Fermion zero-mode influence on neutron-star magnetic field evolution

P. B. Jones*

Department of Physics, University of Oxford, Denys Wilkinson Building, Keble Road, Oxford, OX1 3RH.

ABSTRACT

The quantum phenomena of spectral flow which has been observed in laboratory superfluids, such as \(^3\)He-B, controls the drift velocity of proton type II superconductor vortices in the liquid core of a neutron star and so determines the rate at which magnetic flux can be expelled from the core to the crust. In the earliest and most active phases of the anomalous X-ray pulsars and soft-gamma repeaters, the rates are low and consistent with a large fraction of the active crustal flux not linking the core. If normal neutrons are present in an appreciable core matter-density interval, the spectral flow force limits flux expulsion in cases of rapid spin-down, such as in the Crab pulsar or in the propeller phase of binary systems.

Key words: stars: neutron - stars: magnetic fields - pulsars: general

1 INTRODUCTION

It is now well-established that evolution and decay of the internal magnetic field, rather than rotational energy, is the dominant source of the persistent and burst X-ray luminosities in certain hot young isolated neutron stars. These are observed as the anomalous X-ray pulsars and soft-gamma repeaters, referred to as magnetars (for reviews see Harding & Lai 2006; Woods & Thompson 2006; Pons, Link, Miralles & Geppert 2007). In principle, a large part of the internal magnetic flux could be contained in the liquid core of the star at densities less than \(\sim 2\rho_0\), where \(\rho_0 = 2.8 \times 10^{14} \text{ g cm}^{-3}\) is the nuclear density. In this interval, many calculations of pairing and of the equation of state (Baldo, Burgio & Schulze 1998; Heiselberg & Hjorth-Jensen 2000) show that the core consists of neutrons, protons, relativistic electrons and possibly muons, and that the neutrons and protons are superfluids, pairing in the \(^3\)P\(_2\) – \(^3\)F\(_2\) and \(^1\)S\(_0\) states, respectively. The proton energy gaps found in these calculations indicate a type II superconductor (Baym, Pethick & Pines 1969) However, Baldo & Schulze (2007) have since shown that the inclusion of more refinements, including three-body forces, reduces the energy gap to values consistent with type I superconductivity for certain proton densities. It is also the case that other authors have more recently considered the possibility of type I superconductivity with various motivations (Sedrakian, Sedrakian & Zharkov 1997; Buckley, Metlitski & Zhitnitsky 2004; Sedrakian 2005; Alford, Good & Reddy 2005; Charbonneau & Zhitnitsky 2007; Alford & Good 2008). But it remains a reasonable assumption that type II superconductivity exists for at least part of the density interval we are considering here.

In the type II system, the magnetic flux is quantized in units of \(\phi_0 = \hbar c/2e = 2.07 \times 10^{-7} \text{ G cm}^2\) and is confined to the cores of proton vortices (see, for example, Tinkham 1996). Steady-state motion of the proton vortices is determined by balance of the Magnus force, electromagnetic interactions, external forces including buoyancy forces (Muslimov & Tsygan 1985; Jones 2006) and interaction with vortices of the rotating neutron superfluid (Muslimov & Tsygan 1985; Sauls 1989; Jones 1991; Ruderman, Zhu & Chen 1998), and forces derived from interactions between the quasiparticles of the system. The distinct populations of superfluid quasiparticles concerned are those localized in the cores of proton vortices and those delocalized in the continuum. Previous vortex velocity calculations (Jones 2006) assumed both the internal temperatures \(T < 10^8 \text{ K}\) of typical radio pulsars and negligibly small neutron and proton quasiparticle number densities. Under these conditions, proton vortices and therefore the internal flux distribution could be quickly expelled to the core-crust boundary.

