Axions from SN1987A

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Axion emission from SN1987A by nucleon-nucleon axion bremsstrahlung is considered. On the basis of the neutrino observations the axion luminosity must be $\lesssim 10^{53}$ erg s^{-1}. This occurs if (1) axions couple very weakly: $m_a \lesssim 0.75\times 10^{-3}$ eV; or (2) axions couple strongly enough to be "trapped" and radiated from an "axion sphere" with $T_a \lesssim 8$ MeV: $m_a \gtrsim 2.2$ eV. In general, "axion trapping" occurs for $m_a \gtrsim 0.016$ eV. These mass constraints are at best reliable to within a factor of $\simeq 3$.

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The axion is the (hypothetical) pseudo Nambu-Goldstone boson associated with the spontaneous breakdown of the Peccei-Quinn quasisymmetry. In 1977 Peccei-Quinn (PQ) symmetry was proposed to solve the "strong *CP* problem"; ten years later it is still probably the best solution to this important problem.[1] The original axion with symmetry-breaking scale equal to the electroweak scale was quickly ruled out by laboratory experiment and on astrophysical grounds (axion emission from the sun and red giants[2]). To wit, the "invisible axion" was introduced,[3,4] with symmetry-breaking scale $f_a \gg 300$ GeV. Generically, invisible axions are of two types: Dine-Fischler-Srednicki[3] (DFS) and hadronic.[4] The DFS axion has fundamental couplings to all fermions with strength $\sim m_f/f_a$, while the hadronic axion only has fundamental couplings to quarks, and possibly only to heavy, exotic quarks. Both types of axions couple (through anomalies) to photons and nucleons.

Cosmology and astrophysics set stringent bounds to the axion mass. Cosmologically produced axions contribute excessive mass density today, unless[5]

$$m_a \gtrsim (3.6\times 10^{-6}\,\text{eV})\gamma^{-0.85}[\Lambda_{\rm QCD}/(200\,\text{MeV})]^{-0.6}, \tag{1}$$

where $\Lambda_{\rm QCD}$ is the QCD scale parameter, and $\gamma \gtrsim 1$ accounts for any entropy production after axion production: $\gamma \equiv$(entropy per comoving volume after)/(entropy per comoving volume before). Light axions are emitted from stars, thereby affecting stellar evolution, especially that of red giants.[2,6-8] At present, the most stringent limits are $m_a \lesssim 0.01$ eV (DFS)[6] and $m_a \lesssim 3$ eV (hadronic).[7] The latter limit depends upon the axion's anomalous coupling to two photons. In simple, unified models this coupling is fixed; in exotic models it can be much smaller,[9] increasing this upper bound by up to a factor of 15.

In Ellis and Olive[10] axion emission from SN1987A from electron processes was considered. However, the dominant process is nucleon-nucleon axion bremsstrahlung,[11] which is the process considered here. Since the axion-nucleon coupling arises in large part from axion-pion mixing, it is rather model insensitive,[9,12,13] and so the bounds I derive apply to both hadronic and DFS axions. In order that axion emission not cool the newly born, hot neutron star too rapidly (in a time less than or of order of a few seconds), and thereby quench the emission of thermal neutrinos, neutrinos which were observed in at least two underground detectors,[14,15] I require that the axion luminosity Q_a be less than 10^{53} erg s^{-1}. This can occur in one of two ways: first, if the axion coupling is very small, $m_a \lesssim 0.75\times 10^{-3}$ eV; second, if the axion coupling is sufficiently large so that axions are "trapped" and thermalized in the hot core, and the "axion sphere" has a temperature $\lesssim 8$ MeV: This occurs for $m_a \gtrsim 2.2$ eV.

