CONDITIONS FOR STEADY GRAVITATIONAL RADIATION FROM ACCRETING NEUTRON STARS

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ABSTRACT

The gravitational wave- and accretion-driven evolution of the angular velocity, core temperature, and (small) amplitude of an r-mode of neutron stars in low-mass X-ray binaries and similar systems is investigated. The conditions required for evolution to a stable equilibrium state (with gravitational wave flux proportional to average X-ray flux) are determined. In keeping with conclusions derived from observations of neutron star cooling, the core neutrons are taken to be normal while the core protons and hyperons and the crust neutrons are taken to be singlet superfluids. The dominant sources of damping are then hyperon bulk viscosity (if much of the core is at least 2–3 times nuclear density) and (e⁻e and n-n) shear (and possibly magnetic) viscosity within the core-crust boundary layer. It is found that a stable equilibrium state can be reached if the superfluid transition temperature of the hyperons is sufficiently small (≤2 × 10⁹ K), allowing the gravitational radiation from Sco X-1 and several other neutron stars in low-mass X-ray binaries to be potentially detectable by the second-generation LIGO (and VIRGO) arrays.

Subject headings: dense matter — gravitational waves — stars: neutron — X-rays: binaries

In this Letter we present the results of an investigation of the evolution of rapidly rotating accreting neutron stars under the influence of their emission of gravitational radiation. We employ a modification (Wagoner, Hennawi, & Liu 2001) and extension of the two-component (equilibrium+p-perturbation) model of the star introduced by Owen et al. (1998) and also employed by Levin (1999), but we restrict the analysis to small perturbations. These are assumed to be in the form of r-modes (Andersson 1998; Friedman & Morsink 1998; Lindblom, Owen, & Morsink 1998; Andersson, Kokkotas, & Schutz 1999), which radiate mainly via Coriolis-driven velocity perturbations rather than the density perturbations of the less powerful f-modes. We must allow for large uncertainties in many of the relevant properties of neutron stars, such as the superfluid transition temperatures and the properties of the core-crust boundary layer. Details will be presented in a subsequent paper.

After developing a general formalism, we focus on conditions in which the neutron star angular velocity (and thus gravitational wave frequency) evolves slowly toward an equilibrium state, in which the rate of accretion of angular momentum from the surrounding disk is balanced by its rate of loss via gravitational radiation. If this equilibrium is achieved, the observed flux of gravitational radiation can be shown to be proportional to the observed flux of X-rays from the accretion (Wagoner 1984; Bildsten 1998).

One of our longer term goals is the development of parameterized expressions describing possible time evolutions of the gravitational wave frequency and amplitude to facilitate detection by LIGO, VIRGO, and similar interferometer detectors. The brightest low-mass X-ray binaries (LMXBs) thought to contain a neutron star are the prime targets.

In this exploratory investigation, it is sufficient to consider a Newtonian neutron star in equilibrium (with equatorial radius R) that is perturbed by a nonaxisymmetric infinitesimal fluid displacement ξ = f(r, θ)δρ(r, θ) δ(r, θ) ≈ αR, with α ≪ 1. Based on the work of Friedman & Schutz (1978a) and Levin & Ushomirsky (2001a), the total angular momentum J of the star can be decomposed into its equilibrium angular momentum J_e and a perturbation proportional to the canonical angular momentum J_c. That is, \( J = J_e + (1 - K_e)J_c \), with \( J_e = -\alpha J_c \) M the mass, and \( \Omega \) the (uniform) angular velocity of the equilibrium star. All constants \( K \) are dimensionless, with \( K_e \sim K_c \sim 1 \).

In classical mechanics, the action \( I = E/\omega \) of any normal mode of a set of oscillators (with frequency \( \omega \)) is an adiabatic invariant. For a fluid, the analogous quantity should be \( E \), where \( E \) is the canonical energy of the perturbation in the co-rotating frame and \( \omega = \sigma + m\Omega \) is its frequency in that frame. However, we also have the general relation \( E = -(\omega/m)J_c \) (Friedman & Schutz 1978a). Therefore, following Ho & Lai (2000), we assume that the canonical angular momentum is also an adiabatic invariant and should therefore be unaffected by the slow rate of mass accretion. Thus, it obeys the usual relation (Friedman & Schutz 1978b)

\[
\frac{dJ}{dt} = 2J_e [F_e - F_e (M, \Omega, T_e)],
\]

where \( F_e \) is the gravitational radiation growth rate and \( F_e \) is the viscous damping rate. The latter also depends upon a spatially averaged temperature \( T_e \).

