer Theorem 3 as a consequence of Theorem 1). In Theorem 3.4, this connection allows us to prove that the reflection map is injective, which is the content of Theorem 4. In Theorem 5, which is stated and proven in Sect. 3.5, we construct data which are incoming and compactly supported but, nevertheless, lead to a solution which fails to be in C^1 on the Cauchy horizon.

We end this section with the statement of our two non-existence results. In Sect. 3.6, we formulate Theorem 6, the non-existence of the T energy scattering theory for the Klein–Gordon equation with conformal mass on the interior of (anti-) de Sitter–Reissner–Nordström black holes. The proof of Theorem 6 is given in Sect. 6. Finally, in Theorem 7, stated in Sect. 3.7, we show the non-existence of the T energy scattering map for the Klein–Gordon equation

on the interior of Reissner–Nordström. The proof of Theorem 7 is given in Sect. 7.

3.1. Existence and Boundedness of the T Energy Scattering Map

First, we define the forward (resp. backward) evolution on a dense domain.

Definition 3.1. The domains of the forward and backward evolution are defined as

$$\begin{aligned} \mathcal{D}_{\mathcal{H}}^T := \{ \psi \in C_c^\infty(\mathcal{H}) \subset \mathcal{E}_{\mathcal{H}}^T \text{ s.t. the Cauchy evolution of } \psi \text{ has} \\ \times \text{ compact support on constant } r = \text{const. hypersurfaces} \} \end{aligned} \quad (3.2)$$

and

$$\begin{aligned} \mathcal{D}_{\mathcal{CH}}^T := \{ \psi \in C_c^\infty(\mathcal{CH}) \subset \mathcal{E}_{\mathcal{CH}}^T \text{ s.t. the backward evolution of } \psi \text{ has} \\ \times \text{ compact support on constant } r = \text{const. hypersurfaces} \}, \end{aligned} \quad (3.3)$$

respectively. Here, we consider $r_- < r < r_+$ and note that if ψ is compactly supported on one $\{r = \text{const.}\}$ slice, then, as a direct consequence of the domain of dependence, its evolution will be compactly supported on all other $\{r = \text{const.}\}$ hypersurfaces for $r_- < r < r_+$.

We will prove in Lemma 5.1 in Sect. 5 that $\mathcal{D}_{\mathcal{H}}^T \subset \mathcal{E}_{\mathcal{H}}^T$ and $\mathcal{D}_{\mathcal{CH}}^T \subset \mathcal{E}_{\mathcal{CH}}^T$ are dense domains.

These definitions of the domains are motivated by the following observation.

Remark 3.1. Suppose we are given data in $\mathcal{D}_{\mathcal{H}}^T$ on the event horizon \mathcal{H} . Consider now the unique Cauchy evolution (cf. Proposition 2.1) and observe that its restriction to the Cauchy horizon \mathcal{CH} will lie in $\mathcal{D}_{\mathcal{CH}}^T$. This holds true since we can first smoothly extend the metric beyond the Cauchy horizon \mathcal{CH} and then use the compact support on a constant r_* hypersurface to solve an equivalent Cauchy problem in an appropriate region which extends the Cauchy horizon \mathcal{CH} and includes the support of the solution, but does not include i^+ . The smoothness of the solution up to and including the Cauchy horizon \mathcal{CH} follows now from Cauchy stability.

In view of Remark 3.1, we can define the forward and backward map on the domains $\mathcal{D}_{\mathcal{H}}^T$ and $\mathcal{D}_{\mathcal{CH}}^T$, respectively.

Definition 3.2. Define the forward map $S_0^T : \mathcal{D}_{\mathcal{H}}^T \subset \mathcal{E}_{\mathcal{H}}^T \rightarrow \mathcal{D}_{\mathcal{CH}}^T \subset \mathcal{E}_{\mathcal{CH}}^T$ as the unique forward evolution from data on the event horizon to data on the Cauchy horizon. More precisely, let ψ be the solution to (1.1) arising from initial data $\Psi \in \mathcal{D}_{\mathcal{H}}^T \subset \mathcal{E}_{\mathcal{H}}^T$. Then, define $S_0^T(\Psi)$ as the restriction of ψ to the Cauchy horizon, i.e., $S_0^T(\Psi) := \psi|_{\mathcal{CH}} \in \mathcal{D}_{\mathcal{CH}}^T$.

Similarly, let ϕ be the unique backward evolution of (1.1) arising from $\Phi \in \mathcal{D}_{\mathcal{CH}}^T$. Then, define the backward map by $B_0^T(\Phi) := \phi|_{\mathcal{H}} \in \mathcal{D}_{\mathcal{H}}^T$.

Remark 3.2. Note that by the uniqueness of the Cauchy evolution we have that S_0^T and B_0^T are inverses of each other, i.e., $B_0^T \circ S_0^T = \text{Id}_{\mathcal{D}_{\mathcal{H}}^T}$, $S_0^T \circ B_0^T = \text{Id}_{\mathcal{D}_{\mathcal{CH}}^T}$.

Now, we are in the position to state our main theorem.

Theorem 1. *The map $S_0^T: \mathcal{D}_{\mathcal{H}}^T \subset \mathcal{E}_{\mathcal{H}}^T \rightarrow \mathcal{D}_{\mathcal{CH}}^T \subset \mathcal{E}_{\mathcal{CH}}^T$ is bounded and uniquely extends to*

$$S^T: \mathcal{E}_{\mathcal{H}}^T \rightarrow \mathcal{E}_{\mathcal{CH}}^T, \quad (3.4)$$

called the “scattering map”. The scattering map S^T is a Hilbert space isomorphism, i.e., a bounded and invertible linear map with bounded inverse $B^T: \mathcal{E}_{\mathcal{CH}}^T \rightarrow \mathcal{E}_{\mathcal{H}}^T$ satisfying

$$B^T \circ S^T = \text{Id}_{\mathcal{E}_{\mathcal{H}}^T}, \quad S^T \circ B^T = \text{Id}_{\mathcal{E}_{\mathcal{CH}}^T}. \quad (3.5)$$

Here, $B^T: \mathcal{E}_{\mathcal{CH}}^T \rightarrow \mathcal{E}_{\mathcal{H}}^T$ is the “backward map,” which is the unique bounded extension of B_0^T .

In addition, the scattering map S^T is pseudo-unitary, meaning that for $\psi \in \mathcal{E}_{\mathcal{H}}^T$, we have

$$\int_{\mathcal{H}_A} |T\psi|^2 - \int_{\mathcal{H}_B} |T\psi|^2 = \int_{\mathcal{CH}_B} |TS^T\psi|^2 - \int_{\mathcal{CH}_A} |S^T\psi|^2. \quad (3.6)$$

In more traditional language, Theorem 1 yields existence, uniqueness, and asymptotic completeness of scattering states.

The proof of Theorem 1 is given in Sect. 5. Let us note already that Theorem 1 is a posteriori the physical space equivalent of the uniform boundedness of the scattering coefficients proved in Theorem 2 (see Sect. 3.2). This equivalence is made precise in Theorem 3 (see Sect. 3.3).

Remark 3.3. Note that in general, neither initial data nor scattered data have to be bounded in L^∞ or continuous. Indeed, we have that $\Phi_A(u, \theta, \varphi) = \log(u)\chi_{u \geq 1} \in \mathcal{E}_{\mathcal{CH}_A}^T$, where $\chi_{u \geq 1}$ is a smooth cutoff. Thus, there exist initial data $B^T(\Phi_A) \in \mathcal{E}_{\mathcal{H}}^T$ such that its image under the scattering map is not in L^∞ and not continuous. We emphasize the contrast with the estimates from [17] for which more regularity and decay along the event horizon \mathcal{H} are necessary.

3.2. Uniform Boundedness of the Transmission and Reflection Coefficients

On the level of the o.d.e. (2.36) in the separated picture, the problem of boundedness of the scattering map reduces to proving that the transmission coefficient \mathfrak{T} and the reflection coefficient \mathfrak{R} are uniformly bounded over all parameter ranges of $\omega \in \mathbb{R}$ and $\ell \in \mathbb{N}_0$. This is stated as Theorem 2.

Theorem 2. *The reflection and transmission coefficients $\mathfrak{R}(\omega, \ell)$ and $\mathfrak{T}(\omega, \ell)$ are uniformly bounded, i.e., they satisfy*

$$\sup_{\omega \in \mathbb{R}, \ell \in \mathbb{N}_0} (|\mathfrak{R}(\omega, \ell)| + |\mathfrak{T}(\omega, \ell)|) \lesssim 1. \quad (3.7)$$

Theorem 2 is proved in Sect. 4. As discussed in the introduction, the proof relies on an explicit calculation for $\omega = 0$ together with a careful analysis of the radial o.d.e. (2.36), involving properties of special functions and perturbations thereof.

Let us note that, given Theorem 1, we could infer Theorem 2 as a corollary (using the theory to be described in Sect. 3.3). We emphasize, however, that in the present paper we use Theorem 2 to prove Theorem 1 in Sect. 5.

3.3. Connection Between the Separated and the Physical Space Picture

In this section, we will make the connection of the separated and physical space picture precise.

First, let us note that we have natural Hilbert space decompositions $\mathcal{E}_{\mathcal{H}}^T \cong \mathcal{E}_{\mathcal{H}_A}^T \oplus \mathcal{E}_{\mathcal{H}_B}^T$ and $\mathcal{E}_{\mathcal{CH}}^T \cong \mathcal{E}_{\mathcal{CH}_A}^T \oplus \mathcal{E}_{\mathcal{CH}_B}^T$.

Proposition 3.1. *The Hilbert spaces $\mathcal{E}_{\mathcal{H}}^T$ and $\mathcal{E}_{\mathcal{CH}}^T$ of finite T energy on the event horizon \mathcal{H} and on the Cauchy horizon \mathcal{CH} admit the orthogonal decompositio*

$$\mathcal{E}_{\mathcal{H}}^T \cong \mathcal{E}_{\mathcal{H}_A}^T \oplus \mathcal{E}_{\mathcal{H}_B}^T \text{ and } \mathcal{E}_{\mathcal{CH}}^T \cong \mathcal{E}_{\mathcal{CH}_A}^T \oplus \mathcal{E}_{\mathcal{CH}_B}^T. \quad (3.8)$$

Proof. Clearly, the embedding $i: \mathcal{E}_{\mathcal{H}_A}^T \oplus \mathcal{E}_{\mathcal{H}_B}^T \hookrightarrow \mathcal{E}_{\mathcal{H}}^T$ is well defined and isometric. It remains to show that i is surjective. Let $\psi \in C_c^\infty(\mathcal{H})$. First, we know that we can approximate (in T -energy) $\psi|_{\mathcal{H}_A}$ on \mathcal{H}_A with functions $\psi_\epsilon \in C_c^\infty(\mathcal{H}_A)$ which are supported away from the past bifurcation sphere. On \mathcal{H}_A , choose non-degenerate coordinates $(V, \theta, \varphi) := (V_{\mathcal{H}}, \theta, \varphi)$ as in Sect. 2.1.2 and recall that the past bifurcation sphere is $\{V = 0\}$. Then, for small $\epsilon > 0$, set

$$\psi_\epsilon(V, \theta, \varphi) := \psi(U = 0, V, \theta, \varphi) \chi(-\log(V)), \quad (3.9)$$

where $\chi: \mathbb{R} \rightarrow [0, 1]$ is smooth and such that $\text{supp}(\chi) \subseteq (-\infty, 2]$ and $\chi|_{(-\infty, 1]} = 1$. Then, it is straightforward to check that $\psi_\epsilon \in C_c^\infty(\mathcal{H}_A)$ and

$$\int_{\mathcal{H}_A} J^T[\psi - \psi_\epsilon]_\mu n^\mu d\text{vol} \lesssim \int_{\mathbb{S}^2} \int_0^\infty V \partial_V(\psi - \psi_\epsilon)^2 dV \sin \theta d\theta d\varphi \rightarrow 0 \quad (3.10)$$

as $\epsilon \rightarrow 0$. Analogously, we can do this for \mathcal{H}_B from which the claim follows. \square

We will use this identification to represent the scattering map also in the Fourier picture and show how these pictures connect. To do so, we define the following.

Definition 3.3. For $(\Psi_A, \Psi_B) \in \mathcal{E}_{\mathcal{H}_A}^T \oplus \mathcal{E}_{\mathcal{H}_B}^T$, note that $\partial_v \Psi_A(v, \theta, \phi) \in L^2(\mathbb{R} \times \mathbb{S}^2; \mathbb{C})$ and analogously for Ψ_B . Hence, in mild abuse of notation, we can define the Fourier and spherical harmonics coefficients $\mathcal{F}_{\mathcal{H}_A}(\Psi_A)$ and $\mathcal{F}_{\mathcal{H}_B}(\Psi_B)$ as

$$i \mathcal{F}_{\mathcal{H}_A}(\Psi_A)(\omega, m, \ell) := r_+ \int_{\mathbb{R}} \int_{\mathbb{S}^2} \partial_v \Psi_A(v, \theta, \varphi) e^{-i\omega v} Y_{\ell m}(\theta, \varphi) \sin \theta d\theta d\varphi \frac{dv}{\sqrt{2\pi}} \quad (3.11)$$

and

$$-i\omega \mathcal{F}_{\mathcal{H}_B}(\Psi_B)(\omega, m, \ell) := r_+ \int_{\mathbb{R}} \int_{\mathbb{S}^2} \partial_u \Psi_B(u, \theta, \varphi) e^{i\omega u} Y_{\ell m}(\theta, \varphi) \sin \theta d\theta d\varphi \frac{du}{\sqrt{2\pi}}. \quad (3.12)$$

Similarly, for $(\Phi_A, \Phi_B) \in \mathcal{E}_{\mathcal{CH}_A}^T \oplus \mathcal{E}_{\mathcal{CH}_B}^T$ set

$$-i\omega \mathcal{F}_{\mathcal{CH}_A}(\Phi_A)(\omega, m, \ell) := r_- \int_{\mathbb{R}} \int_{\mathbb{S}^2} \partial_u \Phi_A(u, \theta, \varphi) e^{i\omega u} Y_{\ell m}(\theta, \varphi) \sin \theta d\theta d\varphi \frac{du}{\sqrt{2\pi}} \quad (3.13)$$

and

$$i\omega \mathcal{F}_{\mathcal{CH}_B}(\Phi_B)(\omega, m, \ell) := r_- \int_{\mathbb{R}} \int_{\mathbb{S}^2} \partial_v \Phi_B(v, \theta, \varphi) e^{-i\omega v} Y_{\ell m}(\theta, \varphi) \sin \theta d\theta d\varphi \frac{dv}{\sqrt{2\pi}}. \quad (3.14)$$

Also, recall the Hilbert space decomposition $\mathcal{E}_{\mathcal{H}}^T \cong \mathcal{E}_{\mathcal{H}_A}^T \oplus \mathcal{E}_{\mathcal{H}_B}^T$ and $\mathcal{E}_{\mathcal{CH}}^T \cong \mathcal{E}_{\mathcal{CH}_B}^T \oplus \mathcal{E}_{\mathcal{CH}_A}^T$. Thus, the scattering matrix can be also decomposed as

$$S^T = \begin{pmatrix} S_{BA}^T & S_{BB}^T \\ S_{AA}^T & S_{AB}^T \end{pmatrix}, \quad (3.15)$$

where

$$S_{ij}^T: \mathcal{E}_{\mathcal{H}_j}^T \rightarrow \mathcal{E}_{\mathcal{CH}_i}^T \quad (3.16)$$

is a bounded linear map for $i, j \in \{A, B\}$.⁵

Definition 3.4. Define the Hilbert spaces

$$\hat{\mathcal{E}}_{\mathcal{H}_A}^T := \ell^2(Z; L^2(r_+^{-2} \omega^2 d\omega)), \quad \hat{\mathcal{E}}_{\mathcal{H}_B}^T := \ell^2(Z; L^2(r_-^{-2} \omega^2 d\omega)),$$

$$\hat{\mathcal{E}}_{\mathcal{CH}_A}^T := \ell^2(Z; L^2(r_-^{-2} \omega^2 d\omega)), \quad \hat{\mathcal{E}}_{\mathcal{CH}_B}^T := \ell^2(Z; L^2(r_+^{-2} \omega^2 d\omega)),$$

where $Z = \{(m, \ell) \in \mathbb{Z} \times \mathbb{N}_0 : |m| \leq \ell\}$.

The Hilbert spaces defined in Definition 3.4 are unitary isomorphic to their corresponding physical energy spaces. This is captured in

Proposition 3.2. *The linear maps defined in (3.11)–(3.14)*

$$\mathcal{F}_{\mathcal{H}_A} \oplus \mathcal{F}_{\mathcal{H}_B}: \mathcal{E}_{\mathcal{H}_A}^T \oplus \mathcal{E}_{\mathcal{H}_B}^T \rightarrow \hat{\mathcal{E}}_{\mathcal{H}_A}^T \oplus \hat{\mathcal{E}}_{\mathcal{H}_B}^T \quad (3.17)$$

$$\mathcal{F}_{\mathcal{CH}_B} \oplus \mathcal{F}_{\mathcal{CH}_A}: \mathcal{E}_{\mathcal{CH}_B}^T \oplus \mathcal{E}_{\mathcal{CH}_A}^T \rightarrow \hat{\mathcal{E}}_{\mathcal{CH}_B}^T \oplus \hat{\mathcal{E}}_{\mathcal{CH}_A}^T \quad (3.18)$$

are unitary.

Proof. This follows from the fact that the Fourier transform and the decomposition into spherical harmonics are unitary maps. \square

Now, we will define the scattering map in the separated picture and show that it is bounded.

Proposition 3.3. *The scattering map in the separated picture*

$$\hat{S}^T: \hat{\mathcal{E}}_{\mathcal{H}_A}^T \oplus \hat{\mathcal{E}}_{\mathcal{H}_B}^T \rightarrow \hat{\mathcal{E}}_{\mathcal{CH}_B}^T \oplus \hat{\mathcal{E}}_{\mathcal{CH}_A}^T, \quad (3.19)$$

defined as the multiplication operator

$$\hat{S}^T = \begin{pmatrix} \hat{S}_{BA}^T & \hat{S}_{BB}^T \\ \hat{S}_{AA}^T & \hat{S}_{AB}^T \end{pmatrix} := \begin{pmatrix} \mathfrak{T}(\omega, \ell) & \bar{\mathfrak{R}}(\omega, \ell) \\ \mathfrak{R}(\omega, \ell) & \bar{\mathfrak{T}}(\omega, \ell) \end{pmatrix}, \quad (3.20)$$

is bounded. Moreover, the map \hat{S}^T is invertible with bounded inverse given by

$$\hat{S}^{T^{-1}} = \begin{pmatrix} \bar{\mathfrak{T}}(\omega, \ell) & -\bar{\mathfrak{R}}(\omega, \ell) \\ -\mathfrak{R}(\omega, \ell) & \mathfrak{T}(\omega, \ell) \end{pmatrix}. \quad (3.21)$$

⁵Note that T does not denote the transpose but the fact that it is the scattering map associated with the T vector field.

Proof. Indeed, \hat{S}^T is bounded in view of the uniform boundedness of the transmission and reflection coefficients \mathfrak{T} and \mathfrak{R} (cf. Theorem 2). Also note that $|\mathfrak{T}|^2 = 1 + |\mathfrak{R}|^2$ implies that

$$\det(\hat{S}^T) = 1 \quad (3.22)$$

which shows (3.21). The boundedness of \hat{S}^T is again immediate since the scattering coefficients are uniformly bounded. \square

Using the previous definitions, we obtain the following connection for the scattering map between the physical space and the separated picture.

