Center for Gravitational Physics and Geometry
The Pennsylvania State University
University Park, PA 16802-6300

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Ocean g-Modes on Rotating Neutron Stars

Lars Bildsten and Greg Ushomirsky

Department of Physics, Department of Astronomy and Space Sciences Laboratory
University of California, Berkeley, CA 94720
I: bildsten & gregus@fire.berkeley.edu

and

Curt Cutler

Department of Physics
Pennsylvania State University, University Park, PA 16803
I: curt@phys.psu.edu

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ABSTRACT

We recently argued that neutron stars accreting at rates $\gtrsim 10^{-9} M_{\odot}\ \mathrm{yr}^{-1}$ (the "Z" sources) are covered with massive oceans and conjectured that waves in these oceans might modulate the outgoing X-ray flux at frequencies comparable to what is observed. For slowly rotating neutron stars, we showed that the low radial order ($n \sim 1-2$) g-mode oscillations are in the 5-8 Hz range for $l=1$, in rough agreement with the ubiquitous ~ 6 Hz quasi-periodic oscillations (QPO's) seen in the Z sources. In this *Letter*, we extend the thermal g-mode calculations to the case of rapidly rotating stars. The $m=0$ modes are most relevant, since the $m \neq 0$ modes will necessarily acquire a high frequency in the observer's non-rotating frame. For spin frequencies $\Omega \gg \omega$, we find that the lowest $m=0$ g-mode frequency is $\omega = 2^{1/4}(\Omega\omega_0)^{1/2}$, where ω_0 is the corresponding non-rotating $l=1$ g-mode frequency. In the context of non-radial thermal g-modes, there are two ways to explain the fact that all six Z sources show ~ 6 Hz QPO's: (1) the neutron stars are all spinning at frequencies $\lesssim 6$ Hz and a low-order (i.e. only a few radial nodes) g-mode is responsible for the oscillations, or (2) the g-mode is of higher radial order ($n \sim 10-50$) and is brought to ~ 6 Hz by fast rotation, thus requiring that all six neutron stars have rotation frequencies within a factor of 2-3 of each other.

Subject headings: accretion – stars: neutron – stars: oscillations – stars: rotation – X-rays: stars

1. Introduction

It seems likely that neutron stars accreting at rates in excess of $10^{-9} M_\odot$ yr^{-1} (the so-called "Z" sources) burn the accreting hydrogen and helium in a way less violent than by Type I X-ray bursts. Both the observed lack of Type I X-ray bursts (see van der Klis's 1995 review) and the theoretical understanding of the nature of burning at these high accretion rates (Fushiki & Lamb 1987, Bildsten 1993, 1995) tell us that the hydrogen and helium burn to midweight elements (i.e. C, O, Ne, ...). Fusion of this fuel to iron-group elements does not occur until densities $\gtrsim 10^9$g cm^{-3} are reached, leading to the accumulation of a massive ($\sim 10^{-6} M_\odot$) degenerate liquid ocean (Bildsten & Cutler 1995, hereafter BC). Neutron stars accreting at lower rates burn the accreted hydrogen and helium directly to iron group elements via Type I X-ray bursts (Lewin, van Paradijs & Taam 1992). The larger Coulomb force in iron leads to crystallization at much lower densities and a much smaller ocean mass.

In our previous paper (BC) we showed that, for nonrotating stars, the ocean's thermal buoyancy results in g-modes (basically internal gravity waves) with frequencies of a few Hz for $l = 1$ and one or two radial nodes. The similarity between this frequency and the ~ 6 Hz quasi-periodic oscillations (QPO's) seen in the bright sources (see van der Klis 1995 and references therein) has motivated much of our work. The remarkable uniformity of these QPO's in all six known Z sources led Hasinger (1987) to conjecture that the frequency was set by a resonant density wave above the neutron star surface. Another model is that these QPO's involve oscillations in a coronal accretion flow (extending out to 100 stellar radii) at rates near the Eddington limit (Fortner, Lamb & Miller 1989 and Miller & Park 1995).

Our earlier results were limited to stars rotating slower than the mode frequency. However, one might expect that stars accreting at these high rates will be rotating more rapidly than a few Hz. In particular, if the accreting matter always arrives with the specific angular momentum of a particle orbiting at the stellar radius R, it takes $\sim 10^5$ yr of accretion at $\dot{M} \approx 10^{-8} M_\odot$ yr^{-1} to reach a spin frequency of $\sim$ 10Hz. Even though such spin periods have not been detected in these systems (Vaughan et al. 1994), we feel it is crucial to understand the influence of rotation on the mode frequencies.

