ON THE "FIELD-THEORETICAL APPROACH" TO
THE NEUTRON-ANTINEUTRON OSCILLATIONS IN
NUCLEI

Vladimir Kopeliovich,*
Institute for Nuclear Research of RAS
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Abstract

It is argued that within the correct treatment of analytical properties of the transition amplitudes, in particular, the second order pole structure, characteristic for the \( n - \bar{n} \) transition in nuclei, the "infrared divergences" discussed in some papers, do not appear. Explicit calculation with the help of diagram technique shows that the neutron-antineutron oscillations are strongly suppressed within a deuteron, as well as within arbitrary nucleus, in comparison with the oscillations in vacuum.

1. The neutron-antineutron transition induced by the baryon number violating interaction (\( \Delta B = 2 \)) predicted within some variants of grand unified theories (GUT) has been discussed in many papers since 1970 [1], see [2, 3, 4, 5, 6]. Experimental results of searches for such transition are available, in vacuum (reactor experiments [7], and references therein), in nucleus \(^{16}\)O [8] and in \( Fe \) nucleus [9], see also the PDG tables.

During the later time there have been many speculations that the neutron-antineutron oscillations in nuclei are not suppressed in comparison with the \( n - \bar{n} \) transition in vacuum [10, 11]. The arguments were based on the "true field-theoretical approach" to this problem. The result of [10] has been criticized in a number of papers [12, 13, 14, 15] which used somewhat different approaches (potential, S-matrix, diagram), and general physics arguments.

However, in view of continuing publications [11] containing same statement as in [10], it seems to be necessary to analyze this problem just within the

*e-mail: kopelio@inr.ru
quantum field theory based approach used in [10, 11]. My consideration is close to the approach of paper [14] where the diagram technique has been applied to study neutron-antineutron transition in nuclei, although differs from [14] in some details. More recent realistic calculations of the neutron-antineutron transition in nuclei can be found in [16, 17].

Here we give first some general arguments based on analytical properties of amplitudes in favour of suppression of \( n - \bar{n} \) transition in nuclei (section 3). The simplest example of the deuteron where the final result can be obtained in closed form, is considered in details in section 4, and the result of [14] for the case of the deuteron is reproduced. The analogy with the nucleus formfactor at zero momentum transfer is noted in section 5.

2. To introduce notations, let us consider first the \( n\bar{n} \) transition in vacuum which is described by the baryon number violating interaction (see, e.g. [2, 13, 14]) \( V = \mu_{n\bar{n}}\sigma_1/2 \), \( \sigma_1 \) being the Pauli matrix. \( \mu_{n\bar{n}} \) is the parameter which has the dimension of mass, to be predicted by grand unified theories and to be defined experimentally \(^1\). The \( n - \bar{n} \) state is described by 2-component spinor \( \Psi \), lower component being the starting neutron, the upper one - the appearing antineutron. The evolution equation is

\[
i \frac{d\Psi}{dt} = (V_0 + V)\Psi
\]

with \( V_0 = m_N - i\gamma_n/2 \) in the rest frame of the neutron (\( m_N \) is the nucleon mass, \( \gamma_n \) - the (anti)neutron normal weak interaction decay width, and we take \( \gamma_{\bar{n}} = \gamma_n \), as it follows from \( CP \)-invariance of weak interactions). Eq. (1) has solution

\[
\Psi(t) = \exp[-i (\mu_{n\bar{n}}t \sigma_1/2 + V_0t)] \Psi_0 = \left[ \cos \frac{\mu_{n\bar{n}}t}{2} - i\sigma_1 \sin \frac{\mu_{n\bar{n}}t}{2} \right] \exp(-iV_0t)\Psi_0,
\]

(2)

Here \( \Psi_0 \) is the starting wave function, e.g. \( \Psi_0 = (0, 1)^T \). In this case we have for an arbitrary time

\[
\Psi(\bar{n}, t) = -i \sin \frac{\mu_{n\bar{n}}t}{2} \exp(-iV_0t), \quad \Psi(n, t) = \cos \frac{\mu_{n\bar{n}}t}{2} \exp(-iV_0t),
\]

(3)

which describes oscillation \( n - \bar{n} \). Evidently, for large enough observation times, \( t_{\text{obs}} \gg 1/\mu_{n\bar{n}} \), the average probabilities to observe neutron and antineutron are

\(^1\)There is relation \( \mu_{n\bar{n}} = 2\delta m \) with the parameter \( \delta m \) introduced in [2]. The neutron-antineutron oscillation time in vacuum is \( \tau_{n\bar{n}} = 1/\delta m = 2/\mu_{n\bar{n}} \), see also [17] and references in this paper.
equal if we neglect the natural decay of the neutron (antineutron):

$$W(\bar{n}) = |\Psi(\bar{n})|^2 = |\Psi(n)|^2 = W(n) = 1/2. \quad (4)$$

This case is, however, of academic interest only, since $\gamma_n \gg \mu_{n\bar{n}}^2$. It should be stressed that in the vacuum neutron goes over into antineutron, also discrete localized in space state, which can go over again to the neutron, so the oscillation neutron to antineutron and back takes place.

