Axions from cooling compact stars: pair-breaking processes

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Abstract

Once formed in a supernova explosion, a neutron star cools rapidly via neutrino emission during the first $10^4$-$10^5$ yr of its life-time. Here we compute the axion emission rate from baryonic components of a star at temperatures below their respective critical temperatures $T_c$ for normal-superfluid phase transition. The axion production is driven by a charge neutral weak process, associated with Cooper pair breaking and recombination. The requirement that the axion cooling does not overshadow the neutrino cooling puts a lower bound on the axion decay constant $f_a > 6 \times 10^9 T_{c9}^{-1}$ GeV, with $T_{c9} = T_c / 10^9$ K. This translates into an upper bound on the axion mass $m_a < 10^{-3} T_{c9}$ eV.

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1. Introduction

CP-violation in the strong sector of the Standard Model arises due to a topological interaction term in the QCD Lagrangian

$$\mathcal{L}_\theta = \frac{g^2}{32\pi^2} F_{\mu\nu}^a \tilde{F}_{\mu\nu}^a,$$

where $F_{\mu\nu}^a = \partial_\mu A_\nu - \partial_\nu A_\mu + gf^{abc} A_\mu^b A_\nu^c$ is the gluon field strength tensor, $g$ is the strong coupling constant, $\tilde{F}_{\mu\nu}^a = \epsilon_{\mu\nu\lambda\rho} F^{\lambda\rho a}/2$, $f^{abc}$ are the structure constants of SU(3) group, $\theta$ is the parameter which parametrizes the non-perturbative vacuum states of QCD $|\theta\rangle = \sum_n \exp(-in\theta)|n\rangle$, where $n$ is the winding number characterizing each distinct state, which is not connected to another by any gauge transformation. The QCD action changes by $2\pi$ under the shift $\theta \to \theta + 2\pi$, i.e., $\theta$ is a periodic function with a period of $2\pi$. In presence of quarks the physical parameter is not $\theta$, but

$$\bar{\theta} = \theta + \arg \det m_q,$$

where $m_q$ is the matrix of quark masses. Experimentally, the upper bound on the value of this parameter is $\bar{\theta} \lesssim 10^{-10}$, which is based on the measurements of the electric dipole moment of neutron $d_n < 6.3 \cdot 10^{-26} e$ cm. The smallness of $\bar{\theta}$ is the strong CP problem:

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the Standard Model does not provide any explanation on why this number should not be of order unity.

An elegant solution to the strong CP-problem is provided by the Peccei-Quinn mechanism [3–5]. This solution amounts to introducing a global $U(1)_{PQ}$ symmetry, which adds an additional anomaly term to the QCD action proportional to the axion field $a$. This term acts as a potential for the axion field and gives rise to an expectation value of the axion field $\langle a \rangle \sim -\theta$. The physical axion field is then $a - \langle a \rangle$, so that the undesirable $\theta$ term in the action is replaced by the physical axion field. The axion is the Nambu-Goldstone boson of the Peccei-Quinn $U(1)_{PQ}$ symmetry breaking [4, 5], and its effective Lagrangian has the form

$$\mathcal{L}_a = -\frac{1}{2} \partial_\mu a \partial^\mu a + \mathcal{L}_{int}(\partial_\mu a, \psi),$$

where the second term describes the coupling of the axion to fermion fields ($\psi$) of the Standard Model.

