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Submergence and re-diffusion of the neutron star magnetic field after the supernova

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Abstract. We consider the effect of post core-collapse accretion on the magnetic field of the new-born neutron star. If this accretion is hypercritical then the ram pressure overwhelms the magnetic field pressure and we show that in this case the accretion induced flow time scale in the upper layers of the neutron star is shorter, by orders of magnitude, than the magnetic field ohmic diffusion time scale. This means that the magnetic field is frozen in the matter and any initial magnetic field of the neutron star is rapidly submerged beneath the accreted matter. If the accreting matter is weakly, or non, magnetized, this implies that neutron stars produced by supernovae in which hypercritical accretion occurred are born with weak, or even vanishing, surface magnetic field. Later diffusion of the magnetic field back to the surface could produce a delayed switch-on of a pulsar (Muslimov & Page 1995): We model this re-diffusion in detail for a wide range of submergence depths and discuss the consequences for pulsar observables, such as the period P and its time derivative $\dot{P}$. As a result of the field submergence, the spin-down age τ_{sd} can be much larger than the real age of the pulsar. Moreover, we show that if the field submergence is deep enough the magnetic field will be hidden for many millions of years. This mechanism of field submergence may explain the lack of evidence for the presence of a pulsar in all recent supernovae and may also contribute to the discrepancy between the estimated pulsar birth rate and type Ib+II supernova rate. In particular, we predict that a pulsar will never turn-on in the remnant of SN 1987A.

Key words: magnetic fields – stars: neutron – stars: pulsars: general – stars: supernovae: general – stars: supernovae: individual: SN 1987A

1. Introduction

A paradigm of pulsar (PSR) and neutron star (NS) astronomy is that NSs are born in supernova (SN) explosions, which is sometimes extended into saying that SNe form PSRs. However, the association of PSRs and SNe is a challenging issue. The recent SN 1987A provided a brilliant confirmation that SNe do form NSs by the detection of the neutrinos emitted by the proto-NS. However, direct evidence for the presence of a NS in the supernova remnant (SNR) of SN 1987A is still lacking. The same problem arises with several other recent SNe which show no clear evidence for the presence of a PSR (Chevalier & Franson 1992). More generally, there is a growing discrepancy between the estimated SN rate and the PSR birth rate in the Galaxy (Frail 1998). Complementarily, only few SNR (about 10 to 20) have a PSR associated (Kaspi 1996) a fact which may be another facet of the same *missing pulsars problem*. (For a discrepant opinion see however Gaensler & Johnston 1995.)

However, SNe may also form black holes and, as has been demonstrated recently, do sometimes produce magnetars (Kouveliotou et al. 1998). Spruit & Phinney (1998) proposed a solution to this missing PSR problem by arguing that if the PSR magnetic field is a fossil of the progenitor field many NSs are born with such low rotational periods (but strong magnetic fields) that they are not able to emit radio pulses. This proposition may be seen as a wide extension of the 'injection' scenario of Vivekanand & Narayan (1981), which however seems to be not necessary in the light of recent PSR population studies (Lorimer et al. 1993). Parallely, Allakhverdiev et al. (1997) proposed that there is a large population of low luminosity radio PSRs mostly undetected and probably undetectable.

We propose here that a substantial fraction of NSs are born with low or even vanishing surface magnetic field. We show that if the magnetic field is generated by a dynamo action in the proto-NS (Thompson & Duncan 1993) then the hypercritical accretion expected to occur later on (Chevalier 1989) will submerge it unless the simplified traditional picture of fossil flux amplification can be maintained. Once the accretion has stopped, the submerged field will slowly diffuse back toward the surface, producing a delayed switch-on of a PSR (Muslimov & Page 1995), on a time scale which may range form hundreds to many millions of years.

The paper is organized as follows. In Sect. 2 we describe and discuss the hypercritical accretion expected to take place in some SNe explosions and in Sect. 3 consider the process of magnetic field submergence. Sect. 4 discusses the NS magnetic field existent at the end of the hypercritical accretion phase. The subsequent re-magnetization of the NS due to the back diffusion

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of the submerged field is described as well as its consequences for the delayed switch-on of a PSR in Sect. 5. Finally, Sect. 6 is devoted to a discussion of our results where possible observational consequences of our scenario are presented.

2. Hypercritical accretion in supernovae

Since the early works on the SN explosion mechanism it has been realized that enormous accretion onto the new born NS takes place (Colgate 1971; Zel'dovich et al. 1972) and after SN 1987A interest in this hypercritical accretion has been renewed (Chevalier 1989; Houck & Chevalier 1991; Brown & Weingartner 1994; Fryer et al. 1996; Colpi et al. 1996). Independently of the uncertainties on the exact explosion mechanism, on which the existence of the early accretion crucially depends, a reversed shock wave can be generated when the first expanding shock wave reaches the outer boundary of the helium layer of the progenitor. Chevalier (1989) showed that this reversed shock wave produces a late accretion phase at a hypercritical rate which may last for months. This fall back is initially Bondi-like followed by a transition to a dust-like accretion phase (Chevalier 1989; Colpi et al. 1996). Assuming a spherically symmetric flow, in the particular case of SN 1987A (Chevalier 1989) the Bondi phase starts about $t_0 = 7 \times 10^3$ seconds after the initial explosion, roughly the time at which the falling back matter hits the NS surface, with a rate of the order of

$$\dot{M}_B \approx 2 \times 10^{28} \ (t/t_0)^{-3/2} \ \ \mathrm{g/s} \tag{1}$$

and the transition to dust-like regime occurs around $t_1 \approx 1.7 \times 10^4$ seconds with a rate

$$\dot{M}_D \approx 6 \times 10^{27} \ (t/t_1)^{-5/3} \ \ \mathrm{g/s}. \tag{2}$$