Since the parameter $\mu_{n\bar{n}}$ is small, expansion of $\sin$ and $\cos$ can be made in Eq. (3) at not too large times. In this case the average (over the time $t^{\text{obs}} \ll 1/\mu_{n\bar{n}}$) change of the probability of appearance of antineutron in vacuum is (for the sake of brevity we do not take into account the (anti)neutron natural instability which has obvious consequences)

$$W(\bar{n}; t^{\text{obs}})/t^{\text{obs}} = |\Psi(\bar{n}, t^{\text{obs}})|^2/t^{\text{obs}} \simeq \frac{\mu_{n\bar{n}}^2 t^{\text{obs}}}{4} \quad (5)$$

which has, obviously, dimension of the width $\Gamma$. So, in vacuum the transition $n \rightarrow \bar{n}$ is suppressed if the observation time is small, $t^{\text{obs}} \ll 1/\mu_{n\bar{n}}$. From existing data obtained with free neutrons from reactor the oscillation time is greater than $0.86 \cdot 10^8 \text{sec} \simeq 2.7 \text{ years}$ [7], therefore,

$$\mu_{n\bar{n}} < 1.5 \cdot 10^{-23} \text{ eV}, \quad (6)$$

very small quantity.

Recalculation of the quantity $\mu_{n\bar{n}}$ or $\tau_{n\bar{n}}$ from existing data on nuclei stability [8, 9] is somewhat model dependent, and different authors obtained somewhat different results, within about 1 order of magnitude, see e.g. discussion in [14, 16, 17]. Most recent results for $\mu_{n\bar{n}}$ obtained from nuclear stability data are close to (6) [16, 17], see also next section.

3. In the case of nuclei the $n - \bar{n}$ line with the transition amplitude $\mu_{n\bar{n}}$ is the element of any amplitude describing the nucleus decay $A \rightarrow (A - 2) + \text{mesons}$, where $(A - 2)$ denotes a nucleus or some system of baryons

$$\int_0^\infty |\Psi(n, t)|^2 dt = \frac{2\gamma_n^2 + \mu_{n\bar{n}}^2}{2\gamma_n(\gamma_n + \mu_{n\bar{n}})}, \quad \int_0^\infty |\Psi(\bar{n}, t)|^2 dt = \frac{\mu_{n\bar{n}}^2}{2\gamma_n(\gamma_n + \mu_{n\bar{n}})}$$

for neutron as initial state and for arbitrary, but different from zero $\gamma_n$. The difference between both quantities is obvious, and disappears when $\gamma_n \rightarrow 0$. 

2It is a matter of simple algebra to calculate the integrals over time of the probabilities $|\Psi(n, t)|^2$ and $|\Psi(\bar{n}, t)|^2$: 
with baryonic number $A - 2$, see Fig. 1. The decay probability is therefore proportional to $\mu_{n\bar{n}}^2$, and we can write by dimension arguments

$$\Gamma(A \to (A - 2) + mesons) \sim \frac{\mu_{n\bar{n}}^2}{m_0},$$

(7)

where $m_0$ is some energy (mass) scale. For the result of [10, 11] to be correct the mass $m_0$ should be very small, $m_0 \sim \mu_{n\bar{n}} \sim 10^{-23}$ eV, but we shall argue that $m_0$ is of the order of normal hadronic or nuclear scale, $m_0 \sim m_{hadr} \sim (10-100)$ MeV. We can obtain the same result from the above vacuum formula (5), if we take the time $t_{obs} \sim 1/m_{hadr}$.

Indeed, the matrix element of any diagram containing such transition

$$T(A \to (A - 2) + mesons) \sim$$

$$\sim \mu_{n\bar{n}}(A - Z) \int V(A; n, (A - 1)) \frac{\tilde{T}(\bar{n} + (A - 1) \to (A - 2) + mesons)}{(E_n - E_n^0 + i\delta)^2} dE_n \simeq$$

$$\simeq -2\pi i (A - Z) \frac{d(V \tilde{T})}{dE_n}(E_n = E_n^0),$$

(8)

according to the Cauchy theorem known from the theory of functions of complex variable. $E_n$ is the neutron (antineutron) energy - integration variable, $E_n^0$ is the (anti)neutron on-mass-shell energy $E_n^0 \simeq m_N + \vec{p}^2/2m_N$. The energy-momentum conservation should be taken into account for the vertex $V(A \to n + (A - 1))$ which includes the propagator of the $(A - 1)$ system, and for the annihilation
amplitude $\tilde{T}$. The case of the deuteron considered below is quite transparent and illustrative.

