Thermal Relaxation in Young Neutron Stars

Oleg Y. Gnedin\textsuperscript{1}, Dmitry G. Yakovlev\textsuperscript{2,3}, Alexander Y. Potekhin\textsuperscript{2}

\textsuperscript{1}Institute of Astronomy, Madingley Road, Cambridge CB3 0HA, England; ognedin@ast.cam.ac.uk
\textsuperscript{2}Ioffe Physical Technical Institute, 194021, St Petersburg, Russia
\textsuperscript{3}Institute for Theoretical Physics, University of California, Santa Barbara, CA 93106, USA

ABSTRACT

The internal properties of the neutron star crust can be probed by observing the epoch of thermal relaxation. After the supernova explosion, powerful neutrino emission quickly cools the stellar core, while the crust stays hot. The cooling wave then propagates through the crust, due to its finite thermal conductivity. When the cooling wave reaches the surface (age $10^{-100}$ yr), the effective temperature drops sharply from 250 eV to 30 or 100 eV, depending on the cooling model. The crust relaxation time is sensitive to the (poorly known) microscopic properties of matter of subnuclear density, such as the heat capacity, thermal conductivity, and superfluidity of free neutrons. We calculate the cooling models with the new values of the electron thermal conductivity in the inner crust, based on a realistic treatment of the shapes of atomic nuclei. Superfluid effects may shorten the relaxation time by a factor of 4. The comparison of theoretical cooling curves with observations provides a potentially powerful method of studying the properties of the neutron superfluid and highly unusual atomic nuclei in the inner crust.

Key words: dense matter – stars: neutron – X-rays: stars.

1 INTRODUCTION

Neutron stars are natural astrophysical laboratories of superdense matter. In their cores, at densities above the nuclear matter density $\rho_0 = 2.8 \times 10^{14} \text{ g cm}^{-3}$, the properties of matter such as equation of state and even composition are largely unknown (e.g., Lattimer & Prakash 2000, and references therein). In the absence of exact theory of superdense matter, different theoretical models predict different equations of state and compositions (neutrons, protons and electrons; hyperons; pion or kaon condensates; deconfined quarks).

One of the potentially powerful methods to probe the internal structure of isolated neutron stars is modelling of their cooling (e.g., Pethick 1992; Page & Applegate 1992; Page 1998a,b). The theoretical cooling curves depend on the adopted physical models of the stellar interior, especially the neutrino emission and heat capacity, as well as the superfluidity of neutrons and protons in the core. Confronting theory and observations allows one, for example, to constrain the range of critical temperatures of the superfluidity (e.g., Yakovlev et al. 1999).

Observing thermal emission of the very young neutron stars, $t \lesssim 100$ yr, opens a possibility of studying the properties of the neutron star crust. Soon after a supernova explosion, the young star has large temperature gradients in the inner parts of the crust. While the powerful neutrino emission quickly cools the core, the crust stays hot. The heat gradually flows inward on a conduction timescale, and the whole process can be thought of as a cooling wave propagation from the center towards the surface. During this thermal relaxation the effective temperature stays almost constant at about 250 eV. When the cooling wave reaches the surface, the effective temperature drops sharply by as much as an order of magnitude in the fast cooling scenario, and by a factor of 2 – 3 in the slow cooling scenario. The duration of the relaxation epoch depends mainly on the heat capacity and thermal conductivity of the inner crust (Lattimer et al. 1994).

Although the equation of state of matter at subnuclear density is known sufficiently accurately (Negele & Vautherin 1973, Pethick & Ravenhall 1995), the properties of atomic nuclei are not. The nuclei become unusually neutron-rich, with the smooth proton and neutron distributions. At the bottom of the crust, $\rho \gtrsim 10^{14} \text{ g cm}^{-3}$, the nuclei can be nonspherical and form clusters (Lorenz et al. 1993, Pethick & Ravenhall 1995). The liquid of neutrons dripped from the nuclei may be superfluid, with the critical temperatures that are very model-dependent. The thermodynamic and transport properties of this matter are subject of large theoretical uncertainty.

In this paper we make improved models of the young neutron stars. We obtain the new values of the electron thermal conductivity in the inner crust, based on a realistic...
treatment of the shapes of atomic nuclei. Using a new
numerical code, we calculate the cooling models and determine
the duration of the thermal relaxation epoch. We extend the
analysis of Lattimer et al. (1994) and derive the dependence
of the relaxation time on the microscopic parameters
of the crust. Our preliminary results have been summarized
by Yakovlev et al. (2000).

2 COOLING MODEL

2.1 Equations of thermal evolution

Neutron stars are born very hot in supernova explosions,
with the internal temperature \( T \sim 10^{11} \) K, but gradually
cool down. About twenty seconds after the birth, they be-
come fully transparent for the neutrinos generated in numer-
ous reactions in stellar interiors. We consider the cooling in
the following neutrino-transparent stage. The cooling is re-
alized via two channels, by neutrino emission from the entire
stellar body and by heat conduction from the internal lay-
ers to the surface resulting in thermal emission of photons.
For simplicity, we neglect the possible reheating mechanisms
(frictional dissipation of the rotational energy, Ohmic decay
of the internal magnetic field, or the dissipation associated
with weak deviations from the chemical equilibrium; see,
e.g., Page 1998a).

The internal structure of neutron stars can be re-
garded as temperature-independent (e.g., Shapiro & Teukol-
sky 1983). The relativistic equations of thermal evolution
include the energy and flux equations (Thorne 1977):

\[
\frac{1}{4\pi r^2 c^2 \Phi} \sqrt{1 - \frac{2GM}{c^2 r}} \frac{\partial}{\partial r} \left( e^{\Phi} L_r \right) = -Q_\nu - C_v \frac{\partial T}{\partial t} \quad (1)
\]

\[
L_r = \frac{\kappa}{4\pi r^2} \sqrt{\frac{2GM}{c^2 r}} e^{-\Phi(r)} \left( T e^\Phi \right) \quad (2)
\]

where \( Q_\nu \) is the neutrino emissivity [erg cm\(^{-3}\) \( s^{-1}\)], \( C_v \)
is the specific heat capacity [erg cm\(^{-3}\) K\(^{-1}\)], \( \kappa \) is the thermal
conductivity, and \( L_r \) is the “local luminosity” defined as
the non-neutrino heat flux transported through a sphere of
radius \( r \). The gravitational mass \( m(r) \) and the metric func-
tion \( \Phi(r) \) are determined by the stellar model. After thermal
relaxation, the redshifted temperature \( \tilde{T}(t) \equiv T(r,t) e^{\Phi(r)} \)
becomes constant throughout the interior.

It is conventional (Gudmundsson et al. 1983) to subdi-
vide the calculation of heat transport in the neutron star
interior \( r < R_b \) and in the outer heat-blanketing envelope
\( R_b \leq r \leq R \), where \( R \) is the stellar radius, and the bound-
ary radius \( R_b \) corresponds to the density \( \rho_b = 10^{10} \) g cm\(^{-3}\)
(\( \sim 100 \) meters under the surface). The thermal structure
of the blanketing envelope is studied separately in the sta-
nary, plane-parallel approximation to relate the effective
surface temperature \( T_s \) to the temperature \( T_b \) at the inner
boundary of the envelope. We use the \( T_s - T_b \) relation
obtained by Potekhin et al. (1997) for the envelope composed
mostly of iron.