This *Letter* is a brief report on how the thermal g-mode frequencies and displacements are modified by rapid rotation. The equations for non-radial oscillations in rotating neutron stars are derived in §2. For the case we are considering, the angular and radial equations separate, with the radial equations being identical to those for the non-rotating case. The angular piece is solved for the connecting eigenvalue in §3. We discuss the observational implications of our results in §4.

2. Non-Radial Oscillations in the Presence of Rotation

We consider a neutron star uniformly rotating (see Fujimoto (1993) for a discussion of the angular momentum transport that brings about uniform rotation) at frequency $\Omega \ll (GM/R^3)^{1/2}$ where M and R are the neutron star mass and radius. This allows us to neglect the "centrifugal term" in the equation of motion and treat the unperturbed star as spherical. Since the non-rotating g-mode frequencies are far below the breakup frequency of $\approx$ 1kHz, this approximation still leaves a large range of spin frequencies to explore. The g-modes are confined to the thin ($h \ll R$) ocean on the neutron star surface and do not penetrate into the neutron star's crust (BC). In addition, the small mass of the ocean allows us to neglect the perturbations to the gravitational potential.

We work in the rotating frame, where the fluid is initially at rest, and denote the Lagrangian fluid displacement as $\vec{\xi}$. The momentum balance equation is then

$$i\omega\frac{d\vec{\xi}}{dt} = -\frac{\vec{\nabla}p}{\rho} - g\hat{r} - 2i\omega\vec{\Omega}\times\vec{\xi}, \qquad (1)$$

where the last term is the "Coriolis force", $g = GM/R^2$ is the constant downward acceleration of gravity, and ω is the mode frequency in the rotating frame. We let δ denote the Eulerian variation in some quantity Q, so that the Lagrangian variation is $\Delta Q = \delta Q + \vec{\xi}\cdot\vec{\nabla}Q$. We use spherical coordinates, with $\vec{\Omega}$ parallel to the axis where $\theta = 0$, so that $\vec{\xi} = \xi_r\hat{r} + \xi_\theta\hat{\theta} + \xi_\phi\hat{\phi}$. Since the perturbation equations are linear, and the background is time-independent and spherically symmetric, we decompose solutions into the form $\delta Q(r,\theta)e^{im\phi+i\omega t}$. Frequencies in the observer's inertial frame are then $\omega_I = \omega - m\Omega$.

Given the above assumptions, the perturbed momentum balance equations are

$$\rho\omega^2\xi_r = \frac{\partial\delta p}{\partial r} + g\delta\rho - 2i\Omega\omega\rho\xi_\phi\sin\theta. \qquad (2a)$$

$$-\rho\omega^2\xi_\theta = -\frac{1}{R}\frac{\partial\delta p}{\partial\theta} + 2i\Omega\omega\rho\xi_\phi\cos\theta. \qquad (2b)$$

$$-\rho\omega^2\xi_\phi = -\frac{im\delta p}{R\sin\theta} - 2i\Omega\omega\rho(\xi_\theta\cos\theta + \xi_r\sin\theta), \qquad (2c)$$

where we have replaced $1/r$ with $1/R$ as the ocean is very thin ($h \ll R$). The oscillations are adiabatic at $\rho \gtrsim 10^5$ g cm^{-3} so that the Lagrangian density and pressure variations are related by $\Delta p/p = \Gamma_1\Delta\rho/\rho$, where $\Gamma_1 \equiv (d\ln p/d\ln\rho)_{ad}$ is the adiabatic index. We use this relation to eliminate $\delta\rho$ in equation (2a), thus obtaining

$$\rho(\omega^2 - N^2)\xi_r = \frac{\partial\delta p}{\partial r} + \frac{\delta p}{h\Gamma_1} - 2i\Omega\omega\rho\xi_\phi\sin\theta, \qquad (3)$$

where $h = p/\rho g$ is the local scale height, and $N^2 = -\mathcal{A}g$ is the local Brunt-Väisälä frequency, where $\mathcal{A} \equiv d\log\rho/dr - d\log p/dr(1/\Gamma_1)$ is the convective discriminant.

