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**AXION EMISSION FROM SN 1987A :
NUCLEAR PHYSICS CONSTRAINTS**

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ABSTRACT

We re-examine the energy-loss rate from the possible axion emission during the supernova SN 1987A collapse, paying particular attention to the modification of the emission rate due to the process $NN \rightarrow NN\alpha$ by the nuclear medium. The main corrections are ρ -meson exchange, many-body effects and nucleon effective mass. All of these decrease the previous lower bound for f_a substantially and lead to a new lower bound $f_a \geq (0.7-5) \times 10^{10}$ GeV.

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1. Introduction

The observed rate of neutrino emission from the supernova SN 1987A sets a lower limit for the axion coupling strength to the nucleon. This observation has generated considerable interest [1-3] since it leaves only a rather narrow range of possible values for the coupling constant, when used together with cosmological constraints. However, the discussion of axion emission has so far neglected the influence of short-range processes and of many-body renormalization of the coupling constants and masses in the dense matter. It is the purpose of the present letter to establish their influence on the estimate and the uncertainty.

Among the relic particles that can be radiated during the collapse of the star, the axion is the main possible candidate. This particle was introduced initially by Peccei and Quinn [4] to solve the strong CP problem [5], and modified later on by Dine, Fischler and Srednicki [6] ("invisible axion") in order to account for the experimental upper limits on its mass. In this interpretation, the axion arises as an almost massless Goldstone boson associated with spontaneous breakdown of a global symmetry (the so-called Peccei-Quinn symmetry). The present limitations on its mass m_a come from astrophysical and cosmological considerations. It can be expressed in terms of the axion decay constant f_a [6],

$$m_a \approx 7.2 \times 10^{-5} \left(\frac{10^{12} \text{ GeV}}{f_a} \right) \text{ eV}. \quad (1)$$

The upper bound on f_a comes from cosmological constraints in a standard Big-Bang universe [7]

$$f_a \leq 4 \times 10^{12 \pm 0.34} \left(\frac{\Lambda_{\text{QCD}}}{200 \text{ MeV}} \right)^{0.6} \text{ GeV}. \quad (2)$$

In these expressions, we assume that the number of quarks with Peccei-Quinn charges is six. One of the best previous lower limits on f_a is given by the requirement that the helium cycle in red giant stars (and consequently their temperature) is compatible with experimental observations (luminosity and number of red giants). One thus finds [8,9]

$$f_a \geq 1.4 \times 10^{10} \cos^2 \beta \text{ GeV}, \quad (3)$$

where β is an arbitrary angle depending upon the Higgs sector of the low-energy theory. Another lower bound on f_a , independent of β , can be obtained from neutron-star cooling [9,10]

$$f_a \geq 10^9 \text{ GeV}. \quad (4)$$

The lower bound set by the supernova SN 1987A may be more restrictive by one to two orders of magnitude. During the final collapse of the star, neutrinos are emitted when the density is less than about $2 \times 10^{11} \text{ g/cm}^3$ (we recall that ordinary nuclear matter density is $\rho_0 \approx 0.25 \times 10^{15} \text{ g/cm}^3$). Above this density, the core is opaque to neutrino propagation, but the more weakly coupled axion can still escape from the interior, which is at a higher temperature. In this region, the temperature is about 70 MeV, and the density about three times ordinary nuclear matter density.

Under these conditions, the dominant process for the emission of axions is by far nucleon-nucleon bremsstrahlung[9], as indicated schematically in fig.1. Previous estimates of this contribution, taking only π -exchange in the plane wave approximation, led to a lower bound of about 1×10^{11} to 8×10^{11} GeV, depending on whether the nucleons were degenerate or not [1-3]. We shall come back to this point later on.

