Emission and Detectability of Hadronic Axions from SN 1987A

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We show that axions emitted by SN 1987A, with coupling strengths to nucleons in the range $9\times10^{-7} \lesssim g_{aN} \lesssim 1\times10^{-3}$, would have produced an unacceptably large signal at the Kamiokande proton-decay experiment. Axions with $10^{-10} \lesssim g_{aN} \lesssim 1.4\times10^{-6}$ would have altered the neutrino flux from that observed, so using results from SN 1987A alone, one may nominally rule out the entire range $10^{-10} \lesssim g_{aN} \lesssim 10^{-3}$. Uncertainties in our estimates prevent us from categorically excluding a small window around $g_{aN} \simeq 10^{-6}$, which could have interesting consequences for supernova theory and cosmology.

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The potential importance of axions to particle physics, astrophysics, and cosmology has led to their continued study. Severe constraints are placed on axion properties by arguments concerning their potential emission from stars at various stages of evolution, and particularly from type-II supernova (see Ref. 1, for a review). Soon after the detection of neutrinos[2,3] from SN 1987A it was realized that the $\sim$10-sec neutrino pulse ruled out the existence of "exotic" particles that would have carried off a significant portion of the supernova energy. If axions exist, the SN 1987A neutrino data require axion-nucleon couplings to be[4-6] either (a) $g_{aN} \lesssim 10^{-10}$, or (b) $g_{aN} \gtrsim$ (a number thought to be $\sim 10^{-7}$-10^{-6}). The upper limit (a) represents the strength at which axions produced in the hot core would contribute significantly to its energy losses. The restriction breaks down at a stronger value of the coupling, limit (b), because strongly coupled axions produced in the core of a supernova are reabsorbed before they escape the hot proto-neutron star. Under such conditions, the spectrum of escaping axions is nearly that associated with bosons radiated from a blackbody sphere of radius r_a and temperature $T_a = T(r_a)$. The radius of the axionsphere, r_a, is determined by setting the absorption probability $\tau_a \equiv \int_{r_a}^{\infty} [l_a(r)]^{-1} dr$ to $\frac{2}{3}$. The mean free path for axions, l_a, depends upon their coupling to nucleons as $l_a \sim 1/g_{aN}^2$, and thus the radius of the axion sphere is implicitly a function of g_{aN}. If the luminosity of axions, L_a, dominates that of neutrinos, L_ν, then the corresponding value of g_{aN} is still excluded. As g_{aN} is increased, however, r_a also increases (since less material overburden is required for the same absorption probability) and the luminosity $L_a \sim r_a^2 T_a^4$ decreases. Eventually, $L_a < L_\nu$, and the implied g_{aN} is acceptable.

Turner[5] has used this argument to conclude that axion-nucleon couplings $g_{aN} \gtrsim$ (a few$\times10^{-7}$) are allowed, provided the axion is "hadronic," i.e., its couplings to nonhadronic particles are all small. The proviso allows one to avoid the stellar-evolution restrictions reviewed in Ref. 1. Specific models of hadronic axions, e.g., the Kim-Shifman-Vainshtein-Zakharov axion,[7] have in fact been constructed.

In this paper we show that even hadronic axions with large g_{aN} are problematic because the flux of supernova axions would have been detected at the Kamiokande II proton-decay experiment (KII). Axions can induce nuclear excitations in oxygen via the reaction $a\,^{16}\mathrm{O} \rightarrow {}^{16}\mathrm{O}^*$. The nuclear deexcitation often includes γ rays with energies of 5–10 MeV, which trigger the KII detector with an efficiency similar to that for electrons in the same energy range. To substantiate these statements we have reexamined the strong-coupling cooling argument, calculated the axion absorption cross section on ^{16}O, and estimated the efficiency with which absorption events trigger the KII detector. Together, these ingredients result in an expected signal of KII that, while depending on the parameters in the axion-nucleon Lagrangian and our treatment of SN 1987A, almost always exceeds what was actually seen.

