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RAPID COOLING OF NEUTRON STARS BY HYPERONS AND Δ ISOBARS

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ABSTRACT

We show that direct Urca processes with hyperons and/or nucleon isobars can occur in dense matter as long as the concentration of Λ hyperons exceeds a critical value that is less than 3% and is typically about 0.1%. The neutrino luminosities from the hyperon Urca processes are about 5–100 times less than the typical luminosity from the nucleon direct Urca process, if the latter process is not forbidden, but they are larger than those expected from other sources. These new direct Urca processes provide avenues for rapid cooling of neutron stars which invoke neither exotic states nor the large proton fraction (of order 0.11–0.15) required for the nucleon direct Urca process.

Subject headings: dense matter — stars: neutron

The study of neutron star cooling is potentially an important source of information about the interior constitution of neutron stars. Until recently, the general view was that, if matter consisted of nucleons, cooling would be relatively slow, while if matter were in what is generally referred to as an exotic state, such as a pion condensate, a kaon condensate, or quark matter, cooling would be faster—so fast, in fact, that it would be difficult to observe thermal emission from the neutron star surface (for a review see Tsuruta 1986). Recently this view has been called into question, following the demonstration by Lattimer et al. (1991) that ordinary (nonexotic) matter with a proton/nucleon ratio in excess of some value estimated to lie between 0.11 and 0.15, can cool by the so-called direct Urca process even more rapidly than matter in an exotic state. Uncertainties in our knowledge of nuclear physics at high density so far preclude the prediction of whether or not the nucleon direct Urca process occurs in neutron stars.

Many calculations of the composition of dense matter lead to the conclusion that baryons more massive than the nucleon, namely, hyperons and nucleon isobars, should be present in dense matter, and in this *Letter* we consider neutrino emission from matter containing them. One of our most important conclusions is that even minuscule concentrations of Λ's can give rise to a luminosity in neutrinos exceeding that of exotic states. This situation had, in fact, been anticipated by Boguta (1981), who pointed out that large symmetry energies at high densities would allow rapid cooling via the direct Urca process to take place. This provides additional channels for rapid cooling of neutron stars which does not involve either exotic states or a large proton fraction.

The simplest possible neutrino-emitting processes are of the type

$$B_1 \to B_2 + \ell + \bar{\nu}_\ell \; ; \tag{1}$$

$$B_2 + \ell \to B_1 + \nu_\ell \; , \tag{2}$$

where B_1 and B_2 are baryons, and ℓ is a lepton, either an electron or a muon. These are so-called direct Urca processes, and the one for neutrons and protons has been considered previously (Lattimer et al. 1991). Here we shall consider processes where the baryons can be strange particles (Λ's, Σ's, and Ξ's) or nonstrange nucleon isobars (Δ^0, Δ^-) as well as nucleons.

At temperatures well below typical Fermi temperatures ($T_F \sim 10^{12}$ K), massive fermions of species i participating in the process must have momenta close to their respective Fermi momenta, p_{F_i}. In addition, the neutrino and antineutrino momenta are $\sim kT/c \ll p_{F_i}$. Since matter is very close to beta equilibrium, the chemical potentials of the constituents satisfy the condition

$$\mu_{B_1} = \mu_{B_2} + \mu_\ell \; , \tag{3}$$

which is equivalent to the general condition for chemical equilibrium

$$\mu_i = b_i \mu_n - q_i \mu_\ell \; , \tag{4}$$

where b_i is the baryon number of particle i and q_i is its charge. Thus, the condition of energy conservation is easily satisfied for some states close to the respective Fermi surfaces. The condition for momentum conservation, which plays a decisive role in discussions of neutrino emission, is that it must be possible to construct a triangle from the Fermi momenta of the two baryons and the lepton, and therefore the three triangle inequalities

$$p_{F_i} + p_{F_j} \geq p_{F_k} \; , \tag{5}$$

where i, j, and k are B_1, B_2, and ℓ, and cyclic permutations of them, must be satisfied. When the inequalities are not satisfied, the dominant neutrino emission processes are the modified Urca processes, which differ from equation (1) by the presence in the initial and final states of a bystander particle, whose sole purpose is to make possible conservation of momentum for particles close to the Fermi surfaces. For hyperons these modified Urca processes have been considered by Maxwell (1987).