Chevalier (1989) argued that, when $\dot{M}$ is above $\dot{M}_{cr} \sim 2 \times 10^{22}$ g/s $\sim 3 \times 10^{-4}$ $M_\odot$ /yr, the ram pressure at the NS surface is larger than the pressure of a 10^{12} G magnetic field, a condition largely satisfied at the above rates, which means that the accretion flow is purely hydrodynamical. This $\dot{M}_{cr}$ is also the accretion rate at which hypercritical accretion is stopped by radiation pressure and becomes unstable (Houck & Chevalier 1991; Fryer et al. 1996), at which time about 0.1 $M_\odot$ has been accreted in the case of SN 1987A.

Many factors will affect these simple estimates of $\dot{M}$. Entropy generation due to the decay of ^{56}Ni and ^{56}Co reduces the accretion rates as time goes but does not affect the initial rate (Chevalier 1989). Rotation is potentially significant if the stagnation radius $r_{st} = l^2/GM$, where l is the angular momentum per unit mass of the accreting matter, is outside the NS (Chevalier 1996). In that event the centrifugal barrier can strongly affect accretion. This is however unlikely to happen since $r_{st} > R$ requires that l be larger than 10^{16} cm^2 s^{-1}: according to Langer et al's (1997) models of rotating massive stars the pre-SN core's specific angular momentum is below this value. For comparison, the specific angular momentum of a 1.4 $M_\odot$ NS near break-up is also below 10^{16} cm^2 s^{-1}. However, if the proto-NS has a rotational period of the order of a millisecond, i.e., $l \sim 10^{16}$ cm^2 s^{-1} then an efficient dynamo is expected to occur within it and generate a magnetic field of the order of 10^{15-16} G (Duncan & Thompson 1992): combined with the angular momentum this is likely to impede accretion onto this 'magnetar' unless the accretion rate is extremely high (as estimated for SN 1987A for example).

A large kick velocity v_{kick} of the new-born NS may also be worrisome for the accretion. The Bondi accretion rate, proportional to $c_{s\,\infty}^{-3}$ for an accretor at rest ($c_{s\infty}$ being the speed of sound at infinity), becomes then proportional to $(c_{s\,\infty}^3 + v_{kick}^3)^{-1}$. From typical values for SNe explosions (Bethe 1993) we estimate that $c_{s\,\infty} \sim 400\,\mathrm{km\ s}^{-1}\ (M_{CO}\ v_9\ t_h)^{-1/2}$, for an adiabatic homologous expansion of a radiation dominated CO core with mass M_{CO} (in $M_\odot$); v is the expansion velocity after the passage of the reverse shock, and t_h the time in hours. Here and thereafter we note as X_n for X measured in units of 10^n [cgs]. For SN 1987A, $t_h \sim 2$ when the reverse shock wave reaches the NS surface, but for more typical SNe t_h is easily 10 times larger (Chevalier 1989) and moreover $v_9 \sim 1$ for a weak reverse shock instead of 0.06 after the strong reverse shock of SN 1987A (Woosley 1988). With a kick velocity of the order of 500 km/s (Lorimer et al., 1997) $\dot{M}_B$ is thus seriously reduced.

Finally, the new-born NS may have an enormous rotational energy $E_{rot} \sim 4 \times 10^{52}\ P^{-2}$ ergs, where P is the period in milliseconds. If a substantial torque can act on it and force it to inject its rotational energy loss into the surroundings this may be more than enough to impede accretion. The nature of the torque is not clear since naive in vacuum magneto-dipolar radiation, as usually considered for isolated PSRs, is inappropriate in the dirty environment of a new-born NS. Moreover, a fast spinning hot NS is expected to lose most of its rotational energy by gravitational waves (Anderson et al., 1998).

In summary, unless the late fall-back is enormous there are many possible factors which can reduce it and even suppress it totally. We will thus *assume* that hypercritical accretion is taking place and consider its effect. In the present paper we will treat this hypercritical accretion as being spherically symmetric. Which fraction of new-born NSs actually undergo such an accretion phase is a very delicate issue that we can only leave for future work.