We use the conventions adopted by Raffelt[1] to specify our axion model. The axion-nucleon interaction is given by

$$\mathcal{L}_{aNN} = \frac{1}{2f_a} \bar{\psi}_N \gamma^\mu \gamma^5 (c_0 + c_1 \tau_3) \psi_N \partial_\mu a \,, \tag{1}$$

where $c_0 = \frac{1}{2}(c_p + c_n)$ and $c_1 = \frac{1}{2}(c_p - c_n)$ are isoscalar and isovector coefficients. Axion-nucleon couplings are given by $g_{ap} = c_p m_p / f_a$ and $g_{an} = c_n m_n / f_a$. We define the quantity $g_{aN} \equiv (1/\sqrt{2})(g_{ap}^2 + g_{an}^2)^{1/2}$ as a convenient measure of the total coupling.

For densities and temperatures appropriate to supernova, the dominant axion-emission mechanism is nucleon-nucleon bremsstrahlung $NN \rightarrow NNa$. Similarly, the main absorption process is $NNa \rightarrow NN$. We calculate the rate for this process in the nonrelativistic one-pion-exchange approximation. The thermal averaged absorption rate is

$$\frac{1}{l_a} = \frac{2\pi^{3/2}}{\zeta(3)} \frac{\alpha_\pi^2 \alpha_a \rho_N^2}{T_a^{1/2} m_N^{9/2}} F \,, \tag{2}$$

where $\alpha_\pi = g_{\pi NN}^2/4\pi \simeq 15$, $m_N = 0.939$ GeV, $\alpha_a \equiv 1/4\pi f_a^2$, and ρ_N is the nucleon number density. F is a factor of order unity that depends upon the structure of the supernova and details of the axion couplings to nucleons. It may be written as a sum of terms involving pp, nn, and pn scattering,

$$F = c_p^2 x_p^2 F_p + c_n^2 x_n^2 F_n + 4x_p x_n [c_0^2 F_0 + (c_0^2 + c_1^2) F_2] , \quad (3)$$

where x_p and x_n are the local proton and neutron fractions in the supernova. The individual F's are thermal averages over linear combinations of the expressions $[k^2/(k^2+m_\pi^2)]^2$, $[l^2/(l^2+m_\pi^2)]^2$, $k^2 l^2/(k^2+m_\pi^2)(l^2+m_\pi^2)$, and $(\mathbf{k}\cdot\mathbf{l})^2/(k^2+m_\pi^2)(l^2+m_\pi^2)$, where $\mathbf{k}$ and $\mathbf{l}$ are the three-momentum transfers in the direct- and exchange-scattering diagrams. We refer the reader to Ref. 8 for the exact expressions for the matrix elements but note that we have numerically evaluated the thermal averages, including the effects of a finite pion mass $m_\pi = 138$ MeV.

Having determined l_a as a function of T_a, we obtain the radius of the axion sphere by scaling to the neutrino mean free path and to the neutrino-sphere radius r_ν. Specifically, we set the absorption probability at the axion sphere $\tau_a \equiv \int_{r_a}^\infty l_a^{-1} dr = \int_{r_\nu}^\infty \bar\sigma \rho_N dr \equiv \tau_{\bar\nu_e}$. The effective thermal-averaged cross section that determines the neutrino mean free path is $\bar\sigma \equiv x_p \langle \sigma_{CC} \rangle + \langle \sigma_{NC} \rangle$. The neutral-current scattering cross section σ_{NC} is $\sim \frac{1}{3}$ of the charged-current absorption cross section σ_{CC}. By scaling to the $\bar\nu_e$ flux from a numerical supernova model that provides a reasonable fit to the KII and IMB data (model 56 of Burrows[9]) we hope to eliminate most of the uncertainty intrinsic in our particular choice of supernova parameters.