The low-temperature assumption is not appropriate for magnetars whose X-ray luminosities in the active phase, \(\sim 10^3 - 10^4\) yr after formation, are thought to be a consequence of plastic flow of the crust (with its frozen-in field) under very high magnetic stresses (Harding & Lai 2006; Woods & Thompson 2006). This dissipative process must maintain a large input of thermal energy to the inner crust and core of the star so that internal temperatures are likely to be much higher than those inferred for the isolated radio pulsars. Neglect of continuum superfluid quasiparticles is not justified at these higher temperatures unless it can be es-
established that the neutron energy gap is large at all relevant wavenumbers. Although early calculations gave such values (Amundsen & Østgaard 1985), much work (see the review by Heiselberg & Hjorth-Jensen 2000) finds a small angle-averaged $^{3}$He-$^{3}$F$_{2}$ neutron energy gap, $\Delta_{n} \sim 0.1$ MeV, at the wavenumbers concerned. More recently, Khodel, Khodel & Clark (2001) also obtained similar small values for certain neutron-neutron interactions. But there is also the possibility of pion condensation at a matter density $\sim 2\rho_{0}$. Its effect has been investigated by Khodel et al (2004) and is to increase the gap in the neutron single-particle spectrum. Nevertheless, we regard the existence of a small neutron superfluid gap, or possibly even a normal Fermi liquid, as being not improbable for some wavenumber bands within the density interval considered in this paper.

A new development in condensed-matter physics, the phenomenon of spectral flow described by Volovik (1996; see also the review of Kopnin 2002), has greatly clarified and extended our understanding of the forces acting on moving vortices and has been confirmed by measurements on the $^{3}$He-B phase (Beevan et al 1997). The present paper applies these ideas to show that forces generated by spectral flow limit the velocities of proton vortices in the outer core of hot young neutron stars to an extent that is consistent with a large fraction of the active flux not linking the core.

2 SPECTRAL FLOW

The fermion zero-mode excitations concerned are those quasiparticles localized within an isolated vortex core at energies below $\Delta_{p}$ (referred to the chemical potential as zero). Approximate eigenvalues (Caroli & Matricon 1965; de Gennes 1966) in the limit $p_{F}\xi \gg 1$, where $\xi$ is the superfluid coherence length and $p_{F}$ the Fermi wavenumber, are $E_{pp}$, where $\mu$ is the (half-integer) angular momentum component parallel with the vortex axis $\hat{z}$, and

$$E_{pp} = -\mu \hbar \omega_{0},$$
$$\hbar \omega_{0} = \frac{\hbar^{2}p_{F}}{m_{p}^{*}c_{p_{F}}} = \frac{\pi \Delta_{p}^{2}p_{F}}{2E_{pp}p_{\perp}} \ll \Delta_{p}. \quad (1)$$

(This expression neglects interaction with the core magnetic flux; see Bardeen et al 1969). The proton bare and effective masses are $m_{p}$ and $m_{p}^{*}$ and its Fermi energy is $E_{Fp}$. The components of the proton wave-vector parallel with and perpendicular to the vortex axis are $p$ and $p_{\perp}$. Values of $\mu$ are large at the temperatures of Table 1, so that localized quasiparticles interacting with an external heat-bath (consisting, in this case, of a distribution of continuum quasiparticles) can be treated by a Boltzmann equation (see Kopnin 2002) in which the quantum number $\mu$ is replaced by a continuous variable. In this semi-classical model, the vortex structure is that of a cylindrical hole in the superfluid order parameter and quasiparticle states within this hole are normal particle and hole orbits in the form of chords limited at each end by Andreev reflection (Andreev 1964), the specific process of reflection that occurs at a normal-superfluid interface.