3. Submergence of the initial magnetic field

The evolution of the magnetic field is governed by the MHD induction equation

$$\frac{\partial \boldsymbol{B}}{\partial t} = -\frac{c^2}{4\pi}\nabla \times \left(\frac{1}{\sigma}\nabla \times \boldsymbol{B}\right) + \nabla \times (\boldsymbol{v} \times \boldsymbol{B})\ , \tag{3}$$

where σ is the electrical conductivity and $\boldsymbol{v}$ is the flow velocity of the matter. To get a qualitative impression of the effect of hypercritical accretion let us first consider the time scales resulting from the two terms in the r.h.s. of Eq. 3 for ohmic diffusion

$$\tau_{ohm} = \frac{4\pi\sigma L^2}{c^2}, \tag{4}$$

and matter flow

$$\tau_{flow} = \frac{L}{v_r}, \tag{5}$$

where $v_r = \dot{M}/4\pi r^2 \rho$ is the radial velocity of the accretion flow, L is a typical length scale and c the speed of light, ρ being the matter density and r the radial coordinate. As we show below, given the physical conditions during hypercritical accretion τ_{ohm} is huge in comparison with all other time scales involved in the problem, and in particular with τ_{flow}.

In general terms the electrical conductivity due to electron - ion collisions is given by

$$\sigma = \frac{e^2 n_e}{m^* \nu_{ei}}, \tag{6}$$

where n_e the electron number density, $m^* = (m_e^2 c^2 + p^2)^{1/2}/c$ their effective mass (m_e being the electron rest mass and p their momenta) and

$$\nu_{ei} = n_i v \sigma_t \tag{7}$$

is the electron-ion collision frequency, with n_i the ion number density and v the relative electron-ion velocity. For electron - ion collisions the transport cross section is

$$\sigma_t = \frac{4\pi (Ze^2)^2}{m^{*2} v^4} \times \Lambda, \tag{8}$$

where Z is the ions charge (Landau & Lifschitz 1990). The Coulomb logarithm Λ will be taken as 1 for our estimates.

In the accreting envelope the temperature is of the order of 10^{10} K (Chevalier 1989) and, being a wholly ionized plasma we can take $v \sim c$, $m^* \sim k_B T/c^2$ giving

$$\sigma \sim 1.5 \times 10^{22} \, \frac{1}{Z} \, T_{10} \;\; [\mathrm{s}^{-1}]. \tag{9}$$

Notice that in this regime σ is independent of the density, and the result is only valid for $T \gg 10^9$ K. With this we obtain

$$\frac{\tau_{ohm}}{\tau_{flow}} \sim 10^{14} \, \frac{T_{10} \, L_5 \, (\dot{M}/\dot{M}_{cr})}{r_6^2 \, Z \, (\rho/\mathrm{g\ cm}^{-3})} \tag{10}$$

Thus τ_{ohm} is much larger than τ_{flow} for any reasonable values of L, r and ρ when $\dot{M}$ is above $\dot{M}_{cr}$.

In the NS crust, the conductivity becomes density dependent because of electron degeneracy. At high enough density the electrons are relativistic and then $v = c$ and $m^* = p_F/c$, p_F being the Fermi momentum related to n_e by $n_e = p_F^3/3\pi^2\hbar^3$. In the phase where the ions form a liquid, we then have

$$\sigma \sim 8.6 \times 10^{21} \left(\frac{\rho_6}{AZ^2}\right)^{1/3} \;\; [\mathrm{s}^{-1}] \, . \tag{11}$$

At $T \sim 10^{10}$ K practically the whole crust is liquid. With this we obtain

$$\frac{\tau_{ohm}}{\tau_{flow}} \sim 2 \times 10^{10} \, \frac{L_5 \, (\dot{M}/\dot{M}_{cr})}{r_6^2 \, \rho_6^{(2/3)} (AZ^2)^{1/3}}. \tag{12}$$

Thus τ_{ohm} is much larger than τ_{flow} in the crust too.

In the low density region of the crust, electrons are partially degenerate and there is a smooth transition between the results of Eqs. 9 and 11. Fig. 1 explicitly shows τ_{ohm} and τ_{flow} in the NS crust for different accretion rates taking into account the exact $T - \rho$ dependence of σ.

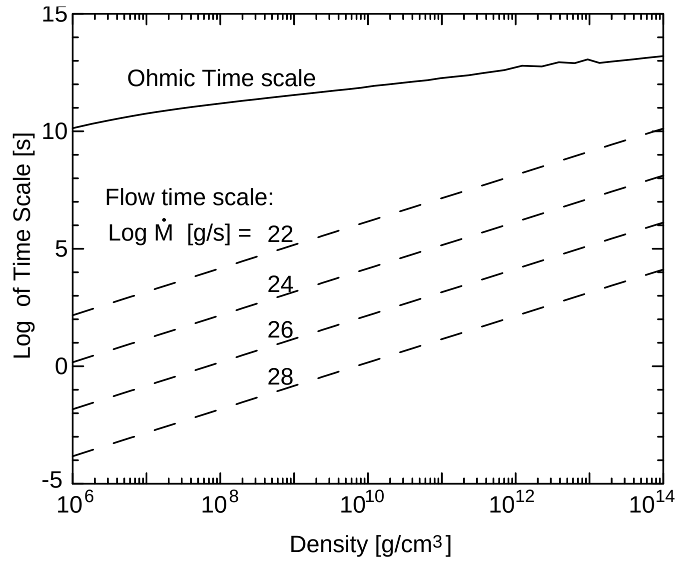

Fig. 1. Comparison of the ohmic and flow time scales in the NS envelope and crust. The scale height L is taken as 1 km, a typical value for the inner crust pressure scale height. We assume here $T = 10^{10}$ K, but since σ is T-independent while the crust is liquid and then increases in case of crystallization of the deeper layers, the τ_{ohm} plotted here is a conservative estimate.

