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Structure, Deformations and Gravitational Wave Emission of Magnetars

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Abstract. Neutron stars can have, in some phases of their life, extremely strong magnetic fields, up to 10^{15-16} G. These objects, named magnetars, could be powerful sources of gravitational waves, since their magnetic field could determine large deformations. We discuss the structure of the magnetic field of magnetars, and the deformation induced by this field. Finally, we discuss the prospects of detection of the gravitational waves emitted by these stars.

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1. Introduction

Magnetars are neutron stars (NSs) whose spin-down and bright emission activity are powered by the stellar magnetic field. The interest of the scientific community in these objects has been growing since 1992, when Thompson and Duncan [1, 2] first proposed a model which explains the spin-down rate and the emission properties of two classes of astrophysical objects, the soft-gamma repeaters (SGRs) and the anomalous X-ray pulsars (AXPs) in terms of strong magnetic fields.

These objects have a very steep spin-down, and a very intense X-ray (and gamma-ray) activity, with periodic bursts of $\sim 10^{41}$ erg/s. Furthermore, in the last decades three giant flares from SGRs have been observed, with luminosities reaching $\sim 10^{47}$ erg/s. The observed spin-down of SGRs and AXPs corresponds (through the well-known dipole emission formula $P\dot{P} \propto B^2$) to surface magnetic fields of the order of $10^{14} - 10^{15}$ G. In the model of Thompson and Duncan, the gamma activity is understood in terms of the evolution of the interior magnetic field, which is as large as the surface field, or even larger; in their model, the field (or a significant fraction of it) is a toroidal field[‡]. This magnetic field has been produced in the early phases of the NS life, just after the supernova explosion, due to flux conservation in the core collapse and/or to dynamo processes, related to convective motion and differential rotation. It is worth noting that, although observed magnetars are slowly rotating, with periods P of the order of 10 s,

[‡] If we define a polar coordinate frame (r, θ, ϕ) about the magnetic axis, the r - and θ -components of the magnetic field are called *poloidal*, the ϕ -component is called *toroidal*.

newly born magnetars could have much higher rotation rates, with periods $P \sim 10^{-3}$ s. Today we know 18 magnetars§, but it is believed that a significant fraction ($\gtrsim 10\%$) of NSs would possibly become magnetars at some stage of their evolution [3].

Due to their extreme properties, magnetars are very interesting objects both for astrophysics and for gravitational wave physics. Quasi-periodic oscillations have been detected in the aftermath of the giant flares of SGRs; this is the first observational evidence of NS oscillations [4]. It has been suggested that magnetars may be the central engine for some gamma-ray bursts [1, 5, 6]. Last but not least, as we discuss below, the magnetic field could produce a deformation much larger than that due to other mechanisms, thus magnetars are also interesting sources of gravitational waves [7]-[12]. We also remark that the present and future observational properties of magnetars could shed light on the internal composition of NSs, and thus on the behaviour of matter at supranuclear densities.

For these reasons, in the last decades magnetars have been widely studied. However, their internal structure is still poorly understood. We do not know, for instance, how strong is the interior magnetic field, and whether the toroidal components prevail on the poloidal ones; we do not know whether the field is mainly dipolar or the higher order multipoles dominate. This information would be very important, to understand the astrophysical processes involving magnetars, and to assess the relevance of these stars as gravitational wave sources.

The magnetar model proposed in [1, 2] is dynamical, and the magnetic field evolves from its birth to its decay [13, 14], through different processes (ambipolar diffusion, Hall drift, Ohmic decay). However, in some phases of the early life of a neutron star it is legitimate to describe a magnetar as a stationary object, using the ideal magnetohydrodynamics (MHD) approximation, as we shall briefly explain.

Let us consider what happens when a strongly magnetized neutron star is born.

- In the first seconds after the supernova explosion, the proto-neutron star is a very complicate and dynamical object, with turbulent and convective motion, differential rotation, and (eventually unstable) oscillations. In this period dynamo processes amplify the stellar magnetic field.
- After few (or few tens of) seconds, convective instability is suppressed, and the matter composing the star can be described by a single, non-superfluid perfect fluid with infinite conductivity (ideal MHD approximation). As shown by numerical simulations in the ideal MHD approximation [15, 16], the fluid is likely to settle down to a stationary configuration on a dynamical timescale of the order of Alfvén's time ($t_A \sim 0.01 - 10$ s).
- After few minutes, matter becomes superfluid and the crust forms; thus the ideal MHD approximation no longer applies. The magnetic field evolves on timescales $\sim t_{decay}$ of the order of thousands of years or more.