On the other hand, conservation of total angular momentum requires that

\[
\frac{dJ}{dt} = 2J_e F_e + J_c (t),
\]

where \( J_c = \dot{J}_c \) is the rate of accretion of angular momentum. The mass is accreted with specific angular momentum \( j_c \) at a rate \( M(t) \).

Combining these equations then gives the dynamical evolution relations

\[
\frac{1}{J_c} \frac{d\alpha}{dt} = F_e - F_e + [K_e F_e + (1 - K_e) F_e] \alpha^2 - \frac{J_e}{J_c^2} \frac{\dot{M}(t)}{t},
\]

\[
\left( \frac{I_c}{J_c} \right) \frac{d\Omega}{dt} = -2K_e F_e + (1 - K_e) F_e) J_c \alpha^2 + \frac{J_e - \dot{J}_c}{J_e} \frac{\dot{M}(t)}{t},
\]

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where \( I_s(M, \Omega) = \partial J_s / \partial \Omega \) and \( j_o(M, \Omega) = \partial J_o / \partial M \). In keeping with the fact that it is sufficient to also work to lowest order in \( \Omega/\Omega_{\text{max}} \) (as well as the relativity parameter \( GM/Rc^3 \)) in this exploratory investigation, we take \( J_s - j_o \approx j_s \) and \( J_o \approx I_o \Omega \). We also note that \( K_s \approx 0.094 \) (Owen et al. 1998). (The value of \( K_s \) is unimportant, since we will see that \( F \) remains very close to \( F_0 \).

Finally, thermal energy conservation for the entire star gives our third evolution equation:

\[
\int \frac{\partial T}{\partial t} c_s \, dV = C(T) \frac{dT}{dt} \approx 2 \tilde{E}_c F_c(T_c) + K_n(M) c^2 - L_c(T_c),
\]

where the corotating frame canonical energy \( \tilde{E}_c = K_n \Omega I_s \alpha_s^2 \), with \( K_n = K_o/3 \). Since all constituents (electrons, nucleons, etc.) are degenerate, the specific heat at constant volume \( c_v \) (and correspondingly \( c_p \)) is essentially the same as that at constant pressure. The terms on the right-hand side of this equation represent viscous heating, pycnonuclear reactions and neutron emissions in the inner crust (proportional to a time-averaged mass accretion rate), and neutrino luminosity. The hydrogen-helium-burning rate is assumed to be balanced by the surface emission of photons (Schatz et al. 1999), especially at the large accretion rate \( M = 10^{-8} M_\odot \text{yr}^{-1} \) (roughly \( \frac{1}{2} \) the Eddington rate, appropriate to our primary targets) that we adopt. The mass accretion rate can be estimated from accretion energy conservation. The photon luminosity arising directly from the accretion is \( L_{\text{acc}} \approx (GM/R)(M/t) \) for a slowly rotating neutron star with a negligible magnetosphere.

From now on we take the perturbation to be due to the dominant \( l = m = 2 \) r-mode, in which case the gravitational wave frequency \( f_{gw} = (4/3)f_{cs} = 20/3 \pi \). In order to facilitate comparison with previous results, we adopt the neutron star model of Owen et al. [1998; \( p \approx \rho^2, \ M = 1.4 \ M_\odot, \text{and} \Omega = (\pi G(p))^{1/2} = 8.1 \times 10^4 \text{rad s}^{-1} \)] for numerical work. Then, the gravitational radiation growth rate of this mode is \( F_s = \Omega^2/\tau_{\gamma r} \), with \( \tau_{\gamma r} \approx 3.26 \text{ s} \) and \( \Omega = \Omega/\Omega_c \).

Comparison of observations of thermal emission from isolated neutron stars with computed cooling histories (Tsuruta et al. 2000) has led Kaminker, Yakovlev, & Gnedin (2002) to conclude regarding the most likely state of the constituents, which we tentatively adopt. Specifically, the maximum values of the (density dependent) superfluid transition temperatures are taken to be (1) \( T_s \leq 10^5 \text{ K} \) for the (triplet) core neutrons (so they will be normal), (2) \( T_o \approx 5 \times 10^5 \text{ K} \) for the (singlet) core protons, and (3) \( T_s \approx 5 \times 10^6 \text{ K} \) for the (singlet) inner crust neutrons. In what follows we also assume that the thermal conductivity timescales \( \tau_{\gamma} \) are short enough to give relations \( T_s(T_c) \) and \( T_o(T_c) \) between these three spatially averaged temperatures that appear in equation (5). Typically, \( \tau_{\gamma} \approx 0.1-1 \text{ yr} \) in the core and \( \tau_{\gamma} \approx 10-100 \text{ yr} \) in the inner crust (Brown 2000; Haensel, & Yakovlev 2001).

