

Dynamical Response of Viscous Objects to Gravitational Waves

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We study the dynamical response of viscous materials to gravitational waves, in the context of a fully relativistic theory of fluid dynamics. For the first time, we calculate oscillation modes and scattering properties of viscous stars. Viscous stars absorb high-frequency radiation, following a dispersion relation introduced by Press. In the extremely large viscosity regime, stars would become reflectors of waves, but this regime appears to be forbidden by causality bounds. In the context of black hole mimickers, we show how maximally viscous stars on the threshold of stability mimic the absorption of a black hole with the same mass. Our results suggest that rotating viscous stars will amplify incoming radiation.

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Introduction—Knowledge of the microscopic structure of matter, or of the large scale structure of our Universe, and all the technology which our society is built upon, are possible thanks to the coupling between light and matter. Reflection and refraction of electromagnetic waves are consequences of Maxwell's equations, which are used to infer properties of the media, or sometimes nuisances that prevent us from understanding phenomena occurring at large distances. Matter is not impervious to light.

The interaction between gravitational waves (GWs) and matter, and its effect on GW generation, is an active field of research [1–20]. On cosmological scales, the weak gravitational coupling, together with the generally low density of matter, turns our Universe practically transparent to GWs [1–4]. Particularly, if matter is modeled as a perfect fluid, GWs passing through it are left completely undamped. Viscosity-driven dissipation in our expanding Universe is expected to be very small [21].

However, for more compact matter distributions, the story is different. In an era where precision GW physics is making its way and where multimessenger astrophysics is a reality, a precise understanding of the interaction between GWs and matter is necessary. We highlight, in particular, the relevance for the following: (i) The physics of accretion disks. GWs can heat them, potentially giving rise to electromagnetic signatures [22].

(ii) The physics of stars, since viscous dissipation plays an important role in spectroscopy [23–25], can saturate the f - and r -mode instabilities [26–30] or may quench the maximum-mass instability, possibly allowing for larger-mass neutron stars, as well as significantly impact the after-

merger GW signal [31–34]. Previous works estimated the damping rate of stellar modes due to internal dissipative processes based on energy balance laws, motivated by the separation of scales between the mode oscillation rate and the viscous dissipation timescale [24]. A self-consistent, relativistic calculation of the impact of viscosity on spectroscopy, in the context of a well-posed relativistic theory is, surprisingly, lacking or treated only partially [35].

(iii) The physics of superradiance. Dissipation mechanisms imply amplification during scattering of low-frequency GWs when objects rotate [36]. Viscous damping might, therefore, lead to superradiant amplification of GWs.

(iv) The black hole paradigm. It has been argued that the rapid relaxation seen in GW events must be associated with horizons or then with a material with abnormally large viscosity [37]. However, if electromagnetism is a good guide, highly viscous materials without horizons should be very good reflectors of GWs, possibly leading to characteristic signatures [38].

Dissipating materials: Zeldovich model—Before turning to a relativistic model of viscosity, we point out that objects that dissipate energy very efficiently tend to behave as perfect reflectors of radiation. We start with a bare-bones model of absorption. Let Ψ be a scalar field propagating in flat space, which is scattered at a planar interface separating vacuum, for positive $z > 0$ (region I), from a medium with some constant absorption α , located at $z < 0$ (region II). The following captures the propagation of Ψ in an absorbing medium at rest [39,40],

$$\square\Psi + \alpha\frac{\partial\Psi}{\partial t} = 0. \quad (1)$$

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For a wave traveling from region I toward region II,

$$\begin{aligned}\Psi &= A_{\text{in}} e^{-i\omega t + ik_x x - ik_z z} + A_{\text{out}} e^{-i\omega t + ik_x x + ik_z z} \quad (\text{region I}), \\ \tilde{\Psi} &= \mathcal{T} e^{-i\omega t + ik_x x - iKz} \quad (\text{region II}).\end{aligned}$$

The equation of motion implies the dispersion relation,

$$\begin{aligned}\omega^2 &= k_x^2 + k_z^2, & z > 0, \\ \omega^2 - k_x^2 - K^2 - i\omega^2 \tilde{\alpha} &= 0, & z < 0,\end{aligned}$$

with $\tilde{\alpha} \equiv \alpha/\omega$.

