Limits on $n\bar{n}$ oscillations from nuclear stability

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Abstract

The relationship between the lower limit on the nuclear stability lifetime as derived from the non disappearance of ‘stable’ nuclei ($T_d \gtrsim 5.4 \times 10^{31}$ yr), and the lower limit thus implied on the oscillation time ($\tau_{n\bar{n}}$) of a possibly underlying neutron-antineutron oscillation process, is clarified by studying the time evolution of the nuclear decay within a simple model which respects unitarity. The order-of-magnitude result $\tau_{n\bar{n}} \approx 2\left(T_d/\Gamma_{\bar{n}}\right)^{1/2} > 2 \times 10^8$ sec, where $\Gamma_{\bar{n}}$ is a typical $\bar{n}$ nuclear annihilation width, agrees as expected with the limit on $\tau_{n\bar{n}}$ established by several detailed nuclear physics calculations, but sharply disagreeing by 15 orders of magnitude with a claim published recently in Phys. Rev. CRAP.

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I. INTRODUCTION

The stability of nuclei, as determined by looking for proton decay \[1,2\], places a lower limit on the ‘disappearance’ lifetime \(T_d > (5.4 \pm 1.1) \times 10^{31}\) yr, and sets a lower limit also on the lifetime of other hypothetical processes such as neutron-antineutron \((n\bar{n})\) oscillations in free space. The lower limit on the free-space oscillation time \(\tau_{n\bar{n}}\) which has emerged from several quantitative nuclear physics calculations using \(n\) and \(\bar{n}\) optical potentials \[3–6\] is approximately given by

\[
\tau_{n\bar{n}} \approx 2\left(\frac{T_d}{\Gamma_{\bar{n}}}\right)^{1/2} > (2.1 \pm 0.2) \times 10^8\text{ sec },
\]  

(1)

where \(\Gamma_{\bar{n}} \sim 100\) MeV is a typical \(\bar{n}\)-nuclear annihilation width. This slightly exceeds the limit

\[
\tau_{n\bar{n}} > 0.9 \times 10^8\text{ sec }
\]  

(2)

set in the ILL-Grenoble reactor experiment \[7\]. A new experiment planned at ORNL will hopefully improve this limit by two orders of magnitude \[8\].

Eq. (1) may be rewritten as

\[
T_d \approx \frac{1}{4}(\Gamma_{\bar{n}}\tau_{n\bar{n}})\tau_{n\bar{n}},
\]  

(3)

showing that, since \(\Gamma_{\bar{n}}\tau_{n\bar{n}} \gg 1\) owing to the huge lifetime distinction between strong interactions (lifetime \(\sim \Gamma_{\bar{n}}^{-1}\)) and superweak interactions (lifetime \(\sim \tau_{n\bar{n}}\)), the nuclear disappearance lifetime \(T_d\) induced by \(n\bar{n}\) oscillations is many orders of magnitude longer than the \(n\bar{n}\) oscillation lifetime in free space. This is equivalent to the common statement that \(n\bar{n}\) oscillations in matter undergo severe suppression, since the neutron and the antineutron feel nuclear potentials which are extremely different from each other, and the mass degeneracy which allows such pure oscillations between a free \(n\) and a free \(\bar{n}\) is thus removed in the nuclear medium.

Nazaruk \[9,10\] has raised objections to the use of nuclear physics potential models, claiming that \(n\bar{n}\) oscillations are not suppressed at all in the nuclear medium. Thus, in the first paper \[9\] he obtained \(T_d \sim \tau_{n\bar{n}}\), resulting in the limit \(\tau_{n\bar{n}} > 10^{31}\) yr which advances the lower limit \(\tau_{n\bar{n}}\) by about 30 orders of magnitude! Dover, Gal, Richard \[11\] and Krivoruchenko \[12\] subsequently pinpointed errors in that paper and reestablished the generally accepted lower limit \(\tau_{n\bar{n}}\). In the second, more recent paper \[10\], Nazaruk ‘rederived’ the \(T_d \sim \tau_{n\bar{n}}\) result, but argued that since the nuclear decay is non exponential, a more careful consideration of the nuclear stability limit translates into \(\tau_{n\bar{n}} > 10^{16}\) yr, which advances the lower limit \(\tau_{n\bar{n}}\) by ‘only’ 15 orders of magnitude. Such far-reaching claims, when published by a first-rate journal, should not go unanswered.

