EVOLUTION OF PROTO–NEUTRON STARS WITH KAON CONDENSATES

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Received 2000 August 24; accepted 2001 January 16

ABSTRACT

We present simulations of the evolution of a proto–neutron star in which kaon-condensed matter might exist, including the effects of finite temperature and trapped neutrinos. The phase transition from pure nucleonic matter to the kaon condensate phase is described using Gibb's rules for phase equilibrium, which permit the existence of a mixed phase. A general property of neutron stars containing kaon condensates, as well as other forms of strangeness, is that the maximum mass for cold, neutrino-free matter can be less than the maximum mass for matter containing trapped neutrinos or that has a finite entropy. A proto–neutron star formed with a baryon mass exceeding that of the maximum mass of cold, neutrino-free matter is therefore metastable, that is, it will collapse to a black hole at some time during the Kelvin-Helmholtz cooling stage. The effects of kaon condensation on metastable stars are dramatic. In these cases, the neutrino signal from a hypothetical galactic supernova (distance ~8.5 kpc) will stop suddenly, generally at a level above the background in the Super-Kamiokande and Sudbury Neutrino Observatory detectors, which have low-energy thresholds and backgrounds. This is in contrast to the case of a stable star, for which the signal exponentially decays, eventually disappearing into the background. We find the lifetimes of kaon-condensed metastable stars to be restricted to the range of 40–70 s and weakly dependent on the proto–neutron star mass, in sharp contrast to the significantly larger mass dependence and range (1–100 s) of hyperon-rich metastable stars. We find that a unique signature for kaon condensation will be difficult to identify. The formation of the kaon condensate is delayed until the final stages of the Kelvin-Helmholtz epoch, when the neutrino luminosity is relatively small. In stable stars, modulations of the neutrino signal caused by the appearance of the condensate will therefore be too small to be clearly distinguished with current detectors, despite the presence of a first-order phase transition in the core. In metastable stars, the sudden cessation in the neutrino signal occurs whether it is caused by kaon condensation, hyperons, or quarks. However, if the lifetime of the metastable star is less than about 30 s, we find that it is not likely to be due to kaon condensation.  

Subject headings: elementary particles — equation of state — stars: interiors — stars: neutron

1. INTRODUCTION

Proto–neutron stars (PNSs) are formed in the aftermath of gravitational collapse supernova, the end state of stars more massive than about 8 $M_\odot$. These objects are prodigious emitters of neutrinos of all types, which, if detected in terrestrial detectors, could reveal details of the supernova mechanism and the properties and composition of dense matter. The observation of neutrinos from supernova (SN) 1987A (Bionta et al. 1987; Hirata et al. 1987) confirmed the standard scenario for the early evolution of PNSs (Burrows & Lattimer 1986). At the beginning, this PNS is very hot and lepton-rich, and after a typical time of several tens of seconds, the star becomes deleptonized and cold: a neutron star has been formed. The detection of neutrinos radiating from the PNS surface is an unequivocal signal of the formation of this kind of object, since the direct collapse of the iron core into a black hole that might occur, perhaps in relatively extremely massive stars because of accretion, would result in a neutrino signal of rather short duration (Burrows 1988).

One of the chief objectives in modeling PNSs is to determine their internal compositions. Many simulations of dense matter predict the appearance of strange matter, in the guise of hyperons, a kaon condensate, or quark matter at supernuclear density (Prakash et al. 1997 and references therein). An important question is whether or not neutrino observations from a supernova could reveal the presence of such matter. One interesting possibility is that the existence of strange matter in neutron stars makes a sufficiently massive PNS metastable, so that after a delay of 10–100 s, the PNS collapses into a black hole (Brown 1994; Brown & Bethe 1994; Thorsson, Prakash, & Lattimer 1994; Prakash, Cooke, & Lattimer 1995; Glendenning 1995; Keil & Janka 1995; Ellis, Lattimer, & Prakash 1996). Such an event might be straightforward to observe as an abrupt cessation of neutrinos when the instability is triggered. Once the star becomes unstable, the collapse to a black hole proceeds on a timescale much shorter than the diffusion timescale, and the neutrinos still trapped in the inner regions cannot escape (Baumgart et al. 1996a; Baumgarte, Shapiro, & Teukolsky 1996b).

A previous paper (Pons et al. 1999, hereafter Paper I) presented calculations of the evolution of PNSs, studying the sensitivity of the results to the initial model, the total mass, the underlying equation of state (EOS), and the possible presence of hyperons. PNS simulations of stars with hyperons, not including hyperon contributions to the opacity, were earlier performed by Keil & Janka (1995). Paper I showed that the major effect on the neutrino signal before the onset of any possible metastability is the PNS mass; larger masses give rise to larger luminosities and gen-
erally higher average emitted neutrino energies. In addition, it was found that mass windows for metastable models could be as large as 0.3 $M_\odot$, ranging from baryon masses $M_B = 1.7$ to 2.0 $M_\odot$. The lifetimes of hyperonic metastable stars decrease with stellar mass and range from a few to longer than 100 s. The detection of neutrinos from SN 1987A over a timescale of 10–15 s is thus consistent with either the formation of a stable PNS or a metastable PNS containing hyperons, as long as its mass was less than about 0.1 $M_\odot$, below the maximum mass for cold, catalyzed hyperonic matter. Larger PNS masses would lead to a collapse to a black hole on a timescale shorter than that observed.

A similar situation could be encountered if the EOS allowed the presence of other forms of "exotic" matter, manifested in the form of a Bose condensate (of pions or kaons) or quarks (Prakash et al. 1995; Ellis, Lattimer, & Prakash 1996; Steiner, Prakash, & Lattimer 2000). As suggested some years ago (Kaplan & Nelson 1986), kaons obtain an effective mass that decreases with density in dense matter. As a consequence, the ground state of hadronic matter at high density might contain a kaon condensate (Brown et al. 1992; Muto & Tatsumi 1992; Maruyama et al. 1994; Thorsson, Prakash, & Lattimer 1994), which would result in a much softer EOS (Glendenning & Schaffner-Bielich 1998; Tatsumi & Yasuhira 1999; Pons et al. 2000, hereafter Paper II). However, a high lepton fraction, which exists in the early PNS evolution, suppresses the formation of kaons just as hyperon formation or the appearance of quarks is impeded (Prakash et al. 1995; Prakash et al. 1997).

