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CONSTRAINTS ON AXIONS FROM SN 1987A

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ABSTRACT

We argue that the detection of neutrinos from SN 1987a implies that axion emission from the collapsing star must not have been significant. The best axion supernova limits come from nucleon bremsstrahlung processes which involve axion-quark couplings rather than the electron or photon couplings used in obtaining axion limits from red giants. The axion-quark couplings are less sensitive to specific axion model parameters and therefore are capable of yielding a model independent lower bound on f_a. Quantitatively this yields an axion coupling limit of $f_a \gtrsim 8\times10^{11}$ GeV (or equivalently $m_a \lesssim 9\times10^{-5}$ eV). This approaches the upper bound on $f_a \lesssim 4\times10^{12}$ GeV from cosmological density arguments. While the bounds do not exclude the axion, if the axion exists, it must dominate the cosmological mass density. This same argument can also be applied to analogous couplings for any particle with $m \lesssim 50$ MeV.

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The detection of neutrinos from SN 1987a by Kamioka[1] and IMB[2] has provided us with a new laboratory for testing exotic processes. In particular, it is known[3] that through neutral currents, all types of neutrinos will be emitted during a gravitational collapse in roughly equal amounts. Thus $\sim 1/6$ come out as detectable $\bar{\nu}_e$'s if there are three generations[4] and if no other kinds of particles are emitted. The total amount of energy available to be radiated as neutrinos is the neutron star binding energy, $\sim 3 \times 10^{53}$ ergs. Thus if more species of neutrinos exist, or if some other species of particle is also emitted from the stellar collapse, then the fraction of binding energy emitted as $\bar{\nu}_e$'s will be reduced. This argument has previously been used to set limits on the number of neutrino flavors.[5,6] In this paper we extend the argument to constrain axions[6] (or other (pseudo) Nambu-Goldstone bosons).

Peccei and Quinn[7] proposed a global (pseudo) symmetry whose spontaneous breakdown led to the existence of an almost massless (pseudo) Goldstone boson, the axion, to solve the strong CP problem.[8] The original axion, however, interacted too strongly and had a mass and lifetime which were excluded by laboratory experiments[9] as well as astrophysical data.[10] Dine, Fischler, and Srednicki (DFS)[11] modified the Peccei-Quinn proposal to yield a much more weakly interacting—"invisible"—axion which was not in violation of laboratory and/or astrophysical[12,13] constraints. The couplings of the DFS invisible axion to electrons and photons were constrained from Red Giants[14] by requiring the axion cooling not to prevent helium ignition.

In what follows, we calculate limits on the axion decay constant f_a, which can be related to the axion mass and coupling[15]

$$m_a \simeq 7.2 \times 10^{-5}(10^{12}\ \mathrm{GeV}/f_a)(\frac{N}{6}) \tag{1}$$

$$g_{ae} = 4(\cos^2\beta)(m_e/f_a) \tag{2}$$

$$g_{aN} = C_{AN} m_N/f_a \tag{3}$$

where $g_{ae(N)}$ is the axion coupling to electrons (nucleons), $C_{AN}(\cos^2\beta)$ contains the details of the quark content of nucleons we discuss in more detail below and β is an arbitrary angle depending on the Higgs sector of the low energy theory. For example, if there are two Higgs doublets H and $\bar{H}$ giving masses to up and down quarks respectively, with vacuum expectation values $v = \langle H \rangle$ and $\bar{v} = \langle \bar{H} \rangle$, $\cos^2\beta = v^2/(v^2+\bar{v}^2)$. One should be aware that in the literature, there are several conventions for the normalization of f_a (e.g. Bardeen and Tye[16] take $f_a \to 2f_a$).

The best limit to date on f_a has been the red giant limit[14]

$$f_a > 1.4 \times 10^{10}\cos^2\beta \tag{4}$$

For comparison one should remember that the cosmological upper bound[17] on f_a is

$$f_a \leq 4 \times 10^{12\pm0.34}\left(\frac{\Lambda_{QCD}}{200\ \mathrm{MeV}}\right)^{0.6}(\frac{N}{6}) \tag{5}$$

(where N is the number of quarks with Peccei-Quinn charges, usually taken as the number of quark flavors; we assume $N = 6$). The bound (5) is for a Standard Big Bang Universe where all angles of the axion field are averaged over. In an inflating universe only one angle is selected since only one primordial horizon gives us the present universe. Thus, depending on the value for this angle, the limit may be increased accordingly. Recently it had been suggested[6] that because of the extreme conditions that were present during the collapse of SN 1987a, one could improve on the lower bound on f_a. Though the result of reference 6 was comparable to the red giant limit, this calculation requires correction and here we will discuss the supernova limit in more detail. One should also note that all of the previous limits depend on the undetermined angle β, whereas our supernova limit below is valid for all values of β.

