Axion Emission From Neutron Stars

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We show that axion emission from neutron stars is the dominant energy-loss mechanism for a range of values of the Peccei-Quinn symmetry-breaking scale (F) not excluded by previous constraints. This gives the possibility of obtaining a better bound on F from measurements of surface temperature of neutron stars.

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The axion, a pseudo-Goldstone boson associated with the Peccei-Quinn symmetry,[1] was introduced as a natural solution to the strong CP problem.[2] The symmetry-breaking scale F, however, is left undetermined in the theory.[3] The present astrophysical[4] and cosmological considerations[5] have placed bounds on the parameter, and are consistent with 10^8 GeV $< F < 10^2$ GeV. Since axions couple very weakly to matter in this parameter range, they qualify as a candidate for dark matter in the universe required in the missing-mass problem. In addition, they could play an important role in galaxy formation,[6] with the axionic domain walls being the seeds for density perturbations in the early universe.[7] With these attractive features of axions, a better determination of the parameter F is needed. At present, unfortunately, the constraint on F from laboratory experiments is far less restrictive than the astrophysical bound, and future experiments for detecting axions by means of axion haloscope, helioscope,[8] and measuring the macroscopic forces which they mediate[9] have been proposed.

In this Letter we consider the effects of axion emission from neutron stars. We first calculate the axion emission rates from neutron-star matter. We point out that for neutron stars the conventional neutrino cooling scenario must be completely modified if the value of F lies close to 10^8 GeV. In such a case, we show that neutron stars would cool predominantly as a result of axion emission rather than neutrino emission. The motivation for considering neutron stars as a source of axions is as follows. Neutrinos couple to matter via vector and pseudovector vertices, while the axion's coupling is pseudoscalar. Neutrino emission either is accompanied by a charged lepton ($e^-, \bar{\mu}, \ldots$) or occurs as a pair ($\nu\bar{\nu}$) as a result of lepton number conservation, while a single axion may be emitted. Now, each degenerate fermion ($n, p, e^-, \ldots$) contributes a factor of T/T_F (which is typically $\sim 10^{-3}$–10^{-4}) from phase space because of the high degeneracy of neutron-star matter, where T is the temperature and T_F an appropriate Fermi temperature. Each relativistic thermal particle, on the other hand, introduces T^3 from phase space. Therefore, the differences in the energy dependence of the vertices as well as in the numbers of degenerate fermions and thermal relativistic particles result in quite a different energy-loss rate for axion emission compared with neutrino emission.

Let us start by looking at the axion emission processes in neutron-star matter. Among the processes examined, we find the following two processes to be the most important: axion bremsstrahlung from neutron-neutron collision,

$$n + n \rightarrow n + n + A, \tag{1}$$

as illustrated in Fig. 1(a), and axion bremsstrahlung by the electrons in the crust,

$$e^- + (Z,A) \rightarrow e^- + (Z,A) + A, \tag{2}$$

as shown in Fig. 1(b). We first outline the calculation of the energy-loss rate from (1). The axion-nucleon interaction Lagrangian density is given by[3] $\mathscr{L}_N = i g_{ANN} \bar{\psi}_n \gamma_5 \psi_n \phi_A$, where $g_{ANN} = C_A m_N / F$ with C_A ($\simeq 1.25$) the axial-vector renormalization constant and m_N the nucleon mass. For neutron-neutron collisions, we use the one-pion–exchange potential derived from the p-wave pion-nucleon pseudovector coupling (with strength $f \simeq 1.00$) in the Born approximation.[10] We assign $(E_1, \vec{p}_1)$ and

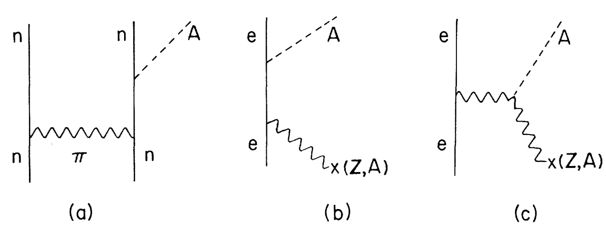

FIG. 1. Axion emission processes: (a) neutron bremsstrahlung, (b) electron bremsstrahlung (Compton-type process), and (c) the Primakoff process.