§ For an up-to-date catalog, see <http://www.physics.mcgill.ca/~pulsar/magnetar/main.html>

We remark that as the crust forms, the magnetic field is likely to freeze in the stationary configuration reached in the previous stage. Therefore, this configuration could be an appropriate description of the stellar magnetic field for timescales $t_A \lesssim t \lesssim t_{decay}$.

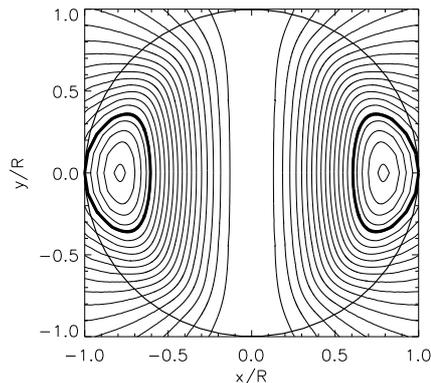


Figure 1. The field lines of a twisted-torus magnetic field configuration, projected in the meridional plane. The toroidal field is confined in the region inside the thick curves.

In the last decade many authors have been developing models of stationary magnetized neutron stars in ideal MHD [17]-[23],[10],[11], including more and more ingredients in order to capture the essential features of the system: poloidal and toroidal fields, general relativity, “realistic” equation of state (EOS). In recent papers on the subject [21]-[23], [11] a *twisted-torus configuration* has been considered, in which the poloidal magnetic field extends throughout the star and in the exterior, whereas the toroidal field is confined into a torus-shaped region inside the star, where the field lines are closed (see Fig. 1). There are different reasons for this choice:

- It has long been known that purely toroidal and purely poloidal magnetic field configurations are unstable [24, 25, 26]; it is expected that a stable configuration should have both components [27], with comparable amplitudes.
- Numerical simulations [15, 16] have shown that the magnetic field tends to a twisted-torus configuration in which the toroidal and poloidal components have comparable amplitudes, for quite generic initial conditions (see also the analysis of [28]). This configuration appears to be stable, at least on a timescale $t_A \lesssim t \ll t_{decay}$. We remark that these simulations have been performed in a Newtonian framework, assuming a polytropic EOS for the stellar fluid.
- The results of [15, 16] can be understood, at least qualitatively, as follows. Let us consider the *magnetic helicity*

$$H_m = \int \mathbf{A} \cdot \mathbf{B} dV \quad (1)$$

where \mathbf{A} , \mathbf{B} are the vector potential and the magnetic field, respectively (note that magnetic helicity can also be defined in a relativistic framework). The following properties hold.

- The magnetic helicity is conserved on a timescale $\ll t_{decay}$.
- It vanishes if the field is either purely poloidal or purely toroidal. Thus, if the field is mixed (poloidal and toroidal) at the beginning, it must remain mixed for a long time.
- The toroidal field is proportional to the electric current, thus, neglecting the stellar magnetosphere, it must vanish outside the star.
- The ratio between the toroidal and poloidal amplitudes can be described by a function ζ , which is constant along each field line [19]. Therefore, a field line which extends outside the star must have $\zeta = 0$, i.e. it must be purely poloidal.

It follows that, as the magnetic field reaches a stationary configuration, it must retain a mixed character, and the toroidal field must be confined inside the star, since the field lines with a non-vanishing toroidal component cannot cross the stellar surface. Such lines cover a torus-shaped region, tangent to the stellar surface at the equator. This is the twisted-torus configuration (see Fig. 1).

In the next Sections we discuss the features of magnetars with twisted-torus magnetic fields; our study is based on a model we have recently developed [22, 11] (see also [10]). In Section 2 we briefly describe our model, and determine the magnetic field structure, discussing the relative amplitude of toroidal and poloidal fields we expect. In Section 3 we determine the stellar deformation induced by the magnetic field, discussing how it depends on the EOS of the matter composing the star. In Section 4 we discuss the possible gravitational emission of magnetars. We remark that, while Sections 2, 3 are based on the results of [22, 11], Section 4 is a new contribution to the debate on the relevance of magnetars as sources for present and future gravitational wave detectors.