As we will see, for gravitational collapse events such as SN 1987a, the dominant axion production is via nucleon bremsstrahlung processes rather than lepton or photon axion production processes. Thus the supernova constrains quark-axion couplings and so complements the red giant limits on electron and photon axion couplings. The previous limits on quark-axion couplings came from neutron-star cooling[18] where a limit of $f_a \gtrsim 6 \times 10^8 - 3 \times 10^9$ GeV was obtained but was sensitive to assumptions about nuclear structure. We will see that a significantly better limit can be obtained from the supernova and that it is stronger than the red giant limit. It should also be remembered that it is conceivable for axion models to be created without electron-axion couplings but that quark-axion couplings of some type are always present.

To obtain our axion limits we have to evaluate axion emission rates in the density-temperature regime appropriate to the final core collapse. This is a regime at far higher temperature ($T \sim 1$ to 70 MeV) than that utilized in red giant calculations ($T \sim 0.1$ MeV), and above the electron-positron rest mass, so that previous techniques did not apply. Also, unlike the red giant case where stellar evolution is dominated by photon emission and diffusion time scales, the later stages of stellar evolution (carbon burning and beyond) are dominated by neutrino emission. Until the final collapse when the density ρ exceeds $\sim 2 \times 10^{11}$ g/cm^3, the star is transparent to the neutrinos which free stream out of the star. Hence for $\rho \lesssim 2 \times 10^{11}$ g/cm^3 if the axion emission rates at a given temperature are below the neutrino emission rates, the neutrino emission fully dominates and axions have little effect. Thus the most restrictive axion limits are expected to come once densities exceed $\sim 2 \times 10^{11}$ g/cm^3 and the core is opaque to neutrino propagation (a "neutrinosphere" forms) but the more weakly coupled axions are still able to escape from the higher temperature interior. (A more quantitative discussion of axion mean-free-paths will be given later.) Red giant and neutron star cooling arguments suggest that the axion, if it exists, must couple sufficiently weakly to be in this regime. The temperature of the neutrinosphere is $T(\nu_e) \sim 3$ MeV, $T(\nu_\mu, \nu_\tau) \sim 6$ MeV, and the central core reaches $T \sim 70$ MeV. Thus neutrino emission is characterized by the rates at temperatures more than an order of magnitude below the axion emission temperatures, and yet we know that neutrino emission dominated for SN 1987a. It is this effect that enables interesting limits to be obtained.

Axion Emission Rates

To determine our limits on the axion decay constant f_a, we will consider the following processes: 1) compton scattering off electrons, $\gamma + e \rightarrow a + e$; 2) electron-positron annihilation, $e^+e^- \rightarrow \gamma + a$; 3) plasmon decay, $\gamma \rightarrow \gamma\prime + a$ and 4) nucleon-nucleon scattering $N + N \rightarrow N + N + a$. As we will see it is the last of these which is dominant during a supernova collapse and therefore gives us the strongest limit on f_a. The limits on the axion-electron couplings are not as restrictive as the previously quoted limits from red giants.

We begin by discussing the electron processes. The rates for compton scattering and annihilation have been calculated. At very high energies the cross sections for these two processes are equal[19]

$$\sigma_{\gamma e\to ae} = \sigma_{ee\to a\gamma} \simeq \frac{2\pi\alpha\alpha_{ae}}{s} ln\, s/m_e^2 \tag{6}$$

for center of mass energy squared $s >> m_e^2$ where $\alpha = 1/137$ is the fine structure constant and $g_{ae} = (4\pi\alpha_{ae})^{1/2}$ is the coupling of axions to electrons.

The cross-sections (6) yield an energy loss rate given by

$$R = \frac{3.2\times 10^{47}}{4\pi^3\rho}\alpha\alpha_a T^5\{2\zeta(3)I_{1,2} + 2I_{1,1}I_{2,1} + \frac{\pi^2}{6}I_{2,2}\}\ \mathrm{erg/gs} \tag{7}$$

where

$$I_{n,m} = \int_0^\infty dx x^n\, ln^m\left(\frac{2x}{z}\right) f_x \tag{8a}$$

$$f_x = (e^x - 1)^{-1}, \quad z = m_e/T \tag{8b}$$

For our supernova collapse with $T \simeq 70$ MeV and $\rho = 10^{15}$ g/cm^3, we find $I_{1,2} \simeq 52$, $I_{1,1} \simeq 9$, $I_{2,1} \simeq 15$, $I_{2,2} \simeq 99$ and $R \simeq 1.8\times 10^{40}\alpha_{ae}$ erg/g s, which for a 0.5 $M_\odot$ central core (at these conditions) gives $L = 1.8\times 10^{73}\alpha_{ae}$ erg/s which, over the duration of the collapse ($\sim$ 10 seconds) yields a total energy output of $E = 1.8\times 10^{74}\alpha_{ae}$ ergs and can be compared to an allowed excess over the neutrino output of $E_{a_{max}} = 3\times 10^{53}$ erg. Thus we arrive at the limit $\alpha_a < 1.7\times 10^{-21}$ or $f_a > 1.4\times 10^7$ GeV $cos^2\beta$.

