ROUTE TO
Fermilab Library
0 1160 0013641 0
FERMILAB
JUL 0 1992
LIBRARY

RAPORT NR 1569/PH

SPIN PROPERTIES OF DENSE MATTER AND THE ORIGIN OF MAGNETIC FIELD OF NEUTRON STAR

MAREK KUTSCHERA
H. Niewodniczański Institute of Nuclear Physics, ul. Radzikowskiego 152
31-342 Kraków, Poland

and

WŁODZIMIERZ WÓJCIK
Institute of Physics, Technical University, ul.Podchorążych 1
30-084 Kraków, Poland

December 1991

ABSTRACT

Possible mechanisms of a spontaneous spin polarization in the dense baryon matter of the neutron star core are discussed. It is shown that proton impurities in a neutron matter are likely to spontaneously polarize at neutron star core densities becoming the source of the magnetic field. The quark core, if present, can be also magnetized due to ordering of quark spins in the pion-condensed phase of a chirally broken quark matter. This model can account for a recent evidence that decay of magnetic field occurs only for neutron stars which accreted matter in their evolution.

1. Introduction

The problem considered here is a question if the spin ordering can occur in the ground state of dense baryon matter which exists in the neutron star interior. One expects the phase transition from the nucleon phase to the quark phase to occur at such densities. Hence the baryon matter in neutron stars can exist in the form of nucleon matter, below the deconfinement density and in the form of quark matter in the deconfined phase. In both cases these are fermion systems in which spin interactions could lead to spin ordering. Our goal here is to identify possible mechanisms which could produce a spin polarization in the nucleon matter and in the quark matter of the neutron star core. The two proposed mechanisms are different but both of them result in a permanently magnetized phases of, respectively, nucleon matter and quark matter in the neutron star core.

[1]) Lecture given by M. Kutschera at the XV INTERNATIONAL SCHOOL OF THEORETICAL PHYSICS, 16-23 September 1991, Szczyrk, Poland

INP-1569 PH

ION The magnetized matter present in the neutron star core can contribute to the magne moment of the star. In such a case magnetic field serves as a long-distance probe wh can transmit information from the neutron star core. This may be a unique possibility observe a direct signal of existence of the most dense baryon states achievable in nat which can not be produced in laboratory collisions of heavy nuclei. Studies on magne fields of neutron stars could thus be of great importance for our understanding of t ground state of dense baryon matter. There is also a purely astrophysical aspect of t problem: The strong magnetic fields of neutron stars, which range from $10^8 G$ to $10^{13} G$, a not well understood, i.e. their origin is not known. Presence of a magnetized core insi neutron stars, which is predicted by the models discussed here, seems to be an interesti possibility which, however, deserves further studies. For example, an important point be clarified is the influence of the outer layers of the neutron star on the magnetic fie produced by the inner core.

The magnetization of the nucleon phase of the neutron star core matter occurs d to a spontaneous polarization of the proton impurities in the neutron matter [1], and it discussed in the next section. Magnetic phase of quark matter, considered in sect. 3, ca exist in the chirally broken quark phase, which develops a pion condensate [2]. Implicatio of the presence of a magnetized core for magnetic fields of neutron stars are discussed sect. 4. In particular accretion-induced magnetic field decay of millisecond pulsars [3] discussed there.

2. Polarized Protons in the Neutron Star Core

If the central density of a neutron star is below the deconfinement density the neutro star core is composed mostly of neutrons with some admixture of protons, electrons an muons which form a uniform liquid of densities exceeding the nuclear saturation densi $n_0 = 0.17 fm^{-3}$. The proton component is required for the beta stability of the system. Th proton fraction $x = n_P/n$ is expected to be of the order of a few percent. At the saturatio density n_0 the proton fraction is $x \approx 0.05$ [4], and it changes with increasing densit Various model interactions give different $x(n)$ [4], however for many realistic interactio the proton component disappears at sufficiently high densities [4]. In the calculatio presented here we use the Ravenhall's parametrization of the Friedman-Pandharipan equation of state (FPR) as given by Lattimer [5]. The proton fraction for the FPR equatio of state decreases from $x \approx 0.05$ at $n = 0.2 fm^{-3}$ to zero at $n = 0.9 fm^{-3}$ (the cur UV14+TNI in Fig.11 in Ref.[4]). The maximum mass of the neutron star for the FP equation of state is close to $1.8 M_\odot$.