The appearance of kaons after deleptonization produces a softening of the EOS, which can destabilize a sufficiently massive star, causing a collapse into a black hole.

In this paper we concentrate on models obtained in the context of kaon condensation and explore their minimal and maximal effects. One important difference between models containing hyperons and those containing a kaon condensate stems from the first-order phase transition that often exists in the latter case. For some models of kaon-nucleon interactions, mixed-phase regions can be present. Previous works (Glendenning 1992; Prakash et al. 1995; Glendenning & Schaffner-Bielich 1999; Paper II) have emphasized that satisfying the Gibbs' phase rules for equilibrium in situations when more than one conserved charge (in our case, baryon number and electric charge, lepton number being constrained by beta equilibrium) is present can radically change the pressure-density relation compared to the case in which a Maxwell construction is employed. Knorren, Prakash, & Ellis (1995) and Schaffner & Mishustin (1996) have emphasized that the threshold density for the appearance of strange particles, hyperons and kaons, is sensitive to the rather poorly known interactions in dense matter. In models in which hyperons appear at lower densities than kaons, the role of kaons can be relatively small. On the other hand, should it turn out that the interactions of the charged hyperons are relatively repulsive, only neutral hyperons may be present, and kaons can play a dominant role.

Although the possibility of PNS metastability was first discovered in the context of kaon condensation in neutron star matter (Thorsson et al. 1994), a full simulation of a PNS evolution with a kaon condensate has not been performed so far. Baumgarte et al. (1996b) have previously studied the hydrodynamical collapse of metastable PNSs containing kaon condensates. They were, however, mostly interested in the dynamical collapse to a black hole, not in the PNS evolution. They employed the kaon EOS from Thorsson et al. (1994), in which the transition from normal matter to the phase containing kaons was described by a Maxwell construction. They also rescaled the neutrino mean free paths by a factor of 1000 to accelerate the quasi-static part (i.e., the Kelvin-Helmholtz phase) of the evolution in order to save computation time. This produces little effect on the collapse to a black hole, since neutrinos are essentially trapped in the short time of collapse. However, their results for the Kelvin-Helmholtz stage can only be considered rough estimates because of this oversimplification.

In this paper, we perform the first evolutionary simulations of PNSs with kaon condensates during the Kelvin-Helmholtz epoch by using Henyey-like evolution and flux-limited diffusion modules as described in Paper I. As in Paper I, we ignore the effects of accretion, which are generally believed to be small after a successful supernova explosion, expected to occur within 1 s of core bounce. We use a field theoretical EOS that includes kaon condensation (Paper II) using boson exchange interactions, and we employ Gibbs' phase rules for multicomponent matter in which baryon number, electric charge, and lepton number are conserved during the evolution. We use the opacities developed by Reddy, Prakash, & Lattimer (1998); Burrows & Sawyer (1998, 1999); and Reddy et al. (1999) in the nucleon sector and account for the modifications that arise from the presence of kaon-condensed matter (Thorsson et al. 1995; Ji & Min 1998; Reddy 1998).

Our presentation is organized as follows: Section 2 provides an outline of the basic features of the EOS of matter containing a kaon condensate. In § 3 we describe the thermodynamical properties of kaon-condensed matter at finite temperature and with varying amounts of leptons likely to be encountered in the evolution of a PNS. Section 4 contains a description of the neutrino opacities in kaon-condensed matter. Results from numerical simulations of the Kelvin-Helmholtz phase of the evolution of a PNS with a kaon condensate are presented in § 5. Here we also explore the sensitivity of the results to the initial mass and consequences associated with metastable PNSs that collapse into black holes. The detectability of black hole formation in the Super-Kamiokande (SuperK) and Sudbury Neutrino Observatory (SNO) detectors from possible future galactic supernovae are investigated in § 6. Our conclusions are given in § 7.

2. EQUATION OF STATE OF KAON-CONDENSED MATTER

The EOS of kaon-condensed matter including the effects of trapped neutrinos and finite temperature for several different field theoretical models for both the nucleon-nucleon and kaon-nucleon interactions was studied in Paper II. Here we summarize the basic relations required for the computation of the EOS. Since our purpose in this paper is to explore whether a discriminating neutrino signal would be observed from kaon-condensed matter, we present results for cases in which a first-order phase transition occurs at zero temperature.

To model the baryonic phase, we use a field theoretical description in which baryons interact via the exchange of $\sigma$, $\omega$, and $\rho$ mesons. Including leptons, the Lagrangian density
is given by (Serot & Walecka 1986)
\[
L = L_{\text{int}} + L_{\ell} + L_{\text{pot}}
\]
\[
= \sum_{m} \bar{B}(-i\gamma^\mu \partial_\mu - g_{\omega N} \omega_n - g_{\rho N} \rho^\mu b_\mu \cdot t - M^*) B
\]
\[-\frac{1}{2} W_{\nu \mu} W_{\rho \delta} + \frac{3}{2} m_\omega^2 \omega_n \omega_\mu - \frac{1}{2} B_{\mu \nu} \rho^\nu + \frac{1}{2} m_\rho^2 b_\mu b_\nu
\]
\[
+ \frac{1}{2} \partial_\mu \sigma^\mu \sigma - \frac{1}{2} m_\sigma^2 \sigma^2 - U(\sigma) + \sum_{\ell} \{ -i\gamma^\mu \delta_{\mu \ell} - m_\ell \} \ell.
\]  
Here $B$ is the Dirac spinor for baryons, $t$ is the isospin operator, and $g$ and $m$ are the couplings and masses of the mesons, respectively. The sums include baryons $B = n, p,$ and leptons, $\ell = e^-, \mu^-$, and $\nu$. The field-strength tensors for the $\omega$ and $\rho$ mesons are $W_{\nu \mu} = \partial_\nu \omega_\mu - \partial_\mu \omega_\nu$ and $B_{\mu \nu} = \partial_\mu \rho_\nu - \partial_\nu \rho_\mu$, respectively. The potential $U(\sigma)$ represents the self-interactions of the scalar field and is taken to be of the form (Boguta & Bodmer 1977)
\[
U(\sigma) = \frac{1}{2} b_m g_m (g_{\sigma N} \sigma^3 + \frac{1}{4} c (g_{\sigma N} \sigma)^4,
\] where $m_b$ is the bare nucleon mass.