Then, in the temperature range of interest (10^8 \text{ K} \leq T \leq 10^9 \text{ K}), the viscous damping rate of this mode is \( F_s = F_{cs}(T_c) + F_{cs}(\Omega, T) + F_{cs}(\Omega, T) \). Even if the neutrons were not normal, the contribution to this damping from the mutual friction between a neutron superfluid and the superconducting proton-relativistic electron fluid (Lindblom & Mendell 2000) would be negligible. Therefore, we have neglected damping due to energy extraction from the mode to magnetic fields, which requires that \( B \approx 10^{11} \text{ G} \) (Rezzolla, Lamb, & Shapiro 2000; Rezzolla et al. 2001).

The first term in the damping rate is produced by the ordinary shear viscosity throughout the core. Adding the contribution of the \( n-n \) scattering (Owen et al. 1998) to that of the \( e-e \) scattering (Lindblom & Mendell 2000) gives \( F_{\text{sh}} = 1/(\tau_0 T_0^2) \), where \( \tau_0 \approx 0.69 \times 10^6 \text{ s} \) and \( T_0 = T/10^{10} \text{ K} \).

The second term is produced by the ordinary shear and magnetic viscosity in the crust-core boundary layer. For a normal fluid, this damping rate has been calculated by Andersson et al. (2000), Bildsten & Ushomirsky (2000), Lindblom, Owen, & Ushomirsky (2000), Bildsten & Ushomirsky (2000), and Wu, Matzner, & Arras (2001) neglecting magnetic effects, while Bildsten & Ushomirsky (2000), Mendell (2001), and Kinney & Mendell (2002) included them. For superfluid neutrons and protons, this rate was calculated by Kinney & Mendell (2002). We approximate the damping rate for our case (normal neutrons and superfluid protons in the boundary layer) by adding the rate of Kinney & Mendell (2002) for a normal uncharged fluid to their superfluid rate when only the proton vortexes (magnetic flux tubes) are pinned to the crust. The magnetic damping is due to the interaction of the field with the effective current density of the proton vortexes, producing cyclotron-vortex waves. Although there is no mutual friction because no neutron vortexes are present, we also neglect the small coupling between these fluids produced by the \( n-e \) scattering. We then obtain (after an approximate algebraic simplification)

\[
F_{\text{sh}} = 3.15 \times 10^{-2} S_{n_s} \frac{\sqrt{\Omega}}{T_s} + 3.71 \times 10^{-8} S^2 \frac{\sqrt{B}}{T_s} \text{ s}^{-1},
\]

where \( B \) (in units of gauss) is the radial magnetic field. The slippage factor \( S \), is the fractional degree of pinning of the vortexes in the crust (Kinney & Mendell 2002), and \( S_{n_s} = (2S_e + S_p)/3 \), with the slippage factor \( S_p \), the fractional difference in velocity of the normal fluid between the crust and the core (Levin & Ushomirsky 2001b). Although both slippage factors were defined to be at most unity, we also let them contain the uncertainties in our model. (If the pinning were dominated by neutron vortexes, the coefficients in eq. [6] would be \( (2-6) \times 10^3 \) times larger, resulting in no growth of the mode unless \( S_p \ll 1 \).) In obtaining most of the numerical results below, we take the magnetic contribution in equation (6) to be negligible, which requires that \( B \approx 0 \times 10^{10} \text{[1 + (2(S_e/S_p))^2}] \Omega T_s^{-2} \text{ G} \).

The third term arises from the bulk viscosity produced by out-of-equilibrium hyperon reactions (which dominate that produced by direct and modified Urca reactions). This has been studied by Jones (2001a, 2001b) and Lindblom & Owen (2002) for normal nuclear matter and by Haensel, Levenfish, & Yakovlev (2002) for superfluids. We employ the results of Lindblom & Owen (2002) for \( n + n \approx p + \Sigma \). However, with our assumptions about superfluidity, the reaction whose rate is least reduced by superfluid phase-space blocking (but is more difficult to calculate) should be \( n + n \approx n + \Lambda \). These hyperons should be present at densities \( \rho \approx (5-8) \times 10^{14} \text{ g cm}^{-3} \), which are achieved over a large fraction of the core of relevant mass (\( M \approx 1.3-1.6 \ M_\odot \)) neutron star models for many nuclear equations of state (which are softened by their presence: Balberg, Lichtenstadt, & Cook 1999). Employing the superfluid reduction factor \( R_{\text{sh}}(T/T_s) \) of Haensel et al. (2002) for normal nucleons and a hyperon with singlet superfluid transition temperature \( T_s \), we obtain