The proton impurities strongly influence the magnetic properties of the neutron st matter. Under certain conditions they are likely to spontaneously polarize [1]. To s this let us compare energy of polarized and normal phase assuming the proton admixtu in the core to be of the order of a few percent. If the proton-neutron spin interaction approximated by the effective contact potential with a strength g^{PN}, we can write t change of the energy per unit volume with respect to the unpolarized phase as follows:

$$\delta\epsilon = \delta\epsilon_N + g^{PN}\delta s_N \delta s_P + \delta\epsilon_P, \qquad ($$

where δs_N and δs_P is the neutron and proton spin excess, respectively. The first and the third term describe respectively the change of the neutron and proton energies due to a small polarization. The second term represents the proton-neutron spin interaction. The main contributions to this interaction come from the one-pion exchange, the ρ-exchange and the second-order tensor interaction. These contributions, calculated in Ref.[6], give $g^{PN} \approx -2fm^2$. The change of the neutron energy density $\delta\epsilon_N$ can be expressed in terms of the Landau Fermi-liquid theory as

$$\delta\epsilon_N = \frac{1}{2N_N}(1 + G_0^{NN})\delta s_N^2. \tag{2}$$

Here $N_N = m_* k_F/\pi^2$ is the density of states at the Fermi level and G_0^{NN} is the spin dependent Landau parameter for pure neutron matter. Calculations of Ref.[7] show that G_0^{NN} depends weakly on density and we take $G_0^{NN} \approx 1$ in the whole density range of interest. This shows, that pure neutron matter does not possess a ferromagnetic phase. Minimizing $\delta\epsilon$, Eq.1, with respect to δs_N we find

$$\delta\epsilon_{min} = -\frac{N_N}{2(1 + G_0^{NN})}(g^{PN})^2\delta s_P^2 + \delta\epsilon_P. \tag{3}$$

This formula shows, that the spin instability ($\delta\epsilon_{min} < 0$) is controlled by the proton term $\delta\epsilon_P$. If this term makes only a small contribution, the system displays a spontaneous proton polarization. There are various ways for this to happen, depending on the values of $g^{PN}, G_0^{NN}, G_0^{PP}$ and the proton effective mass m_P^*. For example if the proton-neutron spin interaction g^{PN} is sufficiently strong and $G_0^{PP} \sim 0$ the spin instability occurs at higher densities [1]. The same effect occurs if the proton effective mass becomes sufficiently high and $G_0^{PP} \sim 0$ since then the change in the proton kinetic energy due to polarization is small. In an extreme case of localized protons $\delta\epsilon_P = 0$. Clearly, the behaviour of proton impurities is crucial for the spin properties of the system. The possibility of localization of proton impurities was considered in Ref.[8].

Assuming that protons are fully polarized we find the magnetization in the form

$$M = [-\frac{g^{PN}N_N}{1 + G_0^{NN}}\mu_N + \mu_P]n_P. \tag{4}$$

The quantity in the parenthesis is the effective magnetic moment per one proton. At low densities it is close to the bare proton magnetic moment μ_P, whereas at higher densities it changes sign and becomes parallel to the neutron magnetic moment μ_N. The reversal of the direction of the effective magnetic moment depends crucially on the sign of g^{PN}, since from minimization of $\delta\epsilon$, Eq.1, one finds $\delta s_N \sim -g^{PN}\delta s_P$. Hence for negative values of g^{PN} the induced neutron spin excess points in the same direction as the proton spin density. In Fig.1 we show the magnetization, Eq.4, as a function of the neutron density for $g^{PN} = -2.5fm^2$ and the FPR proton fraction (solid line) and also for a constant proton fraction $x = 0.05$ (dashed line). The magnetization changes sign at $n_r \approx 0.48fm^{-3}$. For the FPR proton fraction magnetization vanishes above $0.9fm^{-3}$. One can easily imagine that the dipole magnetic field of the neutron star will reflect this behaviour of the magnetization.

Our model contains three parameters. The first one is the proton-neutron spin interaction g^{PN}. It determines the density n_r at which the magnetization vanishes. We use the value $g^{PN} = -2.5fm^2$ which is 25% higher, than the value of Ref.[6]. The two other parameters are the critical densities n_l and n_u which limit the density range in which the magnetic phase exists. As far as the upper limit n_u is concerned a natural choice for the FPR equation of state is $n_u = 0.9fm^{-3}$, a density at which the proton fraction vanishes. This limit can be however lower, since at very low x the core temperature may exceed the Curie temperature. The lower critical density n_l should be determined by the nucleon hamiltonian. Here we shall treat both critical densities as adjustable parameters.

3. Magnetized Quark Matter with Broken Chiral Symmetry

If the central density of a neutron star exceeds the deconfinement density an inner quark core can appear. It is generally expected that the quark matter forms the ground state of baryon matter at densities above the deconfining transition. Deconfinement at zero temperature and high baryon density can be viewed as delocalization of valence quarks due to overlap of neighbouring nucleons, hence it is a geometrical rather than dynamical effect.

There is another fundamental phase transition in QCD: the chiral symmetry restoration, which makes the valence quarks massless (up to small current masses). This is a genuine dynamical effect and it would seem rather accidental if the chiral field vanished at the same density where valence quarks become deconfined. If deconfinement occurs at a lower density than the chiral symmetry restoration, the chirally broken quark matter is the ground state of baryon matter in the density range between the two transitions. Thus, we shall deal with deconfined quarks with "constituent"-like masses generated by the broken chiral symmetry.