The effective temperature determines the photon lu-
iminosity: \( L_\gamma = L_r(r,t) = 4\pi R^2 T_\gamma^4(t) \). Both \( L_\gamma \) and \( T_\gamma \)
refer to the locally-flat reference frame on the surface.
A distant observer would register the “apparent” luminosity

| \( M_\odot \) | \( R \) (km) | \( \rho_c 14 \) | \( M_{\text{crust}} \) (\( M_\odot \)) | \( \Delta R_{\text{crust}} \) (km) | \( M_D \) (\( M_\odot \)) | \( R_D \) (km) |
|---|---|---|---|---|---|---|
| 1.1 | 12.20 | 8.50 | 0.050 | 1.66 | ... | ... |
| 1.2 | 12.04 | 9.52 | 0.044 | 1.45 | ... | ... |
| 1.3 | 11.86 | 10.70 | 0.039 | 1.26 | ... | ... |
| 1.4 | 11.65 | 12.20 | 0.033 | 1.09 | ... | ... |
| 1.44 | 11.54 | 12.98 | 0.031 | 1.02 | 0.000 | 0.00 |
| 1.5 | 11.38 | 14.20 | 0.028 | 0.93 | 0.065 | 2.84 |
| 1.6 | 11.01 | 17.20 | 0.022 | 0.77 | 0.301 | 4.61 |
| 1.7 | 10.37 | 23.50 | 0.016 | 0.59 | 0.685 | 5.79 |
| 1.73 | 9.71 | 32.50 | 0.011 | 0.47 | 0.966 | 6.18 |

\( ^a \) threshold configuration for switching on direct Urca process
\( ^b \) configuration with maximum allowable mass
\( ^c \) central density in 10\(^{14} \) g cm\(^{-3}\)
\( ^d \) \( \Delta R_{\text{crust}} \) is defined as \( R - R_{\text{core}} \)

The internal structure of neutron stars can be re-
garded as temperature-independent (e.g., Shapiro & Teukol-
sky 1983). The relativistic equations of thermal evolution
include the energy and flux equations (Thorne 1977):

\[
L_r = \frac{\kappa}{4\pi r^2} \sqrt{\frac{2GM}{c^2 r}} e^{-\Phi(r)} \left( T e^\Phi \right) \quad (2)
\]

where \( Q_\nu \) is the neutrino emissivity [erg cm\(^{-3}\) \( s^{-1}\)], \( C_v \)
is the specific heat capacity [erg cm\(^{-3}\) K\(^{-1}\)], \( \kappa \) is the thermal
conductivity, and \( L_r \) is the “local luminosity” defined as
the non-neutrino heat flux transported through a sphere of
radius \( r \). The gravitational mass \( m(r) \) and the metric func-
tion \( \Phi(r) \) are determined by the stellar model. After thermal
relaxation, the redshifted temperature \( \tilde{T}(t) \equiv T(r,t) e^{\Phi(r)} \)
becomes constant throughout the interior.

It is conventional (Gudmundsson et al. 1983) to subdi-
vide the calculation of heat transport in the neutron star
interior \( r < R_b \) and in the outer heat-blanketing envelope
\( R_b \leq r \leq R \), where \( R \) is the stellar radius, and the bound-
ary radius \( R_b \) corresponds to the density \( \rho_b = 10^{10} \) g cm\(^{-3}\)
(\( \sim 100 \) meters under the surface). The thermal structure
of the blanketing envelope is studied separately in the sta-
nary, plane-parallel approximation to relate the effective
surface temperature \( T_s \) to the temperature \( T_b \) at the inner
boundary of the envelope. We use the \( T_s - T_b \) relation
obtained by Potekhin et al. (1997) for the envelope composed
mostly of iron.

The effective temperature determines the photon lu-
iminosity: \( L_\gamma = L_r(r,t) = 4\pi R^2 T_\gamma^4(t) \). Both \( L_\gamma \) and \( T_\gamma \)
refer to the locally-flat reference frame on the surface.
A distant observer would register the “apparent” luminosity

\[
L^\infty = L_s (1 - r_g/R) \text{ and the “apparent” effective tempera-
ture } \tilde{T}^\infty = T_s \sqrt{1 - r_g/R} \text{, where } r_g = 2GM/c^2 \text{ is the grav-
itational radius. Typically, } r_g/R = 30 - 40\%.
\]

We have developed a new evolutionary code, based on
the Henyey-type scheme on a grid of spherical shells (Kip-
penhahn et al. 1967). The hydrostatic model of the neutron
star with a given equation of state is calculated separately
and is fixed throughout the calculation. In the initial con-
figuration the star has a constant redshifted temperature
throughout the interior, \( \tilde{T} = 10^{10} \) K, and no heat flux, \( L_r = 0 \). Also, to improve numerical convergence the thermal con-
ductivity in the core is boosted for the initial epoch \( t < 10^{-2} \) yr.
Since the crust is thermally detached from the core at
such small age, this correction has no effect on the cooling
curves. Full details of the new code are available on the in-
ternet (http://www.ast.cam.ac.uk/~ognedin/ns/ns.html).

2.2 Physics input

We use the equation of state of Negele & Vautherin (1973)
in the stellar crust with the smooth composition model of
ground-state matter to describe the properties of atomic nu-
clei (Kaminker et al. 1999). We assume that the nuclei are
spherical throughout the entire crust. The core–crust inter-
fase is placed at \( \rho_c = 1.5 \times 10^{14} \) g cm\(^{-3}\). For simplicity, we
consider the neutron star cores composed of neutrons \( n \),
protons \( p \) and electrons \( e \) and use the moderately stiff
phenomenological equation of state of Prakash et al. (1988)
in the simplified version proposed by Page & Applegate
1992), in agreement with our previous work (Yakovlev et al.
1999 and references therein).

The parameters of the models are summarized in Table
1. It shows the stellar masses, radii, central densities, crust
masses, and crust thicknesses for a number of models. The
maximum mass of the stable neutron stars with this equa-
tion of state is 1.73 \( M_\odot \). We define the crust thickness as
\( \Delta R_{\text{crust}} = R - R_{\text{core}} \), while the proper geometrical thickness
(for \( \Delta R_{\text{crust}} < R \)) is \( \Delta R_{\text{crust}} \sqrt{1 - r_g/R} \). As we increase
the mass \( M \), the radii and crust masses of the stable config-
urations get smaller, i.e., the stars become more compact.

Free neutrons in the inner crust and both neutrons and
protons in the core of a neutron star are likely to be superfluid. We assume the singlet-state pairing of the protons in the core. Superfluidity of free neutrons in the crust and of neutrons in the outermost part of the core is known to be of the singlet-state type, but at higher densities it switches to the triplet-state type. Various microscopic theories predict a large scatter of the critical temperatures of the neutron and proton superfluids, $T_{cn}$ and $T_{cp}$, depending on the nucleon-nucleon potential model and the many-body theory employed (see Yakovlev et al. 1999 for references).

As an example we will use two models, a weak and a strong superfluidity (Fig. 1). The model of strong superfluidity corresponds to the higher $T_c$. It is based on the rather large energy gaps calculated by Elgoyev et al. (1996) for the singlet-state pairing (with the maximum gap of about 2.5 MeV) and by Hoffberg et al. (1970) for the triplet-state pairing. The weak superfluid model makes use of the smaller gaps derived by Wambach et al. (1993) (with the maximum gap of about 1 MeV) for the singlet-state superfluid and by Amundsen & Østgaard (1985) for the triplet-state neutron superfluid. For simplicity, we use the same function $T_c(n)$ to describe the singlet pairing of free neutrons in the crust ($n = n_n$) and of the protons in the core ($n = n_p$).

The cooling of neutron stars is mainly determined by the neutrino emissivity, specific heat capacity, and thermal conductivity. We include all relevant sources of neutrino emission (Yakovlev et al. 1999, 2001): the direct and modified Urca processes, $nn$, $pp$ and $np$ bremsstrahlung in the core; and plasmon decay, $e^-e^+$ pair annihilation, electron-nucleus ($eZ$) and $nn$ bremsstrahlung in the crust. The emissivity of the proton branch of the modified Urca process derived by Yakovlev & Levenfish (1995) has been corrected by Yakovlev et al. (2001), which has almost no effect on the cooling curves. We include the proper reduction of the neutrino reactions by superfluidity and also an additional neutrino emission due to Cooper pairing of superfluid nucleons. The effective nucleon masses in the core and the crust are set equal to 0.7 of their bare masses.