$(E_3, \vec{p}_3)$ for the initial and final four-momenta of a neutron, $(E_2, \vec{p}_2)$ and $(E_4, \vec{p}_4)$ for another neutron, and $(E_5, \vec{p}_5)$ for the axion. Then, the squared matrix element summed over the initial and final neutron spins is[11]

$$|M^{(1)}|^2 = \sum_{\sigma_1\sigma_2\sigma_3\sigma_4} |M_{fi}^{(1)}|^2 = \frac{256}{3} m_n^{*2} \left(\frac{f}{m_\pi}\right)^4 g_{ANN}^2 \left[\frac{|\vec{k}|^4}{(|\vec{k}|^2 + m_\pi^2)^2} + \frac{|\vec{l}|^4}{(|\vec{l}|^2 + m_\pi^2)^2} + \frac{|\vec{k}|^2|\vec{l}|^2}{(|\vec{k}|^2 + m_\pi^2)(|\vec{l}|^2 + m_\pi^2)}\right], \tag{3}$$

where m_π is the pion mass, m_n^* is the neutron effective mass, $\vec{k} \equiv \vec{p}_1 - \vec{p}_3$, and $\vec{l} \equiv \vec{p}_1 - \vec{p}_4$. The energy-loss rate (power per unit volume) due to (1) is

$$\epsilon_{ANN} = \frac{1}{2V} \left(\prod_{i=1}^{5} V \int \frac{d^3p_i}{(2\pi)^3}\right) E_5 W_{fi}^{(1)} n(p_1) n(p_2) [1 - n(p_3)][1 - n(p_4)], \tag{4}$$

where

$$W_{fi}^{(1)} = V(2\pi)^4 \delta^{(4)}(p_3 + p_4 + p_5 - p_1 - p_2) |M^{(1)}|^2 / \prod_{i=1}^{5} 2E_i V$$

is the transition probability, V is the normalization volume, and $n(\vec{p}_i)$ is the Fermi function. The factor $\frac{1}{2}$ is introduced to avoid double counting of the states of identical particles. The factors in the form $1 - n(\vec{p}_i)$ take into account the Pauli blocking due to neutron degeneracy. Unlike photons, or neutrinos in the initial stages of stellar collapse, for F in the currently allowed range axions interact so weakly with matter that their mean free path is always much larger than the size of the star. The calculation of (4) is straightforward and one obtains[12]

$$\epsilon_{ANN} \simeq \frac{31}{1890\pi} \frac{g_{ANN}^2}{\hbar c} \left(\frac{f}{m_\pi}\right)^4 \frac{m_n^{*2} p_F(n)}{\hbar^4 c^6} F(x) (k_B T)^6, \tag{5}$$

where

$$F(x) = 1 - \tfrac{3}{2} x \arctan(1/x) + x^2/2(1 + x^2),$$

with $x = m_\pi c / 2p_F(n)$, $p_F(n)$ is the neutron Fermi momentum, and we have neglected the axion mass, which is very small: $m_A \sim 10^{-5}[F/(10^{12}\ \text{GeV})]$ eV.[3]

The temperature dependence in (5) may be understood as follows. The axions are emitted thermally, so that they have a characteristic energy $\sim k_B T$. The matrix element has no temperature (energy) dependence since a factor T^{-1} from the nonrelativistic neutron propagator cancels one power of T from the pseudoscalar coupling vertex. The phase-space integrals over momenta may be converted into energy integrals. Since the neutrons are degenerate, each phase-space integral is restricted to near the neutron Fermi sphere of thickness $\sim k_B T$ giving a factor of T^4 from four neutron states. The axion, being relativistic, gives a factor T^3 from its phase-space integral. There are one power of T from axion energy, a factor T^{-1} from the energy-conserving delta function, and another factor of T^{-1} from the normalization of the axion field operator. Thus, the axion emission rate from (1) is proportional to T^6.