The emissivity due to the direct Urca process may be derived from Fermi's golden rule. The rate of antineutrino energy emission from baryon (spin $\frac{1}{2}$) beta decay is given by

$$\epsilon = \frac{2\pi}{\hbar} 2 \sum_{p_i} \overline{\mathscr{M}^2} n_1(1 - n_2)(1 - n_3) E_4 \delta^{(4)}(p_1 - p_2 - p_3 - p_4) , \quad (6)$$

where n_i is the Fermi function, the subscripts $i = 1, 2$ refer to the baryons B_1 and B_2, and the subscripts $i = 3, 4$ refer to the lepton and the antineutrino, respectively. The p_i's are 4-momenta and E_4 is the antineutrino energy. Also, $\mathscr{M}$ is the matrix element for the decay of baryon B_1 to B_2, and we have assumed that wave functions are normalized in the usual nonrelativistic way. The bar over $\mathscr{M}^2$ denotes an average over spin states of the initial baryon and a sum over spins of particles in the final state. The sum over states is to be performed over 3-momenta $\boldsymbol{p}_i$, and the prefactor 2 takes into account the initial spins of the baryon B_1. The factors $1 - n_2$ and $1 - n_3$ express the fact that, because of the Pauli principle, the process cannot occur if the final states are occupied.

Assuming SU(3) symmetry, the matrix element for the process $B_1 \to B_2 \ell \bar{\nu}_\ell$ is given by (see, for example, Gaillard & Sauvage 1984)

$$M = (G_F/\sqrt{2}) C \bar{u}_{B_2} [\gamma_\mu f_1 + \gamma_\mu \gamma_5 g_1] u_{B_1} \bar{u}_\ell \gamma^\mu (1 + \gamma_5) u_\nu , \quad (7)$$

where very small tensor terms have been ignored. In the above, G_F is the weak coupling constant; $C = \cos \theta_C$ ($\sin \theta_C$) for a change of strangeness $|\Delta S| = 0$ (1); and θ_C is the Cabibbo angle. In the nonrelativistic limit, the squared matrix element is given by

$$\overline{\mathscr{M}^2} = \overline{M^2} \Big/ \left(\prod_{i=1}^{4} 2E_i \right)$$
$$= G_F^2 C^2 [(f_1^2 + 3g_1^2) + \boldsymbol{v}_\ell \cdot \boldsymbol{v}_\nu (f_1^2 - g_1^2)] , \quad (8)$$

where $\boldsymbol{v} = \boldsymbol{p}/E$ is the velocity. The parameters f_1 and g_1 of the baryon weak decay process are given in Table 1.

The phase space sums in equation (6) may be simply performed using the methods of Fermi-liquid theory. If the neutrino momentum is neglected in the momentum-conserving delta function, the contribution of the term involving $\boldsymbol{v}_\ell \cdot \boldsymbol{v}_\nu$ vanishes on integrating over angles. The energy emission in neutrinos by process (2) of the Urca pair is equal to that in antineutrinos by process (1), so the total emissivity in both neutrinos and antineutrinos is

$$\epsilon = \frac{457\pi}{10{,}080} \frac{G_F^2 C^2 (f_1^2 + 3g_1^2)}{\hbar^{10} c^5} m_{B_1} m_{B_2} \mu_\ell (kT)^6 \Theta_t \quad (9)$$

$$= 4.00 \times 10^{27} \left(\frac{Y_e n}{n_s} \right)^{1/3} \frac{m_{B_1} m_{B_2}}{m_n^2} R T_9^6 \Theta_t \text{ ergs cm}^{-3} \text{ s}^{-1}, \quad (10)$$

where Θ_t is the threshold factor, which is $+1$ if the triangle inequalities are satisfied and zero otherwise. The mass of species i is denoted by m_i, and T_9 is the temperature in units of 10^9 K. The electron fraction is $Y_e = n_e/n$, and $n_s = 0.16 \text{ fm}^{-3}$ is the equilibrium density of nuclear matter. The baryon squared matrix element normalized to the corresponding quantity for neutrons is denoted by R in equation (10). Explicitly,

$$R = \frac{[C^2(f_1^2 + 3g_1^2)]_{B_1 \to B_2 \ell \nu_\ell}}{[C^2(f_1^2 + 3g_1^2)]_{n \to p \ell \nu_\ell}} . \quad (11)$$