Theorem 3. *The following diagram commutes and each arrow is a Hilbert space isomorphism:*

$$\begin{array}{ccc} \mathcal{E}_{\mathcal{H}_A}^T \oplus \mathcal{E}_{\mathcal{H}_B}^T & \xrightarrow{S^T} & \mathcal{E}_{\mathcal{CH}_B}^T \oplus \mathcal{E}_{\mathcal{CH}_A}^T \\ \mathcal{F}_{\mathcal{H}_A} \oplus \mathcal{F}_{\mathcal{H}_B} \downarrow & & \downarrow \mathcal{F}_{\mathcal{CH}_B} \cap \mathcal{F}_{\mathcal{CH}_A} \\ \hat{\mathcal{E}}_{\mathcal{H}_A}^T \oplus \hat{\mathcal{E}}_{\mathcal{H}_B}^T & \xrightarrow{\hat{S}^T} & \hat{\mathcal{E}}_{\mathcal{CH}_B}^T \oplus \hat{\mathcal{E}}_{\mathcal{CH}_A}^T. \end{array}$$

Moreover, the maps S^T and \hat{S}^T are pseudo-unitary satisfying (3.6) and (2.54), respectively. More concretely, for $(\Psi_A, \bar{\Psi}_B) \in \mathcal{E}_{\mathcal{H}_A}^T \oplus \mathcal{E}_{\mathcal{H}_B}^T$, we can write

$$\begin{pmatrix} \Phi_B \\ \Phi_A \end{pmatrix} = \begin{pmatrix} \mathfrak{T} & \mathfrak{R} \\ \bar{\mathfrak{R}} & \bar{\mathfrak{T}} \end{pmatrix} \begin{pmatrix} \Psi_B \\ \Psi_A \end{pmatrix}, \quad (3.23)$$

where $\partial_u \Phi_A \in L^2(\mathcal{CH}_A)$ and $\partial_v \Phi_B \in L^2(\mathcal{CH}_B)$ can be represented by

$$\begin{aligned} \partial_u \Phi_A(u, \theta, \varphi) &= \frac{1}{\sqrt{2\pi}r_-} \int_{\mathbb{R}} \sum_{|m| \leq \ell} -i\omega \mathfrak{T}(\omega, \ell) \mathcal{F}_{\mathcal{H}_A}(\Psi_A)(\omega, m, \ell) Y_{m\ell}(\theta, \varphi) e^{-i\omega u} d\omega \\ &\quad + \frac{1}{\sqrt{2\pi}r_-} \int_{\mathbb{R}} \sum_{|m| \leq \ell} -i\omega \bar{\mathfrak{T}}(\omega, \ell) \mathcal{F}_{\mathcal{H}_B}(\Psi_B)(\omega, m, \ell) Y_{m\ell}(\theta, \varphi) e^{-i\omega u} d\omega \end{aligned} \quad (3.24)$$

and

$$\begin{aligned} \partial_v \Phi_B(v, \theta, \varphi) &= \frac{1}{\sqrt{2\pi}r_-} \int_{\mathbb{R}} \sum_{|m| \leq \ell} i\omega \mathfrak{T}(\omega, \ell) \mathcal{F}_{\mathcal{H}_A}(\Psi_A)(\omega, m, \ell) Y_{m\ell}(\theta, \varphi) e^{i\omega v} d\omega \\ &\quad + \frac{1}{\sqrt{2\pi}r_-} \int_{\mathbb{R}} \sum_{|m| \leq \ell} i\omega \bar{\mathfrak{R}}(\omega, \ell) \mathcal{F}_{\mathcal{H}_B}(\Psi_B)(\omega, m, \ell) Y_{m\ell}(\theta, \varphi) e^{i\omega v} d\omega \end{aligned} \quad (3.25)$$

as well as $\Phi_A \in \mathcal{E}_{\mathcal{CH}_A}^T \cong \dot{H}^1(\mathbb{R}; L^2(\mathbb{S}^2))$, $\Phi_B \in \mathcal{E}_{\mathcal{CH}_B}^T \cong \dot{H}^1(\mathbb{R}; L^2(\mathbb{S}^2))$ can be represented by regular distributions as

$$\begin{aligned} \Phi_A(u, \theta, \varphi) &= \frac{1}{\sqrt{2\pi}r_-} \text{p. v.} \int_{\mathbb{R}} \sum_{|m| \leq \ell} \mathfrak{R}(\omega, \ell) \mathcal{F}_{\mathcal{H}_A}(\Psi_A)(\omega, m, \ell) Y_{m\ell}(\theta, \varphi) e^{-i\omega u} d\omega \\ &\quad + \frac{1}{\sqrt{2\pi}r_-} \text{p. v.} \int_{\mathbb{R}} \sum_{|m| \leq \ell} \bar{\mathfrak{T}}(\omega, \ell) \mathcal{F}_{\mathcal{H}_B}(\Psi_B)(\omega, m, \ell) Y_{m\ell}(\theta, \varphi) e^{-i\omega u} d\omega \end{aligned} \quad (3.26)$$

and

$$\begin{aligned}\Phi_B(v, \theta, \varphi) = & \frac{1}{\sqrt{2\pi r_-}} \text{p.v.} \int_{\mathbb{R}} \sum_{|m| \leq \ell} \mathfrak{T}(\omega, \ell) \mathcal{F}_{\mathcal{H}_A}(\Psi_A)(\omega, m, \ell) Y_{m\ell}(\theta, \varphi) e^{i\omega v} d\omega \\ & + \frac{1}{\sqrt{2\pi r_-}} \text{p.v.} \int_{\mathbb{R}} \sum_{|m| \leq \ell} \mathfrak{R}(\omega, \ell) \mathcal{F}_{\mathcal{H}_B}(\Psi_B)(\omega, m, \ell) Y_{m\ell}(\theta, \varphi) e^{i\omega v} d\omega.\end{aligned}\quad (3.27)$$

Proof. This is a direct consequence of Theorems 1, 2 and (5.30), (5.31) in the proof of Proposition 5.1. \square

From the previous representation of the scattered solution, we can draw a link between the boundedness of the scattering map and the fact that compactly supported incoming data will lead to solutions which vanish in the future bifurcation sphere \mathcal{B}_+ . This is the content of the following

Corollary 3.1. *Let $\Psi = (\Psi_A, 0) \in \mathcal{E}_{\mathcal{H}_A}^T \oplus \mathcal{E}_{\mathcal{H}_B}^T$ be purely incoming smooth data. Assume further that Ψ_A is supported away from the past bifurcation sphere \mathcal{B}_- and future timelike infinity i^+ .*

Then, the Cauchy evolution ψ arising from Ψ_A vanishes at the future bifurcation sphere \mathcal{B}_+ .

On the other hand, if Ψ , as above, led to a solution which does not vanish at the future bifurcation sphere \mathcal{B}_+ , then the scattering map $S^T : \mathcal{E}_{\mathcal{H}}^T \rightarrow \mathcal{E}_{\mathcal{CH}}^T$ could not be bounded.

Proof. The first claim is a direct consequence of (3.27) in Theorem 3.

For the second statement, let Ψ_A be compactly supported data on the event horizon and assume that its Cauchy evolution ψ does not vanish at the future bifurcation sphere \mathcal{B}_+ . Now take data $\tilde{\Psi}_A$ which is supported away from the past bifurcation sphere \mathcal{B}_- and satisfies $T\tilde{\Psi}_A = \Psi_A$. Then, $\tilde{\Psi}_A \in \mathcal{E}^T$ but its Cauchy evolution $\tilde{\psi}$ satisfies $\tilde{\psi}|_{\mathcal{CH}} \notin \mathcal{E}_{\mathcal{CH}}^T$ since

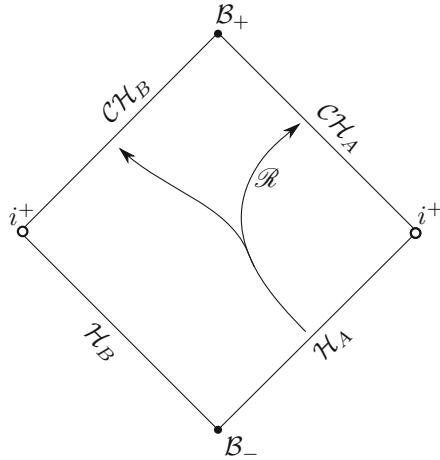
$$\|\tilde{\psi}|_{\mathcal{CH}_B}\|_{\mathcal{E}_{\mathcal{CH}_B}^T}^2 = \int_{\mathbb{R} \times \mathbb{S}^2} |\tilde{\psi}|_{\mathcal{CH}_B}(v, \theta, \varphi)|^2 dv \sin \theta d\theta d\varphi = \infty, \quad (3.28)$$

as $\psi|_{\mathcal{CH}} = \tilde{\psi}|_{\mathcal{CH}_B}$ does not vanish at the future bifurcation sphere \mathcal{B}_+ . By cutting off smoothly, one can construct normalized (in $\mathcal{E}_{\mathcal{H}}^T$ -norm) smooth compactly supported initial data on $\mathcal{E}_{\mathcal{H}}^T$ such that its Cauchy evolution has arbitrary large norm $\mathcal{E}_{\mathcal{CH}}^T$ -norm at the Cauchy horizon. \square

Remark 3.4. For convenience, we have stated the second statement of Corollary 3.1 only for the interior of Reissner–Nordström. However, note that it holds true for more general black hole interiors (e.g., subextremal (anti-) de Sitter–Reissner–Nordström) with Penrose diagram as depicted in Fig. 5.

3.4. Injectivity of the Reflection Map

In this section, we define the reflection operator of purely incoming radiation (cf. Fig. 6) and prove that it has trivial kernel as an operator from $\mathcal{E}_{\mathcal{H}_A}^T \rightarrow \mathcal{E}_{\mathcal{CH}_A}^T$.

FIGURE 6. Reflection \mathcal{R} of purely incoming radiation

Definition 3.5 (*Reflection operator*). For purely incoming radiation $(\Psi_A, 0) \in \mathcal{E}_{\mathcal{H}_A}^T \oplus \mathcal{E}_{\mathcal{H}_B}^T$, define the reflection operator

$$\mathcal{R} : \mathcal{E}_{\mathcal{H}_A}^T \rightarrow \mathcal{E}_{\mathcal{H}_A}^T \quad (3.29)$$

as

$$\mathcal{R}(\Psi_A) = \Phi_A = \text{pr}_A \left(S^T \begin{pmatrix} \Psi_A \\ 0 \end{pmatrix} \right), \quad (3.30)$$

where $\text{pr}_A : \mathcal{E}_{\mathcal{H}_B}^T \oplus \mathcal{E}_{\mathcal{H}_A}^T \rightarrow \mathcal{E}_{\mathcal{H}_A}^T$ is the orthogonal projection.

Theorem 4 *The reflection operator \mathcal{R} defined in Definition 3.5 has trivial kernel.*

Proof. Assume $\mathcal{R}(\Psi_A) = 0$ for some $\Psi_A \in \mathcal{E}_{\mathcal{H}_A}^T$. Then, in view of Theorem 3,

$$\mathcal{R}(\omega, \ell) \mathcal{F}_{\mathcal{H}_A}(\Psi_A)(\omega, m, \ell) = 0 \quad (3.31)$$

for all m, ℓ and a.e. $\omega \in \mathbb{R}$. Moreover, since $\mathcal{R}(\omega, \ell)$ only vanishes on a discrete set (cf. Proposition 2.5), we obtain that $\mathcal{F}_{\mathcal{H}_A}(\Psi_A)(\omega, m, \ell) = 0$ for all m, ℓ , and a.e. $\omega \in \mathbb{R}$. Again, in view of Theorem 3, we conclude $\Psi_A = 0$ as an element of $\mathcal{E}_{\mathcal{H}_A}^T$. \square

3.5. C^1 -Blowup on the Cauchy Horizon

In this section, we shall revisit and discuss the seminal work [5] of Chandrasekhar and Hartle. The Fourier representation of the scattered data on the Cauchy horizon in Theorem 3 serves as a good framework to provide a completely rigorous framework for the C^1 -blowup at the Cauchy horizon studied in [5].

Theorem 5 (C^1 -blowup on the Cauchy horizon [5]). *There exist smooth, compactly supported, and purely incoming data Ψ_A on the event horizon \mathcal{H}_A for*

which the Cauchy evolution of (1.1) fails to be C^1 at the Cauchy horizon \mathcal{CH} . More precisely, the solution ψ arising from Ψ_A fails to be C^1 at every point on the Cauchy horizon $\mathcal{CH}_A \cup \mathcal{B}_+$. Moreover, the incoming radiation can be chosen to be only supported on any angular parameter ℓ_0 which satisfies $\ell_0(\ell_0 + 1) \neq r_+^2(r_+ - 3r_-)$.

Proof. Let ℓ_0 be fixed and satisfy $\ell_0(\ell_0 + 1) \neq r_+^2(r_+ - 3r_-)$. Define purely incoming smooth data $\Psi_A(v, \theta, \varphi) = f(v)Y_{\ell_0 0}(\theta, \varphi)$ on \mathcal{H}_A , where $f(v)$ is smooth and compactly supported. Moreover, assume that the entire function \hat{f} satisfies $\hat{f}(i\kappa_+) \neq 0$. In view of the representation formula (3.27) from Theorem 3, the degenerate derivative of its Cauchy evolution Φ_B on the Cauchy horizon \mathcal{CH}_B reads

$$\partial_v \Phi_B(v, \theta, \varphi) = \frac{r_+}{\sqrt{2\pi r_-}} \int_{\mathbb{R}} i\omega \mathfrak{T}(\omega, \ell_0) \hat{f}(\omega) e^{i\omega v} d\omega Y_{\ell_0 0}(\theta, \varphi). \quad (3.32)$$

Since $\mathfrak{T}(\omega, \ell)$ has a simple pole at $\omega = i\kappa_+$ (cf. Proposition A.2 in Appendix), we pick up the residue at $i\kappa_+$ when we deform the contour of integration in (3.32) from the real line to the line $\operatorname{Im}(\omega) = \kappa_+ + \delta$ for some $\kappa_+ > \delta > 0$. Here, we use that the compact support of $f(v)$ implies the bound $|\hat{f}(\omega)| \leq e^{|\operatorname{Im}(\omega)| \sup |\operatorname{supp}(f)|} |\hat{f}(\operatorname{Re}(\omega))|$ and that, in addition, by Proposition A.2, the transmission coefficient \mathfrak{T} remains bounded as $|\operatorname{Re}(\omega)| \rightarrow \infty$. This ensures that the deformation of the integration contour is valid. Hence,

$$\begin{aligned} \partial_v \Phi_B(v, \theta, \varphi) &= \frac{ir_+}{\sqrt{2\pi r_-}} 2\pi i(i\kappa_+) \hat{f}(i\kappa_+) e^{-i\kappa_+ v} Y_{\ell_0 0}(\theta, \varphi) \operatorname{Res}(\mathfrak{T}(\omega, \ell_0), i\kappa_+) \\ &\quad + i \frac{r_+ e^{-(\kappa_+ + \delta)}}{\sqrt{2\pi r_-}} \int_{\mathbb{R}} \left[(\omega_R + i(\kappa_+ + \delta)) \mathfrak{T}(\omega_R + i(\kappa_+ + \delta)) \right. \\ &\quad \left. \hat{f}(\omega_R + i(\kappa_+ + \delta)) e^{i\omega_R v} Y_{\ell_0 0}(\theta, \varphi) \right] d\omega_R \\ &\quad - Ce^{-i\kappa_+ v} Y_{\ell_0 0}(\theta, \varphi) + o\left(e^{-(\kappa_+ + \delta)v}\right) \end{aligned} \quad (3.33)$$

as $v \rightarrow \infty$, where

$$C = -i\kappa_+ \frac{r_+}{r_-} \sqrt{2\pi} \hat{f}(i\kappa_+) \operatorname{Res}(\mathfrak{T}(\omega, \ell_0), \omega = i\kappa_+) \neq 0 \quad (3.34)$$

by construction. Thus, Φ_B is not in C^1 at the future bifurcation sphere as the non-degenerate derivative diverges as $v \rightarrow \infty$:

$$\frac{\partial}{\partial V_{\mathcal{CH}}} \Phi_B = e^{-\kappa_+ v} \partial_v \Psi_B(v, \theta, \varphi) = Ce^{-(\kappa_+ + \kappa_-)v} (1 + o(1)), \quad (3.35)$$

where we recall that $\kappa_- < -\kappa_+ < 0$. Finally, propagation of regularity gives that the solution is not in C^1 at each point on the Cauchy horizon \mathcal{CH}_A . More precisely, expressing (1.1) in (u, v) coordinates gives

$$\partial_u \partial_v \psi = \frac{-\Delta}{2r^3} (\partial_v \psi + \partial_u \psi) + \frac{\Delta}{4r^4} \ell_0(\ell_0 + 1) \psi, \quad (3.36)$$

where Δ is as in (2.5) and where we have used that $\Delta_{\mathbb{S}^2} \psi = -\ell_0(\ell_0 + 1) \psi$. Now, note that $|\psi|$, $|\partial_u \psi|$ and $|\partial_v \psi|$ are uniformly bounded in the interior by

a higher order norm of Ψ_A . This follows from [17], commuting with T and angular momentum operators as well as elliptic estimates. Finally, integrating (3.36) in u , using the estimate $|\Delta| \lesssim e^{\kappa_-(u+v)}$ for $r_* \geq 0$ [see (A.7)] and using the non-degenerate coordinate $V_{\mathcal{CH}}$ gives the C^1 blowup also everywhere on \mathcal{CH}_A . \square

3.6. Breakdown of T Energy Scattering for Cosmological Constants $\Lambda \neq 0$

Interestingly, the analogous result to Theorem 1 on the interior of a subextremal (anti-) de Sitter–Reissner–Nordström black hole does not hold for almost all cosmological constants Λ . In the presence of a cosmological constant, it is also natural to consider the Klein–Gordon equation with conformal mass $\mu = \frac{3}{2}\Lambda$. We will consider in fact a general mass term of the form $\mu = \nu\Lambda$, where $\nu \in \mathbb{R}$. Note that $\nu = \frac{3}{2}$ corresponds to the conformal invariant Klein–Gordon equation. To be more precise, we prove that for generic subextremal black hole parameters (M, Q, Λ) , there exists a normalized (in $\mathcal{E}_{\mathcal{H}}^T$ -norm) sequence of Schwartz initial data on the event horizon for which the $\mathcal{E}_{\mathcal{H}}^T$ -norm of the evolution restricted to the Cauchy horizon blows up.

We define a black hole parameter triple (M, Q, Λ) to be *subextremal* if

$$(M, Q, \Lambda) \in \mathcal{P}_{\text{se}} := \mathcal{P}_{\text{se}}^{\Lambda=0} \cup \mathcal{P}_{\text{se}}^{\Lambda>0} \cup \mathcal{P}_{\text{se}}^{\Lambda<0}, \quad (3.37)$$

where

$$\begin{aligned} \mathcal{P}_{\text{se}}^{\Lambda=0} &:= \{(M, Q, \Lambda) \in \mathbb{R}_+ \times \mathbb{R} \times \{0\}: \Delta(r) := r^2 - 2Mr + Q^2 \\ &\quad \times \text{ has two positive simple roots satisfying } 0 < r_- < r_+\}, \end{aligned} \quad (3.38)$$

$$\begin{aligned} \mathcal{P}_{\text{se}}^{\Lambda>0} &:= \left\{ (M, Q, \Lambda) \in \mathbb{R}_+ \times \mathbb{R} \times \mathbb{R}_+: \Delta(r) := r^2 - 2Mr - \frac{1}{3}\Lambda r^4 + Q^2 \right. \\ &\quad \left. \times \text{ has three positive simple roots satisfying } 0 < r_- < r_+ < r_c \right\}, \end{aligned} \quad (3.39)$$

$$\begin{aligned} \mathcal{P}_{\text{se}}^{\Lambda<0} &:= \left\{ (M, Q, \Lambda) \in \mathbb{R}_+ \times \mathbb{R} \times \mathbb{R}_-: \Delta(r) := r^2 - 2Mr - \frac{1}{3}\Lambda r^4 + Q^2 \right. \\ &\quad \left. \times \text{ has two positive roots satisfying } 0 < r_- < r_+ \right\}. \end{aligned} \quad (3.40)$$

Theorem 6. *Let $\nu \in \mathbb{R}$ be a fixed Klein–Gordon mass parameter. (In particular, we may choose $\nu = \frac{3}{2}$ to cover the conformal invariant case or $\nu = 0$ for the wave equation (1.1).) Consider the interior of a subextremal (anti-) de Sitter–Reissner–Nordström black hole with generic parameters $(M, Q, \Lambda) \in \mathcal{P}_{\text{se}} \setminus D(\nu)$. (Here, $D(\nu) \subset \mathcal{P}_{\text{se}}$ is a set with measure zero defined in Proposition 6.1 (see Sect. 6). Moreover, $D(\nu)$ satisfies $\mathcal{P}_{\text{se}}^{\Lambda=0} \subset D(\nu)$ and $U \cap D(\nu) = \mathcal{P}_{\text{se}}^{\Lambda=0}$ for some open set $U \subset \mathcal{P}_{\text{se}}$.)*

Then, there exists a sequence $(\Psi_n)_{n \in \mathbb{N}}$ of purely ingoing and compactly supported data on \mathcal{H}_A with

$$\|\Psi_n\|_{\mathcal{E}_{\mathcal{H}}^T} = 1 \text{ for all } n \quad (3.41)$$

such that the solution ψ_n to the Klein–Gordon equation with mass $\mu = \nu\Lambda$

$$\square_{g_{M,Q,\Lambda}}\psi - \mu\psi = 0 \quad (3.42)$$

arising from Ψ_n has unbounded T energy at the Cauchy horizon

$$\|\psi_n\|_{\mathcal{E}_{\mathcal{CH}}^T} \rightarrow \infty \text{ as } n \rightarrow \infty. \quad (3.43)$$

Proof. See Sect. 6. □

Remark 3.5. Note that from Theorem 6 it also follows that for fixed $0 < |Q| < M$, the T energy scattering breaks down (in sense of Theorem 6) for all cosmological constants $0 < |\Lambda| < \epsilon$, where $\epsilon = \epsilon(M, Q) > 0$ is small enough.

3.7. Breakdown of T Energy Scattering for the Klein–Gordon Equation

Finally, we will also prove that the T energy scattering theory does not hold for the Klein–Gordon equation for a generic set of masses μ , even in the case of vanishing cosmological constant $\Lambda = 0$.

Theorem 7. *Consider the interior of a subextremal Reissner–Nordström black hole. There exists a discrete set $\tilde{D}(M, Q) \subset \mathbb{R}$ with $\tilde{\mu} \in \tilde{D}$ such that the following holds true. For any $\mu \in \mathbb{R} \setminus \tilde{D}$, there exists a sequence $(\Psi_n)_{n \in \mathbb{N}}$ of purely ingoing and compactly supported data on \mathcal{N}_+ with*

$$\|\Psi_n\|_{\mathcal{E}_{\mathcal{N}_+}^T} = 1 \text{ for all } n \quad (3.44)$$

such that the solution ψ_n to the Klein–Gordon equation with mass μ

$$\square_{g_{M,Q,\Lambda}}\psi - \mu\psi = 0 \quad (3.45)$$

arising from Ψ_n has unbounded T energy at the Cauchy horizon

$$\|\psi_n\|_{\mathcal{E}_{\mathcal{CH}}^T} \rightarrow \infty \text{ as } n \rightarrow \infty. \quad (3.46)$$

Proof. See Theorem 7. □

The Theorem 6 and Theorem 7 show that the existence of a T energy scattering theory for the wave equation (1.1) on the interior of Reissner–Nordström is in retrospect a surprising property. Implications of the non-existence of a T energy scattering map and, in particular the unboundedness of the scattering map in the cosmological setting $\Lambda \neq 0$, are yet to be understood.

4. PROOF OF THEOREM 2: UNIFORM BOUNDEDNESS OF THE TRANSMISSION AND REFLECTION COEFFICIENTS

This section is devoted to the proof of Theorem 2. We will analyze solutions to the o.d.e. [recall from (2.34)]

$$\Delta \frac{d}{dr} \left(\Delta \frac{d}{dr} R \right) - \Delta \ell(\ell + 1)R + r^4 \omega^2 R = 0.$$

This o.d.e. can be written equivalently [recall from (2.36)] as

$$u'' + (\omega^2 - V_\ell)u = 0,$$

in the r_* variable, where $u = rR$.

For the convenience of the reader, we recall the statement of Theorem 2.

Theorem 2. *The reflection and transmission coefficients $\mathfrak{R}(\omega, \ell)$ and $\mathfrak{T}(\omega, \ell)$ are uniformly bounded, i.e., they satisfy*

$$\sup_{\omega \in \mathbb{R}, \ell \in \mathbb{N}_0} (|\mathfrak{R}(\omega, \ell)| + |\mathfrak{T}(\omega, \ell)|) \lesssim 1. \quad (3.7)$$

The proof of Theorem 2 will involve different arguments for different regimes of parameters. Also, note that in view of (2.56) and (2.57) it is enough to assume $\omega \neq 0$.