Thus, any magnetic field present in the accreting matter and inside the NS at the beginning of fall back is frozen in the inward drifting matter, i.e., we are dealing with an ideal MHD problem. This results in the piling up of the accreted matter on top of the originally magnetized NS matter much faster than the field can diffuse outward. Fig. 2 plots the density of submergence as a function of accreted mass and shows that a small amount of accretion can push the original magnetic field to very high densities. In the case of SN 1987A the initial magnetic field has been pushed into the core of the NS.

To illustrates this submergence we consider an axially symmetric field and solve, in ideal MHD, the induction equation. Fig. 3 shows the resulting evolution of the magnetic field using the accretion parameters estimated for SN 1987A. With our particular initial configuration one sees that after less than two hours the field is already submerged into the core. Notice also the compression of the field in the deep high density layers where v_r becomes very small and B can be amplified by many orders of magnitude. However, the very narrow peaks have such small scale lengths, $\sim$ a few centimeters, that τ_{ohm} becomes comparable or even smaller than τ_{flow} and these peaks are reduced during the accretion phase already. Since this calculation is performed in ideal MHD we cannot calculate this effect but we estimate the length scale L_{ohm} such that $\tau_{ohm}(L_{ohm}) = t$ (= time elapsed since the beginning of accretion) and mark as a dashed line the portion of the peaks with widths comparable to L_{ohm}. At the end of the accretion phase the compressed field, amplified by one to two orders of magnitude, is localized in a layer of thickness of about 100 to 200 meters covered by a layer

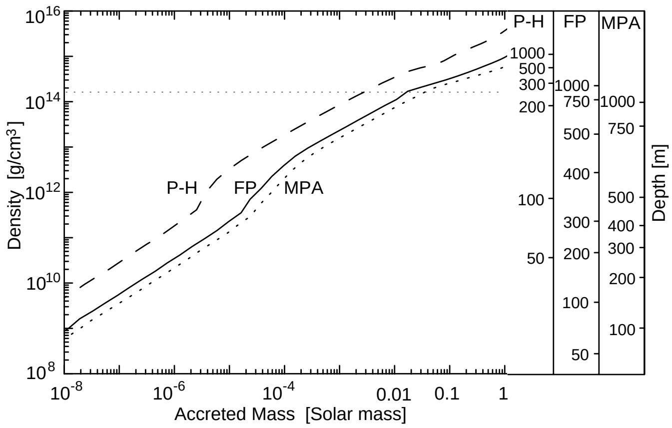

Fig. 2. Density and depth reached by the accreted matter as a function of the accreted mass. The three curves correspond to the standard equation of state FP (Friedman & Pandharipande 1981), a stiff equation of state MPA (Müther et al. 1987), both of which include only nucleons, and a very soft one P-H (Pandharipande 1971) which also includes hyperons. The thin horizontal line shows the crust - core boundary. This figure shows that for a given accreted mass, the density and depth of submergence and hence the delay in the PSR switch-on are sensitive functions of the dense matter equation of state.

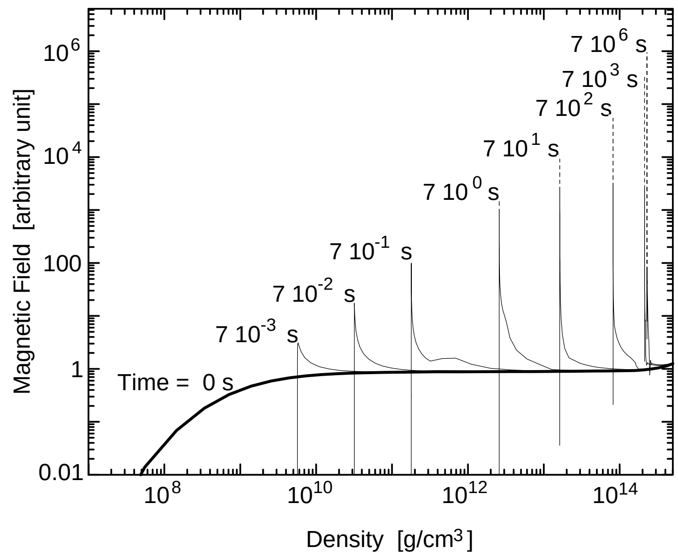

Fig. 3. Distribution of the angle averaged magnetic field strength in the NS as a function of time for the accretion rate estimated for SN 1987A. Initial time corresponds to the beginning of the accretion phase. In less than two hours the initial field is submerged into the core. Notice that the maximum value attained by the field depends on its initial value at low density since the low density region is the most strongly compressed. The calculation assumes ideal MHD, but these zones of highly compressed field have a very small length scale and thus a very small ohmic diffusion time τ_{ohm}: these fields, shown as dashed lines, will eventually be washed out by diffusion when the time becomes comparable to τ_{ohm}.

of a few hundreds of meters of accreted matter. The ohmic decay time of this compressed field is short because of its short length scale and this can be a significant source of heating for the young NS. A 'hidden magnetar' has been produced: this result appears robust since it comes from the compression of the matter and does not depend upon the details of the initial configuration.

