PARITY VIOLATION IN NEUTRINO TRANSPORT AND THE ORIGIN OF PULSAR KICKS

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ABSTRACT

In proto–neutron stars with strong magnetic fields, the neutrino-nucleon scattering/absorption cross sections depend on the direction of neutrino momentum with respect to the magnetic field axis, a manifestation of parity violation in weak interactions. We study the deleptonization and thermal cooling (via neutrino emission) of proto–neutron stars in the presence of such asymmetric neutrino opacities. Significant asymmetry in neutrino emission is obtained because of multiple neutrino-nucleon scatterings. For an ordered magnetic field threading the neutron star interior, the fractional asymmetry in neutrino emission is about 0.006(B/10^14 G), corresponding to a pulsar kick velocity of about 200(B/10^14 G) km s^{-1} for a total radiated neutrino energy of 3 \times 10^{53} ergs.

Subject headings: dense matter — magnetic fields — pulsars: general — radiation transfer — stars: neutron — supernovae: general

1. INTRODUCTION

Recent analyses of pulsar proper motion (Lyne & Lorimer 1994; Lorimer, Bailes, & Harrison 1997; Hansen & Phinney 1994; Lorimer, Bailes, & Harrison 1997; Cordes, Romani, & Lundgren 1993), and studies of pulsar-supernova remnant associations (Frail, Goss, & Whiteoak 1994) indicate that supernova explosions are asymmetric and that the neutron stars receive large kick velocities at birth (250–500 km s^{-1} on average, but possibly with a significant population having velocities greater than 1000 km s^{-1}). Compelling evidence for supernova asymmetry also comes from the detection of geodetic precession in PSR J0045+1913 (Cordes, Wasserman, & Blaskiewicz 1990) and orbital plane precession in the PSR J0045–7319/B star binary (Kaspi et al. 1996). In addition, evolutionary considerations of the neutron star binary population also imply the existence of pulsar kicks (e.g., Fryer, Burrows, & Benz 1998).

Two classes of mechanisms for the natal kicks of pulsars have been suggested. The first class relies on local hydrodynamic instabilities in the proto–neutron stars (e.g., Burrows, Hayes, & Fryxell 1995; Burrows & Hayes 1996; Janka & Müller 1996) as well as global asymmetric perturbations seeded in the presupernova cores (Goldreich, Lai, & Scharling 1996). The asymmetries in matter and temperature distributions lead to asymmetric explosion and/or asymmetric neutrino emission, although the magnitude of the resulting kick velocity is still unclear.

In this Letter, we focus on the second class of models, in which the large pulsar velocities arise from asymmetric neutrino emission induced by strong magnetic fields. Although a number of papers have been written on this subject (Chugai 1984; Dorofeev et al. 1985; Binovatyi-Kogan 1993; Vilenkin 1995; Horowitz & Piekarewicz 1997; Horowitz & Li 1997), they are unsatisfactory for a number of reasons: they either failed to identify the most relevant neutrino emission/interaction processes, or the relevant physical conditions in proto–neutron stars, or stopped at estimating the magnetic field effects on neutrino opacities, and none of them treated neutrino transport in any systematic manner. As a result, their estimates of the kick velocity that may be produced are unreliable. We have carried out systematic studies of neutrino transport in strong magnetic fields (Lai & Qian 1998a, 1998b) in order to pin down the magnitude of the magnetic field effects on neutrino transport and pulsar kicks. Here we report our most important finding on the macroscopic consequence of asymmetric neutrino opacities due to parity violation in weak interactions. More details can be found in our forthcoming publication (Lai & Qian 1998a, hereafter LQ).