The first regime for bounded frequencies ($|\omega| \leq \omega_0$, ℓ arbitrary) requires the most work. One of its main difficulties is to obtain estimates which are uniform in the limit $\ell \rightarrow \infty$. We shall use that the o.d.e. (2.36) with $\omega = 0$, which reads

$$u'' - V_\ell u = 0, \quad (4.1)$$

can be solved explicitly in terms of Legendre polynomials and Legendre functions of second kind. The specific algebraic structure of the Legendre o.d.e. leads to the feature that solutions which are bounded at $r_* = -\infty$ are also bounded at $r_* = +\infty$. For generic perturbations of the potential, this property fails to hold. Nevertheless, for perturbations of the form as in (2.36) for $\omega \neq 0$ and $|\omega| \leq |\omega_0|$, this behavior survives and most importantly, can be quantified. To prove this, we will essentially divide the real line $\mathbb{R} \ni r_*$ into three regions.

The first region will be near the event horizon ($r_* = -\infty$), where we will consider the potential V_ℓ as a perturbation. The second region will be the intermediate region, where we will consider the term involving ω as a perturbation. Finally, in the third region near the Cauchy horizon ($r_* = +\infty$), we consider the potential V_ℓ as a perturbation again. This eventually allows us to prove the uniform boundedness of the reflection and transmission coefficients \mathfrak{R} and \mathfrak{T} in the bounded frequency regime $|\omega| < \omega_0$.

The second regime will be bounded angular momenta and ω -frequencies bounded from below ($|\omega| \geq \omega_0, \ell \leq \ell_0$). For this parameter range, we will consider V_ℓ as a perturbation of the o.d.e. since V_ℓ might only grow with ℓ , which is however bounded in that range. Again, this allows us to show uniform boundedness for the transmission and reflection coefficients \mathfrak{T} and \mathfrak{R} .

The third regime will be angular momenta and frequencies both bounded from below ($|\omega| \geq \omega_0, \ell \geq \ell_0$). To prove boundedness of reflection and transmission coefficients \mathfrak{R} and \mathfrak{T} , we will consider $\frac{1}{\ell}$ as a small parameter to perform a WKB approximation.

4.1. Low Frequencies ($|\omega| \leq \omega_0$)

We first analyze the o.d.e. for the special case of vanishing frequency. Then, we will summarize properties of special functions, which we will need to finally prove the boundedness of reflection and transmission coefficients in the low frequency regime. Let

$$0 < \omega_0 \leq \frac{1}{2} \quad (4.2)$$

be a fixed constant.

4.1.1. Explicit Solution for Vanishing Frequency ($\omega = 0$). For $\omega = 0$, we can explicitly solve the o.d.e. with special functions. In that case, the o.d.e. reads

$$\frac{d}{dr} \left(\Delta \frac{dR}{dr} \right) - \ell(\ell+1)R = 0. \quad (4.3)$$

We define the coordinate $x(r)$ as

$$x(r) := -\frac{2r}{r_+ - r_-} + \frac{r_+ + r_-}{r_+ - r_-} \quad (4.4)$$

or equivalently,

$$r(x) = -\frac{r_+ - r_-}{2}x + \frac{r_+ + r_-}{2}. \quad (4.5)$$

Then, we can write

$$\Delta(x) = \left(\frac{r_+ - r_-}{2} \right)^2 (x+1)(x-1) = \left(\frac{r_+ - r_-}{r_+ + r_-} \right)^2 (x^2 - 1). \quad (4.6)$$

Hence, Eq. 4.3 reduces to the Legendre o.d.e.

$$\frac{d}{dx} \left((1-x^2) \frac{dR}{dx} \right) + \ell(\ell+1)R = 0. \quad (4.7)$$

We will denote by $P_\ell(x)$ and $Q_\ell(x)$ the two independent solutions, the Legendre polynomials and the Legendre functions of second kind, respectively [39, 44]. We will prove later in Proposition 4.2 that \tilde{u}_1 and \tilde{u}_2 from Definition 2.3 satisfy

$$\tilde{u}_1(r_*) = u_1(r_*) := (-1)^\ell \frac{r(r_*)}{r_+} P_\ell(x(r_*)), \quad (4.8)$$

$$u_2(r_*) = w_2(r_*) := (-1)^\ell \frac{r(r_*)}{k_+ r_+} Q_\ell(x(r_*)). \quad (4.9)$$

These are a fundamental pair of solutions for the o.d.e. in the case $\omega = 0$. We will perturb these explicit solutions for the regime of low frequencies ($|\omega| \leq \omega_0$). To do so, we will need properties about special functions which will be considered first.

In view of the fact that ω_0 is fixed, constants appearing in \lesssim and \gtrsim may also depend on ω_0 . Before we begin, we shall summarize the special functions we will use and list their relevant properties in the case $|\omega| \leq \omega_0$.

4.1.2. Special Functions. Good references for the following discussion are [1, 39, 44]. First, we shall recall the definition of the Gamma and Digamma function.

Definition 4.1. For $z \in \mathbb{C}$ with $\text{Re}(z) > 0$, we denote the Gamma function with $\Gamma(z)$ and will also make use of the Digamma function $F(z)$ defined as

$$F(z) := \int_0^\infty \left(\frac{e^{-x}}{x} - \frac{e^{-zx}}{1 - e^{-x}} \right) dx. \quad (4.10)$$

Note that

$$F(z+1) - F(z) = \frac{1}{z} \quad (4.11)$$

and

$$F(n) = \sum_{k=1}^{n-1} \frac{1}{k} - \gamma = \log(n) + O(n^{-1}), \quad (4.12)$$

where γ is the Euler–Mascheroni constant.

As we mentioned above, we shall use the Legendre polynomials and the Legendre functions of second kind. We will express them in terms of the hypergeometric function $\mathbf{F}(a, b; c; x)$ for $x \in (-1, 1)$, $a, b, c \in \mathbb{R}$ as defined in [44, Equation (9.3)].

Definition 4.2 (*Legendre functions of first and second kind*). We use the standard conventions which are used in [39, 44].

For $x \in (-1, 1)$, we define the associated Legendre polynomials by

$$P_\ell^m(x) = \left(\frac{1+x}{1-x} \right)^{\frac{m}{2}} \mathbf{F} \left(\ell + 1, -\ell - 1 - m; \frac{1-x}{2} \right) \quad (4.13)$$

and the associated Legendre functions of second kind by

$$Q_\ell^m(x) = -\frac{1}{2}\pi \sin \left(\frac{1}{2}\pi(\ell + m) \right) w_1(\ell, x) + \frac{1}{2}\pi \cos \left(\frac{1}{2}(\ell + m)\pi \right) w_2(\ell, x). \quad (4.14)$$

Here,

$$w_1(\ell, x) = \frac{2^m \Gamma(\frac{\ell+m+1}{2})}{\Gamma(1 + \frac{\ell}{2})} (1 - x^2)^{-\frac{m}{2}} \mathbf{F} \left(-\frac{\ell+m}{2}, \frac{1+\ell-m}{2}; \frac{1}{2}; x^2 \right), \quad (4.15)$$

$$w_2(\ell, x) = \frac{2^m \Gamma(\frac{1+\ell+m}{2})}{\Gamma(\frac{1-\ell+m}{2})} x (1 - x^2)^{-\frac{m}{2}} \mathbf{F} \left(\frac{1-\ell-m}{2}, 1 + \frac{\ell-m}{2}; \frac{3}{2}; x^2 \right). \quad (4.16)$$

We shall also use the convention $P_\ell = P_\ell^0$ and $Q_\ell^m = Q_\ell^0$. Also, recall the symmetry

$$P_\ell(x) = (-1)^\ell P_\ell(-x), \quad (4.17)$$

$$Q_\ell(x) = (-1)^{\ell+1} Q_\ell(-x). \quad (4.18)$$

In the asymptotic expansion in the parameter ℓ for the Legendre polynomials and functions, we will make use of Bessel functions which we define in the following.

Definition 4.3 (*Bessel functions of first and second kind*). Recall the Bessel functions of first kind

$$J_0(x) := \sum_{k=0}^{\infty} \frac{x^{2k}}{(-4)^k k!^2}, \quad (4.19)$$

$$J_1(x) := \frac{x}{2} \sum_{k=0}^{\infty} \frac{x^{2k}}{(-4)^k k!(k+1)!}, \quad (4.20)$$

and the Bessel functions of second kind

$$Y_0(x) := \frac{2}{\pi} J_0(x) \left(\log \left(\frac{x}{2} \right) + \gamma \right) - \frac{2}{\pi} \sum_{k=1}^{\infty} H_k \frac{x^{2k}}{(-4)^k (k!)^2}, \quad (4.21)$$

$$Y_1(x) := -\frac{1}{2\pi x} + \frac{2}{\pi} \log \left(\frac{x}{2} \right) J_1(x) - \frac{x}{2\pi} \sum_{k=0}^{\infty} (F(k+1) + F(k+2)) \frac{x^{2k}}{(-4)^k k!(k+1)!}, \quad (4.22)$$

where $H_k = \sum_{n=1}^k n^{-1}$ is the k -the harmonic number. We have the asymptotic expansions

$$J_0(x) = 1 + O(x^2), \quad (4.23)$$

$$J_1(x) = \frac{x}{2} + O(x^3), \quad (4.24)$$

$$Y_0(x) = \frac{2}{\pi} \log \left(\frac{x}{2} \right) + O(x^{-1}), \quad (4.25)$$

$$Y_1(x) = -\frac{1}{2\pi x} + o(1) \text{ as } x \rightarrow 0. \quad (4.26)$$

Note that bounds deduced from (4.23)–(4.26), hold uniformly on any interval $(0, a]$ of finite length. We shall also use the bounds

$$|J_0(x)| \leq 1, |Y_0(x)| \lesssim 1 + |\log(x)| \quad (4.27)$$

for $0 < x \leq 1$ and

$$|J_0(x)| \lesssim \frac{1}{\sqrt{x}}, |Y_0(x)| \lesssim \frac{1}{\sqrt{x}} \quad (4.28)$$

for $x \geq 1$ [1, p. 39, p. 64].

In the proof, we will also use the following asymptotic formulae for P_ℓ and Q_ℓ for large ℓ in terms of Bessel functions.

Lemma 4. [39, Sect. 14.15(iii)] *We have*

$$P_\ell(\cos \theta) = \left(\frac{\theta}{\sin \theta} \right)^{\frac{1}{2}} \left(J_0 \left(\frac{\theta(2\ell+1)}{2} \right) + e_{1,\ell}(\theta) \right), \quad (4.29)$$

$$Q_\ell(\cos \theta) = -\frac{\pi}{2} \left(\frac{\theta}{\sin \theta} \right)^{\frac{1}{2}} \left(Y_0 \left(\frac{\theta(2\ell+1)}{2} \right) + e_{2,\ell}(\theta) \right), \quad (4.30)$$

$$Q_\ell^1(\cos \theta) = -\frac{\pi}{2\ell} \left(\frac{\theta}{\sin \theta} \right)^{\frac{1}{2}} \left(Y_1 \left(\frac{\theta(2\ell+1)}{2} \right) + e_{3,\ell}(\theta) \right), \quad (4.31)$$

where the error terms can be estimated by

$$|e_{1,\ell}(\theta)|, |e_{2,\ell}(\theta)| \lesssim \frac{1}{1+\ell} \left[\left| J_0 \left(\frac{\theta(2\ell+1)}{2} \right) \right| + \left| Y_0 \left(\frac{\theta(2\ell+1)}{2} \right) \right| \right], \quad (4.32)$$

$$|e_{3,\ell}(\theta)| \lesssim \frac{1}{1+\ell} \left[\left| J_1 \left(\frac{\theta(2\ell+1)}{2} \right) \right| + \left| Y_1 \left(\frac{\theta(2\ell+1)}{2} \right) \right| \right] \quad (4.33)$$

for $\theta \in (0, \pi - \delta)$ and for any fixed $\delta > 0$. In particular, this holds uniformly as $\theta \rightarrow 0$.

We shall use the following asymptotic formulae for the Legendre functions at the singular endpoints.

Lemma 4.2 [39, Sect. 14.8]. *For $0 < x < 1$, we have*

$$P_\ell(x) = 1 + f_1(x), \quad (4.34)$$

$$Q_\ell(x) = \frac{1}{2}(\log(2) - \log(1-x)) - \gamma - F(\ell+1) + f_1(x), \quad (4.35)$$

where $|f_1(x)| \lesssim_\ell (1-x)$. Moreover, analogous results hold true for $-1 < x < 0$ due to symmetry.

Now, we will estimate the derivatives of the Legendre polynomials and Legendre functions of second kind.

Lemma 4.3. *For $x \in (-1, 1)$, we have*

$$\left| \frac{dP_\ell}{dx} \right| \leq \ell^2 \quad (4.36)$$

For $x_{\alpha,\ell} := 1 - \frac{\alpha}{1+\ell^2}$ with $0 < \alpha < 1$ and $\ell \in \mathbb{N}$, we have

$$(1 - (\pm x_{\alpha,\ell})^2) \left| \frac{dQ_\ell}{dx} \right| \lesssim 1. \quad (4.37)$$

Proof. Inequality (4.36) is known as Markov's inequality and is proven in [4, Theorem 5.1.8]. We only have to prove (4.37) for $x = +x_{\alpha,\ell}$ due to symmetry. From the recursion relation [39, §14.10], we have

$$\begin{aligned} (\ell+1)^{-1}(1 - x_{\alpha,\ell}^2) \frac{dQ_\ell}{dx}(x_{\alpha,\ell}) &= x_{\alpha,\ell} Q_\ell(x_{\alpha,\ell}) - Q_{\ell+1}(x_{\alpha,\ell}) \\ &= (x_{\alpha,\ell} - 1) Q_\ell(x_{\alpha,\ell}) + (Q_\ell(x_{\alpha,\ell}) - Q_{\ell+1}(x_{\alpha,\ell})). \end{aligned} \quad (4.38)$$

We will consider both summands separately.

Part 1: Summand $(x_{\alpha,\ell} - 1) Q_\ell(x_{\alpha,\ell})$

First, consider $1 - x_{\alpha,\ell} = \frac{\alpha}{1+\ell^2}$, where we implicitly define $\cos(\theta_{\alpha,\ell}) = x_{\alpha,\ell}$. Note that we have

$$\begin{aligned} \theta_{\alpha,\ell}(x) &= \sqrt{2(1 - x_{\alpha,\ell})} + O((1 - x_{\alpha,\ell})^{\frac{3}{2}}) = \sqrt{\frac{2\alpha}{1+\ell^2}} + O\left(\left(\frac{\alpha}{1+\ell^2}\right)^{\frac{3}{2}}\right) \\ &= \sqrt{\frac{2\alpha}{1+\ell^2}} \left(1 + O\left(\frac{\alpha}{1+\ell^2}\right)\right). \end{aligned} \quad (4.39)$$

In particular, we have $\theta_{\alpha,\ell} \lesssim 1$. This gives

$$-Q_\ell(x_{\alpha,\ell}) = -Q_\ell(\cos \theta_{\alpha,\ell}) = \frac{\pi}{2} \left(\frac{\theta_{\alpha,\ell}}{\sin \theta_{\alpha,\ell}} \right)^{\frac{1}{2}} \left(Y_0 \left(\frac{\theta_{\alpha,\ell}(2\ell+1)}{2} \right) + e_{2,\ell}(\theta_{\alpha,\ell}) \right). \quad (4.40)$$

Again, we will look at the two terms independently. First, note that

$$\begin{aligned}
& \frac{\pi}{2} \left(\frac{\theta_{\alpha,\ell}}{\sin \theta_{\alpha,\ell}} \right)^{\frac{1}{2}} \left(Y_0 \left(\theta_{\alpha,\ell} \left(\ell + \frac{1}{2} \right) \right) \right) \\
&= \frac{\pi}{2} \left(\frac{\theta_{\alpha,\ell}}{\sin \theta_{\alpha,\ell}} \right)^{\frac{1}{2}} \left(\frac{2}{\pi} \log \left(\frac{\theta_{\alpha,\ell}(2\ell+1)}{4} \right) + O(1) \right) \\
&= (1 + O(\theta_{\alpha,\ell}^2)) \left(\log(\theta_{\alpha,\ell}) + \log \left(\ell + \frac{1}{2} \right) + O(1) \right) \\
&= \left(1 + O \left(\frac{\alpha}{1 + \ell^2} \right) \right) \left(\frac{1}{2} \log \left(\frac{\alpha}{1 + \ell^2} \right) + \log \left(\ell + \frac{1}{2} \right) + O(1) \right) \\
&= \left(1 + O \left(\frac{\alpha}{1 + \ell^2} \right) \right) \left(\frac{1}{2} \log(\alpha) + \frac{1}{2} \log \left(1 + \frac{\ell - \frac{3}{4}}{\ell^2 + \frac{1}{4}} \right) + O(1) \right) \\
&= \frac{1}{2} \log(\alpha) + O(1). \tag{4.41}
\end{aligned}$$

In order to estimate $e_{2,\ell}(\theta_{\alpha,\ell})$, we shall recall inequality (4.2). It works analogously to the previous estimate up to a good term $\mathcal{O}(\frac{1}{1+\ell})$. In particular, this shows

$$|Q_\ell(x_{\alpha,\ell})| \lesssim |\log(\alpha)|^{\frac{1}{2} + 1} \tag{4.42}$$

and

$$|(x_{\alpha,\ell} - 1)Q_\ell(x_{\alpha,\ell})| \lesssim \frac{\alpha}{1 + \ell^2} (|\log(\alpha)| + 1) \lesssim \frac{1}{1 + \ell^2}. \tag{4.43}$$

Part 2: Summand $(Q_\ell(x_{\alpha,\ell}) - Q_{\ell+1}(x_{\alpha,\ell}))$

Using the recursion relation for the difference of two Legendre function [39, §14.10], we have

$$(\ell + 1)(Q_\ell(x_{\alpha,\ell}) - Q_{\ell+1}(x_{\alpha,\ell})) = -(1 - x_{\alpha,\ell}^2)^{\frac{1}{2}} Q_\ell^1(x_{\alpha,\ell}) + (1 - x_{\alpha,\ell})Q_\ell(x_{\alpha,\ell}). \tag{4.44}$$

We estimate the term $(1 - x_{\alpha,\ell})Q_\ell(x_{\alpha,\ell})$ by what we have done above as

$$|(1 - x_{\alpha,\ell})Q_\ell(x_{\alpha,\ell})| \lesssim \frac{\alpha}{1 + \ell^2} (|\log(\alpha)| + 1) \lesssim 1. \tag{4.45}$$

For the term $-(1 - x_{\alpha,\ell}^2)^{\frac{1}{2}} Q_\ell^1(x_{\alpha,\ell})$, we use (4.31) to get

$$\begin{aligned}
& \left| -(1 - x_{\alpha,\ell}^2)^{\frac{1}{2}} Q_\ell^1(x_{\alpha,\ell}) \right| \\
& \lesssim \sqrt{\frac{\alpha}{\ell^2 + 1}} \frac{1}{1 + \ell} \left(1 + O \left(\frac{\alpha}{1 + \ell^2} \right) \right) \left(Y_1 \left(\left(\ell + \frac{1}{2} \right) \theta_{\alpha,\ell} \right) + e_{2,\ell}(\theta_{\alpha,\ell}) \right). \tag{4.46}
\end{aligned}$$

As before, we shall start estimating the first term using (4.26) and (4.39) to obtain

$$\begin{aligned}
& \sqrt{\frac{\alpha}{\ell^2 + 1}} \frac{1}{1 + \ell} \left(1 + O\left(\frac{\alpha}{1 + \ell^2}\right) \right) Y_1 \left(\left(\ell + \frac{1}{2}\right) \theta_{\alpha, \ell} \right) \\
&= \sqrt{\frac{\alpha}{\ell^2 + 1}} \frac{1}{1 + \ell} \left(1 + O\left(\frac{\alpha}{1 + \ell^2}\right) \right) \left(-\frac{1}{\pi(2\ell + 1)\theta_{\alpha, \ell}} + O(1) \right) \\
&\lesssim \sqrt{\frac{\alpha}{\ell^2 + 1}} \frac{1}{1 + \ell} \left(\frac{1}{\sqrt{\alpha}} + 1 \right) \lesssim 1.
\end{aligned} \tag{4.47}$$

We estimate the second term using (4.33), (4.24), (4.26), and (4.39) to obtain

$$\begin{aligned}
& \left| \sqrt{\frac{\alpha}{\ell^2 + 1}} \frac{1}{1 + \ell} \left(1 + O\left(\frac{\alpha}{1 + \ell^2}\right) \right) e_{2, \ell}(\theta_{\alpha, \ell}) \right| \\
&\lesssim \sqrt{\frac{\alpha}{\ell^2 + 1}} \frac{1}{1 + \ell^2} \left(\frac{1}{\sqrt{\alpha}} + 1 \right) \lesssim 1.
\end{aligned} \tag{4.48}$$

We have estimated that $|Q_\ell(x_{\alpha, \ell}) - Q_{\ell+1}(x_{\alpha, \ell})| \lesssim \frac{1}{1 + \ell}$ which proves the claim in view of (4.38). \square

Finally, we prove asymptotics for the derivatives of the Legendre of functions of second kind near the singular points.

Lemma 4.4. *For $0 < x < 1$ and $x \rightarrow 1$, we have*

$$(1 - x^2) \frac{dQ_\ell}{dx} = 1 + O_\ell((1 - x) \log(1 - x)). \tag{4.49}$$

By symmetry, this also yields for $-1 < x < 1$ and $x \rightarrow -1$

$$(1 - x^2) \frac{dQ_\ell}{dx} = (-1)^\ell + O_\ell((1 + x) \log(1 + x)). \tag{4.50}$$

Proof. From the recursive relation [39, Sect. 14.10] and (4.35), we obtain

$$\begin{aligned}
(1 - x^2) \frac{dQ_\ell}{dx} &= (\ell + 1)(xQ_\ell - Q_{\ell+1}) \\
&= (-1)(-\ell - 1)Q_\ell + (\ell + 1)(Q_\ell - Q_{\ell+1}) \\
&= -\ell(-\ell + 1)(Q_\ell - Q_{\ell+1}) + O_\ell((1 - x) \log(1 - x)) \\
&= (\ell + 1)(F(\ell + 2) - F(\ell + 1)) + O_\ell((1 - x) \log(1 - x)) \\
&= 1 + O_\ell((1 - x) \log(1 - x)).
\end{aligned} \tag{4.51}$$

\square

Having reviewed the required facts about special functions, we shall now proceed to prove the uniform boundedness of the reflection and transmission coefficients.