4. The post fall-back magnetic field

The magnetic field present at the surface of the NS when accretion stops will thus be the field which was present in the accreted matter and compressed.

In the standard hypothesis that the PSR magnetic field is a fossil of the progenitor's core field then the accreted matter, coming also from matter previously in the progenitor's core, could bring in a field comparable to the field already present in the NS, i.e., the NS may be born with a strong surface field. However, the hypothesis that the fall-back matter brings in a well-ordered large scale field is questionable since there is still the possibility that this accreting matter has suffered a turbulent episode during which the plasma behaved as a diamagnet (Vaĭnshteĭn & Zel'dovich 1972) and its field could have been severely reduced, which would means that the final surface field of the NS could be weak.

In contradistinction, within the proto-NS dynamo scenario for the origin of PSR magnetic field (Thompson & Duncan 1993) the core of the progenitor is only required to have a small field which will act as a seed for the dynamo action. In this case the field present at the NS surface after accretion will be small (unless the initial spin period of the proto-NS was short enough for the dynamo to produce a magnetar and impeach accretion as mentioned in Sect. 2). The strength of the surface magnetic field of a new-born NS which has undergone hypercritical accretion may thus be very different if its magnetic field is of fossil or proto-NS dynamo origin.

We are going to assume that the accreting matter is non, or weakly, magnetized, and shall explore its consequences. This is natural within the proto-NS dynamo scenario and may also be compatible with the fossil field hypothesis.

5. Neutron star 're-magnetization'

When the accretion stops, the field can start diffusing back toward the surface. This problem has been considered by Muslimov & Page (1995, 1996), in the case of very shallow submergence, $M_{acc} \sim 10^{-5}\ M_\odot$, who showed that after a few hundreds years the surface field strength becomes comparable to the interior one, resulting in a delayed switch-on of the PSR.

As argued above, in many cases the hypercritical accretion is likely to submerge the NS magnetic field to much higher densities, resulting in a much longer re-diffusion time. We hence consider here this re-diffusion for a wider range of submergence depths, using the same method as in Muslimov & Page (1996). We assume the magnetic field to be purely axisymmetric poloidal dipolar and using the Stoke stream function formalism write it as

$$\mathbf{B} = B_0 \left(\frac{S}{r^2} \cos\theta \; \mathbf{e}_r - \frac{1}{2r} \frac{\partial S}{\partial r} \sin\theta \; \mathbf{e}_\theta \right) \quad , \tag{13}$$

where B_0 is some normalization magnitude of the field, $\mathbf{e}_r$ and $\mathbf{e}_\theta$ are unit vectors in the radial and meridional directions, and $S = S(r,t)$ is the stream function. The induction equation, Eq. 3 is then

$$\frac{\partial S}{\partial t} = \frac{c^2}{4\pi\sigma}\left(\frac{\partial^2 S}{\partial r^2} - \frac{2S}{r^2}\right) \quad , \tag{14}$$

where the convective term has been dropped. At the surface ($r = R$) we impose the standard boundary condition that the internal field merges continuously with an external vacuum field. For a dipole field this condition reads

$$R\frac{\partial S}{\partial r} + S = 0 \quad . \tag{15}$$

The inner boundary condition is $S(r)/r^2 \rightarrow 0$ at the center of the star.

The re-diffusion process is determined by the depth of submergence and controlled by the conductive properties of the NS matter. For the electrical conductivity we use Eqs. 6, 7 and 8 and the calculation of the Coulomb logarithm Λ of Yakovlev & Urpin (1980) when matter is in the liquid state. When matter crystallizes we us the results of Itoh et al. (1993) for the electron - phonon contribution and of Yakovlev & Urpin (1980) for the electron - impurity one.

Since the electron - phonon conductivity is temperature dependent the re-diffusion will also be affected by the thermal history of the NS. We restrict ourselves to the case of 'standard cooling' using the cooling curve for the 'FP' model of Van Riper (1991). After about a million years the temperature is low enough for σ to be dominated by impurity scattering only and thus become temperature independent. We fixed the impurity concentration to $Q = 0.01$ and used the chemical composition of the crust for accreted matter as calculated by Haensel & Zdunik (1990).

We consider three different submergence depths, $\rho_{sub} = 10^{12}$, 10^{13}, and 10^{14} g cm^{-3} and locate the currents at the end of the accretion phase around these densities. Fig. 4 shows these initial configurations and their spatial structures at different times. In contradistinction with Fig. 3, we here show only the radial component $\mathbf{B_r}$. We stop the re-diffusion calculation when the surface magnetic field reaches its maximum value: the subsequent evolution would involve field decay and is a different issue not considered here. The initial distribution of the field at $\rho < \rho_{sub}$ reflects the effect of the previous accretion; at $\rho > \rho_{sub}$ it is arbitrary but has no effect on the re-diffusion process and would only influence the subsequent decay phase. Both the decay and the diffusion processes inherent to the induction equation are clearly seen. Due to the much higher conductivity in the deeper layers of the crust the re-diffusion is much slower for deeper submergence.