A very important cooling effect is produced by the pow-

---

**Figure 1.** Density dependence of the critical temperatures of superfluidity (SF) of free neutrons in the inner crust, and neutrons and protons in the core for the strong (solid lines) and weak (dashed lines) superfluid models (see text for details).

**Figure 2.** Density dependence of the neutrino emissivity at $T = 10^9$ K. Solid lines: partial emissivities due to electron-nucleus bremsstrahlung (e-brems) and plasmon decay (plasma) in the crust and the total emissivity produced by direct and modified Urca processes and by nucleon-nucleon (NN) bremsstrahlung in the nonsuperfluid core. Dashed and dash-and-dot lines: the emissivity due to Cooper pairing of neutrons for the models of strong and weak superfluidity in the crust. Vertical dotted lines indicate the neutron drip density and the boundary of the core.

**Figure 3.** Density dependence of the specific heat capacity at $T = 10^9$ K. Solid lines: partial heat capacities of ions ($i$), electrons ($e$) and free neutrons ($n$) in nonsuperfluid crusts, as well as of neutrons, protons ($p$) and electrons in nonsuperfluid cores. Dashed lines: heat capacities of free neutrons in the crust modified by weak or strong superfluidity.
erful direct Urca process (Lattimer et al. 1991). For our equation of state, this process is allowed at densities above the threshold, \( \rho_{\text{crit}} = 1.298 \times 10^{15} \text{ g cm}^{-3} \). If the central density of the model exceeds the threshold, \( \rho > \rho_{\text{crit}} \), the stellar core has a central kernel where the direct Urca process leads to fast cooling. The masses and radii of these kernels, \( M_D \) and \( R_D \), are given in Table 1. The mass of the central kernel increases rapidly with \( M \). In addition, we show the threshold configuration with \( \rho_c = \rho_{\text{crit}} \). It separates the low-mass models, where the direct Urca process is forbidden, from the high-mass models, where the direct Urca is allowed.

For illustration, Fig. 2 shows the emissivity of various neutrino processes versus density at \( T = 10^9 \text{ K} \). The plasmon decay is a powerful neutrino emission mechanism in a hot crust at not very high densities, but it fades away quickly when the temperature decreases below \( 10^9 \text{ K} \). The electron-nucleus bremsstrahlung is efficient throughout the entire crust. In superfluid crusts, the Cooper pairing neutrino emission switches on at temperatures \( T = T_{\text{en}} \), reaches maximum at \( T \) slightly below \( T_{\text{en}} \), and fades away exponentially for \( T < T_{\text{en}} \). In the neutron star core, a large jump of the neutrino emissivity at \( \rho = \rho_{\text{en}} \) is associated with the direct Urca process.

The heat capacity is contributed by neutrons, protons and electrons in the core; and by electrons, free neutrons, and atomic nuclei (vibrations of ions in Coulomb lattice) in the crust. The superfluid effects on the heat capacity of nucleons in the core and of free neutrons in the crust are incorporated according to Levenfish & Yakovlev (1994). In the absence of superfluidity, the neutrons would have the dominant contribution in the core and the inner crust. The effects of neutron superfluidity are illustrated on Fig. 3 in the case of crust superfluidity. When \( T \) falls only slightly below \( T_{\text{en}} \), the superfluidity increases the neutron heat capacity due to the latent heat released at the phase transition. However, for \( T \ll T_{\text{en}} \) the superfluidity exponentially reduces the heat capacity. If the stellar crust was sufficiently colder than in Fig. 3 any neutron superfluid would reduce the contribution of free neutrons to negligibly small values, and the heat capacity would be determined by the ions and electrons. These effects are analogous in the core. The heat capacity in the superfluid core at \( T \ll T_{\text{en}} \) and \( T \ll T_{\text{cp}} \) would be determined by the electrons.

The thermal conductivity in the core is taken as a sum of the conductivities of the electrons (Gnedin & Yakovlev 1995) and neutrons (Baiko et al. 2000). The electron contribution usually dominates. The conductivity in the crust is assumed to be due to the electron scattering off atomic nuclei. In the outer crust the finite size of the proton charge distribution within a nucleus can be neglected. We use the recent results of Potekhin et al. (1999), which include the multiphonon processes in electron–nucleus scattering in Coulomb solid and incipient long-range nucleus–nucleus correlations in strongly coupled Coulomb liquid of atomic nuclei. In the inner crust, we have performed original calculations using the same formalism but taking into account the finite size of the proton core of atomic nuclei (finite-size nuclei). Details of the calculations and the numerical fits are given in the Appendix.

The results are illustrated in Fig. 4, which shows the density profiles of the electron thermal conductivity for \( T = 10^7, 10^8 \) and \( 10^9 \text{ K} \). The conductivity in the core is several orders of magnitude higher than in the crust since there are no such efficient electron scatterers in the core as atomic nuclei. In the crust we plot the thermal conductivity for finite-size nuclei (solid lines) and for point-like nuclei (dashed lines). The finite-size effects are negligible near the neutron drip point but increase the conductivity at the crust base (\( \rho \sim 10^{14} \text{ g cm}^{-3} \)) by a factor of \( 3 \sim 5 \).

Another contribution to the thermal conductivity in the crust may come from the scattering of electrons off charged impurities, the randomly distributed nuclei of different charge number. However, the most important temperature interval for neutron star cooling is \( \sim 10^8 \sim 10^9 \text{ K} \). We have verified that the thermal conductivity is almost unaffected by the impurities in this temperature interval (for a not very impure matter, \( Q_{\text{imp}} \lesssim 1 \), where \( Q_{\text{imp}} \) is the impurity parameter defined in the Appendix), although they can noticeably decrease the conductivity at lower \( T \). Therefore, we neglect the effects of impurities in the present calculations.

3 THERMAL RELAXATION IN NON-SUPERFLUID NEUTRON STARS

First, consider thermal relaxation in a young non-superfluid neutron star. The main features of the process are known from previous cooling simulations (Page & Applegate 1992; Lattimer et al. 1994; Page 1998a, 1998b; and references therein).

Fig. 3 shows cooling models of neutron stars with different masses. In the low-mass models, \( M < 1.44 M_\odot \), the direct Urca process is forbidden. These stars follow the standard cooling scenario and the cooling curves are almost independent of \( M \). The high-mass models go through the fast cooling scenario and demonstrate a spectacular drop of the surface temperature at the end of the thermal relaxation

![Figure 4. Density dependence of the electron thermal conductivity at \( T = 10^7, 10^8 \) and \( 10^9 \text{ K} \) in the core and the crust.](image)
Figure 5. Cooling curves for non-superfluid neutron star models with 1.1, 1.2 \ldots, 1.7 \, M_{\odot}.

Figure 6. Temperature profiles in the 1.3 \, M_{\odot} neutron star without superfluid effects. Numbers next to curves show stellar age. Contours are at 0, 10^{-5}, 10^{-4}, 10^{-3}, 10^{-2}, 10^{-1}, 1, 2, 5, 10, 20, 50, 100, and 1000 yr.

Figure 7. Temperature profiles in the 1.7 \, M_{\odot} neutron star without superfluid effects. Contours are at 0, 10^{-5}, 10^{-4}, 10^{-3}, 10^{-2}, 10^{-1}, 1, 2, 5, 10, 20, and 50 yr.

Figure 8. Ratio of the integrated heat capacities in the crust and the core for the neutron star models of masses 1.1, 1.2, \ldots, 1.7 \, M_{\odot}.
Figure 9. Ratio of the neutrino luminosities in the crust and the core for the neutron star models of masses 1.1, 1.2, ..., 1.7 $M_\odot$.

epoch, $t \sim 50$ yr, due to the emergence of the cooling wave on the surface. The same, although much less pronounced, effect takes place in the case of slow cooling.