2. PARITY VIOLATION AND ASYMMETRIC NEUTRINO OPACITIES

In the presence of an external magnetic field, the differential opacity for neutrino-nucleon scattering, \( \nu_{e,p,r,I} \rightarrow \nu_{e,p,r,I} + N \), can be written in the form

\[
\frac{(d\sigma^{(sc)}_{\nu,E})_{\hat{n} \rightarrow \hat{n}^\prime}}{d\Omega_{\hat{n}^\prime}} = \frac{1}{4\pi} k^{(sc)}_{E\hat{n} \rightarrow \hat{n}^\prime} \times \left[ 1 + \epsilon_{\text{in}} \hat{\Omega} \cdot \hat{B} + \epsilon_{\text{out}} \hat{\Omega}^\prime \cdot \hat{B} + \text{const} (\hat{\Omega} \cdot \hat{\Omega}^\prime) \right],
\]

(1)

where \( \hat{\Omega} \) and \( \hat{\Omega}^\prime \) are the unit momentum vectors for the incoming and outgoing neutrinos, respectively, \( \hat{B} \) is the unit vector along the magnetic field direction, and \( k^{(sc)}_{E\hat{n} \rightarrow \hat{n}^\prime} \) is the total scattering opacity without magnetic field (the subscript \( E \) implies that the opacity depends on the neutrino energy). The presence of the \( \hat{B} \)-dependent terms in equation (1) implies that there is a preferred direction in the opacity. This is a manifestation of parity violation in weak interactions. The coefficients \( \epsilon_{\text{in}} \) and \( \epsilon_{\text{out}} \) are related to the nucleon polarization \( P \) by

\[
\epsilon_{\text{in}} = 2P \frac{c_A (c_A + c_V)}{c_V^2 + 3c_A^2}, \quad \epsilon_{\text{out}} = -2P \frac{c_A (c_A - c_V)}{c_V^2 + 3c_A^2},
\]

(2)

where \( c_V = -1/2 \), \( c_A = -1.15/2 \) for neutrons and \( c_V = 1/2 - 2 \sin^2 \theta_W = 0.035 \), \( c_A = 1.37/2 \) for protons (Raffelt &
Seckel 1995). For nondegenerate nucleons, the spin polarization $P = \mu_p B / kT = 3.15 \times 10^{-5} g_{14} B_{14}$ (10 MeV/T), where $\mu_p$ is the nucleon magnetic moment ($g_e = -1.913$ for neutrons and 2.793 for protons), and $B_{14}$ is the field strength in units of $10^{14}$ G. In the degenerate regime, we have $P = 3\mu_e B / (2\mu_p) = 4.73 \times 10^{-5} g_{14} B_{14} (10 \text{ MeV}/\mu_p)$, where $\mu_N$ is the nucleon chemical potential (or Fermi energy). A general expression for $P$ is given in LQ. In proto–neutron stars, one should average equation (1) over the composition (neutrons and protons). It is also important to note that the differential opacity for $\nu_{e,\mu,\tau}$ scattering on nucleons is obtained by interchanging $\epsilon_{in}$ and $\epsilon_{out}$ in equation (1) as a result of the crossing symmetry of the leading order matrix elements.

The neutrino absorption opacity is also asymmetric, i.e., $\kappa_{\text{abs}}^{(e)} = \kappa_{\text{abs}}^{(e)} (1 + \epsilon_{\text{asym}} \hat{\mathbf{B}} \cdot \hat{\mathbf{B}})$, with $\epsilon_{\text{asym}} \sim P$. However, in the bulk interior of the neutron star, this asymmetry is exactly canceled by the asymmetry in emission (see § 4). One also expects asymmetry in $\nu + e^- \nu + e^+$, however, for relativistic degenerate electrons, the $\nu-e$ opacity is much smaller than the $\nu-N$ opacity, while the electron polarization is $(200 \text{ MeV}/\mu_p)^{-2}$ of the same order as $P$ for nucleons. Therefore, the contribution from $\nu-e$ scattering to asymmetric neutrino transport is small compared with that from $\nu-N$ scattering.