In case the field has been submerged into the NS core, it will be frozen later on into the proton superconductor. The latter forms when T drops below a few times 10^9 K, i.e., after accretion has stopped. When the PSR spins down it is possible that the flux will be expelled from the core (Srinivasan et al. 1990). However, due to the very low external field the spin down is very slow and flux expulsion too. Hence, the re-diffusion time may well be comparable to the Hubble time, but a detailed description of this case is beyond the scope of this paper.

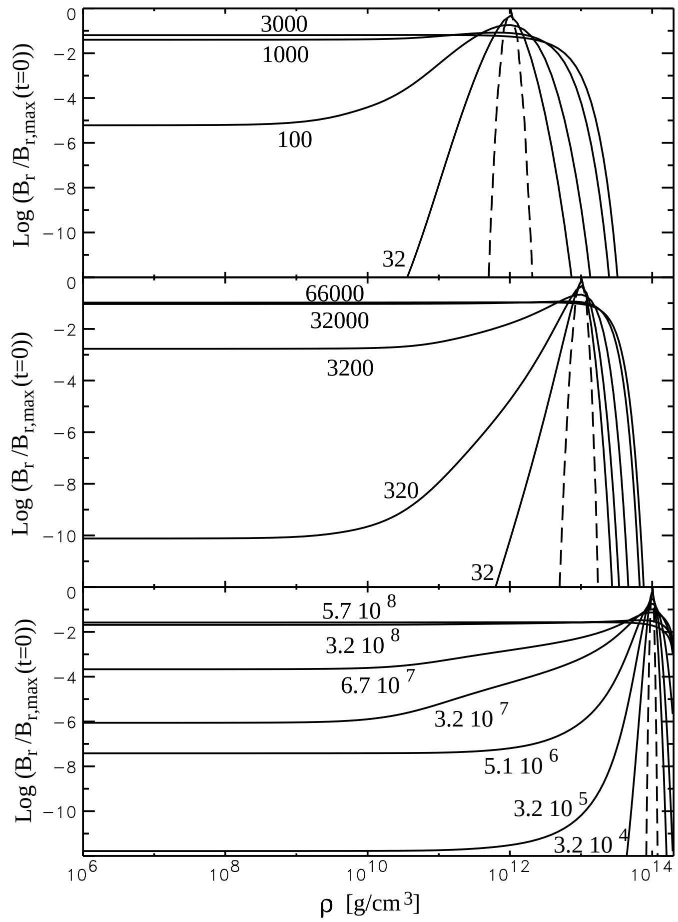

Fig. 4. Time evolution of the interior magnetic field after the hypercritical accretion phase for three different submergence depths. Only the radial component $\mathbf{B_r}$ is plotted in contradistinction with Fig. 3. The assumed initial field location is shown by the dashed lines. The ages of the star are indicated on the lines.

Having the field evolution we can model the rotational history. Assuming magnetic dipole braking the period evolution is given by

$$P^2(t) = P_{in}^2 + \beta \int_0^t B_s^2 \, dt' \ , \tag{16}$$

where $\beta = 16\pi^2 R^6/3c^3 I$, I being the star's moment of inertia and P_{in} the initial spin period. It is now trivial to obtain the active age

$$\tau_{active} \equiv \frac{P}{2\dot{P}} \tag{17}$$

and the spin-down luminosity $L_{spin-down} = 4\pi I\dot{P}/P^3$. Figs. 5, 6 and 7 show the results for the three ρ_{sub} considered. The most important feature seen is that τ_{active} during the re-diffusion phase is *decreasing* with time and is orders of magnitude larger than the real age t. This is simply due to the fact the Eq. 17 is valid only for constant magnetic field and when