In this note I wish to expand the arguments outlined by Dover et al. \[11\], in order to show in detail how Eq. (1) for \(\tau_{n\bar{n}}\) is obtained. In Sec. \[\text{II}\] we consider the \(n\bar{n}\) mass matrix in \(\bar{n}\)-annihilating matter, in order to derive in the most economical way the eigen-lifetimes in a toy seesaw model. In Sec. \[\text{III}\] we study in detail the time evolution of the \(n\bar{n}\) ‘oscillating’ system in matter. In the concluding Sec. \[\text{IV}\] we have tried to identify the error in Nazaruk’s logic.
II. OSCILLATIONS — THE SEESAW MECHANISM

A common approach in problems of oscillations, be it for CP symmetry or in neutrino mass generation problems, is to diagonalize the mass matrix in order to find out the physical states. For \( n\bar{n} \) oscillations induced by a coupling \( \epsilon \), the simplest form of the in-medium mass matrix is

\[
\begin{pmatrix}
  m & \epsilon \\
  \epsilon & m - i \frac{\Gamma_{\bar{n}}}{2}
\end{pmatrix}
\]

where \( m \) is the joint value of the \( n \) and \( \bar{n} \) masses, and \( \Gamma_{\bar{n}} \sim 100 \text{ MeV} \) is the annihilation width of the \( \bar{n} \) in the nuclear medium. For \( \Gamma_{\bar{n}} = 0 \), i.e. in free space, the two eigenmasses are \( m \pm \epsilon \), differing only by a tiny \( \epsilon \) from each other. The reality of these eigenmasses gives rise to a purely oscillatory behavior \( n \leftrightarrow \bar{n} \), with lifetime \( \tau_{n\bar{n}} = \frac{\epsilon}{\Gamma_{\bar{n}}} \). In the nuclear medium, when \( \Gamma_{\bar{n}} \neq 0 \), the eigenmasses of (4) are given by

\[
m_{\text{eigen}} = m - \frac{1}{4} i \Gamma_{\bar{n}} \pm \frac{1}{4} i \Gamma_{\bar{n}} \left( 1 - \frac{16 \epsilon^2}{\Gamma_{\bar{n}}^2} \right)^{1/2},
\]

which to leading order in \( \epsilon^2/\Gamma_{\bar{n}}^2 \) assume the values

\[
m_n = m - \frac{1}{2} i \frac{4 \epsilon^2}{\Gamma_{\bar{n}}} + \ldots, \quad m_{\bar{n}} = m - \frac{1}{2} i \Gamma_{\bar{n}} + \ldots,
\]

where the dots stand for higher order terms. The degeneracy of the real parts of \( m_{\text{eigen}} \) can easily be removed by introducing \( n \)-nucleus and \( \bar{n} \)-nucleus real optical potentials which differ substantially from each other. In discussing decay modes, however, only \( \text{Im} m_{\text{eigen}} \) matters; adding real potentials does not change the splitting between, and the order of magnitude of the two values for \( \text{Im} m_{\text{eigen}} \) in Eq. (4). Therefore, for the sake of simplicity, we ignore these real potentials. Identifying the eigenwidth \( \gamma \) with \(-2 \text{ Im} m_{\text{eigen}} \), one obtains from Eq. (4) two eigenwidths:

\[
\gamma_n = \frac{4 \epsilon}{\Gamma_{\bar{n}}} \epsilon, \quad \gamma_{\bar{n}} = \Gamma_{\bar{n}}.
\]