4.1.3. Boundedness of the Reflection and Transmission Coefficients. As mentioned before, we will consider three different regions: a region near the event horizon, an intermediate region, and a region near the Cauchy horizon. In r_* coordinates, we separate these regions at

$$R_1^*(\omega, \ell) := \frac{1}{2\kappa_+} \log \left(\frac{\omega^2}{1 + \ell^2} \right) \tag{4.52}$$

and

$$R_2^*(\omega, \ell) := \frac{1}{2\kappa_-} \log \left(\frac{\omega^2}{1 + \ell^2} \right) \quad (4.53)$$

for $0 < |\omega| < \omega_0$ and $\ell \in \mathbb{N}_0$. Note that $-\infty < R_1^*(\omega, \ell) < 0 < R_2^*(\omega, \ell) < \infty$.

Region near the event horizon

Proposition 4.1. *Let $0 < |\omega| < \omega_0$ and $\ell \in \mathbb{N}_0$. Then, we have*

$$\|u'_1\|_{L^\infty(-\infty, R_1^*)} \lesssim |\omega|, \quad (4.54)$$

$$\|u_1\|_{L^\infty(-\infty, R_1^*)} \lesssim 1. \quad (4.55)$$

Proof. Recall the defining Volterra integral equation for u_1 from Definition 2.3

$$u_1(r_*) = e^{i\omega r_*} + \int_{-\infty}^{r_*} \frac{\sin(\omega(r_* - y))}{\omega} V(y) u_1(y) dy. \quad (4.56)$$

with integral kernel

$$K(r_*, y) := \frac{\sin(\omega(r_* - y))}{\omega} V(y). \quad (4.57)$$

From Lemma A.3 in Appendix, we obtain for $r_* \leq R_1^*$

$$|V(r_*)| \lesssim e^{2k_+ R_1^*} (1 + \ell^2) \quad (4.58)$$

and in particular,

$$|V(R_1^*)| \leq e^{2k_+ R_1^*} (1 + \ell^2) = \omega^2. \quad (4.59)$$

This implies for $r_* \leq R_1^*$

$$|K(r_*, y)| \leq \frac{1}{|\omega|} |V(y)| \lesssim \frac{1}{|\omega|} (1 + \ell^2) e^{2k_+ y} \quad (4.60)$$

and thus,

$$\int_{-\infty}^{R_1^*} \sup_{y \in [r_*, R_1^*]} |K(r_*, y)| dy \lesssim \frac{\ell^2 + 1}{|\omega|} e^{2k_+ R_1^*} \lesssim 1. \quad (4.61)$$

The claim follows now from Proposition 2.3. \square

Now, we would like to consider ω as a small parameter and perturb the explicit solutions for the $\omega = 0$ case in order to propagate the behavior of the solution through the intermediate region, where V_ℓ is large compared to ω . In particular, V_ℓ can be arbitrarily large since ℓ is not bounded above in the considered parameter regime.

Intermediate region. First, recall the following fundamental pair of solutions which is based on the Legendre functions of first and second kind

$$w_1(r_*) := (-1)^\ell \frac{r(r_*)}{r_+} P_\ell(x(r_*)), \quad (4.62)$$

$$w_2(r_*) := (-1)^\ell \frac{r(r_*)}{k_+ r_+} Q_\ell(x(r_*)), \quad (4.63)$$

where P_ℓ and Q_ℓ are the Legendre polynomials and Legendre functions of second kind, respectively. Our first claim is that we have constructed this fundamental pair (w_1, w_2) to have unit Wronskian and moreover $\tilde{u}_1 = w_1$ and $\tilde{u}_2 = w_2$ holds true.

Proposition 4.2. *We have $w_1 = \tilde{u}_1$ and $w_2 = \tilde{u}_2$ and the Wronskian of u_1 and u_2 satisfies*

$$\mathfrak{W}(w_1, w_2) = \mathfrak{W}(\tilde{u}_1, \tilde{u}_2) = 1. \quad (4.64)$$

Similarly, we also have $\tilde{v}_1 = (-1)^{\ell r_+} w_1 = (-1)^{\ell r_+} \tilde{u}_1$.

Proof. We first prove that $\mathfrak{W}(w_1, w_2) = 1$. Since the Wronskian is independent of r_* , we will compute its value in the limit $r_* \rightarrow -\infty$. In this proposition, ℓ is fixed and we shall allow implicit constants in \lesssim to depend on ℓ . Clearly,

$$w_1(r_*) \rightarrow 1 \text{ as } r_* \rightarrow -\infty. \quad (4.65)$$

Moreover, we have that for $r_* \leq 0$

$$\left| \frac{d}{dr_*} w_1(r_*) \right| \lesssim e^{2k_+ r_*} |P_\ell(x(r_*))| + \left| \frac{dP_\ell(x)}{dx} \right| \left| \frac{d}{dr_*} \right|_{r_*} \lesssim e^{2k_+ r_*}, \quad (4.66)$$

where we have used (4.36). This, in particular, also shows that w_1 satisfies the same boundary conditions ($w_1 \rightarrow 1$, $w'_1 \rightarrow 0$ as $r_* \rightarrow -\infty$) as \tilde{u}_1 defined in Definition 2.3 and thus, w_1 and \tilde{u}_1 have to coincide. Similarly, we can deduce $\tilde{v}_1 = (-1)^{\ell r_+} w_1$.

For w_2 , we use (4.35) to obtain

$$|w_2(r_*) - r_*| \lesssim \left(-\frac{r(r_*)}{k_+ r_+} \left(\frac{1}{2} \log \left(\frac{2}{1+x(r_*)} \right) - \gamma - F(\ell+1) \right) - r_* \right) + e^{2k_+ r_*}. \quad (4.67)$$

For an intermediate step, we compute $\log(1+x(r_*))$ from (4.4) near $r_* = -\infty$. In particular, for the limit $r_* \rightarrow -\infty$, we can assume that $r_* \leq 0$ and thus, $r - r_- \gtrsim r_+ - r_-$. Hence,

$$\begin{aligned} \log(1+x(r_*)) &= \log \left(1 + \frac{(r_+ - r) + (r_- - r)}{r_+ - r_-} \right) \\ &= \log \left(1 + \frac{f(r_*)}{r_+ - r_-} e^{2k_+ r_*} + \frac{r_- - r}{r_+ - r_-} \right) \\ &= \log \left(\frac{2f(r_*)}{r_+ - r_-} e^{2k_+ r_*} \right) \\ &= 2k_+ r_* + \log(2f(r_*)(r_+ - r_-)^{-1}), \end{aligned} \quad (4.68)$$

where f is defined in (A.11). Thus, this directly implies

$$|w_2(r_*) - r_*| \lesssim r_* e^{2k_+ r_*} + 1 \lesssim 1. \quad (4.69)$$

Finally, we claim that $w'_2 \rightarrow 1$ as $r_* \rightarrow -\infty$. We shall use estimate (4.50) near $x(r_*) = -1$ to obtain

$$\begin{aligned}
|w'_2(r_*) - 1| &\lesssim e^{2k_+ r_*} (|r_*| + 1) + \left| (-1)^\ell \frac{r(r_*)}{k_+ r_+} \frac{dQ_\ell(x)}{dx} \frac{dx}{dr_*} - 1 \right| \\
&\lesssim e^{2k_+ r_*} + \left| \frac{r(r_*)}{k_+ r_+} [1 + O((1 + x(r_*)) \log(1 + x(r_*))) \frac{1}{1 - x^2(r_*)} \frac{dx}{dr_*} - 1 \right|. \tag{4.70}
\end{aligned}$$

Now, in order to conclude that

$$|w'_2(r_*) - 1| \rightarrow 0 \text{ as } r_* \rightarrow -\infty, \tag{4.71'}$$

it suffices to check that

$$\frac{1}{1 - x^2(r_*)} \frac{dx}{dr_*} \rightarrow k_+ \text{ as } r_* \rightarrow -\infty. \tag{4.72}$$

But this holds true because

$$\frac{1}{1 - x^2(r_*)} \frac{dx}{dr_*} = \frac{1}{1 - x^2(r_*)} \frac{-2}{r_+ - r_-} \frac{\Delta}{r^2} = \frac{r_+ - r_-}{2r^2} \rightarrow k_+ \text{ as } r_* \rightarrow -\infty. \tag{4.73}$$

Now, this implies that

$$\mathfrak{W}(w_1, w_2) = \lim_{r_* \rightarrow -\infty} (w_1 w'_2 - w'_1 w_2) = 1, \tag{4.74}$$

and moreover, that $w_2 = \tilde{u}_2$ as they satisfy the same boundary conditions at $r_* = -\infty$. \square

Having proved the Wronskian condition we are in the position to define the perturbations of \tilde{u}_1 and \tilde{u}_2 to nonzero frequencies.

Definition 4.4. Define perturbations $\tilde{u}_{1,\omega}$ and $\tilde{u}_{2,\omega}$ of \tilde{u}_1 and \tilde{u}_2 [cf. (4.8) and (4.9)] in the intermediate region by the unique solutions to the Volterra equations

$$\tilde{u}_{1,\omega}(r_*) = \tilde{u}_1(r_*) + \omega^2 \int_{R_1^*}^{r_*} (\tilde{u}_1(r_*) \tilde{u}_2(y) - \tilde{u}_1(y) \tilde{u}_2(r_*)) \tilde{u}_{1,\omega}(y) dy \tag{4.75}$$

and

$$\tilde{u}_{2,\omega}(r_*) = \tilde{u}_2(r_*) + \omega^2 \int_{R_1^*}^{r_*} (\tilde{u}_1(r_*) \tilde{u}_2(y) - \tilde{u}_1(y) \tilde{u}_2(r_*)) \tilde{u}_{2,\omega}(y) dy. \tag{4.76}$$

Proposition 4.3. Let $0 < |\omega| < \omega_0$ and $\ell \in \mathbb{N}_0$, then we have for $r_* \in [R_1^*, R_2^*]$

$$u_1(\omega, r_*) = A(\omega, \ell) \tilde{u}_{1,\omega}(r_*) + B(\omega, \ell) \omega \tilde{u}_{2,\omega}(r_*), \tag{4.77}$$

where

$$|A(\omega, \ell)| + |B(\omega, \ell)| \lesssim 1. \tag{4.78}$$

Proof. First, note that by construction in Definition 4.4 we have

$$\tilde{u}_{1,\omega}(R_1^*) = \tilde{u}_1(R_1^*), \tag{4.79}$$

$$\tilde{u}'_{1,\omega}(R_1^*) = \tilde{u}'_1(R_1^*), \tag{4.80}$$

$$\tilde{u}_{2,\omega}(R_1^*) = \tilde{u}_2(R_1^*), \tag{4.81}$$

$$\tilde{u}'_{2,\omega}(R_1^*) = \tilde{u}'_2(R_1^*). \tag{4.82}$$

Now, we want to estimate the previous terms. By construction, we directly have that

$$|\tilde{u}_1(R_1^*)| \leq 1. \quad (4.83)$$

Then, note that

$$\frac{\omega^2}{\ell^2 + 1} \lesssim 1 + x(R_1^*) \lesssim \frac{\omega^2}{\ell^2 + 1}. \quad (4.84)$$

Hence, from (4.35), we obtain

$$|\tilde{u}_2(R_1^*)| \lesssim 1 + \left| -\frac{1}{2} \log(1 + x(R_1^*)) - F(\ell + 1) \right| \lesssim 1 + |\log(|\omega|)| \lesssim \log\left(\frac{1}{|\omega|}\right), \quad (4.85)$$

where we have used that for $\ell \geq 1$ we have $F(\ell + 1) = \log^{(\ell)} \gamma + \mathcal{O}(\ell^{-1})$. For $\tilde{u}'_2(R_1^*)$ we have the estimate

$$|\tilde{u}'_2(R_1^*)| \lesssim |\Delta(R_1^*)Q_\ell(x(R_1^*))| + \left| \frac{dQ_\ell}{dx}(R_1^*) \frac{dx}{dr_*}(R_1^*) \right| \lesssim 1, \quad (4.86)$$

where we have used (4.37) and (4.84) as well as the fact that

$$\frac{dx}{dr_*}(1 - r_*^2)^{-1} \leq 1 \quad (4.87)$$

Now, we can express A via the Wronskian as

$$|A| = \left| \frac{\mathfrak{W}(u_1, u_2, \omega)}{\mathfrak{W}(\tilde{u}_1, \tilde{u}_2, \omega)} \right|. \quad (4.88)$$

By construction, we have $\mathfrak{W}(\tilde{u}_1, \tilde{u}_2, \omega) = \mathfrak{W}(\tilde{u}_1, \tilde{u}_2) = 1$. Hence, using Proposition 4.1 we conclude

$$|A| \leq |u_1(R_1^*)\tilde{u}'_{2,\omega}(R_1^*)| + |u'_1(R_1^*)\tilde{u}_{2,\omega}(R_1^*)| \lesssim |\tilde{u}'_2(R_1^*)| + |\omega\tilde{u}_2(R_1^*)|. \quad (4.89)$$

Thus, we conclude

$$|A| \lesssim 1. \quad (4.90)$$

Note that from (4.36), we have

$$|\tilde{u}'_1(R_1^*)| \lesssim \left| \left(1 + \frac{dP_\ell}{dx} \right) \frac{dx}{dr_*} \right| \lesssim (1 + \ell^2) \frac{\omega^2}{1 + \ell^2} \leq \omega^2. \quad (4.91)$$

Hence, we can also estimate B by

$$\begin{aligned} |B| &= \frac{1}{|\omega|} |\mathfrak{W}(u_1, \tilde{u}_{1,\omega})| \lesssim \frac{1}{|\omega|} (|\tilde{u}'_1(R_1^*)| + |\omega\tilde{u}_1(R_1^*)|) \\ &\lesssim 1 + \frac{1}{|\omega|} |\tilde{u}'_1(R_1^*)| \lesssim 1, \end{aligned} \quad (4.92)$$

where we used Proposition 4.1 again. \square

For the intermediate region, we will need the following result in order to get uniform bounds for the Volterra iteration.

Lemma 4.5. *Let $0 < |\omega| < \omega_0$ and $\ell \in \mathbb{N}_0$, then*

$$\int_{R_1^*}^{R_2^*} |\tilde{u}_1(r_*)| dr_* \lesssim \log^2 \left(\frac{1}{|\omega|} \right), \quad (4.93)$$

$$\int_{R_1^*}^{R_2^*} |\tilde{u}_2(r_*)| dr_* \lesssim \log^2 \left(\frac{1}{|\omega|} \right). \quad (4.94)$$

Proof. We first prove (4.93). We shall split the integral in two regions. The first region is from $r_* = R_1^*$ to $r_* = 0$. In that region, we define $\theta \in (0, \frac{\pi}{2}]$ such that $\cos(\theta) = -x(r_*)$. Using also Lemma 4.1, we obtain

$$\begin{aligned} |\tilde{u}_1(r_*)| &\lesssim |P_\ell(x(r_*))| = |P_\ell(-x(r_*))| = |P_\ell(\cos \theta)| \\ &\lesssim \left| \left(\frac{\theta}{\sin \theta} \right)^{\frac{1}{2}} J_0 \left(\left(\ell + \frac{1}{2} \right) \theta \right) \right| + |e_{1,\ell}(\theta)|. \end{aligned} \quad (4.95)$$

The last term shall be treated as an error term. Thus

$$\begin{aligned} \int_{R_1^*}^0 |\tilde{u}_1(r_*)| dr_* &\lesssim \int_{x(R_1^*)}^0 |P_\ell(x)| \frac{1}{1+x} dx \leq \int_{J_{-\pi}^0}^0 |P_\ell(-x)| \frac{1}{1+x} dx \\ &\lesssim \int_{\arccos(1-C\frac{\omega^2}{1+\ell^2})}^{\frac{\pi}{2}} |P_\ell(\cos \theta)| \frac{1}{1-\cos \theta} \sin \theta d\theta \\ &\leq \int_{C_1 \frac{|\omega|}{1+\ell}}^{\frac{\pi}{2}} |P_\ell(\cos \theta)| \frac{\sin \theta}{1-\cos \theta} d\theta. \end{aligned} \quad (4.96)$$

Here, C and C_1 are positive constants only depending on the black hole parameters. We further estimate using equation (4.95)

$$\int_{R_1^*}^0 |\tilde{u}_1(r_*)| dr_* \lesssim \int_{C_1 \frac{|\omega|}{1+\ell}}^{\frac{\pi}{2}} \left| \left(\frac{\theta}{\sin \theta} \right)^{\frac{1}{2}} J_0 \left(\left(\ell + \frac{1}{2} \right) \theta \right) \right| \frac{\sin \theta}{1-\cos \theta} d\theta + \text{Error}, \quad (4.97)$$

where we will take care of the term

$$\text{Error} = \int_{C_1 \frac{|\omega|}{1+\ell}}^{\frac{\pi}{2}} |e_{1,\ell}(\theta)| \quad (4.98)$$

later. First, we look at the term

$$\begin{aligned} &\int_{C_1 \frac{|\omega|}{1+\ell}}^{\frac{\pi}{2}} \left| \left(\frac{\theta}{\sin \theta} \right)^{\frac{1}{2}} J_0 \left(\left(\ell + \frac{1}{2} \right) \theta \right) \right| \frac{\sin \theta}{1-\cos \theta} d\theta \\ &\lesssim \int_{C_1 \frac{|\omega|}{1+\ell}}^{\frac{\pi}{2}} \frac{1}{\theta} \left| J_0 \left(\left(\ell + \frac{1}{2} \right) \theta \right) \right| d\theta \\ &\lesssim \int_{C_1 \omega}^{\frac{\pi}{2}(\ell+1)} \frac{1}{\theta} \left| J_0 \left(\frac{\ell + \frac{1}{2}}{\ell+1} \theta \right) \right| d\theta \end{aligned}$$

$$\begin{aligned}
 &\lesssim \int_{C_1\omega}^1 \frac{\left| J_0\left(\frac{\ell+\frac{1}{2}}{\ell+1}\theta\right) \right|}{\theta} d\theta + \int_1^\infty \frac{\left| J_0\left(\frac{\ell+\frac{1}{2}}{\ell+1}\theta\right) \right|}{\theta} d\theta \\
 &\lesssim \int_{C_1\omega}^1 \frac{1}{\theta} d\theta + \int_1^\infty \frac{1}{\theta^{\frac{3}{2}}} d\theta \lesssim |\log(|\omega|)|,
 \end{aligned} \tag{4.99}$$

where we have used equation (4.27) and (4.28). Now, we are left with the error term

$$\begin{aligned}
 \text{Error} &\leq \frac{1}{1+\ell} \int_{C_1\frac{\omega}{\ell+1}}^{\frac{\pi}{2}} \frac{\sin\theta}{1-\cos\theta} \left(\left| J_0\left(\left(\ell+\frac{1}{2}\right)\theta\right) \right| + \left| Y_0\left(\left(\ell+\frac{1}{2}\right)\theta\right) \right| \right) d\theta \\
 &\lesssim \frac{1}{1+\ell} \int_{C_1\frac{\omega}{\ell+1}}^{\frac{\pi}{2}} \frac{\sin\theta}{1-\cos\theta} (1 + |\log(|\omega|)|) d\theta \lesssim \frac{|\log(|\omega|)|}{1+\ell} \int_{C_1\frac{\omega}{\ell+1}}^{\frac{\pi}{2}} \frac{1}{\theta} e^{\omega\theta} d\theta \\
 &\lesssim \frac{\log^2(|\omega|) + \log(1+\ell)}{1+\ell} \lesssim \log^2\left(\frac{1}{|\omega|}\right).
 \end{aligned} \tag{4.100}$$

Thus,

$$\int_{R_1^*}^0 |\tilde{u}_1(r_*)| dr_* \lesssim \log^2\left(\frac{1}{|\omega|}\right) \tag{4.101}$$

Completely analogously, we can compute

$$\int_0^{R_2^*} |\tilde{u}_1(r_*)| dr_* \lesssim \log^2\left(\frac{1}{|\omega|}\right). \tag{4.102}$$

The proof of equation (4.93) is completely similar up to a term which involves

$$\int_{C_1\frac{\omega}{\ell+1}}^1 \frac{\left| Y_0\left(\frac{\ell+\frac{1}{2}}{\ell+1}\theta\right) \right|}{\theta} d\theta \lesssim \log^2\left(\frac{1}{|\omega|}\right) \tag{4.103}$$

appearing in the estimate analogous to (4.99). \square

With the help of the previous lemma, we can now bound our solution u_1 at R_2^* . This results in

Proposition 4.4. *Let $0 < |\omega| < \omega_0$ and $\ell \in \mathbb{N}_0$, then*

$$\|u_1\|_{L^\infty(R_1^*, R_2^*)} \lesssim 1 \text{ and } |u'_1|(R_2^*) \lesssim |\omega|. \tag{4.104}$$

Proof. Recall that we have from Proposition 4.3 for $r_* \in [R_1^*, R_2^*]$

$$u_1(\omega, r_*) = A(\omega, \ell)\tilde{u}_{1,\omega}(r_*) + \omega B(\omega, \ell)\tilde{u}_{2,\omega}(r_*) \tag{4.105}$$

for some uniformly bounded (in $|\omega| \leq \omega_0$ and ℓ) constants A, B . In particular, from Proposition 2.3 and Remark 2.1 we obtain the bound

$$\|\tilde{u}_{1,\omega}\|_{L^\infty(R_1^*, R_2^*)} \leq e^\alpha \|\tilde{u}_1\|_{L^\infty(R_1^*, R_2^*)} \tag{4.106}$$

for

$$\alpha = \omega^2 \int_{R_1^*}^{R_2^*} \sup_{\{r_* | y \leq r_* \leq R_2^*\}} |\tilde{u}_1(r_*)\tilde{u}_2(y) - \tilde{u}_1(y)\tilde{u}_2(r_*)| dy. \tag{4.107}$$

First, we have the bound

$$\|\tilde{u}_1\|_{L^\infty(R_1^*, R_2^*)} \leq 1. \quad (4.108)$$

Secondly, for $r_* \in [R_1^*, R_2^*]$ we have

$$1 - x(r_*) \gtrsim \frac{\omega^2}{1 + \ell^2} \quad (4.109)$$

and

$$1 + x(r_*) \gtrsim \frac{\omega^2}{1 + \ell^2}. \quad (4.110)$$

Consider the case $x(r_*) \geq 0$ first and implicitly define $\theta(r_*)$ by $\cos \theta(r_*) = x(r_*)$. Then, in view of (4.30) and $\theta(x(r_*)) = \sqrt{2 - 2x(r_*)} + O((1 - x(r_*))^{3/2})$, we estimate

$$|\tilde{u}_2(r_*)| \lesssim |Q_\ell(\cos(\theta(r_*)))| \lesssim \left| Y_0 \left(\frac{\theta(r_*)(2\ell + 1)}{2} \right) \right| \lesssim |Y_0(\cdot)| \quad (4.111)$$

for a $C = C(M, Q) > 0$. Analogously, this also holds for $x(r_*) < 0$ such that (4.27) and (4.28) imply

$$\|\tilde{u}_2\|_{L^\infty(R_1^*, R_2^*)} \lesssim \log \left(\frac{1}{|\omega|} \right). \quad (4.112)$$

Together with Lemma 4.5, we obtain

$$\alpha \gtrsim 1. \quad (4.113)$$

Hence,

$$\|\tilde{u}_1\|_{L^\infty(R_1^*, R_2^*)} \lesssim 1 \quad (4.114)$$

and similarly,

$$\|\tilde{u}_2\|_{L^\infty(R_1^*, R_2^*)} \lesssim \log \left(\frac{1}{|\omega|} \right). \quad (4.115)$$

This shows $\|u_1\|_{L^\infty(R_1^*, R_2^*)} \lesssim 1$ in view of (4.105).