Equation (3) exhibits a competition between the radial Coriolis force (coming from the transverse motion) and the buoyancy force ($\propto \rho N^2\xi_r$). We estimate the importance of this new term by using the slow rotation solutions as our first guess, in which case $\xi_\phi/\xi_r \sim R/h$ and $N^2 \gg \omega^2$ (BC), as the thermal buoyancy in the deep ocean gives $N \sim 1-5$ kHz (BC). We can thus neglect the Coriolis term in equation (3) when $|N^2| \gg R\Omega\omega/h$ (which limits the spin frequencies to $\lesssim$200 Hz; see end of §3) so this equation becomes

$$\rho(\omega^2 - N^2)\xi_r = \frac{\partial\delta p}{\partial r} + \frac{\delta p}{h\Gamma_1}. \qquad (4)$$

In addition, $\xi_r \ll \xi_\theta$ for the low l oscillations, so that we omit the term $\propto \xi_r$ in equation (2c).

These two approximations (referred to as the "traditional approximation" in the geophysical literature) dramatically simplify the solution of the perturbation equations as the radial and angular parts completely separate (i.e. $Q = Q_r(r)Q_\theta(\theta)$). To show this, we first combine equations (2c) and (2b) to find ξ_ϕ and ξ_θ in terms of δp and ξ_r. We then substitute these into the continuity equation to obtain

$$-\frac{\partial\xi_r}{\partial r} - \frac{1}{\Gamma_1}\left(\frac{\delta p}{p} - \frac{\xi_r}{h}\right) = \frac{gh}{\omega^2R^2}L_\mu\left[\frac{\delta p}{p}\right], \qquad (5)$$

where L_μ is the following angular operator that depends on the spin parameter $q \equiv 2\Omega/\omega$,

$$L_\mu \equiv \frac{\partial}{\partial\mu}\left(\frac{1-\mu^2}{1-q^2\mu^2}\frac{\partial}{\partial\mu}\right) - \frac{m^2}{(1-\mu^2)(1-q^2\mu^2)} - \frac{qm(1+q^2\mu^2)}{(1-q^2\mu^2)^2}, \qquad (6)$$

with $\mu \equiv \cos\theta$. We thus have two equations [(4) and (5)] to solve for $\delta p/p$ and ξ_r. The resulting equations are separable: we first solve eigenvalue equation $L_\mu(\delta p/p) = -\lambda\delta p/p$, and then solve the purely radial equations (4)and (5), with the right hand side of equation (5) replaced by $-gh\lambda/(\omega^2 R^2)$. These radial equations are identical to those for a nonrotating star, in which case λ is just $l(l+1)$. In some sense, one can view $\lambda^{1/2}$ as the effective transverse wavenumber (i.e. $k_{tr}^2 = \lambda/R^2$).

We have solved (in BC) the radial equations for g-modes in a degenerate (i.e. $\rho \gtrsim 10^7$ g cm^{-3}) ocean of liquid ions with mass Am_p at temperature T, where the restoring force is thermal buoyancy, so we can simply write down the eigenvalue and eigenfunctions. The g-mode frequencies are

$$\begin{aligned} f_{\lambda,n} &= 7.23 \text{ Hz} \left(\frac{\lambda}{2}\frac{T}{10^8 \text{ K}}\frac{12}{A}\right)^{1/2} \\ &\times \left(\frac{10 \text{ km}}{R}\right)\left(\frac{1}{1+0.47n^2}\right)^{1/2}, \end{aligned} \quad (7)$$

where n is the number of radial nodes in ξ_r. This exhibits the familiar g-mode property of having the frequency decrease as n increases, and it goes to $f \propto 1/n$ for large n. We now solve for the eigenvalues λ of the angular operator L_μ, which, using equation (7), will immediately give us the mode frequencies.

3. Numerical Solution of The Angular Equation

The angular eigenvalue equation

$$L_\mu(H) = -\lambda H \quad (8)$$

is known as Laplace's tidal equation, and its solutions are called Hough functions. The standard method (see Chapman & Lindzen 1970) of solving Laplace's tidal equation has been to expand H in spherical harmonics (i.e. Y_{lm} for a fixed m). This yields recursion relations for the expansion coefficients, which are truncated at some large but finite l. We have found it simpler to find the eigenfunctions and eigenvalues of this ODE by direct numerical integration; the only "trick" is to take appropriate care at the singular points ($\mu = \pm 1$ and $\mu = \pm 1/q$).