3. TOY PROBLEM: EFFECT OF MULTIPLE SCATTERINGS

The asymmetry coefficients ($\epsilon_{in}$ and $\epsilon_{out}$) in the opacity are extremely small even at $B_{14} = 10$. However, the cumulative effect due to multiple scatterings can enhance the asymmetry in neutrino emission. To demonstrate this effect, we consider a simple toy problem in which neutrinos leak out of a slab (with an infinite extent perpendicular to the $z$-axis that coincides with the magnetic axis) via biased diffusion; i.e., the neutrino scattering probability has the form proportional to $(1 + \epsilon \cos \theta)$, with $\epsilon \ll 1$ and $\theta$ being the angle between the $z$-axis and the direction of neutrino propagation after each scattering. The evolution of the neutrino number density $n(z,t)$ is governed by the Fokker-Planck equation:

$$\frac{\partial n}{\partial t} = \frac{1}{3} c \nabla^2 n - \frac{1}{3} c \nabla \cdot (\epsilon_{\text{asym}} \hat{\mathbf{B}}),$$

where $\lambda$ is the (constant) mean free path of the neutrino. The surfaces of the slab are located at $z = \pm \lambda \tau$, (so $2\lambda \tau$ is the total optical depth of the slab). Let us suppose that we start with a uniform distribution $n = n_0$ at $t = 0$ (so that there is no initial asymmetry). As neutrinos gradually diffuse out, the profile $n(z,t)$ becomes increasingly asymmetric, and the flux asymmetry $\Delta F/F$ grows (where $F$ is the sum of the fluxes from both surfaces and $\Delta F$ is the difference). For $t \gg \tau_0$ ($\lambda/c$), the “cooling” front reaches the center of the slab, and $\Delta F / F$ saturates to constant value, $\Delta F / F = \epsilon \tau / 2$ (valid for $\epsilon, \ll 1$). This constant value can be derived analytically from a quasi–steady state analysis. The fractional asymmetry in the total neutrino radiation energy is $0.34\epsilon\tau$. In proto–neutron stars, the neutrino optical depth $\tau$ is large. A rough estimate can be made by setting the diffusion time $\sim 3 R^2 / (k \lambda)$ ($R$ is the neutron star radius) to 1 s (the spread in the arrival time of SN 1987A neutrinos was 10 s), which gives $\tau \sim 10^4$. This would imply a large enhancement in neutrino emission asymmetry. Several complications arise when one applies the result of the toy problem to real neutron stars: as shown in § 4, the asymmetry in neutrino flux is determined by the coefficient $\epsilon_{\text{out}} - \epsilon_{\text{in}}$ and the local neutrino energy distribution function. Due to the crossing symmetry, we have $(\epsilon_{\text{out}} - \epsilon_{\text{in}}) = - (\epsilon_{\text{in}} - \epsilon_{\text{out}})$. Because $\varphi_{(e)}^{(e)}$ and $\varphi_{(\mu)}^{(e)}$ have the same local energy distribution function, the “drift fluxes” (see eq. [6]) associated with $\varphi_{(e)}^{(e)}$ and $\varphi_{(\mu)}^{(e)}$ exactly cancel. The “drift fluxes” of $\nu_{\mu}$ and $\nu_{\tau}$ do not cancel because a newly formed neutron star is lepton rich, i.e., it contains more $\nu_{\mu}$ than $\nu_{\tau}$, and the typical energy of $\nu_{\mu}$ is also different from that of $\nu_{\tau}$. As the neutron star is depleptonized, the difference between $\nu_{\mu}$ and $\nu_{\tau}$ becomes smaller, and we expect the asymmetry in the total neutrino flux to diminish gradually. Another issue that is not clear from the toy problem concerns the role of neutrino absorption: for electron-type neutrinos, absorption is more important than scattering (by a factor of a few) in the depleptonization phase. Does the neutrino absorption completely wipe out the asymmetry associated with multiple scatterings?

4. THERMAL EVOLUTION OF PROTO–NEUTRON STARS WITH ASYMMETRIC NEUTRINO OCAPITIES

We now describe our calculation of the thermal evolution of proto–neutron stars in the presence of the asymmetric neutrino opacities discussed in § 2. We start from the general transport equation for the spectral intensity $I_E = I_E^d (r, \mathbf{B}, t)$ of a given neutrino species:

$$\frac{dI_E}{dt} + \hat{\mathbf{B}} \cdot \nabla I_E = \rho \kappa_{\text{abs}}^{(e)} (I_E^{(FD)} - I_E) + \rho (1 - f_E) \times \int \frac{dE^{(sc)}}{d\Omega} \frac{d}{d\Omega} [f_E^d (1 - f_E^d)],$$

where $I_E = I_E^d (r, \mathbf{B}, t)$, $f_E = (2\pi)^3 I_E / E^3$ is the neutrino occupation number, and $I_E^{(FD)}$ is the Fermi–Dirac distribution function for neutrinos:

$$I_E^{(FD)} (T) = \frac{c^3}{4\pi} U_E^{(FD)} (E) \frac{1}{(2\pi)^3} e^{E - \mu - \mu_{(\nu)}^{(FD)} T} + 1$$