We study here quark matter in the effective chiral quark-meson model of Ref.[9]. This model has been rather successful in reproducing properties of the lowest baryons [9]. Its dynamics is governed by the σ-model lagrangian with quark and meson degrees of freedom, which hopefully can be derived from QCD. It reads

$$\mathcal{L} = i\bar{\Psi}\gamma^\mu\partial_\mu\Psi - g\bar{\Psi}(\sigma + i\gamma_5\vec{\tau}\cdot\vec{\pi})\Psi - U(\sigma,\vec{\pi}) + \frac{1}{2}(\partial_\mu\sigma)^2 + \frac{1}{2}(\partial_\mu\vec{\pi})^2, \tag{5}$$

where Ψ is the quark field, σ and $\vec{\pi}$ are the meson fields, g is the quark-meson coupling constant, and U is the "Mexican Hat" potential, which leads to spontaneous chiral symmetry breaking. For the case of exact chiral symmetry ($m_\pi = 0$), the potential is:

$$U(\sigma,\vec{\pi}) = \frac{1}{4}\lambda^2(\sigma^2 + \vec{\pi}^2 - F_\pi^2)^2, \tag{6}$$

where $\lambda^2 = m_\sigma^2/2F_\pi^2$, and m_σ is the mass of the σ-meson. The vacuum expectation value of σ is $\sigma_0 = F_\pi$, where $F_\pi = 93$ MeV is the pion decay constant. The effective mass of the quark in this model is $m = g\sigma$, which in the vacuum becomes $m_0 = gF_\pi$. The values of the quark mass of the order of 500 MeV give good fits to baryon properties in model calculations of Ref.[9]. The value of the mass of the σ meson was taken above 1 GeV [9], however the nucleon fits were not sensitive to this number as long as it was large.

We find that at low densities the system has broken chiral symmetry, i.e. the expectation value of the chiral field does not vanish. One can distinguish two different phases with broken chiral symmetry: the normal quark matter with nonzero sigma field and vanishing pion field, and the pion-condensed phase with space-dependent pion and σ fields. The latter state has nontrivial spin properties as we show below.

In the normal phase (i.e. with no pion condensates) the pion field vanishes and the σ-field is constant in space, $\sigma = \bar{\sigma}$. The quarks acquire a mass $m = g\bar{\sigma}$. In the mean field approximation the system behaves as an ideal gas of fermions whose mass depends on density. The energy density is

$$\epsilon = \frac{1}{(2\pi)^3}\gamma \int_0^{k_F} d^3k(k^2 + g^2\bar{\sigma}^2)^{1/2} + \frac{1}{4}\lambda^2(\bar{\sigma}^2 - F_\pi^2)^2, \tag{7}$$

with

$$\gamma = \gamma_s N_f N_c, \tag{8}$$

where γ_s is the spin degeneracy, N_f is the number of flavours and $N_c = 3$ is the number of colours. For isospin-symmetric normal quark matter (i.e. with equal number of up and down quarks) $\gamma_s = N_f = 2$ and $\gamma = 12$. The quark Fermi momentum is $k_F = (6\pi n_q/\gamma)^{1/3}$, where n_q is the quark density; the baryon density is $n = n_q/3$. The last term in Eq.7 is the energy density of the σ-field resulting from the potential (6). The equilibrium value of the σ-field is obtained by minimizing the energy per baryon, $E = \epsilon/n$, with respect to $\bar{\sigma}$.

The pion-condensed phase has a nonvanishing expectation value of the pion field. We consider here the neutral pion condensate, for which only the π_3-component is nonzero. Following Dautry and Nyman [10] we assume the following ansatz for the σ and π_3 fields

$$\sigma(\vec{r}) = \bar{\sigma}cos(\vec{q}\cdot\vec{r}), \quad \pi_3(\vec{r}) = \bar{\sigma}sin(\vec{q}\cdot\vec{r}), \tag{9}$$

and $\pi_1 = \pi_2 = 0$. The chiral field wave vector $\vec{q}$, and the amplitude $\bar{\sigma}$ will be treated as variational parameters.

From the lagrangian (5) one derives the Dirac equation for quarks moving in the σ and π_3 fields (9). It has the form

$$(-i\vec{\alpha}\cdot\vec{\nabla} + \beta m exp(i\gamma_5\tau_3\vec{q}\cdot\vec{r}))\phi(\vec{r}) = E\phi(\vec{r}), \tag{10}$$

where $m = g\bar{\sigma}$ and ϕ is the Dirac spinor. This Dirac equation can be solved exactly by replacing ϕ by ϕ^R defined through

$$\phi(\vec{r}) = exp(-\frac{1}{2}i\gamma_5\tau_3\vec{q}\cdot\vec{r})\phi^R(\vec{r}). \tag{11}$$

The transformed quark spinors satisfy the equation

$$(-i\vec{\alpha}\cdot\vec{\nabla} - \frac{1}{2}\vec{\alpha}\cdot\vec{q}\gamma_5\tau_3 + \beta m)\phi^R(\vec{r}) = E\phi^R(\vec{r}). \tag{12}$$

which is independent of $\vec{r}$. It admits the plane wave solutions

$$\phi^R(\vec{r}) = exp(i\vec{k}\cdot\vec{r})\chi(\vec{k}), \tag{13}$$

with the spinor $\chi(\vec{k})$ satisfying the equation

$$(\vec{\alpha}\cdot\vec{k} - \frac{1}{2}\vec{\alpha}\cdot\vec{q}\gamma_5\tau_3 + \beta m)\chi(\vec{k}) = E(\vec{k})\chi(\vec{k}). \quad (14$$