2. Structure

We consider (see [22, 11] for more details) a stationary, axisymmetric magnetized NS in the framework of general relativity. We neglect stellar rotation (note that, as shown in [21], twisted-torus configurations are not significantly affected by stellar rotation) and the effect of the magnetosphere. Furthermore, we assume that the stellar matter is described by a single perfect fluid with infinite conductivity (ideal MHD approximation). The magnetic field is treated as a perturbation of a spherically symmetric background with metric

$$ds^2 = -e^{\nu(r)} dt^2 + e^{\lambda(r)} dr^2 + r^2(d\theta^2 + \sin^2 \theta d\phi^2) \quad (2)$$

(ν, λ solutions of the unperturbed Einstein's equations describing the stellar structure) and four-velocity $u^\mu = (e^{-\nu/2}, 0, 0, 0)$. We choose two EOSs, named APR2 [30] and GNH3 [31], to model stars with large and small compactnesses, respectively; indeed, a NS with mass $M = 1.4 M_\odot$ has radius $R = 11.58$ km (APR2 EOS) or $R = 14.19$ km (GNH3 EOS).

The background is perturbed by a stationary, axisymmetric electromagnetic tensor $F_{\mu\nu} = A_{\mu,\nu} - A_{\nu,\mu}$, associated to a current j^μ , an electric field $E_\mu = F_{\mu\nu} u^\nu$ and

a magnetic field $B_\alpha = \frac{1}{2}\epsilon_{\alpha\beta\mu\nu}u^\beta F^{\mu\nu}$. The equations of ideal MHD are the baryon number conservation $(nu^\mu)_{;\mu} = 0$ (n baryon density), the relativistic Euler equation $(\rho + p)a_\mu + p_{,\mu} + u_\mu u^\nu p_{,\nu} - f_\mu = 0$ (ρ mass-energy density, p pressure, $f_\mu = F_{\mu\nu}j^\nu$ Lorentz force), and the vanishing of the electric field $E_\mu = 0$.

With an appropriate gauge choice, the vector potential can be written as

$$A_\mu = (0, e^{(\lambda-\nu)/2}\Sigma, 0, \psi) \quad (3)$$

where the ‘‘flux function’’ $\psi(r, \theta)$ describes the poloidal field, and the function $\Sigma(r, \theta)$ describes the toroidal field. Neglecting higher order terms in the perturbation ($O(B^4)$), a remarkable property holds: the quantity $\sin\theta\Sigma_{,\theta}$ only depends on the flux function ψ (which is constant along each field line). We can then define a function β as

$$\sin\theta\Sigma_{,\theta} \equiv \beta = \beta(\psi). \quad (4)$$

Then, once we impose a form for $\beta(\psi)$, the magnetic field configuration is entirely determined by the flux function $\psi(r, \theta)$, which can be found by solving the *relativistic Grad-Shafranov equation*:

$$\begin{aligned} & -\frac{e^{-\lambda}}{4\pi} \left[\psi'' + \frac{\nu' - \lambda'}{2} \psi' \right] - \frac{1}{4\pi r^2} [\psi_{,\theta\theta} - \cot\theta\psi_{,\theta}] \\ & - \frac{e^{-\nu}}{4\pi} \beta \frac{d\beta}{d\psi} = (\rho + P)r^2 \sin^2\theta [c_0 + c_1\psi] \end{aligned} \quad (5)$$

with c_0, c_1 arbitrary constants. This equation follows from the ideal MHD equations. By expanding the flux function $\psi(r, \theta)$ in Legendre polynomials as

$$\psi(r, \theta) = \sum_{l=1}^{\infty} a_l(r) \sin\theta P_{l,\theta}(\cos\theta), \quad (6)$$

Eq. (5) gives a coupled system of ordinary differential equations for the functions $\{a_l(r)\}_{l=1,2,\dots}$. These equations admit two particular sets of solutions: the symmetric (with respect to the equatorial plane) solutions, with vanishing even-order components ($a_{2l} \equiv 0$) and the antisymmetric solutions, with vanishing odd-order components ($a_{2l+1} \equiv 0$). It is reasonable to expect that the actual field configuration of these stars is, with a good approximation, symmetric with respect to the equatorial plane. Indeed, the magnetic field has a non-vanishing dipole ($l = 1$) component outside the star, and the antisymmetric solutions have vanishing magnetic helicity, therefore symmetric solutions are energetically favoured with respect to the others. Furthermore, an antisymmetric solution would likely be unstable on a dynamical timescale, since two opposite magnetic field loops could annihilate each other. This would be in some sense similar to the Flowers-Ruderman instability of purely poloidal fields [25] (see also [6, 26]).

