Stars as resonant absorbers of gravitational waves

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ABSTRACT

Quadrupole oscillation modes in stars can resonate with incident gravitational waves (GWs), and grow non-linear at the expense of GW energy. Stars near massive black hole binaries (MBHB) can act as GW-charged batteries, discharging radiatively. Mass-loss from these stars can prompt MBHB accretion at near-Eddington rates. GW opacity is independent of amplitude, so distant resonating stars can eclipse GW sources. Absorption by the Sun of GWs from Galactic white dwarf binaries may be detectable with second-generation space-based GW detectors as a shadow within a complex diffraction pattern.

Key words: gravitational waves–stars: interiors–stars: oscillations–opacity–galaxies: active

1 INTRODUCTION

Supermassive black holes (SMBH) with masses in the range \(10^6 - 10^9 M_\odot\) are present in the nuclei of most, perhaps all, nearby galaxies (see e.g. the recent review by Kormendy & Ho 2013). Mergers between galaxies should result in supermassive black hole binaries; indeed active SMBH binaries have been directly resolved at 0.1-1kpc separations in X-rays (Komossa et al. 2013; Comerford & Greene 2014), and at \(\sim\)10 pc separation in the radio (Rodriguez et al. 2006). Most of the binding energy of a merging massive binary is radiated as gravitational waves (Thorne & Braginsky 1976). As the binary approaches merger, the gravitational wave (GW) frequency \(\nu_{GW}\) increases in a chirp, passing through quadrupolar \((\ell = 2)\) oscillation frequencies \(\nu_\ell\) of stars and stellar remnants, resonating whenever \(\nu_{GW} \sim \nu_\ell\). The interaction of GWs with matter has been considered in various contexts, (e.g. Hawking 1966; Kocsis & Loeb 2008, Li, Kocsis & Loeb 2013); the latter suggesting that viscous heating of Sun-like stars by GW from a nearby merging massive black hole binary can reach \(\sim L_\odot\). However, resonant interactions of GWs with normal, as opposed to compact stars (similar to a bar detector), has not received very much attention (Misner, Thorne & Wheeler 1973; Chandrasekhar & Ferrari 1991, 1992; Siegel & Roth 2010, 2011). It has been shown that GW can do work on stellar oscillations leading to potential observable effects on the oscillations (e.g. Fabian & Gough 1984; Kojima & Tanimoto 2005). After this manuscript was submitted, a pre-print appeared on arXiv.org by Lopes & Silk 2014, considering the resonant interaction of GWs with stars, as well as assessing the feasibility of detecting the induced stellar oscillations through astroseismological measurements. In this Letter, we discuss the possibility of GW absorption lines at resonant frequencies in stars, eclipses of GW sources by foreground stars (including the Sun) and the possible use of stars in galactic nuclei as electromagnetic detectors of resonating GW from nearby massive black hole binaries. In the latter case, we show that resonant heating of a single mode in a Sun-like star can be up to \(\sim 11\) orders of magnitude larger than the viscous heating in Li, Kocsis & Loeb 2013.

2 GWS FROM A BINARY RESONATING WITH STELLAR OSCILLATIONS.

A circularized binary with individual BH masses \(M_1\) and \(M_2\) and physical separation \(a_{bin}\) emits GWs at frequency

\[
\nu_{GW} = \frac{2}{t_{orb}} = \frac{G^{1/2}M_{bin}^{1/2}}{\pi a_{bin}^{3/2}} = 2 M_\odot^{-1} a_0^{-3/2} \nu_G W, \text{1 mHz},
\]

for a characteristic duration

\[
t_{GW} = \frac{\Delta a_{bin}}{\dot{a}_{bin}} = 0.8 \eta^{-1} \dot{M}_e^{-5/3} \nu_{GW}^{-8/3} \text{yr},
\]