Since the neutrino emissivity of the direct Urca process is several orders of magnitude larger than that of the modified Urca processes, the fast cooling regime is established even if the central kernel, where the direct Urca process is allowed, occupies a small fraction of the stellar core (Page & Applegate 1992). For the high-mass stars ($M > 1.44 M_\odot$) the cooling curves, again, depend weakly on the mass. The change of the slope of the cooling curves at $t \sim 10^5\sim 10^6$ yr manifests the transition from the neutrino to the photon cooling stage.

The surface temperature at the initial cooling stage (the first 100 years) is rather independent of the equation of state, stellar mass, or the core neutrino luminosity. The surface temperature is mainly determined by the physical properties of matter in the crust. The core and the crust are thermally decoupled, and the effective surface temperature does not reflect the thermal state of the stellar core.

In contrast, the evolution of the central temperature, $T(0, t)$, is drastically different for the slow and fast cooling scenarios. In all low-mass models, $T(0, t) \propto t^{-1/6}$ throughout the neutrino cooling era, $t \lesssim 10^3$ yr, with a small offset in normalization. This follows from the simple power-law temperature dependence of the heat capacity ($C_v \propto T$) and the standard neutrino emissivity ($Q_\nu \propto T^3$). In the models with fast cooling, where the dominant neutrino process is the direct Urca ($Q_\nu \propto T^8$), the scaling relation is $T(0, t) \propto t^{-1/4}$ for the initial period $t < 10^{-2}$ yr. But then until $t \lesssim 10$ yr, the central temperature remains almost constant at $10^8$ K as the heat flows from the warmer outer core, in which direct Urca process is prohibited, into the inner core. During the thermal relaxation epoch, $10^2 < t < 10^3$ yr, the central temperature falls again by a factor of several. After the thermalization until the end of the neutrino era, once again, $T(0, t) \propto t^{-1/4}$.

Figs. 9 and 9 illustrate the effects of thermal relaxation on the internal temperature profiles in the slow and fast cooling scenarios, respectively. Until the age of about 1 yr, the neutron star core, the inner and the outer crusts form almost independent thermal reservoirs. The region around $4 \times 10^{11}$ g cm$^{-3}$, where free neutrons appear in the crust, seems to be the most effective at cooling owing to the powerful neutrino emission (see below). The outer crust cools to $10^{10}$ K in less than a month, while the inner parts remain much hotter. The core also cools independently but is unable to affect the inner crust due to the slow thermal conduction. During the first years the central kernel of the 1.7 $M_\odot$ model in Fig. 9 remains much colder than the outer core. This is because the kernel is cooled by the powerful direct Urca process and thermal conduction is still unable to establish thermal relaxation throughout the core. Almost full core relaxation is achieved in 10 years.

After the first year, the crust temperature profiles of the slow and fast cooling scenarios start to differ. In the former, the temperature gradient between the core and the crust is slowly eroded, as the cooling wave from the center reaches the surface. In the latter, the temperature gradient continues to grow until it reaches a maximum at $t \sim 10$ yr. Then a huge amount of heat releases from the crust and leads to a spectacular drop of the surface temperature by an order of magnitude (which lowers the photon luminosity by four orders of magnitude). At $t = 50$ yr, the entire star is already isothermal. Note, that despite larger temperature gradients, thermal relaxation proceeds overall quicker in the fast cooling scenario.

Prior to thermal relaxation, the contributions of the neutron star crust to the integrated heat capacity and neutrino luminosity are significant (Figs. 8 and 8). In the slow-cooling models, the heat capacity in the crust ranges from 10% to 20% of that in the core, with the larger fraction in the low-mass models (where crusts occupy larger fraction of the volume). In the fast-cooling models, the ratio of the crust to core heat capacities reaches a maximum of 55% at $t \sim 10$ yr before dropping to under 10% after the relaxation. Similarly, the integrated neutrino luminosity of the crust is about 15%–40% of that of the core at $t \sim 1$ yr, and then it drops to a tiny fraction at later times.

The importance of the individual neutrino mechanisms for the crust cooling varies at different epochs. First, for $t < 10^{-2}$ yr in the fast cooling scenario or for $t < 3 \times 10^{-3}$ yr in the slow cooling scenario, the $e^- e^+$ pair emission controls the crust temperature. As the temperature drops below $5 \times 10^8$ K, this process quickly fades away. The next epoch is controlled by plasmon decay. It dominates for $10^{-2} < t \lesssim 10$ yr (fast cooling) or $3 \times 10^{-3} < t \lesssim 10$ yr (slow cooling). The last epoch of thermal relaxation lasts for the period $10 \lesssim t < 100$ yr (fast cooling) or $10 \lesssim t < 1000$ yr (slow cooling), where either electron-nucleus or neutron-neutron bremsstrahlung is important. In fact, both neutrino processes give almost identical cooling curves in the absence of superfluidity. However, free neutrons in the crust are thought to be in a superfluid state which strongly
suppresses \( n \) bremsstrahlung. Therefore, electron-nucleus bremsstrahlung is likely to be the dominant neutrino mechanism in this last epoch.

4 RELAXATION TIME

The duration of the thermal relaxation epoch is potentially interesting from the observational point of view. This problem has been studied in a number of papers, with the most detailed and thorough work by Lattimer et al. (1994; also see references therein). Those authors considered thermal relaxation for the fast cooling and defined the relaxation time \( t_w \) as the moment of the most negative slope of the cooling curve, \( \ln T_c(\ln t) \), of a young neutron star. This is a typical time for the cooling wave to reach the surface. We will mainly focus on the case of rapid cooling, where the relaxation effects are more pronounced.

According to Lattimer et al. (1994), the relaxation time of rapidly cooling neutron stars of various masses is determined mainly by the crust thickness \( \Delta R_{\text{crust}} \) and is given by a simple scaling relation

\[
t_w \approx \alpha t_1, \quad \alpha \equiv \left( \frac{\Delta R_{\text{crust}}}{1 \text{ km}} \right)^2 (1 - r_g/R)^{-3/2}.
\]  

(3)

Here, \( t_1 \) is the normalized relaxation time which depends solely on the microscopic properties of matter, such as the thermal conductivity and heat capacity. In superfluid neutron stars, \( t_1 \) is sensitive to the magnitude and density dependence of the critical temperature of neutron superfluidity in the crust, as we discuss later. It is important that \( t_1 \) appears to be almost independent of the neutron star model, its mass \( M \) and radius \( R \). We have verified that this scaling holds also for the slow-cooling non-superfluid models.

For the non-superfluid stars with the core–crust interface placed at \( \rho_0 / 2 \), which is close to our value \( \rho_{cc} = 1.5 \times 10^{14} \text{ g cm}^{-3} \), Lattimer et al. (1994) obtained \( t_1 \approx 26 \) yr. Our rapidly cooling models show similar scaling, \( t_1 = 28.4 \pm 0.2 \) yr, whereas in the slowly cooling models, \( t_1 = 32.9 \pm 1.2 \) yr.

The dependence of \( t_w \) on the thermal conductivity \( \kappa \) and heat capacity \( C_v \) follows from a simple estimate of the thermal relaxation time in a uniform slab of width \( l \):

\[
t_w \sim C_v l^2 / \kappa.
\]  

(4)

The proper width of a thin crust (\( \Delta R_{\text{crust}} \ll R \)), taking into account the effects of general relativity, is \( l = \Delta R_{\text{crust}}/\sqrt{1 - r_g/R} \). This gives \( t_w \propto 1/(1 - r_g/R) \) in equation (3). An additional factor \( 1/\sqrt{1 - r_g/R} \) accounts for the gravitational dilution of time intervals.