At very high densities we might expect plasmon decay to produce an even higher energy loss rate as is the case with neutrino emission. However, unlike neutrino emission or axion emission at low temperatures by plasmon decay the largest energy scale in the problem is the temperature. We expect that in the limit of interest where the plasma frequency $\omega >> m_e$ as well as $T >> m_e$ the order of magnitude of the matrix element is given by the conventional anomalous triangle diagram with the electron treated as a massless fermion. Therefore, we have approximated the matrix element for the decay by

$$|M|^2 = \frac{e^4\cos^4\beta}{4\pi^4 f_a^2}\epsilon_{\mu\nu\alpha\beta}\,\epsilon^{\mu\nu\gamma\delta}p^\alpha k^\beta p_\gamma k_\delta \tag{9}$$

where p, k are the plasmon and axion 4-momenta respectively. From equation(9), we find a total decay rate $\Gamma_a = \frac{\alpha^2\omega}{16\pi^3 f_a^2}$ and an energy loss rate

$$R = \frac{4\alpha^2\omega^4T^3\cos^4\beta}{(2\pi)^5 f_a^2\rho}\int_{\omega/2T}^\infty dx f_x \int_{\omega/2T}^\infty z dy f_y$$

$$\simeq 3.65\times 10^{30}\cos^4\beta(1GeV/f_a)^2 \mathrm{erg/gs} \tag{10}$$

where $\omega \simeq 28$ MeV (at T=70 MeV and $\rho = 10^{15}$g/cm^3). This corresponds to a total energy output (assuming $M = .5M_\odot$ and $\Delta t = 10$ seconds) $E = 3.7\times 10^{64}\cos^4\beta(1\ \mathrm{GeV}/f_a)^2$ergs and the limit (taking 3×10^{53} ergs in axions) of $f_a > 3.5\times 10^5$ GeV $\cos^2\beta$ which is not competitive.

Because of the extreme nucleon density during collapse, it is important to consider the emission of axions due to nucleon-nucleon scatterings. The axion-nucleon couplings can be written down by starting from the axion current[15]

$$j_\mu^a = j_\mu^{PQ} - \frac{N}{(1+z+w)}(\bar{u}\gamma_\mu\gamma_5 u + z\bar{d}\gamma_\mu\gamma_5 d + w\bar{s}\gamma_\mu\gamma_5 s) \qquad (11)$$

where

$$j_\mu^{PQ} = f_a\partial_\mu a + \sum_i x_i\bar{q}_i\gamma_\mu\gamma_5 q_i \qquad (12)$$

In these equations,

$$z \equiv \frac{m_u}{m_d}, w \equiv \frac{m_u}{m_s} \qquad (13)$$

which we estimate to be $z = 0.565$, $w = 0.029$ respectively, N is the number of quarks coupling to the axion, assumed here to be 6, and in the generic models discussed here

$$x_u + x_d = 2,\ \ x_d = x_s,\ \ x_u - x_d = -2\cos 2\beta. \qquad (14)$$

Gathering terms, we can rewrite (11) as

$$\begin{aligned} j_\mu^a = f_a\partial_\mu a &+ (x_u - \frac{6}{1+z+w})\bar{u}\gamma_\mu\gamma_5 u \\ &+(x_d - \frac{6z}{1+z+w})\bar{d}\gamma_\mu\gamma_5 d \\ &+(x_s - \frac{6w}{1+z+w})\bar{s}\gamma_\mu\gamma_5 s. \end{aligned} \qquad (15)$$

To determine the couplings to nucleons, we need the matrix elements $\langle p|\bar{q}_i\gamma_\mu\gamma_5 q_i|p\rangle \equiv s_\mu\Delta q_i$ where s_μ is the proton spin, and the isospin relations $\langle p|\bar{u}\gamma_\mu\gamma_5 u|p\rangle = \langle n|\bar{d}\gamma_\mu\gamma_5 d|n\rangle$, $\langle p|\bar{d}\gamma_\mu\gamma_5 d|p\rangle = \langle n|\bar{u}\gamma_\mu\gamma_5 u|n\rangle$ and $\langle p|\bar{s}\gamma_\mu\gamma_5 s|p\rangle = \langle n|\bar{s}\gamma_\mu\gamma_5 s|n\rangle$. One combination of the $\langle p|q_i\gamma_\mu\gamma_5 q_i|p\rangle$ (Δq_i) is fixed by neutron β-decay:

$$\Delta u - \Delta d = g_A = 1.25 \qquad (16)$$

Another combination is fixed by hyperon β-decay data and flavor SU(3) symmetry for the baryon octet. This assumption fits the available data on hyperon β-decay very well and is normally used to extract from them the value of the Cabibbo angle.[20] It leads to

$$\Delta u + \Delta d - 2\Delta s = 3F - D = 0.682 \qquad (17)$$

where we have inserted the experimental value $F/D = 0.63$ which differs slightly from the naive quark model (NQM) value often assumed, and used $g_A = F + D$. An extra assumption is needed to determine completely the three $\langle p|\bar{q}_i\gamma_\mu\gamma_5 q_i|p\rangle$. Conventionally it has been[21] assumed that $\langle p|\bar{s}\gamma_\mu\gamma_5 s|p\rangle = 0$ which has the physical meaning that strange quarks carry none of the nucleon spin. In this case one would have

$$NQM : \left\{ \begin{array}{l} \Delta u = 0.966 \\ \Delta d = -0.284 \\ \Delta s = 0 \end{array} \right\} \qquad (18)$$

