Implications of magnetar non-precession

K. Glampedakis\(^1\) & D.I. Jones\(^2\)

\(^1\) Theoretical Astrophysics, University of Tuebingen, Auf der Morgenstelle 10, Tuebingen, D-72076, Germany
\(^2\) School of Mathematics, University of Southampton, Southampton SO17 1BJ, UK

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ABSTRACT

The objects known as anomalous X-ray pulsars and soft gamma repeaters are commonly identified with magnetars, neutron stars with ultrastrong magnetic fields. The rotational history of these objects has, so far, revealed no evidence of free precession. At the same time these objects do not generally appear to have magnetic axes nearly parallel or orthogonal to their spin axes. In this paper we show that the combination of these two observations, together with simple rigid-body dynamics, leads to non-trivial predictions about the interior properties of magnetars: either (i) elastic stresses in magnetar crusts are close to the theoretical upper limit above which the crustal matter yields or (ii) there is a “pinned” superfluid component in the magnetar interior. As a potentially observable consequence of these ideas we point out that, in the case of no pinned superfluidity, magnetars of stronger magnetic field strength than those currently observed would have to be nearly aligned/orthogonal rotators.

Key words: stars: magnetars – stars: neutron – stars: rotation

1 INTRODUCTION

The magnetars are a subset of the neutron star population with ultrastrong magnetic fields, typically greater than $10^{14}$ G, and long spin periods, typically $1 - 10$ seconds. They manifest themselves observationally as anomalous X-ray pulsars (AXPs) and soft gamma repeaters (SGRs); see Woods & Thompson (2004) for a review. The latter class of objects are of particular interest, as periodicities have been observed in the tails of bursts from some objects, possibly providing the first ever evidence for excitation of neutron star normal modes. The theory of magnetars as very strongly magnetised neutron stars was developed by Thompson & Duncan (see e.g. Thompson & Duncan (1995, 1996, 2001)), in which the decaying magnetic field powered the outbursts.

In this paper we will concern ourselves with something that the magnetars don’t do—they do not seem to undergo free precession. More precisely, we will base our analysis on the following two working assumptions, both motivated directly by observations: (i) magnetars do not precess, and (ii) their magnetic pulsation axes are neither aligned nor orthogonal to the spin axes. On the basis of these two observational results and simple rigid body dynamics we can then deduce that either (a) the stellar crust is very highly strained, right at or close to the limit predicted theoretically by detailed modelling, or (b) the magnetars contain a pinned superfluid component. We believe that either conclusion is interesting.

The first assumption, the lack of precession, follows from timing studies. Precession would result in a periodic modulation in the time of arrival (and possibly also shape) of the magnetar pulsations (Jones & Andersson 2001). Indeed, early observations of timing irregularities prompted Melatos to suggest precession was a generic feature of the AXP population (Melatos 1999). However, further observation of the AXPs (Kaspi, Chakrabarty & Steinberger 1999; Kaspi et al. 2001; Gavriil & Kaspi 2002) and SGRs (Woods et al. 2000) has ruled out precession at an amplitude level above the rms amplitude of magnetar timing noise. Small-amplitude precession, ‘buried’ inside the timing noise, is still a possibility and future timing data analysis should attempt to address this issue. (A strict lack of precession is not essential for our arguments: as will become clear, a small precession amplitude would imply that the magnetar spins about an axis close to a principal axis or the superfluid pinning axis; this is the essential point for our modelling).

Our second assumption concerns the magnetar inclination angles, i.e. the angle $\theta_B$ between their spin axes and magnetic axes; we assume that magnetars are neither aligned ($\theta_B \approx 0$) nor orthogonal ($\theta_B \approx \pi/2$). By talking of a ‘magnetic axis’, we are implicitly assuming there exists a well defined axis about which the internal magnetic field is approximately symmetric; numerical simulations indicate this is likely to be the case (Braithwaite 2009). Then, to make any observationally-motivated statement about the location of this axis, we also assume that this axis coincides with the axis defined by the magnetar pulsations, which is
It is a reasonable assumption given the known importance of the magnetic field in determining the surface temperature profile (Oz{\text{"}el}, Psaltis & Kaspi 2001).