The equation is symmetric under $\mu \to -\mu$, so we restrict attention to the range $0 \le \mu \le 1$ and look for solutions with definite parity: the "odd" ("even") solutions have $H(0) = 0$ ($H'(0) = 0$). For all $q \neq 1$, the tidal equation has a regular singular point at $\mu = 1$ (see Bender & Orszag 1978). If we substitute a Frobenius series $H(\mu) = \sum_{n=0}^{\infty} a_n(\mu-1)^{n+\alpha}$ into the tidal equation, expand the equation in powers of $(\mu - 1)$ about $\mu = 1$, and equate coefficients, we find that the regular (non-singular) solution has $\alpha = |m|/2$. The equation becomes more tractable numerically if we factor out this behavior. With the substitution

$$H(\mu) \equiv (1-\mu^2)^{|m|/2} y(\mu), \quad (9)$$

the tidal equation becomes

$$\begin{aligned} (1 &- \mu^2)y''(\mu) \\ &+ \left(\frac{2q^2\mu(1-\mu^2)}{1-q^2\mu^2} - 2\mu(1+|m|)\right) y'(\mu) \\ &+ \left(\lambda(1-q^2\mu^2) - (|m|+m^2)\right. \\ &- \left.\frac{2|m|q^2\mu^2 + qm(1+q^2\mu^2)}{1-q^2\mu^2}\right) y(\mu) = 0, \end{aligned} \quad (10)$$

where $y(\mu)$ is analytic at $\mu = 1$.

For $q > 1$, there is an additional regular singular point at $\mu = 1/q$. Careful analysis via Frobenius series shows that the two linearly independent solutions of (10) are analytic at $1/q$, and thus can be represented by a Taylor series. We find the appropriate behavior of y near the singularity by writing $y = y_0 + y_1 x + y_2 x^2/2 + y_3 x^3/6 + \ldots$ where $x \equiv (\mu - 1/q)$, substituting into the tidal equation (10), and collecting terms of same order in x. The first two orders, $O(x^{-1})$ and $O(x^0)$, yield the same condition:

$$y_1 + \frac{q(|m| + mq)}{1-q^2} y_0 = 0, \quad (11)$$

while $O(x^1)$, upon using (11), reduces to

$$\begin{aligned} y_3 &- q\frac{(mq-q^2-3(|m|+1))}{1-q^2}y_2 \\ &- q^3\left[\frac{(|m|+mq)(1+2m^2-mq-q^2+3|m|)}{(1-q^2)^2}\right. \\ &- \left.\frac{4\lambda(1-q^2)}{(1-q^2)^2}\right]y_0 = 0. \end{aligned} \tag{12}$$

Subsequent orders yield recursion relations for y_n, for $n \geq 4$, in terms of y_{n-1}, y_{n-2}, etc. From this, we see that y_0 and y_2 can be considered "unconstrained" or "free" data, which completely determine y_1, y_3, y_4, etc. Thus, to integrate away from $\mu = 1/q$ towards either 1 or 0, we start at $\mu = 1/q \pm \epsilon$, with $y(1/q \pm \epsilon) = y_0 \pm y_1\epsilon + \frac{1}{2}y_2\epsilon^2$ and $y'(1/q \pm \epsilon) = y_1 \pm y_2\epsilon$, where y_1 is determined from y_0 by (11). The "shape" of the solution is set by the ratio y_2/y_0: y_0 just determines the overall scale.

For $q < 1$ we shoot from $\mu = 1-\epsilon$ (relating $y'(1)$ to $y(1)$ by demanding that $(1-\mu^2)y''(\mu) = 0$ as $\mu \to 1$ in equation (10)) and vary the eigenvalue λ to find the odd and even modes. For $q > 1$, we: (1) pick a matching point $\mu_{\rm ma}$ in the interval $1/q < \mu_{\rm ma} < 1$, (2) choose trial values for y_2/y_0 and λ, and (3) shoot towards $\mu = \mu_{\rm ma}$ from both $\mu = 1-\epsilon$ and $\mu = 1/q+\epsilon$. The function y can always be made continuous at $\mu_{\rm ma}$ by scaling y_0. It remains to match y' at $\mu_{\rm ma}$, while satisfying either $y(0) = 0$ or $y'(0) = 0$ (for odd or even solutions, respectively). Thus there are two conditions to satisfy, with two free variables. Ushomirsky & Bildsten (1995) contains a complete discussion of the solutions of this equation for all modes (including the r-modes). Here, we briefly summarize the properties of the relevant g-modes.

As is evident from the transverse momentum equations [(2b)and (2c)], the Coriolis force strengthens as one moves towards the pole. For $q > 1$, this introduces a critical angle $\mu = 1/q$.