In the mean field approximation, the thermodynamic potential per unit volume is (Serot & Walecka 1986)
\[
\frac{\Omega_{\ell \nu}}{V} = \frac{1}{2} m_\sigma^2 \sigma^2 + U(\sigma) - \frac{1}{2} m_\omega^2 \omega_\sigma - \frac{1}{2} m_\rho^2 b_\sigma
\]
\[-2T \sum_{\nu, \rho} \int \frac{d^3 k}{(2\pi)^3} \ln \left[ 1 + e^{-\beta(E - \nu + \rho)} \right].
\]  
Here the inverse temperature is denoted by $\beta = 1/T$ and $E^* = (k^2 + M^*)^{1/2}$, with $M^* = m_b + g_b \sigma$ denoting the nucleon effective mass. The chemical potentials are given by $\mu_\rho = \nu_\rho + g_\rho \omega_\rho + 1/2 g_b b_\rho$, $\mu_b = \nu_b + g_b \omega_b - 1/2 g_b b_b$. (4)

Using $\Omega_{\ell \nu}$, the thermodynamic quantities can be obtained in the standard way.

The leptons are included in the model as noninteracting particles, since their interactions give negligible contributions compared to those of their free Fermi gas parts.

For the kaon sector, we use a Lagrangian that contains the usual kinetic energy and mass terms along with the meson interactions (Glendenning & Schaffner-Bielich 1999; Paper II). Kaons are coupled to the meson fields through minimal coupling; specifically,
\[
L_K = \mathcal{D}_\mu K^+ \mathcal{D}^\mu K^- - m_{K^\pm}^2 K^+ K^-,
\] where the vector fields are coupled via the standard form
\[
\mathcal{D}_\mu = \partial_\mu + ig_{\omega K} \omega_\mu + ig_{\rho K} \rho_\mu b_\mu t,
\] and $m_{K^\pm} = m_K - \frac{1}{2} g_{\omega K} \sigma$ is the effective kaon mass.

In the mean field approach, the thermodynamic potential per unit volume in the kaon sector is (Paper II)
\[
\frac{\Omega_K}{V} = \frac{1}{2} (f^0)^2 [m_{K^\pm}^2 - (\mu + X_0)^2] + \int_0^{\infty} \frac{d^3 p}{(2\pi)^3}
\]
\[
\times \left[ \ln \left[ 1 - e^{-\beta(\omega - \mu)} \right] + \ln \left[ 1 - e^{-\beta(\omega + \mu)} \right] \right],
\] where $X_0 = g_{\omega K} \omega_\sigma + g_{\rho K} b_\sigma$, the Bose occupation probability $f_0(x) = (e^{\beta(x)} - 1)^{-1}$, $\omega^2 = (p^2 + m_K^2)^2 + X_0$, $f = 93$ MeV is the pion decay constant, and the condensate amplitude, $\theta$, can be found by extremization of the partition function. This yields the solution $\theta = 0$ (no condensate), or if a condensate exists, the equation
\[
m_{K^\pm} = \mu_K + X_0,
\] where $\mu_K$ is the kaon chemical potential. In beta-stable stellar matter the conditions of charge neutrality,
\[
\sum_{\ell} q_{\ell} n_{\ell} - n_e - n_K = 0,
\] and chemical equilibrium,
\[
\mu_\ell = b_\ell \mu_e - q_{\ell} (\mu_{\ell} - \mu_e),
\] $\mu_K = \mu_e - \mu_p$
are also fulfilled.

The kaon condensate is assumed to appear by forming a mixed phase with the baryons satisfying Gibbs’ rules for phase equilibrium (Gibbs 1876). Matter in this mixed phase is in mechanical, thermal, and chemical equilibrium, so that
\[
p^I = p^II, \quad T^I = T^II, \quad \mu^I = \mu^II,
\] where the superscripts “I” and “II” denote the nucleon and kaon condensate phases, respectively. The conditions of global charge neutrality and baryon number conservation are imposed through the relations
\[
\chi^I q^I + (1 - \chi) q^II = 0,
\] $\chi^I p_{\ell}^I + (1 - \chi) n_{\ell}^II = n_{\ell}^I$, (13)
where $\chi$ denotes the volume fraction of nucleonic phase, $q$ the charge density, and $n_{\ell}$ the baryon density. We ignore the fact that the phase with the smallest volume fraction forms finite-size droplets; in general, this would tend to decrease the extent of the mixed-phase region. Further general consequences of imposing Gibbs’ rules in a multi-component system are that the pressure varies continuously with density in the mixed phase and that the charge densities must have opposite signs in the two phases to satisfy global charge neutrality. We note, however, that not all choices of nucleon-kaon and kaon-nucleon interactions permit the Gibbs’ rules to be satisfied (for an example of such an exception, see Paper II). The models chosen in this work do allow the Gibbs’ rules to be fulfilled at zero and finite temperatures and in the presence of trapped neutrinos.

The nucleon-kaon couplings are determined by adjusting them to reproduce the properties of equilibrium nucleonic matter at $T = 0$. We use the numerical values used by Glendenning & Moszkowski (1991), i.e., equilibrium density $n_0 = 0.153 \text{ fm}^{-3}$, equilibrium energy per particle of symmetric nuclear $E/A = -16.3$ MeV, effective mass $M^* = 0.78M$, compression modulus $K_0 = 240$ MeV, and symmetry energy $a_{sym} = 32.5$ MeV. These values yield the coupling constants $g_{\omega N}/m_\omega = 3.1507$ fm, $g_{\rho N}/m_\rho = 2.1954$ fm, $g_{\omega K}/m_\omega = 2.1888$, $b = 0.008659$, and $c = -0.002421$.