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where $t_{\text{orb}}$ is the orbital period, $M_{\text{bin}} = M_1 + M_2$ is the total binary mass, $r_{\text{eq}} = GM_{\text{bin}}/c^2$ is the gravitational radius of the binary, $a_1 \equiv a_{\text{in}}/(10^9)$, $M_6 \equiv M_{\text{bin}}/10^6 M_\odot$, $\eta \equiv M_1 M_2/(M_1 + M_2)^2$ is the symmetric mass ratio, $\eta_{-3} \equiv \eta/10^{-3}$, $\nu_{\text{GW},1} \equiv \nu_{\text{GW}}/1$ mHz and where the orbital decay $\dot{a}_{\text{bin}}$ is driven by the (quadrupolar) GW emission (Peters & Mathews 1963). The resulting GW strain amplitude averaged over directions is given by

$$h = \sqrt{\frac{32}{5}\frac{G^2}{c^4} \frac{M_{\text{bin}} \dot{a}_{\text{bin}}}{D_{a_{\text{in}}}}} = 1.6 \times 10^{-7} \tau^{2/3}_{\text{stellar}} M_6^{-1/3} \mu_3 D_{*,r_1}^{-1}$$

where $M_{\text{bin}}$ is the reduced mass, $D_*$ is the resonating star’s distance from the binary, $\mu_3 \equiv M_{\text{bin}}/10^3 M_\odot$ and $D_{*,r_1} \equiv D_*/10^3 r_\odot$. Sun-like stars have $\ell = 2$ oscillation modes with frequencies $\omega_{*2} = 2n\nu_c$ spanning $10\mu$Hz–$0.1$Hz (Aerts et al. 2010), which can match the frequency of GWs from a binary source. Tens of low-radial-order $f$, $g$, and $p$-modes with overlap integrals $\gtrsim 10^{-3} M_\odot$ span $\sim 0.1$–$1$ mHz in solar models (Aerts et al. 2010).

We follow the approach and definitions of Rathore et al. (2005) in representing GW-driven oscillations of a stellar mode by a driven damped harmonic oscillator, whose displacement $x(t)$ is the solution to

$$\ddot{x} + \frac{\dot{x}}{\tau_d} + \omega^2 x = F(t),$$

where $\tau_d$ is the damping time of the stellar mode and $F(t) = F_{\text{GW}}(t)$ is the driving force. Low-order g-modes in the linear regime damp radiatively on the timescale $\tau_d \sim 10^8$ yr (Kumar & Goodman 1996). However, in the nonlinear regime (with energy in the mode $E_{\infty} \gtrsim 10^{35}$ erg), coupling to high-degree g-modes reduces this timescale to $\tau_d \approx 50 E_{\infty}^{1/2}$ day, (here $E_{\infty} = E_{\infty}/10^{42}$ erg). For $f$-modes in convective stars, linear dissipation through turbulent viscosity takes $\sim 10^4$ yr (Ray, Kembhavi & Antia 1987), but in the nonlinear regime, dissipation via high order p-modes (of degree $\ell = 0, 2, 4$) occurs on timescale $\tau_d = 3 \times 10^4 E_\nu^{1/2}$ day (Kumar & Goodman 1996). Multiplying eq. (4) by $\dot{x}$ and integrating over time yields the familiar form, expressing conservation of energy

$$\dot{E}_m + \dot{Q} = \dot{W},$$

where $E_m$ is the mechanical energy, $Q$ is the energy lost via dissipation, and $W$ is the work done by the driving force,

$$E_m = \frac{\dot{x}^2}{2} + \frac{\omega^2_x x^2}{2}, \quad Q = \int_{t_0}^{t} \dot{x} \ddot{x} dt, \quad W = \int_{t_0}^{t} F(t) \dot{x} dt,$$

respectively. All three quantities are per unit mass in a single mode. Excitation of non-radial oscillations in stars and stellar remnants due to tidal capture is relatively well-studied (e.g. Press & Teukolsky 1977; Reisenegger & Goldreich 1994; Rathore et al. 2005), compared to oscillation excitation due to incident GWs. In the latter case, the effective driving force per unit mass due to GWs is nearly sinusoidal, with a characteristic amplitude $|F_{\text{GW}}| = \omega^2_{\text{GW}} h R_\odot$. (Miser, Thorne & Wheeler 1973; Khosroshahi & Sobouti 1997; Siegel & Roth 2010, 2011). The frequency of $|F_{\text{GW}}|$ during the inspiral phase of a binary evolves slowly (i.e. the number of orbits at $\nu_{\text{GW}}$ is $N = \nu_{\text{GW}}/\nu_{\text{GW}} > 1$). In this limit, $|F_{\text{GW}}|$ is a sinusoid of nearly constant amplitude but slowly increasing frequency $\nu_{\text{GW}} \approx 0 + \nu_{\text{GW}} dt$. It can be shown that in the absence of damping ($\tau_d \to \infty$), the effective duration of the resonant forcing, while the source drifts across a resonance, is $t_F \approx 1/\sqrt{4\nu_{\text{GW}}}$ (e.g. Rathore et al. 2005), yielding