Imposing continuity of Ψ and its derivative along the interface yields

$$\left| \frac{A_{\text{out}}}{A_{\text{in}}} \right|^2 \rightarrow \frac{\omega^4 \tilde{\alpha}^2}{16(\omega^2 - k_x^2)^2} + \mathcal{O}(\tilde{\alpha}^4), \quad \tilde{\alpha} \rightarrow 0, \quad (2)$$

$$\rightarrow 1 - \frac{2\sqrt{2(\omega^2 - k_x^2)}}{\omega\sqrt{\tilde{\alpha}}} + \mathcal{O}(\tilde{\alpha}^{-1}), \quad \tilde{\alpha} \rightarrow \infty. \quad (3)$$

At small $\tilde{\alpha}$ (analogous to low-viscosity fluids), the reflection coefficient is small, and the wave is nearly perfectly transmitted from region I to region II, as might be anticipated. In the high damping limit ($\tilde{\alpha} \rightarrow \infty$), the situation is reversed, and almost all the incident wave gets reflected back.

For spherical boundaries, similar features and scaling with α are recovered [40]. One can also study sound waves in planar interfaces separating an ideal from a viscous fluid [41,42]. The result is similar: in the high-dissipation regime ($\tilde{\alpha} \rightarrow \infty$), the reflection coefficient tends to unity. Highly absorbing materials are good reflectors of radiation. This procedure is akin to studying reflection of electromagnetic waves off a perfect conductor, a material that easily absorbs and interacts with incoming radiation. Perfect conductors, too, are perfect reflectors [43]. We turn now to viscous stars and GWs.

Formalism: Axial GWs—we consider axial perturbations of a self-gravitating, viscous fluid in a fully relativistic approach introduced by Bemfica *et al.* [45–47] and Kovtun [44] (see also Refs. [48,49]), which we refer to as BDNK in the following, which we refer to as BDNK in the following. The perturbation equations—summarized below—are discussed in detail elsewhere [Eqs. (38)–(40) in Ref. [50]]. Static, spherically symmetric geometries describing a non-rotating star can be written as

$$ds^2 = -e^\Phi dt^2 + e^\Lambda dr^2 + r^2 d\Omega^2, \quad (4)$$

where $\Phi = \Phi(r)$, $\Lambda = \Lambda(r)$, and $d\Omega^2$ is the line element of the unit sphere. We can define the gravitational mass inside a shell of radius r as $M(r) = r(1 - e^{-\Lambda})/2$. All the viscous terms vanish in the background, and the equilibrium equations reduce to the usual Tolman-Oppenheimer-Volkoff equations,

$$\frac{dM}{dr} = 4\pi r^2 \mathbf{e}, \quad (5)$$

$$\frac{d\Phi}{dr} = \frac{2e^\Lambda}{r^2} (M + 4\pi r^3 \mathbf{p}), \quad \frac{d\mathbf{p}}{dr} = -\frac{\mathbf{e} + \mathbf{p}}{2} \frac{d\Phi}{dr}, \quad (6)$$

where $\mathbf{p}$ is the fluid pressure, related to the energy density $\mathbf{e}$ via the equation of state, with associated sound speed $c_s^2 = \partial\mathbf{p}/\partial\mathbf{e}$. Neutron stars may have crusts with their own torsional [51–53] and shear modes [54,55]. For simplicity, we will ignore these and other aspects present when there are phase transitions and focus on polytropic equations of state,

$$\mathbf{p} = \kappa \mathbf{e}^{1+1/n}. \quad (7)$$

For concreteness, we take $\kappa = 700 \text{ km}^{-2.5}$, $n = 4/5$ and choose a rather compact star with central density $\rho_c = 3 \times 10^{15} \text{ g cm}^{-3}$ resulting in an object with mass $M_S = 1.54 M_\odot$, radius $R_S = 8.78 \text{ km}$, and compactness $M_S/R_S = 0.259$. Although the star is relatively compact, these parameters render the star stable against radial perturbations [56] and are approximately compatible with current constraints.