The twisted-torus configurations are those for which $\beta(\psi)$ is continuous and has the form

$$\beta(\psi) \sim \Theta(|\psi/\bar{\psi}| - 1), \quad (7)$$

where $\bar{\psi} \equiv \psi(R, \pi/2)$ is the value of the function ψ on the stellar surface at the equator, and Θ is the Heaviside step function. This can be understood by looking at Fig. 1.

The magnetic field lines are also lines of constant ψ , and the thick line corresponds to $\psi = \bar{\psi}$. The toroidal region inside the thick line has $\psi > \bar{\psi}$, and Eq. (7) implies that $\beta \neq 0$, i.e. the toroidal field is non-vanishing, only in this region.

In [11] we solved the relativistic Grad-Shafranov equation, expanded in Legendre polynomials (with l odd), assuming a quite general parametrization for $\beta(\psi)$ (compatible with the twisted-torus condition (7)) and employing the two EOSs APR2 and GNH3, which span a wide range of stellar compactness. A remarkable result we have found is that the toroidal field never contributes to more than 13% of the total magnetic energy of the star. This is due to the fact that, if we enhance the amplitude of the toroidal field (roughly speaking, by making β larger), the region where the toroidal field is non-vanishing shrinks. Note however that, in this region, the toroidal field can be larger than the poloidal field. Similar results have been obtained in [23], using a polytropic EOS in a Newtonian framework. We remark that this result, if confirmed, would challenge an assumption often used in magnetar models [32, 9], i.e. that the toroidal field prevails onto the poloidal inside the star.

The main open issue regarding these configurations is their stability. Indeed, they are stationary by construction, but may be unstable. Actually, in [28] it has been found that magnetic field configurations in which the toroidal field accounts for less than 20% of the total magnetic energy appear to be unstable (in the framework of Newtonian gravity and assuming a polytropic EOS). However, recent stability analyses of purely poloidal magnetic field configurations (see [29] and references therein) show that the onset of the instability is localized along the “neutral line”, which is the circle in the equatorial plane threading the closed field lines inside the star (see Fig. 1); as argued in [29], a strong toroidal component along this line, like in the twisted-torus configurations, could suppress the instability even when the overall energy of the toroidal field is small.

3. Deformations

Once the magnetic field configuration has been determined with the perturbative approach outlined above, it is possible to compute (perturbatively) the corresponding stellar deformation by solving Einstein’s equations $\delta G_{\mu\nu} = \frac{8\pi G}{c^4} \delta T_{\mu\nu}$ [10, 11]. The quadrupole ellipticity

$$\epsilon_Q = \frac{Q}{I} \quad (8)$$

(Q mass-energy quadrupole moment, I mean momentum of inertia) is the most relevant quantity encoding the stellar deformation: it depends on the distribution of matter throughout the entire star (note that the gravitational wave emission depends on ϵ_Q). The mass-energy quadrupole moment Q can be extracted by the far field limit of the metric

$$g_{00} \rightarrow \dots - 2Q \frac{e^\nu}{r^3} P_2(\cos \theta) \quad (9)$$

and in the weak field limit it reduces to $Q \simeq \int_V \rho(r, \theta) r^2 P_2(\cos \theta) dV$.

The poloidal field tends to make the star oblate, which corresponds to $\epsilon_Q > 0$. The toroidal field, instead, tends to make it prolate, i.e. with $\epsilon_Q < 0$. The determination of the sign of ϵ_Q is important, because if $\epsilon_Q < 0$ a “spin flip” mechanism, suggested by Jones and Cutler [7, 8], could take place: the angle between the rotation axis and the magnetic axis would grow until they become orthogonal. This process would be associated to a large gravitational emission. However, as discussed in Section 2, the stationary twisted-torus configurations seem to be mainly poloidal, and indeed the corresponding deformations always have $\epsilon_Q > 0$. Therefore, the twisted-torus configurations seem not to be compatible with the Jones-Cutler mechanism.