\[
F_{\text{sh}} = f_{\text{sh}} \frac{t_s^2 \tau(T)}{1 + [2\Omega T(T)/3]^{2}}, \quad \tau(T) = \frac{t_s T_s^{-2}}{R_{\text{sh}}(T/T_s)},
\]

where \( 1/\tau \) is the reaction relaxation rate. The constants \( t_0 \approx \)
and hyperon bulk viscosity (solid curve) are balanced by the neutrino emission (Brown 2000). The vanishing of equation (5), with the nuclear heating in the inner crust balanced by the neutrino bremsstrahlung and Cooper pairing of neutron-neutron neutrinos and (4) of evolution of $\alpha(t)$ and $\tilde{\Omega}(t)$, we can consider the thermal evolution (eq. [5]). With the neutrinos normal, they dominate the specific heat, giving $C(T) \approx 1.5 \times 10^{38} T_s \text{ ergs K}^{-1}$. The electron contribution is about 15 times less. We also take the nuclear heating constant $K_n = 1 \times 10^{-3}$ (Brown 2000).

The neutron luminosity is taken to be $L_s = L_{\text{int}} T_s^6 \times R_{\text{int}}(T), L_{\text{int}} = L_{\text{int}}(T) + R_{\text{int}}(T) + L_{\text{cr, int}} T_s^6 + L_{\text{m, int}} T_s^8 + L_{\text{cp, int}} T_s^8$. The proton superfluid reduction factors for the direct Urca reactions ($R_{\text{int}}$) and the modified Urca reactions ($R_{\text{int}}$), as well as most of the following constants, are obtained from the review of Yakovlev, Levinsh, & Shibanov (1999), using $T_s = 5 \times 10^9$ K. The other terms represent (inner crust) electron-ion and (core) neutron-neutron neutrino bremsstrahlung and Cooper pairing of (inner crust) neutrons. We take $L_{\text{int}} = f_{\text{int}} T_s^6$, where $f_{\text{int}}$ is the fraction of the neutron star that is above the density threshold for the direct Urca reactions (comparable to the above threshold for the hyperon reactions). The constants $L_{\text{int}} \approx 1.0 \times 10^{32} \text{ ergs s}^{-1}$ and $L_{\text{cr, int}} \approx 9.1 \times 10^{38} \text{ ergs s}^{-1}$ are obtained by fitting the results of Brown (2000; who did not consider the other processes) for normal and superfluid neutron stars. Finally, $L_{\text{m, int}} \approx 0.01 T_s^6$, while $L_{\text{cp, int}} \approx 9 \times 10^{31} \text{ ergs s}^{-1}$ is proportional to the scale length of variation of the neutron superfluid transition temperature within the inner crust (taken to be 100 m; Yakovlev, Kaminker, & Gnedin 2001).

We are interested in the evolution of neutron stars after they have been spun up to the point where the gravitational radiation growth rate has become equal to the viscous damping rate: $F_c \Omega_s = F_\text{cr} = F_\text{cr}$. This equality defines our initial state. Before that time, we see from equation (1) that any intrinsic perturbation could not grow from its (infinitesimal) value $\alpha_{\text{init}}$. The initial temperature $T_0$ is then determined by the vanishing of equation (5), with the nuclear heating in the inner crust balanced by the neutrino emission (Brown 2000).

This temperature is most sensitive to the direct Urca factor $f_{\text{int}}$, but varies by only 14% over the range $0 \leq f_{\text{int}} \leq 1$. Consistent with our assumption that the hyperon reactions are operative, we adopt $f_{\text{int}} = 1$. For a typical value $\tilde{\Omega}_s = 0.3$, the timescale $\tau_0 \equiv 1/f_{\text{int}} \approx 10^4$ s. Two other key timescales are those due to cooling and accretion: $\tau_a \equiv (C(T_s) T_s)^{1/2} / (L_s(T_s) T_s) \approx 5 \times 10^9$ yr and $\tau_r \equiv j_s / (j_s(M) T_s) \approx 10^7$ yr. Since we are concerned with timescales $\Delta_i \ll M_j / (M_i) \approx 10^{15}$ yr, we can take $M(t) \equiv M_i$.

In contrast to the initial state, the equilibrium state $X_i$ of our dynamical variables $X(t) = \{\alpha, \tilde{\Omega}, T\}$ is defined by the vanishing of the evolution equations (2), in addition to equations (1) and (5). From equation (3) or (4), one can see that the equilibrium amplitude is given by $\alpha_0 = (2K \tau_c F_\text{cr})^{1/2} \approx 10^{-6}$ to $10^{-3}$. These values of $\alpha_0$ are much less than the saturation amplitude of the $r$-mode instability (Arras et al. 2002).

