

Magnetic Fields in High-Density Stellar Matter

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Abstract. I briefly review some aspects of the effect of magnetic fields in the high density regime relevant to neutron stars, focusing mainly on compact star structure and composition, superconductivity, combustion processes, and gamma ray bursts.

THE “MAGNETIC FIELD” OF “NEUTRON STARS”

A large body of evidence now identifies the presence of strong magnetic fields at (or very near) the surface of neutron stars. Except for a few cases, all the estimations of the magnetic fields of these objects come from the timing measurements of pulsar slow-down. The pulsar is usually modelled as a rapidly rotating neutron star with a dipolar magnetic field configuration in which the rotation axis is not aligned with the dipole axis. Since the period of the pulsed emission is associated with the rotation period P , accurate observations of P and its time derivatives \dot{P} , \ddot{P} , and even the third derivative of P , have been used to shed light on pulsar dynamics. Within this framework the dynamic equation reads $\dot{\Omega} = K\Omega^3$, with $K \equiv (R^6 B_p^2 \sin^2 \alpha)/(6c^3 I)$, and $\Omega = 2\pi/P$. Therefore, some properties of pulsars are plausible of interpretation in terms of timing data, e.g.:

- the component of the B -field *at* the magnetic pole *along* the rotation axis

$$B_{\perp} = B_p \sin \alpha = \sqrt{\frac{6c^3 I}{R^6} \frac{P \dot{P}}{(2\pi)^2}} \quad (1)$$

- the pulsar age t (or alternatively the rotational period at birth P_i),

$$t = \tau \left(1 - \frac{P_i^2}{P^2} \right) \quad (2)$$

where $\tau \equiv -P/2\dot{P}$ is the so called characteristic age, I is the moment of inertia of the neutron star, R its radius, and α is the angle between the rotation axis and the dipole axis. Note that, for determining B_{\perp} the values of I and R must be assumed based mostly on theoretical arguments.

A test of the reliability of these estimates, can be obtained from an independent determination of the age of the pulsar, which is possible for pulsars associated with supernova remnants (SNRs). Assuming that a given pulsar and a given SNR were born together, we can identify t_{SNR} (the SNR age) with the pulsar age t in Eq. (2). Since within the dipole model τ is an upper limit for t , a comparison of t_{SNR} with τ provides a relevant

test for pulsar dynamic models. Another interesting feature is that also P_i can be inferred from timing observations through Eq. (2). For some associations it is found $t_{SNR} \leq \tau$. The straightforward interpretation in these cases is that the pulsar was born with a small P_i (probably close to the mass shed limit P_{break}) and therefore $t \leq \tau$. The cases for which $t_{SNR} \ll \tau$ are usually understood by assuming that the pulsar was born with P_i very close to the present value of P . The associations with $t_{SNR} > \tau$ cannot be explained within the dipolar model. This is the case of, e.g.,

- PSR J1846-0258 with a characteristic age $\tau = 700$ yrs while the associated SNR Kes75 has an inferred age $t_{SNR} > 2000$ yrs [7],
- the association PSR B1757-24 / SNR G5.4-1.2 [8].

These discrepancies are usually attributed to different causes:

- the association is false, e.g. the pulsar could be a foreground object unrelated to the SNR, or the estimated SNR age is no reliable (see [8]),
- the pulsar spin-down torque is not due to pure magnetic dipole radiation, but other physical agents contribute to the spin down.

Another test is possible if B is known independently. A case is 1E1207.4-5209 [6] for which $B \sim 8 \times 10^{10}$ G has been determined from the observation of cyclotron lines (in disagreement with the dipole model that predicts $B_{\perp} = 2 - 3 \times 10^{12}$ G). Note that the characteristic age of 1E1207.4-5209 is $\tau \approx 4.5 \times 10^5$ yrs while the age of the associated SNR PKS 1209-51/52 is only $t_{SNR} < 10^4$ yrs [6, 13]. Then, from the point of view of the age determination the association does not appear as contradictory with the dipole model since the age can be adjusted assuming $P_i \sim P_0$. This stresses the obvious fact that $t < \tau$ is a necessary condition to believe in the dipole model, but not a sufficient one.