Intuitively, since \( \gamma_{\bar{n}} = \Gamma_{\bar{n}} \) is the nuclear annihilation width of an antineutron, \( \gamma_n \) must stand for the decay rate of the neutron (in units where \( \hbar = 1 \)). It is clear that this decay rate undergoes a huge suppression factor, \( 4 \epsilon / \Gamma_{\bar{n}} \ll 1 \), with respect to the free-space oscillation rate \( \epsilon \). The huge disparity between \( \gamma_n \) and \( \gamma_{\bar{n}} \) is a good demonstration of the seesaw mechanism encountered in the discussion of neutrino mass generation problems; note that the product \( \gamma_n \gamma_{\bar{n}} = 4 \epsilon^2 \) is of the same order of magnitude as in free space, which is a necessary condition for the seesaw mechanism. These qualitative arguments need to be explored quantitatively by studying the time evolution of \( n\bar{n} \) ‘oscillations’ in the nuclear medium. This is done in the next section.

III. EXPLICIT TIME EVOLUTION

For simplicity, following Ref. [11], we write the time dependent coupled Schrödinger equations for zero momentum neutron and antineutron in nuclear matter (\( \hbar = 1 \)):
\[ i \partial_t \psi_n = \varepsilon \psi_n, \quad i \partial_t \psi_{\bar{n}} = \varepsilon \psi_{\bar{n}} - \frac{i \Gamma_{\bar{n}}}{2} \psi_{\bar{n}}. \] (8)

No nuclear (real) potentials \( U_n \) and \( U_{\bar{n}} \) appear in the present discussion, which is focussed on the interplay and competition between the free-space oscillation rate \( \varepsilon = \tau_n^{-1} \) and the \( \bar{n} \) decay rate \( \Gamma_{\bar{n}} \) due to the strong-interaction \( \bar{n} \) nuclear annihilation, and their effect will be briefly discussed later on. The coupled first order equations (8) give rise to the following second order differential equation for each one of \( \psi_n, \psi_{\bar{n}} \):

\[
\left( \partial_t^2 + \frac{1}{2} \Gamma_{\bar{n}} \partial_t + \varepsilon^2 \right) \psi = 0. \] (9)

Seeking eigensolutions of the form \( \psi_j = \exp(i \omega_j t/2) \), the eigenfrequencies \( \omega_j \) satisfy a quadratic equation

\[
\omega^2 - i \Gamma_{\bar{n}} \omega - 4\varepsilon^2 = 0, \] (10)

yielding two solutions:

\[
\omega_{\bar{n}} = i \Gamma_{\bar{n}} \left\{ 1 - \frac{1}{2} \left[ 1 - \left( 1 - 16 \frac{\varepsilon^2}{\Gamma_{\bar{n}}^2} \right)^{1/2} \right] \right\} = i \Gamma_{\bar{n}} \left( 1 - 4 \frac{\varepsilon^2}{\Gamma_{\bar{n}}^2} + \ldots \right), \] (11)

\[
\omega_n = i \Gamma_n \frac{1}{2} \left[ 1 - \left( 1 - 16 \frac{\varepsilon^2}{\Gamma_n^2} \right)^{1/2} \right] = i \frac{4\varepsilon^2}{\Gamma_n} \left( 1 + 4 \frac{\varepsilon^2}{\Gamma_n^2} + \ldots \right), \] (12)

where the expansion in even power of the extremely small quantity \( \varepsilon^2/\Gamma_{\bar{n}}^2 \ll 1 \) is indicated. These eigenfrequencies are related to the eigenwidths in Eq. (7) by \( \omega = i \gamma \). The eigenfrequency \( \omega_{\bar{n}} \) gives rise to exponential decay of \( |\psi| \) with a rate \( \Gamma_{\bar{n}} \), to leading order. Dover et al. [11] noted that \( \omega_{\bar{n}} \) should be discarded since it corresponds to the rate of disappearance of an antineutron when its oscillation coupling \( \varepsilon \) to the neutron is neglected. It was further noted that the eigenfrequency \( \omega_n \) gives rise to a rate of disappearance \( 4(\varepsilon/\Gamma_n)\varepsilon \) which was then interpreted as that for a neutron, due to its oscillation coupling \( \varepsilon \) to the antineutron. This latter rate is equivalent to the estimate given in Eq. (3) for \( T_d \). We note that since \( \omega_n \omega_{\bar{n}} = -4\varepsilon^2 \), the huge disparity between the frequencies \( \omega_n \) and \( \omega_{\bar{n}} \) is a good demonstration of the seesaw effect encountered in the discussion of scales of masses in high energy physics.