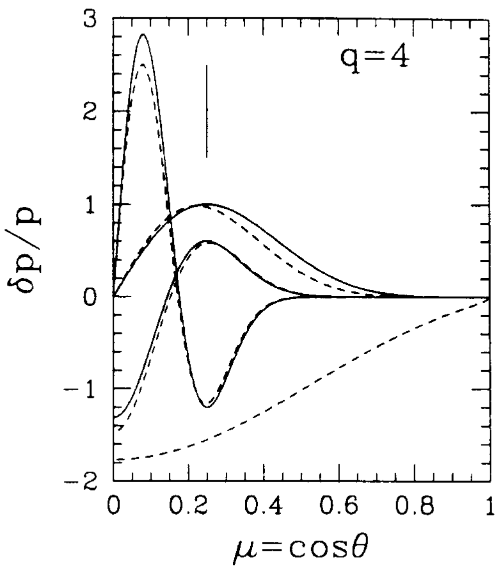

Fig. 1.— Angular eigenfunctions of equation (8) for $q \equiv 2\Omega/\omega = 4$. The solid lines show the first two odd eigenfunctions (with $\lambda = 16.54$ and 394.5) and the first even eigenfunction ($\lambda = 142.5$) for $m = 0$. The dashed line with large amplitude at the pole is the $m = -1$ even parity mode with $\lambda = 1.135$. The other three dashed lines are analogs of the $m = 0$ modes in the angular coordinate. In particular, we plot the two lowest $m = -1$ odd parity eigenfunctions ($\lambda = 25.93$ and $\lambda = 405.1$) and the lowest $m = 1$ even parity solution ($\lambda = 135.9$).

Poleward of this angle, the Coriolis force dominates, while equatorward of this angle, the pressure term dominates. This confines the propagating g-modes to the equatorial region. Figure 1 shows some eigenfunctions with $\lambda > 0$ for $q = 4$; note they evanesce above the critical angle $\mu = 1/q$ (marked by the vertical arrow). The solid lines show the first three $m = 0$ modes, while the dashed lines show eigenfunctions for $m = \pm 1$.

The discussion of the $\lambda > 0$ eigenfunctions is simplified for very fast rotation ($q \gg 1$). In this limit, almost all of the motion is confined to a thin band about the equator of dimension $L \approx 2R/q$ (i.e. most of the motion is in the latitudes $-1/q < \mu < 1/q$). The eigenvalues then arise from "fitting" the required number of oscillations into this angular band, giving the scaling

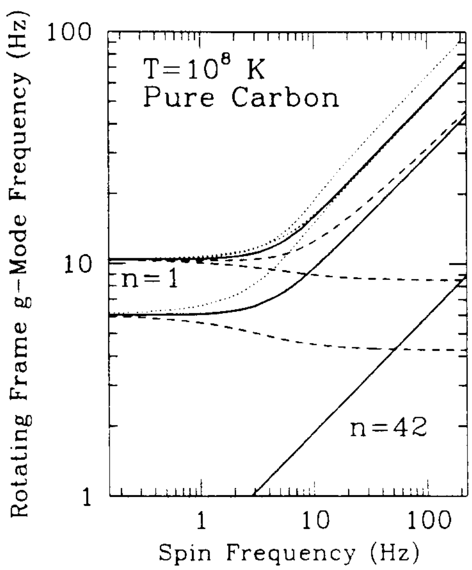

Fig. 2.— The thermal g-mode frequencies for $n = 1$ *in the rotating frame* for a pure carbon ocean at $T = 10^8$ K. The $m = 0$ modes are shown by the heavy solid line and have the same frequency in the observer's frame. The $m < 0$ modes are shown by the dashed lines. The $m = -1$ and $m = -2$ even parity modes go to a low and constant frequency in the rotating frame for arbitrarily high spins. The positive m modes are shown by the dotted curves. The lower solid curve is the $m = 0$, $l_\mu = 1$ mode for $n = 42$, which gives 6 Hz at a 100 Hz spin frequency.

$\lambda \propto q^2$. If we denote l_μ as the number of zero crossings between $-1/q$ and $1/q$, we find that the $m = 0$ eigenvalues asymptote to $\lambda = (2l_\mu - 1)^2 q^2$ when $q \gg 1$. Hence the lowest mode is $l_\mu = 1$ (odd) which gives $\lambda \approx q^2$ for $q \gg 1$. For non-zero m there is an additional possibility, which is that the angular eigenfunction is nearly constant and has no zero crossings. In this case, the eigenvalues stay low and approach m^2 as $q \to \infty$. This eigenfunction is shown in the lowest dashed line on Figure 1. The mode frequencies are the two dashed curves in Figure 2 that asymptote to constant values at high spin frequencies.

