Enhanced Supernova Axion Emission and its Implications

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We calculate the axion emission rate from reactions involving thermal pions in matter encountered in supernovae and neutron star mergers, identify unique spectral features, and explore their implications for astrophysics and particle physics. We find that it is about 2 – 5 times larger than nucleon-nucleon bremsstrahlung, which in past studies was considered to be the dominant process. The axion spectrum is also found to be much harder. Together, the larger rates and higher axion energies imply a stronger bound on the mass of the QCD axion, and better prospects for direct detection in a large underground neutrino detector from a nearby galactic supernova.

The axion, a hypothetical particle initially introduced to explain the smallness of the observed CP-violating interactions in QCD [1, 2], is a well-motivated dark matter (DM) candidate [3–5]. Axions produced during inflation would account for the totality of the dark matter in the universe if their mass is in the range from a few µeV to a few 10 µeV [6, 7], the exact value depending on unknown initial conditions. While this observation has motivated ongoing experimental searches for axions in the mass range 2 µeV ≤ m_a ≤ 25 µeV [8, 9], there is interest in axions with higher masses and experimental proposals to discover them [10–12] for two main reasons. First, recent work shows that if DM axions are produced after inflation, their mass needs to be considerably larger to account for DM. When the contribution of topological defects to the axion production is properly accounted for in post-inflationary scenarios studies find that m_a ≥ 25 µeV (see, e.g., [13] and references therein). Recent investigations suggest masses as high as 0.5 – 3.5 meV [14, 15], or even 15 meV [16], depending on the specific axion model. Second, axion masses m_a ≥ 1 – 10 meV are particularly interesting for astrophysics, since these axions can have a noticeable impact on stellar evolution, supernovae, and the cooling of white dwarfs and neutron stars [17–23].

The principle finding of this Letter is that the pion induced axion emission from supernovae (SNe) provides new opportunities to either discover or constrain meV scale axions. We find that it strengthens the SN bound on axions and improves the prospect for both direct and indirect detection of SN axions in the parameter range of interest for particle physics, cosmology, and astrophysics.

The detection of about 20 neutrinos from the core-collapse SN in the Large Magellanic Cloud in 1987, called SN 1987A, continues to provide one of the most stringent bounds on the properties of the QCD axion. Pioneering work in Ref. [24] found that the axion emission due to nucleon-nucleon bremsstrahlung NN → NNa could dramatically alter the early cooling of neutron star born with a fiducial temperature T ≃ 30 MeV and change its neutrino luminosity. Subsequent improvements in the description of the axion emissivity from a SN core, over several years, demonstrated that the suppression of the neutrino luminosity due to axion emission would discernibly alter the observed neutrino events from SN 1987A to provide stringent bounds on the axion nucleon couplings [25–33]. This bound excludes QCD axions with masses in the range 15 meV ≤ m_a ≤ 10 keV [33].

In all of these studies, the nucleon-nucleon bremsstrahlung reaction NN → NNa was assumed to be the dominant channel for the axion production in a SN core. The role of the pion induced reaction, π^-p → na was first discussed in Refs. [28, 29], and in Ref. [30] was found to make the dominant contribution for a sufficiently high pion abundance. However, initial estimates suggested that the thermal pion population was too small for the pion reaction to be competitive [17]. For this reason, pions and reactions involving pions in SNe have been largely ignored.

A recent study demonstrated that the strong interactions enhance the abundance of negatively charged pions [34]. The study found that this enhancement can be reliably calculated for a wide range of density and temperature encountered in the SNe core using the virial expansion. Motivated by this result, and by the large suppression of the bremsstrahlung rate found in [33], we revisit the calculation of the axion emissivity due to the reaction π^-p → na to assess its impact. We find that for pion densities predicted by the virial expansion, the pion induced reaction dominates over the nucleon bremsstrahlung process over a wide range of ambient conditions and has important implications for the axion bounds derived from SN 1987A and direct detection in next-generation experiments. The enhanced emission also has implications for astrophysics of both core-collapse and neutron star mergers. In what follows we describe our finding and these