There are ongoing experimental searches for the axion and the cosmology and astrophysics provide strong complementary constraints. Because axions can be effectively produced in the interiors of stars they act as an additional sink of energy. The requirements that the energy loss from a star is consistent with the astrophysical observations place lower bounds on the coupling of axions to the Standard Model particles, and hence on the Peccei-Quinn symmetry breaking scale [6, 7]. The latter limit translates into an upper limit on the axion mass. Such arguments have been applied to the physics of supernova explosions [8–12] and white dwarfs [13]. In the case of supernova explosions the dominant energy loss process is the emission of an axion in the nucleon ($n$) bremsstrahlung $n + n \rightarrow n + n + a$. The same process was considered earlier by Iwamoto as a cooling mechanism for mature neutron stars, i.e., neutron stars with core temperature in the range $10^8 - 10^9$ K [14]. The implications of the axion emission by the modified nucleon bremsstrahlung, as calculated in Ref. [14], on the cooling of neutron stars were briefly discussed in Ref. [15]. However, it is now well established through cooling simulations of compact stars [16–18], that their neutrino cooling era, which spans the time period $t \leq 10^4 - 10^5$ yr after birth, is strongly affected by the neutrino emission from its superfluid phases due to the process of neutrino emission by Cooper-pair breaking [19–24].

In this article we compute the rate at which the superfluid phases of a neutron star lose their energy by axion emission via the processes of Cooper pair breaking and recombination. This work concentrates on the baryonic interiors of compact stars and considers for the sake of simplicity $S$-wave superfluids. The inner crusts and the baryonic core of a neutron star features iso-triplet spin-0 $S$-wave superfluids; in addition the core may contain spin-1, $P$-wave neutron superfluid [25, 27] and at high densities, iso-singlet, spin-1, neutron-proton $D$-wave superfluid [28]. The extension of the present work to $P$-wave and $D$-wave superfluids is straightforward. We use natural units, $\hbar = c = k_B = 1$. 

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2. Currents, matrix elements and emissivity

The coupling of axion to baryonic fields is described by the following interaction Lagrangian

\[ \mathcal{L}_{\text{int}} = \frac{1}{f_a} B^\mu L_\mu, \]  

(4)

where \( f_a \) is the axion decay constant, the baryon and axion currents are given by

\[ B^\mu = C_a \bar{\psi} \gamma^\mu \gamma_5 \psi, \quad L_\mu = \partial_\mu \phi, \]  

(5)

where \( C_a \) are model dependent coupling constants of order of unity. In the case of a multicomponent baryonic system the baryon current contains a sum over all components. The squared matrix element for the process of axion emission is then given by

\[ |\mathcal{M}_a|^2 = \frac{1}{2} f_a^{-2} (B^\mu B^\nu)(L_\mu L^\nu). \]  

(6)

The energy radiated per unit time in axions (axion emissivity) is given by the phase-space integral over the probability of the process of emission

\[ \epsilon_a = f_a^{-2} \int \frac{d^3q}{(2\pi)^32\omega} \omega g(\omega)q_\mu q_\nu \text{Im}\Pi^{\mu\nu}_a(q), \]  

(7)

where \( q \) and \( \omega \) are the axion momentum and energy. Here we defined the polarization tensor of baryonic matter

\[ \text{Im}\Pi^{\mu\nu}_a(\omega, \vec{q}) = \frac{1}{2} \sum_n (B_{\mu}^n B^{\nu\dagger}_n) \delta^4(q - \sum_i p_i), \]  

(8)

where the \( i \) sum is over the four-momenta of the baryons. Upon carrying out the angular integral in Eq. (7) we write the emissivity in terms of a one-dimensional integral

\[ \epsilon_a = f_a^{-2} \int_0^{\infty} d|\vec{q}| \ q^0 g(\omega) \kappa_a(q), \]  

(9)

where the contraction of the axion currents with the baryonic polarization tensor is given by

\[ \kappa_a(q) = q_\mu q_\nu \text{Im}\Pi^{\mu\nu}_a(q). \]  

(10)

So far the expression for the axion emissivity is completely general; we will need to compute the polarization tensor of baryonic matter for the process of interest.