This equation can be easily diagonalized giving the following spectrum [10]

$$E_\pm(\vec{k}) = \{m^2 + \vec{k}^2 + \frac{1}{4}(\vec{q})^2 \pm [m^2(\vec{q})^2 + (\vec{q}\cdot\vec{k})^2]^{1/2}\}^{1/2}. \quad (15$$

One can notice that the interaction term in the Dirac hamiltonian (10) can be writte as $-\frac{1}{2}\vec{\Sigma}\cdot\vec{q}\tau_3$, where $\vec{\Sigma} = \gamma_5\gamma_0\vec{\gamma}$ is the relativistic spin operator. Quasiparticles diagonalizin the interaction hamiltonian are thus the up ($\tau_3 = +1$) and down ($\tau_3 = -1$) quarks wit the spin polarized along the wave vector $\vec{q}$.

The ground state wave function of the system in the mean field approximation i assumed to be a single Slater determinant composed of quasiparticle wave functions tha satisfy Eq.10. The baryon density is

$$n = \frac{1}{3(2\pi)^3}\gamma\left[\int d^3k\Theta(E_F - E_-(\vec{k})) + \int d^3k\Theta(E_F - E_+(\vec{k}))\right], \quad (16$$

where the spin-flavour-colour degeneracy γ, Eq.8, is $\gamma = 6$ since for isospin-symmetric u and down quark matter both branches of the spectrum (15) have spin-flavour degenerac $\gamma_s N_f = 2$. The above integrals can be calculated analytically in cylindrical coordinate with k_z parallel to $\vec{q}$.

The energy density of the system which follows from the lagrangian (5) and Eq.9 is

$$\epsilon = \frac{1}{(2\pi)^3}\gamma\left[\int d^3kE_-(\vec{k})\Theta(E_F - E_-(\vec{k})) + \int d^3kE_+(\vec{k})\Theta(E_F - E_+(\vec{k}))\right]$$
$$+\frac{1}{2}\bar{\sigma}^2(\vec{q})^2 + \frac{1}{4}\lambda^2(\bar{\sigma}^2 - F_\pi^2)^2. \quad (17$$

The first term in brackets is the energy of the occupied Fermi seas of - and + quasiparticle The next term is the kinetic energy density of the chiral field and the last term is tl potential energy of the chiral field. The integrals in Eq.17 are also analytic.

At a fixed value of the baryon number density, the ground state corresponds to tl values of $q = |\vec{q}|$ and $\bar{\sigma} = m/g$ minimizing the grand potential Ω which for $T = 0$ is

$$\Omega = \epsilon - 3E_F n, \quad (18$$

where $3E_F = \mu$ is the baryon chemical potential, and ϵ and n are given by Eqs.17 and 1 respectively.

We show in Fig.2 the energy per baryon as a function of the baryon density f $m_0 = 500$ MeV and $m_\sigma = 1200$ MeV for the pion condensed phase (solid line with the lab C). The pion condensed phase has lower energy than the normal phase (N) in the whc density range shown. There is a strongly self-bound state in this phase with the energy p baryon of 1170 MeV at the "saturation density", $n_0 = 0.5fm^{-3}$. The thermodynamical

stable phase exists only for densities above n_0. We also show results (dashed lines) for the nonlinear sigma model [11].

The transformation (11) does not change the spin properties of the quasiparticles because the spin operator $\Sigma_z = \gamma_5\gamma_0\gamma^3$ commutes with the chiral rotation $\exp(-i\gamma_5\tau_3\vec{q}\cdot\vec{r}/2)$. In the rest frame of the quasiparticle, we have

$$[h^R(\vec{k}=0), \Sigma_z] = 0. \tag{19}$$

Here h_R is the rotated Dirac hamiltonian in the plane wave basis.

In the nonrelativistic case the relation (19) holds in any frame, and in the nonrelativistic limit all spins of the given flavour are aligned as was pointed out by Dautry and Nyman [10]. In this case the interaction hamiltonian is $-\frac{1}{2}\vec{\sigma}\cdot\vec{q}\tau_3$, where $\vec{\sigma}$ is the spin operator. Configuration obtained by filling only the $E_-(\vec{k})$ states has the same number of up and down quarks and zero spin density. There is, however, nonzero net magnetization since the magnetic moments of up and down quarks are different. Assuming Dirac magnetic moments we have $\mu_u = -2\mu_d$ where μ_u and μ_d are the up and down quark magnetons, respectively. The magnetization is

$$M = g(\mu_u s_u + \mu_d s_d), \tag{20}$$

where $s_u = \frac{1}{2}(n_{u\uparrow} - n_{u\downarrow})$ and $s_d = \frac{1}{2}(n_{d\uparrow} - n_{d\downarrow})$ are spin densities of up and down quarks, respectively, and $g = 2$ for Dirac particles.