The perturbation of the Boltzmann distribution function caused by vortex motion consists of a displacement of the Fermi sphere in the plane perpendicular to $\hat{z}$ which also contains the vortex velocity $\mathbf{v}_{L}$. It relaxes through binary collisions with either localized (vortex-core) quasiparticles or with continuum quasiparticles. Thus the relaxation rate is given by the transport relaxation times, $\tau^{-1} = \tau_{loc}^{-1} + \tau_{c}^{-1}$, but only the latter process transfers momentum between the vortex and the continuum. The existence of these two distinct relaxation processes means that, within the semi-classical model, the localized quasiparticles mediating the force between the continuum heat-bath with bulk velocity $\mathbf{v}_{T}$ and a vortex of velocity $\mathbf{v}_{L}$, have a steady-state bulk velocity $\mathbf{v}_{loc} = \mathbf{v}_{L}/\tau_{loc} + \mathbf{v}_{T}/\tau_{c}$. This lags behind $\mathbf{v}_{L}$ in general, and in the limits $\tau_{c}$ or $\tau_{loc} \rightarrow \infty$ tends, as expected, to $\mathbf{v}_{L}$ or $\mathbf{v}_{T}$. Thus spectral flow (Volovik 1996) depends on the velocity difference

$$\mathbf{v} = \mathbf{v}_{L} - \mathbf{v}_{loc} = (\mathbf{v}_{L} - \mathbf{v}_{T})\tau/\tau_{c}. \quad (2)$$

In Andreev reflection, a particle (hole) is retroreflected as a hole (particle), so maintaining the orientation of the chord. It has angular momentum $\mu \equiv (r - \mathbf{v}t) \times \mathbf{p}_{\perp},$ where $r - \mathbf{v}t$ is its coordinate in the vortex rest frame. The retroreflection is not exact owing to the motion of the superfluid so that the chord precesses about the vortex axis with angular frequency $\omega_{0}$ in a direction opposite to that of the circulating superfluid (Stone 1996). Motion of the vortex induces a finite rate of change of orbital angular momentum, given by the classical variable $\dot{\mu} = -[\mathbf{v} \times \mathbf{p}_{\perp}] \cdot \hat{z}$, that drives spectral flow, which is the movement of quasiparticles both to and from the negative energy continuum $E_{pp} < 0$ along the continuous dispersion relation given by equation (1). For $-\mu \sim 10$ or more, we rely on the correspondence principle in assuming that classical-orbit quasiparticles give an adequate description of spectral flow. In the case of the physical discrete spectrum, with half-integer $\mu$, the process of spectral flow relies on the existence of sufficient level-broadening to allow quasiparticles to move along the dispersion relation by hopping from state to state. We refer to Stone (1996) and Kopnin (2002) for details of the relationship between the continuum and discrete-$\mu$ cases. Solution of the Boltzmann equation, followed by integration of the rate of change of quasiparticle momenta over $\mu$ and $\mathbf{p}_{\perp}$, gives the force acting on the vortex, including the Magnus and spectral flow components. We do not give further details here because the review by Kopnin (2002; pp. 1654 - 1660) contains a complete and very transparent account.

On this basis, the force-balance equation per unit length of vortex is

$$\mathbf{f}_{B} + \mathbf{f}_{V} + \mathbf{F}_{M} + \mathbf{F}_{I} + \mathbf{F}_{s} = 0, \quad (3)$$

in which $\mathbf{f}_{B,V}$ are the external forces caused by buoyancy and interaction with neutron vortices. The Iordanskii force (Thouless, Ao & Niu 1996; Sonin 1997) is unimportant here and has been neglected in equation (3). The Magnus force is expressed in terms of the proton electromagnetic current density $J_{p}^{e}$ obtained from the entrainment coefficients (Andreev & Bashkin 1975; Alpar, Langer & Sauls 1984; Jones 1991), $\mathbf{v}_{L}$ and the superfluid bulk velocities,