3. Polarization tensor of superfluid baryon matter

At sufficiently low densities and temperatures baryonic matter forms a \(^1S_0\) pair condensate. In compact stars this is the case for all baryons except neutrons, which may form \( P \)-wave superfluid at densities at and above the saturation density. To describe the response
Figure 1: The two diagrams contributing to the polarization tensor of baryonic matter, which defines the axion emissivity. The “normal” baryon propagators for particles (holes) are shown by single-arrowed lines directed from left to right (right to left). The double arrowed lines correspond to the “anomalous” propagators $F$ (two incoming arrows) and $F^+$ (two outgoing arrows). The horizontal dashed lines represent the axion $a$.

of baryonic matter to the axion field we use the methods developed for the description of neutrino interactions in Refs. \[22, 24\] (see also \[23, 29\]). The imaginary-time momentum-space correlators are given by the $2 \times 2$ Nambu-Gor’kov matrix

$$G_{\sigma\sigma'}(i\omega_n, \vec{p}) = \begin{pmatrix} \hat{G}_{\sigma\sigma'}(i\omega_n, \vec{p}) & \hat{F}_{\sigma\sigma'}(i\omega_n, \vec{p}) \\ \hat{F}_{\sigma\sigma'}^+(i\omega_n, \vec{p}) & \hat{G}_{\sigma\sigma'}^+(i\omega_n, \vec{p}) \end{pmatrix},$$

(11)

The elements of the matrix are time-ordered correlators of the baryon field $\psi_B$ and $\psi_B^\dagger$; in the frequency-momentum domain these are given by

$$\hat{G}_{\sigma\sigma'}(i\omega_n, \vec{p}) = \delta_{\sigma\sigma'} \left( \frac{u_p^2}{i\omega_n - \epsilon_p} + \frac{v_p^2}{i\omega_n + \epsilon_p} \right),$$

(12)

$$\hat{F}_{\sigma\sigma'}(i\omega_n, \vec{p}) = -i\sigma_y u_p v_p \left( \frac{1}{i\omega_n - \epsilon_p} - \frac{1}{i\omega_n + \epsilon_p} \right),$$

(13)

where $F_{\sigma\sigma'}^+(i\omega_n, \vec{p}) = F_{\sigma\sigma'}(i\omega_n, \vec{p})$, $\omega_n = (2n + 1)\pi T$ is the fermionic Matsubara frequency, $\sigma_y$ is the $y$ component of the Pauli-vector in spin space, $u_p^2 = (1/2) (1 + \xi_p/\epsilon_p)$ and $u_p^2 + v_p^2 = 1 \epsilon_p = \sqrt{\epsilon_p^2 + \Delta_p^2}$ is the single particle energy in the paired state, with $\Delta_p$ being the (generally momentum- and energy-dependent) gap in the quasiparticle spectrum and $\xi_p = v_F (p - p_F)$ is the single-particle spectrum in the normal state, where $v_F$ and $p_F$ are the (effective) Fermi velocity and momentum. The propagator for the holes is defined as $\hat{G}_{\sigma\sigma'}^+(i\omega_n, \vec{p}) = G_{\sigma\sigma'}(i\omega_n, -\vec{p})$. For an $S$-wave condensate we have $\hat{G}_{\sigma\sigma'}(i\omega_n, \vec{p}) = \delta_{\sigma\sigma'} G(i\omega_n, \vec{p})$, $\hat{F}_{\sigma\sigma'}(i\omega_n, \vec{p}) = -i\sigma_y F(i\omega_n, \vec{p})$, etc. The polarization tensor of a superfluid obtains contributions from four distinct diagrams that can be formed from the normal and anomalous propagators with four distinct effective vertices. However, for the axial vector perturbations the vertices are not renormalized in the medium and, therefore, one proceeds with the bare vertices, in which case the number of the distinct contributions to the polarization tensor reduces to a sum of two admissible bare loops (see Fig. 1)

$$\bar{A}^T(q) \equiv \Pi_{GG}(q) - T \Pi_{FF}(q),$$

(14)

where $T = \pm 1$ is the time reversal operator and

$$\Pi_{GG}(q) = T \int \frac{d\vec{p}}{(2\pi)^3} \sum_{ip_n} G(p)G(p + q) = \frac{\nu(0)}{4} \int_{-1}^{1} dx (G * G),$$