In comparison with the nucleon direct Urca process, the emissivities of hyperon processes are suppressed because they have smaller matrix elements, as shown in the last column of Table 1, which gives R for the various processes. Note that characteristic hyperon direct Urca emissivities, although smaller than that for nucleons, are larger by a factor of order $(T_F/T)^2 \sim 10^6 T_9^2$ than those of hyperon-modified Urca processes (Maxwell 1987). Interactions will affect neutrino emission in a number of ways, as was discussed by Lattimer et al. (1991), but they are not expected to reduce the emission by more than a factor of 10. Superconductivity and superfluidity of particles participating in an Urca process will reduce neutrino emission rates in a manner similar to that described by Lattimer et al. (1991) for the nucleon process.

Among other particles that we have not considered up to now are the Δ isobars, which are nonstrange spin 3/2 baryons with masses of about 1232 MeV, only a little greater than that of the Σ^-. Should any process with Δ's be kinematically allowed, it would result in neutrino energy emission comparable to the characteristic rate for the nucleon direct Urca process, and much greater than those for reactions in which there is a change of strangeness. Detailed rates for Δ reactions will be considered separately.

Throughout our discussion we have generally assumed that the lepton is the electron. Apart from the threshold factor Θ, the luminosity for Urca processes with muons is the same as that for the corresponding processes with electrons, since in equilibrium $\mu_e = \mu_\mu$. While the muon has a higher mass, and therefore a lower Fermi momentum, than the electron, it is not excluded that some processes with muons might have a lower threshold density than the corresponding electron process, since Fermi momenta in multicomponent interacting systems can depend on density in a complicated way.

The threshold concentration for an Urca process to occur is determined by the triangle inequalities in equation (5), and one general consequence of these is that, under conditions where neutrons are the most abundant baryons, the threshold concentrations for reactions in which neutrons participate will generally be higher than those for reactions without neutrons. For the first process listed, the one for nucleons, Lattimer et al. (1991) determined the threshold proton/nucleon ratio to lie in the range $x_c = 0.11$–0.15, the exact value of x_c depending upon the ratio of muons to electrons. However, it is not known whether this condition is ever satisfied in neutron stars,

TABLE 1

WEAK PROCESSES FOR NUCLEONS AND HYPERONS

Transition	C	f_1	g_1	R
$n \to p\ell\bar{\nu}_\ell$	$\cos \theta_C$	1	$F+D$	1
$\Lambda \to p\ell\bar{\nu}_\ell$	$\sin \theta_C$	$-\sqrt{3/2}$	$-\sqrt{3/2}(F + D/3)$	0.0394
$\Sigma^- \to n\ell\bar{\nu}_\ell$	$\sin \theta_C$	-1	$-(F-D)$	0.0125
$\Sigma^- \to \Lambda\ell\bar{\nu}_\ell$	$\cos \theta_C$	0	$\sqrt{2/3}D$	0.2055
$\Sigma^- \to \Sigma^0\ell\bar{\nu}_\ell$	$\cos \theta_C$	$\sqrt{2}$	$\sqrt{2}F$	0.6052
$\Xi^- \to \Lambda\ell\bar{\nu}_\ell$	$\sin \theta_C$	$\sqrt{3/2}$	$\sqrt{3/2}(F-D/3)$	0.0175
$\Xi^- \to \Sigma^0\ell\bar{\nu}_\ell$	$\sin \theta_C$	$\sqrt{1/2}$	$(F+D)/\sqrt{2}$	0.0282
$\Xi^0 \to \Sigma^+\ell\bar{\nu}_\ell$	$\sin \theta_C$	1	$F+D$	0.0564
$\Xi^- \to \Xi^0\ell\bar{\nu}_\ell$	$\cos \theta_C$	1	$F-D$	0.2218