We emphasize in this letter that a more precise estimate of the axion decay constant from SN 1987A requires a detailed analysis of the conditions under which the axions are emitted. Given that the density of the medium is rather large compared to ordinary nuclear matter density, several corrections well-known in nuclear physics must be included. We shall concentrate in the following on three different aspects. First of all, the elementary amplitude given in fig.1 involves the complete nucleon-nucleon (NN) potential. This is particularly important here, since the dominant contribution comes from the pion-dominated tensor interaction. It is well established from the NN interaction that it is weakened and regularized inside distances of 1 to 1.5 fm, and that this phenomenon is quantitatively described by ρ -exchange (or two-pion exchange in the $J=T=1$ channel). This effect plays an important rôle at high densities. Moreover, many-body effects also renormalize the meson propagators as well as meson-nucleon coupling constants. For π - and ρ -exchange the most important renormalization comes from Δ -hole excitations. Finally, the question of the nucleon effective mass at this high density will be addressed. We believe that all these corrections dominate over various other uncertainties in the numerical evaluation of the energy loss by axions.

In order to discuss these effects at a quantitative level, we implicitly assume that in the last stage of the collapse, the core is still made up of individual nucleons interacting via meson exchange. This is the usual assumption made to derive the nuclear equation of state in the physical conditions of the collapse of the star. This sets constraints on various quantities. The first and second ones concern the transition to a new phase of matter in which the quarks are deconfined or in which chiral symmetry is restored. These two limits may a priori be different. In a simple geometrical model based on the bag model, the critical density for deconfinement is given by $\rho_c = (4\sqrt{2} R_c^3)^{-1}$, where R_c is the confinement radius of the quarks in the nucleon. This corresponds to the situation where the bags just touch. Inverting this relation, we therefore require that the quark core radius of the nucleon be smaller than the confinement radius corresponding to the characteristic density in the medium, $\rho_c \approx 3\rho_0$, i.e. smaller than 0.75 fm. This is actually not a very strong constraint. On the other hand, the pion degree of freedom (and hence the NN potential) can only exist in the region where chiral symmetry is spontaneously broken. As we approach the critical density for chiral

symmetry restoration, the pion decay constant goes to zero while its mass and coupling constant remain almost unchanged [11]. We assume here that this transition occurs at a density larger than $3\rho_0$. The last constraint concerns the occurrence of pion condensation in the nuclear medium from polarization effects. As will be discussed later on, this is a much more important constraint, and determines to a large extent the corrections from Δ -excitations in the medium via pion rescattering.

2. Elementary process

In the elementary axion bremsstrahlung process indicated in fig.1, the simplest mechanism is associated with π -exchange between uncorrelated nucleons. So far it has usually been assumed that this mechanism is the dominant one, leading to the following matrix element squared [10] for axion bremsstrahlung from nn or pp pairs

$$|M_\pi|^2 = \frac{256}{3} m_N^2 \left(\frac{f_{\pi N}^2}{m_\pi^2} \right)^2 g_{aN}^2 \frac{k^4}{(m_\pi^2 + k^2)^2} + \text{E.T.} , \quad (5)$$

where $f_{\pi N}$ is the pseudo-vector πNN coupling constant, $f_{\pi N}^2/4\pi=0.08$, g_{aN} the axion-nucleon coupling constant, m_N and m_π are the nucleon and pion mass respectively. The momentum k is that of the exchanged meson. The exchange terms (denoted by E.T.) are not given explicitly here for simplicity. In the limit where the axion momentum is small, the exchange terms amount to about twice the direct term [12]. This is a rather good approximation for small temperatures (degenerate limit) [3]. We shall take this simple estimate for the nn and pp transition amplitude in the subsequent developments. One should note that as long as one sticks to nucleon intermediate states in fig.1, the π -exchange matrix element squared (5) is identical for pseudoscalar and pseudovector axion-nucleon and π -nucleon vertices, and is consistent with chiral symmetry for s-wave πN scattering.

In the evaluation of the π -exchange contribution, only the tensor component V_T of the NN potential contributes. It is obtained from the π -exchange interaction

$$\begin{aligned} V^\pi(\mathbf{k}) &= - \frac{f_{\pi N}^2}{m_\pi^2} \tau_1 \cdot \tau_2 \frac{\boldsymbol{\sigma}_1 \cdot \mathbf{k} \boldsymbol{\sigma}_2 \cdot \mathbf{k}}{m_\pi^2 + k^2} \\ &= - \frac{1}{3} \frac{f_{\pi N}^2}{m_\pi^2} \tau_1 \cdot \tau_2 \frac{k^2}{m_\pi^2 + k^2} \left[\boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2 + (3 \boldsymbol{\sigma}_1 \cdot \mathbf{k} \boldsymbol{\sigma}_2 \cdot \mathbf{k} - \boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2) \right] \\ &\equiv V_C^\pi(\mathbf{k}) + V_T^\pi(\mathbf{k}) . \end{aligned} \quad (6)$$