Here we wish to proceed in greater detail and rigor, using the eigenfrequencies (11, 12), to construct the wavefunctions \( \psi_j \) with the appropriate temporal boundary conditions for the neutron (\( j = n \)) and for the antineutron (\( j = \bar{n} \)). In considering the nuclear disappearance lifetime \( T_d \), these boundary conditions are

\[
\psi_n(t = 0) = 1, \quad \psi_{\bar{n}}(t = 0) = 0. \] (13)

The corresponding linear combinations of the eigensolutions are then given by

\[
\psi_n(t) = \frac{\omega_{\bar{n}}}{\omega_n - \omega_{\bar{n}}} \exp \left( i \omega_{\bar{n}} t/2 \right) - \frac{\omega_n}{\omega_n - \omega_{\bar{n}}} \exp \left( i \omega_n t/2 \right), \] (14)
\[
\psi_n(t) = \frac{2\varepsilon}{\omega_n - \omega_n} \left( \exp(i\omega nt/2) - \exp(i\omega _n t/2) \right).
\]

(15)

Note that since \(\omega_n/\omega_n \approx 4\varepsilon^2/\Gamma_n^2 \ll 1\), \(\psi_n(t)\) is dominated at all times by the eigensolution with eigenfrequency \(\omega_n\). It is also clear, by inspecting Eq. (11), that the \(\bar{n}\) wavefunction which evolves from zero at \(t = 0\) always remains very small with respect to the neutron wavefunction:

\[
|\psi_n(t)/\psi_n(t)|^2 \sim \frac{4\varepsilon^2}{\Gamma_n^2},
\]

(16)

this approximate relationship becoming exact for times \(t \gg \Gamma_n^{-1}\) when the second exponent on the right-hand side of Eqs. (14, 15) may be safely dropped out. For such ‘long’ times, the time dependence of both \(\psi_n\) and \(\psi_{\bar{n}}\) is given by

\[
|\psi(t)|^2 \sim \exp \left( -4\frac{\varepsilon^2}{\Gamma_n} t \right),
\]

(17)

so that the rate of disappearance \(\Gamma_d\) may be read off this exponential decay:

\[
\Gamma_d = \frac{4\varepsilon^2}{\Gamma_n}.
\]

(18)

Recalling that \(\varepsilon = \tau_{nn}^{-1}\), we obtain for the disappearance lifetime

\[
T_d = \frac{1}{\frac{1}{4} \left( \Gamma_n \tau_{nn} \right) \tau_{\bar{n}n}},
\]

(19)

which agrees precisely with Eq. (3) of Sec. I.

The minimal model of Eq. (8) for the time evolution of \(\psi_n\) and \(\psi_{\bar{n}}\) may be extended to include real potentials \(U_n\) and \(U_{\bar{n}}\), and spatial structure:

\[
i\partial_t \psi_n = -\frac{\Delta}{2m} \psi_n + U_n(r) \psi_n + \varepsilon \psi_{\bar{n}},
\]

(20)

\[
i\partial_t \psi_{\bar{n}} = -\frac{\Delta}{2m} \psi_{\bar{n}} + (U_{\bar{n}}(r) - iW(r)) \psi_n + \varepsilon \psi_n,
\]

(21)

with the same \(t = 0\) boundary conditions specified in Eq. (13). In order to obtain the decay rate of the neutron, we multiply Eq. (20) for \(\psi_n\) by \(\psi_n^*\), and the corresponding equation for \(\psi_{\bar{n}}^*\) by \(\psi_n\), subtract from each other and integrate over space. The result is

\[
- \partial_t \int |\psi_n|^2 \, d^3r = 2\varepsilon \int \text{Im} (\psi_n^* \psi_n) \, d^3r.
\]

(22)