The kaon-meson couplings $g_{\omega K}$ and $g_{\rho K}$ are related to the magnitude of the kaon optical potential $U_K$ at the saturation density $n_0$ of isospin symmetric nuclear matter:
\[
U_K(n_0) = -g_{\omega K} \sigma (n_0) - g_{\rho K} \omega (n_0).
\]  
Fits to kaonic atom data have yielded values in the range $-50$ to $-200$ MeV (Friedman, Gal, & Batty 1994; Waas & Weise 1997; Friedman et al. 1999; Baca, Garcia-Recio, & Nieves 2000; Ramos & Oset 2000). We use $g_{\omega K} = g_{\omega N}/3$ and $g_{\rho K} = g_{\rho N}/2$ on the basis of simple quark and isospin counting. Given the uncertainty in the magnitude of $|U_K|$, consequences for several values of $|U_K|$ were explored in Paper II. Moderate values of $|U_K|$ generally produce a second-order phase transition and, therefore, lead to mod-
erate effects on the gross properties of stellar structure. Values in excess of 100 MeV were found necessary for a first-order phase transition to occur; in this case kaon condensation occurs at a relatively low density with an extended mixed-phase region, which leads to more pronounced effects on the structure due to a significant softening of the EOS.

3. COMPOSITION AND STRUCTURE OF PROTO–NEUTRON STARS WITH KAON CONDENSATES

In Figure 1 we show the variation of pressure with baryon density for three different choices of entropy per baryon, $s$, and lepton fraction $Y_L = (n_e + n_\mu + n_\tau)/n_B$. The optical potential $U_K$ is set to $-120$ MeV. In an evolving PNS, $s$ and $Y_L$ are not constant, but these three choices are reasonable approximations of the ambient conditions before $(s = 1, Y_L = 0.35)$ and after $(s = 2, Y_L = 0)$ the deleptonization stage as well as for a cold, neutrino-free $(s = 0, Y_L = 0)$ neutron star. The dotted line corresponds to lepton-rich matter, the dashed line to hot, neutrino-free matter, and the solid line to cold, neutrino-free matter. Results obtained from Gibbs’ construction are shown by heavy lines, while those for the pure phases are shown by thin lines.

Applying Gibbs’ rules eliminates the region of negative compressibility (this is more evident in the case of cold matter) and also decreases the critical density for the onset of kaon condensation. Above the threshold density for condensation, the kaon concentration builds up rapidly, which results in a significant softening of the EOS. The effects of trapped neutrinos are similar to those obtained for hyperons in Prakash et al. (1997) and Paper I. The most notable feature due to the presence of trapped neutrinos is that the critical density for kaon condensation is much higher in the case of neutrinos being trapped than in the case of neutrinos having left the star. Notice that the extent of a Maxwell construction would also be largely reduced in neutrino-trapped matter, and the differences between both kinds of mixed phase, Maxwell and Gibbs, are very small.

The phase boundaries of the different phases are displayed in Figure 2 in a $Y_L$-$n_B$ plane for an optical potential $U_K$ of $-100$ MeV (left-hand panel) and $-120$ MeV (right-hand panel). The nucleonic phase, the pure kaon matter phase, and the mixed phase are labeled I, II, and III, respectively. Solid lines mark the phase transition at zero temperature, and dashed lines mark the phase transition at an entropy per baryon of $s = 1$. Note that finite entropy effects are small and do not affect significantly the phase transition density. The dash-dotted line shows the electron fraction $Y_e$ as a function of density in cold, catalyzed matter (for which $Y_L = Y_e$), which is the final evolutionary state. The region to the left of this line corresponds to negative neutrino chemical potentials and cannot be reached during normal evolutions. The solid and dashed lines, which separate the pure phases from the mixed phase, vary roughly linearly with the lepton fraction. Also notice the large, and nearly constant, densities of the boundary between the mixed phase III and the pure kaon phase II. These densities, for the cases shown, lie above the central densities of the maximum mass stars, so that region II does not generally exist in proto–neutron stars (see Fig. 5 below). The effect of increasing the lepton number is to reduce the size of the mixed phase (which in fact shrinks to become a second-order phase transition for $Y_L > 0.4$ and $U_K = -100$ MeV) and shift the critical density to higher densities. A similar effect is produced by decreasing the magnitude of the optical potential.

Figure 3 shows the baryon (upper panel) and lepton (lower panel) chemical potentials for the three ambient conditions employed in Figure 1. In neutrino-free matter, we observe a rapid decrease of the electron chemical potential produced by the gradual substitution of electrons by kaons. When neutrinos are trapped, $\beta$-equilibrium prevents the electron chemical potential from decreasing because the neutrino chemical potential increases with density. In
Figure 3.—Baryon (top) and lepton (bottom) chemical potentials for the physical conditions shown in Fig. 1.

Figure 4.—Particle concentrations as a function of baryon density $n_B$ for the physical conditions shown in Fig. 1.

In Figure 5 we show the baryon mass ($M_B$: upper panels) and gravitational mass ($M_G$: lower panels) versus central baryon density for $U_K = -120$ (left-hand panels) and $-100$ MeV (right-hand panels), respectively. The onset of kaon condensation is shown by the diamond on each curve. In each case, the maximum mass of the cold, neutrino-free star is below the maximum mass of the hot star, irrespective of whether it has neutrino trapping. This figure indicates the existence of a mass window for metastability ($1.7 M_\odot < M_B < 2.0 M_\odot$ for $U_K = -120$ MeV and $2.0 M_\odot < M_B < 2.2 M_\odot$ for $U_K = -100$ MeV) in which a PNS will become unstable at the end of its Kelvin-Helmholtz stage. A similar effect was found in Paper I due to the presence of hyperons in matter. It should be emphasized that this figure uses constant $s$ and $Y_L$ profiles within the star. In fact, these quantities are not constant in the interior, so the precise mass window depends on the initial profile and the evolution. We have found, however, that the mass window exhibited in this figure is a reasonable approximation.