(15)
where $T$ is the temperature, $p \equiv (ip_n, \vec{p})$ with $p_n$ being the fermionic Matsubara frequency, $\nu(0) = m^* p_F/\pi^2$ is the density of states of neutrons with $m^*$ being their effective mass, which may include the wave function renormalization (the so-called $E$-mass), $x$ is the cosine of the angle formed by the vectors $\vec{q}$ and $\vec{p}$. The convolution is defined as

$$(G * G) = T \int_{-\infty}^{\infty} d\xi_p \sum_{ip_n} G(p) G(p + q),$$

with a similar expression for $(F * F)$. Carrying out the Matsubara sums appearing in the convolutions we find

$$A^\pm(q) = -2\nu(0) \left[ \frac{x\delta}{1 - x\delta} + \frac{(1 \pm 1)}{2} \right] (F * F^+),$$

where $(\ldots) \equiv \frac{1}{2} \int_{-1}^{1} dx(\ldots)$, $x \equiv \hat{q} \cdot \hat{p}$ and $\delta = |q| v_F/\omega$ is a small parameter of the theory. To leading order in $\delta$ parameter the imaginary part of the $(F * F)$ convolution can be computed analytically

$$\text{Im}(F * F^+)_0 = \frac{\pi \Delta^2 \theta(\omega - 2\Delta)}{\omega} \tanh \left( \frac{\omega}{4T} \right),$$

whereas the real part follows from the dispersion relation

$$\text{Re}(F * F^+)_0 = \frac{1}{\pi} \int_{-\infty}^{\infty} \frac{d\omega'}{\omega - \omega'} \text{Im}(F * F^+)_0(\omega').$$

For the contraction (10) we obtain

$$\kappa_\alpha(q) = \omega^2 \text{Im}\Pi^{\alpha\alpha}(q) - q_i \omega \text{Im}\Pi^{0\alpha}(q) - \omega q_j \text{Im}\Pi^{0j}(q) + \frac{q^2}{3} \text{Im}\Pi^{ij}(q),$$

where, to the accuracy we are working, $\text{Im}\Pi^{0\alpha}(q) = A^+ v_F^2$, $\text{Im}\Pi^{ij}(q) = 3A^-(q)$, $q_j \text{Im}\Pi^{0j}(q) = \tilde{A}^- q v_F$ and $q_i \text{Im}\Pi^{0i}(q) = \tilde{A}^+ q v_F$. Here $A^\pm \equiv A^\pm(\hat{q} \cdot \hat{p})$ are the first moments of $A^\pm$ integrals with respect to the cosine of the angle formed by the vectors $\vec{q}$ and $\vec{p}$. The leading order in $\delta$ expansion of the contraction gives

$$\kappa_\alpha(q) = -2\nu(0) \left[ \omega^2 v_F^2 \left( 1 + \frac{\delta^2}{3} \right) - \frac{2\omega^2 \delta^2}{3} + \frac{\bar{q}^2 \delta^2}{3} \right] \text{Im}(F * F^+)_0$$

$$= -\frac{4\nu(0)}{3} \bar{q}^2 v_F^2 \text{Im}(F * F^+)_0 + O(\delta^4).$$

Eq. (21) completes our evaluation of the baryonic polarization tensor and its contraction with axionic current, which is the key input for the computation of the axion emissivity. Note that to $O(\delta^4)$ accuracy the $(F * F^+)$ convolution appears at its lowest ($\delta = 0$) order. More accurate evaluation is not required, because of the major uncertainty in the coupling strength of axions to the Standard Model fermions as well as other large uncertainties in the physics of cooling neutron stars.
4. Axion emissivity

The axion emissivity is obtained on substituting Eqs. (18) and (21) in Eq. (9)

\[ \epsilon_a = \frac{8}{3\pi} f_a \nu(0) v_F^2 T^5 I_a, \]  