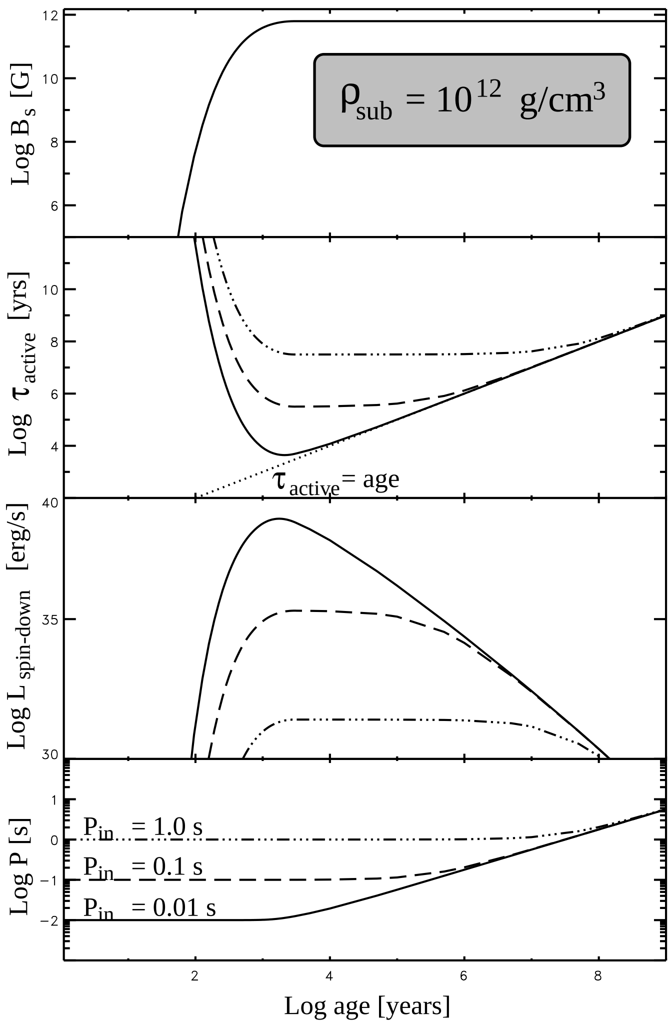

Fig. 5. Evolution the surface field B_s, and active time τ_{active}, spin-down luminosity $L_{spin-down}$ and rotational period P for three initial periods P_{in}. The initial field submergence density is $\rho_{sub} = 10^{12}\,\mathrm{g\,cm^{-3}}$ (see Fig. 4).

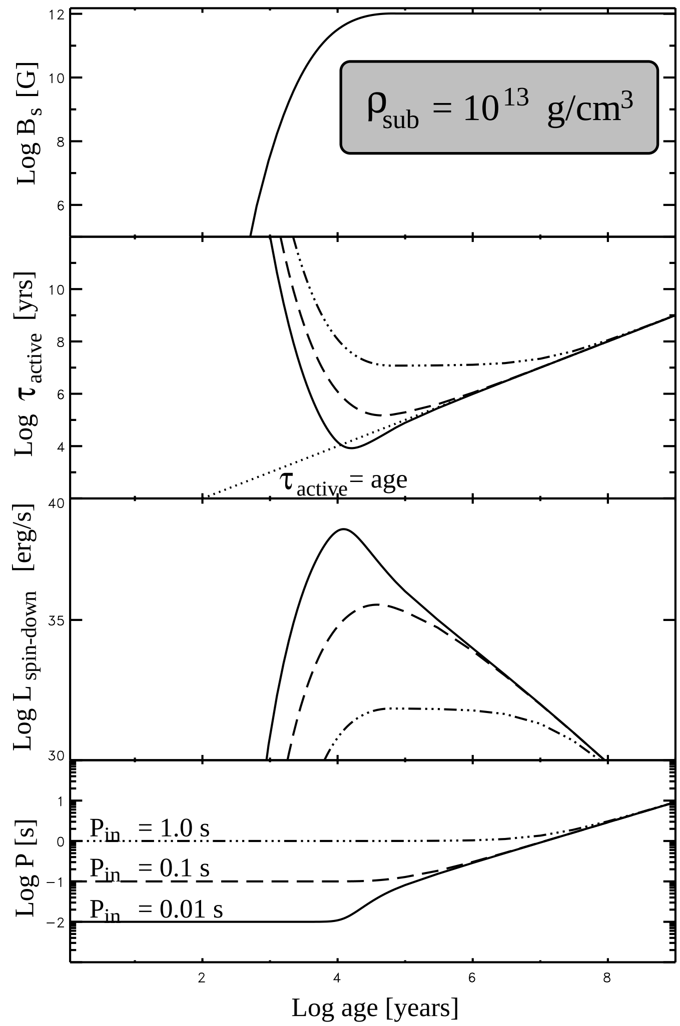

Fig. 6. Same as Fig. 5 but with $\rho_{sub} = 10^{13}\,\mathrm{g\,cm^{-3}}$.

$P_{in} \ll P$. $L_{spin-down}$ is initially very small but reach observationally significant values at an age when τ_{active} is still a wild overestimate of t. Notice that for $\rho_{sub} = 10^{14}\,\mathrm{g\,cm^{-3}}$ most of the re-diffusion occurs when the star is cool enough that the conductivity is dominated by electron - impurity collisions, which produces a two phase re-diffusion.

6. Discussion

As we have shown here, one should seriously consider the possibility that SNe make 'de-magnetized' NSs which may, or may not, turn on as PSRs later. The submergence of the field by any amount of accreted matter can easily explain the absence of PSR activity in recent SNe (Muslimov & Page 1995) but the very existence of PSRs leads to the conclusion that in some fraction of SNe the fall back accretion is weak enough that it allows a re-diffusion of the field in $\sim 10^3$–10^4 years. Given the relation between submergence depths, total amount of accreted matter (Fig. 2) and the rediffusion times (Figs. 5, 6, 7), this requires that the total accreted mass is much less than 0.01 $M_\odot$. If the SN Ib+II rate and PSR birthrate are comparable then this upper bound on the accreted mass should apply to most SNe, while if the former rate is significantly higher than the latter then in a large fraction of SNe the total accreted mass could be higher. This mechanism may contribute to the solution of the missing PSR problem but it is difficult to estimate how significant it is. The fraction of NSs which will never turn-on as PSRs depends crucially on how many SNe undergo heavy accretion. This unfortunately depends on the poorly understood late evolution of the envelope of massive stars, but the nature of the progenitor of SN 1987A shows that this fraction is not vanishingly small (Woosley 1988).