In a particular gauge, axial metric perturbations can be written as

$$\delta^{(1)} g_{\mu\nu} = \begin{pmatrix} \mathbf{0} & e^{\Phi/2} k_u^{\ell m} X_\theta^{\ell m} & e^{\Phi/2} k_u^{\ell m} X_\phi^{\ell m} \\ \text{sym} & e^{\Lambda/2} k_n^{\ell m} X_\theta^{\ell m} & e^{\Lambda/2} k_n^{\ell m} X_\phi^{\ell m} \\ & 0 & 0 \\ & 0 & 0 \end{pmatrix}, \quad (8)$$

where $X_{\theta/\phi}^{\ell m}$ are vector spherical harmonics, and we follow the definition and conventions of Ref. [57]. Here, $k_{u,n} = k_{u,n}(t, r)$. We drop the ℓm label from now on, assuming we deal with a single mode. The perturbation of the fluid four-velocity is written as

$$\delta^{(1)} u_\mu = \beta [0, 0, X_\theta, X_\phi], \quad (9)$$

with $\beta = \beta(t, r)$. The perturbed Einstein equations lead to the following coupled wave equations [50]:

$$-\frac{\partial^2 \psi}{\partial t^2} + \frac{\partial^2 \psi}{\partial r_\star^2} - \mathcal{V}\psi = 16\pi\eta e^{\Phi/2} \left(\frac{\partial \psi}{\partial t} + \mathbf{a}\beta \right), \quad (10)$$

$$\begin{aligned} & -\tau \frac{\partial^2 \beta}{\partial t^2} + \frac{\eta}{\mathbf{e} + \mathbf{p}} \frac{\partial^2 \beta}{\partial r_\star^2} + \eta \left(\mathbf{b}_1 \frac{\partial \beta}{\partial r_\star} + \mathbf{b}_2 \frac{\partial \beta}{\partial t} + \mathbf{b}_3 \beta \right) \\ & = \mathbf{c}_1 \frac{\partial^2 \psi}{\partial t \partial r_\star} + \mathbf{c}_2 \frac{\partial \psi}{\partial t} + \mathbf{c}_3 \frac{\partial \psi}{\partial r_\star} + \mathbf{c}_4 \psi, \end{aligned} \quad (11)$$

where the tortoise coordinate $dr_\star = e^{(\Lambda-\Phi)/2} dr$, and the master function ψ is the usual Regge-Wheeler master

variable, related to the metric perturbation k_n via $r\psi = e^{\Phi/2}k_n$, as well as the Regge-Wheeler potential,

$$\mathcal{V} = e^{\Phi} \left(\frac{\ell(\ell+1)}{r^2} - \frac{6M}{r^3} + 4\pi(\mathbf{e} - \mathbf{p}) \right). \quad (12)$$

The coefficients $\mathbf{a}$, $\mathbf{b}_{1,2,3}$, and $\mathbf{c}_{1,2,3,4}$, as well as the transport equation for k_n , can be found in Supplemental Material [58]. The variables that we solve with, ψ and β , respectively, describe propagating gravitational waves and differential rotation inside the star. The equations are written in terms of the shear viscosity η and the transport coefficient τ (“ τ_Q ” in Refs. [47,50]). We work in the family of causal and stable frames, but note that hydrodynamic modes are invariant under frame redefinitions [59]. We will choose a particular dependence of these coefficients in terms of the properties of the star and dimensionless, constant numbers, in the following manner:

$$\eta = \mathbf{p} R_S \hat{\eta}, \quad \tau = \frac{\mathbf{p}}{\mathbf{e}} R_S \hat{\tau}, \quad (13)$$

where R_S is the radius of the star. This parametrization makes the principal part of the equation for β regular across the star. In order for the theory to be causal, the dimensionless parameters $\hat{\eta}$ and $\hat{\tau}$ must satisfy certain constraints [47,60]. A simple analysis shows that, in order for the characteristic speed of β to be smaller than unity, $\hat{\eta} \leq \hat{\tau}$. Causality enforces stricter constraints, see Supplemental Material [58]. In particular, $\hat{\eta} \leq 3/4$ and $\hat{\tau} > 1 + C\hat{\eta}$, where C is a number that only depends on the equation-of-state parameters and the sound speed inside the star.