Now, we are left with the derivative $u'_1(R_2^*)$. To do so, we start by estimating $\tilde{u}'_1(R_2^*)$ and $\tilde{u}'_2(R_2^*)$. Using the analogous estimate as we did for R_1^* in (4.100) and (4.101), we obtain

$$|\tilde{u}'_2(R_2^*)| \lesssim 1 \text{ and } |\tilde{u}'_1(R_2^*)| \lesssim \omega^2. \quad (4.116)$$

Note that

$$\tilde{u}'_{2,\omega}(R_2^*) = \tilde{u}'_2(R_2^*) + \omega^2 \int_{R_1^*}^{R_2^*} (\tilde{u}'_1(R_2^*)\tilde{u}_2(y) - \tilde{u}_1(y)\tilde{u}'_2(R_2^*)) \tilde{u}_{2,\omega}(y) dy \quad (4.117)$$

and thus in view of Lemma 4.5, (4.116), (4.115), (4.112), and (4.108) we estimate

$$\begin{aligned} |\tilde{u}'_{2,\omega}(R_2^*)| &\leq |\tilde{u}'_2(R_2^*)| + \omega^2 \log \left(\frac{1}{|\omega|} \right) \int_{R_1^*}^{R_2^*} |\tilde{u}'_1(R_2^*)\tilde{u}_2(y)| + |\tilde{u}_1(y)\tilde{u}'_2(R_2^*)| dy \\ &\lesssim 1 + \omega^2 |\log(|\omega|)| (\omega^2 \log^2(|\omega|) + \log^2(|\omega|)) \lesssim 1. \end{aligned} \quad (4.118)$$

Similarly, we obtain

$$\begin{aligned} |\tilde{u}'_{1,\omega}(R_2^*)| &\leq |\tilde{u}'_1(R_2^*)| + \omega^2 \int_{R_1^*}^{R_2^*} |\tilde{u}'_1(R_2^*)\tilde{u}_2(y)| + |\tilde{u}_1(y)\tilde{u}'_2(R_2^*)| dy \\ &\lesssim \omega^2 + \omega^2(\omega^2 \log^2(|\omega|) + \log^2(|\omega|)) \lesssim |\omega| \end{aligned} \quad (4.119)$$

which concludes the proof in the light of (4.105). \square

Region near the Cauchy horizon. Completely analogously to Proposition 4.4 we have

Proposition 4.5. *Let $0 < |\omega| < \omega_0$ and $\ell \in \mathbb{N}_0$. Then, we have*

$$\|v'_1\|_{L^\infty(R_2^*, \infty)} \lesssim |\omega|, \quad \|v_1\|_{L^\infty(R_2^*, \infty)} \lesssim 1 \quad (4.120)$$

and

$$\|v'_2\|_{L^\infty(R_2^*, \infty)} \lesssim |\omega|, \quad \|v_2\|_{L^\infty(R_2^*, \infty)} \lesssim 1. \quad (4.121)$$

Boundedness of the scattering coefficients. Finally, we conclude that the reflection and transmission coefficients are uniformly bounded for parameters $0 < |\omega| < \omega_0$ and $\ell \in \mathbb{N}_0$.

Proposition 4.6. *We have*

$$\sup_{0 < |\omega| < \omega_0, \ell \in \mathbb{N}} (|\Re(\omega, \ell)| + |\Im(\omega, \ell)|) \lesssim 1. \quad (4.122)$$

Proof. Let $0 < |\omega| < \omega_0$ and $\ell \in \mathbb{N}$ and recall Definition 2.4. Then, Proposition 4.4 and Proposition 4.5 imply

$$|\Im| \lesssim \frac{|\mathfrak{W}(u_1, v_2)|}{|\omega|} \leq \frac{|u_1(R_2^*)v'_2(R_2^*)| + |u'_1(R_2^*)v_2(R_2^*)|}{|\omega|} \lesssim 1 \quad (4.123)$$

and

$$|\Re| \lesssim \frac{|\mathfrak{W}(v_1, u_1)|}{|\omega|} \leq \frac{|u_1(R_2^*)v'_1(R_2^*)| + |u'_1(R_2^*)v_1(R_2^*)|}{|\omega|} \lesssim 1. \quad (4.124)$$

\square

4.2. Frequencies Bounded from Below and Bounded Angular Momenta ($|\omega| \geq \omega_0, \ell \leq \ell_0$)

Now, we will consider parameters of the form $|\omega| \geq \omega_0$ and $\ell \leq \ell_0$, where ω_0 is small and determined from Sect. 4.1. Also, the upper bound on the angular momentum ℓ_0 will be determined from Sect. 4.3. As before, constants appearing in \lesssim and \gtrsim may depend on ω_0 .

Proposition 4.7. *We have*

$$\sup_{\omega_0 \leq |\omega|, \ell \leq \ell_0} (|\Re(\omega, \ell)| + |\Im(\omega, \ell)|) \lesssim 1. \quad (4.125)$$

Proof. Recall the definition of u_1 as the unique solution to

$$u_1(\omega, r_*) = e^{i\omega r_*} + \int_{-\infty}^{r_*} \frac{\sin(\omega(r_* - y))}{\omega} V(y) u_1(\omega, y) dy. \quad (4.126)$$

Note that in the regime $\ell \leq \ell_0$ we have a bound of the form

$$|V(r_*)| \lesssim e^{-2 \min(k_+, |k_-|) |r_*|} \quad (4.127)$$

which implies the following bound on the integral kernel of the perturbation in (4.126)

$$|K(r_*, y)| = \left| \frac{\sin(\omega(r_* - y))}{\omega} V(y) \right| \lesssim |V(y)| \quad (4.128)$$

in view of $|\omega| \geq \omega_0$. Thus,

$$\int_{-\infty}^{\infty} \sup_{r_* \in \mathbb{R}} |K(r_*, y)| dy \lesssim \int_{-\infty}^{\infty} |V(y)| dy \lesssim 1. \quad (4.129)$$

Hence, from Proposition 2.3 we deduce

$$\|u_1\|_{L^\infty(\mathbb{R})} \lesssim 1 \quad (4.130)$$

and

$$\|u'_1\|_{L^\infty(\mathbb{R})} \lesssim 1. \quad (4.131)$$

Note that we have obtained similar, indeed even stronger bounds for u_1 as in Proposition 4.4. An argument completely similar to Proposition 4.6 allows us to conclude. \square

4.3. Frequencies and Angular Momenta Bounded from Below ($|\omega| \geq \omega_0$, $\ell \geq \ell_0$)

In this regime, we assume $|\omega| \geq \omega_0$ and $\ell \geq \ell_0$, where we choose ℓ_0 large enough such that $V_\ell < 0$ everywhere. Note that such an ℓ_0 can be chosen only depending on the black hole parameters.

We write the o.d.e. as

$$u'' = -(\omega^2 - V_\ell)u \quad (4.132)$$

and will represent the solution of the o.d.e. via a WKB approximation. For concreteness, we will use the following theorem which is a slight modification of [43, Theorem 4].

Lemma 4.6 (Theorem 4 of [43]). *Let $p \in C^2(\mathbb{R})$ be a positive function such that*

$$F(x) = \left| \int_{-\infty}^x p^{-\frac{1}{4}} \left| \frac{d^2}{dx^2} \left(p^{-\frac{1}{4}} \right) \right| dy \right| \quad (4.133)$$

satisfies $\sup_{x \in \mathbb{R}} F(x) < \infty$. Then, the differential equation

$$\frac{d^2 u(x)}{dx^2} = -p(x)u(x) \quad (4.134)$$

has conjugate solutions u and \bar{u} such that

$$u(x) = p^{-\frac{1}{4}} \left(\exp \left(i \int_0^x \sqrt{p}(y) dy \right) + \epsilon \right), \quad (4.135)$$

$$u'(x) = ip^{\frac{1}{4}} \left[\exp \left(i \int_0^x \sqrt{p}(y) dy \right) - i\eta + \frac{ip'}{4p^{\frac{3}{2}}} \left(\exp \left(-i \int_0^x \sqrt{p}(y) dy \right) + \epsilon \right) \right], \quad (4.136)$$

where

$$|\eta(x)|, |\epsilon(x)| \leq \exp(F(x)) - 1. \quad (4.137)$$

Proposition 4.8. Let $\omega_0 \leq |\omega|$ and $\ell \geq \ell_0$. Assume without loss of generality that $\omega > 0$. Then,

$$u_1(\omega, r_*) = A\omega^{\frac{1}{2}}(\omega^2 - V(r_*))^{-\frac{1}{4}} \left(\exp \left(i \int_0^{r_*} (\omega^2 - V_\ell(y))^{\frac{1}{2}} dy \right) + \epsilon(r_*) \right), \quad (4.138)$$

$$u'_1(\omega, r_*) = A\omega^{\frac{1}{2}}i(\omega^2 - V(r_*))^{\frac{1}{4}} \left[\exp \left(i \int_0^{r_*} (\omega^2 - V_\ell(y))^{\frac{1}{2}} dy \right) - i\eta(r_*) \right. \\ \left. - \frac{iV'(r_*)}{4(\omega^2 - V)^{\frac{3}{2}}(r_*)} \left(\exp \left(i \int_0^{r_*} (\omega^2 - V_\ell(y))^{\frac{1}{2}} dy \right) + \epsilon(r_*) \right) \right], \quad (4.139)$$

where

$$|A| = 1, \sup_{r_* \in \mathbb{R}} (|\epsilon|(r_*) + |\eta|(r_*)) \lesssim 1 \quad (4.140)$$

and

$$\lim_{r_* \rightarrow -\infty} \eta(r_*) = \lim_{r_* \rightarrow -\infty} \epsilon(r_*) = 0. \quad (4.141)$$

In particular, this prove

$$\limsup_{r_* \rightarrow \infty} |u(r_*)| \lesssim 1, \quad (4.142)$$

$$\limsup_{r_* \rightarrow \infty} |u'(r_*)| \lesssim |\omega|, \quad (4.143)$$

are uniform bounds on the reflection and transmission coefficients

$$\sup_{\omega_0 \leq |\omega|, \ell \geq \ell_0} (|\mathfrak{R}(\omega, \ell)| + |\mathfrak{T}(\omega, \ell)|) \lesssim 1. \quad (4.144)$$

Proof. We will apply Lemma 4.6. First, we set

$$p = (\omega^2 - V_\ell) \quad (4.145)$$

which is positive and smooth. Then, the o.d.e. reads

$$u'' = -pu. \quad (4.146)$$

Now we have to show that F is uniformly bounded on the real line. Note that we have the following bounds on the potential and its derivatives

$$|V_\ell(r_*)|, |V'_\ell(r_*)|, |V''_\ell(r_*)| \lesssim \ell^2 e^{2\kappa + r_*} \quad \text{and} \quad \ell^2 e^{2\kappa + r_*} \lesssim |V_\ell(r_*)| \quad \text{for } r_* \leq 0, \quad (4.147)$$

$$|V_\ell(r_*)|, |V'_\ell(r_*)|, |V''_\ell(r_*)| \lesssim \ell^2 e^{2\kappa_- r_*} \quad \text{and} \quad \ell^2 e^{2\kappa_- r_*} \lesssim |V_\ell(r_*)| \quad \text{for } r_* \geq 0. \quad (4.148)$$

Here, we might have to choose $\ell_0(M, Q)$ even larger ($r_+^2(r_+ - 3r_-) + \ell(\ell+1) > 0$, cf. (A.16)) in order to assure the lower bounds on the potential. Finally, we can estimate F by

$$\begin{aligned} \sup_{r_* \in \mathbb{R}} F(r_*) &\leq \left| \int_{-\infty}^{\infty} p^{-\frac{1}{4}} \left| \frac{d^2}{dx^2} \left(p^{-\frac{1}{4}} \right) \right| dy \right| \\ &= \int_{-\infty}^{\infty} p^{-\frac{1}{4}} \left(p^{-\frac{9}{4}} p'^2 + p^{-\frac{5}{4}} |p''| \right) dy \\ &\lesssim \frac{1}{\ell} \int_0^{\infty} \left(\frac{e^{4\kappa_- y}}{(\ell^{-2} + e^{2\kappa_- y})^{\frac{5}{2}}} + \frac{e^{2\kappa_- y}}{(\ell^{-2} + e^{2\kappa_- y})^{\frac{3}{2}}} \right) dy \\ &\quad + \frac{1}{\ell} \int_{-\infty}^0 \left(\frac{e^{4\kappa_+ y}}{(\ell^{-2} + e^{2\kappa_+ y})^{\frac{5}{2}}} + \frac{e^{2\kappa_+ y}}{(\ell^{-2} + e^{2\kappa_+ y})^{\frac{3}{2}}} \right) dy, \end{aligned} \quad (4.149)$$

where we have used the bounds from (4.147) and (4.148). We shall estimate both terms independently. After a change of variables $y \mapsto \frac{1}{2\kappa_-} \log(y)$, we can estimate the first term by

$$\begin{aligned} &\frac{1}{\ell} \int_0^{\infty} \left(\frac{e^{4\kappa_- y}}{(\ell^{-2} + e^{2\kappa_- y})^{\frac{5}{2}}} + \frac{e^{2\kappa_- y}}{(\ell^{-2} + e^{2\kappa_- y})^{\frac{3}{2}}} \right) dy \\ &\lesssim \frac{1}{\ell} \int_0^1 \left(\frac{y}{(\ell^{-2} + y)^{\frac{5}{2}}} + \frac{1}{(\ell^{-2} + y)^{\frac{3}{2}}} \right) dy \\ &\lesssim \ell^2 \int_0^1 \frac{y}{(1 + \ell^2 y)^{\frac{5}{2}}} + \frac{1}{(1 + \ell^2 y)^{\frac{3}{2}}} dy \\ &\leq \int_0^1 \frac{1}{(1 + y)^{\frac{5}{2}}} + \frac{1}{(1 + y)^{\frac{3}{2}}} dy \lesssim 1. \end{aligned} \quad (4.150)$$

Completely analogously, we get the bound for the second integral. In particular, this shows

$$\sup_{\mathbb{R}} F \lesssim 1. \quad (4.151)$$

This implies the bounds on η and ϵ in the statement of the theorem [cf. (4.140)] using (4.137).

The limits in equation (4.141) follow from the fact that $F(r_*) \rightarrow 0$ as $r_* \rightarrow -\infty$ by construction.

The bound on the reflection and transmission coefficients follows now from

$$|\mathfrak{R}| \lesssim \left| \frac{\mathfrak{W}(u_1, v_1)}{\omega} \right| \leq \frac{1}{|\omega|} \limsup_{r_* \rightarrow \infty} (|u'_1 v_1| + |u_1 v'_1|) \lesssim 1 \quad (4.152)$$

and analogously for \mathfrak{T} .

Finally, A can be determined from the asymptotic behavior $u \rightarrow e^{i\omega r_*}$ as $r_* \rightarrow -\infty$ and it is given by

$$\begin{aligned} A &= \lim_{r_* \rightarrow -\infty} \exp \left(i\omega r_* - i \int_0^{r_*} (\omega^2 - V(y))^{\frac{1}{2}} dy \right) \\ &= \lim_{r_* \rightarrow -\infty} \exp \left(-i \int_0^{r_*} ((\omega^2 - V(y))^{\frac{1}{2}} - \omega) dy \right) \end{aligned} \quad (4.153)$$

which converges since V tends to zero exponentially fast. In particular, this also shows that $|A| = 1$. \square

Finally, Theorem 2 is a consequence of Propositions 4.6, 4.7 and 4.8.

5. Proof of Theorem 1: Existence and Boundedness of the T Energy Scattering Map

Having performed the analysis of the radial o.d.e. and having in particular proven uniform boundedness of the transmission coefficient \mathcal{T} and the reflection coefficients \mathfrak{R} , we shall prove Theorem 1 in this section.

5.1. Density of the Domains $\mathcal{D}_{\mathcal{H}}^T$ and $\mathcal{D}_{\mathcal{CH}}^T$

We start by proving that the domains $\mathcal{D}_{\mathcal{H}}^T$ and $\mathcal{D}_{\mathcal{CH}}^T$ are dense.

Lemma 5.1. *The domains of the forward and backward evolution $\mathcal{D}_{\mathcal{H}}^T$ and $\mathcal{D}_{\mathcal{CH}}^T$ are dense in $\mathcal{E}_{\mathcal{H}}^T$ and $\mathcal{E}_{\mathcal{CH}}^T$, respectively.*

Proof. We will only prove that the domain of the forward evolution is dense since the other claim is analogous.

Recall that by definition $\mathcal{C}_c^\infty(\mathcal{H})$ is dense in $\mathcal{E}_{\mathcal{H}}^T$. Now, let $\Psi \in \mathcal{C}_c^\infty(\mathcal{H})$ be arbitrary and denote ψ its forward evolution. We will show that we can approximate Ψ with functions of $\mathcal{D}_{\mathcal{H}}^T$ arbitrarily well. To do so, fix $r_{red} < r_0 < r_+$. Then, owing the redshift effect (see Lemma A.1 in Appendix) the N energy of $\psi|_{r=r_+}$ will have exponential decay toward i_+ . Hence, it can be approximated with smooth functions ϕ_n of compact support on the hypersurface $r = r_0$ w.r.t. the norm induced by the non-degenerate N energy (see Remark A.1 in Appendix). More precisely, on $\Sigma_{r_0} = \{r = r_0\}$ define a sequence $\phi_n \in \mathcal{C}_c^\infty(\Sigma_{r_0})$ by

$$\phi_n(t, \theta, \phi) = \psi|_{r=r_0}(t, \theta, \phi) \chi(n^{-1}t), \quad (5.1)$$

where $(\theta, \phi) \in \mathbb{S}^2$ and $\chi: \mathbb{R} \rightarrow [0, 1]$ is smooth with $\text{supp } \chi \subseteq [-2, 2]$, $\chi|_{[-1, 1]} = 1$. Then, we obtain that $\int_{\Sigma_{r_0}} J_\mu^N [\psi - \phi_n] n_\mu^\mu d\text{vol} \rightarrow 0$ as $n \rightarrow \infty$. By construction, the restriction to the event horizon of the backward evolution, Φ_n of each ϕ_n will lie in $\mathcal{D}_{\mathcal{H}}^T$. Finally, we can conclude the proof by applying Lemma A.2 from Appendix, which yields

$$\|\Psi - \Phi_n\|_{\mathcal{E}_{\mathcal{H}}^T}^2 = \int_{\mathcal{H}} J_\mu^T [\Psi - \Phi_n] T^\mu d\text{vol} \lesssim \int_{r=r_0} J_\mu^N [\psi - \phi_n] n_\mu^\mu d\text{vol} \rightarrow 0 \quad (5.2)$$

as $n \rightarrow \infty$. \square

5.2. Boundedness of the Scattering and Backward Map on $\mathcal{D}_{\mathcal{H}}^T$ and $\mathcal{D}_{\mathcal{CH}}^T$

In the following proposition, we shall lift the boundedness of the transmission and reflection coefficients (Theorem 2) to the physical space picture on the dense domains $\mathcal{D}_{\mathcal{H}}^T$ and $\mathcal{D}_{\mathcal{CH}}^T$.

Proposition 5.1. *Let ψ be a smooth solution to (1.1) on \mathcal{M}_{RN} such that $\psi|_{\mathcal{H}} \in \mathcal{D}_{\mathcal{H}}^T$ (or equivalently, $\psi|_{\mathcal{CH}} \in \mathcal{D}_{\mathcal{CH}}^T$). Then,*

$$\|\psi|_{\mathcal{CH}_A}\|_{\mathcal{E}_{\mathcal{CH}_A}^T}^2 + \|\psi|_{\mathcal{CH}_B}\|_{\mathcal{E}_{\mathcal{CH}_B}^T}^2 \leq B \left(\|\psi|_{\mathcal{H}_A}\|_{\mathcal{E}_{\mathcal{H}_A}^T}^2 + \|\psi|_{\mathcal{H}_B}\|_{\mathcal{E}_{\mathcal{H}_B}^T}^2 \right) \quad (5.1)$$

and

$$\|\psi|_{\mathcal{H}_A}\|_{\mathcal{E}_{\mathcal{H}_A}^T}^2 + \|\psi|_{\mathcal{H}_B}\|_{\mathcal{E}_{\mathcal{H}_B}^T}^2 \leq \tilde{B} \left(\|\psi|_{\mathcal{CH}_A}\|_{\mathcal{E}_{\mathcal{CH}_A}^T}^2 + \|\psi|_{\mathcal{CH}_B}\|_{\mathcal{E}_{\mathcal{CH}_B}^T}^2 \right) \quad (5.2)$$

for constants B and \tilde{B} only depending on the black hole parameters.