Recently, an alternative estimate of the $\langle p|\bar{q}_i\gamma_\mu\gamma_5 q_i|p\rangle$ has become possible,[22] based on the EMC collaboration[23] measurement of the spin-dependent muoproduction structure function $g_1^p(x,Q^2)$. In the framework of QCD,

$$\int_0^1 dx g_1^p(x,Q^2) \simeq 1/2 \sum_i Q_i^2 \Delta q_i \tag{19}$$

where the operators are renormalized at the momentum scale Q in the case of the EMC data. If one neglects any dependence on Q^2 between the scales $Q^2 << 1GeV^2$ of neutron and hyperon β-decay and the scale $Q^2 \simeq 10GeV^2$ of the EMC data,* one can use the EMC measurement $\int_0^1 dx g_1^p(x,Q^2) = 0.113 \pm 0.012 \pm 0.025$ together with equations (16) and (17) to obtain

$$EMC : \begin{Bmatrix} \Delta u = 0.708 \\ \Delta d = -0.542 \\ \Delta s = -0.257 \end{Bmatrix} \tag{20}$$

We regard the differences between the estimates (18,20) as indicative of the uncertainties in estimating the matrix elements $\langle p|\bar{q}_i\gamma_\mu\gamma_5 q_i|p\rangle$. The effective axion-nucleon couplings are given in terms of these by

$$\begin{aligned} \mathcal{L} = \left(\frac{2m_N}{f_a}\right) a\{\bar{p}\gamma_5 p[(x_u - 3.76)\Delta u \\ +(x_d - 2.13)\Delta d + (x_s - 0.11)\Delta s] \\ +\bar{n}\gamma_5 n[(x_u - 3.76)\Delta d + (x_d - 2.13)\Delta u \\ +(x_s - 0.11)\Delta s]\} \end{aligned} \tag{21}$$

or in terms of $\cos 2\beta$ (14):

$$\begin{aligned} \mathcal{L} = 2\left(\frac{2m_N}{f_a}\right) a\{\bar{p}\gamma_5 p[(-2.76 - \cos 2\beta)\Delta u \\ +(-1.13 + \cos 2\beta)\Delta d + (0.89 - \cos 2\beta)\Delta s] \\ +\bar{n}\gamma_5 n[(-2.76 - \cos 2\beta)\Delta d + (-1.13 + \cos 2\beta)\Delta u \\ +(0.89 - \cos 2\beta)\Delta s]\} \end{aligned} \tag{22}$$

where we have used the numerical values of z, ω in the above. Finally, we can collect terms using both the NQM and EMC values for Δq_i to determine the couplings g_{an} and g_{ap}, using equation 3.

$$C_{ap} = \begin{Bmatrix} -4.7 - 2.5\cos 2\beta \text{ ; NQM} \\ -3.1 - 2.0\cos 2\beta \text{ ; EMC} \end{Bmatrix} \tag{23}$$

$$C_{an} = \begin{Bmatrix} -.61 + 2.5\cos 2\beta \text{ ; NQM} \\ .93 + 3.0\cos 2\beta \text{ ; EMC} \end{Bmatrix} \tag{24}$$

* This assumption has been questioned by Jaffe[24] but no large source of Q^2-dependence has yet been identified.

One should note that these expressions for C_A can lead to significant differences in the final results with respect to the naive choice of $C_A = 1.25$.

Because axion production by nucleon scattering is more significant than the previous processes, we will discuss it in more detail. Iwamoto[18] had calculated the energy loss rate for axion emission using neutron-neutron bremsstrahlung. To calculate the matrix element he assumed a one-pion exchange graph with a pseudo-vector (derivative) coupling of pions to nucleons. Iwamoto[18] estimated the matrix element for this process to be

$$|M|^2 = 256 m_N^2 \left(\frac{4\pi f^2}{m_\pi^2}\right)^2 g_{an}^2 \frac{|k|^4}{(|k|^2 + m_\pi^2)^2} \tag{25}$$

where $4\pi f^2 \approx 1$ is related to the pion-proton coupling and k is the momentum transfer between neutrons. Using a pseudo-scalar coupling of axion to nucleons Pantziris and Kang[25] find

$$|M|^2 = 2048 g_{an}^2 \left(\frac{4\pi f^2}{m_\pi^2}\right)^2 m_N^4 \frac{|k|^2}{(|k|^2 + m_\pi^2)^2} \tag{26}$$

which would lead to an enhancement of the emission rate by a factor $O(10^2)$. Both of these estimates are uncertain due to more complicated nuclear effects. However, we believe that because the pseudo-vector coupling at low energies gives better agreement with experimental cross-sections, and it is in any case more conservative we use the rate of Iwamoto. A more detailed calculation of this rate using low energy N-N scattering directly is in progress.[26]