Boundary and junction conditions—At the center of the star ($r \sim 0$), regular solutions $\psi, \beta \sim a_{\psi,\beta} r^{\ell+1}$ (irregular solutions behave as $\psi, \beta \sim r^{-\ell}$). At the surface, we use Israel’s junction conditions for the Einstein tensor $[[G]]_{r\mu} = 0$ where $[[X]] = X(R_+) - X(R_-)$ denotes the jump of the quantity X at the surface of the star [61]. In the notation of [50], for the perturbation of the stress energy tensor, $\delta^{(1)}T_{r\mu} \propto \theta_n$, which, since the pressure vanishes at the surface, yields

$$\eta \left[2r^2 e^{\Lambda/2} \frac{d\psi}{dt} + e^{\Phi/2} \left(2r \frac{d\beta}{dr} + \beta(e^{\Lambda} - 5) \right) \right] = 0. \quad (14)$$

We find that this is equivalent to the nonrotating limit of the junction condition (40) in Refs. [62,63]. Regularity conditions also impose continuity of ψ and its radial derivative. Finally, in the exterior of the star, $\beta = 0$ and far away monochromatic solutions behave as

$$\psi \sim A_{\text{in}} e^{-i\omega t - i\omega r_*} + A_{\text{out}} e^{-i\omega t + i\omega r_*}. \quad (15)$$

Since we choose to work with polytropic equations of state for the fluid, Eq. (14) is automatically satisfied at the

surface of the star for our choice of η (13). However, some of the coefficients of the lower order terms in the equation for β diverge at the surface of the star. Therefore, regularity demands the condition

$$\begin{aligned} z_S \frac{\partial \beta}{\partial t} + \frac{M_S \hat{\eta}}{R_S} \frac{\partial \beta}{\partial r_*} + \frac{M_S(5M_S - 2R_S) \hat{\eta}}{R_S^3} \beta \\ = - \frac{M_S \hat{\eta}}{z_S} \frac{\partial \psi}{\partial t} - R_S \frac{\partial \psi}{\partial r_*} - z_S^2 \psi, \end{aligned} \quad (16)$$

where $z_S^2 = 1 - 2M_S/R_S$. We have solved Eqs. (10) and (11) with two independent approaches. One is a direct integration in the frequency domain that “shoots” for the a_{ψ}/a_{β} ratio [64,65], and the other is a time-domain integration routine. Results from both codes agree, and we recover with very good accuracy the quasinormal mode (QNM) frequencies in the perfect fluid case [66]. Details on the numerical implementation of both approaches are given in Supplemental Material [58].

Scattering—We can now address important problems in the context of gravitational radiation. The first concerns scattering of GWs from viscous stars. The problem of scattering reduces to calculating

$$\mathcal{R}^2(\omega) = \frac{|A_{\text{out}}|^2}{|A_{\text{in}}|^2}, \quad (17)$$

for a given frequency ω , with a solution that satisfies the appropriate boundary conditions. Given that the system of equations is linear and homogeneous, the only free parameter at the origin is the ratio a_{ψ}/a_{β} , chosen such that the solution is well behaved at the surface.

Our results are summarized in Figs. 1 and 2. When the radiation wavelength is large (small frequency), it barely interacts with the star and the reflection coefficient is unity. For inviscid materials, the same happens at *any* frequency: the wave penetrates the star but there is no dissipation, so it exits with a phase shift and no absorption.