The central part V_C cancels out when one takes the various time orderings. As we already mentioned in the introduction, the tensor part is strongly reduced at high momentum transfer (short distances) by the ρ -exchange interaction. Indeed, the relevant contribution of the ρ -meson to the potential writes

$$\begin{aligned}
 V^\rho(\mathbf{k}) &= - \frac{g_{\rho N}^2}{4m_N^2} (1+\kappa_\rho)^2 \tau_1 \cdot \tau_2 \frac{\sigma_1 \times \mathbf{k} \cdot \sigma_2 \times \mathbf{k}}{m_\rho^2 + k^2} \\
 &= - \frac{1}{3} \frac{g_{\rho N}^2}{4m_N^2} (1+\kappa_\rho)^2 \tau_1 \cdot \tau_2 \frac{k^2}{m_\rho^2 + k^2} \left[2 \sigma_1 \cdot \sigma_2 - (3 \sigma_1 \cdot \mathbf{k} \sigma_2 \cdot \mathbf{k} - \sigma_1 \cdot \sigma_2) \right] \\
 &\equiv V_C^\rho(\mathbf{k}) + V_T^\rho(\mathbf{k}),
 \end{aligned} \tag{7}$$

so that the tensor component arises with a minus sign compared to π -exchange. In Eq.(7), $g_{\rho N}$ is the vector ρNN coupling constant, and κ_ρ is the ratio of the tensor to vector coupling.

The matrix element squared is then

$$|M_{\pi+\rho}|^2 = \frac{256}{3} m_N^2 \left(\frac{f_{\pi N}^2}{m_\pi^2} \right)^2 g_{a N}^2 k^4 \left[\frac{1}{(m_\pi^2 + k^2)} - \frac{C_\rho}{(m_\rho^2 + k^2)} \right]^2 + \text{E.T.}, \tag{8}$$

where

$$C_\rho = \left(\frac{g_{\rho N} (1+\kappa_\rho)}{2m_N} \frac{m_\pi}{f_{\pi N}} \right)^2. \tag{9}$$

With the parameters deduced from the Paris potential [13], $C_\rho = 1.67$. Note that the ρ -meson mass is chosen in such a way that it simulates the exact 2π -exchange in the $J=T=1$ channel, which leads to $m_\rho \approx 600$ MeV [13]. The correction to the elementary process from ρ -exchange is expected to be the most important one, all other meson exchange cancelling out in leading order when the different time orderings and exchange terms are included [12].

In order to have a more quantitative estimate, several corrections must now be considered. First of all, one should calculate the transition amplitude with correlated wave functions instead of plane waves, since we know that the NN potential is strongly repulsive at short distances. This leads to a very small value for the correlated wave

function for relative distances less than about 0.5 fm. One can easily take this correction into account by incorporating the correlation function into the definition of the NN potential itself according to

$$\langle \varphi^C | V(r) | \varphi^C \rangle \equiv \langle \varphi^O | V^C(r) | \varphi^O \rangle, \quad (10)$$

where φ^C and φ^O are the correlated and uncorrelated two-body wave functions respectively, with

$$\varphi^C(r) \equiv \varphi^O(r) f(r), \quad (11.a)$$

$$V^C(r) \equiv V(r) f^2(r). \quad (11.b)$$

The correlation function is given by $f(r)$ in Eq.(11). For practical reasons, one takes [12]

$$f^2(r) = 1 - j_0(q_c r), \quad (12)$$

where q_c fixes the scale of the correlations.

In practice, one includes in the phenomenological function $f^2(r)$ also the effects of the various renormalizations of the meson-nucleon vertices (form factors) from the finite size of the nucleon and multi-meson configurations. In the following, we shall take $q_c=800$ MeV/c. It is important to note here that the modification of the potential due to the correlation function is rather small because of the cancellation between π - and ρ -exchange in Eq.(8). It is therefore not necessary to have a more accurate treatment of the correlations.