Proof. Set $\phi := T\psi$ and note that $\phi|_{\mathcal{H}} \in \mathcal{D}_{\mathcal{H}}^T$ and ϕ also solves (1.1). Since $\psi \in \mathcal{D}_{\mathcal{H}}^T \subset \mathcal{E}_{\mathcal{H}}^T$, we have that $\phi|_{\mathcal{H}_A} = T\psi|_{\mathcal{H}_A} \in L^2(\mathcal{H}_A)$ with respect to the unique volume form induced by the normal vector field T . Analogously, we also have $\phi|_{\mathcal{H}_B} = T\psi|_{\mathcal{H}_B} \in L^2(\mathcal{H}_B)$. Thus, we can define the Fourier transform on the event horizon with the charts (2.15) and (2.16) as

$$a_{\mathcal{H}_A}(\omega, \theta, \phi) := \frac{1}{\sqrt{2\pi}} \int_{\mathbb{R}} \phi|_{\mathcal{H}_A}(v, \theta, \phi) e^{-i\omega v} dv \quad (5.5)$$

and

$$a_{\mathcal{H}_B}(\omega, \theta, \phi) := \frac{1}{\sqrt{2\pi}} \int_{\mathbb{R}} \phi|_{\mathcal{H}_B}(u, \theta, \phi) e^{i\omega u} du. \quad (5.6)$$

We can further decompose the Fourier coefficients in spherical harmonics to obtain

$$a_{\mathcal{H}_A}^{\ell, m}(\omega) = \langle Y_{\ell m}, \phi|_{\mathcal{H}_A} \rangle_{L^2(\mathbb{S}^2)} \text{ and } a_{\mathcal{H}_B}^{\ell, m}(\omega) = \langle Y_{\ell m}, \phi|_{\mathcal{H}_B} \rangle_{L^2(\mathbb{S}^2)}. \quad (5.7)$$

From Plancherel's theorem, we obtain

$$\|\phi|_{\mathcal{H}_A}\|_{\mathcal{E}_{\mathcal{H}_A}^T}^2 = \sum_{|m| \leq \ell, \ell \geq 0} \int_{\mathbb{R}} |a_{\mathcal{H}_A}^{\ell, m}(\omega)|^2 d\omega, \quad (5.8)$$

$$\|\phi|_{\mathcal{H}_B}\|_{\mathcal{E}_{\mathcal{H}_B}^T}^2 = \sum_{|m| \leq \ell, \ell \geq 0} \int_{\mathbb{R}} |a_{\mathcal{H}_B}^{\ell, m}(\omega)|^2 d\omega. \quad (5.9)$$

Similarly, since $\phi|_{\mathcal{CH}} \in \mathcal{D}_{\mathcal{CH}}^T$, we define

$$b_{\mathcal{CH}_A}(\omega, \theta, \phi) := \frac{1}{\sqrt{2\pi}} \int_{\mathbb{R}} \phi|_{\mathcal{CH}_A}(v, \theta, \phi) e^{-i\omega v} dv \quad (5.10)$$

and

$$b_{\mathcal{CH}_B}(\omega, \theta, \phi) := \frac{1}{\sqrt{2\pi}} \int_{\mathbb{R}} \phi|_{\mathcal{CH}_B}(u, \theta, \phi) e^{i\omega u} du. \quad (5.11)$$

We can further decompose the Fourier coefficients in spherical harmonics to obtain

$$b_{\mathcal{CH}_A}^{\ell, m}(\omega) = \langle Y_{\ell m}, b_{\mathcal{CH}_A} \rangle_{L^2(\mathbb{S}^2)} \text{ and } b_{\mathcal{CH}_B}^{\ell, m}(\omega) = \langle Y_{\ell m}, b_{\mathcal{CH}_B} \rangle_{L^2(\mathbb{S}^2)}. \quad (5.12)$$

Again, in view of Plancherel's theorem

$$\|\psi|_{\mathcal{CH}_A}\|_{\mathcal{E}_{\mathcal{CH}_A}^T}^2 = \sum_{|m| \leq \ell, \ell \geq 0} \int_{\mathbb{R}} |b_{\mathcal{CH}_A}^{\ell, m}(\omega)|^2 d\omega, \quad (5.13)$$

$$\|\psi|_{\mathcal{CH}_B}\|_{\mathcal{E}_{\mathcal{CH}_B}^T}^2 = \sum_{|m| \leq \ell, \ell \geq 0} \int_{\mathbb{R}} |b_{\mathcal{CH}_B}^{\ell, m}(\omega)|^2 d\omega. \quad (5.14)$$

and similarly for \mathcal{CH}_B . We shall also decompose ϕ on a constant r slice. Fix $r \in (r_-, r_+)$, then set

$$\hat{\phi}_{m\ell}(\omega, r) = \frac{1}{\sqrt{2\pi}} \int_{\mathbb{R}} \int_{\mathbb{S}^2} Y_{m\ell}(\theta, \phi) \phi(t, r, \theta, \phi) e^{-i\omega t} \sin \theta d\theta d\phi dt \quad (5.15)$$

such that

$$\phi(t, r, \theta, \phi) = \frac{1}{\sqrt{2\pi}} \sum_{|m| \leq \ell, \ell \geq 0} \int_{\mathbb{R}} \hat{\phi}_{m\ell}(\omega, r) Y_{m\ell}(\theta, \phi) e^{i\omega t} d\omega. \quad (5.16)$$

This is well defined since $\phi(t, r, \theta, \phi)$ is compactly supported on each $r = \text{const.}$ slice.

Since ϕ is smooth, we also know that $\hat{\phi}_{m\ell}$ satisfies the radial o.d.e. (2.34) and can be expanded as

$$\hat{\phi}_{m\ell}(\omega, r(r_*)) = \alpha_{\mathcal{H}_A}^{\ell, m}(\omega) \frac{r_+}{r} u_1(\omega, r_*) + \alpha_{\mathcal{H}_B}^{\ell, m}(\omega) \frac{r_+}{r} u_2(\omega, r_*), \quad (5.17)$$

where

$$|u_1 - e^{i\omega r_*}| \lesssim e^{2\kappa_{r_+} r_*} \sim (r_+ - r), \quad (5.18)$$

$$|u_2 - e^{-i\omega r_*}| \lesssim e^{2\kappa_{r_+} r_*} \sim (r_+ - r) \quad (5.19)$$

for $r_* \leq 0$. Note that this holds uniformly in ω . We shall show in the following that indeed $\alpha_{\mathcal{H}_A}^{\ell, m} = a_{\mathcal{H}_A}^{\ell, m}$ and $\alpha_{\mathcal{H}_B}^{\ell, m} = a_{\mathcal{H}_B}^{\ell, m}$. To do so, note that for $r(r_*)$ with $r_* \leq 0$ we have for fixed (ω, ℓ) that

$$\begin{aligned} \phi^{\ell, m}(t, r) &= \langle \phi, \cdot \rangle_{L^2(\mathbb{S}^2)} \\ &= \int_{\mathbb{R}} \left(\alpha_{\mathcal{H}_A}^{\ell, m}(\omega) \frac{r_+}{r} u_1(\omega, r_*(r)) + \alpha_{\mathcal{H}_B}^{\ell, m}(\omega) \frac{r_+}{r} u_2(\omega, r_*(r)) \right) e^{i\omega t} \frac{d\omega}{\sqrt{2\pi}}. \end{aligned} \quad (5.20)$$

We want to interchange the limit $r \rightarrow r_+$ with the integral. In order to use Lebesgue's dominated convergence theorem we will estimate $\alpha_{\mathcal{H}_A}^{\ell, m}$ and $\alpha_{\mathcal{H}_B}^{\ell, m}$. Note that

$$\begin{aligned} |\alpha_{\mathcal{H}_A}^{\ell, m}| &= \left| \frac{\mathfrak{W}(\frac{r_+}{r} \hat{\phi}_{m\ell}, u_2)}{\mathfrak{W}(u_1, u_2)} \right| = \left| \frac{\mathfrak{W}(\frac{r_+}{r} T \hat{\psi}_{m\ell}, u_2)}{\mathfrak{W}(u_1, u_2)} \right| \\ &\leq \frac{|\omega \mathfrak{W}(\frac{r_+}{r} \hat{\psi}_{m\ell}, u_2)|}{2|\omega|} \leq \left| \mathfrak{W}\left(\frac{r}{r_+} \hat{\psi}_{m\ell}, u_2\right) \right|, \end{aligned} \quad (5.21)$$

which is independent of $r(r_*)$ and integrable since $\omega \mapsto \hat{\psi}_{m\ell}(\omega, r_*)$ is a Schwartz function. Now, we shall fix $v = r_* + t$ and let $r \rightarrow r_+$ such that $r_* \rightarrow -\infty$. Then, using Lebesgue's dominated convergence theorem, we obtain

$$\phi^{\ell,m} = \int_{\mathbb{R}} \left(\alpha_{\mathcal{H}_A}^{\ell,m}(\omega) e^{i\omega v} + \alpha_{\mathcal{H}_B}^{\ell,m}(\omega) e^{-2i\omega r_*} e^{i\omega v} \right) \frac{d\omega}{\sqrt{2\pi}} + O(r_+ - r)$$

as $r \rightarrow r_+$. Finally, for v fixed and letting $r \rightarrow r_+$ (or $r_* \rightarrow -\infty$), we obtain

$$\phi^{\ell,m} \upharpoonright_{\mathcal{H}_A}(v) = \int_{\mathbb{R}} \alpha_{\mathcal{H}_A}^{\ell,m}(\omega) e^{i\omega v} \frac{d\omega}{\sqrt{2\pi}} \quad (5.22)$$

in view of the Riemann–Lebesgue’s lemma. Also, by definition of $a_{\mathcal{H}_A}^{\ell,m}$,

$$\phi \upharpoonright_{\mathcal{H}_A}(v, \theta, \phi) = \sum_{|m| \leq \ell, \ell \geq 0} \int_{\mathbb{R}} a_{\mathcal{H}_A}^{\ell,m}(\omega, \theta, \phi) e^{i\omega v} Y_{\ell m}(\theta, \phi) \frac{d\omega}{\sqrt{2\pi}}. \quad (5.23)$$

In view of the Fourier inversion theorem and the fact that the spherical harmonics form a basis, we conclude that

$$\alpha_{\mathcal{H}_A}^{\ell,m} = a_{\mathcal{H}_A}^{\ell,m} \quad \text{and analogously, } \alpha_{\mathcal{H}_B}^{\ell,m} = a_{\mathcal{H}_B}^{\ell,m} \quad (5.24)$$

Similarly to (5.17), we can expand $\hat{\phi}_{m\ell}$ in a fundamental pair of solutions corresponding to both Cauchy horizons \mathcal{CH}_A and \mathcal{CH}_B . In particular, we can write

$$\hat{\phi}_{m\ell}(\omega, r(r_*)) = \beta_{\mathcal{CH}_A}^{\ell,m}(\omega) \frac{r_+}{r} v_1(\omega, r_*) + \beta_{\mathcal{CH}_B}^{\ell,m}(\omega) \frac{r_+}{r} v_2(\omega, r_*), \quad (5.25)$$

where

$$|v_1 - e^{-i\omega r_*}| \lesssim_{\ell} e^{2\kappa_{-}r_*} \sim (r - r_-), \quad (5.26)$$

$$|v_2 - e^{i\omega r_*}| \lesssim_{\ell} e^{2\kappa_{-}r_*} \sim (r - r_-). \quad (5.27)$$

for $r_* \geq 0$. Similarly to (5.24), we can prove

$$\frac{r_+}{r_-} \beta_{\mathcal{CH}_A}^{\ell,m}(\omega) = b_{\mathcal{CH}_A}^{\ell,m}(\omega) \text{ and } \frac{r_+}{r_-} \beta_{\mathcal{CH}_B}^{\ell,m}(\omega) = b_{\mathcal{CH}_B}^{\ell,m}(\omega). \quad (5.28)$$

Moreover, from the uniform boundedness of the reflection and transmission coefficients (cf. Theorem 2), we have the estimate

$$\begin{aligned} |\beta_{\mathcal{CH}_A}^{\ell,m}(\omega)| + |\beta_{\mathcal{CH}_B}^{\ell,m}(\omega)| &= \frac{r_+}{r_-} |\beta_{\mathcal{CH}_A}^{\ell,m}(\omega)| + \frac{r_+}{r_-} |\beta_{\mathcal{CH}_B}^{\ell,m}(\omega)| \\ &= \frac{r_+}{r_-} \left(|\Re \alpha_{\mathcal{H}_A}^{\ell,m} + \bar{\Im} \alpha_{\mathcal{H}_B}^{\ell,m}| + |\bar{\Re} \alpha_{\mathcal{H}_B}^{\ell,m} + \Im \alpha_{\mathcal{H}_A}^{\ell,m}| \right) \\ &\leq C(|\alpha_{\mathcal{H}_A}^{\ell,m}(\omega)| + |\alpha_{\mathcal{H}_B}^{\ell,m}(\omega)|) = C(|a_{\mathcal{H}_A}^{\ell,m}(\omega)| + |a_{\mathcal{H}_B}^{\ell,m}(\omega)|) \end{aligned} \quad (5.29)$$

for a constant C which only depends on the black hole parameters. Here, we have used the fact that

$$\begin{pmatrix} \beta_{\mathcal{CH}_B}^{\ell,m} \\ \beta_{\mathcal{CH}_A}^{\ell,m} \end{pmatrix} = \begin{pmatrix} \Im & \bar{\Re} \\ \Re & \bar{\Im} \end{pmatrix} \begin{pmatrix} \alpha_{\mathcal{H}_A}^{\ell,m} \\ \alpha_{\mathcal{H}_B}^{\ell,m} \end{pmatrix}. \quad (5.30)$$

In view of $1 = |\Im|^2 - |\Re|^2$, we also have

$$\begin{pmatrix} \alpha_{\mathcal{H}_A}^{\ell,m} \\ \alpha_{\mathcal{H}_B}^{\ell,m} \end{pmatrix} = \begin{pmatrix} \bar{\Im} & -\bar{\Re} \\ -\Re & \Im \end{pmatrix} \begin{pmatrix} \beta_{\mathcal{CH}_B}^{\ell,m} \\ \beta_{\mathcal{CH}_A}^{\ell,m} \end{pmatrix} \quad (5.31)$$

from which we deduce

$$|a_{\mathcal{H}_A}^{\ell,m}(\omega)| + |a_{\mathcal{H}_B}^{\ell,m}(\omega)| \lesssim |b_{\mathcal{CH}_A}^{\ell,m}(\omega)| + |b_{\mathcal{CH}_B}^{\ell,m}(\omega)|. \quad (5.32)$$

Estimate (5.29) and (5.32) show the claim in view of (5.8), (5.9), (5.13), and (5.14). Finally, in view of the Fourier inversion theorem, note that the previous also justifies the Fourier representation of scattering map (3.20), and the Fourier representations (3.24) and (3.25). \square

5.3. Completing the Proof

Having proven Lemma 5.1 and Proposition 5.1, we can finally show Theorem 1 in the following.

Proof of Theorem 1. Since $\mathcal{D}_{\mathcal{H}}^T \subset \mathcal{E}_{\mathcal{H}}^T$ is dense (Lemma 5.1) and $S^T: \mathcal{D}_{\mathcal{H}}^T \subset \mathcal{E}_{\mathcal{H}}^T \rightarrow \mathcal{D}_{\mathcal{CH}}^T \subset \mathcal{E}_{\mathcal{CH}}^T$ is a bounded injective map (Remark 3.2, Proposition 5.1), we can uniquely extend S_0^T to the bounded injective scattering map

$$S^T: \mathcal{E}_{\mathcal{H}}^T \rightarrow \mathcal{E}_{\mathcal{CH}}^T. \quad (5.33)$$

Analogously, in view of Proposition 2.2, Remark 3.1–3.2, and Proposition 5.1, we can uniquely extend the bounded injective map $B_0^T: \mathcal{D}_{\mathcal{CH}}^T \subset \mathcal{E}_{\mathcal{CH}}^T \rightarrow \mathcal{D}_{\mathcal{CH}}^T \subset \mathcal{E}_{\mathcal{CH}}^T$ to the bounded injective backward map $B^T: \mathcal{E}_{\mathcal{CH}}^T \rightarrow \mathcal{E}_{\mathcal{H}}^T$ (Lemma 5.1).

Since $B_0^T \circ S_0^T = \text{Id}_{\mathcal{D}_{\mathcal{H}}^T}$ and $S_0^T \circ B_0^T = \text{Id}_{\mathcal{D}_{\mathcal{CH}}^T}$ on dense sets, it also extends to $\mathcal{E}_{\mathcal{H}}^T$ and $\mathcal{E}_{\mathcal{CH}}^T$ from which (3.5) follows. Similarly, it suffices to check (3.6) for $\psi \in \mathcal{D}_{\mathcal{H}}^T$. Indeed, (3.6) holds true for $\psi \in \mathcal{D}_{\mathcal{H}}^T$ in view of the T energy identity. \square

6. Proof of Theorem 6. Breakdown of T Energy Scattering for Cosmological Constants $\Lambda \neq 0$

In the presence of a cosmological constant Λ , the situation regarding the T energy scattering problem is changed radically. In this section we will consider the subextremal (anti-) de Sitter–Reissner–Nordström black hole interior $(\mathcal{M}_{\text{int}}, g_{Q,M})$ which is completely analogous to $(\mathcal{M}_{\text{RN}}, g_{Q,M})$. We will assume that $(\mathcal{M}, Q, \Lambda) \in \mathcal{P}_{\text{se}}$ as defined in Sect. 3.6. Also, recall that in the presence of a cosmological constant it is natural to look at the Klein–Gordon equation

$$\square_g \psi - \mu \psi = 0 \quad (6.1)$$

with mass $\mu = \frac{3}{2}\Lambda$ for the conformal invariant equation or more general $\mu = \nu\Lambda$ for fixed $\nu \in \mathbb{R}$.

This section is devoted to prove Theorem 6 which relies on the fact that solutions of the corresponding radial o.d.e. in the vanishing frequency limit $\omega = 0$ generically map bounded solutions at $r_* = -\infty$ to unbounded solutions at $r_* = +\infty$. More precisely, for $\Lambda \neq 0$ we obtain—after separation of variables for (6.1) and setting $dr_* = h^{-1}dr$ —the o.d.e.

$$-u'' + V_{\ell,\Lambda} u = \omega^2 u \quad (6.2)$$

for $u(r_*) = r(r_*)R(r_*)$, where

$$V_{\ell,\Lambda} = h \left(\frac{hh'}{r} + \frac{\ell(\ell+1)}{r^2} - \mu \right) = h \left(\frac{\frac{dh}{dr}}{r} + \frac{\ell(\ell+1)}{r^2} - \mu \right) \quad (6.3)$$

and

$$h = \frac{\Delta}{r^2} = 1 - \frac{2M}{r} - \frac{1}{3}\Lambda r^2 + \frac{Q^2}{r^2}. \quad (6.4)$$

Here, consider $r(r_*)$ as a function r_* and recall that $'$ denotes the derivative with respect to r_* . The presence of the mass and the cosmological constant leads to a modification of the potential $V_{\ell,\Lambda}$.

Nevertheless, the potential $V_{\ell,\Lambda}$ still decays exponentially at $\pm\infty$ and we can define asymptotic states $u_1^{(\Lambda)}, u_2^{(\Lambda)}$, and $v_1^{(\Lambda)}, v_2^{(\Lambda)}$ for $\omega \neq 0$ and $\tilde{u}_1^{(\Lambda)}, \tilde{u}_2^{(\Lambda)}$, and $\tilde{v}_1^{(\Lambda)}, \tilde{v}_2^{(\Lambda)}$ for $\omega = 0$ just as in the case where $\Lambda = \mu = 0$ in Definition 2.3. In particular, $\tilde{u}_1^{(\Lambda)}$ and $\tilde{v}_1^{(\Lambda)}$ remain bounded as $r_* \rightarrow -\infty$ and $r_* \rightarrow +\infty$, respectively. In contrast to that, $\tilde{u}_2^{(\Lambda)}$ and $\tilde{v}_2^{(\Lambda)}$ grow linearly in their respective limits. The next proposition states that in the presence of a cosmological constant, solutions to (6.1) in the case $\omega = 0$ which are bounded at $r_* = -\infty$ do not need to be bounded at $r_* = +\infty$.

Proposition 6.1. *Fix $\nu \in \mathbb{R}$ (e.g., $\nu = \frac{3}{2}$ for the conformal invariant mass) and fix subextremal black hole parameters $(M, Q, \Lambda) \in \mathcal{P}_{\text{se}}$. Assume moreover that $(M, Q, \Lambda) \notin D(\nu)$, where $D(\nu) \subset \mathcal{P}_{\text{se}}$ is defined in the proof and has measure zero. Then, there exists an $\ell_0 = \ell_0(\nu) \in \mathbb{N}_0$ such that we have*

$$\tilde{u}_1^{(\Lambda)} = A(\ell, \Lambda, M, Q) \tilde{v}_1^{(\Lambda)} + B(\ell, \Lambda, M, Q) \tilde{v}_2^{(\Lambda)}, \quad (6.5)$$

with $B = B(\ell_0, \Lambda, M, Q) \neq 0$. Moreover, $\mathcal{P}_{\text{se}}^{\Lambda=0} \subset D(\nu)$ for all $\nu \in \mathbb{R}$ and there exists an open subset U in $\mathcal{P}_{\text{se}}^{\Lambda=0} \subset U \subset \mathcal{P}_{\text{se}}$ and $\mathcal{P}_{\text{se}} \cap U = \mathcal{P}_{\text{se}}^{\Lambda=0}$.