The main effect of lowering $|U_K|$ is a reduction in the size of the metastability window and its displacement to
The opacity for the neutral current-scattering processes involving nucleons in the presence of a charged kaon condensate are (Thorsson et al. 1996; Muto, Tatsumi, & Iwamoto 2000) of the growth rate for a kaon condensate indicate a timescale of less than or about $10^{-4}$ s at the temperatures of interest (above a few MeV). Since this is much shorter than the Kelvin-Helmholtz timescale (several seconds), we will assume that matter reaches chemical equilibrium instantaneously. However, for hydrodynamical studies, such as those involving collapse to a black hole or pulsations due to core quakes, the kaon condensation formation timescale might have to be considered.

Initial entropy per baryon, $s$, and lepton fraction, $Y_e$, profiles are taken to be the same as in Paper I, where we studied the evolution of PNSs containing either pure nucleons or nucleon/hyperon mixtures. The numerical code used to perform the simulations was also described in Paper I; the code is based on a Heneyy-like scheme in which the structure and transport equations are solved in sequential steps and then corrector steps are taken until convergence is reached.

For $M_B < 1.6 M_\odot$, the central density does not exceed the critical value for kaon condensation (for the set of parameters employed), and the evolution is identical to that of pure nucleon (np) models described in Paper I. For stars with $1.6 M_\odot < M_B < 2.05 M_\odot$, a kaon condensate will form during the evolution. We can classify the PNSs with a kaon condensate core in two subclasses: (1) stable, those

![Diagram](image-url)

**Fig. 6.—** Maximum baryon mass for the physical conditions shown in Fig. 1 as a function of the optical potential $U_K$. Metastable stars cannot exist when the cold, catalyzed maximum mass exceeds that of neutrino trapped matter. The triangle indicates the binary pulsar gravitational mass constraint $M_{\text{gw}} = 1.44 M_\odot$.

higher values, and this is shown in Figure 6. The value of $U_K$ must be limited to values greater than about $-126$ MeV in order that the binary pulsar mass constraint, $1.44 M_\odot$, not be violated. For values of $U_K$ greater than about $-80$ MeV, it is also clear that metastable models cannot exist. Another interesting feature is that the maximum masses of hot, neutrino-free models are nearly the same (for $U_K = -120$ MeV) or even larger (for $U_K = -100$ MeV) than those of lepton-rich configurations, preventing metastable stars from collapsing on deleptonization and delaying the onset of instability until the end of the cooling stage. This means that kaon-induced instabilities cannot occur very quickly, as can be the case with hyperon induced collapses.

4. NEUTRINO OPACITIES

For neutral- and charged-current processes involving nucleons and leptons, we utilize the recent advances made in the works of Reddy et al. (1998, 1999) and Burrows & Sawyer (1998, 1999). In these works, neutrino cross sections for arbitrary degeneracies were calculated, including the effects of interactions and collective excitations in beta-stable isospin asymmetric matter. These opacities form the basis of our baseline calculations against which modifications to the opacities due to the presence of a kaon condensate are assessed.

In bulk equilibrium and the degenerate limit, the factors by which the normal reaction rates must be modified in the presence of a charged kaon condensate are (Thorsson et al. 1995)

\[
\begin{align*}
\nu_e + n(K) &\rightarrow n(K) + e^- : \frac{1}{2} \sin^2 \theta \ tan^2 \theta_C , \\
\nu_e + p(K) &\rightarrow p(K) + e^- : \sin^2 \theta \ tan^2 \theta_C , \\
\nu_e + n(K) &\rightarrow p(K) + e^- : \cos^2 (\theta/2).
\end{align*}
\]

Here, $n(K)$ denotes an excitation that is a superposition of a neutron and a proton and reduces to a free neutron or proton in the absence of a charged kaon (or pion) condensate, $\theta$ denotes the amplitude of the condensate, and $\theta_C$ is the Cabibbo angle ($\sin \theta_C = 0.23$). The factors in this equation are all less than unity. The opacity for the neutral current-scattering processes involving nucleons in the presence of a kaon condensate has also been investigated recently by Ji & Min (1998) and Reddy (1998), who included the SU(3) singlet contribution. The general result is that a significant reduction in the cross section is caused by the presence of a kaon condensate formed in a bulk medium.

However, the situation is radically different for the opacity due to a kaon condensation in a mixed phase, in which finite-size effects may be important (Glendenning & Schaffner-Bielich 1998; Christiansen, Glendenning, & Schaffner-Bielich 2000; Reddy, Bertsch, & Prakash 2000). Reddy et al. (2000) found that coherent scattering of neutrinos off kaon-condensed droplets in the mixed phase increases the neutrino cross section, relative to bulk nucleonic matter by factors of 10–20. The overall enhancement, however, is sensitive to the droplet size, which depends on the surface tension.

We consider two extreme models, one in which the cross section or opacity is significantly smaller (20 times) than that of the neutrino-nucleon processes, to account for a condensate in bulk matter, and one in which the cross section is significantly larger (20 times), to account for possible effects of a droplet phase. We apply these factors to the overall opacity, while in fact the reduction factor for the bulk condensate should apply only to the volume occupied by the kaon-condensed phase. The mixed phase occupies only a small fraction ($\sim 0.2$) of the total volume even in the most massive PNSs, which mitigates considerably the large reductions in the opacity in this case. We wanted to ensure that our models will unmistakably bracket the true situation. Our results conclusively demonstrate that even large modifications to the opacity caused by kaon condensation do not have a pronounced effect on neutrino signals.

5. EVOLUTION

We now proceed to the evolution of a PNS in which a kaon condensate forms. We will focus on the case in which $U_K = -120$ MeV, for which the effects of a kaon condensation are relatively pronounced. Calculations (Prakash et al. 1996; Muto, Tatsumi, & Iwamoto 2000) of the growth rate for a kaon condensate indicate a timescale of less than or about $10^{-4}$ s at the temperatures of interest (above a few MeV).
with $M_B < 1.75 M_\odot$, which is the maximum mass of cold, neutrino-free neutron stars (see Fig. 5); and (2) metastable, those configurations with $1.75 M_\odot < M_B < 2.05 M_\odot$.