From equation (25), Iwamoto[18] derives an expression for the energy loss rate due to axion production in neutron-neutron scattering

$$\dot{\epsilon}_n = \frac{31}{3780\pi} \left(\frac{4\pi f^2}{m_\pi^2}\right)^2 (kT)^6 g_{an}^2 m_N^2 p_F(n) F(x_n)$$

$$= 8.2 \times 10^{43} (1\mathrm{GeV}/f_a)^2 C_{an}^2 (\rho_{12} X_n)^{1/3} T_{MeV}^6 F(x_n) \ \mathrm{ergs/cm^3 s} \tag{27}$$

(we have divided Iwamoto's result by a factor of two to take into account the identical particles in both the initial and final states) where $p_F(n)$ is the neutron fermi momentum

$$p_F(n) = (3\pi^2 n_n)^{1/3}$$

$$= 50.9 (\rho_{12} X_n)^{1/3} \ \mathrm{MeV} \tag{28}$$

and n_n is the number density of neutrons, $X_n = \frac{m_N n_n}{\rho}$, $\rho_{12} = \rho/10^{12}$ g/cm^3, $x_n = m_\pi / 2p_F(n)$ and $F(x) = 1 - \frac{3}{2} x \tan^{-1}\left(\frac{1}{x}\right) + x^2/2(x^2+1)$. For the energy loss due to proton-proton scattering, we can use charge independence and

$$\dot{\epsilon}_p = 8.2 \times 10^{43} (1 \ \mathrm{GeV}/f_a)^2 C_{ap}^2 (\rho_{12} X_p)^{1/3} T_{MeV}^6 F(x_p) \ \mathrm{ergs/cm^3 s} \tag{29}$$

For the energy loss due to neutron-proton scattering, we must correct the above rates. First there is a factor of 4 because n and p are not identical particles, but there is also a factor of 1/2 due to the normalization of the superposition of np and pn. More importantly, there are two isospin channels for n-p scattering rather than one for n-n or p-p scattering.

In addition, it is simple to check that the $I = 0$ channel is 9 times more important thus resulting in an additional factor of 10. Collecting these factors we find that

$$\dot{\epsilon}_{np} = 1.6 \times 10^{45} (1\mathrm{GeV}/f_a)^2 ((C_{ap} + C_{an})/2)^2 (\rho_{12})^{1/3}$$
$$\times ((X_n^{1/3} + X_p^{1/3})/2) T_{MeV}^6 F(x_{np}) \text{ ergs/cm}^3/s \tag{30}$$

where $x_{np} = m_\pi/(p_F(n) + p_F(p))$.

Before calculating the limit on the decay constant f_a, we must check the conditions under which axions produced in the collapse actually escape. If we consider therefore the scattering of axions by nucleons, for $T << m_N$ we can use the cross-section for $a+e \to \gamma+e$ with the appropriate substitution of $m_e \to m_N$ and $g_{ae} \to g_{an}$ so that[19]

$$\sigma = \frac{e^2 {g_{an}}^2}{8\pi {m_N}^2} \left[\frac{m_N}{q} \log\left(1 + \frac{2q}{m_N}\right) - 2\frac{(1 + 3q/m_N)}{(1 + 2q/m_N)^2} \right] \tag{31}$$

where q is the axion energy. The mean free path for this process has been calculated[13]

$$l^{-1} = (5.1 \times 10^{10} \text{ cm}^{-1}) \frac{64 e^2 {g_{an}}^2}{48\pi^3 {m_N}^6} \frac{\zeta(5)}{\zeta(3)} (kT)^7 I(z,\ \lambda) \tag{32}$$

where

$$I(z,\ \lambda) = \int \frac{x^2(4x^2 + 3z^2)dx}{e^{(x^2+z^2)^{1/2} - \lambda} + 1} \tag{33}$$

and $\lambda = \mu_n/kT$. We find that $l = 5 \times 10^{-8} g_{an}^{-2}$ cm and that $l \gtrsim 10^6$ cm for $g_{an} < 2 \times 10^{-7}$ or $f_a \gtrsim O(10^7)$ GeV being the condition for axions to escape due to scatterings off nucleons. We can also compute the mean free path for the three body process $a + N + N \to N + N$ given by

$$l^{-1} \simeq \dot{\epsilon}\pi^2/(2.7)T^4 \tag{34}$$

where $\dot{\epsilon} = \dot{\epsilon}_n + \dot{\epsilon}_p + \dot{\epsilon}_{np}$ and $2.7T$ is the average axion energy. We find in this case that $l \simeq 4.64 \times 10^{-15} f_a^2$ and that $l \sim 10^6$ cm for $f_a \sim 10^{10}$ GeV, but that axions continue to escape unless $f_a \lesssim 10^8$ GeV.