The integrals defining τ_a and $\tau_{\bar\nu_e}$ require the density and temperature profiles of the supernova. Following Turner,[5] we choose particularly simple expressions $\rho_N = \rho_N(r_\nu)(r_\nu/r)^p$, $T_N = T(r_\nu)[\rho_N/\rho_N(r_\nu)]^{1/3}$. The radius r_a is then fixed by the relation

$$\frac{r_a}{r_\nu} = \left(\frac{2\pi^{3/2}}{\zeta(3)} \frac{\alpha_\pi^2}{T_a^{1/2} m_N^{9/2}} \frac{\rho_N(r_\nu)}{\bar\sigma} \frac{p-1}{2p-1} \alpha_a F \right)^{6/(11p-6)} , \quad (4)$$

which we solve iteratively to handle the temperature dependence in F. We use supernova parameters $p=5$, $T_{\bar\nu_e} = 3.87$ MeV, $\rho_N(r_{\bar\nu_e}) = 2.1\times10^{11}$ g cm^{-3}, and $x_p = 0.3$, which yield a neutrino luminosity $L_{\bar\nu_e} = 7\times10^{51}$ ergs/sec. This corresponds to $L_{\bar\nu_e}$ at 1 sec in Burrows's model 56. Since about half the events occur in the first second, it seems reasonable to take these values as time averages over the full supernova.

Once r_a/r_ν is determined, T_a/T_ν and the ratio $R \equiv L_a/L_\nu \approx L_a/6L_{\bar\nu_e}$ follow. We obtain the axion flux ϕ_a by scaling to the neutrino flux, which we take to be $\phi_{\bar\nu_e} = 5.2\times10^9$ cm^{-2} for SN 1987A. The axion flux is then

$$\frac{d\phi_a}{dE_a} = \frac{2}{3\zeta(3)} \frac{\phi_{\bar\nu_e}}{1+R} \frac{r_a^2 T_a^2}{r_\nu^2 T_\nu^3} \frac{y^2}{e^y - 1} , \quad (5)$$

where $y = E_a/T_a$ and the factor $1/(1+R)$ ensures that the energy budget of the supernova model is respected.

Having determined the axion flux, we define a "cooling limit" by imposing the restriction $L_a/L_{\bar\nu_e} < 6$. This allows an axion flux on the order of the total flux from three neutrino flavors. For the case $c_0 = c_1$, this constraint results in the limit $g_{aN} \gtrsim 1.4\times10^{-6}$. This value exceeds Turner's[5] for several reasons: The evaluation of F in Eq. (3) reduces the absorption rate, our scaling procedure requires a factor of 2 lower luminosity, and the nuclear density is a factor of ~ 3 smaller in our supernova model.

We turn now to the signal at Kamiokande. The number of expected axion events is

$$N = N_w \int dE \frac{d\phi_a(E)}{dE} \sum_i S_i \delta(E - E_i) \epsilon_i , \quad (6)$$

where S_i is the strength for absorbing an axion of energy E_i during a transition to nuclear state $|i(J_i^\pi, T_i)\rangle$, ϵ_i is the efficiency with which the decay of this state is detected, and N_w is the number of water molecules in the detector. To obtain the absorption strength we start from Eq. (1) and, following, e.g., Ref. 10, arrive at the expression

$$S_i = \frac{4\pi^2 E_i}{f_a^2} |\langle i(J_i^\pi, T_i) || M_{J_i}^5(E_i) + L_{J_i}^5(E_i) || 0^+, T=0 \rangle|^2 , \quad (7)$$

where $M_{JM}^5(E)$ and $L_{JM}^5(E)$ are the usual Coulomb and longitudinal multipole projections of the nuclear axial-vector current. As a consequence of the derivative coupling in Eq. (1) and the small axion mass, only nuclear eigenstates of unnatural parity can be excited in Eq. (7).

We obtain our eigenstates from an SU(3)-based shell-model calculation that includes components with up to $5\hbar\omega$ of excitation energy. These wave functions have been shown to reproduce a number of sensitive observables and are described in detail in Ref. 11. The predicted giant dipole strength distribution—a quantity of particular relevance to us because of the similarity between photon and axion absorption—is quite accurate. In evaluating Eq. (7) we multiply all $J^\pi T = 1^+0$ matrix elements by a factor 1/1.25 to incorporate the phenomenological quenching of Gamow-Teller strength and scale the $T=1$ piece of $M_{0^-}^5(E)$ by 1.7 to approximate the effects of meson exchange[12] on the isovector "axial charge."