The large discrepancy between the active age, τ_{active}, and the real age of a PSR during the re-magnetization phase raises the possibility that a relatively young object may be very old - looking. One way to identify such a PSR would be through an association with a SNR. A superficial comparison of the Princeton PSR Catalogue (Taylor et al., 1993) with the Green's SNR catalogue (Green, 1998) reveals more than 50 PSRs appear-

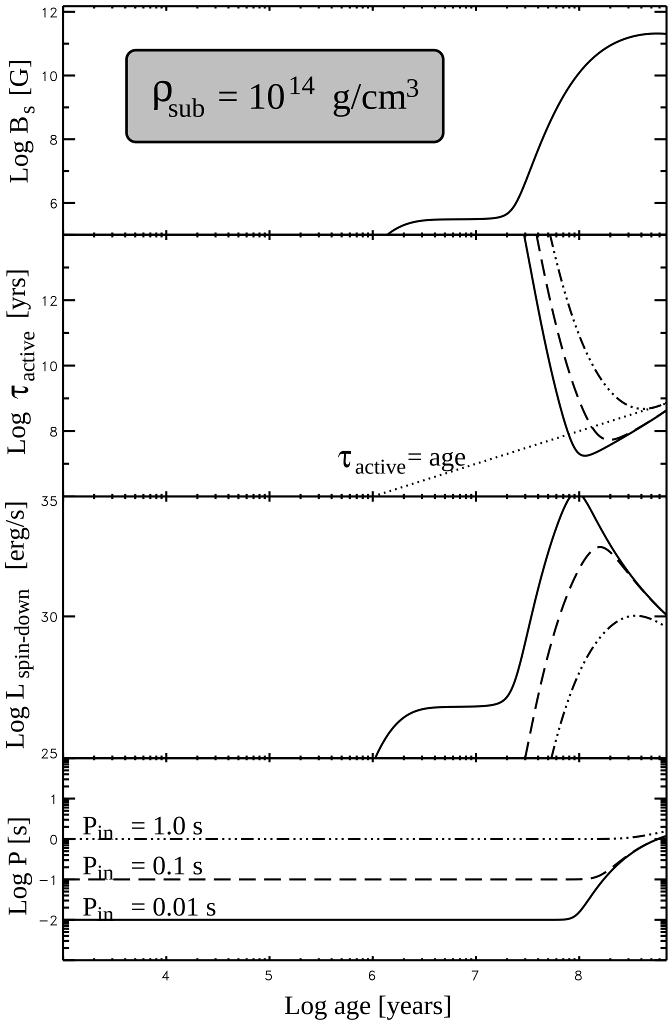

Fig. 7. Same as Fig. 5 but with $\rho_{sub} = 10^{14}\,\mathrm{g\,cm^{-3}}$.

ing within a SNR, more than 30 of them having τ_{active} above 10^5 yrs. Searches for association have so far explicitly excluded old - looking PSRs, but in the light of the results presented above some of these may be associated with the SNR. Unfortunately identifying them is a very difficult task.

As emphasized in Sect. 3, the proposed scenario relies upon the assumption that the accreted matter is weakly magnetized, either because the progenitor's core had a very weak magnetic field or because the explosion and/or accretion process demagnetized it. Moreover, there remains also the possibility that immediately after the fall back a mechanism generates a strong field in the very surface layers. An example of such an effective mechanism is based on a thermoelectric instability (Blandford et al. 1983) driven by the strong temperature gradient existent in the outer crust. As long as T_{s6}^4/g_{s14}, where T_s is the surface temperature and g_s the surface gravity, is larger than ≈ 100 (Wiebicke & Geppert 1996) thermoelectric field amplification acts. This relation may be maintained for several centuries if the NS cools according to the so called 'standard model' of NS cooling (Page 1998) and the delayed switch-on of PSRs may constitute an independent tool for the study of NS interiors.

X-ray observations may be our best chance to identify demagnetized NSs: a strong magnetic field induces a non uniform surface temperature distribution in a cooling NS, which potentially results in modulation of its thermal emission, observable in the X-ray band (Page 1995). An excellent candidate for a non magnetized NS is offered by RX J185635-3754 (Walter et al. 1996) for which surface thermal emission has been clearly detected with no modulation at a 1% level. Another possible candidate could be the young NS observed in the SNR PKS1209-52 (Zavlin et al. 1998) which also shows no signs of modulation, but with much less stringent limits (24%) on the modulation amplitude. PSR activity can also be detected through the presence of a plerion and weak radio emission has been detected around the NS in PKS1209-52 which indicates the presence of a plerion and thus possibly some weak PSR activity: Vasisht et al. (1997) interpret it as due to a strongly magnetized, slowly rotating NS, but it could also come from a fast rotating, weakly magnetized one. Future X-ray observatories like AXAF, Spectrum-X and XMM will detect more cooling NSs. The spectral resolution of these telescopes will open the possibility to estimate the surface field strength directly from spectral fits (Pavlov et al. 1995), particularly when supplemented by optical and ultraviolet observations (Pavlov et al. 1996), and allow us to estimate the fraction of non, or weakly, magnetized ones.

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