Iwamoto's calculation was appropriate for the degenerate nucleon limit of neutron stars. However, for stellar collapse, this limit must be looked at with more care. We find from our collapse calculation[4] that $\frac{\mu_B}{T} \approx 1$ in the region where $\rho > 10^{13}$ g/cm^3 and $\mu_B = \mu_n - m_N$. To check on the validity of using the degenerate formula we took Iwamoto's general formula

$$\epsilon_{aNN} = \frac{1}{2V} \left[\prod_{i=1}^{5} V \int \frac{d^3 p_i}{(2\pi)^3} E_a W_{f_i}^{(1)} n(p_1) n(p_2)(1 - n(p_3))(1 - n(p_4)) \right] \tag{35}$$

where

$$W = V(2\pi)^4 \delta^4(\Sigma P) |M|^2 / \prod_{i=1}^{5} 2E_i V \tag{36}$$

and evaluated the integrals assuming $\frac{\mu_B}{T} << 1$. We find, neglecting blocking factors and taking $|M|^2 = 256 m_N^2 g_{aN}$, that

$$\dot{\epsilon} = 3.1 \times 10^{44} T_{MeV}^3 (1\mathrm{GeV}/\mathrm{f_a})^2 [C_{an}^2 {\rho_{12}}^2] \text{ ergs/cm}^3 \text{ s} \tag{37}$$

for neutron-neutron bremsstrahlung. By evaluating $\dot{\epsilon}$ throughout the ρ, T plane assuming $X_n = 1$, $X_p = 0$ for both degenerate and non-degenerate limits, we find that the degenerate formula is a good approximation in regions near $\frac{\mu_B}{T} = 1$. We find that the degenerate formula smoothly merges to the non-degenerate formula when $\frac{\mu_B}{T} < 1$.

For a stellar model to use in numerical evaluation of the axion energy loss we use the $1.64 M_\odot$ model of Mayle and Wilson[27]. Using the temperature, density and composition structure versus time for the 1.64 $M_\odot$ model, we calculated the energy loss due to axion production for this core by integrating equations (27) to (30) over time and the volume of the core. Note that this is not a self-consistent calculation of axion production as the energy loss to axions is not fed back into the evolution of the core. However, it does serve as a first estimate of the value for f_a, and allows the dependence of f_a on β to be seen without doing many different self-consistent calculations for different β, which is computationally expensive. We determine a value for f_a by requiring $E_{axion} \le E_\nu = 3 \times 10^{53}$ ergs, in which case the neutrino signal from SN 1987a would have been very different from the signal that was seen. Typical values for the abundances are $X_n \simeq 0.88$ and $X_p \simeq 0.12$ which vary with time and over the structure of the core. In this case we find

$$f_a \ge \left\{ \begin{array}{c} .76(1 + .078\cos 2\beta + .082\cos^2 2\beta)^{1/2} \times 10^{12}\text{GeV} \ ; \ \text{NQM} \\ .34(1 - .16\cos 2\beta + .64\cos^2 2\beta)^{1/2} \times 10^{12}\text{GeV} \ ; \ \text{EMC} \end{array} \right\} \tag{38}$$

Normalizing the NQM value of f_a at $\beta = 0$ to unity, we display f_a as a function of β in figure 1.

Notice that all possible values of β yield constraints on f_a. This is in contrast to previous limits on f_a utilizing only the electron coupling to axions and occurs because the sum of pp, nn, and np rates compensate for variations in any individual rate due to β dependences.

As previously mentioned, the above estimate of f_a suffers from the fact that the energy loss to axions was not fed back into the evolution of the core. To remedy this, we chose to recalculate the evolution for the 1.64 $M_\odot$ core including the axion production rate self-consistently. We took the NQM values of C_{an} and C_{ap} for $\beta = 0$, along with equations (27) to (30), and did two new evolutionary calculations for $f_a = .4 \times 10^{12}$ GeV and $f_a = .8 \times 10^{12}$ GeV. The results of the calculations can be seen in figures 2 and 3. The duration of the electron antineutrino signal is greatly shortened when $f_a = .4 \times 10^{12}$ GeV, terminating in about 5 seconds (a 12 second duration was seen at the Kamiokande-II detector). When $f_a = .8 \times 10^{12}$ GeV, the signal is about 7 seconds in duration. Thus, if the core studied here was anything like the collapsed core of SN 1987a, we can say that $f_a > .8 \times 10^{12}$ GeV for the above axion coupling model, or the neutrino signal would have been very different from that seen.

The variation of f_a with β for self-consistent calculations should not be much different from that seen in figure 1, and taking the critical value of f_a to be $.8 \times 10^{12}$ GeV for the $\beta = 0$, NQM case using Iwamoto's matrix element (26), we show in table 1, the sensitivity of the critical f_a value to other assumptions, such as using the Pantziris and Kang[25] matrix element instead of that calculated by Iwamoto,[18] or using the EMC rather than the NQM values of C_{an} and C_{ap}. All are more restrictive than the neutron star cooling limits and the red giant limits and all approach the cosmological density limits of $\lesssim 4 \times 10^{12}$ GeV. Note that a lower limit of $f_a \gtrsim 8 \times 10^{11}$ GeV corresponds to an upper limit on the mass of the axion of 9×10^{-5} eV.