The evolution of stars containing kaon condensates (npK) deviates from those of np stars when the mixed phase with a condensate forms at the core. The size of this mixed phase grows on a diffusion timescale, i.e., the average time it takes for the neutrinos produced during kaonization to escape the inner core. In Figure 7, we show profiles of particle fractions as a function of radius for a stable star with $M_B = 1.7 M_\odot$ toward the end of the Kelvin-Helmholtz stage ($t = 60$ s). At this time, the star is at the beginning of its long-term cooling phase, during which its final catalyzed state is achieved. Notice that the size of the core is about 3 km, but it contains only a small fraction ($\approx 0.05 M_\odot$) of the star’s total mass.

In Figure 8, the evolutions of stable ($1.7 M_\odot$) and metastable stars are compared. Both the central baryon number densities and kaon fractions $Y_K$ are displayed. In each case, the time at which kaon condensation occurs is indicated by a diamond. Asterisks mark the times at which the evolution of metastable stars could not be followed further in our simulations, i.e., when a configuration in hydrostatic equilibrium could not be found. At this time, the PNS is unstable to gravitational collapse into a black hole. For the stable star, kaons appear after about 40 s. Thereafter, the star’s central density increases in a short interval, about 5 s, until a new stationary state with a mixed phase is reached. The evolution of the metastable stars is qualitatively different, inasmuch as the central density increases monotonically from the time the condensate appears to the time of gravitational collapse.

It is interesting that the lifetimes in all cases shown lie in the narrow range of 40–70 s (see Fig. 8). They decrease mildly with increasing $M_B$. Further insight is obtained by examining Figures 9 and 10. In Figure 9 the evolution of the central baryon density and kaon fraction are shown as functions of the central value of $Y_L$. The initial contraction of the star occurs at a nearly constant lepton fraction, as can be seen from the vertical trajectories of the central density. This is followed by a stage in which the lepton fraction decreases while little variation in the central density occurs.
until kaon condensation takes place (marked by diamonds in the figures). Thereafter, both the central density and the kaon fraction increase until the PNS reaches its final stage, either the cold-deleptonized configuration or the unstable configuration that collapses to a black hole.

Although the densities and lepton fractions for which condensation occurs are different for stars with different masses, the typical timescales are similar, as is evident from Figure 10, in which the evolution of the electron and neutrino concentrations, $Y_e$ (top panel) and $Y_\nu$ (middle panel), respectively, and the temperature (bottom panel) at the star’s center are shown. For all cases, the lepton concentrations exhibit a plateau stage until about 30–40 s and remain high enough to prevent the formation of a kaon condensate. Once the condensate is formed, $Y_\nu$ may increase as kaons gradually replace electrons in the inner core.

Notice that the plateau stage is also present in the evolution of stars containing only nucleons, as was shown in Paper I. Until a kaon condensate appears, the evolution of stars considered in this work proceeds similarly to that of a nucleons-only star because the number of thermal kaons present is rather small. The plateau is caused by an increase in neutrino opacity due to high temperatures, which reduces the leakage of neutrinos from the core, in turn maintaining a high lepton fraction. Consequently, kaonization of matter is delayed until the end of this epoch because of the dependence of the threshold density on $Y_\nu$. This explains why kaon-condensed stars of different masses become unstable at approximately the same time. This is qualitatively different from metastable hyperon-rich stars studied in Paper I, for which the lifetime is a much more sensitive function of mass. This is displayed in Figure 11, in which the lifetimes as a function of $M_B$ for stars containing hyperons (npH) and npK stars are compared. In both cases, the larger the mass, the shorter the lifetime. For kaon-rich PNSs, however, the collapse is delayed until the final stage of the Kelvin-Helmholtz epoch, while this is not necessarily the case for hyperon-rich stars.

Strangeness appearing in the form of a mixed phase of strange quark matter also leads to metastability. Although quark matter is also suppressed by trapped neutrinos (Prakash et al. 1995; Steiner et al. 2000), the transition to quark matter can occur at lower densities than the most optimistic kaon case, and the dependence of the threshold density on $Y_\nu$ is less steep than that for kaons. Thus, it is an expectation that metastability due to the appearance of quarks, as for the case of hyperons, might be able to occur relatively quickly. Calculations of PNS evolution with a mixed phase of quark matter, including the possible effects of quark matter superfluidity (Carter & Reddy 2000) are currently in progress and will be reported separately.

An interesting question concerns how short the lifetime of a kaon-condensed metastable star can be. For the case of $U_K = -120$ MeV, this time is constrained by the fact that metastability disappears for $M_B \approx 2.05 M_\odot$. For smaller magnitudes of the optical potential, the lifetimes are larger since the critical densities are larger; these densities are reached only after longer times. Conversely, increasing the magnitude of $U_K$ decreases the lifetime. The magnitude of the optical potential is limited, however, by the binary pulsar mass constraint that the maximum mass must exceed $1.44 M_\odot$ or $|U_K| < 126$ MeV for our interaction. In this case, the lifetime is found to be about 40 s. This leads to the interesting conclusion that should metastability from a
PNS with a timescale of, say, 30 s or less be observed, it would be evidence for the existence of hyperonic or quark matter rather than a kaon condensate.

6. LUMINOSITIES AND SIGNALS IN TERRRESTRIAL DETECTORS

The time dependence of the neutrino signal from a supernova is the only observable means through which the physical processes occurring in the high-density inner core of a PNS may be inferred. Our study is more applicable to times longer than approximately 1 s after core bounce, after which effects of dynamics and accretion become unimportant. Studies of the neutrino signal during the first second, during which approximately one-third of the neutrino energy is emitted, require more accurate techniques for neutrino transport (Messer et al. 1998; Yamada, Janka, & Suzuki 1999; Burrows et al. 2000) coupled with recent developments in supernova simulations (Rampp & Janka 2000; Mezzacappa et al. 2001).