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is the contribution in the absence of interactions, and the second virial coefficients \( b_n^{\pi-} \) and \( b_p^{\pi-} \) include the contributions due to \( \pi^- \) interactions with neutrons and protons, respectively. It is adequate to retain only leading term in the virial expansion when the fugacity of pions \( z_{\pi} \) and neutrons and protons denoted by \( z_n = \exp[\beta(\mu_n - m_n)] \) and \( z_p = \exp[\beta(\mu_p - m_p)] \) are small compared to unity. For a wide range of typical conditions encountered in a SN where \( z_{\pi} \ll 1 \) and \( z_p \ll 1 \), Eq. (1) provides a reliable estimate of the pion number density. For typical conditions encountered in the SN, the pion fraction \( Y_\pi = n_\pi / n_B \) where \( n_\pi \) is the pion number density and \( n_B \) is the baryon density was found to be in the range 1-5\% for \( n_B \lesssim n_{sat} \).

To describe reactions involving thermal pions it is necessary to define the relation between the pion energy and its momentum given by

\[
E_\pi(p) = \sqrt{m^2 + p^2} + \Sigma(p), \tag{3}
\]

where \( \Sigma(p) \) is the self-energy of pions at finite temperature and density, and incorporates the effects of pion interactions with nucleons. We employ a model in which the effective interaction between pions and nucleons is directly related to the measured pion-nucleon phase shifts (often called the pseudo-potential) to calculate \( \Sigma(p) \). The model is calibrated to reproduce the model-independent results obtained in the virial expansion and its use in calculating reactions is described in detail in Ref. [34].

The number of axions emitted per unit volume and per unit of time and energy is given by [40]

\[
\frac{d\dot{n}_a}{d\omega_a} = \int \frac{2d^3p_p}{(2\pi)^32m_N} \frac{d^3p_\pi}{(2\pi)^32m_\pi} \frac{2d^3p_n}{(2\pi)^32m_N} \frac{4\pi\omega_a^2}{(2\pi)^4\delta^4(p_f - p_i)|\mathcal{M}|^2 f_p f_\pi (1 - f_n) . \tag{4}
\]

The squared transition matrix element in Eq. (4) is averaged over both initial and final nucleon spins and given by

\[
|\mathcal{M}|^2 = 4g_{aN}^2\gamma_{st}(\omega_a) \left( \frac{g_A}{2F_\pi} \right)^2 |\mathbf{p}_\pi|^2 , \tag{5}
\]

where \( \mathbf{p}_\pi \) is the pion momentum, \( g_A = 1.26 \) the axial coupling and \( F_\pi = 92.4 \) MeV the pion decay constant. The effective axion-nucleon coupling \( g_{aN} \) is defined as

\[
g_{aN}^2 = g_a^2 \left[ \frac{1}{2}(C_{ap}^2 + C_{an}^2) + \frac{1}{3}C_{an}C_{ap} \right] , \tag{6}
\]

where \( g_a = m_N/f_a \), \( m_N \) being the nucleon mass and \( f_a \) the Peccei-Quinn scale. The \( C_{ui} \) are the model dependent \( O(1) \) dimensionless axion-fermion couplings. The couplings have been recently calculated for the KSVZ [41, 42] and the DFSZ [43, 44] models in Ref. [45].

\[1\) We notice a discrepancy in Eq. (6) with respect to the result [28, 30], i.e. a minus sign in front of the \( 1/3C_{an}C_{ap} \) term.
(see [46] for a discussion of these parameters in a large class of axion models). The function $\gamma_{sf}(\omega_\pi) = \omega_\pi^2/\omega_\pi^2 + (\Gamma/2)^2$ in Eq. (5) is a simple ansatz suggested in Refs. [29, 38] to account for the finite lifetime of the nucleon spin due to scattering in the dense medium, and $\Gamma$ is the nucleon spin fluctuation rate. At a fiducial temperature $T = 30\text{ MeV}$ and mass density $\rho = 10^{14}\text{ g/cm}^3$, the calculations in [33, 47] indicate that $\Gamma \simeq 35\text{ MeV}$.