Using the form (11.b) for the correlated potential together with Eq.(12), the matrix element is modified to

$$|M_{\pi\rho}^C|^2 = \frac{256}{3} m_N^2 \left(\frac{f_{\pi N}^2}{m_\pi^2} \right)^2 g_{aN}^2 k^4 \times \quad (13)$$

$$\left[\left(\frac{1}{(m_\pi^2 + k^2)} - \frac{1}{(m_\pi^2 + q_c^2 + k^2)} \right) - C_\rho \left(\frac{1}{(m_\rho^2 + k^2)} - \frac{1}{(m_\rho^2 + q_c^2 + k^2)} \right) \right]^2 + \text{E.T.} .$$

Finally, one must include the effect of the meson rescattering on the NN potential in the nuclear medium. This phenomenon occurs already at the three-nucleon level, and gives rise to nuclear three-body forces. The prototype process is the isobar excitation [14] indicated in fig.2. In nuclear matter, and after summation over the intermediate particle in fig.2, one can rewrite this correction as a polarization of the π - and ρ -propagators from Δ -hole (Δ -h) excitations [15]. Its physics has been extensively discussed in the literature and is closely related to the pionic Lorentz-Lorenz effect [16]. The corresponding correction is indicated in fig.3. Other contributions to the three-body force associated with the s-wave πN amplitude (for instance σ -term or correction from excitations of nucleons above the Fermi sea) may also contribute sizeably. As far as the latter correction is concerned, the density dependence is rather weak, and therefore it does not dominate at high density in contrast to the case for normal nuclear matter density [16]. For simplicity, and given the limited knowledge of many-body interactions at high density, we restrict ourselves in the following to the polarization mechanism from Δ -h excitations, as indicated in figs.2 and 3. It gives rise to the following polarization function for the pion

$$\Pi_o(k^2, \rho) \equiv -k^2 \chi_o(k^2, \rho), \quad (14)$$

where

$$\chi_o(k^2, \rho) = \frac{8}{9} \frac{f_{\pi N \Delta}^2}{m_\pi^2} \frac{\rho}{m_\Delta - m_N} F_{\pi N \Delta}^2(k^2), \quad (15)$$

and $F_{\pi N \Delta}$ is the $\pi N \Delta$ form factor written in a monopole form

$$F_{\pi N \Delta}(k^2) = \frac{\Lambda_\pi^2 - m_\pi^2}{\Lambda_\pi^2 + k^2}. \quad (16)$$

with $\Lambda_\pi \approx 1.1$ GeV/c. The $\pi N \Delta$ coupling constant $f_{\pi N \Delta}$ is taken equal to twice $f_{\pi N}$ [15], and m_Δ is the Δ -mass. The pion propagator is thus modified according to

$$\frac{1}{m_\pi^2 + k^2} \rightarrow \frac{1}{m_\pi^2 + \Pi_o(k^2, \rho) + k^2}. \quad (17)$$

Higher order corrections (or many-body effects for $A > 3$) can easily be incorporated in the polarization function by summing the Δ -hole bubbles to all orders [15,16], according to the equation indicated in fig.4. In this

figure, the Δ -hole- Δ -hole interaction $V_{\sigma\tau}$ is the irreducible interaction in the spin-isospin channel of the pion, i.e. the part of the NN interaction which is not already included in the iteration of the simple bubble of fig.3. It is usually assumed to be of zero range, with

$$V_{\sigma\tau} = g' \sigma_1 \cdot \sigma_2 \tau_1 \cdot \tau_2 \quad (18)$$

with g' fixed from various experimental results. We take here $g'=0.7$ [15,16]. The polarization function writes to all orders

$$\Pi(k^2, \rho) = \frac{\Pi_o(k^2, \rho)}{1 + g' \chi_o(k^2, \rho)} \equiv -k^2 \chi(k^2, \rho). \quad (19)$$

As a consequence of many-body effects, the π NN and ρ NN vertices are also renormalized, as can be seen in fig.5. Therefore, the resulting π -exchange matrix element squared is now

$$|M_\pi|^2 = \frac{256}{3} m_N^2 \left(\frac{f_{\pi N}^2}{m_\pi^2} \right)^2 g_{\pi N}^2 \left(\frac{\chi^\pi}{\chi_o^\pi} \right)^4 \frac{k^4}{\left(m_\pi^2 + (1 - \chi^\pi) k^2 \right)^2} + \text{E.T.} , \quad (20)$$