The stellar deformation induced by twisted-torus magnetic field configurations depends on the EOS: less compact stars have larger deformations. Furthermore, if one changes the magnetic field configuration (i.e. changes the choice of $\beta(\psi)$ satisfying (7), see Section 2), the stellar deformation changes less than 10%. Note that, since the poloidal and toroidal fields have competing effects and the poloidal field prevails, it follows that larger toroidal fields correspond to smaller deformations.

It is possible to summarize the deformations of these magnetized NSs as follows:

$$\epsilon_Q \simeq k \left[\frac{B_{pole}(G)}{10^{16}} \right]^2 \times 10^{-4}. \quad (10)$$

(B_{pole} is the amplitude of the dipolar surface magnetic field at the stellar pole). Here k is a coefficient which depends on the EOS: $k \sim 4$ for the APR2 EOS, $k \sim 9$ for the GNH3 EOS. We do not expect other EOSs to give results very different from these (the results of [23] for a polytropic EOS are also similar).

The ellipticities (10) are larger than the bounds derived in [33, 34] (see also [35]), $|\epsilon| \lesssim 10^{-5} - 10^{-6}$, by evaluating the maximal strain that the crust can sustain. However, these bounds do not apply necessarily to our case. Indeed, here we consider a fluid star which is deformed by the magnetic field before the crust forms. In this scenario, the equilibrium configuration of the crust would be its initial, non-spherical shape, and the limits derived in [33, 34] may be violated. However, we do not know how long the crust would remain in a non-spherical shape: in order to understand the evolution and the persistence of the stellar deformation, a dynamical study of the magnetic field evolution on longer time-scales would be needed.

4. Gravitational wave emission

If an axisymmetric NS with quadrupole ellipticity ϵ_Q induced by a magnetic field, rotates about an axis forming an angle α with the magnetic axis, it emits gravitational waves. If α is small, gravitational radiation is mainly emitted at the same frequency ν as the rotation rate, with amplitude

$$h_0 \simeq \frac{4G}{rc^4} (2\pi\nu)^2 I |\epsilon_Q| \sin \alpha. \quad (11)$$

We remark that the best available estimate of the “wobble angle” α of a neutron star is $\alpha = 3^\circ$ for PSR B1828-11 [36]. In the Jones-Cutler process, which takes place as

$\epsilon_Q < 0$, the wobble angle increases towards $\alpha = 90^\circ$, with a great enhancement of the gravitational radiation. However, as discussed in Section 3, this is not the case for the twisted-torus configurations.

The detectability of gravitational emission from magnetically deformed NS, described by Eq. (11), depends both on the overall magnetic field amplitude, which determines ϵ_Q , and on the duration of the emission process. Indeed, different dissipative processes tend to reduce both the wobble angle and the rotation frequency, then reducing the time the emission frequency spends in the bandwidth of ground-based interferometers (from few tens to few hundreds of Hertz).

- As discussed in [37], the wobble angle of an oblate ($\epsilon_Q > 0$) star with rotation period P would decay, due to internal dissipation, in a timescale

$$\tau_d \sim \frac{nP}{\epsilon_Q} \quad (12)$$

where the parameter n is unknown, since we do not have a clear understanding of the damping mechanism. However, all the possible damping processes which have been considered so far (see [8, 37] and references therein) lead to typical values for n in the range $10^2 - 10^5$. This would correspond to a damping timescale ranging from few months to few tens of years if $B_{pole} \sim 10^{15}$ G. Therefore, after at most few tens of years the rotation and symmetry axis would become nearly parallel, and the gravitational emission would become negligible, unless some pumping mechanism [37] takes place which increases the wobble angle.