The distribution functions of the different interacting species are the usual Fermi-Dirac or Bose-Einstein distribution,

$$f_i(E) = \frac{1}{e^{(E_i(p_i) - \mu_i)/T} + 1},$$

where the $+$ sign applies to fermions while the $-$ is for bosons, and $\mu_i$ are the chemical potentials for $i = p, n, \pi$. Corrections to the dispersion relations $E_i(p_i)$ of nucleons are incorporated through the equation

$$E_i = m_N + |p_i|^2/2m_N + U_i,$$

where the nucleon effective mass $m_N^*$ and single-particle potentials $U_i$ are obtained from Ref. [34]. The modification to the pion dispersion relation due to its interactions with nucleons is incorporated through Eq. (3) with $\Sigma(p)$ obtained consistently as described in Ref. [34].

The differential axion number luminosity, which is defined to be the total number of axions emitted in a specified energy range per unit time from the SN is obtained by integrating Eq. (4) over the SN volume and is given by

$$\frac{dN_a}{d\omega_a} = \int d^3r \frac{d\hat{n}_a}{d\omega_a}.$$

The energy radiated in axions per unit volume and time, called the axion emissivity, can be calculated directly from Eq. (4) as

$$Q_a = \int d\omega_a \omega_a \frac{d\hat{n}_a}{d\omega_a},$$

where the phase-space integrals in the previous equation can be performed to obtain a simpler expression for pionic processes

$$Q_a^\pi = \frac{g_A^2\sqrt{2m_NT}}{\sqrt{2mN\pi^3}} \left( \frac{g_A}{2F_\pi} \right)^2 \frac{z_\pi z_p}{1 - z_\pi} \int dx_p \frac{x_p^2}{e^{x_p^2 + 2 \pi}} \int dx_\pi \frac{x_\pi^2 e_\pi}{e^{x_\pi + (\Gamma/2T)^2}},$$

with $x_p = |p_p|^2/\sqrt{2m_N\pi^2}$, $x_\pi = |p_\pi|^2/T$, $y_\pi = m_\pi/T$, and $e_\pi = E_\pi/T$. The fugacities $z_\pi$ and $z_p$ were defined earlier. Finally, the total axion energy luminosity is given by

$$L_a = \int d^3r Q_a(r).$$

The enhancement of the axion emission rate due to the pion reaction relative to the bremsstrahlung calculated in [33] can be gauged from Table I where we compare the $\pi N$ and $NN$ axion emissivity at different post-bounce times using ambient conditions taken from the SN model described in [33] at a specific radial location $r = 10\text{ km}$. We estimate the total axion emissivity $L_a$ by assuming average values for $T$ and $\rho$ within the region $r < 12\text{ km}$. This is shown in the last column of the Table. We realize that the axion emissivity is increased by factor of about 4 due to pionic reactions at $t_{pb} = 1\text{ s}$. At later times the pion contribution is less important, the total emissivity is only a factor 2 larger than the one from NN process for at $t_{pb} = 6\text{ s}$.

The more stringent bound on the axion mass implied by the larger emissivity can be estimated using an observation made by Raffelt [48] who found that for

$$\frac{Q_a}{\rho} > 10^{19}\text{ erg g}^{-1}\text{s}^{-1},$$

simulations predicted a significant shortening of the SN 1987A neutrino signal. The axion emissivity is typically calculated at a fiducial density $\rho = \rho_{sat}, T = 30\text{ MeV}$, and proton fraction $Y_p = 0.3$. In Table II we show the bounds derived for the KSVZ axion obtained using the fiducial densities $\rho = \rho_{sat}$ and $\rho = \rho_{sat}/2$ at temperature $T = 30\text{ MeV}$ and proton fraction $Y_p = 0.3$. Since the rates are $\propto m_a^2$, the factor of 4 enhancement in the rate strengthens the axion mass bound by a factor 2. We caution the reader that while this simple estimate captures that trend and the relative importance of the pion reaction, detailed SN simulations with pions will be needed to derive a robust bound.