To evaluate the detection efficiency ϵ_i, we first estimate the number and energies of photons produced in the decay of each state $|i\rangle$. A few of the excited states lie below nucleon-emission threshold (12.13 MeV). The most important of these, the 10.96-MeV $J^\pi T = 0^- 0$ state, decays exclusively by two-photon emission. The more highly excited states, however, will emit either a proton or a neutron, leaving a residual ^{15}N or ^{15}O nucleus, which may be left in an excited state and subsequently decay via photon emission. Because the negative-parity (mostly $J^\pi = 0^-$ or 2^-) states in ^{16}O are similar in character to the 1^- states produced in photoabsorption, we take branching ratios to states in ^{15}N and ^{15}O directly from measured photodisintegration cross sections.[13] The positive-parity states of ^{16}O pose a greater problem; we distribute their decays somewhat arbitrarily among the low-lying positive-parity states in ^{15}N and ^{15}O, which introduces an uncertainty of perhaps 20% in the final event rate.

The emitted photons will trigger a detection event with an efficiency $\kappa_\gamma(E)$ that rises with energy beginning at about 5 MeV. The KII Collaboration has not published $\kappa_\gamma(E)$, but it has provided efficiencies for the detection of energetic electrons. Both electrons and photons produce electromagnetic cascades, but since only electrons emit Čerenkov light, they will be detected with somewhat higher efficiency. Here we set $\kappa_\gamma(E) = \kappa_e(\beta E)$, with $\beta = 0.85$ at all energies E. For comparison, when calibrating their detector with a γ-ray source —average energy $\langle E_\gamma \rangle = 8.4$ MeV—the KII Collaboration[14] found a value $\beta = 0.96$. The Sudbury group[15] examined the same issue using Monte Carlo techniques for 6-MeV γ's and found $\beta = 0.73$.

We are now in a position to apply Eq. (6). Figure 1 shows the number of KII events as a function of g_{aN} for

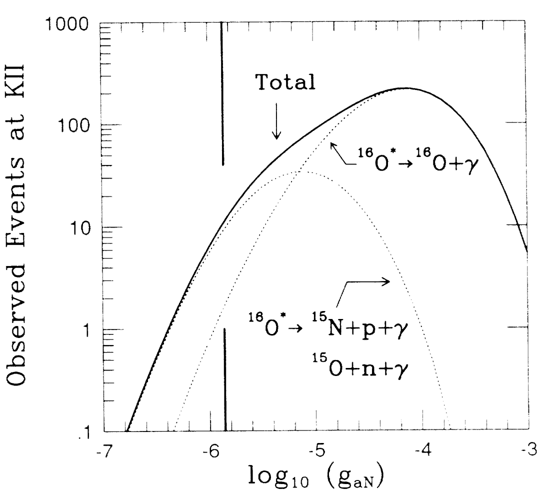

FIG. 1. Expected number of observed events at Kamiokande, as a function of the axion-nucleon coupling g_{aN}. We show the total signal (solid curve) as well as separate contributions (dotted curves) from different decay processes of the excited oxygen nuclei. We estimate that to the left of the heavy vertical line axion energy losses dominate the cooling of the supernova. A small window around $g_{aN} \sim 10^{-6}$ may not be excluded by either direct detection or rapid cooling.

the axion-nucleon couplings $c_0 = c_1$ and the supernova parameters derived from Burrows's model 56. We recall that, for this model, the cooling constraint $L_a < L_\nu$ ruled out $g_{aN} < 1.4 \times 10^{-6}$. At this value we predict 11.1 events at KII, and the number grows as g_{aN} increases. If we assume that the KII data imply fewer than 5 axion detections,[16] then we may conclude from the figure that all values of g_{aN} in the range $9 \times 10^{-7} \lesssim g_{aN} \lesssim 1.0 \times 10^{-3}$ are ruled out. Since this range overlaps that from the supernova cooling argument, it follows that the various constraints from SN 1987A preclude all axion models between those tested in accelerator experiments to the edge of those allowed by cosmological considerations —irrespective of the coupling to electrons or photons.