The parameter \( t_1 \) can be roughly estimated using equation (3) with the values of the specific heat capacity and thermal conductivity in the crust near the crust–core interface, for instance, at \( \rho = 10^{14} \text{ g cm}^{-3} \). Fig. 5 shows that the relevant temperature at the interface is \( T \approx 2 \times 10^{8} \text{ K} \). At those temperature and density according to our physics input, \( C_v = 1.9 \times 10^{20} \text{ erg cm}^{-3} \text{ K}^{-1} \) and \( \kappa = 1.8 \times 10^{20} \text{ erg cm}^{-3} \text{ s}^{-1} \text{ K}^{-1} \). This gives \( t_1 \approx (1 \text{ km})^2 C_v/\kappa \approx 34 \) yr, in qualitative agreement with the model value.

Figure 4 shows the sensitivity of the relaxation time to test variations of physical properties of the neutron star crust in the fast cooling scenario of the 1.5 \( M_\odot \) star. The corresponding values of \( t_w \) and \( t_1 \) are listed in Table 2. They may differ from the mean values (Table 3) within the errorbars. We present also the results for the superfluid models discussed in the next section.

Switching off the neutrino emission from the crust, while keeping the heat capacity, slows down the thermalization epoch by almost an order of magnitude, from 50 to 250

| test crust model                             | \( t_w \) (yr) | \( t_1 \) (yr) | \( t_1^b \) (yr) |
|---------------------------------------------|----------------|----------------|------------------|
| real model, no SF                           | 52.4           | 28.8           | 33.9             |
| no crust neutrinos                          | 253.5          | 139.2          | 134.9            |
| only plasmon decay neutrinos                | 67.6           | 37.1           | 41.6             |
| only \( eZ \) neutrino bremsstr.            | 58.3           | 32.0           | 34.5             |
| no neutron heat capacity                    | 15.3           | 8.4            | 6.7              |
| cond. for point-like nuclei                 | 131.8          | 72.4           | 102.3            |
| real model, weak crust SF                   | 20.2           | 11.1           | 3.3              |
| no Cooper neutrinos                         | 29.0           | 15.9           | 19.0             |
| weak core+crust SF                          | 22.3           | 12.2           | 25.7             |
| real model, strong crust SF                 | 15.0           | 8.2            | 6.7              |
| no Cooper neutrinos                         | 15.5           | 8.5            | 6.9              |
| strong core+crust SF                        | 10.7           | 5.9            | 5.8              |

\( ^a \) for the 1.5 \( M_\odot \) model, with \( \alpha = 1.821 \)
\( ^b \) for the 1.3 \( M_\odot \) model, with \( \alpha = 2.875 \)
Turning on the plasmon decay alone would lead to the cooling curve not very different from the real one, although the thermalization epoch would be delayed by about 30% ($t_w = 68$ yr). Turning off the plasmon decay and switching on the electron-nucleus bremsstrahlung instead would give a hotter neutron star before the relaxation but almost correct relaxation time, $t_w = 58$ yr. Combined together, plasmon decay and electron-nucleus bremsstrahlung would reproduce accurately the real cooling curve.

Restoring the neutrino emissivity in the crust but switching off the heat capacity of crustal neutrons ($C_v = 0$) we obtain much faster relaxation, which lasts about 15 years. This numerical experiment imitates the suppression of the neutron heat capacity by strong neutron superfluidity (discussed in Sect. 3). On the other hand, had we neglected the quantum suppression of the heat capacity of the nuclei below the Debye temperature, the latter heat capacity would have become important in older neutron stars, $t \gtrsim 10^4$ yr, strongly delaying the cooling.

Finally, the relaxation time depends on the thermal conductivity in the inner crust. For instance, a neglect of the strong neutron superfluidity strips” of opacity, thermal conductivity, or neutrino emissivity have been calculated so far for the phase of non-spherical nuclei. Note that the thermal conductivity has not been calculated for the electron-nucleus bremsstrahlung instead would give a hotter neutron star before the relaxation but almost correct relaxation time, $t_w = 58$ yr. Combined together, plasmon decay and electron-nucleus bremsstrahlung would reproduce accurately the real cooling curve.

We have run additional cooling models with $M = 1.5 M_\odot$, varying the heat capacity and thermal conductivity within the crust (at $\rho = \rho_0$) by a fixed factor 1/8, 1/4, 1/2, 2, 4, 8. The results confirm that $t_w$ and $t_1$ are indeed quite sensitive to the variations of $C_v$ and $\kappa$, in agreement with the qualitative estimate, equation (3), and the results of Lattimer et al. (1994). It is important that these variations do not invalidate the scaling relation for the relaxation time, equation (3). Moreover, if $C_v$ is increased and $\kappa$ is decreased, the dependence of $t_1$ on the values of the heat capacity and thermal conductivity is described by a simple scaling relation $t_1 \propto C_v/\kappa^{0.8}$.

Figure 11. Sensitivity of the crust relaxation time to the variations of the heat capacity, thermal conductivity, and neutrino emissivity in various density regions of the 1.5 $M_\odot$ neutron star model with the non-superfluid core. In each 0.2 dex of log $\rho$, $Q_v$, and $\kappa$ are reduced by a factor of 8 (solid lines, no superfluidity), and $C_v$ is enhanced by a factor of 8 without superfluid effects (dashes) and with weak superfluidity (dots).

We have also done sensitivity tests of the relaxation time, analogous to those performed by Epstein et al. (1983). In each density region, 0.2 dex of log $\rho$, either the heat capacity, thermal conductivity, or neutrino emissivity have been changed by a factor 1/8, 1/4, 1/2, 2, 4, 8. Figure 11 shows the variations of $t_w$ when $Q_v$ and $\kappa$ are reduced by a factor of 8, and $C_v$ is enhanced by a factor of 8. The relaxation time depends mostly on the values of $C_v$ and $\kappa$ in the crust in the density range $10^{13} - 1.5 \times 10^{14}$ g cm$^{-3}$ near the crust-core interface, being rather insensitive to the variations of $Q_v$. Variations of the physical parameters in the core affect the crustal relaxation much less strongly (at least for non-superfluid models). The density range where $t_w$ is most affected by the variations of $\kappa$ is narrower than the density range where it is affected by the variations of $C_v$. The most important temperature range which influences $t_w$ is $10^8 \lesssim T \lesssim 10^9$ K. The properties of matter in these “sensitivity strips” of $\rho$ and $T$ are very model-dependent. For instance, the nuclei may be strongly non-spherical (rods, plates, etc.) at $\rho > 10^{14}$ g cm$^{-3}$, which is not included in our calculations. Note that the thermal conductivity has not been calculated so far for the phase of non-spherical nuclei.

However, if $C_v$ were noticeably lower or $\kappa$ noticeably higher than in our basic non-superfluid models, the decrease of the crust relaxation time would saturate at $t_w \approx 13$ yr. This is the time it takes the inner core with the direct Urca emission to equilibrate thermally with the outer core (cf Fig. 2). More generally, this is a core relaxation time $t_{core}$, which is almost independent of the parameters of the crust. It can be estimated using the same formalism of heat diffusion, equation (5), through a slab of material between the direct-Urca-allowed kernel and the boundary of the core, with $l = R_{core} - R_D$. The relativistic factors appear as $(e^{-3\Phi_{cc}})^3$, where $\Phi_{cc}$ is the metric function at the core-crust interface.

We have calculated the core relaxation time by setting the crust conductivity very high for the fast cooling models of non-superfluid neutron stars with $M = 1.5, 1.6, 1.7 M_\odot$. The least-squares fit gives

$$t_{core} = t_2 \left( \frac{R_{core} - R_D}{10 \text{ km}} \right)^2 e^{-3\Phi_{cc}},$$

with $t_2 = (9.1 \pm 0.8)$ yr. Thus, the fast cooling models may have two distinct relaxation times, in the core and in the crust, and the latter is typically longer, at least for non-superfluid models.