Given these assumptions, the result that the magnetars are not aligned rotators follows trivially from the fact that we see magnetar pulsations at all. That the two axes are not orthogonal can be argued on the basis of the harmonic content of the pulse profiles. An orthogonal rotator with two nearly identical anti-podal hot spots would, regardless of the observer’s location, display a strong second harmonic in its pulse power spectrum, something which has not been reported to be generic (but see the data for 1E 2259.1+586 in Fig. 2 of Gavriil & Kaspi 2002, which does display strong harmonic structure). Therefore it seems that the majority of the magnetars are neither nearly aligned nor nearly orthogonal. However, one caveat needs to be added: the argument above, while safe for, say, a radio pulsar, where the emission comes from high magnetospheric altitudes, is less secure for a magnetar. As described in Oz{\text{"}el}, Psaltis & Kaspi (2001), if the magnetar radiation is produced at the stellar surface itself, the combination of beaming and relativistic light bending can result in even a single hot spot profile giving rise to a pulse with multiple peaks, making statements concerning the inclination angle very difficult to make. We regard the assumption of non-orthogonality between the spin and magnetic axes as the weakest link in our argument.

Ours is not the first analysis of the precessional dynamics of magnetars. As noted above, Melatos (1999) attempted to interpret the irregular spin-down of AXPs as free precession, but sustained observations did not confirm his model (Kaspi, Chakrabarty & Steinberger 1999; Kaspi et al. 2001; Gavriil & Kaspi 2002). More closely related to our work is that of Wasserman (2003), who applied the same techniques as us (albeit without accounting for the presence of a superfluid component) to work out conditions under which radio pulsars would not precess, and used these to made statements about the conditions under which precession is likely, and how the star’s shape may change plastically over long timescales. As outlined above, the problem we address in this paper is different and perhaps simpler. The observations indicate that the magnetars do not precess, which we will make the perfect pinning assumption, i.e. the pinned neutron vortices are physically immobilised with respect to the crust which in turn implies a superfluid angular velocity \( \Omega^c = \Omega^s \hat{n}^c \) fixed in the crust frame.

We will use to parameterise the sizes of these contributions to the moment of inertia tensor. If \( I_0 = I_c + I_s \) denotes the spherical moment of inertia of the whole star, we can define

\[
\epsilon_c = \frac{\Delta I_c}{I_0}, \quad \epsilon_B = \frac{\Delta I_B}{I_0}, \quad \epsilon_s = \frac{I_s}{I_0} \tag{3}
\]

The deformation \( \epsilon_c \) is sourced by strains in the solid crust. The recent calculation by Haskell, Jones & Andersson (2006) place upper limits on this of

\[
(\epsilon_c)_{\text{max}} \approx 10^{-5} \left( \frac{\nu_{br}}{10^{-3}} \right), \tag{4}
\]

where we have parameterised the crustal breaking strain \( \nu_{br} \) in terms of the value found in the state-of-the-art calculations of Horowitz & Kadau (2009).

For the magnetic deformation a somewhat crude estimate (but consistent with more rigorous calculations, e.g. Akg{"u}n & Wasserman 2008; Colaiuda et al. 2008; Lander & Jones (2009)) is

\[
\epsilon_B \approx 10^{-6} \left( \frac{H}{10^{15} \text{ G}} \right) \left( \frac{\tilde{B}}{10^{15} \text{ G}} \right), \tag{5}
\]

where \( \tilde{B} \) is a volume average of the internal magnetic field and \( H \approx 10^{15} \text{ G} \) if the core sustains type II proton superconductivity, otherwise \( H = \tilde{B} \). The deformation \( \epsilon_B \) can be either positive or negative depending on the relative strength between the poloidal and toroidal magnetic field components (Mestel & Takhar 1972; Cutler 2002).

Clearly, given the previous estimates, the magnetic and elastic deformations are small numbers; the same is not necessarily true for \( \epsilon_s \). If pinning occurs only in the inner crust, only a few percent of the stellar superfluid contributes to \( I_s \), giving \( \epsilon_s \approx 10^{-2} \). However, if vortex pinning is efficient in the neutron star core due to the interaction with the magnetic fluxtubes then \( \epsilon_s \) could even be of order unity. Either way, \( \epsilon_s \) is likely to be much greater than either of \( \epsilon_c \) and \( \epsilon_B \).