In the relativistic case the spin densities s_u and s_d are somewhat modified. This is because relativistically the spin operator does not commute with the boost operator. The mean value of the spin is reduced in the relativistic case as compared with the nonrelativistic one. The spin densities s_u and s_d are calculated relativistically as

$$s_{u(d)} = \frac{1}{4}\langle\bar{\Psi}(1 \pm \tau_3)\gamma_0\Sigma_z\Psi\rangle. \tag{21}$$

One finds that these spin densities can also be expressed analytically [2].

In Fig.3 we show the ratio of relativistic and nonrelativistic spin densities as a function of the baryon density for $m_0 = 500$ MeV (solid lines) and for $m_0 = 300$ MeV (dashed lines). The ratio is less than unity and decreases with increasing density. Magnetization is still given by the formula (20). As one can see the relativistic effects reduce the magnetization of the system.

We have shown that the quark matter with neutral pion condensate exhibits spin ordering, which leads to nonzero magnetization of the system. This is a result of the fact, that the up and down quark Fermi seas are oppositely polarized along the direction of the pion field wave vector $\vec{q}$. The electric charges of up and down quarks have opposite sign and thus the magnetic moments of all the up and down quarks point in the same direction. The net magnetization of the system is however reduced as compared with the nonrelativistic case. If the quark core is present inside a neutron star it is composed of the chirally broken quark matter. Our results indicate that this is most likely the pion-condensed quark matter which possesses a net magnetization. The quark core could thus contribute to the magnetic moment of the star.

4. Accretion-Induced Magnetic Field Decay and Magnetic Core in a Neutron Star

The origin and evolution of the pulsar magnetic field remain two unsolved problems of the neutron star physics. Recently, however, a new coherent picture of the evolution of the field is emerging [12,13]. The evidence is accumulating the widely accepted view that the pulsar magnetic field decays on a time scale of 10^7 years may be incorrect. Only neutron stars that have been recycled in binaries show clear evidence for magnetic field decay [12]. To this category belong both millisecond pulsars and binary radio pulsars, which have $B_P < 10^{10.6}G$. These recycled pulsars probably highly contaminate the general population of single pulsars giving an impression that the magnetic fields of isolated neutron stars decay on such a relatively short time scale [14]. The recent observation of cyclotron lines from γ-ray-burst sources which are old neutron stars indicates that their magnetic fields have not decayed [15,16].

The analysis, summarized in Ref.[12], has important implications for models of the neutron star magnetic field. It strongly indicates that the magnetic field of an isolated neutron star has the value $B_P \sim 10^{12} - 10^{13}G$ and is permanent. On the other hand, the magnetic field of recycled neutron stars is diminished with respect to these values by an amount, which is directly related to the total amount of accretion that took place during the X-ray binary phase [17], i.e. the decay of the magnetic field is monotonically related to the mass increase of the accreting neutron star. We show, that these features are quite naturally accounted for in a polarized proton core model which we described in sect 2. In this model the dipole magnetic field of a neutron star is produced by a permanently magnetized matter in the dense core of the star. The magnetization results from the ferromagnetically ordered proton spins with density-dependent magnetic moments. For the heaviest neutron stars a quark core can appear which, if magnetized, can contribute to the total magnetic moment.

The magnetic field at the surface of the star is obtained by solving the Tolman-Oppenheimer-Volkoff equations for the FPR equation of state and integrating the magnetization over the polarized core. Using the radius R of a given star and its magnetic moment $\vec{M}$ we calculate the surface magnetic field at the magnetic pole $B_P = 2 \mid \vec{M} \mid / R^3$. Neutron stars with central densities below n_l are nonmagnetic. Neutron stars with central densities $n_l < n_c < n_r$ have magnetic moments increasing with the mass. If the central density passes n_r, there appears an oppositely magnetized inner core, and the total magnetic moment of the star decreases and so does the surface magnetic field. The maximum for all curves occurs at $M_r \sim 0.94M_\odot$, which corresponds to $n_c = n_r$. The field decreases until the mass reaches the value M_z, which corresponds to the polarized core with zero total magnetic moment and the magnetic field vanishes for M_z. This mass depends on n_l. For higher neutron star masses the inner core contributes more and the magnetic moment of the star starts to increase again, albeit in the opposite direction. For still heavier stars magnetized quark core appears and the magnetic moment rises dramatically (dashed line in Fig.4).

One can notice that the abrupt change of the magnetic field from $10^{12}G$ to zero occurs over a narrow neutron star mass range $M_z - M_r$ which is of the order of $0.1M_\odot$. This behaviour is, generally, consistent with the above picture of the evolution of the neutron

star magnetic field. If the initial pulsar mass corresponds to the field a few times $10^{12}G$ then accretion of a few tenths of solar mass leads to a significant decay of the magnetic field. The pulsar then evolves exactly as in the model of Shibazaki et al. [17]. One should however notice that if the accretion exceeds a certain amount, which for the initial mass equal to M is $\Delta M = M_x - M$, the magnetic field starts to increase again very rapidly. One can argue that this is exactly the case of the neutron star in an X-ray binary 4U 1627-27 which, as was pointed out by Verbunt et al. [18], accreted a lot of matter but still retains a strong magnetic field.