Throughout I will follow Srednicki,[12] but using the normalization conventions of Kaplan[9] and Sikivie.[13] [Note that $(f_a/N)_{\rm Srednicki} = 2(f_a/N)_{\rm Kaplan,Sikivie} \equiv f_a/N$, where N is the color anomaly of the PQ symmetry.] The axion mass and symmetry-breaking scale are related by $m_a \simeq (0.62\text{ eV})/[(f_a/N)/(10^7\text{ GeV})]$. The effective interaction Lagrangean of the axion with electrons, nucleons, and photons is

$$\mathcal{L}_{\rm int} = ig_{aee}(\bar e\gamma_5 e)a + ig_{aNN}(\bar n\gamma_5 n)a + ig_{aPP}(\bar p\gamma_5 p)a + g_{a\gamma\gamma}a\mathbf{E}\cdot\mathbf{B}, \tag{2}$$

where a is the axion field. The couplings are

$$g_{aee} = \{X'_e/N + (3\alpha^2/4\pi)[E\ln(f_a/m_e)/N - 1.93\ln(\Lambda_{\rm QCD}/m_e)]\}m_e/(f_a/N),$$

$$g_{a\gamma\gamma} = (\alpha/2\pi)(N/f_a)(E/N - 1.93),$$

$$g_{ann} = [(-F_{A0}-F_{A3})(X'_d/2N - 0.18) + (-F_{A0}+F_{A3})(X'_u/2N - 0.32)][m/(f_a/N)],$$

$$g_{app} = [(-F_{A0}-F_{A3})(X'_u/2N - 0.32) + (-F_{A0}+F_{A3})(X'_d/2N - 0.18)][m/(f_a/N)].$$

Here X_i' are the PQ charges of the electron and u and d quarks, $\alpha \simeq 1/137$, E is the electromagnetic anomaly of the PQ symmetry ($=8N/3$ in simple unified models), m_e is the electron mass, m is the nucleon mass, and F_{A0} and F_{A3} are the axial-vector isoscalar and isovector pion-nucleon couplings. Experiment suggests that $F_{A3} \simeq -1.25$ and theory that $F_{A0}=0.6F_{A3} \simeq -0.75$ (see Ref. 12, and Ref. 29 therein). For the DFS axion, $N=6$, $X_e'/N=\cos^2\beta/3$, $X_u'=1-\cos2\beta$, and $X_d'=1+\cos2\beta$, where β parameterizes the relative sizes of the "up" and "down" PQ vacuum expectation values.[12] For the hadronic axion, $X_e'=X_u'=X_d'=0$. The coupling of the hadronic axion to the electron arises only through radiative corrections, and the couplings to nucleons only through axion-pion mixing.

For axion emission from the supernova we will only be interested in the axion-nucleon couplings, $g_{ann} \simeq [(X_d'/N - X_u'/4N) - 0.20]m/(f_a/N)$ and $g_{app} \simeq [(X_u'/N - X_d'/4N) - 0.55]m/(f_a/N)$. Lacking precise knowledge of X_u'/N, X_d'/N, F_{A0}, and F_{A3}, where necessary I take $g_{app} \simeq g_{ann} \simeq 0.5m/(f_a/N)$, for both types of axions. [I note that for $X_u'/N=0.64 \simeq \frac{2}{3}$ and $X_d'/N=0.36 \simeq \frac{1}{3}$ both g_{ann} and g_{app} actually vanish.]

SN1987A confirmed astrophysicists' most cherished belief about type-II supernovae[16]—that most of the $\simeq 3\times10^{53}$ ergs of gravitational binding energy released during core collapse is carried off by neutrinos. If we assume that all neutrino species were emitted in roughly equal numbers, the detection of $\bar{\nu}_e$'s by the Kamiokande II[14] and Irvine-Michigan-Brookhaven[15] (IMB) collaborations indicates that thermal neutrinos with a temperature of $\simeq 4$ MeV carried off roughly a few times 10^{53} ergs from the supernova.[17] Since the neutrino pulse lasted a few seconds, the inferred neutrino luminosity is $\simeq 10^{53}$ erg s^{-1}.

According to the generally accepted, and now basically confirmed, theory of core collapse,[18] a type-II supernova is initiated when the $\simeq 1.4M_\odot$ Fe core of a massive star collapses (on a time scale of milliseconds). The collapse is halted when the core reaches a few times nuclear density $[(6\text{–}10)\times10^{14}$ g cm$^{-3}]$. The hydrodynamic shock resulting from the core bounce propagates outward, eventually leading to the optical fireworks. Because of the very high densities, neutrinos are trapped in the hot core ($T \simeq 30$–70 MeV), and are radiated from a "neutrinosphere" $[R \simeq (2\text{–}3)\times10^6$ cm$]$ where the density is $\simeq 10^{12}$ g cm^{-3}, and the temperature is $\simeq 4$ MeV. Neutrino emission cools the core in a few seconds.