The reliability of these conclusions depends on their sensitivity to the model parameters. We note first that the two parameters c_0 and c_1 in Eq. (1) do not contribute equally to the detection cross sections—even though they have comparable effects in the supernova. The 10.96-MeV $J^\pi T = 0^- 0$ state of ^{16}O is detected with high efficiency ($\sim 65\%$, as opposed to $\sim 7\%$ for the particle-emission states), and this magnifies the importance of the isoscalar coupling c_0, particularly for large values of g_{aN}.[17] To illustrate this, the figure shows detector events due to $^{16}O^* \to {}^{16}O + \gamma$ and those involving particle emission separately. The particle-emission curve sums over many states and is insensitive to c_0 or c_1 separately. These events dominate the signal for small g_{aN} (large f_a) where the temperature of the axion flux is high enough to excite the 20–25-MeV states without significant thermal suppression. At lower temperatures detection is dominated by the strength below the particle-emission threshold, i.e., by the 10.96-MeV state. Thus, if the isoscalar coupling is made small, the upper bound on g_{aN} from detection moves to around 10^{-5}.

There is also a question of how appropriate it is to use a single set of parameters to describe a dynamic situation. Even if this is acceptable, we might have chosen a slightly different set of parameters. For example, keeping the density gradient fixed, we see from Eq. (4) that we maintain the value of $R \equiv L_a/L_\nu$ by keeping the quantity $\rho_N(r_{\bar\nu_e})\alpha_a F/\bar\sigma$ constant. Different choices for the density or temperature at the neutrinosphere require corresponding changes in α_a (or, equivalently, in g_{aN}^2) and hence in the detection rate. Allowing $\rho_N(r_{\bar\nu_e})$ to be a factor of 3 larger than our fiducial case lowered the number of events at KII to ~ 4, at the new cooling limit.

A more severe problem is our use of a supernova model that does not explicitly include energy losses and transport by axions. As long as $L_a \ll L_\nu$, our scaling procedure should be adequate; however, when $L_a \sim L_\nu$ the use of an "unperturbed" supernova model is suspect. As a result, we cannot exclude a small region of g_{aN} right at the cooling limit. A detailed numerical model of axion transport in the supernova core seems necessary (but may not be sufficient) to settle this issue.[18]

If a small window really does exist between the cooling and detection limits, then the results presented above may have further consequences for supernova theory and cosmology. The cross sections for axion absorption are larger by a factor of 10–100 than those expected from neutral-current interactions of neutrinos.[19] This leads us to speculate that axionic energy absorbed in front of the shock wave could help in creating successful supernova explosions. It also suggests that axion absorption could play an important role in nucleosynthesis, in analogy to the "ν process" explored in Ref. 20.

Another possibility arises if we push parameters so that $f_a = 3\times 10^5$ GeV is still allowed. For example, taking $\beta = 0.73$ as a low value for photon detection efficiency at KII, $T_{\bar{\nu}_e} = 3.5$ MeV, $\rho_N(r_\nu) = 4\times 10^{11}$ g cm^{-3}, $c_0 = 0.35$, and $c_1 = 0.17$, we predict two 10.96-MeV absorption events. This value of f_a yields an axion mass $m_a = 20$ eV, corresponding to an 8.6-sec delay time for 10.96-MeV axions absorbed into the 0^- state. Thus, it becomes possible to attribute the $\sim$8-sec gap at KII to a real delay instead of a moderately unlikely statistical aberration. We note that thermally produced 20-eV axions would also be contribute significantly to the closure density of the Universe, although limits on the UV background[21] at 10 eV require an axion lifetime $\tau_a \gtrsim 10^{24}$ sec, or, equivalently, an anomalously small axion-photon coupling $c_{a\gamma\gamma} \equiv g_{a\gamma\gamma} f_a \lesssim 10^{-6}$.

All such speculation aside, our main point is that "hadronic" axions are directly detectable in water Čerenkov detectors and that the SN 1987A observations at Kamiokande place severe constraints on such models.

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[17]Note that for larger g_{aN} the axion mass increases and so does the time dispersion of the axion flux. This will tend to wash out the signal from the axions absorbed into the broad resonances above the particle emission threshold; however, the 10.96-MeV state is quite narrow. Although that signal will be delayed, it will not be dispersed, and so it should still be detectable against background.

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