The *braking index* defined by $n \equiv \Omega \ddot{\Omega} / \dot{\Omega}^2$ is the quantity most usually employed to study the pulsar braking behavior to second order. Within the dipole model n should be equal to 3 at any time. Accurate determinations of n_{obs} have been reported in the literature only for five pulsars. The observed values are all smaller than 3, i.e. in contradiction with the dipole model. A number of scenarios have been proposed to explain the discrepancies. In general, they are attributed to a time variation of $K \equiv (R^6 B_p^2 \sin^2 \alpha) / (6c^3 I)$ associated with: a decoupling of a portion of the star's liquid interior from the external torque [9], or the time evolution of the moment of inertia I due to phase transitions in the core of a rotating neutron star [10], variation in the angle α between the magnetic and rotation axes [11], or magnetic field decay [12]. There are also many pulsars for which large (and in some cases negative) measured values of $\ddot{\Omega}$ and n_{obs} have been obtained. These disagreements are thought to be caused by timing noise and the recovery from glitches [2]. Although a definitive explanation for the disagreement between theory and observations is still lacking, it seems clear that other physical agents than pure dipolar emission govern the pulsar dynamics (at least) to second order.

Another interesting fact is that there are a few radio pulsars with *inferred* B -fields above the quantum critical field $B_{ce} = m_e^2 c^3 / e \hbar \simeq 4.41 \times 10^{13}$ G, above which some models [15] predict that radio emission should not occur [4, 14]. This group of “high-field” radio-pulsars may actually constitute a larger fraction of the total pulsar population since there are selection effects that make difficult the detection of long period radio

pulsars. This represents a challenge for pulsar emission models (see Melrose in this volume, and references therein). Alternatively, the high inferred B -fields could be simply the result of the very schematic picture offered by a purely dipole model. In this sense, other dynamic descriptions (e.g., pulsar disks, propellers [16]) may gain new interest.

EFFECTS OF B ON NEUTRON STAR STRUCTURE AND COMPOSITION

Since the fields in the interior of neutron stars can be even larger than those believed to exist at the surface, it is necessary to explore the properties of bulk matter in the presence of large and also huge magnetic fields. An upper limit to the neutron star magnetic-field strength follows from the virial theorem of magneto-hydrostatic equilibrium: $2T + 3\Pi + W + \mathcal{M}$. Since $2T + 3\Pi$ is always positive, the magnetic energy of the neutron star can never exceed its gravitational binding energy. For a star of mass M and radius R this gives $(4/3\pi R^3)(B^2/8\pi) \sim 3GM^2/5R$, or

$$B_{max} = 1.4 \times 10^{18} \left(\frac{M}{M_\odot} \right) \left(\frac{R}{10\text{km}} \right)^{-2} \text{ G.} \quad (3)$$

Concerning the equation of state, the essential distinction between zero and non-zero magnetic fields comes from the different way of counting the energy states of the particles due to the modification of the phase space introduced by B . Magnetic fields quantize the orbital motion of the charged particles into Landau levels. For a non-relativistic free particle of charge q and mass m in a constant magnetic field (along the z axis), the kinetic energy of the transverse motion becomes $E_\perp \rightarrow (n_L + 1/2)\hbar\omega_c$ with $n_L = 0, 1, 2, \dots$, where $\omega_c = |q|B/(mc)$ is the cyclotron frequency. The transverse motion of the electrons becomes relativistic when $\hbar\omega_c > m_e c^2$, i.e. for magnetic fields larger than the critical value:

$$B_{ce} = \frac{m_e^2 c^3}{q_e \hbar} = 4.41 \times 10^{13} \text{ G.} \quad (4)$$

For such large fields the free-particle energy has the following relativistic expression:

$$E = [c^2 p_z^2 + m^2 c^4 (1 + 2n_L \beta)]^{1/2} \quad (5)$$

with $\beta = B/B_{ce}$ [17].