Recently new measurements of the mass of two millisecond binary pulsars were reported: Wolszczan [19] has determined the mass of the 37.9ms pulsar PSR1534+12 to be $1.32 \pm 0.03 M_\odot$ and Ryba and Taylor [20] have determined the mass of the 5.36ms pulsar PSR 1855+09 to be $1.27^{+0.23}_{-0.15} M_\odot$. This together with the very precise determination of the mass of the pulsar PSR1913+16, which is $1.442 \pm 0.003 M_\odot$ [21] allows for a more precise comparison of our model with the data. Magnetic fields of these pulsars are respectively $10^{9.98}G$, $10^{8.48}G$ and $10^{10.35}G$. In Fig.4 we show these values together with an (arbitrary) error bar for the magnetic field amounting to the factor of three uncertainty which should account for the unknown deviations of the pulsar radiation law from the magnetic rotator formula. The curves correspond to $n_l = 0.43 fm^{-3}$ and $n_u \approx 0.54 fm^{-3}$. For the curve 1 the FPR proton fraction is used. We also show results corresponding to scaling down the magnetization of Fig.1 by a factor of 0.6 (curve 2). One should notice that in order to move closer to the data points we used a rather thin shell of magnetized matter. When the magnetic quark core appears the field starts to increase much faster since magnetization of quark matter is much higher as shows the dashed line in Fig.4.

This work is partially supported by the KBN grant no. 2.0204.91.01. One of the authors (MK) wishes to thank Professor R. Mańka for inviting him to Szczyrk Summer School.

5. References

1. M. Kutschera and W. Wójcik, Phys. Lett. **223B**, 11 (1989).
2. M. Kutschera, W. Broniowski and A. Kotlorz, Nucl. Phys. **A516**, 566 (1990).
3. M. Kutschera and W. Wójcik, Cracow preprint INP 1555/PH, 1991.
4. R.B. Wiringa, V. Fiks and A. Fabrocini, Phys. Rev. **C38**, 1010 (1988).
5. J.M. Lattimer, Ann. Rev. Nucl. Part. Sci. **31**, 337 (1981).
6. S.O. Bäckman, G.E. Brown and J.A. Niskanen, Phys. Rep. **124**, 1 (1985).
7. S.O. Bäckman, C.G. Källman and O. Sjöberg, Phys. Lett. **43B**, 263 (1973).
8. M. Kutschera and W. Wójcik, Acta Phys. Pol. **B21**, 823 (1990).
9. M. C. Birse and M. K. Banerjee, Phys. Lett. **136B**, 284 (1984); Phys. Rev. **D31**, 118 (1985).
10. F. Dautry and E. M. Nyman, Nucl. Phys. **A319**, 323 (1979).
11. M. Kutschera, W. Broniowski and A. Kotlorz, Phys. Lett. **237B**, 159 (1990).

12. E.P.J. van den Heuvel, to appear in "Neutron Stars: Theory and Observation", eds. J. Ventura and D. Pines, Kluwer Academic Publishers, Dordrecht 1991.
13. D. Bhattacharya and E.P.J. van den Heuvel, Phys. Rep. **203**, 1 (1991).
14. M. Bailes, Astrophys. J. **342**, 917 (1989).
15. T. Murakami et al., Nature **335**, 234 (1988).
16. E.E. Fenimore et al., Astrophys. J. Lett. **335**, L71 (1988).
17. N. Shibazaki et al., Nature **342**, 656 (1989).
18. F. Verbunt, R.A.M.J. Wijers and H. Burm, Astron. Astrophys. **234**, 195 (1990).
19. A. Wolszczan, Nature **350**, 688 (1991).
20. M.F. Ryba and J.H. Taylor, Astrophys. J. **371**, 739 (1991).
21. J.H. Taylor and J.M. Weisberg, Astrophys. J. **345**, 434 (1989).