The corresponding modifications to the ρ -propagator and ρ NN vertex can be deduced as well, and incorporated in Eq.(13) in a similar way [17]. For the numerical calculation, we have evaluated the form factor in Eq.(15) at an average momentum transfer to get rid of the momentum dependence of χ . The vertex and propagator corrections at $\rho=3 \rho_o$ in Eq.(20), with $g'=0.7$, are $(\chi^\pi/\chi_o^\pi)^4 = 0.39$, $1 - \chi^\pi = 0.12$. From a quantitative point of view, the two corrections (modification of the propagators and vertices) go in opposite directions, the net effect being an overall reduction of the matrix element squared.

A first consequence of the modification (17) of the π - propagator is that at sufficiently high density, the propagator will pick up a pole at a critical momentum k_C defined by

$$m_\pi^2 + k_C^2 [1 - \chi(k_C^2, \rho)] = 0 . \quad (21)$$

This is the condition for pion condensation in the nuclear medium. In order to be consistent with our assumption that the medium at $3 \rho_0$ is of the same nature as at normal nuclear matter density the parameter g' must lead to pion condensation only at very high density. Otherwise the equation of state would be completely different to the one assumed here. Because of the form factor in Eq.(15), the critical momentum transfer cannot be more than about 500 MeV/c. This condition gives a lower limit on g' , since a smaller value for g' increases the polarization and lowers the density at which pion condensation occurs. To get pion condensation at $3 \rho_0$, one would have to take an unrealistic value of the order of 0.4 for g' . The value we use here ($g'=0.7$) is thus on the safe side and is consistent with values found in many other investigations [15,16]. For this value, the π -propagator is enhanced by a maximum factor of 4 at about 500 MeV/c as compared to zero density.

3. Energy-loss rate

The evaluation of the energy-loss rate from the transition amplitude $|M|^2$ has already been discussed in detail in both the degenerate and non-degenerate limit for the nucleons [1-3]. In the physical conditions under which the axions are expected to be emitted, typically $\rho=3\rho_0$ and $T=70$ MeV, the nuclear medium is in fact in the transition region between these two cases. In the following we give detailed results for the degenerate limit, and we discuss only qualitatively the non-degenerate limit, which is calculated in ref.[3].

The degenerate limit has been given in refs.[10,1-3], leading to the following rate for the simple π -exchange amplitude (5) on nn,pp and np pairs

$$\epsilon_{n,p}(D) = \frac{31}{3780\pi} m_N^2 T^6 g_{an,p}^2 \left(\frac{f_{\pi N}}{m_\pi} \right)^4 p_F(n,p) F(x_{n,p}^\pi), \quad (22.a)$$

$$\epsilon_{np}(D) = \frac{31}{405\pi} m_N^2 T^6 \frac{g_{an}^2 + g_{ap}^2}{2} \left(\frac{f_{\pi N}}{m_\pi} \right)^4 \frac{p_F(n) + p_F(p)}{2} F(x_{np}^\pi), \quad (22.b)$$

with

$$F(x) = 1 - \frac{3}{2} x \operatorname{arctg} \frac{1}{x} + \frac{x^2}{2(x^2 + 1)}, \quad (23)$$

and

$$x_{n,p}^\pi = \frac{m_\pi}{2p_F(n,p)}; \quad x_{np}^\pi = \frac{m_\pi}{p_F(n) + p_F(p)}. \quad (24)$$