Uncertainty in our calculations also arises from the initial model for SN 1987a and the assumed supra-nuclear density equation of state employed. Mayle and Wilson[27] studied three different models for SN 1987a. They found only one of these models (a 1.64 $M_\odot$ core) gave good agreement with the observed neutrino detections. One of the other models studied by Mayle and Wilson, a 1.27 $M_\odot$ core, emitted neutrinos for about 6 seconds before the luminosity fell to values unmeasurable at the Earth. The peak temperature reached in this model was 35 MeV, and although this core would not explain the neutrino detection, it does give an estimate of the range in temperatures seen in collapsing cores, when compared with the peak temperature of 70 MeV seen in the 1.64 $M_\odot$ core. Iwamoto's matrix element (26) really should be written using an effective nucleon mass, m_N^*, rather than with m_N, the low-density nucleon mass ($\sim$ 940 MeV). Most of the axion emission comes from regions of the core where densities are about four times normal nuclear matter density (2.5×10^{14} g/cc). In order to see the importance of neglecting the difference between the effective nucleon mass and the low-density nucleon mass, we compared the equation of state used in the study reported in Mayle and Wilson[27] with an equation of state that included a value of m_N^* given by $m_N^* = m_N/(1 + bp)$ with $b > 0$, along with a relativistic Fermi gas treatment of the nucleons (Δ particles and pions were also included). For a given energy density and matter density, the temperature increases when the effective nucleon mass is used ($m_N^* < m_N$). It is found that this rise in temperautre half compensates for the smaller m_N^* in the formula for the energy loss rate (i.e. $m_N^* T^6 \sim$ constant). Thus equation of state uncertainties may contribute about a factor of 2 uncertainty in axion energy loss.

We have shown that the detection of neutrinos from SN 1987a constrains axion couplings to values approaching their upper bounds from cosmological density arguments. These results make it unlikely that any axion search experiment sensitive only to stronger couplings will succeed. It also means that axion models must be tuned to escape simultaneously our lower bounds and the cosmological density limit upper bound.

The supernova limits come from the nucleon-axion couplings as opposed to the electron and photon axion couplings used in the red giant limits.[13] However most axion models have quark-axion couplings since the axion is supposed to be related to the strong CP problem whereas not all conceivable models have electron-axion couplings. However, non-standard axion models may have couplings with $N \neq 6$ in which case our limits should be scaled by $N/6$. The arguments can be extended to any other scalar or pseudo-scalar particle of mass $\lesssim$ 50 MeV which couples to quarks and so would be emitted during the collapse.

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Table I

β	Model	$g_{NN\pi}$	f_a (10^{12} GeV)
0	NQM	I	.80
60°	NQM	I	.76
90°	EMC	I	.45
40°	EMC	I	.33
0	NQM	PK	5.0
60°	NQM	PK	4.8
90°	EMC	PK	2.8
40°	EMC	PK	2.1

Lower limits on f_a scaled from the first entry which uses the results shown in figures 1 and 2. Two values of β, different coupling models (NMQ and EMC), and pion coupling to nucleons (pseudo-vector as assumed by Iwamoto or pseudo-scalar as assumed by Pantziris and Kang) are displayed.

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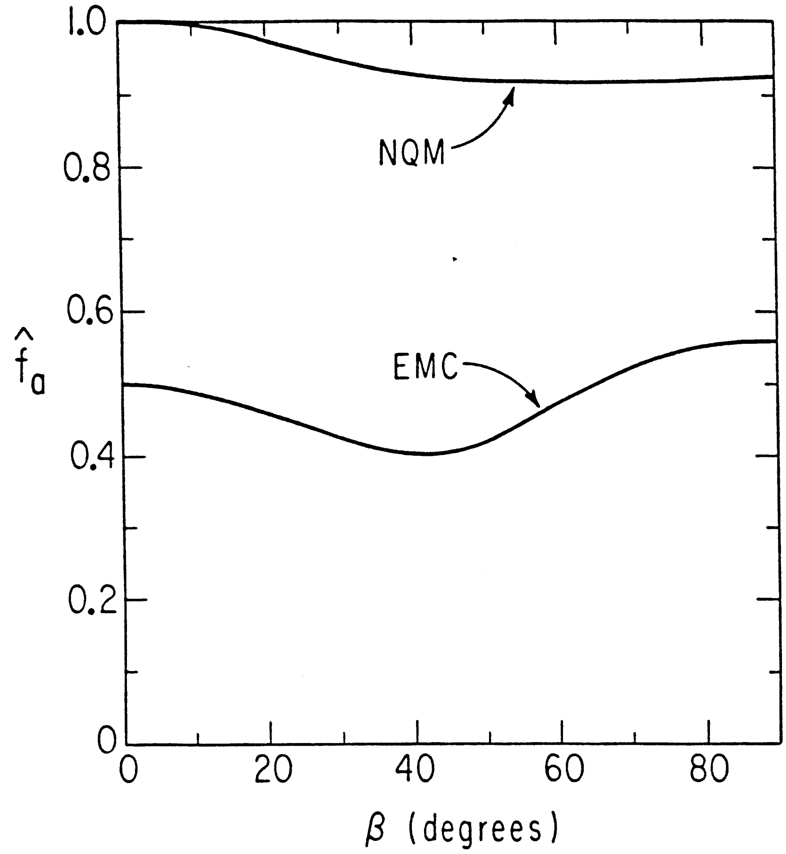
Figure Captions