NOTE.—The quantity $R = [C^2(f_1^2 + 3g_1^2)]_{B_1 \to B_2 \ell \nu_\ell} / [C^2(f_1^2 + 3g_1^2)]_{n \to p \ell \nu_\ell}$ was calculated using the central values of $\sin \theta_C = 0.231 \pm 0.003$, $F = 0.477 \pm 0.012$, and $D = 0.756 \pm 0.011$.

because the proton concentration is determined by the symmetry energy of matter above nuclear density, and this is poorly known. In models, such as relativistic mean field ones, in which the symmetry energy increases linearly with density, the condition is generally satisfied in neutron stars with masses exceeding 1 $M_{\odot}$. However, if the symmetry energy increases less rapidly with density, the critical density (if it exists at all) may never be attained in a given neutron star.

We turn now to hyperons and nucleon isobars. The composition of matter is determined by solving the chemical equilibrium conditions (4), subject to the conditions of charge neutrality and baryon number conservation. One expects that Λ, with a mass of 1116 MeV, the Σ^-, with a mass of 1197 MeV, and Δ^-, with a mass of 1232 MeV, first appear at roughly the same density, because the somewhat higher masses of the Σ^- and Δ^- are compensated by the presence of the electron chemical potential in the equilibrium condition for the Σ^- and the Δ^-. More massive, and more positively charged, particles than these appear only at higher densities, if they do so at all. Most model calculations show that positively charged particles appear at densities in excess of those found in 1.4 $M_{\odot}$ neutron stars, and in the following discussion we shall assume that they are not present. In that case, the individual concentrations of Σ^- and Δ^- will always be less than that of protons.

The question of whether any of the hyperon Urca processes are allowed when the nucleon Urca processes are forbidden is most conveniently addressed in terms of the proton-to-nucleon ratio $x = x_p/(x_n + x_p)$, where the x_i are the number fractions of species i. Recall that nucleon Urca processes are forbidden for $x < x_c$. If $x > x_c$, the Urca process will be permitted with nucleons, unless the concentrations of hyperons and/or nucleon isobars is large enough to reduce substantially the electron concentration. *In any event, if the nucleon Urca process is forbidden, the Urca processes* $\Sigma^- \rightarrow n + \ell + \bar{\nu}$ *and* $\Delta^- \rightarrow n + \ell + \bar{\nu}$ *are also forbidden, because* x_{Σ^-}, $x_{\Delta^-} < x_p$. Under these circumstances, the only Urca processes that could still be allowed are those involving Λ's: $\Lambda \rightarrow p + \ell + \bar{\nu}$, $\Sigma^- \rightarrow \Lambda + \ell + \bar{\nu}$, and $\Delta^- \rightarrow \Lambda + \ell + \bar{\nu}$.

The Urca threshold concentration of Λ's, x_{Λ_c}, can be roughly estimated from the triangle inequalities (5) by its value when the abundances of Σ^- and Δ^- are zero, namely, $x_{\Lambda_c} = |x_p^{1/3} - x_e^{1/3}|^3$. This threshold concentration is zero if the numbers of protons and electrons are equal, as they are if electrons and protons are the only charged particles present. Including the presence of muons but no additional particles, an upper limit for x_{Λ_c} when $x_{\Sigma^-} = x_{\Delta^-} = 0$ is found to be

$$x_{\Lambda_c} = \frac{x(2^{1/3} - 1)^3}{2 + x(2^{1/3} - 1)^3} \leq \frac{(2^{1/3} - 1)^3}{6(2^{1/3} + 1)} \simeq 0.0013 , \quad (12)$$

where the numerical value is obtained assuming $x < x_c$ (recall that if $x \geq x_c$, the more efficient nucleon Urca process is permitted when no hyperons or nucleon isobars are present). A more realistic estimate of x_{Λ_c} is obtained by relaxing the condition that $x_{\Sigma^-} = x_{\Delta^-} = 0$. When Σ^-'s and/or Δ^-'s are present, x_{Λ_c} may be larger or smaller than the value in equation (12) depending on the relative concentrations, but it is possible to show that an upper limit still exists:

$$x_{\Lambda_c} < \frac{x}{(1 + x^{1/2})(2 + 4x^{1/2} + 5x + 4x^{3/2} + 2x^2)} < 0.032 . \quad (13)$$