$$\mathbf{F}_{M} = \frac{1}{c}J_{p}^{e} \times \phi_{0}. \quad (4)$$

A similar expression gives $\mathbf{F}_{I}$, the electromagnetic interaction with electrons and muons, in terms of the lepton current density $J_{l}$. The spectral flow force has components parallel
with and perpendicular to $v_L - v_T$,
$$F_{sf\perp} = -C \frac{1}{1 + \frac{\omega_0^2}{\omega_T^2} \tau_c} \hat{z} \times (v_L - v_T),$$
(5)
$$F_{sf\parallel} = -C \frac{\omega_T}{1 + \frac{\omega_0^2}{\omega_T^2} \tau_c} (v_L - v_T),$$
(6)
(Volovik 1996; Kopnin 2002), with
$$C = \frac{\pi \hbar \rho_p}{m_p} \tanh \left( \frac{\beta \Delta_p}{2} \right),$$
(7)
where $\beta^{-1} = k_B T$. The total proton density is $\rho_p = \rho_p^s + \rho_p^n$, but at the temperatures considered here, the normal component $\rho_p^n$ is negligibly small compared with the superfluid $\rho_p^s$.

The simple expressions given by equations (5) and (6) are valid at temperatures such that $\omega_0^2 \tau_c \gg 1$, which we shall find to be always satisfied here. The definition of relaxation times follows Kopnin and Lopatin (1997).

At this point, it is interesting to note that frictional forces on vortices in astrophysical superfluids have been estimated previously from the scattering of continuum quasiparticles by core quasiparticles, with neglect of spectral flow. Following this procedure, Sedrakian (1998) found an expression for the frictional force on neutron vortices arising from the scattering of continuum proton quasiparticles. In the present context, the order of magnitude of such a force would be
$$F_{\text{coll}}^\parallel = \left( \frac{m_p v_L}{\tau_c} \right) \left( \frac{p F_k b T}{\pi \hbar \omega_0} \right).$$
(8)

This is smaller than the spectral flow component, equation (6), by a factor of the order of $k_B T / E_F p$, where $E_F p$ is the proton Fermi energy, and so is negligible by comparison. It follows that the frictional force produced by the scattering of electrons or muons is also negligible because the transport relaxation time for this process is several orders of magnitude longer than the strong-interaction values of $\tau_c$.

3 APPLICATION TO PROTON VORTEXES

Spectral flow is a remarkable quantum phenomenon whose recognition has enabled the understanding of laboratory measurements of the force acting on moving $^3$He-B vortices (see, for example, Bevan et al 1997). Its predictions are transparent in the semi-classical modelling of what would otherwise be an extremely cumbersome problem. The ratio of vortex radius to average interparticle spacing is given by $p F_k \xi$, which, for $^3$He-B, has values in the interval $10^{-2} - 10^{-3}$. Typical values in the proton superfluid are perhaps an order of magnitude smaller. For example, from the energy gap obtained by Baldé & Schulze (2007; the curve in Fig. 2 of that paper containing all corrections) we find $p F_k = 7.5$ at $\rho_0 \equiv 0.16 \text{ fm}^{-3}$, but this increases rapidly thereafter as a function of matter density. Use of the Boltzmann equation does require that the number of quasiparticles be large, and this condition is satisfied if we accept that the effective number is that contained within a length of vortex at least an order of magnitude greater than $\xi$. The circumstances existing in the core of a neutron star also differ from those, for example, of Bevan et al (1997) in relation to the problem of defining the heat bath and its velocity. Apart from collective excitations, the heat bath contains the lepton-photon system and continuum quasiparticle excitations of the neutron and proton superfluids. It is also possible that the neutrons may exist as a normal Fermi liquid within limited intervals of density. All these components interact electromagnetically or strongly with the nuclei embedded in the solid crust of the star and, in the absence of vortex motion, would be stationary with respect to it. But any component more strongly coupled with moving vortices than with the crust would be expected to flow with them. The neutrons are the most important component and in their case we shall follow the argument given by Kopnin (2002; Sect. 9 of that paper) in relation to the analogous problem in $^3$He-B and assume that the continuum quasiparticle rest frame is stationary with respect to the crust provided their mean free path for scattering by other continuum quasiparticles is shorter than the mean free path for scattering by localized proton quasiparticles. In a steady state of vortex motion, the force $f_d + f_V$ is transferred to the neutrons and must be balanced by a neutron chemical potential gradient. But this is very small, of the order of $10^{-6} - 10^{-5} \text{ eV cm}^{-1}$.