The second process we consider occurs in the crust of the neutron star, which is composed of a lattice of ions and degenerate relativistic electrons. There, an electron collides with an ion of charge $Z|e|$ and mass number A and emits an axion [process (2)]. We find the Compton-type process [Fig. 1(b)] to be more important than the Primakoff process [Fig. 1(c)]. The interaction Lagrangian density is $\mathscr{L}_e = i g_{Aee} \bar{\psi}_e \gamma_5 \psi_e \phi_A$, where $g_{Aee} = m_e / F$ with m_e the electron mass.[3] The squared matrix element summed over the initial and final spins of the electron is

$$|M^{(2)}|^2 = \sum_{\sigma_1\sigma_2} |M_{fi}^{(2)}|^2 = \frac{2Z^2\alpha^2 g_{Aee}^2}{(|k|^2 + q_{FT}^2)} \left(\frac{1 + \beta\cos\theta_{23}}{1 - \beta\cos\theta_{13}} + \frac{1 + \beta\cos\theta_{13}}{1 - \beta\cos\theta_{23}} + \left[\frac{2\beta^2(\cos\theta_{13}\cos\theta_{23} - \cos\theta_{12})}{(1 - \beta\cos\theta_{13})(1 - \beta\cos\theta_{23})}\right]\right), \tag{6}$$

where the indices $i = 1, 2, 3$ denote the initial and final electrons and the axion, respectively, $\beta = [1 + (m_e/|\vec{p}_1|)^2]^{-1/2}$, and $\cos\theta_{13} = \vec{p}_1 \vec{p}_3 / |\vec{p}_1||\vec{p}_3|$, etc. The factor $Z|e|/(|\vec{k}|^2 + q_{FT}^2)$ comes from the screened static Coulomb field of the ion with the Thomas-Fermi wave number $q_{FT} = (4\alpha/\pi)^{1/2} p_F(e)$, where

$\alpha = e^2/4\pi \simeq \frac{1}{137}$. The energy-loss rate (power) is

$$\epsilon'_{Ae} = V^{-1}\left[\prod_{i=1}^{3} V\int \frac{d^3p_i}{(2\pi)^3}\right]\int \frac{d^3k}{(2\pi)^3} W_{fi}^{(2)} E_3 n(\vec{p}_1)[1-n(\vec{p}_2)], \tag{7}$$

where

$$W_{fi}^{(2)} = V(2\pi)^4\delta^{(4)}(p_2+p_3-p_1-k)|M^{(2)}|^2/\prod_{i=1}^{3} 2E_i V$$

with $k=(0,\vec{k})$. Performing the phase-space integrals in (7), one finds the energy-loss rate (power per unit volume) to be

$$\epsilon_{Ae} = \epsilon'_{Ae} n_b/A \simeq \frac{\pi^2}{120}\frac{Z^2\alpha}{A}\frac{g_{Aee}^2}{\hbar c}\frac{n_b(k_B T)^4}{\hbar[cp_F(e)]^2}[2\ln(2\gamma)-1], \tag{8}$$

where γ is the Lorentz factor of the electron with Fermi momentum $p_F(e)$. Again, the temperature dependence of (8) may be explained easily. The matrix element (6) is temperature (energy) independent since a factor of T^{-1} from the electron propagator cancels one power of T from the pseudoscalar coupling vertex. The initial and final relativistic degenerate electrons give one power of T from each phase-space integral, respectively. The axion gives T^3 from its phase-space integral, one power of T comes from axion energy, a factor of T^{-1} from the energy-conserving delta function, and another factor of T^{-1} from the boson normalization. Altogether, the axion emission rate from (2) is proportional to T^4.

These two processes are to be compared with the following neutrino emission processes: the modified URCA process[13]

$$n+n \to n+p+e^- +\bar{\nu}_e, \quad n+p+e^- \to n+n+\nu_e, \tag{9}$$

and the neutrino bremsstrahlung by the electrons in the crust[14]

$$e^- + (Z,A) \to e^- + (Z,A) + \nu + \bar{\nu}. \tag{10}$$

The processes (9) and (10) have temperature dependences $\propto T^8$ and $\propto T^6$, respectively. From Eqs. (5) and (8) we find that the axion emission rates in fact have temperature dependences different from the neutrino case. Axion emission tends to dominate over neutrino emission at lower temperatures, as a result of its milder temperature dependence.