We will now easily find the non-precessional solutions. The total angular momentum of our star is given by the sum of the crustal and superfluid contributions

\[
J^i = I_c^i \hat{\Omega}_c + I_s^i \hat{\Omega}_s. \tag{6}
\]

Let \( \theta_c \) and \( \theta_B \) denote the angles that the crustal and magnetic deformations make with the angular velocity vector, so that \( \hat{n}_c^i \hat{\Omega}_c = \cos \theta_c \) and \( \hat{n}_B^i \hat{\Omega}_s = \cos \theta_B \). The angular momentum can then be written

\[
J^i = \Omega \left( I_c \hat{n}_c^i \Omega_c + \Delta I_c \cos \theta_c \hat{n}_c^i + \Delta I_B \cos \theta_B \hat{n}_B^i \right) + I_s \Omega_s \hat{n}_s^i. \tag{7}
\]

The Euler equations of rigid body dynamics then give

\[
\dot{J}^i + \epsilon^{ijk} \Omega_j J_k = 0, \tag{8}
\]
where the time derivative is evaluated in the rotating crust frame. For a non-precessional and therefore stationary solution we must have \( \dot{J} = 0 \) which implies that \( J' = 0 \). We therefore obtain the relationship
\[
(I_c - \lambda)\hat{n}_i^c + \Delta I_c \cos \theta_c \hat{n}_i^c + \Delta I_B \cos \theta_B \hat{n}_i^B + I_\Omega \frac{\Omega}{\Omega} \hat{n}_i^s = 0
\] (9)
between the four unit vectors \( \hat{n}_c, \hat{n}_B, \hat{n}_s, \hat{n}_s \).

Equation (9) is crucial—we can easily derive our two key results from it. Firstly, consider the case where there is no pinned superfluid. Setting \( I_s = 0 \) we obtain
\[
\frac{I_c - \lambda}{I_0} \hat{n}_i^c + \epsilon_c \cos \theta_c \hat{n}_i^c + \epsilon_B \cos \theta_B \hat{n}_i^B = 0.
\] (10)
This shows that the three unit vectors \( \hat{n}_c, \hat{n}_B, \hat{n}_s \), are coplanar. We can therefore make the decomposition into a pair of basis vectors \( \hat{n}_1, \hat{n}_2 \) such that
\[
\hat{n}_1 = \sin \theta_B \hat{n}_1^c + \cos \theta_B \hat{n}_1^s, \quad \hat{n}_2 = \sin \theta_c \hat{n}_1^c + \cos \theta_c \hat{n}_1^s.
\] (11, 12)
Inserting into equation (10) and projecting along \( \hat{n}_1 \) leads to the non-precession condition
\[
\epsilon_c \sin 2\theta_c = -\epsilon_B \sin 2\theta_B. \quad (14)
\]
Thus, in the case of zero pinning, non-precession implies a simple relation between the crustal and magnetic deformations. The same result was obtained by Wasserman (2003).

Now consider the case where there is pinning. We define the angle \( \alpha = \hat{n}_s^c \hat{n}_s^s = \cos \theta_s \). Taking the vector product of equation (9) with \( \hat{n}_s^c \) and squaring the result leads to
\[
\left(2\epsilon_c \frac{\Omega}{\Omega} \sin \theta_s \right)^2 = (\epsilon_c \sin 2\theta_c)^2 + (\epsilon_B \sin 2\theta_B)^2 + 8\epsilon_c \epsilon_B \cos \theta_c \cos \theta_B \cos \beta - \cos \epsilon_c \cos \theta_B). \quad (15)
\]
This is a rather cumbersome relation between the angles that the superfluid, the crustal deformation, and the magnetic axis make with the rotation axis. However, from this we see one important thing: the star rotates about an axis very close to the superfluid pinning axis, with the angle of misalignment being of order of the larger of \( \Delta I_c/I_s \) and \( \Delta I_B/I_s \). In the notation of (3), we therefore have
\[
\theta_s \sim \max \left( \frac{\epsilon_c}{\epsilon_s}, \frac{\epsilon_B}{\epsilon_s} \right) \ll 1. \quad (16)
\]
Note that in the pinning scenario, the magnetar non-precession does not place any constraint on the crustal deformation. Also, it is possible for the star to spin about the crustal principal axis (as defined by eqn. (14)) and have a pinned superfluid, provided the pinning axis lies exactly along the spin axis (\( \theta_s = 0 \)). This would correspond to both our scenarios being realised. However, there is no apriori reason to expect the symmetry axis of the pinning sites to coincide with the spin axis, so this is a clearly special case.

It is also interesting to note the well known result that the non-precessing state is the one that minimises the kinetic energy at fixed \( J \) (Landau & Lifshitz 1977; Shahani 1972). Therefore, an isolated star, dissipating kinetic energy, would evolve to such a state on the dissipation timescale.