Shortly after collapse the hot core is almost isentropic with an entropy per baryon $s/n_b \simeq 2$. For $\rho_{14} \gtrsim 2$, $s/n_b \propto T\rho_{14}^{-2/3}$, while for $\rho_{14} \lesssim 2$, $s/n_b \propto T\rho_{14}^{-1/3}$ [throughout, $\rho_{14} \equiv \rho/(10^{14}$ g cm$^{-3})$, $\hbar = c = k_B = 1$]. The inner core, which contains most of the mass, has roughly constant density. During the first few seconds, before the star's lepton number is carried off by ν_e's, the core should have roughly equal numbers of neutrons and protons, with number density $n \simeq (1.8\times10^{-3}$ GeV$^3)[\rho_{14}(0)/8]$. On the basis of the above, I adopt the following simple (and transparent) model for the newly formed neutron star: mass $\simeq 1.4M_\odot$; central density $\rho_{14}(0) \simeq 8$; radius of the constant-density region $R \simeq (10^6$ cm$)[\rho_{14}(0)/8]^{-1/3}$; $T \simeq (20$ MeV$)c_1\rho_{14}^{1/3}$ ($\rho_{14} \lesssim 2$), $\simeq (15$ MeV$)c_1\rho_{14}^{2/3}$ ($\rho_{14} \gtrsim 2$). As we will see, axion emission depends on the temperature and central density; $\rho_{14}(0)$ and c_1 allows me to display the dependence of the results upon the central density and core adiabat.

If axions exist, both axion and thermal neutrino emission will cool the core. Under the conditions that pertain the dominant axion emission process is nucleon-nucleon axion bremsstrahlung. Since neutrinos were observed to have come from the supernova over a time interval of 5–10 s, axions alone should not cool the core in a time less than this. If the axion luminosity were, say, $\gtrsim 10^{54}$ erg s^{-1}, axions would cool the core in less than a second—clearly inconsistent with the neutrino observations. On the other hand, if the axion luminosity were, say, $\lesssim 10^{52}$ erg s^{-1}, axion emission would have only a slight effect on the cooling of the core. At the intermediate luminosity of $\simeq 10^{53}$ erg s^{-1}, axions should affect the cooling significantly, perhaps enough to be inconsistent with the observation of neutrinos. I shall use 10^{53} erg s^{-1} as the maximum permissible axion luminosity, and note that the constraints scale only as the square root of this luminosity. I will not explicitly consider the "back reaction" of axion cooling upon the model of core cooling.

Axion emission from hot neutron stars through nucleon-nucleon axion bremsstrahlung has been calculated in the degenerate limit by Iwamoto.[19] In our case the Fermi momentum is $p_F \simeq (190$ MeV$)\rho_{14}^{1/3}$; with the T-ρ relation, it follows that $\epsilon_F/(3T/2) \simeq 0.8c_1^{-1}$ ($\rho_{14} \gtrsim 2$), $\simeq 0.6c_1^{-1}\rho_{14}^{1/3}$ ($\rho_{14} \lesssim 2$): A newly born neutron star is not strongly degenerate. Using the matrix element computed by Iwamoto,[19] I have calculated the axion bremsstrahlung cross section in the nonrelativistic (NR), nondegenerate limit:

$$\langle\sigma\rangle = (3/80\pi^3)(T/m)^2 f^4 g_i^2 m^2/m_\pi^4 \simeq (1.2\times10^{-27}\,\text{cm}^2)g^2[T/(1\,\text{GeV})]^2, \qquad (3)$$

where m_π is the pion mass and $f \simeq 1$ is the pion-nucleon coupling. The cross section has been averaged both thermally and over initial spins, and a factor of $\frac{1}{4}$ has been included to account for identical particles in the initial and final states. I have also made the approximation $3mT \gg m_\pi^2$. Here g_i^2 is the appropriate axion-nucleon coupling squared: g_{ann}^2 for $n+n \rightarrow n+n+a$; g_{app}^2 for $p+p \rightarrow p+p+a$; and $\simeq 2(g_{ann}^2+g_{app}^2)$ for $p+n \rightarrow p+n+a$ (here, the extra factor of 4 to "undo" the previous factor of $\frac{1}{4}$). From the cross section it is simple to compute the axion luminosity from the core[20]:

$$Q_a = n^2\langle\sigma|v_1-v_2|\rangle E_a V \simeq (5.1\times10^{72}\,\text{erg s}^{-1})[\rho_{14}(0)/8]^{10/3}c_1^{7/2}g^2, \qquad (4)$$

where $E_a \simeq 3T$ is the average energy per axion emitted, $V \simeq (4.9\times 10^{59}\ \mathrm{GeV}^{-3})[8/\rho_{14}(0)]$ is the core volume, and $g^2 = 3(g_{ann}^2 + g_{app}^2)$ accounts for all the processes mentioned above.

On the basis of this simple model and with use of the limit to the axion luminosity, $Q_a \lesssim 10^{53}$ erg s^{-1}, we obtain the limit $g \lesssim 1.4\times 10^{-10}[\rho_{14}(0)/8]^{-5/3}c_1^{-7/4}$. Furthermore, if we assume that $g_{ann} \simeq g_{app} \simeq \frac{1}{2} m/(f_a/N)$, so that $m_a \simeq 5.4g$ MeV, we obtain the bound

$$m_a \lesssim 0.75\times 10^{-3}[\rho_{14}(0)/8]^{-5/3}c_1^{-7/4}\ \mathrm{eV} \qquad (5)$$

or $f_a/N \gtrsim 8\times 10^9$ GeV. While there is consensus as to the postcollapse adiabat $(s/n_b \simeq 2 \Longrightarrow c_1 \simeq 1)$,[18] the central density depends upon the equation of state at supernuclear densities. And the central density determines the central temperature, $T(0) \propto \rho_{14}(0)^{2/3}$. Changing $\rho_{14}(0)$ by a factor of 2 changes $T(0)$ by a factor of 1.6, and our bound by a factor of 3. Changing our criterion for the maximum axion luminosity by a factor of 2 only modifies our bound by a factor of $\sqrt{2}$.

Next, we must consider axion reabsorption to check our implicit assumption that once emitted, axions just "stream out." The axions produced should have an approximately thermal spectrum. From the Boltzmann equation it follows that the mean free path l for a thermal distribution of axions is

$$l^{-1} \simeq n^2 \langle \sigma | v_1 - v_2 | \rangle / (T^3/\pi^2). \qquad (6)$$

To assess the importance of reabsorption I compare l to the size of the core, $d \simeq V^{1/3} \simeq 2\times 10^6$ cm: $d/l \simeq 1.1\times 10^{17} g^2 [\rho_{14}(0)/8]^4 c_1^{7/2}$; I find that $d/l \gtrsim 1$ for $g \gtrsim 3.0\times 10^{-9}$. That is, for $m_a \gtrsim 0.016$ eV axions become "trapped" and thermalize in the core.

In the "trapped regime" ($m_a \gtrsim 0.016$ eV) there will be an "axiosphere" with temperature T_a and radius r_a determined by the condition $\tau_a \simeq \frac{2}{3}$, where the axion "optical depth" τ_a is given by $\tau_a = \int_\infty^{r_a} dr/l$. To compute τ_a one needs to know $\rho(r)$ and $T(r)$ outside the constant-density inner core. In the spirit of our simple model I assume $\rho_{14} = (r/r_{14})^{-n}$ ($r_{14} \simeq 1.5\times 10^6$ cm, $n \simeq 3$–7), and as before, $T \simeq (20\ \mathrm{MeV}) c_1 \rho_{14}^{1/3}$. Then it follows that

$$\tau_a \simeq 1.8\times 10^{15} c_1^{-6} g^2 \frac{[T_a/(20\ \mathrm{MeV})]^{5.5-3/n}}{11n/6 - 1},$$

or[21]

$$T_a \simeq (3.1\times 10^{-2}\ \mathrm{MeV}) c_1^{12/11} g^{-4/11} \zeta^{-2/11}, \qquad (7)$$