The normalized crustal relaxation time $t_1$ shows very small variations with the neutron star mass, $\pm 0.2$ yr for fast cooling and $\pm 1.2$ yr for slow cooling (see Table 5). Even combining the models with all masses we find that the scal-
ing relation, eq. (3), holds remarkably well, $t_1(\text{all}) = 31 \pm 2$ yr. Therefore, the cooling of the crust prior to thermal relaxation is indeed insensitive to the thermal evolution of the core. The slow cooling makes the relaxation much less pronounced but it does not change its basic features. This property owes to the very high thermal conductivity in the core (Fig. 4), which makes the core relaxation time typically much shorter than the crust one. Accordingly, the cooling curves do not depend on the exact values of the conductivity in the core as long as it is much higher than in the crust.

5 EFFECTS OF SUPERFLUIDITY

Free neutrons in the crust and both neutrons and protons in the core of a neutron star are likely to be in the superfluid state. We implement the effects of superfluidity as discussed in Sect. 2.2. It turns out that superfluidity in the crust affects the cooling curves at the thermal relaxation stage, while superfluidity in the core affects cooling at later stages. First, consider the effects of neutron superfluidity in the crust. If the temperature $T$ falls much below the critical temperature $T_{cn}$, the superfluidity reduces strongly the neutron heat capacity and $\nu\nu$ neutrino bremsstrahlung. While the latter is compensated by electron-nucleus bremsstrahlung, the former effect leads to a faster thermal relaxation. In addition, a new neutrino process is allowed in the superfluid state, the neutrino emission due to Cooper pairing of free neutrons. It further accelerates cooling and thermal relaxation of the crust.

Figure 12 illustrates how the weak neutron superfluidity

![Figure 13. Temperature profiles (solid lines) in the interior of the 1.5 $M_\odot$ model with strong superfluidity both in the crust and the core. Numbers next to curves show the stellar age. The contours are at 0, $10^{-4}$, $10^{-2}$, 1, and 20 yr. Dotted lines show the temperature profiles of the non-superfluid star.

![Figure 14. Superfluid effects in the crust of the 1.5 $M_\odot$ neutron star model with the non-superfluid core. Dashed line is for the case of weak neutron superfluidity, while dash-and-dot line is for the case of strong superfluidity. Dotted line is obtained neglecting the Cooper-pair neutrino emission for the weak-superfluid model. Solid line is the cooling curve for the non-superfluid crust.]
in the crust carves out the temperature profiles in the standard cooling scenario, \( M = 1.3 M_\odot \). For the first \( 10^{-4} \) yr, when the temperature is above \( T_{cn} \), the thermal structure is identical to that of the non-superfluid model. Later, the region where the critical temperatures are the highest cools much faster than the neighboring layers. The acceleration of cooling is mainly due to the reduction of the heat capacity and switching on the Cooper-pairing neutrino emission. A sequence of points in Fig. 12 in which solid lines start to deviate from the dotted ones reproduces the density profile of \( T_{cn} \) shown in Fig. 1. As the temperature falls further, wider density regions become affected, producing shells of cool matter surrounded by hotter layers on both sides. After the cooling wave from the core reaches the outer crust, the star settles into almost the same isothermal state as the non-superfluid model, but faster.

The effects are much stronger if the superfluidity is allowed for in the stellar core. Figure 4 shows the combined effect of the strong superfluidity in the core and crust. Both neutrons and protons are superfluid in the core of this \( 1.5 M_\odot \) neutron star. In addition to the trough in the crust layers, the core develops a complex structure. All the sources of neutrino emission and the nucleon heat capacity in the core are affected by the superfluidity, while the electron heat capacity is not and becomes dominant in highly superfluid cores. As soon as the temperature drops significantly below the critical temperatures \( T_{cn} \) or \( T_{cp} \), Cooper-pairing neutrino emission comes into play. It starts in the inner part of the core and drives the temperature down. At \( t \approx 10^{-2} \) yr that region is even cooler than in the non-superfluid model, while the other parts of the core are slightly hotter. By the age of 1 yr, this cool region includes all of the core except for the inner kernel. As a result, thermal relaxation proceeds on a shorter timescale, and at \( t \lesssim 20 \) yr the stellar interior is isothermal. Thus, for the large assumed values of \( T_{cn} \) and \( T_{cp} \), the neutrino emission due to Cooper pairing becomes so strong that, instead of slowing down, the presence of the core superfluidity accelerates the cooling.

Table 3. Normalized relaxation times in the crust \( t_1 \), eq. (4), and in the core \( t_2 \), eq. (5).

|                | fast cooling | slow cooling |
|----------------|--------------|--------------|
| crust \( t_1 \) yr | 28.4 ± 0.2  | 32.9 ± 1.2   |
| weak crust SF   | 10.3 ± 0.2  | 3.4 ± 0.3   |
| weak core+crust SF | 11.5 ± 0.2  | 25.0 ± 0.4  |
| strong crust SF | 7.0 ± 0.9   | 6.8 ± 1.0   |
| strong core+crust SF | 6.2 ± 0.4  | 5.7 ± 0.2   |
| core \( t_2 \) yr |             |              |
| no SF           | 9.1 ± 0.8   |              |
| weak core SF    | 11.2 ± 0.5  |              |
| strong core SF  | 3.6 ± 0.2   |              |

The large difference of the values of \( t_1 \) for the slowly and rapidly cooling stars with weakly-superfluid crusts indicates that, generally, superfluidity may violate the simple scaling relations for the relaxation time as well as the strict thermal decoupling of the crust and the core prior to the relaxation. This happens because superfluidity makes the heat capacity (and other properties of matter) a strong function of density.

The shortening of the thermal relaxation phase in a rapidly cooling star, due to the superfluid reduction of the crustal heat capacity, was emphasized by Lattimer et al. (1994). Using a model of superfluidity, which is closer to our model of strong superfluidity in the crust, Lattimer et al. find that the relaxation time becomes three times as short, \( t_r = 8.4 ± 2.0 \) yr, in qualitative agreement with our results. Our calculations indicate that the effect is sensitive to the model of neutron superfluidity in the crust and thus, it can be used to test such models.

When thermal relaxation is over, the cooling is mainly regulated by the neutrino luminosity and heat capacity of the core. The processes in the crust cease to play a significant role except for the very low-mass neutron stars with large crusts. We find, however, that for some superfluid models the neutrino luminosity of the crust may affect the cooling for a short period of time during the transition from the neutrino cooling era to the photon era. For instance, cooling of superfluid neutron stars after the thermal relaxation is discussed in detail by Yakovlev et al. (1999).

6 CONCLUSIONS

We have studied thermal evolution of young neutron stars using a new cooling code and updated physics input. The effective surface temperature of isolated neutron stars in the first 100 yr is determined by the properties of the crust. After the cooling wave from the core reaches the surface, the effective temperature drops from 250 eV to 50 eV or lower.
We confirm the conclusion of Lattimer et al. (1994) that the duration of the relaxation epoch in a fast cooling scenario depends sensitively on the heat capacity and thermal conductivity of the crust. The relaxation time scales with the size of the crust, eq. (3), and the normalization is $t_1 = 28$ yr without superfluid effects or $t_1 = 7$ yr with strong superfluid effects. We find that the same conclusion holds in the case of slow cooling, in which thermal relaxation is less pronounced.

We also investigate the effects of various neutrino emission mechanisms in the crust, superfluidity of free neutrons, heat capacity, and thermal conductivity on the relaxation time (as summarized in Table 2 and Fig. 11). The relaxation time is most sensitive to the variations of $\kappa$ and $C_v$ in the density range $10^{13} - 1.5 \times 10^{14}$ g cm$^{-3}$ near the crust-core interface.