Proof. Let $\nu \in \mathbb{R}$ be fixed. In the case $\Lambda = 0$, we can represent \tilde{u}_1 with Legendre polynomials and in particular we have that $B(\ell, \Lambda = 0, M, Q) = 0$ for all ℓ and $0 < |Q| < M$. Note that we can write B as

$$B(\Lambda, \ell, M, Q) = \frac{\mathfrak{W}(\tilde{v}_2^{(\Lambda)}, \tilde{u}_1^{(\Lambda)})}{\mathfrak{W}(\tilde{v}_1^{(\Lambda)}, \tilde{v}_2^{(\Lambda)})} = \mathfrak{W}(\tilde{v}_2^{(\Lambda)}, \tilde{u}_1^{(\Lambda)}) \quad (6.6)$$

for all Λ such that $(M, Q, \Lambda) \in \mathcal{P}_{\text{se}}$.

Step 1: $\mathcal{P}_{\text{se}} \subset \mathbb{R}^3$ is open and has two connected components where either $Q > 0$ or $Q < 0$. For the sake of completeness, we will give a proof of Step 1, although this seems a quite well-known fact. Note that $\mathcal{P}_{\text{se}} = \mathcal{P}_{\text{se}}^{\Lambda>0} \cup \mathcal{P}_{\text{se}}^{\Lambda<0} \cup \mathcal{P}_{\text{se}}^{\Lambda=0}$ is open which can be inferred from its definition.

For the second statement, first note that $\{Q = 0\} \cap \mathcal{P}_{\text{se}} = \emptyset$. We will now show that $\{Q > 0\} \cap \mathcal{P}_{\text{se}}$ is connected. In Proposition A.3 in Appendix, we show that $\mathcal{P}_{\text{se}}^{\Lambda>0} \cap \{Q > 0\}$ and $\mathcal{P}_{\text{se}}^{\Lambda<0} \cap \{Q > 0\}$ are path-connected. To conclude, note that for every $(M_0, Q_0, \Lambda_0 = 0) \in \mathcal{P}_{\text{se}}^{\Lambda=0}$, there exist paths from (M_0, Q_0, Λ_0) to both $(M_0, Q_0, \epsilon) \in \mathcal{P}_{\text{se}}^{\Lambda>0}$ and $(M_0, Q_0, -\epsilon) \in \mathcal{P}_{\text{se}}^{\Lambda<0}$ for

some $\epsilon(M_0, Q_0) > 0$. Together with the fact that $\mathcal{P}_{\text{se}}^{\Lambda=0} \cap \{Q > 0\}$ is path-connected, this shows that $\{Q > 0\} \cap \mathcal{P}_{\text{se}}$ is path-connected and similarly that $\{Q < 0\} \cap \mathcal{P}_{\text{se}}$ is path-connected which proves the claim.

Step 2: $\mathcal{P}_{\text{se}} \ni (M, Q, \Lambda) \mapsto B(\ell, \Lambda, M, Q)$ is real analytic To show *Step 2* we first express (6.5) in r coordinates. Note that for $(M, Q, \Lambda) \in \mathcal{P}_{\text{se}}$ equation (6.5) is equivalent to

$$\frac{r_+}{r_-}(-1)^\ell P_\ell^{(\Lambda)}(x(r)) = A(\ell, \Lambda) \tilde{P}_\ell^{(\Lambda)}(x(r)) + B(\ell, \Lambda) \tilde{Q}_\ell^{(\Lambda)}(x(r)), \quad (6.7)$$

where $r \in (r_-, r_+)$,

$$x(r) := -\frac{2r}{r_+ - r_-} + \frac{r_+ + r_-}{r_+ - r_-}, \quad (6.8)$$

$$r(x) = -\frac{r_+ - r_-}{2}x + \frac{r_+ + r_-}{2} \quad (6.9)$$

and $0 < r_- < r_+$. Now, note that $\mathcal{P}_{\text{se}} \ni (M, Q, \Lambda) \mapsto r_-$ and $\mathcal{P}_{\text{se}} \ni (M, Q, \Lambda) \mapsto r_+$ are real analytic. Moreover, we can write $\Delta = (r - r_-)(r - r_+)p(r)$ for a second order polynomial $p(r)$, where $\mathcal{P}_{\text{se}} \ni \Lambda \mapsto p(r)$ is also real analytic for fixed r . Now, $P_\ell^{(\Lambda)}$, $\tilde{P}_\ell^{(\Lambda)}$, and $\tilde{Q}_\ell^{(\Lambda)}$ appearing in (6.7) are defined as the unique solutions of

$$\frac{d}{dx} \left((1 - x^2)p(r(x)) \frac{dR}{dx} \right) + \gamma(\ell + 1, \Delta - r(x)^2\nu\Lambda R) = 0 \quad (6.10)$$

satisfying

$$P_\ell^{(\Lambda)} = (-1)^\ell + O_{\ell, \nu}(x) \text{ as } x \rightarrow -1, \quad (6.11)$$

$$\frac{dP_\ell^{(\Lambda)}}{dx} = O_\ell(1) \text{ as } x \rightarrow -1, \quad (6.12)$$

$$\tilde{P}_\ell^{(\Lambda)} = -1 + \gamma_\nu(1 - x) \text{ as } x \rightarrow 1, \quad (6.13)$$

$$\frac{d\tilde{P}_\ell^{(\Lambda)}}{dx} = \gamma_\ell(1) \text{ as } x \rightarrow 1, \quad (6.14)$$

$$\tilde{Q}_\ell^{(\Lambda)} = -\frac{1}{2} \log(1 - x) + O_\ell(1) \text{ as } x \rightarrow 1, \quad (6.15)$$

$$\frac{d\tilde{Q}_\ell^{(\Lambda)}}{dx} = \frac{1}{2(1 - x)} + O_\ell((1 - x) \log(1 - x)) \text{ as } x \rightarrow 1. \quad (6.16)$$

Note that (6.10) depends real analytically on $(M, Q, \Lambda) \in \mathcal{P}_{\text{se}}$ such that $P_\ell^{(\Lambda)}(x)$, $\tilde{P}_\ell^{(\Lambda)}(x)$, $\tilde{Q}_\ell^{(\Lambda)}(x)$ are real analytic functions of $(M, Q, \Lambda) \in \mathcal{P}_{\text{se}}$ for $x \in (-1, 1)$. Hence, $\mathcal{P}_{\text{se}} \ni (M, Q, \Lambda) \mapsto B(\ell, \Lambda, M, Q)$ is real analytic.

Step 3: $B(\ell_0(\nu), \Lambda, M, Q)$ only vanishes on a set $D(\nu) \subset \mathcal{P}_{\text{se}}$ of measure zero. The claim follows from

$$\frac{\partial B(\ell, \Lambda, M_0, Q_0)}{\partial \Lambda} \Big|_{\Lambda=0} \neq 0 \quad (6.17)$$

for some $0 < |Q_0| < M_0$. Throughout *Step 2*, we fix $0 < |Q_0| < M_0$ and avoid writing their explicit dependence. First note that for $\Lambda = 0$ we obtain

the Legendre functions of first and second kind, i.e., $P_\ell^{(0)} = \tilde{P}_\ell^{(0)} = P_\ell$ and $\tilde{Q}_\ell^{(0)} = Q_\ell$ and $B(0, \ell) = 0$. Now, define coefficients $\tilde{A}(\ell, \Lambda)$ and $\tilde{B}(\ell, \Lambda)$ to satisfy

$$P_\ell^{(\Lambda)} = \tilde{A}(\ell, \Lambda) \tilde{P}_\ell^{(\Lambda)} + \tilde{B}(\ell, \Lambda) \tilde{Q}_\ell^{(\Lambda)}, \quad (6.18)$$

and note that (6.17) is equivalent (use that $B(\ell, 0) = \tilde{B}(\ell, 0) = 0$) to

$$\frac{\partial \tilde{B}(\ell, \Lambda)}{\partial \Lambda} \Big|_{\Lambda=0} \neq 0. \quad (6.19)$$

By construction, $P_\ell^{(\Lambda)}$ solves (6.10). Multiplying

$$\frac{d}{dx} \left((1 - x^2) p(r(x)) \frac{dP_\ell^{(\Lambda)}}{dx} \right) + \ell(\ell + 1) P_\ell^{(\Lambda)} - r(x)^2 \nu \Lambda P_\ell^{(\Lambda)} = 0 \quad (6.20)$$

by $P_\ell^{(0)}$ and integrating from $x = -1$ to $x = 1$ yields

$$0 = \int_{-1}^1 P_\ell^{(0)} \left(\frac{d}{dx} \left((1 - x^2) p(r(x)) \frac{dP_\ell^{(0)}}{dx} \right) + \ell(\ell + 1) P_\ell^{(0)} - r(x)^2 \nu \Lambda P_\ell^{(0)} \right) dx. \quad (6.21)$$

Using the expansion (6.18) and the properties (6.11)–(6.16) at the end points $x = -1$ and $x = 1$ gives after an integration by parts

$$0 = \int_{-1}^1 P_\ell^{(0)} \left(\frac{d}{dx} \left((1 - x^2) p(r(x)) \frac{dP_\ell^{(0)}}{dx} \right) + \ell(\ell + 1) P_\ell^{(0)} - r(x)^2 \nu \Lambda P_\ell^{(0)} \right) dx + p(r(1)) \tilde{B}(\ell, \Lambda). \quad (6.22)$$

Now, taking $\partial_\Lambda|_{\Lambda=0}$ and integrating by parts once again yields

$$\begin{aligned} & p(r(1)) \partial_\Lambda|_{\Lambda=0} \tilde{B}(\ell, \Lambda) \\ &= \int_{-1}^1 \left[\left| \frac{dP_\ell^{(0)}}{dx} \right|^2 (1 - x^2) \partial_\Lambda|_{\Lambda=0} (p(r(x))) + \left| P_\ell^{(0)} \right|^2 \partial_\Lambda|_{\Lambda=0} (\nu r(x)^2 \Lambda) \right] dx \\ &= \int_{-1}^1 \left[\left| \frac{dP_\ell^{(0)}}{dx} \right|^2 (1 - x^2) \partial_\Lambda|_{\Lambda=0} (p(r(x))) + \nu \left| P_\ell^{(0)} \right|^2 r(x)^2|_{\Lambda=0} \right] dx. \end{aligned} \quad (6.23)$$

Recall that we are in the subextremal range which guarantees that $p(r(1)) \neq 0$. We will now distinguish two cases, $\nu = 0$ and $\nu \neq 0$.

Part I: $\nu = 0$. In the case $\nu = 0$, we have

$$p(r(1)) \partial_\Lambda|_{\Lambda=0} \tilde{B}(\ell, \Lambda) = \partial_\Lambda|_{\Lambda=0} \int_{-1}^1 \left| \frac{dP_\ell}{dx} \right|^2 (1 - x^2) p(r(x)) dx \quad (6.24)$$

In the case $\nu = 0$, we will choose $\ell = 1$ such that

$$\begin{aligned}
& p(r(1))\partial_\Lambda|_{\Lambda=0}\tilde{B}(1, \Lambda) \\
&= \partial_\Lambda|_{\Lambda=0} \int_{-1}^1 (1-x^2)p(r(x))dx \\
&= \partial_\Lambda|_{\Lambda=0} \int_{-1}^1 -\Delta(r(x)) \frac{4}{(r_+ - r_-)^2} dx \\
&= x\partial_\Lambda|_{\Lambda=0} \left(\frac{-8}{(r_+ - r_-)^3} \int_{r_-}^{r_+} \Delta(r) dr \right) \\
&= -8\partial_\Lambda|_{\Lambda=0} \left(\frac{\frac{r_+^3 - r_-^3}{3} - M_0(r_+^2 - r_-^2) + Q_0^2(r_+ - r_-) - \frac{1}{15}\Lambda(r_+^5 - r_-^5)}{(r_+ - r_-)^3} \right) \\
&= \frac{8(r_+^5 - r_-^5)}{15(r_+ - r_-)^3} \Big|_{\Lambda=0} \\
&\quad + 8 \frac{\frac{r_+^3 - r_-^3}{3} - M_0(r_+^2 - r_-^2) + Q_0^2(r_+ - r_-)}{(r_+ - r_-)^5} (r_+^4 - r_-^4) \Big|_{\Lambda=0} \\
&\quad - \frac{8}{3} \frac{r_+^6 + r_-^6 - 2M_0(r_+^5 + r_-^5) + Q_0^2(r_+^4 + r_-^4)}{(r_+ - r_-)^4} \Big|_{\Lambda=0} \\
&= \frac{-8}{15} (3r_+^3 + 3r_-^2 + 4r_+ r_-) \Big|_{\Lambda=0} \\
&= x \frac{-8}{15} (6M_0^2 - Q_0^2) < -24M_0^2
\end{aligned}$$

The last step is a long but direct computation using that $\Delta = r^2 - 2M_0r + Q_0^2 - \frac{\Lambda}{3}r^4$ and $r_\pm|_{\Lambda=0} = M_0 \pm \sqrt{M_0^2 - Q_0^2}$, i.e., $Q_0^2 = r_+ r_-|_{\Lambda=0}$ and $2M_0 = r_+|_{\Lambda=0} + r_-|_{\Lambda=0}$. Moreover, in view of the inverse function theorem we have

$$\partial_\Lambda|_{\Lambda=0} r_+ = \frac{r_+^4}{3(r_+ - r_-)} \Big|_{\Lambda=0} \quad (6.25)$$

and

$$\partial_\Lambda|_{\Lambda=0} r_- = -\frac{r_-^4}{3(r_+ - r_-)} \Big|_{\Lambda=0}. \quad (6.26)$$

Let $\Lambda \cdot \nu = 0$. In this case, we choose $\ell = 0$ such that $P_\ell^{(0)} = 1$ and $\frac{d}{dr} \frac{\ell^{(0)}}{r} = 0$. Hence,

$$\begin{aligned}
p(r(1))\partial_\Lambda|_{\Lambda=0}\tilde{B}(\ell, \Lambda) &= \partial_\Lambda|_{\Lambda=0} \int_{-1}^1 r(x)^2 \nu \Lambda dx \\
&= \nu \partial_\Lambda|_{\Lambda=0} \int_{-1}^1 \left(-\frac{r_+ - r_-}{2} x + \frac{r_+ + r_-}{2} \right)^2 \Lambda dx \\
&= \nu \left(\frac{1}{6}(r_+ - r_-)^2 + \frac{1}{2}(r_+ + r_-)^2 \right) \Big|_{\Lambda=0} \neq 0. \quad (6.27)
\end{aligned}$$

This shows that $\mathcal{P}_{\text{se}} \ni (M, Q, \Lambda) \mapsto B(\ell_0(\nu), M, Q, \Lambda)$ is a non-trivial real analytic function which zero set $D(\nu)$ has zero measure. The proof also shows that $\mathcal{P}_{\text{se}}^{\Lambda=0} \subset D(\nu)$ and that there exists an open set $U \subset \mathcal{P}_{\text{se}}$ with $\mathcal{P}_{\text{se}}^{\Lambda=0} \subset U$ and $D(\nu) \cap U = \mathcal{P}_{\text{se}}^{\Lambda=0}$. \square

Proposition 6.2. *Let $\nu \in \mathbb{R}$ be fixed. Let $\omega \neq 0$, $(M, Q, \Lambda) \in \mathcal{P}_{\text{se}}$, and $\ell \in \mathbb{N}_0$. Then, define completely analogously to Definition 2.4 transmission and reflection coefficients $\mathfrak{T}(\omega, \ell, \Lambda)$ and $\mathfrak{R}(\omega, \ell, \Lambda)$ as the unique coefficients such that*

$$u_1^{(\Lambda)} = \mathfrak{T}(\omega, \ell, \Lambda) v_1^{(\Lambda)} + \mathfrak{R}(\omega, \ell, \Lambda) v_2^{(\Lambda)} \quad (6.28)$$

holds.

Now, assume further that $(M, Q, \Lambda) \in \mathcal{P}_{\text{se}} \setminus D(\nu)$, where $D(\nu)$ is defined in Proposition 6.1. Then, there exists an $\ell_0 = \ell_0(\nu)$ such that

$$\lim_{\omega \rightarrow 0} |\mathfrak{R}(\omega, \ell_0)| = \lim_{\omega \rightarrow 0} |\mathfrak{T}(\omega, \ell_0)| = +\infty. \quad (6.29)$$

This shows that \mathfrak{T} and \mathfrak{R} have a simple pole at $\omega = 0$.

Proof. Fix $\ell_0 = \ell_0(\nu)$ from Proposition 6.1 and $(M, Q, \Lambda) \in \mathcal{P}_{\text{se}}$ such that $B(\ell_0, \Lambda, M, Q) \neq 0$. Now, note that the o.d.e. implies that $\frac{d}{d\omega} \ln(\bar{u}u') = 0$ which shows that $1 = |\mathfrak{T}|^2 - |\mathfrak{R}|^2$. In particular, either $|\mathfrak{T}|$ and $|\mathfrak{R}|$ are both bounded or both unbounded as $\omega \rightarrow 0$. Also note that as $\omega \rightarrow 0$, we have that $u_1^{(\Lambda)} \rightarrow \tilde{u}_1^{(\Lambda)}$ pointwise.

Now, assume for a contradiction that there exists a sequence $\omega_n \rightarrow 0$ such that $|\mathfrak{T}(\omega_n)|$ and $|\mathfrak{R}(\omega_n)|$ remain bounded. Thus,

$$\begin{aligned} \limsup_{\omega_n \rightarrow 0} \|u_1^{(\Lambda)}\|_{L^\infty(\mathbb{R})} &\leq \limsup_{\omega_n \rightarrow 0} \|u^{(\Lambda)}\|_{L^\infty((-r_0^*, r_0^*))} \\ &\quad + \limsup_{\omega_n \rightarrow 0} \|\mathfrak{R}v_1^{(\Lambda)} + \mathfrak{T}v_2^{(\Lambda)}\|_{L^\infty((0, \infty))} \leq C \end{aligned} \quad (6.30)$$

for some constant $C > 0$. Now, using that $B(\ell_0, \Lambda, M, Q) \neq 0$ in Proposition 6.1, we can choose a $r_0^* \in \mathbb{R}$ such that $|\tilde{u}_1^{(\Lambda)}(r_0^*)| > C$ which contradicts the fact that $u_1^{(\Lambda)} \rightarrow \tilde{u}_1^{(\Lambda)}$ pointwise as $\omega_n \rightarrow 0$. \square

Finally, this allows us to prove Theorem 6 which we restate in the following for the convenience of the reader.

Theorem 6. *Let $\nu \in \mathbb{R}$ be a fixed Klein–Gordon mass parameter. (In particular, we may choose $\nu = \frac{1}{2}$ to cover the conformal invariant case or $\nu = 0$ for the wave equation (1.1).) Consider the interior of a subextremal (anti-) de Sitter–Reissner–Nordström black hole with generic parameters $(M, Q, \Lambda) \in \mathcal{P}_{\text{se}} \setminus D(\nu)$. (Here, $D(\nu) \subset \mathcal{P}_{\text{se}}$ is a set with measure zero defined in Proposition 6.1 (see Sec. 6.1).) Moreover $D(\nu)$ satisfies $\mathcal{P}_{\text{se}}^{\Lambda=0} \subset D(\nu)$ and $U \cap D(\nu) = \mathcal{P}_{\text{se}}^{\Lambda=0}$ for some open set $U \subset \mathcal{P}_{\text{se}}$.)*

Then, there exists a sequence $(\Psi_n)_{n \in \mathbb{N}}$ of purely ingoing and compactly supported data on \mathcal{H}_A with

$$\|\Psi_n\|_{\mathcal{E}_{\mathcal{H}}^T} = 1 \text{ for all } n \quad (3.41)$$

such that the solution ψ_n to the Klein–Gordon equation with mass $\mu = \nu\Lambda$

$$\square_{g_{M, Q, \Lambda}} \psi - \mu \psi = 0 \quad (3.42)$$

arising from Ψ_n has unbounded T energy at the Cauchy horizon

$$\|\psi_n|_{\mathcal{CH}}\|_{\mathcal{E}_{\mathcal{CH}}^T} \rightarrow \infty \text{ as } n \rightarrow \infty. \quad (3.43)$$

Proof. Fix $\ell_0 = \ell_0(\nu)$ from Proposition 6.2 such that the reflection and transmission coefficients blow up as $\omega \rightarrow 0$. Define a sequence of compactly supported functions Ψ_n on \mathcal{H}_A by $\Psi_n(v, \theta, \varphi) = f_n(v)Y_{0\ell}(\theta, \varphi)$, such that $f_n \in C_c^\infty(\mathbb{R})$,

$$\int_{\mathbb{R}} \omega^2 |\hat{f}_n(\omega)|^2 d\omega = 1 \text{ and } \int_{-\frac{1}{n}}^{\frac{1}{n}} \omega^2 |\hat{f}_n(\omega)|^2 d\omega \geq \epsilon \int_{\mathbb{R}} \omega^2 |\hat{f}_n(\omega)|^2 d\omega = \epsilon \quad (6.31)$$

for some $\epsilon > 0$.⁶ Imposing vanishing data on \mathcal{H}_B , this gives rise to a unique smooth solutions ψ_n up to but excluding the Cauchy horizon. Arguments completely analogous to those given in the proof of Proposition 5.1 show that

$$\|\psi_n|_{\mathcal{CH}}\|_{\mathcal{E}_{\mathcal{CH}}^T}^2 = \frac{r_+^2}{r_-^2} \int_{\mathbb{R}} \omega^2 (|\mathfrak{R}(\omega, \ell)|^2 + |\mathfrak{T}(\omega, \ell)|^2) |\hat{f}_n(\omega)|^2 d\omega. \quad (6.33)$$

Thus,

$$\begin{aligned} \|\psi_n|_{\mathcal{CH}}\|_{\mathcal{E}_{\mathcal{CH}}^T}^2 &\geq \frac{r_+^2}{r_-^2} \int_{-\frac{1}{n}}^{\frac{1}{n}} \omega^2 (|\mathfrak{R}(\omega, \ell)|^2 + |\mathfrak{T}(\omega, \ell)|^2) |\hat{f}_n(\omega)|^2 d\omega \\ &\geq \epsilon \frac{r_+^2}{r_-^2} \inf_{\omega \in [-\frac{1}{n}, \frac{1}{n}]} (|\mathfrak{R}|^2 + |\mathfrak{T}|^2). \end{aligned} \quad (6.34)$$

Since $|\mathfrak{R}|, |\mathfrak{T}| \rightarrow \infty$ as $\omega \rightarrow 0$, also $\inf_{\omega \in [-\frac{1}{n}, \frac{1}{n}]} |\mathfrak{R}| \rightarrow \infty$ and $\inf_{\omega \in [-\frac{1}{n}, \frac{1}{n}]} |\mathfrak{T}| \rightarrow \infty$ as $n \rightarrow \infty$. Thus, as $n \rightarrow \infty$, we have

$$\|\psi_n|_{\mathcal{CH}}\|_{\mathcal{E}_{\mathcal{CH}}^T}^2 \rightarrow \infty. \quad (6.35)$$

□

7. Proof of Theorem 7: Breakdown of T Energy Scattering for the Klein–Gordon Equation

In this last section, we will prove that for a generic set of Klein–Gordon masses, there does not exist a T scattering theory on the interior of Reissner–Nordström for the Klein–Gordon equation. For the convenience of the reader, we have restated Theorem 7.