- A NS with dipolar field at the pole B_{pole} and wobble angle α spins down with a period derivative \dot{P} given by [3, 38]

$$B_{pole} \simeq 6.4 \times 10^{19} \frac{\sqrt{P\dot{P}}}{\sin \alpha} \text{ G} \quad (13)$$

(note that many authors consider the average surface magnetic field, which is $\frac{1}{2}B_{pole}$ [39, 40]). Since Eq. (13) implies that $\dot{P} = K/P$ (where K is a constant), the star slows down from an initial period P_{in} to a period P_{fin} in the characteristic time

$$\tau_c = \frac{1}{2K}(P_{fin}^2 - P_{in}^2) \simeq \frac{P_{fin}^2}{2K}. \quad (14)$$

Therefore, a NS with B_{pole} of the order of 10^{15} G and a small wobble angle could lie in the bandwidth of ground based interferometers for a time ranging from few months to few years. If the magnetic field is larger the star spins down more rapidly, making detection more difficult. The detection is also unlikely if the parameter n in (12) is much smaller than the upper limit $\sim 10^4$, since in this case the wobble angle would rapidly decay.

In Fig. 2 we show the signal emitted from a NS with $B_{pole} = 10^{15}$ G and wobble angle $\alpha = 3^\circ$, at a distance of 10 kpc (i.e. in our galaxy), computed by Eqns. (10), (11). The star initially rotates with $\nu = 700$ Hz, which is close to the largest rotation frequency of known pulsars; note that magnetars are believed to rotate rapidly at birth [1]. This

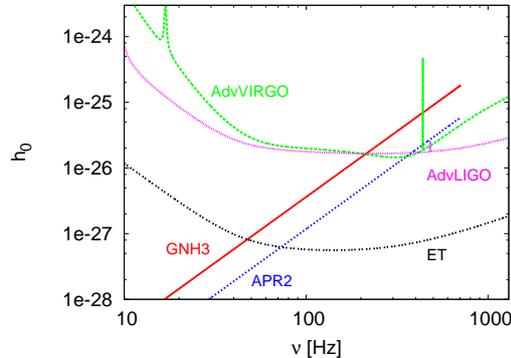


Figure 2. Gravitational signal emitted by a rotating NS at a distance of 10 kpc, deformed by a twisted-torus magnetic field, with $B_{pole} = 10^{15}$ G, $\alpha = 3^\circ$, $M = 1.4 M_\odot$ and EOS GNH3 (solid line) and APR2 (dotted line). The sensitivity curves of advanced LIGO, advanced VIRGO (with integration time of three months) and ET (with integration time of one year) are shown for comparison.

signal is compared with the sensitivity curves of the advanced detectors LIGO, VIRGO (assuming an integration time of three months) and of the third generation detector ET (assuming one year integration time) \parallel . An estimate of the spin-down time τ_c by Eq. (14) shows that, if the wobble angle decay is not too fast, the signal lies in the bandwidth of advanced LIGO/VIRGO for a few months, and it lies in the bandwidth of ET for a few years, consistently with the integration times we have employed (see also [41]).

Figure 2 shows that the signal could be well detected by ET, and marginally detected by advanced LIGO/VIRGO. However, one should also take into account the event rate of the process generating the gravitational wave signal. In our scenario, a NS could maintain a strong, twisted-torus magnetic field and the corresponding deformation for several years (in the most optimistic case, up to thousands of years), thus there may be several NS in our galaxy with large deformation. However, only few of them would rotate rapidly enough to be detected by ground based interferometers. If there is no spin-up process, the rate of the events described in Fig. 2 would be at most the same as NS birth rate, i.e. few per century in our galaxy. On the other hand, accretion from a companion star could spin-up the star and increase the wobble angle; in this case the event rate may be significantly larger. However, it should be remarked that no evidence of a companion has been found for any of the observed magnetars so far. We also mention that, as discussed in [12], the stochastic background of gravitational waves from magnetars could be detectable by the third generation detector ET.

Finally, we mention that LIGO and Virgo set an upper limit of the order of $\sim 10^{-4}$ on the deformations of known pulsars (Crab, J0537-6910 and J1952+3252) [42]. Indeed, larger deformations would have produced signals strong enough to be detected. These limits are stronger than current limits arising from spin-down [38]. We remark that the deformations considered in the analysis of [42] are different from those considered here

\parallel <http://www.ligo.caltech.edu>; <http://www.ego-gw.it>; <http://www.et-gw.eu>

and in current literature on magnetars. Indeed, we consider axially symmetric stars inclined by an angle α with respect to the rotation axis (which yield gravitational waves at frequency ν). In the data analysis carried on in [42], instead, tri-axial deformations (without inclination) have been considered, which yield gravitational waves at frequency 2ν .

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