The first order correction to Eq.(22.b) from the difference in neutron and proton density is small and decreases the energy-loss rate by about 20%. We do not consider it here, given the limited knowledge of the axion coupling constant (see below). We have also assumed in Eq.(22.b) that the axion-neutron and axion-proton coupling constants are equal, and have replaced them by their mean squared value [3] [†]. In these equations $p_F(n,p)$ is the neutron and proton Fermi momentum respectively,

$$p_F(n,p) = (3\pi^2 X_{n,p} \rho)^{1/3}, \quad (25)$$

where X_n and X_p are the neutron and proton mass fractions. We take here $X_n=0.88$ and $X_p=0.12$ [1]. The function $F(x)$ in Eqs.(22) comes from the integration of $k^4/(m_\pi^2+k^2)^2$. The corresponding function G for the terms of the form $k^4/[(m_1^2+k^2)(m_2^2+k^2)]$ which appears when ρ -exchange and correlated wave functions are taken into account is [12]

$$G(x_1, x_2) = 1 + \frac{x_1^3}{x_2^2 - x_1^2} \operatorname{arctg} \frac{1}{x_1} + \frac{x_2^3}{x_1^2 - x_2^2} \operatorname{arctg} \frac{1}{x_2}. \quad (26)$$

with obvious notations.

The various corrections to the single π -exchange contribution considered up to now lead to a modification of the function $F(x)$ in Eqs.(22). In the degenerate limit, these corrections depend only on the density and not on the (effective) nucleon mass and temperature in the medium. In the limit of very high density (i.e. $p_F \gg m_\pi$), the energy-loss rate is simply proportional to $F(0)$ which is unity. This approximation is equivalent to neglecting the pion mass in the matrix element squared (5). For the physical conditions we are interested in, this assumption is not very justified, since the relevant momentum transfers are of the order of a few times the pion mass. Keeping the exact density dependence of the process, one gets a factor of $F(x_{np}^\pi)=0.62$ instead of unity for the dominant process $n+p \rightarrow n+p+a$. As we already indicated, ρ -exchange contributions to the NN tensor interaction reduce this further to 0.12. These effects can be calculated quite reliably. As we mentioned above, the polarization corrections on π -exchange mechanism are very sensitive to the interaction $V_{\sigma\pi}$ in this density domain. With our choice $g'=0.7$, we find a factor of 0.36 compared to the bare interaction result of 0.62. Approximately the same reduction factor of 2 applies when ρ -exchange corrections are considered, leading to a final result of 0.069, when both correlation functions and polarization effects are included. We therefore get an overall reduction factor of about 15 as compared to the simplest estimate $F(0)=1$ and a factor of 9 as compared to the bare π -exchange contribution.

[†] Note that the energy-loss rate from $n+p \rightarrow n+p+a$ is about a factor of two less than in ref.[1]. This difference is probably due to the treatment of the exchange terms [3].

The non-degenerate limit has been derived in ref.[3] in the case where $|M|^2$ can be approximated by a constant (very high density limit). As we mentioned above, this is not a very good approximation, in particular when polarization effects or heavy-meson exchange are considered. Moreover, a simple estimate of the energy-loss rate in this limit is not possible due to the strong cancellation between π - and ρ -exchange contributions. A numerical calculation of the phase space integrals, along the lines discussed in ref.[3], and keeping track of the exact momentum dependence is certainly needed. This is beyond the scope of the present study. One should note that the temperature dependence of the energy-loss rate is quite different in the degenerate and non-degenerate limits, varying as T^6 and $T^{7/2}$ respectively. Because of this difference, the degenerate limit (valid at small T) overestimates the energy-loss rate for temperatures larger than about 60-90 MeV depending upon the nucleon mass in the nuclear medium (one half to one third the free nucleon mass, with the scale given by $m_N T$ in the non-degenerate limit). Since the temperature at which the axion is supposed to be predominantly emitted is about 70 MeV, this corresponds to the transition region between these two regimes. In this region, the expression for the degenerate limit is still a realistic estimate, given the limitations in the determination of the energy-loss rate.

The expressions (22) also involve the nucleon mass and the axion-nucleon coupling constant (g_{aN}). All other quantities have been discussed previously. The modification of the free nucleon mass in the medium is relatively well under control at moderate densities. At normal nuclear matter density, the effective nucleon mass is about 0.7 of the free mass and is a decreasing function of the density. At higher density, the determination of the nuclear equation of state is much more uncertain. In relativistic calculations one gets for instance an effective mass of about one third the free mass at three times normal density [18]. This is certainly a lower bound since relativistic corrections are expected to saturate at high density [19]. We shall take as a starting point $m_N = 1/2 m_N^0$, where m_N^0 is the free nucleon mass.