Finally, note that this last point has repercussions for the so-called Mestel-Jones spin-flip mechanism (Mestel & Takhar 1972; Jones 1975; Cutler 2002), in which a star with a dominantly prolate magnetic deformation of its inertia tensor (i.e. \( \epsilon_B < 0 \)) is driven to a configuration with the magnetic axis orthogonal to the spin axis. Clearly, the presence of a pinned superfluid completely suppresses this, assuming that the magnetic and pinning axes remain fixed with respect to each other (as in our model). According to the above discussion, the minimum energy state for a star containing a pinned superfluid has \( \theta_s \approx 0 \), i.e. the pinning axis must remain close to the spin axis, preventing the migration of the magnetic axis to an orthogonal location.

### 3 Extensions to the Basic Model

The previous model can be extended in terms of realism by incorporating additional physical effects. These include (i) rotational “bulges” in the moments of inertia, and (ii) the action of a non-dissipative torque due to the external dipole field (the so-called “anomalous” torque (Goldreich 1970)).

With the rotational deformations included the previous moment of inertia tensors become
\[
\begin{align*}
I_{ij}^c &= I_i \delta_{ij}^{c} + \Delta I_i \hat{n}_i^c \hat{n}_j^c + \Delta I_s \hat{n}_i^c \hat{n}_j^s + \Delta I_B \hat{n}_i^B \hat{n}_j^B, \quad (17) \\
I_{ij}^s &= I_i \delta_{ij}^{s} + \Delta I_s \hat{n}_i^s \hat{n}_j^s.
\end{align*}
\] (18)

For a slowly rotating object like a magnetar we would expect \( \Delta I_O \ll I_c \) and \( \Delta I_s \ll I_s \). The total angular momentum is
\[
J = \Omega \left\{ (I_c + \Delta I_O) \hat{n}_O^c + \Delta I_c \cos \theta_c \hat{n}_c^c + \Delta I_B \hat{n}_B^B \right\} + (I_s + \Delta I_s) \Omega \hat{n}_s^s.
\] (19)

This expression shows that the centrifugal deformations are not expected to play any significant role since they are simply “absorbed” in the pre-existing spherical pieces.

The second modification to the basic model, the anomalous torque \( N^a_\alpha \), will appear as a source term in the Euler equation (S). If the exterior space is assumed perfect vacuum and the magnetic field dipolar this torque can be written as
\[
N^a_\alpha = I_\alpha \Omega (\hat{n}_1^c \hat{m}_1) e^{ijk} \Omega_j \hat{m}_k,
\] (20)
where \( m^i \) is the external magnetic dipole moment and \( \epsilon_\alpha \) is an effective “deformation”
\[
\epsilon_\alpha = \frac{m^2}{I_\alpha R c^2}.
\] (21)
and \( R \) is the stellar radius. Expressed in terms of the field strength \( B_4 \) on the magnetic pole we obtain the estimate
\[
\epsilon_\alpha \approx 10^{-7} \left( \frac{B_4}{10^{15} G} \right)^2. \quad (22)
\]
Clearly, the impact of the anomalous torque on the stellar dynamics is important only for magnetar-strength fields. This was the central idea in the so-called radiative precession model advocated by Melatos (1999).

With the inclusion of the anomalous torque the equation of motion (S) takes the form
\[
J^i + \epsilon^{ijk} \Omega_j [J_k - \epsilon_\alpha I_\alpha \Omega (\hat{n}_1^c \hat{m}_1) \hat{m}_k] = 0.
\] (23)
Precession will not occur provided
\[
J^i = \lambda^i \hat{n}_1^c + \epsilon_\alpha I_\alpha \Omega (\hat{n}_1^c \hat{m}_1) \hat{m}_i.
\] (24)
It is natural to assume that the interior and exterior magnetic fields share the same symmetry axis (see discussion in the next section). Setting $\hat{n}_I = \hat{n}_B$ in (24) we recover the previous results (14)-(16) with the recalibrated deformations $\epsilon_B \rightarrow \epsilon_B - \epsilon_A$ and $\epsilon_A \rightarrow \epsilon_A + \Delta I/I_0$.

4 DISCUSSION

First consider equation (14), which applies in the limit of no pinning. The magnetic ellipticity $\epsilon_B$ is at least as large as would be estimated by substituting the external magnetic field strength into equation (4), i.e. for a typical magnetar $\epsilon_B \sim 10^{-6}$ for $\dot{B} = 10^{15}$ G. In reality this number could be significantly larger if we take at face value the most recent calculations of stable magnetic equilibria in neutron star models which suggest that the interior field $\dot{B}$ could be much stronger than the exterior field $B_e$ (Braithwaite 2009).