Young cooling neutron stars can serve as astrophysical laboratories of matter at subnuclear densities. The sensitive dependence of the relaxation time on the microscopic properties of matter in the deep inner crust provides a possibility to study these properties by observing the emergence of the cooling wave. In order to realize this method, one needs to detect thermal emission from a very young neutron star in the range from 50 to 250 eV. Such stars have not been found so far, but may be detected in young supernova remnants by Chandra and XMM.

To improve the cooling theory of young neutron stars, it would be desirable to continue theoretical studies of microscopic properties of matter in the inner crust, in particular, the possible unusual phases of nonspHERical nuclei at $\rho > 10^{14}$ g cm$^{-3}$, as well as the thermal conductivity, heat capacity, and neutrino emission of such matter. It is also important to refine our knowledge of matter containing spherical nuclei at $\rho < 10^{14}$ g cm$^{-3}$. In this paper we present (Appendix) and use the results of new calculations of the electron thermal conductivity due to the scattering of electrons off spherical atomic nuclei, taking into account the effects of finite sizes of proton nucleus cores.

Acknowledgements

We are grateful to P. Haensel and A. D. Kaminker for stimulating discussions and useful suggestions, and to D. Page for his suggestion to optimize the performance of the cooling code. This work was supported in part by RFBR (grant No. 99-02-18099), INTAS (grant No. 96-0542), NSF (grant No. PHY99-07949), and PPARC.

APPENDIX A: ELECTRON CONDUCTIVITY OF THE INNER CRUST

In the inner crust, the electrons are strongly degenerate. The heat and charge are transported mainly by the electrons with energies in a narrow thermal band near the Fermi level $\epsilon_F$. The coefficients of the electrical and thermal conductivities can be written as

$$\sigma = \frac{e^2 n_e \tau_e}{m_e^*}, \quad \kappa = \frac{\pi^2 k_B^2 T n_e \tau_e}{3 m_e^*},$$

(A1)

where $m_e^* \equiv \epsilon_F/e^2$, $n_e$ is the electron number density, $k_B$ is the Boltzmann constant, and $\tau_e$ and $\tau_n$ are the effective relaxation times. In the presence of several electron scattering mechanisms, the cumulative relaxation time is determined by the Matthiessen rule (e.g., Ziman 1960). In the envelopes of neutron stars, one traditionally considers the electron (ei), electron-electron, and electron-impurity scatterings. We focus on the main mechanism, the ei scattering.

A modern theoretical treatment of the conductivities due to the ei scattering in strongly coupled Coulomb plasmas of ions (atomic nuclei) has been proposed by Baiko et al. (1998) and applied to the outer envelopes of neutron stars by Potekhin et al. (1999). In the latter work, practical formulae for $\sigma$ and $\kappa$ incorporating all three scattering mechanisms have been derived; the corresponding Fortran code is available electronically at http://www.ioffe.rssi.ru/astro/condc.

For the inner crust these expressions must be modified. First, the atomic nuclei cannot be considered as point-like. Second, the ei scattering in the inner crust at temperatures much below $10^9$ K changes its character: so called umklapp processes cease to dominate and the normal processes (with electron momentum transfer within one Brillouin zone) become more important.

A1 Effects of finite size of ions

The effective ei relaxation time is

$$\tau_{ei} = \frac{3\pi\hbar}{4Z e^2 \alpha_e^2 \Lambda_{ei}(\epsilon_F)}$$

(A2)

where $\alpha_e \equiv e^2/(\hbar c)$ is the fine structure constant, and $\Lambda_{ei}$ is the Coulomb logarithm. For a classical plasma of ions in the Born approximation (e.g., Potekhin et al. 1999),

$$\Lambda_{ei} = \int_0^{2k_F} dq \, q^3 |\phi(q)|^2 S(q) \left[ 1 - \beta_e^2 \left( \frac{q}{2k_F} \right)^2 \right],$$

(A3)

$$\phi(q) \equiv \frac{F(q)}{q^2 - \epsilon(q)},$$

where $k_F = (3\pi^2 n_e)^{1/3}$ is the electron Fermi wave number, $\beta_e = v_F/c$, $v_F = h k_F/m_e^*$ is the Fermi velocity, $\epsilon(q)$ is the static dielectric function, $S(q)$ is the static structure factor of the ions (more precisely, its inelastic part -- see Baiko et al. 1998), and $F(q)$ is the form factor of the ions. The latter three functions allow for the electron polarization, ion-ion correlations, and the finite size of the ions, respectively.

The quantization of ionic motion becomes important at $T \ll T_p$, where $T_p = h \sqrt{4\pi n_e Z^2 e^2 / M / k_B}$ is the ion plasma temperature, $n_i = n_e / Z$ is the ion number density, $M = A m_i$ is the ion mass, and $m_i$ is the atomic mass unit. In this case, in equation (A3) one should use an effective structure factor, different for $\sigma$ and $\kappa$. These effective factors have been derived by Baiko et al. (1998) and fitted by Potekhin et al. (1999) for the case where the umklapp processes dominate.

A1.1 Uniform charge distribution in atomic nuclei

At $\rho \ll \rho_c$ a good approximation for the charge density within a nucleus is a step function. The corresponding form factor is

$$F(q) = \frac{3}{(q r_{nucl})^2} \left[ \sin(q r_{nucl}) - q r_{nucl} \cos(q r_{nucl}) \right],$$

(A4)

where $r_{nucl}$ is the proton core radius.
In general, $\Lambda_{ei}$ depends on $\rho$, $T$, $Z$, $A$, $r_{\text{nuc}}$, and it appears to be almost independent of the type of Coulomb lattice. It is convenient, however, to introduce the dimensionless parameters: the electron relativity parameter $x_e = h k_F / (m_e c)$; the ion size parameter $x_{\text{nuc}} = r_{\text{nuc}} / a_i$, where $a_i = (4\pi n_i / 3)^{-1/3}$ is the ion-sphere radius; the Coulomb screening parameter of ions $\Gamma = Z^2 e^2 / (a_i k_B T)$; the inverse ion quantum parameter $t_p \equiv T / T_p$; and the electron screening parameter $s_e = k_{\text{e}}^2 (2 k_{\text{e}}^2)^2 = a_i / (\pi \beta_e)$, where $k_{\text{e}}$ is the electron screening wave number. In accordance with the previous results, it is also convenient to introduce an auxiliary parameter $s = (2 k_{p} r_{p})^{-2}$, where $r_{p} = a_i / \sqrt{3} \pi$ is the Debye screening length for the ideal plasma of ions. Finally, the basic parameters that characterize the Coulomb lattice are the normalized first- and second-negative moments of the phonon frequency, $u_{-1} \approx 3$ and $u_{-2} \approx 13$.

We have calculated numerically the effective Coulomb logarithms $\Lambda_{ei}^{\text{eff}}$ for about a hundred pairs $(\rho, T)$, varying over orders of magnitude in the domain of strongly degenerate electrons and strongly coupled ions, for eight ion species from $Z = 12$ to $Z = 60$, and for five values of $x_{\text{nuc}}$ from 0 to 0.4. We have included the non-Born corrections in the same manner as Yakovlev (1987). The results can be fitted using the effective screening function

$$|\phi_{\text{eff}}(q)|^2 = \frac{1}{(\xi^2 + \eta^2)^2} \left[ 1 - e^{-s_{\phi} q^2} \right] e^{-s_{\gamma} q^2} G_{s,\gamma} D,$$  

instead of $|\phi(q)|^2 S(q)$ in equation (A3). Here, the first term corresponds to the Debye screening with the effective screening wave number $q_{\phi}$, the factor in square brackets corrects for the ion corrections, and the functions $G$ and $D$ describe ionic quantum effects, as in Potekhin et al. (1999). An additional factor $e^{-s_{\gamma} q^2}$ plays a role of the effective form factor. The numerical values of $\Lambda$ obtained from the accurate theory are reproduced (within several percent in the most important $\rho - T$ range) if we use the effective screening function (A4) with the following parameters:

$$s \equiv q_{\phi}^2 / (2 k_{F})^2 = (s_i + s_e) e^{-\beta_e}$$
$$\beta_e = \pi \alpha_e \beta_i$$
$$w \equiv (2 k_{F})^2 s_0 = (u_{-2} / s_0) (1 + \beta_e / 3)$$
$$w_1 \equiv (2 k_{F})^2 s_0 = 14.73 x_{\text{nuc}}^2 (1 + \sqrt{x_{\text{nuc}}} Z / 13) (1 + \beta_e / 3)$$
$$G_{s} = (1 + 0.0361 Z^{-3/4} / t_p^2)^{-1/2} (1 + 0.122 \beta_e^2 / 3)$$
$$G_{\gamma} = G_{s} + \frac{0.0105 t_p}{(t_p^2 + 0.0081)^{3/2}} \left[ 1 + \beta_i^2 \beta_e^2 \right] (1 - Z^{-1})$$
$$D = \exp \left[ -0.42 \sqrt{x_i / (AZ)} u_{-1} \exp(-9.1 t_p) \right]$$

where $x_{\text{nuc}} = 0$, these equations reduce to those presented in Potekhin et al. (1999).

With the effective screening function in the form (A4), we can integrate equation (A3) analytically and obtain

$$\Lambda_{ei}^{\text{eff}} = [\Lambda_0(s, w + w_1) - \Lambda_0(s, w_1)] G_{s,\gamma} D,$$  

where

$$\Lambda_0(s, w) = \Lambda_1(s, w) - \beta_e^2 \Lambda_2(s, w),$$
$$2 \Lambda_1(s, w) = \ln \frac{s + 1}{s} + \frac{s}{s + 1} (1 - e^{-w})$$

$$- (1 + sw) e^{-w} [E_1(sw) - E_1(sw + w)]$$
$$2 \Lambda_2(s, w) = \frac{e^{-w} - 1 + w}{w} - \frac{s^2}{s + 1} (1 - e^{-w}) - 2s \ln \frac{s + 1}{s}$$
$$+ s (2 + sw) e^{-w} [E_1(sw) - E_1(sw + w)],$$

and $E_1(x) = \int_{x}^{\infty} y^{-1} e^{-y} dy$ is the standard exponential integral. Note that using equations (A14) and (A15) directly may result in large numerical round-off errors in the limiting cases $s \ll 1$, $w \ll 1$, or $w \gg 1$. In these cases, explicit asymptotic expressions of Potekhin et al. (1999) should be used.

A1.2 Realistic charge distribution

Near the bottom of the crust, the boundary of the proton core in nuclei becomes fuzzy and the above results should be modified. Oyamatsu (1993) approximated the proton charge distribution by a function proportional to $(1 - r / r_m)^8$ (at $r < r_m$), where the power index $b$ controls the “sharpness” of the charge profile. The parameters $b$ and $r_m$ have been described by the simple functions of mass density in the smooth composition model of crust matter (Kaminker et al. 1999). Using that model we have calculated the electrical and thermal conductivities for $\rho = 10^9 - 10^{14.1}$ g cm$^{-3}$ and $T = 10^7 - 10^{10}$ K, and compared the results with the fitting formulae (A12) - (A13). We have found that the electrical conductivity is reproduced (within $\approx 5-10\%$), if we define $x_{\text{nuc}}$ in equation (A8) so as to reproduce the same mean-square proton-core radius as in the approximation of uniform charge distribution:

$$x_{\text{nuc}} = \frac{r_m}{a_i} \frac{1 - 15 / (5 + b) + 15 / (5 + 2b) - 5 / (5 + 3b)}{1 - 9 / (3 + b) + 9 / (3 + 2b) - 1 / (1 + b)}.$$  

In equation (A10) which determines the thermal conductivity at $T \ll T_p$, one should additionally multiply $x_{\text{nuc}}$ by a factor $b/(0.5 + b)$. The maximum fit error of $\kappa$ does not exceed 13%.

A2 Normal processes at very low temperatures

Near the boundaries of the Brillouin zones the dispersion relation of electrons differs from the free-electron case, and at the boundaries the electron energy spectrum contains gaps. The gaps $\Delta \epsilon$ can be estimated in the weak coupling approximation (e.g., Kittel 1986) as $\Delta \epsilon \sim \phi(k_F) = 4\pi Z e^2 n_i k_F^2 = 4 e^2 / (3\pi k_F)$. The effect of gaps is most significant if the deviation of the electron momentum from the intersection line between the Fermi surface and the Brillouin zone boundary does not exceed $\Delta k \sim \Delta \epsilon / (hv_F) \sim \frac{4}{\pi} (\alpha_e / \beta_i) k_F \ll k_F$. However, with decreasing temperature the strips of the Fermi surface, between which the umklapp processes proceed effectively, become narrower and closer to these intersection lines. When the widths of the strips, $\sim t_p (6\pi^2 n_i)^{1/3}$, become smaller than $\Delta k$, the umklapp processes are frozen out and the normal processes prevail. The above estimates indicate that this happens when the temperature falls below

$$T_u \sim T_p Z^{1/3} \alpha_e / (3\beta_i).$$

In this case, the formalism used in Sec. A1 becomes invalid. The asymptotes for the $\epsilon_i$ scattering rates at $T \ll T_u$ have been obtained by Raikh & Yakovlev (1982). In our notation, they yield:
Figure A1. Density dependence of the electron thermal conductivity at $T = 10^7$, $10^8$ and $10^9$ K in the neutron star crust. Solid and dot-dashed lines: ground state composition (Negele & Vautherin 1973), rectangular profile of nuclear charge (solid lines: pure crystal or liquid; dot-dashed line: including impurities; bottom of the hatched region: neglecting freezing of umklapp processes at $T = 10^9$ K); dashed lines: smooth-composition model, realistic profile of nuclear charge, no impurities.

\[
\begin{align*}
\Lambda_{ei}^{\sigma,\kappa} & = \frac{a_\zeta}{A^{1/2} Z} \times \left\{ \left( \frac{4}{3} \frac{\alpha_i}{\beta_i} \right) t_p^3 \right\}, \\
\end{align*}
\]

where

\[
a_\zeta = 180 \left( \frac{3}{\pi \eta m_u} \right)^{1/2} \zeta(5) \approx 50,
\]

and $\zeta(5) = 1.0369$ is the value of the Riemann zeta function.

Now we interpolate between the high-temperature Coulomb logarithm $\Lambda_{ei,\text{high}}^{\sigma,\kappa}$ given by equation (A12) and the low-temperature asymptote $\Lambda_{ei,\text{low}}^{\sigma,\kappa}$ given by equation (A18):

\[
\Lambda_{ei}^{\sigma,\kappa} = \Lambda_{ei,\text{high}}^{\sigma,\kappa} \exp(-T_u/T) + \Lambda_{ei,\text{low}}^{\sigma,\kappa} \left[ 1 - \exp(-T_u/T) \right].
\]

Figure A1 shows the thermal conductivities for three values of the temperature. In the neutron star envelope, the plasma is in the solid phase except for the segment to the left of the asterisk at the highest $T = 10^9$ K. At the lowest $T = 10^7$ K, the change in the scattering from umklapp to normal processes leads to a significant increase in the conductivity shown by the hatched region. However, this increase can be compensated by the scattering off impurities. For example, the dashed line shows the conductivity for $Q_{\text{imp}} \equiv (Z_{\text{imp}} - Z) n_{\text{imp}} / n_1 = 0.2$, where $n_{\text{imp}}$ and $Z_{\text{imp}}$ are the impurity number density and charge number, respectively. For comparison, dashed lines show the thermal conductivity in the smooth-composition model with smooth proton charge profile.

© RAS, MNRAS