At high densities, the magnetic-field effect on the bulk equation of state is expected to become negligible as more and more Landau levels are filled. However, where the density is not too high, the electrons reside mainly in low-lying Landau levels and important modifications are expected with respect to the zero-field case. The critical density ρ_B below which the effect of Landau quantization is important can be roughly estimated by considering a degenerate electron gas in a constant magnetic field. The electron Fermi momentum p_F is obtained from $n_e = (eB/hc)(2p_F/h)$. The electrons occupy only the ground Landau level when the Fermi energy $p_F^2/(2m_e)$ is less than the cyclotron energy $\hbar\omega_c$. From this we find $\rho_B = 7.09 \times 10^3 Y_e^{-1} B_{12}^{3/2} \text{ gcm}^{-3}$, being

$Y_e = Z/A$ the number of electrons per baryon. In general, we can use the condition $\rho_B > \rho$, or $B_{12} > 27(Y_e \rho_6)^{2/3}$, to estimate the critical value of B above which Landau quantization will affect physics at a density ρ .

From this simple analysis we see that in neutron stars, magnetic fields are important for understanding the behavior of the surface layers. The surface layers are expected to be composed of ^{56}Fe (residual from the supernova explosion that originated the neutron star), and also of lighter elements that can fall onto the NS surface from the supernova remnant, from the surrounding interstellar medium, or from a binary companion. This, together with eventual nuclear reactions and weak interactions at the surface during accretion, may result in a quite complex composition. On the other hand, strong gravity separates the accreted elements [18]. The lightest ones go up, and a hydrogen-helium envelope is likely to form on the top of the surface unless it has completely burnt out. The superficial layers can be in a gaseous state (atmosphere) or a condensed state (liquid or solid surface), depending on surface temperature, magnetic field and chemical composition [20]. For a hydrogen-helium envelope, a condensed surface forms at low temperatures ($T < 10^6$ K) and very strong magnetic fields ($B > 10^{15}$ G) while for iron it depends on the cohesive properties of the iron condensate [19].

On the other hand, it has long been recognized that neutron stars are sources of soft X rays during the $\sim 10^5 - 10^6$ years of the cooling phase after their birth in supernova explosions. A systematic study of X-ray emission from isolated neutron stars has shown that in general there are two different components of the X-ray emission, thermal and non-thermal. The non-thermal component with a power-law spectrum, is believed to originate from the pulsar's magnetosphere, while the thermal component is emitted from the surface layers of the neutron star. Magnetic fields affect dramatically the emergent spectra of the thermal emission reflecting the changes in the atomic structure, and may also produce a strong polarization of the emission. Also, the geometry of the magnetic field may introduce an angular dependence of opacities and equations of state that lead to an anisotropic cooling and consequently a strong anisotropy of the thermal emission [21, 22]. An accurate knowledge of the modifications introduced by the magnetic field in the thermal component is particularly interesting because this radiation contains important information about neutron star properties such as its surface temperature, magnetic field, and chemical composition, as well as the NS radius and mass, that are ultimately related to the internal composition (see [20] for a detailed analysis).

The inelastic scattering off electrons which are quantized in the strong magnetic fields of neutron stars produce cyclotron lines in the spectra providing a tool for direct measurements of neutron star magnetic fields. Cyclotron resonances have been seen in a handful of neutron stars in binary systems, the first one being Hercules X-1. Only recently the first direct measure of an isolated neutron star magnetic field has been possible [6]: A long X-ray observation of 1E1207.4-5209 with XMM-Newton has shown three distinct features, regularly spaced at 0.7, 1.4 and 2.1 keV which are interpreted as cyclotron absorption from electrons, resulting a magnetic field strength of 8×10^{10} G [6]. As mentioned in the Introduction, this field strength differs by two orders of magnitude from the value obtained from timing measurements of the pulsar slowdown, challenging the most accepted models of pulsar dynamic behavior.

It is interesting to check whether magnetic fields are relevant near the center of the