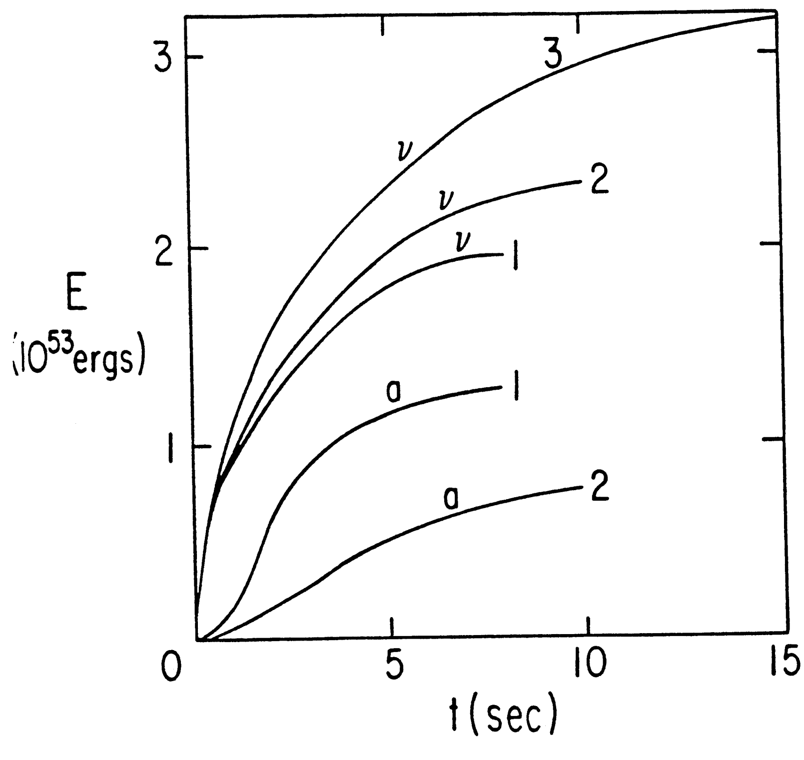
Figure 1. The relative values of $\hat{f}_a$ (normalized so that $\hat{f}_a$ for $\beta = 0$ for the NQM model with Iwamoto's matrix element is unity, i.e. $\hat{f}_a = f_a/0.82 \times 10^{12}$ GeV) from equation (38) required for equal energy loss to axions and neutrinos, versus angle β for NQM and EMC models using Iwamoto's matrix element.

Figure 2. Emitted energies in neutrinos (ν) and axions (a) versus time for three models– model 1: $f_a = .4 \times 10^{12}$, model 2: $f_a = .8 \times 10^{12}$, model 3: $f_a = \infty$. All the above used $\beta = 0$ in the NQM model with Iwamoto's matrix element.

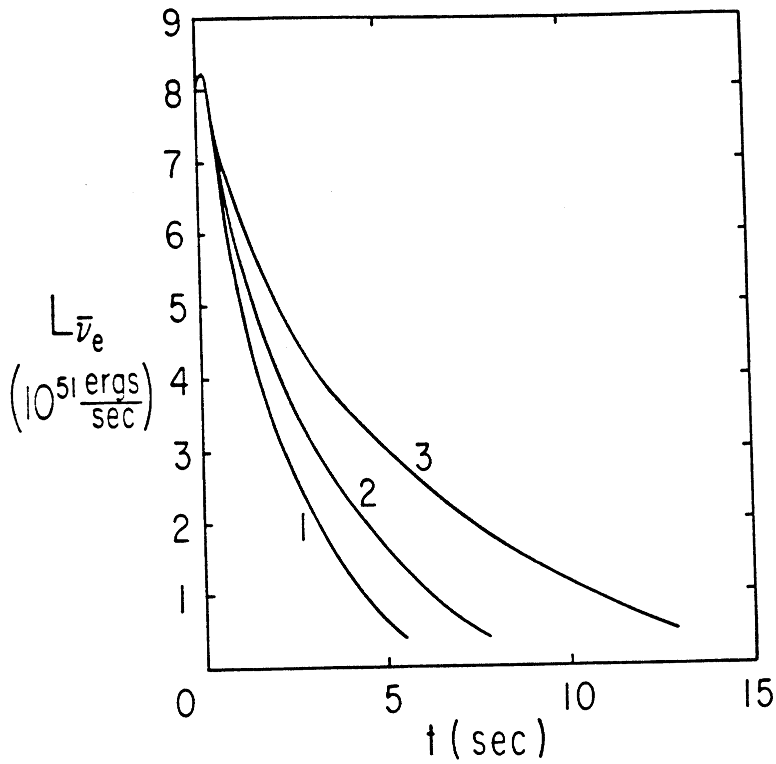
Figure 3. Antineutrino luminosities vesus time for the three models of Figure 2. Observationally, neutrinos were emitted for 12 seconds.



- Figure 1 -



- Figure 2 -



9
8
7
6
5
4
3
2
1
0
5
10
15
$L_{\bar{\nu}_e}$
$\left(10^{51}\frac{ergs}{sec}\right)$
1
2
3
t (sec)

- Figure 3 -