Theorem 7. *Consider the interior of a subextremal Reissner–Nordström black hole. There exists a discrete set $\tilde{D}(M, Q) \subset \mathbb{R}$ with $0 \in \tilde{D}$ such that the following holds true. For any $\mu \in \mathbb{R} \setminus \tilde{D}$, there exists a sequence $(\Psi_n)_{n \in \mathbb{N}}$ of purely ingoing and compactly supported data on \mathcal{H}_A with*

$$\|\Psi_n\|_{\mathcal{E}_A^T} = 1 \text{ for all } n \quad (3.44)$$

⁶ Such a function can be constructed by setting $f_n(v) := \frac{c}{\sqrt{n}} f(\frac{v}{n})$ for smooth $f: \mathbb{R} \rightarrow [0, 1]$ with $\text{supp}(f) \subset [-2, 2]$, $f|_{[-1, 1]} = 1$ and some normalization constant $c > 0$. Indeed,

$$\int_{-\frac{1}{n}}^{\frac{1}{n}} \omega^2 |\hat{f}_n(\omega)|^2 d\omega = \int_{-\frac{1}{n}}^{\frac{1}{n}} \omega^2 |\sqrt{n} \hat{f}(n\omega)|^2 d\omega = \int_{-1}^1 \omega^2 |\hat{f}(\omega)|^2 d\omega =: \epsilon > 0 \quad (6.32)$$

in view of $\hat{f}(0) = \int_{\mathbb{R}} f(v) dv > 0$.

such that the solution ψ_n to the Klein–Gordon equation with mass μ

$$\square_{g_{M,Q,\Lambda}} \psi - \mu \psi = 0 \quad (3.45)$$

arising from Ψ_n has unbounded T energy at the Cauchy horizon

$$\|\psi_n|_{\mathcal{CH}}\|_{\mathcal{E}_{\mathcal{CH}}^T} \rightarrow \infty \text{ as } n \rightarrow \infty. \quad (3.46)$$

Proof. The proof of this statement is easier than and similar to the proof of Theorem 6 and the proofs of the propositions leading up to it. More precisely, similar to Sect. 6 we define asymptotic states $\tilde{u}_1^{(\mu)}$, $\tilde{v}_1^{(\mu)}$ and $\tilde{v}_2^{(\mu)}$ and define $A(\ell, \mu)$ and $B(\ell, \mu)$ by $\tilde{u}_1^{(\mu)} = A(\ell, \mu)\tilde{v}_1^{(\mu)} + B(\ell, \mu)\tilde{v}_2^{(\mu)}$. As in Sect. 6, $\mathbb{R} \ni r_* \mapsto B(\ell, \mu)$ is real analytic and from the o.d.e. $-u'' + V_{\ell, \mu}u = 0$ we obtain

$$\frac{\partial B(\ell, \mu)}{\partial \mu} \Big|_{\mu=0} = \int_{-\infty}^{\infty} \frac{\partial V_{\ell, \mu}}{\partial \mu} \Big|_{\mu=0} \tilde{u}_1^2 dr_*, \quad (7.1)$$

where

$$V_{\ell, \mu} = h \left(\frac{hh'}{r} + \frac{\ell(\ell+1)}{r^2} - \mu \right) = h \left(\frac{\frac{\alpha}{\alpha'}}{r} - \frac{\ell(\ell+1)}{r^2} - \mu \right) \quad (7.2)$$

and

$$h = 1 - \frac{\alpha'}{r} + \frac{\alpha'^2}{r^2} \quad (7.3)$$

as in (2.5). Now, note that

$$\frac{\partial V_{\ell, \mu}}{\partial r_*} \Big|_{\mu=0} = -h > 0 \quad (7.4)$$

which is manifestly positive from which we can infer, by analyticity, that $B(\ell, \mu) \neq 0$ for all $\mu \in \mathbb{R} \setminus \tilde{D}$, where $\tilde{D} = \tilde{D}(M, Q) \subset \mathbb{R}$ is a discrete set. This proves the analogous statements to Proposition 6.1 and Proposition 6.2. The claim of Theorem 7 follows now as in the proof of Theorem 6. \square

Acknowledgements

The authors would like express their gratitude to Mihalis Dafermos for many valuable discussions and helpful remarks. The authors also thank Igor Rodnianski, Jonathan Luk, and Sung-Jin Oh for useful conversations. CK Acknowledges support from the EPSRC and thanks Princeton University for hosting him as a VSRC. YS acknowledges support from the NSF Postdoctoral Research Fellowship under Award No. 1502569.

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Appendix A. Additional Lemmata

Energy Estimates in the Interior

Lemma A.1. *Let $\Psi \in C_c^\infty(\mathcal{H})$ and denote by ψ its evolution in the interior. Then, the non-degenerate N energy of Ψ decays exponentially toward i^+ in every $\{r = r_0\}$ hypersurface for $r_{red} < r_0 < r_+$. Here, r_{red} only depends on the black hole parameters.*

Proof. This argument is very similar to [17, Proposition 4.2]. We only prove it for the right component of i^+ and clearly only have to look at a neighborhood of i^+ . First, recall the existence of the celebrated redshift vector field N satisfying $K^N[\psi] \geq b J_\mu^N[\psi] n_v^\mu$ for $r_+ \geq r \geq r_{red}$, where n_v is the normal to a $v = \text{const.}$ hypersurface.⁷

We set

$$E(v_0) = \int_{v=v_0, r=r_{red}}^{v_0, r=r_+} J_\mu^N n_v^\mu \text{dvol}, \quad (\text{A.1})$$

and apply the energy identity with the redshift vector field N in the region $\mathcal{R} = \{r \in [r_{red}, r_+], v \in [v_0, v_1]\}$, where v_0 is large enough such that $v_0 > \sup \text{supp}(\Psi)$. This gives in view of the coarea formula that

$$E(v_1) - E(v_0) + \tilde{b} \int_{v_0}^{v_1} E(v) \text{d}v \leq 0 \quad (\text{A.2})$$

for every $v_1 \geq v_0 > \sup \text{supp}(\Psi)$. Inequality (A.2), smoothness of $v \mapsto E(v)$ and a further application of the energy identity in the region $\{v \geq v_0, r_+ \geq r \geq r_{red}\}$ finally show

$$\int_{v \geq v_0, r=r_{red}} J_\mu^N n_r^\mu \text{dvol} \leq C \exp(-\tilde{b} v_0), \quad (\text{A.3})$$

where C is a constant depending on Ψ . This concludes the proof. \square

Remark A.1. By cutting off smoothly, we can clearly approximate Ψ on a $\{r = \text{const.}\}$ hypersurface with compactly supported functions for any fixed $r \in (r_{red}, r_+)$.

Lemma A.2. *Let ψ be a smooth solution of the wave equation on \mathcal{M}_{RN} such that its restriction to the event horizon has compact support and let $r_0 \in (r_{red}, r_+)$. Then,*

$$\int_{\mathcal{H}} J_\mu^T n^\mu \text{dvol} \lesssim \int_{\{r=r_0\}} J_\mu^N n^\mu \text{dvol}. \quad (\text{A.4})$$

⁷The normal is fixed by making a choice of a volume form on the null hypersurface.

Proof. We shall use the vector field $S = r^{-2}\partial_{r_*}$. By potentially making r_{red} larger, we can assure that the bulk term $K^S := \nabla^\mu J_\mu^S$ of the vector field S has a fixed negative sign in $r_0 \in (r_{red}, r_+)$. This current is analogous to the current introduced in [17, par. 4.1.3.2]. Moreover, applying the energy identity in the region $\mathcal{R} = \{r_0 \leq r \leq r_+\}$ and noting that $J^N[\psi]_\mu n^\mu|_{r=r_0} \sim J^S[\psi]_\mu n^\mu|_{r=r_0}$ as well as $J^T[\psi]_\mu n^\mu|_{\mathcal{H}} \sim J^S[\psi]_\mu n^\mu|_{\mathcal{H}}$ yields

$$\int_{\{r=r_0\}} J^N[\psi]_\mu n^\mu d\text{vol} + \int_{\mathcal{R}} K^S d\text{vol} \gtrsim \int_{\mathcal{H}} J_\mu^T n^\mu d\text{vol}. \quad (\text{A.5})$$

This concludes the proof. \square

Analytic properties of the potential and the scattering coefficients. In the following, we would like to summarize analytic properties of the potential $V_\ell(r)$ and u_1, u_2, v_1 and v_2 as functions of ω . This is similar to parts of [5].

First, however, we will show the exponential decay of the potential V_ℓ as $r_* \rightarrow \pm\infty$.

Lemma A.3. *We have*

$$|\Delta(r_*)| \lesssim e^{2k_+ r_*} \text{ for } r_* \leq 0 \quad (\text{A.6})$$

and

$$|\Delta(r_*)| \lesssim e^{-2k_+ r_*} \text{ for } r_* \geq 0. \quad (\text{A.7})$$

Moreover, we have

$$|V_\ell(r_*)|, |V'_\ell(r_*)|, |V''_\ell(r_*)| \lesssim (1 + \ell(\ell + 1))e^{2k_+ r_*} \text{ for } r_* \leq 0 \quad (\text{A.8})$$

and

$$|V_\ell(r_*)|, |V'_\ell(r_*)|, |V''_\ell(r_*)| \lesssim (1 + \ell(\ell + 1))e^{2k_- r_*} \text{ for } r_* \geq 0. \quad (\text{A.9})$$

Proof. Note that

$$r_+ - r = \tilde{C}(r - r_-)^{\frac{k_-}{k_+}} e^{-2k_+ r} e^{2k_+ r_*} \quad (\text{A.10})$$

for a constant \tilde{C} only depending on the black hole parameters. Thus, for $r_* \leq 0$, we have

$$r_+ - r(r_*) = f(r_*)e^{2k_+ r_*} \quad (\text{A.11})$$

for a smooth function $f(r_*)$, which is uniformly bounded below and above for $r_* \leq 0$. Moreover, we have $f'(r_*), f''(r_*) \rightarrow 0$ exponentially fast as $r_* \rightarrow -\infty$. The estimates (A.8) and (A.9) are now straightforward applications of the chain rule and the fact that $\frac{dr}{dr_*} = \frac{\Delta}{r^2}$ and $\Delta = (r - r_-)(r - r_+)$. \square

Proposition A.1. *The potential V_ℓ can be expanded as*

$$V_\ell(r_*) = \sum_{m \in \mathbb{N}} C_m e^{2\kappa_+ m r_*}, \quad (\text{A.12})$$

where $|C_m| \lesssim_\ell e^{-\sigma m}$ for a $\sigma > 0$.

Proof. Define the variable

$$z(r) := e^{2\kappa_+ r_* (r)} = C e^{2\kappa_+ r} (r_+ - r)(r - r_-)^{\frac{\kappa_+}{\kappa_-}}, \quad (\text{A.13})$$

where $C > 0$ is such that $z(\frac{r_+ + r_-}{2}) = 1$. From the inverse function theorem, it follows that $V_\ell(z) = V_\ell(r(z))$ can be analytically continued in a neighborhood of $z = 0$ and thus, there exists a Taylor expansion around $z = 0$ such that

$$V_\ell(z) = \sum_{n=1}^{\infty} C_m z^m. \quad (\text{A.14})$$

Hence,

$$V_\ell(r_*) = \sum_{n=1}^{\infty} C_m e^{2\kappa_+ m r_*}, \quad (\text{A.15})$$

where

$$\begin{aligned} C_1 &= \frac{dV_\ell}{dz} \Big|_{z=0} = \frac{dV_\ell}{dr} \Big|_{r=r_+} \frac{dr}{dz} \Big|_{z=0} \\ &= \frac{r_+ - r_-}{r_+^4} (r_+^2 (r_+ - 3r_-) + \ell \cdot \rho + 1). \end{aligned} \quad (\text{A.16})$$

Note that the coefficients C_m decay exponentially fast in m . To see this, remark that we can redefine $\tilde{r}_* := r_* - \rho$ for some constant $\rho > 0$. Similarly to (A.15), we expand V_ℓ as

$$V_\ell = \sum_{m=1}^{\infty} D_m e^{2\kappa_+ m \tilde{r}_*} \quad (\text{A.17})$$

which shows $C_m = D_m e^{-2\kappa_+ m \rho}$. By analyticity, we have $|D_m| \leq |\tilde{C}|^{m+1}$ for some $\tilde{C} > 0$ and thus,

$$|C_m| \lesssim \ell e^{-\sigma m} \quad (\text{A.18})$$

for a fixed $\sigma > 0$. \square

Proposition A.2. *Let $\ell \in \mathbb{N}$ be fixed. Then,*

$$\sup_{\{|\operatorname{Re}(\omega)| > 1\}} |\Re(\omega, \ell)| + |\Im(\omega, \ell)| \lesssim \ell. \quad (\text{A.19})$$

Moreover, $\Im(\omega, \ell)$ has a pole of order one at $\omega = i\kappa_+$ given that $\ell(\ell + 1) \neq r_+^2(r_+ - 3r_-)$.

Proof. Recall, that u_1 is the unique solution to

$$u_1(r_*) = e^{i\omega r_*} + \int_{-\infty}^{r_*} \frac{\sin(\omega(r_* - y))}{\omega} V(y) u_1(y) dy. \quad (\text{A.20})$$

In [5], it is shown that the Volterra iteration has the form

$$u_1(r_*) = e^{i\omega r_*} \left(1 + \sum_{n=1}^{\infty} u_1^{(n)}(r_*) \right), \quad (\text{A.21})$$

where

$$u_1^{(n)}(r_*) = \sum_{\substack{m_n, \dots, m_1 \in \mathbb{N} \\ m_n > \dots > m_1}} C_{m_n - m_{n-1}} C_{m_{n-1} - m_{n-2}} \dots C_{m_1} d_{m_n} \dots d_{m_1} e^{2\kappa + m_n r_*} \quad (\text{A.22})$$

with $d_m = -(4m\kappa_+ (m\kappa_+ + i\omega))^{-1}$. Note that in view of the bound in (A.18) one can check that the Volterra iteration for u_1 converges on $\omega \in \mathbb{C} \setminus \{im\kappa_+ : m \in \mathbb{N}\}$ and moreover,

$$\sup_{\{|\operatorname{Re}(\omega)| > 1\}} |u_1(r_* = 0)| \lesssim_\ell 1, \quad (\text{A.23})$$

$$\sup_{\{|\operatorname{Re}(\omega)| > 1\}} |u_1'(r_* = 0)| \lesssim_\ell |\omega|. \quad (\text{A.24})$$

Analogously, we have that v_1 is analytic on $\omega \in \mathbb{C} \setminus \{im\kappa_- : m \in \mathbb{N}\}$ and v_2 is analytic on $\omega \in \mathbb{C} \setminus \{-im\kappa_- : m \in \mathbb{N}\}$. Moreover,

$$\sup_{\{|\operatorname{Re}(\omega)| > 1\}} |v_1(r_* = 0)| \lesssim_\ell 1. \quad (\text{A.25})$$

$$\sup_{\{|\operatorname{Re}(\omega)| > 1\}} |v_1'(r_* = 0)| \lesssim_\ell |\omega|. \quad (\text{A.26})$$

and

$$\sup_{\{|\operatorname{Re}(\omega)| > 1\}} |v_2(r_* = 0)| \lesssim_\ell 1, \quad (\text{A.27})$$

$$\sup_{\{|\operatorname{Re}(\omega)| > 1\}} |v_2'(r_* = 0)| \lesssim_\ell |\omega|. \quad (\text{A.28})$$

This finally shows (A.19) in view of the definition of the transmission and reflection coefficients \mathfrak{T} and \mathfrak{R} using Wronskians, cf. Definition 2.4.

Now, we prove that $\mathfrak{T}(\omega, \ell)$ has a pole of order one at $\omega = i\kappa_+$ assuming that $\ell(\ell + 1) \neq r_+^2 (r_+ - 3r_-)$. First, note that

$$u_1^{(1)}(r_*) = \sum_{m_1 \in \mathbb{N}} C_{m_1} d_{m_1} e^{2\kappa + m_1 r_*} \quad (\text{A.29})$$

has a pole of order one at $\omega = i\kappa_+$ since $C_1 \neq 0$, see (A.16). Since for $n \neq 1$ there is no term of the form $e^{2\kappa r_*}$ in (A.22) as $m_n \geq n$, the pole at $\omega = i\kappa_+$ cannot be canceled by the other terms and must occur in u_1 . Moreover, this pole of u_1 at $\omega = i\kappa_+$ is not of higher order than one since d_1 does not occur at higher powers than one in the Volterra iteration. This implies that $\mathfrak{T}(\omega, \ell)$ has a pole of order one at $\omega = i\kappa_+$. \square

Connectedness of the Subextremal Parameter Range

Proposition A.3. *Let the subextremal parameter space $\mathcal{P}_{\text{se}}^{\Lambda > 0}$ and $\mathcal{P}_{\text{se}}^{\Lambda < 0}$ be defined as in (3.39) and (3.40), respectively. Then, $\mathcal{P}_{\text{se}}^{\Lambda > 0} \cap \{Q > 0\}$, $\mathcal{P}_{\text{se}}^{\Lambda < 0} \cap \{Q > 0\}$, $\mathcal{P}_{\text{se}}^{\Lambda > 0} \cap \{Q < 0\}$ and $\mathcal{P}_{\text{se}}^{\Lambda < 0} \cap \{Q < 0\}$ are path-connected.*

Proof. The claim follows for $\mathcal{P}_{\text{se}}^{\Lambda > 0} \cap \{Q > 0\}$ and $\mathcal{P}_{\text{se}}^{\Lambda > 0} \cap \{Q < 0\}$ from the following continuous parameterizations

$$\mathcal{P}_{\text{se}}^{\Lambda>0} \cap \{Q > 0\} = \left\{ (M, Q, \Lambda) \in \mathbb{R} \times \mathbb{R} \times \mathbb{R} : \begin{aligned} \Lambda &= 3(r_+^2 + r_-^2 + r_c^2 + r_+r_c + r_c r_- + r_+r_-)^{-1}, \\ 6M &= \Lambda(r_+ + r_-)(r_+ + r_c)(r_- + r_c), \\ Q &= \left(\frac{\Lambda}{3}(r_+ + r_- + r_c)(r_- r_+ r_c) \right)^{\frac{1}{2}} \\ \text{for } 0 < r_- < r_+ < r_c \end{aligned} \right\} \quad (\text{A.30})$$

and

$$\mathcal{P}_{\text{se}}^{\Lambda<0} \cap \{Q > 0\} = \left\{ (M, Q, \Lambda) \in \mathbb{R} \times \mathbb{R} \times \mathbb{R} : \begin{aligned} \Lambda &= 3 \left(\frac{3}{4}(r_+ + r_-)^2 - r_+r_- - \xi_i \right)^{-1}, \\ 6M &= -\Lambda \left(\frac{1}{4}(r_+ + r_-)^2 + \xi_i - r_+r_- \right) (r_+ + r_-), \\ Q &= \left(-\frac{\Lambda}{3}r_+r_- \left(\frac{3}{4}(r_+ + r_-)^2 + \xi_i \right) \right)^{\frac{1}{2}} \\ \text{for } 0 < r_- < r_+ \text{ and } \xi_i > \left(\frac{3}{4}(r_+ + r_-)^2 - r_+r_- \right)^{\frac{1}{2}} \end{aligned} \right\} \quad (\text{A.31})$$

in view of the fact that $\{0 < r_- < r_+ < r_c\}$ and $\{0 < r_- < r_+, \xi_i > (\frac{3}{4}(r_+ + r_-)^2 - r_+r_-)^{\frac{1}{2}}\}$ are path-connected as subsets of \mathbb{R}^3 . In the following, we will show (A.30) and (A.31).

First, in the case $\Lambda > 0$, note that (A.30) follows from comparing coefficients of

$$\frac{-3}{\Lambda}(r^2 - 2Mr + Q^2 - \frac{1}{3}\Lambda r^4) = (r - r_-)(r - r_+)(r - r_c)(r - r_0)$$

for $r_0 < 0 < r_- < r_+ < r_c$. Indeed, we obtain $r_0 = -(r_- + r_+ + r_c)$ and (A.30) can be deduced.

In the case $\Lambda < 0$, note that $\frac{-3}{\Lambda}(r^2 - 2Mr + Q^2 - \frac{1}{3}\Lambda r^4)$ only has two real roots $0 < r_- < r_+$ such that we compare coefficients of

$$\frac{-3}{\Lambda}(r^2 - 2Mr + Q^2 - \frac{1}{3}\Lambda r^4) = (r - r_-)(r - r_+)(r - \xi)(r - \bar{\xi})$$

with $\xi = \xi_r + i\xi_i$. We obtain $2\xi_r = -(r_+ + r_-)$ and $\xi_i > (\frac{3}{4}(r_+ + r_-)^2 - r_+r_-)^{\frac{1}{2}}$ to guarantee $\Lambda < 0$. Now, a direct computation shows (A.31).

Completely analogously, we can show path-connectedness for $\mathcal{P}_{\text{se}}^{\Lambda>0} \cap \{Q < 0\}$ and $\mathcal{P}_{\text{se}}^{\Lambda<0} \cap \{Q < 0\}$. \square

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Communicated by Mihalis Dafermos.

Received: July 1, 2018.