neutron star core since huge field strengths may be expected in this region. Near the center, the density is well above the nuclear saturation density ($\rho_0 = 2.7 \times 10^{14} \text{g cm}^{-3}$) and matter may be assumed to be composed mainly by neutrons and protons, with a substantial amount of other charged and uncharged baryons, such as the hyperons Λ^0 , Σ^+ , Σ^0 , Σ^- , Ξ^- , and Ξ^0 (but see below). These baryons are in β -equilibrium with negatively charged leptons (electrons and muons). In this context, Landau quantization leads to a softening of the EOS relative to the case in which magnetic fields are absent, and increases proton and lepton abundances (above $5 \times 10^{18} \text{G}$). Besides Landau quantization, it is also important the effect of polarizing the spin (or magnetic moments) of the nucleons and hadrons. This produces an overall stiffening of the equation of state and a further suppression of hyperons. In strong magnetic fields, contributions from the anomalous magnetic moments of nucleons and hadrons must also be considered. All these effects can be incorporated in a relativistic mean-field equation of state including not only nucleons but also hyperons [50]. General relativistic magneto-hydrostatic calculations with this equations of state show that maximum average fields within a stable neutron star are limited to values $B \sim 1 - 3 \times 10^{18} \text{G}$, in agreement with the simple estimates given above. Predictably, this field strength is not large enough to influence particle compositions or the matter pressure considerably.

SUPERFLUIDITY AND SUPERCONDUCTIVITY

There are two qualitatively different ideas about the magnetic field inside neutron stars, namely: a) it is present in the entire star (core and crust), b) it is located mainly in the crust. This depends strongly on the mechanism responsible for the generation of B (see Reisseneger in this Volume) and on the appearance of superconducting phases in neutron star matter (this Section).

Since neutron star interiors behave in most cases as highly degenerate Fermi systems, ($k_B T \ll \mu$, being μ the chemical potential of neutron star matter), the presence of any attractive channel in the particle interactions is expected to originate a superfluid (and a superconductor in the case of charged particles). There are two different classes of superconductors, which affect the magnetic field differently. Type I superconductors exhibit a very sharp transition to a superconducting state and perfect diamagnetism, the ability of expelling an already present magnetic field from the superconductor when it is cooled below the critical temperature (the so called Meissner effect). Type II superconductors differ from type I in that their transition from a normal to a superconducting state is gradual across a region of partially superconducting behavior (characterized by two critical fields H_{c1} and H_{c2}). For magnetic fields smaller than H_{c1} there exists a Meissner effect. For $H_{c1} < H < H_{c2}$, type II superconductors show a partially superconducting behavior. In this range of H there exist some rather novel “mesoscopic” phenomena like flux-lattice vortices, widely observed and studied in terrestrial fluids. Flux-lattice vortices are tubes of electrical current induced by an external magnetic field where superconductivity is suppressed. For $H > H_{c2}$ the vortices are too dense to maintain the condensate, and the normal phase is once again restored.

From a theoretical point of view, the criterion that determines whether a superconductor is of type I or type II is the Ginzburg-Landau parameter κ , obtained from the Ginzburg-Landau equations [24]. In their original paper, Ginzburg and Landau [24] analyzed the energy of the interface between a normal and a superconducting phase, kept in equilibrium in the bulk by an external magnetic field at the critical value. They found analytically that the surface energy vanishes at $\kappa = \kappa_c = 1/\sqrt{2}$. The physical meaning of this critical value was clarified further by Abrikosov [25]: it represents the demarcation line between type I ($\kappa < 1/\sqrt{2}$) and type II ($\kappa > 1/\sqrt{2}$) superconductors.

In neutron stars, several types of baryon pairing are believed to appear due to the attractive component of the nuclear force which allows binding into Cooper pairs through the well-known BCS mechanism. In the inner crust region the 1S_0 partial wave of the nucleon-nucleon interaction is attractive and neutrons will form 1S_0 pairs. This happens in the range $10^{-3}\rho_0 < \rho < 0.7\rho_0$, where $\rho_0 = 2.7 \times 10^{14} \text{g cm}^{-3}$ is the saturation density of symmetric nuclear matter [26]. In the core region having $\rho > 0.7\rho_0$, the 3P_2 partial wave of the nucleon-nucleon interaction becomes attractive enough leading to 3P_2 neutron pairing. In contrast, the 1S_0 partial wave would become repulsive and the neutrons would cease to pair in the 1S_0 state. Instead, the 1S_0 proton pairs are predicted to appear owing to its fraction smaller than neutrons. At much higher baryon density $\rho > 2\rho_0$ various hyperons may emerge; some kinds of them possibly form pairs in the same way as nucleons do.