For simplicity, we limit our study to the signal produced by electron antineutrinos absorbed onto protons in a pure water detector, for which the cross section is \( \sigma(E) = \sigma_0 E^2 \), where \( E_\nu \) is the \( \nu_e \) energy in MeV and \( \sigma_0 = 9.3 \times 10^{-44} \) cm\(^2\). This was the major reaction observed during the SN 1987A neutrino burst, and it is by far the process with the highest rate in the newer detectors. Assuming that the neutrinos leaving the PNS have a Fermi-Dirac spectrum with zero chemical potential, knowledge of the neutrino luminosity and average temperatures permit the signal from a supernova to be calculated according to (Paper I)

\[
\frac{dN}{dt} = \frac{\sigma_0 n_\nu}{4\pi D^2} \mathcal{M} G_4(E_{th}, T_\nu^*) F_4(0) T_\nu^* L_\nu. \tag{16}
\]

In this equation, \( L_\nu \) is the antineutrino luminosity (assumed to be one-sixth of the total neutrino luminosity), \( n_\nu = 6.7 \times 10^{28} \) free protons per kiloton of water, \( D \) is the distance to the supernova, \( \mathcal{M} \) is the fiducial mass of the detector, and \( T_\nu^* \) is the redshifted neutrino temperature. Also, \( F_4(0) = 7\pi^4/120 \) is an ordinary Fermi integral, and \( G_4(E_{th}, T) \) denotes a modified, truncated Fermi integral:

\[
G_4(E_{th}, T) = \int_{E_{th}}^\infty dz z^4 W(zT)(1 + e^z)^{-1}, \tag{17}
\]

with \( E_{th} \) being the detector threshold and \( W(E) \) the detector efficiency. We studied the signal observed in the Kamiokande II (KII) and Irving-Michigan-Brookhaven (IMB) detectors, whose masses, thresholds, and efficiencies are taken to be the same as in Lattimer & Yahil (1989), except that 6 ktons is used for the IMB mass, as discussed by Schramm & Brown (1990). We have assumed thresholds and fiducial masses of 5.6 (2) MeV and 22.4 (1.7) ktons for the SuperK (SNO) detectors, respectively, and we have also assumed a detector efficiency of unity above the threshold for both. Although we assumed that PNS neutrinos have a zero chemical potential spectrum, equation (16) does not depend very sensitively on this assumption (Lattimer & Yahil 1989).

The luminosity \( L_{BG} \) corresponding to a background rate below which the signal is lost, \( (dN/dt)_{BG} \), can be obtained from inverting equation (16):

\[
L_{BG} = 105.0 \left( \frac{dN}{dt} \right)_{BG} \frac{(D/10 \text{ kpc})^2}{(-\mathcal{M}/\text{kton})} F(T_\nu^*), \tag{18}
\]

where \( F(T_\nu^*) \) also depends on the threshold and efficiency of each detector. We have obtained the following analytical fits

\[
F(T) = \begin{cases} 
1 + \exp[2.0(1.2 - T)]/25.5(T - 0.75) & \text{KII, } T > 0.75 \\
(-5.27T + 2.57T^2)^{-1} & \text{IMB, } T > 2.05 \\
1 + \exp[2.25(1.2 - T)]/23.3T & \text{SuperK} \\
1 + \exp[6.0(0.4 - T)]/23.3T & \text{SNO},
\end{cases}
\]

where \( T \) is in MeV. In general, we will assume that the limiting background rate is about 0.2 Hz (one count every 5 s), since the corresponding time between counts is then about 10% of the entire signal time. A better estimate would require an in-depth statistical analysis that is beyond the scope of this paper.

In Figure 12 the evolution of the total neutrino energy luminosity is shown for different models. Notice that the drop in the luminosity for the stable star (solid line), associated with the end of the Kelvin-Helmholtz epoch, occurs at approximately the same time as for the metastable stars with somewhat higher masses. In all cases, the total luminosity at the end of the simulations is below \( 10^{51} \) ergs s\(^{-1}\). The two upper shaded bands correspond to SN 1987A...
detection limits with KII and IMB, and the lower bands correspond to detection limits in SNO and SuperK for a future galactic supernova at a distance of 8.5 kpc. The width of the bands represents the uncertainty in the $\bar{\nu}_e$ average energy because of the limitations of the flux-limited diffusion approximation, as discussed in Paper I. The times when these limits intersect the model luminosities indicate the approximate times at which the count rate drops below the background rate $(dN/dt)_{bg} = 0.2$ Hz.

The poor statistics in the case of SN 1987A precluded a precise estimate of the PNS mass. Nevertheless, had a collapse to a black hole occurred in this case, it must have happened after the detection of neutrinos ended. Thus the SN 1987A signal is compatible with a late kaonization-induced collapse, as well as a collapse due to hyperonization or to the formation of a quark core. More information would be extracted from the detection of a galactic SN with the new generation of neutrino detectors.

In SNO, about 400 counts are expected for electron antineutrinos from a supernova located at 8.5 kpc. The statistics would therefore be improved significantly compared to the observations of SN 1987A. A sufficiently massive PNS with a kaon condensate becomes metastable, and the neutrino signal terminates before the signal decreases below the assumed background. In SuperK, however, up to 6000 events are expected for the same conditions (because of the larger fiducial mass), and the effects of metastability due to condensate formation in lower mass stars would be observable.

It is interesting to assess the role of neutrino opacities in the presence of kaon condensates. For this purpose, we calculated the luminosities with two extreme assumptions concerning the opacities, namely, that the opacity is (1) 20 times smaller than the baseline opacity, corresponding to the bulk case (Thorsson et al. 1995; Ji & Min 1998), and (2) that it is 20 times larger than the baseline opacity, corresponding to finite-size effects in the mixed phase (Reddy et al. 2000).

In Figure 13, we show the temporal evolution of the central baryon density (top panel) and the total luminosity (bottom panel) for models with and without the correction to the opacity. In addition, we also show results corresponding to the case of pure np matter, in which the kaon condensate is not allowed. The magnitude of the opacity in the mixed phase chiefly controls the growth rate of the central density (top panel), and hence of the condensate, with larger opacities causing longer delays. The effect of reduced opacities on the luminosities (bottom panel) is barely distinguishable from the baseline case. In the case of enhanced opacities for the stable star the luminosity is first reduced by about 10%–20% for 10 s following the appearance of kaons at $t \sim 45$ s, but eventually exceeds that of the baseline case. Nevertheless, there is only a small change in the lifetime produced by even large opacity changes. The metastable star in Figure 13 with $2 \ M_\odot$ collapses before much information about the physics in the core reaches the neutrinosphere. It is also remarkable that the largest differences in the luminosities between the models with and without a condensate, but for the same opacity, are smaller than about 2%.