We now make a more quantitative comparison of axion and neutrino emission. For this purpose, we must specify the equation of state of neutron-star matter. We allow for the theoretical uncertainties in nuclear physics by choosing two types of equations of state: a medium-soft equation of state by Bethe and Johnson (BJ)[15] and a stiff equation of state by Pandharipande and Smith (PS).[16] Soyeur and Brown[17] have calculated the neutrino luminosity in detail by using these two equations of state. They find that, without nucleon superfluidity, the modified URCA process[11] is the most important energy-loss mechanism at temperatures down to $\sim 2\times 10^8$ K ($\equiv T_1$); below that the photon radiation from the surface takes over. Comparing (5) and the modified URCA rate,[13] one finds that axion emission will dominate neutrino emission unless[10,18] $g_{ANN} < 2.1\times 10^{-9}[T/(10^9\ \mathrm{K})]$. Taking $T = T_1$, one obtains the condition for the axions not to interfere with the neutrino cooling, $g_{ANN} < 4\times 10^{-10}$ or $F > 3\times 10^9$ GeV.

In the presence of nucleon superfluidity,[19,20] the condition becomes less restrictive. Proton superfluidity is expected to set in first at $T_p \sim 4.6\times 10^9$ K.[19] Then, all the processes involving protons are suppressed as a result of the energy gap, estimated to be $\Delta_p \sim 0.7$ MeV.[19] Next, the neutrons become superfluid at $T_n \sim 7\times 10^8$ K.[20] For this reason, the rates of neutrino emission by nucleons drop sharply below $\sim 10^9$ K (BJ) or 4×10^9 K (PS) ($\equiv T_2$), where the neutrino bremsstrahlung by the electrons in the crust (10) becomes the major energy-loss mechanism.[17] We compare the neutrino rate from (10) at the crust density $n_c = 0.5 n_0$ with (5) at central densities and obtain $g_{ANN} < 4\times 10^{-10}$ (BJ) and 6×10^{-10} (PS) or $F > 3\times 10^9$ GeV (BJ) and 2×10^9 GeV (PS) in order for the axions not to dominate the neutrino rate at $T < T_2$. The photon radiation eventually takes over at temperatures below $\sim 3\times 10^8$ K (BJ) or 2×10^8 K (PS) ($\equiv T_3$). At $T = T_3$ we can compare the rate (8) with the neutrino bremsstrahlung rate (10), and obtain the condition for the axions not to dominate photon cooling at $T < T_3$: $g_{Aee} < 9\times 10^{-13}$ (BJ) and 6×10^{-13} (PS) or $F > 6\times 10^8$ GeV (BJ) and 9×10^8 GeV (PS). We note that in spite of the uncertainties in the equation of state, axion emission could be the dominant cooling mechanism if $10^8 < F \lesssim 3\times 10^9$ GeV. This parameter range of F has not been ruled out from any previous considerations.[4,5]

Finally, we discuss the implications of our analysis in the light of observations. The Einstein Observatory found three x-ray emitting point sources in the centers of supernova remnants (Crab, Vela, and RCW 103).[21] The cooling calculations based on the standard neutrino scenario are found to be consistent with the observed x-ray flux from these objects within the theoretical uncertainties.[22] If one tentatively concludes that there should not be any extra cooling mechanisms other than neutrino emission, one may obtain a bound on the axion parameters. In order to do so one must keep the following in mind. The observation of the x-ray spectra is needed to assure that these x rays are in fact thermal. In addition, one must make sure that there are no heating mechanisms at work inside neutron stars. With these provisions, we stress the potential importance of the x-ray observations of neutron-star surface temperatures as a means to obtain a better bound on F.

In summary, the standard scenario of neutron-star cooling due to neutrino emission must be modified unless the Peccei-Quinn symmetry-breaking scale for axions $F > 6\times10^8$–3×10^9 GeV. The uncertainty comes from our lack of knowledge of the neutron-star matter equation of state. The region $10^8 < F \lesssim 3\times10^9$ GeV has not been ruled out from the previous considerations. We have also pointed out the potential importance of the data of x-ray flux from the neutron stars in placing a better bound on F.

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[10]We expect the many-body effects on nucleon-nucleon interactions in (1) to be similar to those of the corresponding process for neutrino emission. We shall employ the simplest interaction and compare these processes on an equal footing.

[11]The axion line may be attached to any of the four neutron lines in Fig. 1(a), which gives four distinct direct diagrams. In addition, there are four exchange diagrams. These eight amplitudes are to be summed, giving the squared matrix element 64 terms.

[12]For simplicity, we have approximated the third term in the square bracket of Eq. (3) to be the same as the first two terms. This approximation becomes exact when $m_\pi c/p_F(n) \rightarrow 0$.

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