Within a conventional picture, the proton superconductor is believed to be type II supporting a lattice of magnetic flux tubes that carry the magnetic flux through the neutron star. For some values of the proton correlation length and the London penetration depth, the distant proton vortices repel each other leading to formation of a stable vortex lattice. In addition, the rotation of the neutron star causes a lattice of quantized vortices to form in the superfluid neutron state, similar to the observed vortices that form when superfluid Helium is rotated. However, according to recent work [27], a detailed analysis of both the proton and neutron Cooper pair condensates, indicates that the superconductor may in fact be type I with critical temperature $H_c \sim 10^{14} \text{G}$ and would exhibit the Meissner effect. This is in agreement with astrophysically based arguments indicating that the conventional picture of a neutron star as a type-II superconductor may have to be reconsidered [28].

If the central densities of compact stars exceed the density of deconfinement phase transition to a quark matter phase, the deconfined quark matter can settle in a color superconducting phase if there exist any attractive channel of the quark-quark interaction. This subject was already addressed in the early 1980s [31] but came back a few years ago since the realization that the typical superconducting gaps in quark matter may be much larger than those predicted in the early works (Δ as high as $\sim 100 \text{MeV}$) [32]. There exist many phases of color superconducting quark matter. Whereas at the lower densities (near ρ_0) the phase structure is thought to be very rich, it seems clear that at very large densities the lowest energy state is the so called Color-Flavor-Locked (CFL) phase [39] where quarks of all flavors and colors pair. The effect of the pairing gap on the equation of state and on the structure of compact stars may be extremely important at densities near the zero pressure point (i.e. near the surface of the star). Estimations show that there is much room for paired quark matter being more stable than nuclear

matter and atomic nuclei [29], implying that there may exist “neutron stars” formed by CFL strange matter from the center up to the surface [30].

According to early calculations [31], the quark superconductor was thought to be marginally type I with a zero-temperature critical field $H_c \sim 10^{16}$ G. This picture has been recently reconsidered [35, 33, 34, 36, 37]. Quark matter would make the gluons massive (there is a color Meissner effect) but allows a “modified” photon (a combination of electromagnetic and some gluon fields), which remains massless in the color superconducting phase. Therefore, a color superconductor would be penetrated by an external magnetic field without restricting them to quantized flux tubes.

EFFECTS OF B ON HYPOTHETICAL COMBUSTION PROCESSES

Flames, instabilities, and asymmetries

Combustion processes are common in astrophysics. A classic example is type Ia supernovae in which a thermonuclear explosion completely disrupts a white dwarf. Combustion processes maybe also relevant for neutron stars. In some situations (see e.g. [40, 44]), part or the whole neutron star can convert into quark matter. This conversion is believed to proceed as a combustion process [40, 42, 43]. The combustion starts as a laminar deflagration (slow combustion) propagating outward, and then enters a regime of turbulent deflagration due to the action of several hydrodynamic instabilities, such as the Landau-Darrieus (LD) and the Rayleigh-Taylor (RT) instability, in which the flame wrinkles and accelerates. For the largest scales (those much larger than the thickness of the flame), the RT instability is expected to dominate over LD in the astrophysical case. These effects can be understood within a fractal model in which the flame behaves like a fractal with an area $\bar{A} \propto \bar{R}^D$ (with $2 \leq D < 3$), where D is the fractal dimension of the surface, and \bar{R} is the mean radius of the wrinkled surface [45]. Numerical simulations [45] and laboratory experiments involving different gas mixtures [46] show that the fractal growth actually increases the velocity of the combustion front because of the change in the transport mechanism from a laminar to a fully turbulent burning.