It is important to note, however, that by the time significant effects due to a kaon condensate become visible ($t > 50$ s), our treatment of neutrino transport becomes suspect since neutrinos in the interior are entering the semitransparent regime. This is reflected in the location of the neutrinosphere, which begins to fall well below the surface of the star for these times, and in the total optical depth of the star, which is no longer large. These affect our estimates of both the neutrino signal and the effective detector background.

The expected neutrino signals in the SuperK detector from a supernova at 8.5 kpc are compared in Figure 14 for a sample of stable PNSs with different compositions and masses. To show the effects of compositional changes, stars with $M_b = 1.7 \ M_\odot$ for np, npH, and npK cases are shown, and to show the effects of small changes in the stellar mass, an $M_b = 1.8 \ M_\odot$ for the case of npK is also displayed. The upper panel shows the count rate, and the lower panel shows the integrated number of counts. It is clear that, for the same stellar mass and stable models, the effect of having a kaon condensate core is smaller than the effect of hyperons: the np case is indistinguishable from the npK case. A change in the assumed mass of the PNSs of about 0.1 $M_\odot$ has about twice the effect on the number of counts as the

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Fig. 13.—Evolution of the central baryon density (top) and total neutrino luminosity (bottom) for two stars with $M_b = 1.7$ and $2.0 \ M_\odot$ for different assumptions concerning composition and opacity. Solid lines refer to the baseline model of a kaon-condensed star that assumes opacities corresponding to pure nucleonic matter. Dotted lines refer to stars containing only nucleons. Dashed (dot-dashed) lines refer to models of kaon-condensed stars in which the opacity is increased (decreased) by a factor 20. Asterisks indicate when metastable models become unstable to gravitational collapse. The shaded bands are as in Fig. 12.
librium, which permit a mixed phase. In the models, the phase is described by means of Gibbs’ rules for phase equilibrium from pure nucleonic matter to the kaon condensate at finite temperature and neutrino trapping. The phase transition is treated by employing an EOS that includes the effects of kaons, hyperons, or quarks.

The main conclusions arising from these comparisons is that the presence of a kaon condensate would be difficult to establish in an unambiguous fashion. For stars that are metastable, the small modulations in the time dependence of the neutrino luminosity or count rate due to kaon condensation can be easily masked by differences in composition (i.e., the presence of hyperons) or a small difference in the assumed PNS mass compared to a nucleons-only case. For stars that are metastable, the sudden cessation in the neutrino signal is similar for instabilities whether caused by kaons, hyperons, or quarks.

7. SUMMARY AND CONCLUSIONS

We have studied the effect of kaon condensation on PNS evolution by employing an EOS that includes the effects of finite temperature and neutrino trapping. The phase transition from pure nucleonic matter to the kaon condensate phase is described by means of Gibbs’ rules for phase equilibrium, which permit a mixed phase. In the models explored here, the central densities are not large enough to allow a pure condensed phase to exist.

We can classify stars of different masses in three main groups: (1) stars in which the central density does not exceed the critical value for kaon condensation, (2) stars that can form a mixed-phase core at the end of the Kelvin-Helmholtz epoch but remain stable, and (3) kaon-condensed metastable stars that become unstable and collapse to a black hole at the end of the Kelvin-Helmholtz epoch. For stars in which the effects of kaon condensation are small (this could be either because the star’s mass is low enough to permit only a very small region of the mixed phase or because condensation occurs as a weak second-order phase transition), the differences of the predicted neutrino signal compared to PNSs composed of pure nucleonic matter are very small.

Our calculations show that the variations in the neutrino light curves caused by the appearance of a kaon condensate in a stable star are small and apparently insensitive to large variations in the opacities assumed for them. Relative to a star containing only nucleons, the expected signal differs by an amount that is easily masked by an assumed PNS mass difference of 0.01–0.02 $M_\odot$. This is in spite of the fact that, in some cases, a first-order phase transition appears at the star’s center. The manifestations of this phase transition are minimized because of the long neutrino diffusion times in the star’s core and the Gibbs’ character of the transition. Both act in tandem to prevent either a “core quake” or a secondary neutrino burst from occurring during the Kelvin-Helmholtz epoch.

Observable signals of kaon condensation occur only in the case of metastable stars that collapse to a black hole. In this case, the neutrino signal for a star closer than about 10 kpc is expected to suddenly stop at a level well above that of the background in a sufficiently massive detector with a low-energy threshold such as SuperK. This is in contrast to the signal for a normal star of similar mass for which the signal continues to fall until it is obscured by the background. The lifetime of kaon-condensed metastable stars has a relatively small range, of the order of 50–70 s for the models studied here, which is in sharp contrast to the case of hyperon-rich metastable stars for which a significantly larger variation in the lifetime (a few to over 100 s) was found. This feature of kaon condensation suggests that stars that destabilize rapidly cannot do so because of kaons.

We determined the minimum lifetime for metastable stars with kaons to be about 40 s by examining the most favorable case for kaon condensation, which is obtained by maximizing the magnitude of the optical potential. The maximum optical potential is limited by the binary pulsar mass constraint, which limits the star’s maximum gravitational mass to a minimum value of 1.44 $M_\odot$. Therefore, should the neutrino signal from a future supernova abruptly terminate sooner than 40 s after the birth of the PNS, it would be more consistent with a hyperon- or quark-induced instability than one due to kaon condensation.

It is important to note that the collapse to a black hole in the case of kaon condensation is delayed until the final stages of the Kelvin-Helmholtz epoch, because of the large neutrino diffusion time in the inner core. Consequently, to distinguish between stable and metastable kaon-rich stars through observations of a cessation of a neutrino signal from a galactic supernovae is only possible using sufficiently massive neutrino detectors with low-energy thresholds and...
low backgrounds, such as the current SNO and SuperK and future planned detectors.

This work was supported in part by the Spanish DGCYT grant PB 97-1432 and the US Department of Energy under contract numbers DOE/DE FG02 87 ER-40317 (J. A. P. and J. M. L.) and DOE/DE FG02 88 ER-40388 (M. P.). J. A. P. thanks J. M. Ibáñez for useful discussions and encouragement.

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