where $\zeta = [T_a/(20\ \mathrm{MeV})]^{-3/n}/(11n/6 - 1)$ and is $\simeq 0.1$–0.7. In the trapped regime ($m_a \gtrsim 0.016$ eV), thermal axions (of temperature T_a) are radiated from the axiosphere. Our criterion $Q_a \lesssim 10^{53}$ erg s^{-1} translates to $T_a \lesssim 8$–10 MeV (depending upon the radius of the axiosphere). This constrains g to be $g \gtrsim 4.0\times 10^{-7} c_1^3 (3\zeta)^{-1/2}$ (taking $T_a \lesssim 8$ MeV). That is, for 2.2 eV $\gtrsim m_a \gtrsim 0.016$ eV axions thermalize in the core and are radiated from the axiosphere with luminosity greater than 10^{53} erg s^{-1}. Because of their strong trapping in the core and correspondingly lower axiosphere temperature, axions with $m_a \gtrsim 2.2c_1^3$ eV have a luminosity $\lesssim 10^{53}$ erg s^{-1} and are permissible. For the DFS axion $m_a \gtrsim 2.2$ eV is certainly ruled out.[6] However, for the hadronic axion, $m_a \gtrsim 2.2$ eV may just be allowed (see the introduction of this Letter), especially when the uncertainties of this and the red-giant limit[7] are taken into account. In addition, relic hadronic axions from the "big bang" of about this mass may actually be detectable from their decays into two photons.[22]

In summary, I find that axion emission from SN1987A restricts the axion mass to be $m_a \lesssim 0.75\times 10^{-3}$ eV, or $m_a \gtrsim 2.2$ eV. However, I must emphasize the multitude of uncertainties underlying this bound: (1) Axion emission from the hot core necessarily depends upon a *theoretical* model for the hot core, which itself depends critically upon the equation of state at supernuclear densities. In particular, as noted, the limit varies as $\rho_{14}(0)^{-5/3}$. (2) The matrix element for nucleon-nucleon axion bremsstrahlung was computed by our assuming NR nucleons and one-pion exchange, and ignoring many-body effects (see Ref. 19, and Ref. 11 therein). Since the pion-nucleon coupling is $\simeq 1$ and the densities are supernuclear, these assumptions are questionable. (3) The criterion $Q_a \lesssim 10^{53}$ erg s^{-1} is also subject to question. In principle, one should take into account the "back reaction" of axion emission on a "standard model" of core cooling and neutrino emission, and then compare the resulting neutrino event rate with the observations.[14,15] The lack of a standard model makes this impractical. In any case, as noted, the limits only depend upon $Q_a^{1/2}$. In sum, I believe that the limits are uncertain by at least a factor of 3, and perhaps an order or magnitude.

Raffelt and Seckel[11] and Mayle *et al.*[23] have also considered axion emission from SN1987A by nucleon-nucleon axion bremsstrahlung. Using Iwamoto's rates and a more detailed cooling model which takes into account the effect of axion emission on the core itself, Raffelt and Seckel[11] obtain a similar bound: $f_a/N \gtrsim 10^{10}$ GeV, in the "free-streaming regime." Using very detailed collapse models and Iwamoto's rates, Mayle *et al.*[23] obtain a slightly more restrictive limit for the DFS axion: $m_a \lesssim 0.9\times 10^{-4}$ eV. While these authors[11,23] have noted the possibility of an allowed mass range where axions are trapped, they have not entirely addressed this issue quantitatively.

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Note added.—We have now computed axion emission for arbitrary degeneracy, and find that the nondegenerate rate used here is a good approximation.[24]

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[20]Correcting Iwamoto's axion emission rates by a factor of $\frac{1}{2}$ for identical particles in the final state, I find that his rates are about a factor of 20 higher than mine for $\rho_{14} \simeq 8$.

[21]In addressing the issue of trapping and in computing the temperature of the axiosphere, I have used my NR, nondegenerate cross section, and assumed that the nuclear composition is equal numbers of free neutrons and protons. At these densities and temperatures ($\rho_{14} \simeq 10^{-2}$, $T \simeq 10$ MeV), $\epsilon_F/(3T/2) \simeq 0.09$, so that the nondegenerate regime applies. Shortly after core bounce ($\gtrsim 0.5$ s), the nuclear composition at the densities of interest, $\rho_{14} \gtrsim 10^{-2}$, should be free nucleons (see Lattimer *et al.*, Ref. 18), with free neutrons outnumbering protons as time goes on. In the spirit of my simple model I will neglect this last fact.

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