In the presence of viscosity, the wave interacts with the star material and gets partially absorbed. Higher frequencies induce larger gradients and hence a more effective absorption, resulting in smaller reflection. This mimics the reflection properties of black holes (BHs) (shown as a dashed black line in Fig. 1), for which the reflection coefficient vanishes when $\omega M_S \gtrsim 1$. This is also achieved by increasing the star compactness. In Supplemental Material [58], we show that we recover this behavior even while freezing the fluid modes, in the inverse Cowling approximation (ICA) [68].

In fact, Fig. 1 shows that stars close to the threshold of radial stability have a reflectivity curve similar to that of a nonrotating BH with the same mass. The large frequency reflectivity of viscous stars can be explained simply by recovering the dispersion relation satisfied by GWs passing through a viscous medium [38],

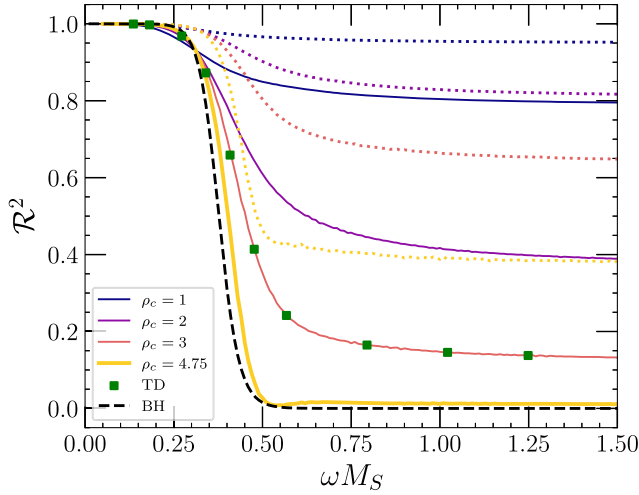


FIG. 1. Reflection coefficient $\mathcal{R}^2$ of Eq. (17) for a quadrupolar ($\ell = 2$) axial GW scattered off a compact, viscous star with $\kappa = 700 \text{ km}^{2.5}$, $n = 4/5$ and varying central density, in units of $10^{15} \text{ g cm}^{-3}$, cf. (7) and legend. We set the transport coefficients to $\hat{\eta} = 0.7$ (solid lines) and $\hat{\eta} = 0.15$ (dotted lines), keeping $\hat{\tau} = 10$. The black dashed line is reflectivity of a nonspinning black hole with mass M_S [36,67]. Green squares correspond to reflectivity extracted from time-domain (TD) simulations for $\rho_c = 3$, showing excellent agreement between our two independent methods.

$$k^2 = \frac{\omega^2}{c^2} \left(1 + i \frac{16\pi G \eta}{\omega} \right). \quad (18)$$

We recover this exact relation by taking the master equation for ψ in the dilute, flat space limit. At large frequencies $\omega \gg G\eta$, the reflectivity $\mathcal{R}^2 = \exp(-16\pi R_S \eta)$, with an accurate description of results in Fig. 2. In Supplemental Material [58], we compute the reflectivity in the dilute, Newtonian regime, finding very good agreement with the fully generic calculation shown here, for small viscosities. Notice that these absorption features are different than the ones expected for models of exotic compact objects, based on a partially absorbing surface [69].

Based on well-known properties of electric conductors and on the examples we discussed concerning large dissipation [Eq. (3)], one is tempted to conclude that very viscous fluids reflect GWs. In fact, Press used this behavior to hypothesize a material, “respondium,” able to reflect GWs or act as a waveguide [38]. Increasing the viscosity past $\hat{\eta} > 3/4$, we observe a transition, which manifests itself in two ways: (a) low-frequency waves are amplified, $\mathcal{R}^2 > 1$, typical of systems with superluminal motion [36]; (b) at large frequencies, waves are no longer absorbed efficiently, but instead are better reflected at the surface of the star as η increases (cf. Fig. 2). Further evidence for efficient reflection is provided by the wave function ψ penetrating less efficiently into the star once viscosity