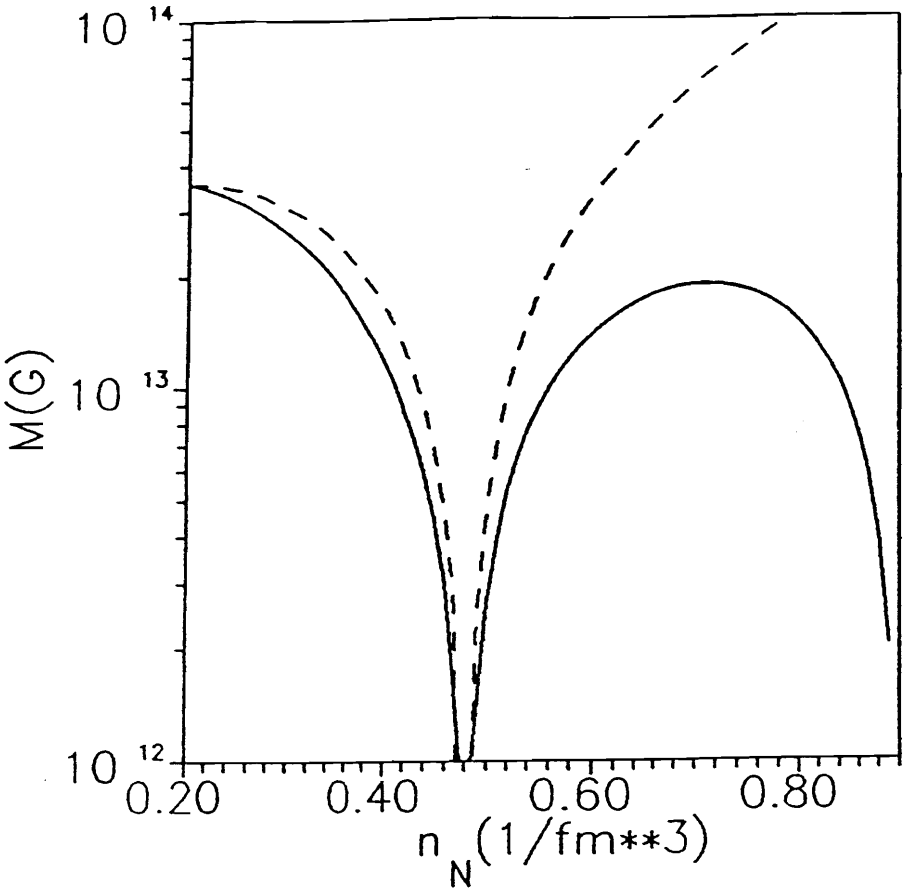

Fig.1 The magnetization of the neutron star matter as a function of neutron density. The solid and dashed lines correspond respectively to the FPR proton fraction and a constant proton fraction x=0.05.

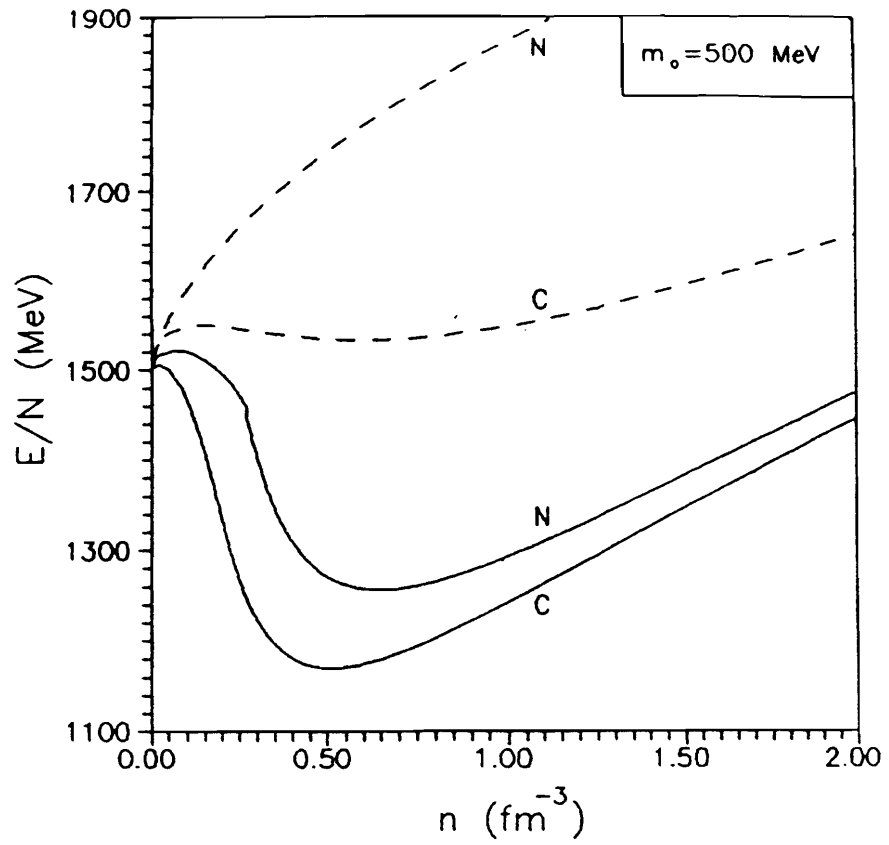

Fig.2 Quark matter energy per baryon as a function of baryon number density for isospin-symmetric quark matter. Solid lines correspond to m_σ = 1200 MeV and dashed lines correspond to the nonlinear σ-model. N and C label the normal and pion-condensed phases, respectively.