Less is known about the variation of g_{aN} with the density. The axion-nucleon coupling constant is usually expressed in terms of f_π using the Goldberger-Treiman relation

$$g_{aN} = m_N \frac{C_A}{f_\pi}, \quad (27)$$

where C_A is the analogue of the axial decay constant for the pion (f_π). It can be extracted at zero density from the EMC experiments or calculated from the non-relativistic quark model [1]. In analogy with the πN and axial coupling constants, one may expect that C_A is a decreasing function of the density, as is the case for the axial coupling constant. As far as the density dependence of f_π is concerned, we make two different choices so as to test the sensitivity of the results with respect to our assumptions. The first one is to take f_π constant, while the second one is to assume that it varies approximately like the nucleon mass as suggested in a Skyrmin model for f_π [20]. Since C_A scales like the nucleon mass, g_{aN} is proportional to its square in the first case, while it is linear in m_N in the second one. Collecting

all the corrections on m_N and g_{aN} , one finally gets a renormalization factor of $(m_N/m_N^0)^6$ or $(m_N/m_N^0)^4$, respectively, in the medium for the energy-loss rates (22).

A good theoretical derivation of the total energy removed by axions during the collapse requires an evaluation of the supernova evolution, which in practice means the use of a numerical evolution code. Especially the time evolution of the temperature and density of the nuclear medium at various stages of the collapse must be known in order to have a good estimate of the emission rate for axions. This is particularly so in the light of the corrections we mentioned above, since these are strongly density dependent. However, it is easy to get a rough lower bound for f_a from our previous considerations: by analogy with Eq.(22), we assume that the lower bound on f_a is simply proportional to the square root of the energy-loss rate ϵ , and take as a reference the calculations of ref.[1]. This has thus the following consequence:

As compared to the results of ref.[1], our estimate for the energy-loss rate is changed by several reduction factors: a factor of 9 from ρ -exchange and polarization corrections, a factor of 2 from the evaluation of the dominant $n+p \rightarrow n+p+a$ amplitude. Finally an additional factor of 2^6 or 2^4 from the modification of m_N and g_{aN} , depending upon the assumption made for the density dependence of f_a (constant or varying as m_N , respectively) with an effective nucleon mass of $1/2 m_N$. A less realistic value for the effective mass of the nucleon ($1/3 m_N$) gives a reduction factor between 3^6 and 3^4 , respectively. The major uncertainties clearly lie in the modifications of m_N and g_{aN} in the nuclear medium, and these are the main limits to the precision of the lower bound on f_a .

With these limitations in mind, we get the following estimate for the lower bound on f_a

$$f_a \geq (0.7 - 5) \times 10^{10} \text{ GeV}. \quad (28)$$

In the derivation of the results (28) we have assumed the largest value for C_A from the non-relativistic quark model. It may be reduced by up to a factor of 2 when C_A is extracted from the EMC experiments [1].

In this letter, we have investigated the various modifications which occur in the axion emission rate for nucleon-nucleon bremsstrahlung process at high density, relatively to the simple estimate in terms of one-pion exchange. Many of these corrections are well established in nuclear systems at normal density: ρ -exchange contribution to the tensor interaction, polarization effects from many-body forces and density dependence of the nucleon mass. In the physical conditions relevant for the axion emission in the supernova SN 1987A, these corrections are large and cannot be fixed with accuracy. The lower bound for the axion decay constant which results from our analysis is much lower than previous estimates, and uncertain by about a factor of 5. We do not believe that the bound can be fixed with better precision, given our actual knowledge of the many-body problem under these conditions of temperature and density.

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Figure captions

Figure 1 : Axion emission by nucleon-nucleon bremsstrahlung. All possible time orderings, as well as the interchange between the two particles in the initial and final states are implicitly understood.

Figure 2 : Typical contribution to the three-body force in nuclei.

Figure 3 : First order polarization of π - and ρ -meson propagators in the medium from Δ -hole excitations.

Figure 4 : Summation to all orders of π and ρ polarization in the medium. $V_{\sigma\tau}$ represents the irreducible particle-hole interaction from π - and ρ -exchange (see text).

Figure 5 : Vertex renormalization in the medium from Δ -hole excitations.

Fig.1

Fig.2

Fig.3

Fig.4

Fig.5