In addition to increasing the total axion emissivity, the reaction involving pions produces axion with a harder energy spectrum. This is to be expected as these reactions harness the rest mass energy of the pion in the initial state. Fig. 1 compares the axion number luminosity obtained from pionic reactions (solid curve) to those from nucleon bremsstrahlung (dashed curve) for our benchmark axion model at a post-bounce time $t_{pb} = 1\text{ s}$.

The larger axion energies, especially axion in the range $200 - 300 \text{ MeV}$ are particularly interesting for detection in neutrino underground experiments. This is because at these energies we expect a resonant enhancement of the axion-nucleon cross-section due to the $\Delta$ intermediate state. These high energy axions can produce neutral and charged pions in water Cherenkov detectors due to the reactions $a + p \rightarrow p + \pi^0, a + p \rightarrow n + \pi^+, a + n \rightarrow p + \pi^-$, and $a + n \rightarrow p + \pi^-$. The operator structure that describes axion coupling to nucleons is nearly identical to the pion-nucleon coupling, but with $f_a$ replaced by $f_a$. This observation has been used earlier to suggest that the cross-section for the reaction $a + p \rightarrow N + \pi$, $\sigma_{aN} \simeq (F_\pi/f_a)^2\sigma_{\pi N}$ where $\sigma_{\pi N}$ is the cross section for $\pi^0 + p \rightarrow p + n^0$ [18]. In the resonance region, which can be accessed when the axion energy $E_a \simeq 200 - 300\text{ MeV}$, the cross-section $\sigma_{aN} \approx 100$ millibarn. For $f_a = 10^9\text{ GeV}$ ($m_a = 5.7\text{ meV}$), an order of
TABLE I: Axion emissivities $Q_a$ in units of $10^{32} \text{erg cm}^{-3} \text{s}^{-1}$ and luminosities $L_a$ in units of $10^{51} \text{erg s}^{-1}$ for KVSZ model ($C_{ap} = -0.47$; $C_{an} = 0$) and $g_a = m_N/f_a = 10^{-9}$, for different post-bounce times.

| $t_{pb}$ (s) | $\rho$ (10$^{14} \text{g/cm}^3$) | $T$ (MeV) | $Y_e$ | $Q_a^{NN}$ (10$^{32} \text{erg cm}^{-3} \text{s}^{-1}$) | $Q_a^\pi$ (10$^{32} \text{erg cm}^{-3} \text{s}^{-1}$) | $Q_a^{tot}/Q_a^{NN}$ | $L_a$ (10$^{51} \text{erg s}^{-1}$) |
|--------------|-------------------------------|-----------|-------|-------------------|-------------------|-------------------|-------------------|
| 1            | 1.45                          | 37.07     | 0.011 | 1.37              | 4.63              | 4.38              | 4.0               |
| 2            | 2.08                          | 38.93     | 0.016 | 3.28              | 8.87              | 3.70              | 8.10              |
| 4            | 3.10                          | 40.56     | 0.027 | 9.08              | 15.87             | 2.75              | 16.63             |
| 6            | 3.65                          | 39.91     | 0.034 | 12.92             | 14.99             | 2.16              | 18.61             |

Magnitude estimate obtained using the axion luminosity in Table I suggests that about 1000 pions will be produced in a megaton water Cherenkov detector for a SN at 1 kpc.

This intriguing prospect for direct detection of axions from a galactic SN warrants further studies and our work identifies several directions for future research. Most importantly, it motivates rigorous calculations of the cross-section for the process $a+p \to N+\pi$ as this is critical for the pion production in water Cherenkov detectors. Such calculations will also address possible resonant enhancement of the inverse reaction $\pi^-+p \to n+a$ in the SN environment. Further work, which goes beyond the virial expansion in Ref. [34], is needed to assess how the pion abundances increase with density in the SN core. Although our initial estimates suggest an exponential increase of the pion thermal population with density, reliable calculations that can accommodate Bose-Einstein condensation of pions at finite temperature will be needed in this context (for a discussion of meson condensation in SN matter, see Ref. [49]). Ultimately, advanced SN simulations that incorporate the pion contribution to both thermodynamics and reactions will be essential to fully assess the impact of the enhanced axion luminosity and energies that we discuss in this Letter.

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