That the magnetars do not seem to be either aligned or orthogonal rotators means that the angle $\theta_B$ is neither close to zero nor close to $\pi/2$, and so the trigonometric factor on the right hand side of (14) is of order unity. It then follows that the crustal deformation must be of the order of the magnetic one: $\epsilon_c \sim \epsilon_B \gtrsim 10^{-6}$, or possibly even larger, if $\epsilon_B$ is small. The value $\epsilon_c \sim 10^{-6}$ is rather significant – it is close to the maximum allowed deformation predicted by detailed modelling of the crust, i.e. equation (1).

It follows that, in the absence of superfluid pinning, magnetar crusts are maximally, or close to maximally, strained. This idea sits well with the standard model of magnetar activity (Thompson & Duncan 1995, 1996, 2001) where the energy budget of the flaring activity is supplied by the evolving magnetic field. The crust itself acts as a gate, with crustal cracking acting as a trigger for the magnetic reconnection events. The maximally strained crust that applies in this scenario makes such events perfectly natural.

One might worry that one could evade our conclusion of a strongly strained crust by relaxing the assumption that the internal and external magnetic fields are aligned (i.e. $\hat{n}_I \neq \hat{n}_B$). Indeed, if the internal magnetic axis lies close to the rotation axis (making the angle $\theta_B$ small), equations (14) allow for a $\epsilon_c$ well below the maximum. However, a misalignment between the internal and external magnetic fields will generate tangential magnetic stresses, which can only be balanced by elastic ones. If both the internal and external fields are of strength $\sim B$, and if the field transitions from one geometry to the other over a shell of thickness $\Delta R < R$, the tangential magnetic stresses will be the order of $B^2/\Delta R$. This stress is greater than the $B^2/R$ that the crust needs to balance in our original scenario with aligned internal and external fields. Clearly, relaxing this assumption does not allow one to avoid our conclusion of a highly strained crust. Additional evidence supporting the alignment between the interior and exterior magnetic fields comes from simulations of stable magnetic equilibria (Braithwaite 2009).

Now consider the case with superfluid pinning (equations (15) and (16)). In this case no constraint is placed on the crustal deformation. In fact, $\epsilon_c$ could be set to zero, with no qualitative change in the non-precessional solution. The only constraint on stellar parameters that the pinning scenario would imply is that the pinning force itself is strong enough to sustain the small misalignment between the crustal and superfluid angular velocity vectors given by equation (15). As discussed by Link & Cutler (2002) the relative velocity between the superfluid and the crust generated by this misalignment produces a Magnus force on the pinned vortices. Making a simple estimate, the relative velocity generated by the misalignment is of order $\Delta v \sim \Omega R\theta_c$. Parameterising with magnetars in mind gives

$$\Delta v \sim 60 \text{ cm/s} \left( \frac{\Omega}{0.63 \text{ Hz}} \right) \left( \frac{\Delta I/I_0}{10^{-6}} \right) \left( \frac{10^{-2}}{\epsilon_c} \right).$$

This is to be compared with the critical velocity for unpinning. The estimates of this have been collected in Jones (2010) who finds a lower unpinning threshold $\sim 10^6 \text{ cm/s} \gg \Delta v$. Thus, the finite strength of vortex pinning is no obstacle to building our non-precessing magnetar models. Furthermore, the weak $\Delta v$ also ensures that the non-precession configuration is immune to the superfluid instability discussed in Glampedakis, Andersson & Jones (2009).

There is one special case that deserves some comment: when the crustal and magnetic axes are aligned, so that $\hat{n}_I = \hat{n}_B$. In this case the star is biaxial rather than triaxial, and precession would correspond to the magnetic axis rotating in a cone of half-angle $\approx \theta_B$ about the fixed vector $J$, with a rotation about $\hat{n}_B$ at the precession frequency superimposed. As is well known, such a precession would have no signature in the pulsations (Jones & Andersson 2001). Such a seemingly non-precessing star could have no pinned superfluid and any amount of crustal strain (even zero – in this case the star is still biaxial), and so would not fit into the scheme described above. How can we be sure that magnetars do not fall into this special case? Firstly, the absence of precessional modulation in the pulsations would require the hot-spot itself to be exactly axisymmetric about the magnetic axis. Secondly, it would be surprising if the crustal strains were symmetric about the magnetic axes, as the crustal strain field is likely to have retained some memory of a more oblate shape inherited from earlier in its life, before it was spun-down to the long periods of the magnetars. Thirdly, such a large-amplitude precession would be expected to damp due to internal dissipative processes (Jones & Andersson 2001). We therefore feel it is very unlikely that this special case is common in the magnetar population.