Equation (4) has been expressed as shown to emphasize that $F_M + F_\ell = 0$ because screening of the electromagnetic current in the whole volume of the neutron-star core must be presumed for a type II superconductor (Jones 1991, 2006) except for the microscopic circulating currents of the vortices. The distribution of leptonic current density is present throughout core and crust prior to the superconducting transition and in non-superconducting regions, can change only on the same long time-scales as the magnetic flux density. Thus the superconducting transition must be to a state of finite supercurrent density such as to satisfy the condition $J_{\rho} + J_1 = 0$ everywhere in the core. In effect, the lepton-proton system is an insulator and the moving-vortex induction field induces no large-scale current density. But $J_{\rho}^0$ is confined to the core of the star, unlike $J_1$, and the return currents that complete the circuit flow as sheets on the surface of the spherical superconducting volume and exclude the magnetic flux present in adjacent non-superconducting regions. In having volume screening, neutron stars differ from laboratory superconductors for which screening is primarily a surface phenomenon (see Tinkham 1996).

In the low temperature limit, $\omega_0^2 \tau_c \gg 1$, only a small longitudinal component of the spectral flow force, given by equation (6), remains. Owing to the cancellation of the Magnus force, equation (3) reduces to $f_d + f_V + F_{sf\parallel} = 0$, showing that movement of proton vortices, and hence expulsion of the field, occurs easily under the influence of buoyancy forces or of interaction with outward-moving neutron vortices as the rotation of the star slows (Jones 2006). At higher temperatures, the existence of spectral flow radically changes this behaviour.

It is convenient to take account of both longitudinal and transverse components of the spectral flow force by obtaining an expression for the component of $v_L - v_T$ parallel with $f_d + f_V$ (5) and (6) are valid, this component is independent of $\tau$ and is,
$$\left( v_L - v_T \right)_\parallel = \frac{\omega_0^2 \tau_c}{C} \left( f_d + f_V \right).$$
(9)
Neutrons are the most important heat-bath component. The relaxation time $\tau_c$, for interaction with superfluid neutrons,
in which \( \bar{\sigma} \) is chosen at a laboratory energy equal to the sum of the Fermi energies, \( E_{p_n} + E_{F_p} \) (for this, the angle-averaged centre-of-mass energy in the medium equals that for the free-particle scattering). The function \( W \) is \( W = 1 \) in the case of a normal Fermi system, but for an isotropic superfluid is given by,

\[
\frac{1}{W} = \frac{8}{\pi^2} \int_0^\infty dx \int_0^\infty dy \frac{X^2}{1 + e^{-x} 1 + e^{-1} 1 + e^{-y}},
\]

in which \( X^2 = x^2 + (\beta \Delta_n)^2 \), with an identical definition of \( Y \). The numerical values of \( W \) obtained in the interval appropriate for Table 1 are satisfactorily fitted by \( W = 0.8(\beta \Delta_n)^{-1} \exp(\beta \Delta_n) \).