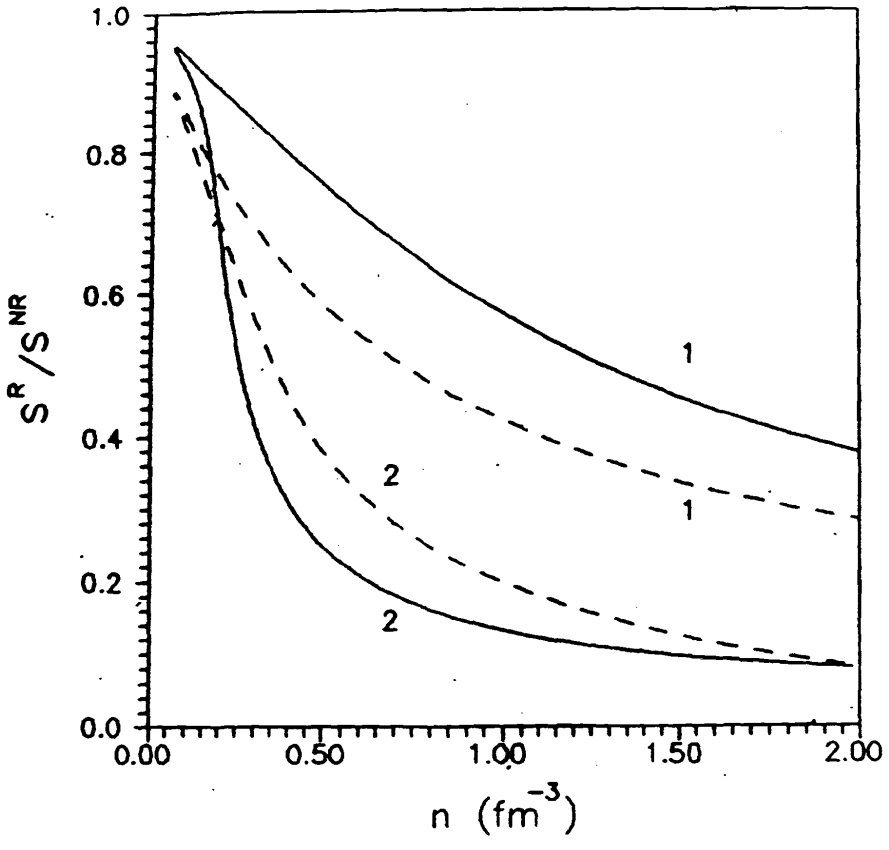

Fig.3 The ratio of relativistic and nonrelativistic spin densities as a function of the baryon number density for $m_0 = 500$MeV (solid curves) and $m_0 = 300$MeV (dashed curves). The labels 1 and 2 correspond to the nonlinear σ-model and to $m_\sigma = 1200$MeV, respectively.

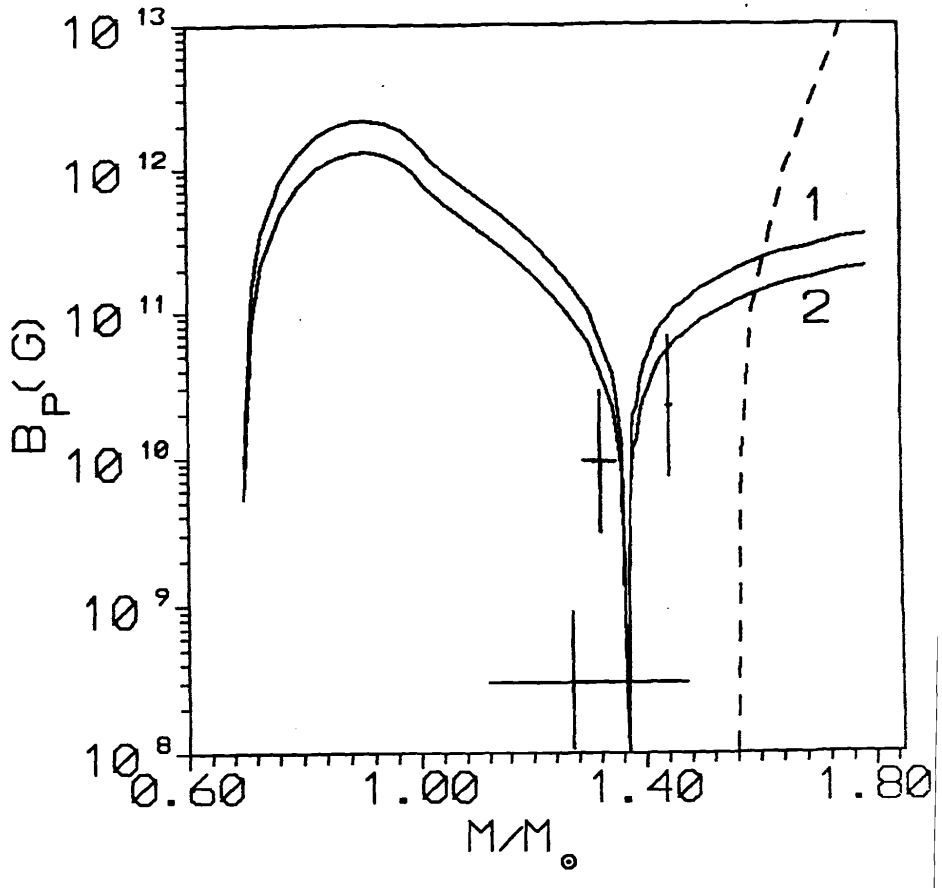

Fig.4 The surface magnetic field as a function of the neutron star mass for the FPI equation of state and for $n_l = 0.43 fm^{-3}$ and $n_u = 0.54 fm^{-3}$. The solid lines correspon to magnetized nucleon matter and the dashed line represents the magnetized quark core

IFJ - 115/91 - 130 egz.