Having arrived at two different non-precession possibilities, described by equations (14) and (15), it is natural to ask the question of how one might hope to distinguish between them. As already noted, the no pinning case, with its maximally strained crust, sits naturally with the magnetar flare model, as quake events in the crust are required to trigger magnetic reconnection events. However, this is hardly decisive - diffusion of the magnetic field may, in either scenario, play a key role in triggering the bursts.

Nevertheless, one intriguing possibility does suggest itself. In the scenario with no superfluid pinning the crustal strain predicted just happens to lie at the upper end of what is believed to be possible on the basis of detailed crust modelling. Suppose this is no coincidence. It may be the case that there exist even more strongly magnetised stars; their magnetic deformations would necessarily be larger than their crustal ones. By eqn. (14) it follows that such stars, provided they are oblate ($\epsilon_B > 0$), would be nearly aligned rotators. There would be a selection effect against seeing such magnetars. So, observation of a tail of strongly magnetised but
nearly aligned rotators would argue in favour of there being no superfluid pinning in magnetars. Similarly, if the magnetic deformation were prolate ($\epsilon_\theta < 0$), strongly magnetised stars would be orthogonal rotators, although as noted above, the current observations do not suggest this to be the case. Note, however, regardless of the sign of $\epsilon_\theta$, the absence of such aligned/orthogonal populations would not necessarily imply that superfluid pinning occurs in magnetars; Nature may simply not provide neutron stars with magnetic field strengths higher than those currently observed.

The presence of a pinned superfluid is also related to the timing profile of magnetars. Indeed, several AXPs are known to have undergone one or more glitches with a fractional $\Omega$-jump comparable to the Vela-type glitches in radio pulsars (for a review see Rim, Kaspi & Gavriil (2008)). The occurrence of large glitches by itself suggests the presence of vortex unpinning. So far no glitches have been detected in SGRs. This should not come as a surprise given that these objects, when in their crust if some mechanism could cause vortex unpinning. As a result, the phase-no glitches have been detected in SGRs. This should not come as a surprise given that these objects, when in quiescence, are actually fainter than the AXPs as well as “louder” in terms of timing noise. As a result, the phase-coherent timing of SGRs required for the identification of glitches becomes a highly problematic procedure. We can, nevertheless, speculate on physical grounds as to why SGRs may not glitch. One possibility could be the absence of vortex pinning in the core, if the core is the only region where pinning could in principle occur. The interior magnetic field in SGRs could be sufficiently strong as to exceed the critical field $H_{c2} \approx 10^{16}$ G above which proton superconductivity is suppressed (Baym, Pethick & Pines 1969). The destruction of superconductivity would obviously imply the absence of magnetic fluxtubes on which the vortices could pin. If this scenario were to be true then the known SGRs would obey the non-precession condition (14) and would have maximally strained crusts.

It is interesting to consider the impact that a hypothetical discovery of magnetar precession would have on our principal conclusions. As we discussed in the introduction, current observations do not exclude the possibility that magnetars could be undergoing small-amplitude precession, masked by their timing noise. This precession could be either fast or slow (i.e. with a period, respectively, comparable to or much longer than the spin period) depending on whether there is a pinned superfluid. A detection of the former type of precession would amount to clear evidence of the existence of a pinned superfluid component within the star. Alternatively, detection of slow precession would mean that (i) there is no pinned superfluid, and (ii) that the star’s rigid body dynamics deviates only slightly from the non-precession condition (13), if the precession amplitude is sufficiently small. Thus, even in this case we would be able to make the same predictions for the magnetar structure as before.

To sum up, we have argued that the high magnetic strains that must necessarily exist in magnetars have interesting implications for their structure, telling us something about the level of crustal strain or the superfluid state of matter in their interior. Our modelling assumes a lack of significant free precession in the magnetars, and that their magnetic axes, as traced by the pulsation axes, are neither nearly aligned nor nearly orthogonal to their spin axes. We therefore conclude by pointing out the importance of better observational data to either confirm or reject these assumptions. In particular, our work shows that there would be a great deal of interest in attempting to construct detailed models of the geometry of magnetar emission, despite the technical difficulties involved. There would also be interest in searching for an ultra-highly magnetised but nearly aligned or orthogonal tail to the magnetar population. Clearly, more detailed observations of magnetar timing and pulse geometry could give us a unique insight into neutron star interiors.

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