This decay rate vanishes for \(\varepsilon \to 0\), but the precise power of \(\varepsilon\) by which it occurs depends on the ratio \(|\psi_{\bar{n}}/\psi_n|\). Acting similarly on Eq. (21) for \(\psi_n\), the decay rate of the antineutron is obtained:

\[
- \partial_t \int |\psi_{\bar{n}}|^2 \, d^3r = \int 2W |\psi_{\bar{n}}|^2 \, d^3r - 2\varepsilon \int \text{Im} (\psi_n^* \psi_{\bar{n}}) \, d^3r.
\]

(23)
We note that the $2\varepsilon \int$ terms in Eqs. (22, 23) appear with opposite signs; what’s lost out of the neutron intensity due to the oscillation coupling $\varepsilon$ is precisely gained by the $\bar{n}$ intensity, and vice versa, on top of the $\bar{n}$ annihilation decay rate. The total decay rate is obtained by adding up Eqs. (22, 23):

$$-\partial_t \int \left( |\psi_n|^2 + |\psi_{\bar{n}}|^2 \right) d^3r = \int 2W |\psi_{\bar{n}}|^2 d^3r , \quad (24)$$

showing that if the antineutron did not annihilate in the nuclear medium ($W = 0$), there would be no loss of intensity from the combined $n, \bar{n}$ space. A rough estimate of $|\psi_{\bar{n}}|^2/|\psi_n|^2$ may be made by inspecting Eq. (21) and noticing that, for the boundary conditions Eq. (13), the evolved $(U_{\bar{n}} - iW)\psi_{\bar{n}}$ term on the right-hand side of Eq. (21) should be of the same order of magnitude as the source term $\varepsilon \psi_n$ which generates it. Since $U_{\bar{n}}$ and $W \sim \Gamma_{\bar{n}}/2$ are all of the same order of magnitude, we obtain

$$\frac{|\psi_{\bar{n}}|^2}{|\psi_n|^2} \sim \frac{\varepsilon^2}{W^2} \sim \frac{4\varepsilon^2}{\Gamma_{\bar{n}}^2} , \quad (25)$$

which agrees with Eq. (16). The properly normalized decay rate $\gamma$ is then given, using Eq. (24), by

$$\gamma \equiv -\partial_t \int \left( |\psi_n|^2 + |\psi_{\bar{n}}|^2 \right) d^3r \bigg/ \int \left( |\psi_n|^2 + |\psi_{\bar{n}}|^2 \right) d^3r = \frac{\int 2W |\psi_{\bar{n}}|^2 d^3r}{\int \left( |\psi_n|^2 + |\psi_{\bar{n}}|^2 \right) d^3r} , \quad (26)$$

which is approximately of order

$$\gamma \sim 2W \frac{|\psi_{\bar{n}}|^2}{|\psi_n|^2} \sim \frac{4\varepsilon^2}{\Gamma_{\bar{n}}} , \quad (27)$$

agreeing with the order of magnitude of $\Gamma_d$, Eq. (18). We note that Eq. (26) expresses the loss of $n$ and $\bar{n}$ intensities to the unspecified final nuclear debris products.

IV. DISCUSSION

We outlined in the previous sections, using a simple and transparent potential model, how the free-space $n\bar{n}$ oscillation period $\tau_{n\bar{n}}$ gets tremendously prolonged in the nuclear medium, by a factor $\Gamma_{\bar{n}}\tau_{n\bar{n}}/4$, to yield the corresponding nuclear decay lifetime $T_d$ of Eq. (19). More detailed calculations \[4,5\], which treat the nuclear medium as a dynamical entity, confirm this order-of-magnitude estimate of $T_d$. It is then perplexing to encounter occasionally claims that $n\bar{n}$ oscillations are not suppressed in the nuclear medium. Below we counter the most recent claim of this sort made by Nazaruk \[10\].