Thus the Urca process *must* occur either when $x \geq x_c$ and no hyperons or isobars are present, or, for any x (and independent of x_{Σ^-} and x_{Δ^-}), when $x_{\Lambda} \geq x_{\Lambda_c}$, which is, typically, a few parts per mille. Note that these conditions have been derived completely independently of any knowledge of nucleon and hyperon interactions. *It is important to appreciate that the* x_{Λ_c}*'s in equation (12) and equation (13) are exceedingly small.* Once a hyperon is present, its concentration increases rapidly with density, and consequently the threshold density for the Urca process for Λ differs little from the threshold density for its appearance. As a result of such low Λ Urca thresholds, at high temperatures there could be regions in the star where the process could occur with the Λ's being nondegenerate. The emissivity then varies as $T^{9/2}$.

Pioneering studies of hyperon concentrations were made by Cameron (1959), Ambartsumyan & Saakyan (1960), Tsuruta & Cameron (1966), and Cohen & Cameron (1971), all of whom found rather low threshold densities and rather high concentrations. Calculations with improved nucleon-nucleon and nucleon-hyperon interactions (Pandharipande 1971; Pandharipande & Garde 1972; Bethe & Johnson 1974) led to the conclusion that hyperons should be less ubiquitous. For example, Pandharipande & Garde found that only Σ^- hyperons would ever appear, and then only for a limited density range, $\sim 2.5n_s$–$4n_s$, while Bethe & Johnson found the hyperon threshold density to lie in the range $\sim 3n_s$–$6n_s$. In relativistic mean field models (Glendenning 1989; Glendenning & Moszkowski 1991; Ellis, Kapusta, & Olive 1991), hyperons appear at a density of about $2n_s$, with their abundance growing rapidly beyond their threshold densities. We have verified that, in these models, the triangle inequalities in equation (5) are met at Λ concentrations that are consistent with the upper limit in equation (13) and also that the direct Urca process is permitted soon after its appearance. However, in these models, the proton fractions always exceed x_c at the densities where hyperons are abundant, and therefore hyperon Urca processes only add to the more effective direct nucleon Urca process. As for nucleon isobars, Pandharipande (1971) found large concentrations, while more recent calculations suggest that their abundances are small, or even zero.

The basic reason that hyperon and isobar concentrations in dense matter are even more uncertain than proton abundances is that hyperon-nucleon and isobar-nucleon interactions at nuclear densities are less well determined from experiment than are nucleon-nucleon interactions, and information about them at higher densities is even more limited. In addition, essentially all recent estimates of hyperon abundances have been made using relativistic mean field models, in which correlations are not calculated in detail, and they consequently do not include aspects of physics known to be important in accounting for the properties of laboratory nuclei. The time is now ripe for taking a fresh look at the composition of dense matter, using available data on hyperon-nucleon and isobar-nucleon interactions with the inclusion of three-body terms (Bodmer, Usmani, & Carlson 1984), our knowledge of the properties of nuclear matter, and recent advances in techniques for calculating the energy of interacting systems.

The calculations above demonstrate that there are a number of Urca processes involving hyperons and isobars which could result in rapid cooling of neutron stars. Consequently, the number of possible ways of achieving rapid cooling is now quite large: the ones considered in this *Letter*, in addition to others, such as by pion condensates, kaon condensates, quark matter, and the nucleon direct Urca process considered earlier. The sensitivity of neutrino emission rates to details of the composition of dense matter provides impetus for a fresh look at this composition, using the latest developments in nuclear and hypernuclear physics and many-body theory.

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