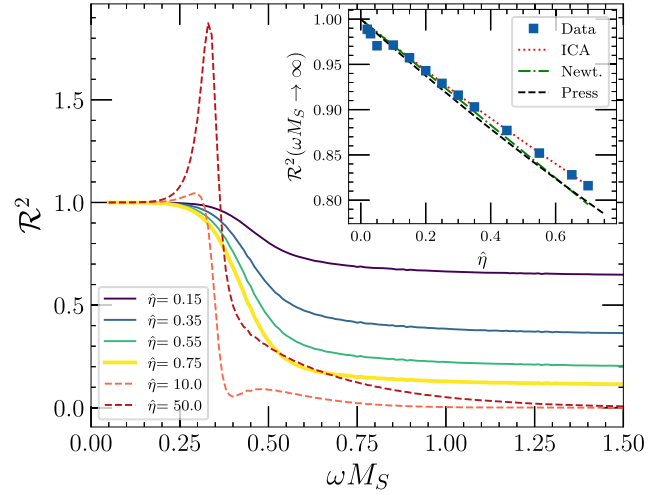


FIG. 2. Reflectivity of the star for axial GWs for different values of the dimensionless shear viscosity $\hat{\eta}$, keeping $\hat{\tau} = 10$. Dashed orange and red lines correspond to configurations that violate the BDNK causality bounds and, as a consequence, result in amplification, $\mathcal{R}^2 > 1$. Inset: the value of the reflectivity at large frequencies $\omega M_S \gg 1$, showing good agreement with the ICA, as well as with the Newtonian limit (see Supplemental Material [58]), and the dispersion relation (18).

increases past the causal bound. When attempting to reproduce these results from time-domain scattering experiments, we find an instability, hinting toward the presence of unstable modes in the system. Therefore, highly viscous stars are, effectively, some kind of respondium material and behave following the examples we worked out above. However, these properties are manifest only when violating the BDNK causality bounds, which are not present in the simplified examples. Therefore, exploring alternative parametrizations of the viscosity or numerical simulations of the full nonlinear problem is necessary to understand whether respondiumlike behavior is possible, while respecting causality.

Axial modes—We also calculated axial QNMs, i.e., eigenfrequencies $\omega = 2\pi f - i/\tau$ for which the wave function is regular at the center and at the surface and for which $A_{\text{in}} = 0$. We identified a spacetime mode [70] and cross-checked our result with time-domain simulations, as described in Supplemental Material [58]. Our results are summarized in Table I. Increasing $\hat{\eta}$ results in a smaller real frequency and a longer damping time. This trend is reproduced by the inverse Cowling approximation. The dependence on η is weak, and the variation of frequency and damping relative to the perfect fluid case is no more than a few percent. Therefore, accurately modeling viscosity is important in order to perform highly precise GW asteroseismology, but it is unlikely to largely bias the inference, even for large viscosities. This is mainly due to the fact that spacetime modes are dominated by the compactness of the star.

TABLE I. Fundamental spacetime mode written as (f, τ) , with f in kilohertz and τ in microseconds, for $\hat{\tau} = 5$ (second column) and in the ICA (third column), for different values of $\hat{\eta}$. The last column shows the combined, relative difference between the QNM frequency and the one of the same mode for a perfect fluid, in percent value, for the two previous columns.

$\hat{\eta}$	$\hat{\tau} = 5$	ICA	$\delta\omega$ (%)
0.05	(9.44,46.5)	(9.49,46.7)	(0.7,0.4)
0.25	(9.24,46.7)	(9.43,47.5)	(2.9,2.3)
0.45	(8.97,47.0)	(9.36,48.3)	(5.7,4.1)
0.65	(8.78,48.1)	(9.29,49.0)	(8.4,5.9)

Additionally, in Supplemental Material [58], we compute the spectrum of the fluid modes in the decoupling limit, i.e., freezing the gravitational perturbations. We show the appearance of a family of modes with damping times of ~ 0.05 – 0.15 ms and oscillation frequencies ranging between 1 and 2 kHz for the fundamental mode of this family, up to, e.g., 35 kHz for the 20th overtone with $\hat{\eta} = 0.7$. This is in good agreement with our numerical time evolutions, which show some evidence for rapidly oscillating modes.