The effective component of the vortex drift velocity is approximated by the spatially-averaged density of localized proton quasiparticles and \( \Delta_n \) is assumed isotropic. It is limited by a critical value \( \omega_0 \tau_c \) at which neutron quasiparticle rest frames move with the vortices rather than being stationary with respect to the crust can be answered, for the temperatures \( T \) and \( \Delta_n \) of the Table, by noting that \( \sigma_{np} \) and the proton-proton cross section \( \sigma_{pp} \) at the equivalent laboratory energy equal to \( 2E_{F_p} \) differ only by a factor of order unity. (See Yao et al 2006; pp. 339-340. For example, \( \sigma_{pp} \approx 150 \) mb at 20 MeV, whereas \( \sigma_{np} \approx 100 \) mb at 70 MeV.) The Kopnin condition is approximately satisfied provided the ratio of the neutron continuum quasiparticle density to the spatially-averaged density of localized proton quasiparticles is much less than unity. The effect is to be emphasized that there is no exponential dependence on \( \omega_0 \) and \( \Delta_n \) is consistent with recent calculations (Heiselberg & Hjorth-Jensen 2004). It is limited by a critical value \( f_{v_v} \) at which neutron vortices cut through the proton-vortex lattice, whose magnitude is at most of the same order as the buoyancy term (Jones 1991). If the outward radial velocity of neutron vortices during spin-down exceeds the velocity given by substituting \( f_B + f_{v_v} \) in equation (9), they will move by cutting through the proton vortex lattice in those cases where a relative sliding motion is possible.

General predictions are, of course, impossible to make owing to the exponential dependence of spectral flow on \( \Delta_n \). Together with \( \Delta_n \), this function of density also controls the core neutrino emissivity and hence the internal temperature of isolated neutron stars less than \( \sim 10^7 \) yr in age (see for example, Fig. 5 of Yakovlev & Pethick 2004) which lack any substantial source of internal dissipation. It is quite possible that values of \( \Delta_n \) larger than those assumed in Table 1 exist in certain density intervals, but it is the regions of small \( \Delta_n \) that have the greater effect on average drift velocities. The choice of \( ^3P_2 - ^3F_2 \) neutron energy gaps in the Table is consistent with recent calculations (Heiselberg & Hjorth-Jensen 2000; Dean & Hjorth-Jensen 2003). At \( k_F \sim 2.0 \) fm\(^{-1} \) and above, higher partial waves make little contribution to attractive forces and it is even possible that there are significant intervals in which the neutrons are not superfluid.
4 CONCLUSIONS

It is of some interest that a new quantum phenomenon recognized in relation to laboratory studies of vortices in $^3$He should have some bearing on neutron-star magnetism. Subject to the reservations described in Sect. 3, the conclusions of this paper that follow from the spectral flow phenomenon are as follows. The presence of charged or neutral baryons as a normal Fermi system within some interval of core-matter density leads to $\omega_0\tau_c$ values of the order of those in the second column of Table 1. These are some orders of magnitude smaller than the $\omega_0\tau_c \approx 5 \times 10^4$ needed for velocities $\sim 3 \times 10^{-7} \text{ cm s}^{-1}$ that give significant magnetic flux transfer from core to crust within the $10^4$ yr magnetar phase. The absence of substantial movement of flux from core to crust in this time interval would be consistent with a field configuration in which the distribution of the active flux is confined to the crust at or close to the time of neutron star formation, its ability to move and evolve being unconstrained by linkage with the core. This is also true, even in the absence of normal neutrons, for the $\Delta_p$ of Table 1 because the higher temperatures there may well be maintained by internal dissipation during the active magnetar phase. Comparison of the continuum and localized quasiparticle densities shows that there can be dragging of the heat-bath frame $v_T$ by $v$, much reducing spectral flow, particularly for very high fields, $B \sim 10^{15}$ G, and for $\Delta_p$ near the minimum consistent with type II superconductivity, but our broad conclusions are unchanged.

There are other circumstances in which spectral flow may be of relevance, such as the case in which the density-dependence of neutron-pairing allows both normal and superfluid neutrons within appreciable intervals of matter density. The values of $\omega_0\tau_c$ given in the second column of Table 1 for modest internal temperatures $T \approx 10^8$ K are small enough to constrain the proton vortex drift velocities that are possible in cases of rapid spin-down such as the Crab pulsar, and the propeller phase of binary systems. The outward movement of neutron vortices in the superfluid region then occurs by sliding relative to the proton vortices or by intersecting them, and is unable to force the outward movement of magnetic flux.

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