Equation (7) lets us write the mode frequency in the rotating frame simply as $\omega^2 = \omega_{l,0}^2(\lambda/l(l+1))$, where $\omega_{l,0}$ is the g-mode frequency for the corresponding l'th mode in the *non-rotating* star. In the limit of $q \gg 1$, this gives

$$\omega^2 = 2\Omega\omega_{l,0}\left[\frac{(2l_\mu - 1)^2}{l(l+1)}\right]^{1/2}, \qquad (13)$$

which approaches $\omega = 2(\Omega\omega_{l,0})^{1/2}$ in the WKB limit of $l = l_\mu \gg 1$, in agreement with Papaloizou & Pringle's (1978) WKB result. Figure 2 shows the frequencies *in the rotating frame* for the $n = 1$ thermal g-mode in a 10^8 K pure Carbon ocean. The heavy solid lines show the $l_\mu = 1, 2$, $m = 0$ modes, which have the same frequencies in the observer's frame. The dashed (dotted) lines show the modes for negative (positive) m, which split away from the $m = 0$ mode as the spin increases, even in the rotating frame. The frequency in the observer's frame is $\omega_I = \omega - m\Omega$.

We now use our solutions to check our earlier approximation of neglecting the Coriolis term (relative to $\rho N^2 \xi_r$) in the radial momentum equation. For the $l_\mu = 1$ thermal g-mode, we find that the Coriolis term is negligible when $q^2 \lesssim nR/2h$, which translates into $\Omega/2\pi \lesssim N/8\pi \approx 200$ Hz, regardless of n.

4. Discussion

We have shown that rapid rotation plays an important role in setting the g-mode frequencies. Our results are best explained in the rotating frame, where we found the simple scaling law of equation (13) for the mode frequencies in the limit of rapid rotation. This scaling was earlier found by Papaloizou & Pringle (1978) in the WKB limit (i.e. $l, n \gg 1$) and we have extended it to the low values of l which are observationally relevant. For rapid rotation ($\Omega \gg \omega$) the $m = 0$, $l_\mu = 1$ mode frequency (same in either the observer's frame or the rotating frame) is $\omega = 2^{1/4}(\Omega\omega_0)^{1/2}$, where ω_0 is the $l = 1$ g-mode frequency for no rotation. Hence, even at very high spin frequencies, the $m = 0$ modes can be observed at frequencies considerably less than the spin frequency. The high spin frequency, low n thermal g-modes coincide with the 15-50 Hz

QPO's seen from three of the Z sources. However. the observed frequency changes are more rapid than could be accommodated by heating or cooling in the ocean, making it seem unlikely that the 15-50 Hz QPO's arise from thermal g-modes.

If we wish to explain the observed $\sim$ 6 Hz QPO as an $m = 0$ non-radial g-mode supported by *thermal buoyancy*, we have two choices: (1) either all six known Z sources rotate at $\lesssim$6 Hz and the observed g-mode is $n \sim 1$, or (2) the mode responsible for the 6 Hz oscillations is of much higher radial order and rapid rotation brings the observed frequency to 6 Hz. To keep the frequencies within the observed range ($5-8$ Hz) requires spin frequencies within a factor of $\approx 2-3$ of each other. For example. for a neutron star rotating at 100 Hz. the non-rotating mode frequency would need to be 0.25 Hz to give an observed frequency of 6 Hz. This must be a thermal mode of high radial order. roughly $n \approx 10 - 50$ depending on the exact temperature and elemental composition. The lower solid line in Figure 2 exhibits such a mode. The mode structure would be "dense" in this rapid rotation limit. as the observed frequencies scale like $n^{-1/2}$. These modes have all of their amplitudes at the equator. which might help in modulating the accretion rate from the disk and hence the X-rays.

The way the mode frequencies vary with l or l_μ (for $m = 0$) for rotating stars is almost the same as in the non-rotating case. In the limit of no rotation. the ratios are $\omega_{l=1}^2 : \omega_{l=2}^2 : \omega_{l=3}^2 = 1 : 3 : 6$. while the limit of rapid rotation gives these ratios as $1 : 3 : 5$. We emphasize again that observing frequencies with these ratios is an important test of our hypothesis. Since $\omega_I = \omega - m\Omega$. the $m \neq 0$ modes are shifted to high frequencies by the rotation. Measuring a few of these modes would measure the spin frequency.

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