(22)

where

\[ I_a = z^5 \int_1^\infty dy \frac{y^3}{\sqrt{y^2 - 1}} f_F(zy)^2, \]

(23)

\[ z = \Delta(T)/T \] and \( f_F(x) = [1 + \exp(x)]^{-1} \) is the Fermi distribution function. The \( T^5 \) scaling of the emissivity is understood as follows. The integration over the phase space of neutrons carries a power of \( T \), since for degenerate neutrons the phase-space integrals are confined to a narrow strip around the Fermi surface of thickness \( T \). The axion is emitted thermally and being relativistic contributes a factor \( T^3 \) to the emissivity. The one power of \( T \) from the energy of the axion and the inverse one power of \( T \) from the energy conserving delta function cancel. The transition matrix element is proportional to the combinations of \( u_p \) and \( v_p \) amplitudes, which are dimensionless, but contain implicit temperature dependence due to the temperature dependence of the gap function. This dependence is not manifest in Eq. (22), i.e., was absorbed in the definition of the integral \( I_a \). Thus, the explicit temperature dependence of the axion emission rate Eq. (22) is \( T^5 \). In the cgs units the axion emissivity Eq. (22) is

\[ \epsilon_a = 1.06 \times 10^{21} \left( \frac{10^{10}\text{GeV}}{f_a} \right)^2 \left( \frac{m^2}{m} \right)^2 \left( \frac{v_F}{c} \right)^3 \left( \frac{T}{10^9 K} \right)^5 I_a \text{erg cm}^{-3} \text{s}^{-1}, \]

(24)

where two powers of \( v_F/c \) arise from the small momentum transfer expansion and one power - from the density of states. At temperatures of order the critical temperature \( T_c \simeq 10^9 \) K the superfluid cools primarily by emission of neutrinos via the pair-breaking processes driven by the axial-vector currents (we continue to assume that potential fast cooling via direct Urca processes is prohibited). The emissivity of these processes in the case of \( 1S_0 \)-wave superfluid is given by [19, 20, 23]

\[ \epsilon_\nu = \frac{4G_F^2 g_A^2}{15\pi^3} \zeta_A \nu(0) v_F^2 T^7 I_\nu, \]

(25)

where \( G_F \) is the weak Fermi coupling constant, \( g_A = 1.25 \) is the axial-vector current coupling constant, \( \zeta_A = 6/7 \) and

\[ I_\nu = z^7 \int_1^\infty dy \frac{y^5}{\sqrt{y^2 - 1}} f_F(zy)^2. \]

(26)

We now require that the axion luminosity does not exceed the neutrino luminosity, i.e.,

\[ \frac{\epsilon_a}{\epsilon_\nu} = \frac{10\pi^2}{f_a^2 G_F^2 g_A^2 \zeta_A} I_a I_\nu < 1. \]

(27)
Substituting the the free-space value of the axial vector coupling $g_A = 1.25$ and introducing $r(z) \equiv z^2 (I_a/I_\nu)$ we transform Eq. (27)

$$\frac{\epsilon_a}{\epsilon_\nu} = \frac{59.2}{f_a^2 G_F^2 \Delta(T)^2} r(z).$$

(28)

Not far from the critical temperature $\Delta(T) \simeq 3.06 T_c \sqrt{1 - t/T}$, which translates into $z = 3.06 t^{-1} \sqrt{1 - t}$, where $t = T/T_c$. Numerical evaluations of the integrals provides the following values $r(0.5) = 0.07$, $r(1) = 0.26$, $r(2) = 0.6$ and asymptotically $r(z) \to 1$ for $z \gg 1$. Substituting the value of the Fermi coupling constant $G_F = 1.166 \times 10^{-5}$ GeV$^{-2}$ in Eq. (28) and noting that $r(z) \leq 1$, we finally obtain

$$f_a > 5.92 \times 10^9 \text{GeV} \left[ \frac{0.1 \text{ MeV}}{\Delta(T)} \right]$$

(29)

which translates into an upper bound on the axion mass

$$m_a = 0.62 \times 10^{-3} \text{eV} \left( \frac{10^{10} \text{GeV}}{f_a} \right) \leq 1.05 \times 10^{-3} \text{eV} \left[ \frac{\Delta(T)}{0.1 \text{ MeV}} \right].$$

(30)

The bound (29) can be written in terms of the critical temperature by noting that $\Delta(T) \simeq T_c$ in the temperature range $0.5 \leq t < 1$ of most interest.