Nazaruk argued that potential models involve double counting by allowing for $\bar{n}$-nucleus elastic and inelastic scattering, and that this is manifested within such models by a calculated nuclear decay probability which is linear in the time $t$ instead of the (correctly anticipated) $t^2$ dependence for ‘short’ oscillation times. Our response is that the only nuclear property of the antineutron included in the present potential model is its nuclear decay width $\Gamma_{\bar{n}}$, ...
unrelated in this model to any underlying $\bar{n}$ scattering process. Furthermore, we do recover a leading $t^2$ dependence by expanding Eqs. (14, 15) in powers of time $t$:

\[
1 - |\psi_n|^2 = \varepsilon^2 t^2 - \frac{1}{6} \varepsilon^2 \Gamma_{\bar{n}} t^3 + \ldots ,
\]

(28)

\[
|\psi_{\bar{n}}|^2 = \varepsilon^2 t^2 - \frac{1}{2} \varepsilon^2 \Gamma_{\bar{n}} t^3 + \ldots ,
\]

(29)

so that the neutron depletion rate would appear to be $\varepsilon$, as in free space — unaffected by the nuclear medium, in particular by the $\bar{n}$ annihilation width $\Gamma_{\bar{n}}$, and agreeing with Nazaruk’s thesis. However, this expansion is useful only for extremely short times, such that $\Gamma_{\bar{n}} t \ll 1$, before the $\bar{n}$ annihilation can have its act and exercise its toll. For relevant times which satisfy $\Gamma_{\bar{n}} t \gg 1$, the expansion (28, 29) becomes useless. As argued here in Sec. III, however, for such (still short) times, the $\omega_{\bar{n}}$ exponent in Eqs. (14, 15) may safely be neglected, resulting in the following exponential decays

\[
|\psi_n(t)|^2 \to \exp \left( -4 \frac{\varepsilon^2}{\Gamma_{\bar{n}}} t \right) ,
\]

\[
|\psi_{\bar{n}}(t)|^2 \to \frac{4 \varepsilon^2}{\Gamma_{\bar{n}}^2} \exp \left( -4 \frac{\varepsilon^2}{\Gamma_{\bar{n}}} t \right) ,
\]

(30)

with a rate of disappearance $\Gamma_d = 4 \varepsilon^2 / \Gamma_{\bar{n}}$ (Eq. (18)) which is directly read off this exponential decay.

The trouble with Nazaruk’s arguments may be traced as follows. Using S-matrix manipulations, for $t \gg \Gamma_{\bar{n}}^{-1}$ he reaches the expression

\[
w(t) \approx \varepsilon^2 t^2 w_n(t)
\]

(31)

for the ‘probability’ $w(t)$ of nuclear decay in terms of the $\bar{n}$ decay probability $w_n(t)$ (see Eqs. (11, 18, 19) of Ref. [10]). He then assumes that, since over this time all the antineutrons got already annihilated, $w_{\bar{n}}(t) \approx 1$. This is wrong: the $\bar{n}$ decay probability is a conditional one, depending on how many antineutrons were there in the first place with respect to neutrons. Unitarity prevents the accumulation of too many $\bar{n}$, and as seen from Eq. (30)

\[
w_{\bar{n}} \left( \Gamma_{\bar{n}}^{-1} \ll t \ll \Gamma_{\bar{n}} / \varepsilon^2 \right) \approx \frac{4 \varepsilon^2}{\Gamma_{\bar{n}}^2} .
\]

(32)

Hence

\[
w(t) \approx \frac{4 \varepsilon^4}{\Gamma_{\bar{n}}^2} t^2 ,
\]

(33)

so that the neutron disappearance rate is of order $2 \varepsilon^2 / \Gamma_{\bar{n}}$, in agreement with Eq. (27) and with our result Eq. (18) for $\Gamma_d$.

To summarize, we have confirmed the well known lower limit obtained on $n\bar{n}$ oscillations in free space from the stability of nuclei, as given by the estimate of Eq. (1) which states that $\tau_{n\bar{n}} \approx 2 \times 10^8$ sec. This estimate agrees to within better than a factor of two with the series of quantitative calculations by Alberico et al. [4] and Dover et al. [5]. There is good reason then to plan reactor experiments [8] which promise to push up this lower limit by perhaps as much as two orders of magnitude.

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