$$t_F = 6.6 \left( \frac{M_{\text{ch}}}{M_\odot} \right)^{-5/6} \left( \frac{\nu_{\text{GW}}}{1 \text{mHz}} \right)^{-11/6} \text{yr.}$$

Here $M_{\text{ch}} = \eta^{3/5} M_{\text{bin}}$ is the chirp mass. Analytic solutions to eq. (4) can be found in two limiting cases: the saturated/steady-state case with constant forcing frequency ($t_F \gg \tau_d$) (Miser, Thorne & Wheeler 1973), and the undamped case ($t_F \ll \tau_d$) (Rathore et al. 2005). Expressing the damping time of a given stellar oscillation mode in terms of the “quality factor” $q_t = \omega_t \tau_d / \pi$, the saturation condition $t_F = \tau_d$ implies that steady-state is reached approximately for

$$M_{\text{ch}} \lesssim 0.35 \left( \frac{\nu_{\text{GW}}}{1 \text{mHz}} \right)^{-1} \left( \frac{q_t}{10^7} \right)^{-6/5} M_\odot.$$  

Steady-state will not be reached for $q_t \gg 10^6$. However, if the star is close to the GW source such that $|F_{\text{GW}}|$ is sufficiently large, the oscillations can grow nonlinear before reaching the steady-state limit. If so, the mode coupling to higher-order modes prohibits further growth, and the effective quality factor is greatly decreased.

3 SATURATED/STEADY-STATE LIMIT ($t_F \gg \tau_d$)

Assuming stationary GW forcing at a constant frequency ($|F| = |F|e^{i\omega t}$), the maximum steady-state displacement $x_{\text{max}}$ is

$$x_{\text{max}} = \frac{|F|}{\sqrt{\omega^2 - \omega^2_{\text{GW}}}} = \frac{8 \pi G}{c^3} \sum_{\ell = 0} q_r^2 \omega^2_{\ell, r} \nu_{\text{GW}}$$

or $x_{\text{max}} \approx |F|/\tau_d / \omega_{*} = \pi R_\odot q_t h$ in the limit $(\omega^2 - \omega^2_{\text{GW}})^2 \ll (\omega / \tau_d)^2$. In the steady-state solution, $E_m = 0$ and the cycle-averaged power of the external forcing ($W = \langle F \dot{x} \rangle$) equals the rate of heating ($\dot{Q} = \langle \dot{x}^2 \rangle / \tau_d$). Taking the limit $\omega_{*} \approx \omega_{**}$, the rate of work done in the steady-state case is $W_\text{s} = \langle F^2 \rangle / \tau_d$

$$W_\text{s} \approx \frac{\pi}{2} R_\odot^2 h q_t \omega^3_{*}.$$  

The cross-section for absorbing GWs is given by

$$\sigma_{\text{GW}} = \frac{M_{\text{bin}}(W_\text{s})}{4 \pi c^3} \approx \frac{8 \pi G}{c^3} M_{\text{bin}} R_\odot^2 \omega^2_{*, r}$$

where $\Phi_{\text{GW}} = (\epsilon^3/16\pi G)^2 h^2$ is GW flux incident on the star and $M_{\text{bin}}$ is the overlap with the normal mode expressed as a measure of the mass involved in the mode such that (Khosroshahi & Sobouti 1997)

$$M_{\text{bin}, x} \equiv \left( \int \eta \nu_h \mathcal{V} d^3 x \right) / \int \rho \xi^2 d^3 x$$