($U_E^{(FD)}$ is the corresponding energy density), with matter temperature $T$ and neutrino chemical potential $\mu$. In equation (4), the absorption opacity $\kappa_{\text{abs}}^{(e)}$ already includes the effect of the stimulated absorption of neutrinos. Multiplying equation (4) by $\hat{\mathbf{B}}$ and then integrating over $\hat{\mathbf{B}}$, we obtain the first-order moment equation:

$$F_E = - \frac{c}{3 \rho \kappa_{\text{abs}}^{(e)}} (\nabla U_E + \epsilon_{\text{asym}} \kappa_{\text{abs}}^{(e)} (U_E^{(FD)} - U_E) \hat{\mathbf{B}}$$

$$+ (\epsilon_{\text{out}} - \epsilon_{\text{in}}) \kappa_{\text{abs}}^{(e)} c U_E (1 - f_E) \hat{\mathbf{B}}.$$
As noted earlier, the drift flux associated with $\mu$-type neutrinos is zero. The equations governing the thermal evolution of a proto-neutron star can be written as

$$\frac{\partial Y_{\nu}}{\partial t} = - \frac{1}{n} \nabla \cdot (S_{\nu} - S_{\nu}^d),$$

$$\frac{\partial U}{\partial t} = - \nabla \cdot (F_{\nu} + F_{\nu}^d + F_{\nu}^F).$$

Here $Y_{\nu} = n_{\nu}/n = Y_e + Y_\mu = (Y_e - Y_\mu) + (Y_\mu - Y_\nu)$ is the lepton number fraction ($n$ is the baryon number density), and $U$ is the internal energy density of the medium (a sum of contributions from neutrinos, electrons, positrons, photons, and nucleons). Note that inside the star, $\beta$-equilibrium holds to a good approximation; thus, we have $\mu_e + \mu_\mu = \mu_\mu + \mu_\nu - 1.293$ MeV ($\mu_e$ and $\mu_\mu$ do not include rest mass). For simplicity, we have neglected gravitational contraction and have adopted a Newtonian treatment.

We assume a uniform magnetic field in the $z$-direction and similarly for other relevant quantities. The angular dependence can then be factored out. We solve the evolutionary equations numerically using an explicit finite-difference code.

Our model neutron star has a mass $M = 1.38 M_\odot$, and a radius $R = 11$ km, with a density profile resembling a $T = 3$ polytrope (the central density is $\rho_c = 8 \times 10^{14}$ g cm$^{-3}$). The initial temperature and $Y_{\nu}$ profiles are shown in Figure 1. In the stellar interior, we have $Y_e = 0.28$, $Y_\mu = 0.366$. The temperature is 8 MeV in the inner core ($M \simeq 0.7 M_\odot$) and reaches 22 MeV in the outer core because of shock heating. These initial conditions are typical of those found in core collapse simulations. The initial chemical potentials of various particles at the center are $\mu_e = 97.7$ MeV, $\mu_\mu = 52$ MeV, $\mu_\nu = 313$ MeV, and $\mu_\nu = 266$ MeV. We assume that there is no asymmetry at $t = 0$. We adopt the zero boundary condition $T = 0$ at the stellar surface. Since the neutrino flux asymmetry mainly accumulates in the interior of the star, our crude treatment of the surface boundary condition is adequate.