The effect of the magnetic field on flames can also be understood within a fractal model [47]. For the sake of simplicity, we shall assume hereafter a dipolar magnetic field geometry and thus the flame propagating in two particularly representative directions, one parallel to the B-field lines (in the polar direction) and the other one perpendicular to them (in the equatorial direction). Within the fractal description, the effective velocity of the flame is given by $v_f = v_{lam}(L/l)^{D-2}$ where L and l are the maximum and minimum length scales of perturbations unstable to the Raleigh-Taylor instability. The characteristic velocity of the Raleigh-Taylor growing modes must be $\geq v_{lam}$, i. e. $ln_{RT}(l) = v_{lam}$, where n_{RT} is the inverse of the characteristic RT time. In the absence of a B-field $n_{RT} = (1/2\pi)\sqrt{gk\Delta\rho/2\rho}$ and so $l = 4\pi\rho v_{lam}^2/g\Delta\rho$. However, in the equatorial direction, the presence of B is essential in modifying the dispersion relation for the RT instability which reads $n_{RTB} = 1/2\pi\sqrt{gk(\Delta\rho/2\rho - kB^2/4\pi g\rho)}$ where $k = 2\pi/\lambda$, λ is the wavelength of the perturbation, and $\Delta\rho = \rho_u - \rho_d$ is the density difference

between the upstream and downstream parts of the flame front [47]. Taking into account the velocity of the flame, the minimum cut-off length is given by the condition $l_e n_{RTB}(l_e) = v_{lam}$, from which we derive the minimum scale in the equatorial direction $l_e = 8\pi(B^2/8\pi + \rho v_{lam}^2/2)/g\Delta\rho$. So, from the above definition of the fractal velocity we have $v_e = v_{lam}(L/l_e)^{D-2}$ and $v_p = v_{lam}(L/l_p)^{D-2}$. Note that L is not modified by B and has the same value in both directions. Then, the ratio between the equatorial and polar velocities is [47]:

$$\xi = \frac{v_p}{v_e} = \left[1 + \frac{B^2}{4\pi\rho v_{lam}^2} \right]^{D-2}. \quad (6)$$

That is, although the magnetic pressure $B^2/(8\pi)$ is not relevant compared to the degeneracy pressure of neutron star matter, the B -field is important for the combustion because it quenches the growth of RT instabilities in the equatorial direction acting as a “surface tension”, while it is innocuous in the polar one where in average we have $\vec{v}_p \times \vec{B} = 0$. More specifically B modifies the minimum RT instability scale and since the turbulent flame velocity is related to RT-growth this results in a different velocity of propagation along each direction. Notice that although this is based solely on a linear analysis, 3D numerical simulations of the development of the magnetic RT instability [49] reinforce the results of asymmetry here reported and also reveals a tendency for its amplification in the non-linear regime. This has interesting effects for thermonuclear supernovae [47, 48] and neutron stars [43, 57].

Gamma-ray bursts from neutron star phase transitions

It has long been recognized that the transition to quark matter inside neutron stars provides a suitable inner engine for GRBs since some key requirements of observations are fulfilled (e.g. energy released [51, 55], timescale of the gamma emission [56], small baryon loading [54], rate of events [53]). The total energy released is $\sim 10^{53}$ ergs [51], and can be even larger if the final state is paired quark matter, due to the liberation of the gap energy. The timescale of gamma emission in this model has been calculated to be ~ 0.2 s [52, 56, 55], i.e. compatible with short gamma ray bursts and giant flares in soft gamma repeaters [57]. Another interesting feature is that the surface of a quark star can guarantee an environment with a very small baryon load, automatically generating conditions needed for an ultra-relativistic fireball [54].

The magnetic field is essential for producing a collimated gamma emission in this class of models due to the effect in the hydrodynamics of the conversion process [56]. The basics of the proposed mechanism are as follows [56]. A seed of quark matter may become active or form following the standard supernova bounce. The interface must then propagate outwards powered by the energy release of converted neutrons, much in the same way as a laboratory combustion. It seems reasonable to assume the combustion to begin as a laminar deflagration, which quickly reaches a regime of turbulent deflagration. The magnetic field generates a strong acceleration of the flame in the polar direction. This results in a (transitory) strong asymmetry in the geometry of the just formed core of hot quark matter. While it lasts, this geometrical asymmetry gives rise to a bipolar

emission of the thermal neutrino-antineutrino pairs produced in the process of quark matter formation. This is because almost all the thermal neutrinos generated in the process of quark matter formation will be emitted in a free streaming regime through the polar cap surface, and not in other directions due to the large opacity of the matter surrounding the cylinder. Further annihilation into electron-positron pairs just above the polar caps, gives rise to a relativistic fireball. Since the timescale of gamma emission is ~ 0.2 s, this provides a suitable explanation for the inner engine of short gamma ray bursts.

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