where $\xi = \xi_{\text{geom}}(r)$ is the displacement for a normal mode and $\mathcal{V} = \mathcal{V}(x^2 - y^2)$. The fractional energy flux removed from the incident GW, corresponds to a resonant ‘optical depth’ ($e^{-1}$). The “effective opacity” seen by the GWs in the steady-state limit is $\tau_{\text{eff}, r} \equiv \sigma_{\text{GW}} / \pi R_\odot^2$ or

$$\tau_{\text{eff}, r} \approx \frac{8 \pi G}{c^3} M_{\text{bin}}(q_t) \omega^3_{*} \approx 0.8 (\nu_{\text{GW}}/1 \text{mHz})^{-1} \left( \frac{q_t}{10^7} \right)^{-6/5} M_\odot.$$  

In general, computing overlap integrals ($M_{\text{bin}}$) between stellar modes and the GW forcing for realistic stellar structure models will be difficult, and will also be very sensitive to the details of stellar structure. A full investigation is beyond the scope.
of this Letter. However, overlap calculations exist for the somewhat similar case of Newtonian tidal forcing by a nearby point-source, both for simplified polytropes (e.g. Press & Teukolsky 1977; Reisenegger & Goldreich 1994) and for more realistic stellar models (Aerts et al. 2010). In general, these show that the lowest-order modes have large overlap integrals, between O(0.1)–O(1) for polytropes (see, e.g. Table 1 in Press & Teukolsky (1977)), but also that simple polytrope models are insufficient to estimate the excitation of g-modes in Sun-like stars (e.g. Weinberg et al. 2012). At higher citation of g-modes in Sun-like stars (e.g. Weinberg et al. 2012), that simple polytrope models are insufficient to estimate the ex-

 overlap integral between forcing of polytropes, showing that the fundamental mode has an overlap integral between 20%–40% for simple polytropic fluid models, with polytrope index 1.5 < n < 2.5 (see their Table 1, where the fundamental f-mode is labeled as p1). We find this result intuitively unsurprising, since the angular part of the overall integral, for ell=2 modes, matches the quadrupolar pattern of the GWs, and the radial integral is over the product of a non-oscillatory eigenmode and a slowly-varying GW forcing function. We conclude that the overlap integral for a number of low-radial-order ell=2 modes is likely to be significant, i.e. close O(0.1)–O(1) in at least a few cases, depending on the details of stellar structure. More sophisticated stellar modelling is needed to compute the overlap for g-modes, and also for non-solar type giant stars, where GW wavelength is closer to the stellar radius and where higher-order modes may have substantially greater overlap integrals. (Lopes & Silk 2014) calculate Solar models beyond a simple polytrope and find that values substantially greater overlap integrals (Lopes & Silk 2014) each with q1 ∼ 100–400, we find odds of ~ 1%–6% for each of the modes that a WDB could lie in the ecliptic plane at that frequency in a 3.3µHz wide bin. The overwhelming majority of WDBs will lie off the ecliptic, but the orbits of future space-based GW detectors may be chosen to allow the most promising eclipses to be observed. We estimate that eLISA will have an O(10%) chance of identifying a WDB near a Solar resonance with a deep O(0.1)–O(1) transit depth that a future space-based GW detector could observe. The chances increase significantly for higher-order modes with much broader resonances (Sist et al. 1993) but owing to their low overlap integrals and/or low q1, these modes will likely produce much shallower transits. GW absorption could also be detected as a result of transits by blotted stars in the ∼ 35,000 galactic nuclei within 50Mpc or extreme mass-ratio inspirals around Sgr A* (Amaro-Seoane et al. 2007). Conservatively assuming a 1% AGN rate, ∼ 350 active galaxies lie within the LISA search window for transits. From McKernan et al. 2012 2013, most of these AGN host SMBH-IMBH or SMBH-sBH (EMRI) binaries. For ∼ 10 such binaries in the LISA frequency window (∼ 0.1–few mHz) we estimate a probability of 0.01–1 that we would see one transit in a 10-year mission, assuming 0.01%–1% chance of a transit/AGN/yr (Beky & Kocsis 2013). EM study of such transits would be challenging but potentially detectable (e.g. McKernan & Yaqoob 1998; Turner & Miller 2009). However, the resonant driving by these systems lasts for only tF ≪ day (eq. 17) so the chance of GW absorption coinciding with EM transits will be negligible.