### 5. Discussion

The neutrino cooling era of compact stars, which spans the time-period $t \leq 10^4 - 10^5$ yr after their birth in supernova explosions is a sensitive probe of the particle physics of their interiors. If one assumes that there are no rapid channels of cooling in neutron stars, i.e., deconfined quarks, above Urca threshold fractions of protons or hyperons (all of which lead to a rapid Urca cooling), then neutron stars cool primarily by neutrino emission in Cooper pair-breaking processes in baryonic superfluids. Here we have shown that if axions exist in Nature, the neutron stars must cool via axion emission in Cooper pair-breaking processes, whose axion emission rate scales as $T^5$. This scaling differs from the $T^7$ scaling of the counterpart neutrino processes. The difference arises from the different phase spaces required for an axion and a pair of neutrinos and is independent of the baryonic polarization tensor. Note also that the rate of axion emission from a $P$-wave superfluid will differ from the $S$-wave rate, derived above, by a factor $O(1)$ and, therefore, will not change quantitatively the obtained bounds on the axion parameters.

If the Standard Model physics provides a consistent explanation of the data on the cooling of neutron stars, one can place a lower bound on the axion decay constant (the breaking scale of the Peccei-Quinn symmetry). Our calculations show that this bound is given by $f_a > 5.92 \times 10^9 T_c^{-1}$ GeV, where $T_c$ is the magnitude of the critical temperature in units of $10^9$ K. This translates into an upper bound on the axion mass $m_a \leq 10^{-3} T_c$ eV. Similar bounds were obtained previously by Iwamoto [14] ($f/10^{10} \text{GeV} > 0.3$) from comparison of the rates of axion bremsstrahlung and modified Urca neutrino emission by mature neutron stars,
and by Umeda et al [15] \((f/10^{10}\text{GeV} > 0.1 - 0.2)\) from fits of cooling simulations to the PSR 0656+14 data. Our lower bound on \(f_a\) is somewhat larger than the one that follows form the requirement that the axions do not “drain” too much energy from supernova process so that it fails \([6, 8-12]\). Furthermore, our bound sensitively depends on the pairing gap in baryonic superfluids, whose magnitude and density dependence are not well-known. The physical implications of the bound \((29)\) can be fully explored with numerical simulations of neutron star cooling. Targeted fits to a specific object exhibiting slow cooling would be more useful than fits to the entire population of neutron stars with measured X-ray fluxes. Examples, of such fits were carried out, e. g., in the case of the neutron star in CAS A \([30, 31]\). The accuracy of the predictions will be limited by the uncertainties in the physical input in cooling simulations and uncertainties inherent to the interpretation of the data. For stars in the neutrino cooling era the dissipative heating processes are unimportant. The potential sources of uncertainty are well documented in the literature \([16-18]\) and include (a) the rate of neutrino/axion emission, in particular, its dependence on the magnitude of the pairing gap; (b) the composition of the surface layers; (c) the influence of the star’s \(B\)-field on the photon and neutrino emission processes; (d) the estimate of the age of any given object. We anticipate that the error bars on the bounds should remain within a factor of few and the accuracy of the bounds can be improved with the help of numerical simulations of cooling compact stars.

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Note that Eq. (9) of this work is incorrect, which could be the reason why the limits on the axion’s mass reported in their Figs. 2-4 are by an order of magnitude larger compared to those derived in this work.
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