The linearized analysis of Wagoner et al. (2001) shows that stability of the equilibrium requires that

$$\left( \frac{1}{L_s} \frac{\partial L_s}{\partial \tilde{\Omega}_s} \right) > \left( \frac{1}{F_\text{cr}} \frac{\partial F_\text{cr}}{\partial \tilde{\Omega}_s} \right) > 0,$$

assuming that $|\partial / \partial T| \sim 1/T$. If so, oscillations of frequency $f_r \sim K_\text{cr} \alpha_0 F_\text{cr}$ are damped at a rate $f_a \sim K_\text{cr} \alpha_0 F_\text{cr}$, where $F_\text{cr}$ is the equilibrium value of $F_\text{cr} = F_\text{cr}$ and $K_\text{cr} = 2K \tau_c F_\text{cr} C(T) \sim 10^3$ is the ratio of rotational to thermal energy. We have also used the fact that the viscous heating term in equation (5) is typically at least 10 times greater than the nuclear heating term. In practice, the first inequality in equation (8) is always satisfied, if the key requirement $\partial F_\text{cr} / \partial \tilde{\Omega}_s > 0$ holds (which is the case if $F_\text{cr}$ is dominated by the hyperon bulk viscosity, as shown in Fig. 1). This requirement applied to the initial state also guarantees slow evolution, as we see below.

In Figure 2 we show the critical curve $\tilde{\Omega}(T_s)$ given by $T_s = F_\text{cr}$ for three choices of the key parameters $T_i$ and $S_{\text{cr}}$. Also shown are the initial state $\tilde{\Omega}_i, T_i/10^9$ K and the equilibrium state $\tilde{\Omega}_s, T_i/10^9$ K, with $T_i > T_s \approx 3.32 \times 10^9$ K. For $T_i \approx 3 \times 10^9$ K, these results are independent of $T_s$. For $T_i \ll 1 \times 10^9$ K, these results are independent of the slippage factors $S_i$ and $S_s$ (and $\tilde{\Omega}_s$ does not increase greatly as $T_s$ decreases in this range).

It can be shown that the evolution from the initial state is also controlled by the sign of $(\partial F_\text{cr} / \partial \tilde{\Omega}_s)$, which is equal to the sign of the slope of the critical curve. If the slope is negative (case
1), there will be a thermal runaway with a growth rate \( f_r \) (given above) that is of the same magnitude as found by Levin (1999). If the slope is close to zero (case 2), there will initially be overstable oscillations of the type found by Wagoner et al. (2001). If the slope is positive (case 3), oscillations of the growing amplitude are damped out on the timescale \( T \), after which it slowly increases to its equilibrium value \( f_e \) (for the parameters of case 3), along with \( \Omega \) and \( T \). The time required to reach equilibrium is \( \Delta t \approx (\Delta \Omega/\Omega)f_e \). Throughout, \( f_r \) remains very close to \( f_e \), so the evolution is along the critical curve.

The maximum rotation rate (due to shedding) of our chosen neutron star model is \( f_{\text{max}} \approx (2/3)\Omega_f/2\pi = 856 \text{ Hz} \). Coherent oscillations in type 1 X-ray bursts have been observed at frequencies \( f \lesssim 590 \text{ Hz} \) (van der Klis 2000), which would then correspond to \( \Omega \lesssim 0.46 \) if they represented the rotation rate. There is evidence that some of these are the first harmonic, in which case the highest spin frequency is 350 Hz (van der Klis 2000). (For isolated neutron stars, \( f \leq 642 \text{ Hz} \).

Observations of the luminosity of LMXBs in their quiescent (low accretion rate) phase constrain the \( r \)-mode amplitude if they have reached the equilibrium state we have considered (Brown & Ushomirsky 2000), where the \( r \)-mode viscous heating balances the neutrino energy loss. This will also be addressed in a subsequent investigation.

Our main conclusion is that evolution to a stable equilibrium state can occur if (1) a significant fraction of the core of the neutron star is above the threshold for hyperons, (2) their superfluid transition temperature \( T \lesssim 2 \times 10^9 \text{ K} \), (3) the core neutrons near the crust are not a superfluid whose vortices are strongly pinned to the crust, and (4) the magnetic field is not too strong in that core-crust boundary layer. If Sco X-1 has been spun up by accretion to such a stable equilibrium state, it should be detectable by the second generation LIGO detectors. (However, its spin period remains unknown.) When signal recycling ("narrow banding") is employed, a few additional LMXB's may also be detectable (Cutler & Thorne 2002).

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