5 ECLIPSES OF GW SOURCES BY THE SUN OR STARS IN THE MBHB HOST GALAXY.

The opacities τeff,u(ω) are independent of D, but g-, f-, and p-mode frequencies for the Sun are coincidentally in the sensitivity band of the proposed eLISA instrument (Cutler 1998). The Sun could therefore annually eclipse eLISA GW sources located in the ecliptic plane – in particular, white dwarf binaries (WDBs) (Crowder & Cornish 2007). The efect would be a “shadow” within a complex GW diffraction pattern near the resonant frequency (since the GW wavelength exceeds 1AU). Using the Monte Carlo simulations of (Tumapu et al. 2008), we estimate (from their Fig. 12) that between 20 (SNR ≥ 5) and 5(SNR≥ 10) individually resolvable WDBs with LISA in a 1 year observation will lie in a 3.3µHz bin (corresponding to width dν/ν 1/100 around log(ν)/3.5 Hz), i.e. near prominent Solar modes. Given that the plane of the ecliptic is 1/360th of the sky, this gives a ~ 2%–6% chance that one such WDB will be occulted annually by a low-order p-mode of the Sun annually. Using the 3 largest mass Solar p-modes (listed in Table 1 in Cutler & Lindblom 1996), each with q1 ∼ 100–400, we find odds of ∼ 1%–6% for each of the modes that a WDB could lie in the ecliptic plane at that frequency in a 3.3µHz wide bin. (2013) calculate Solar models beyond a simple polytrope and find that values substantially greater overlap integrals. (Lopes & Silk 2014) calculate Solar models beyond a simple polytrope and find that values substantially greater overlap integrals. (Lopes & Silk 2014) each with q1 ∼ 100–400, we find odds of ∼ 1%–6% for each of the modes that a WDB could lie in the ecliptic plane at that frequency in a 3.3µHz wide bin. The overwhelming majority of WDBs will lie off the ecliptic, but the orbits of future space-based GW detectors may be chosen to allow the most promising eclipses to be observed. We estimate that eLISA will have an O(10%) chance of identifying a WDB near a Solar resonance with a deep O(0.1)–O(1) transit depth that a future space-based GW detector could observe. The chances increase significantly for higher-order modes with much broader resonances (Sist et al. 1993) but owing to their low overlap integrals and/or low q1, these modes will likely produce much shallower transits. GW absorption could also be detected as a result of transits by blotted stars in the ∼ 35,000 galactic nuclei within 50Mpc or extreme mass-ratio inspirals around Sgr A* (Amaro-Seoane et al. 2007). Conservatively assuming a 1% AGN rate, ∼ 350 active galaxies lie within the LISA search window for transits. From McKernan et al. 2012 2013, most of these AGN host SMBH-IMBH or SMBH-sBH (EMRI) binaries. For ∼ 10 such binaries in the LISA frequency window (∼ 0.1–few mHz) we estimate a probability of 0.01–1 that we would see one transit in a 10-year mission, assuming 0.01%–1% chance of a transit/AGN/yr (Beky & Kocsis 2013). EM study of such transits would be challenging but potentially detectable (e.g. McKernan & Yaqoob 1998; Turner & Miller 2009). However, the resonant driving by these systems lasts for only tF ≪ day (eq. 17) so the chance of GW absorption coinciding with EM transits will be negligible.