The properties of these fluid modes beyond the decoupling limit requires further scrutiny, including an extension of the formalism to the even parity sector [50] (note that viscosity is known to add structure to the spectrum of nontrivial flows [71,72]).

Conclusions—We have derived the equations governing the propagation of GWs through viscous matter (in particular, a star) in the context of a fully relativistic hydrodynamics theory. We find that viscosity induces absorption of GWs, and we have calculated the precise response of compact stars to impinging waves. Compactness also enhances the absorption of GWs. In fact, stars with the maximum compactness for stability, and with the maximum viscosity allowed by causality constraints, reflect GWs in a manner very similar to a black hole with the same mass. When increasing the viscosity past the causal bounds, we find amplification, showing similar features as over-reflection in the Poiseuille flow due to the Orr mechanism [71,72]. Given our results, it is possible that rotating stars can superradiate or even develop superradiant instabilities, which we aim to explore in the future. In the large frequency limit, we recover the behavior predicted by the dispersion relation introduced by Press [38]. Our findings are consistent with past remarks [38]: a perfect reflector of GWs that does not violate energy conditions would be more compact than a black hole. An interesting open question is whether this requirement can be bypassed through quantum collective dynamics of matter reacting to GWs [73].

Absorption of GWs leads to heating of the star. For the maximum allowed $\hat{\eta} = 3/4$, $\eta_{\max} = 1.5 \times 10^{31}$ g/cm/s is comparable to the effective viscosity

$\eta_{\text{BH}} = 1.8 \times 10^{31}$ g/cm/s [37]. This is much larger than the usual shear viscosity expected for a neutron star [37]. However, the absorption is suppressed by geometrical factors. For a supermassive BH coalescence of total mass $M = 10^6 M_\odot$ and a neutron star orbiting the binary at radius $10^3 M$, the fraction of GW energy released by the binary that impinges on the neutron star is $\sim 10^{-16}$. It is very low-frequency radiation (in the low-frequency portion of Fig. 2), hence the energy absorbed is $\ll 10^{-14} M_\odot$. However, during the inspiral stage of binary neutron star mergers, dissipative effects due to viscosity will alter the gravitational waveform through the dynamical tidal deformabilities [74–77], opening the possibility of constraining the viscosities of neutron stars from gravitational wave observations.

We can extrapolate our results in the high-frequency regime to non-self-gravitating, thin circumbinary accretion disks around a BH binary, for which $\eta = \alpha \rho \Omega H^2$ [6,78,79], with ρ the local disk density, Ω the Keplerian orbital velocity, and H the disk height. Now,

$$M\eta = 5 \times 10^{-12} \frac{\alpha}{10^{-2}} \frac{M^2 \rho}{10^{-6}} \frac{M\Omega}{0.05} \left(\frac{H}{0.1M} \right)^2. \quad (19)$$

We are now in the $\omega \gg \eta$ regime. Hence, we find that the ratio between the absorbed GW luminosity $\mathcal{L}_{\text{GW}}^{\text{absorbed}} \sim \mathcal{L}_{\text{GW}}^{\text{peak}} M\eta$ to the total disk luminosity $\mathcal{L}_{\text{accr}} \sim \dot{M} = f_{\text{Edd}} \dot{M}_{\text{Edd}}$, with $\dot{M}_{\text{Edd}}$ the Eddington accretion rate,

$$\frac{\mathcal{L}_{\text{GW}}^{\text{absorbed}}}{\mathcal{L}_{\text{accr}}} \sim 10^4 \frac{10^{-4}}{f_{\text{Edd}}} \frac{10^6 M_\odot}{M} \frac{\alpha}{10^{-2}} \frac{M^2 \rho}{10^{-6}} \frac{M\Omega}{0.05} \left(\frac{H}{0.1M} \right)^2,$$

where the peak GW luminosity for equal mass mergers $\mathcal{L}_{\text{GW}}^{\text{peak}} \sim 10^{23} \mathcal{L}_\odot$ [80]. GWs from mergers can heat up accretion disks to an important level, consistent with the findings in Ref. [5].

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