Figure 1 depicts the $T_0$ and $Y_{\nu 0}$ profiles and the asymmetric perturbations ($\Delta T$ and $\Delta Y_{\nu}$) as functions of time for a magnetic field of strength $B_{\mu B} = 5$. The evolution of the $T_0$ and $Y_{\nu 0}$ profiles is similar to that found by Burrow & Lattimer (1986): the deleptonization phase ($t \lesssim 10$ s) is dominated by the lepton number diffusion. In this early phase, a heat wave moves inward, raising the core temperature. This is followed by a thermal cooling phase in which the core temperature decreases gradually because of energy diffusion. The asymmetries in the temperature and lepton number profiles increase during the deleptonization phase, reaching a maximum at $t = 8$ s (this behavior is analogous to that found in the toy problem discussed in § 3). This growth of asymmetry is also reflected by the increase in the ratio $|\Delta L_{\nu}/L_{\nu}|$ (where $L$ is the luminosity and $\Delta L_{\nu}/L_{\nu}$ is the net $z$-momentum carried away by the neutrinos per unit time), as shown in Figure 2b. As the neutron star becomes more lepton depleted, the difference (in number and energy) between $\nu_e$ and $\bar{\nu}_e$ diminishes, and the asymmetry gradually declines. Figure 2c shows the asymmetry in the integrated neutrino luminosity $[\mathcal{E}(t) = \int_0^t dt L(t)]$ as a function of time.
The asymmetry in the total radiated neutrino energy is \( \alpha = |\Delta \varepsilon_e / \varepsilon_e| = 0.006 B_{\perp} \). This result is smaller than suggested by the toy problem (§ 3) mainly because the energy carried away by electron neutrinos during the deleptonization phase is only about 10% of the total neutrino energy released. In addition, geometric factors and the cancellation of asymmetries associated with neutrons and protons tend to reduce the total asymmetry in neutrino emission.

5. CONCLUSION

In this Letter, we have identified the most efficient way of generating natal pulsar kicks based on asymmetric neutrino emission induced by magnetic fields. The key point of this mechanism is parity violation in weak interactions. The resulting pulsar kick velocity is

\[
V_{\text{kick}} \approx 1000 \left( \frac{\alpha}{0.084} \right) \left( \frac{E_{\text{tot}}}{10^{55} \text{ ergs}} \right) \left( \frac{1.4 M_\odot}{M} \right) \\
= 214 \left( \frac{\langle B_\perp \rangle}{10^{15} \text{ G}} \right) \left( \frac{E_{\text{tot}}}{3 \times 10^{55} \text{ ergs}} \right) \\
\times \left( \frac{1.4 M_\odot}{M} \right) \text{ km s}^{-1},
\]

where \( E_{\text{tot}} \) is the total energy released by neutrinos (of all species) from the proto–neutron star, and \( \langle B_\perp \rangle \) is the (averaged) ordered component of the magnetic field in the star. In reality, rotation (if it is misaligned with the magnetic axis) tends to reduce the kick velocity. We also note that the asymmetric \( \nu_e \) and \( \bar{\nu}_e \) emission from the proto–neutron star would give rise to asymmetric neutrino heating behind the stalled shock in the delayed supernova mechanism. This asymmetric heating could lead to asymmetric supernova explosions.

Another related (but distinct) kick mechanism relies on the fact that \( \nu_e \) and \( \bar{\nu}_e \) absorption opacities near the neutrinosphere depend on the local magnetic field strength (because of the quantization of \( e^- \) and \( e^+ \) energy levels). If the magnitude of the magnetic field in the north pole is different from that in the south pole, then asymmetric neutrino flux can be generated. We show in Lai & Qian (1998b) that this kick mechanism requires a much larger field strength than the mechanism considered in this Letter.

It is plausible that magnetic fields stronger than \( 10^{15} \text{ G} \) can be generated in proto–neutron stars (Thompson & Duncan 1993), and several pieces of evidence (albeit tentative) have been suggested in the literature (e.g., Vasisht & Gotthelf 1997). While hydrodynamically driven kick mechanisms may still be viable (depending on the magnitude of asymmetry seeded in the presupernova core), the possibility that a significant number of pulsars have kick velocities greater than 1000 km s\(^{-1}\) (e.g., Cordes & Chernoff 1998) may well require large magnetic fields (\( 10^{15} \text{ G} \)) to be present in proto–neutron stars.

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