6 RESONANT GW-HEATING OF STARS.

A star orbiting near a merging MBHB (within ∼1 pc) can absorb a significant amount of resonant GW energy. The average undamped heating rate of a single mode during the passage through resonance is MmQa = 1/2 Mm FGW tF / 2. As an illustrative example, we consider an ML = 10^4 M⊙ IMBH separated by ≈ 15 r_g from Sgr A* (M1 = 4 × 10^8 M⊙), and a Sun-like star 10^3 r_g away, resonating with the GWs (at ν ≈ 0.3 mHz):

\[ M_m Q_a = 400 L_\odot \left( \frac{M_\odot}{M_m} \right)^{3/2} \left( \frac{R_\odot}{R_m} \right)^{-2} \left( \frac{D_*}{10^3 r_g} \right)^{-2} \times \left( \frac{M_\odot}{4 \times 10^6 M_\odot} \right)^{-1} \left( \frac{\mu}{10^4 M_\odot} \right)^{3/2} \left( \frac{\nu}{3.0 \text{mHz}} \right)^{7/2}. \]
The heating rate is high, but lasts only for $t_F \sim 1.5$ days for the fiducial parameters for a single mode. The corresponding energy dumped into the mode is

$$E_m = M_m Q_u = 10^{44} \text{erg} \left( \frac{M_m}{M_{\odot}} \right) \left( \frac{R_a}{R_{\odot}} \right)^2 \left( \frac{D_4}{10^{13} \text{yr}} \right)^{-2} \times \left( \frac{M_{\text{bin}}}{4 \times 10^6 M_{\odot}} \right)^{-4/3} \left( \frac{\nu}{10^4 \text{mHz}} \right)^{5/3}. \quad (16)$$

Fig. 1 shows the total energy $E_m$ deposited in a single resonant mode of a star for the fiducial values of $M_m = M_{\odot}, R_a = R_{\odot}, D_4 = 10^{13} \text{yr}$ for three different MBHBs. From Fig. 1 for a large overlap integral, up to $10^{45}$ erg can be deposited into a single mode of a star near an equal mass ($10^6 M_{\odot}, 10^6 M_{\odot}$) MBHB. This is $\sim 11$ orders of magnitude larger than the expected viscous heating of stars (Li, Kocsis & Loeb 2013). If this much energy can emerge on short timescales, the resonating star can act as a prompt electromagnetic signpost of incident GWs.

### 7 RATE OF DISCHARGE OF GW-CHARGED BATTERIES.

Stars can release $E_m$ either electromagnetically (EM) or via GW emission (equivalent to elastic scattering of incident GWs). The GW timescale $\tau_{\text{GW}} \approx 5\nu^2/(GE_m c^4)$ is EM timescales. Energy thermalized in radiation zones emerges on the thermal timescale $\tau_{\text{th}} \approx 10^7 \text{yr}$ for Sun-like stars, implying that massive stars will brighten on this long timescale. The fractional luminosity increase is limited to $E_m / E_*= \lesssim 1$. However, energy deposited in the convection zone emerges on timescale $\tau_{\text{conv}} \sim 10^6 \text{s}$, which may cause significant brightening (see below).

### 7.1 Resonant destruction of stars by GWs.

A star is completely disrupted when the total energy dumped into the star (eq. 10) becomes greater than the binding energy of the star $E_*=GM^2/2R_*$. In the limit $M_{\text{bin}} \sim M_*$, this happens at a radius $D_{\text{tid}} > r_a$ where

$$\frac{D_{\text{tid}}}{D_{\text{tid}}} = 0.03 \left( \frac{M_*}{M_{\odot}} \right)^{-1/6} \left( \frac{R_a}{R_{\odot}} \right)^{1/2} \left( \frac{\nu}{10^{13} \text{mHz}} \right)^{5/6} \left( \frac{\mu}{10^3 M_{\odot}} \right)^{1/2} \quad (17)$$

where $D_{\text{tid}} = R_c (M/M_*)^{1/3}$ is the tidal disruption radius. Solar type stars near MBHBs are thus disrupted by Newtonian tides well before destruction due to resonant GW absorption.

### 7.2 Near-field destructive effects.

Far from MBHBs ($r \gg a_{\text{bin}}$), $F(t)$ in eqn. 11 is dominated by GWs. Close to the MBHB, tidal forcing at frequency $\nu_{\text{GW}} = 2\nu_{\text{bin}}$ is added to $F(t)$ via the Newtonian quadrupole potential $F_{\text{NQ}} \sim G\mu_{\text{bin}}^2 R_c D_{\text{tid}}^{-5}$ and relativistic current dipole force $F_{\text{CD}} \sim (G/\mu_{\text{bin}}) R_c D_{\text{tid}}^{-4}$ (Misner, Thorne & Wheeler 1973) (see eqns. 2.15 in (Alu 2000) and eqn. 6.4 in Johnson-McDaniel et al. 2009). Compared to GW forcing,

$$\frac{F_{\text{NQ}}}{F_{\text{GW}}} \approx \left( \frac{\nu_{\text{GW}}}{0.02 \text{mHz}} \right)^{-4} \left( \frac{D_4}{10^{13} \text{yr}} \right)^{-4} \quad (18)$$

and

$$\frac{F_{\text{CD}}}{F_{\text{GW}}} \approx \left( \frac{\nu_{\text{GW}}}{0.02 \text{mHz}} \right)^{-3} \left( \frac{D_4}{10^{13} \text{yr}} \right)^{-3} \quad (19)$$

Tidal forcing on a star $10^3 D_4 \text{yr}$ from an MBHB is dominated by near-field effects at $\lesssim 0.02 D_{\text{tid}}^{-4}$ mHz. The heating from eqn. (15) scales in the near field as $|F_{\text{NQ}}|^2 / |F_{\text{GW}}|^2$ and $|F_{\text{CD}}|^2 / |F_{\text{GW}}|^2$.

### 8 ELECTROMAGNETIC OBSERVABLES:

GW heating of stars with a large radiative core causes modest structural changes and increase in luminosity, since $E_m \ll E_*$. For fully (or mostly) convective stars (e.g. M-stars), $E_m$ is transferred to high-degree modes concentrated in the outer convective skin, with small mass $M_{\text{out}} \ll M_*$ (Kumar & Goodman 1996). If $E_m > (M_{\text{out}}/M_*) E_*$, the binding energy of the surface skin, the skin can expand (Podsiadlowski 1996), provided $E_m$ is thermalized faster than $\tau_{\text{conv}}$, i.e. for $E_m \gtrsim 10^{35}$ erg. If $10^{-3} M_0$ is shed from a Sun-like star and subsequently accreted onto a $10^6 M_\odot$ MBHB over a $\sim$-year (or $\approx 10$ stellar orbits at $10^3 r_a$), the MBHB is fuelled at 0.01-0.1 its Eddington rate (for 10% radiative efficiency) (Hayasaki, Stone & Loeb 2013; Dai, Escala & Corna 2013). During accretion, the MBHB period appears as see-saw variability in the wings of broad emission lines (McKernan et al. 2013), possibly preceding tidal disruption events by MBHBs.

If $E_m \lesssim 10^{15}$ erg is thermalized, the star may not blow but the luminosity $L'_* = E_m / \tau_{\text{th}}$ can be large. For the $M_1 = M_2 = 10^6 M_\odot$ MBHB in Fig. 1, at $\approx 0.3$ mHz, a star at $D_4 \approx 10^3 r_a$ is heated for $t_F \approx 6$ hr, and $E_m \approx 2 \times 10^{43}$ erg emerges over the nonlinear dissipation timescale $\tau_{n} \approx 4$ yr. During this period $L'_* \approx 45 (M_m/M_*)^2 L_\odot$. Moreover, $\nu_{\text{GW}}$ sweeps through a large number of resonant modes $N_{\text{res}}$ between 40 mHz-12 mHz in the final 4 yr before merger (Aerts et al. 2010). If $N_{\text{res}} \sim 10$ modes can be driven resonantly within a dissipation timescale, then $L'_* \sim \times 10^{23} (N_{\text{res}}/10) (M_m/M_*)^{2} L_\odot$. 

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9 CONCLUSIONS.

Quadrupolar oscillation modes in stars can resonate with incident GWs, reaching non-linear amplitudes at the expense of GW energy. The opacity to GWs is distance-independent, so the Sun can eclipse GW sources (e.g. WDBs) in the ecliptic plane, imprinting absorption lines in GW spectra. Stars near MBHBs act as GW-charged batteries, discharging via a brief, significant luminosity increase in convective stars. Mass loss from the outer skin of stars yields bursts of near-Eddington accretion onto nearby MBHB. Detailed numerical studies (including models of stellar structure, effects of rotation) are needed for more quantitative predictions.

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