The amplitude $\tilde{T}$ which describes the annihilation of the antineutron, and the vertex function $V$ are of normal hadronic or nuclear scale and cannot, in principle, contain a very small factors in denominator (or very large factors, of the order of $10^{15}$, in the numerator). By this reason we come to the above Eq. (7), and the resulting decay width of the nucleus is very small,

$$\Gamma(A \to (A - 2) + mesons) < 10^{-30}\mu_{n\bar{n}},$$

(9)

at least 30 orders of magnitude smaller than the inverse time of neutron-antineutron oscillation in vacuum $\mu_{n\bar{n}}$. From Eq. (7) or (9) we obtain

$$\mu_{n\bar{n}} \sim \sqrt{\Gamma(A \to (A - 2) + mesons)m_0},$$

(10)

and when one tries to get restriction on $\mu_{n\bar{n}}$ from data on nuclei stability [8, 9] the result is close to that from vacuum experiment [7], somewhat smaller, within one order of magnitude [5, 6, 14]. The result of [16] differs from that of [14] for heavier nuclei, and the authors [16] come to the conclusion, that experiments with free neutrons from reactor could provide stronger restriction on the neutron-antineutron transition parameter than experiments on stability of nuclear matter.\(^3\)

According to [10, 11] the probability of the nucleus decay is proportional to $W(t_{obs}) \sim \mu_{n\bar{n}}^2 (t_{obs})^2$ (the process proceeds similar to the vacuum case), where $t_{obs}$ is the large observation time, of the order of $\sim 1$ year or greater. By this reason the extracted value of $\mu_{n\bar{n}}$ is smaller than that given by Eq. (10), by about 15 orders of magnitude. Technical reason for strange result obtained in [10, 11] is the wrong interpretation of the second order pole structure of any amplitude containing the $n - \bar{n}$ transition. Instead of using the well developed Feynman diagram technique, the author [10, 11] tries to construct the space-time picture of the process by analogy with the vacuum case, which is misleading, see also discussion in conclusions.

\(^3\)There is, in fact some kind of competition between both methods, and final result will depend on the progress to be reached in both branches of experiments — with free neutrons and with neutrons bound in nuclei. Friedman and Gal [17] obtained the restriction $\tau_{n\bar{n}} > 3.3 \times 10^8 sec$ from the latest datum on $^{16}O$ stability and using the potential approach. Experiments with ultracold neutrons in a trap have been proposed and discussed in [3, 18], but not performed till now.
4. We continue our consideration with the case of the deuteron which is quite simple and instructive, and can be treated using standard diagram technique \(^4\). The point is that in this case there is no final state containing antineutron — it could be only the \(p\bar{n}\) state, by charge conservation. But this state is forbidden by energy conservation, since the deuteron mass is smaller than the sum of masses of the proton and antineutron. Therefore, if the \(n - \bar{n}\) transition took place within the deuteron, the final state could be only some amount of mesons. The amplitude of the process is described by the diagrams of the type shown in Fig. 2 and is equal to

\[
T(d \rightarrow mesons) = ig_{dnp}m_N\mu_{n\bar{n}} \int \frac{T(\bar{n}p \rightarrow mesons)}{(p^2 - m_N^2)((d - p)^2 - m_N^2)^2(2\pi)^4} \, d^4p. \tag{11}
\]

The constant \(g_{dnp}\) is normalized by the condition [20, 21, 22]

\[
\frac{g_{dnp}^2}{16\pi} = \frac{\kappa}{m_N} = \sqrt{\frac{\epsilon_d}{m_N}}, \tag{12}
\]

which follows, e.g. from the deuteron charge formfactor normalization \(F_d(t = 0) = 1\), see next section. \(\kappa = \sqrt{m_N\epsilon_d}, \epsilon_d \approx 2.22 \text{ MeV}\) being the binding energy of the deuteron. For the vertex \(d \rightarrow np\) we are writing \(2m_Ng_{dnp}\) to ensure the correct dimension of the whole amplitude.

The integration over internal 4-momentum \(d^4p\) in (11) can be made easily taking into account the nearest singularities in the energy \(p_0 = E\), in the non-relativistic approximation for nucleons. As we shall see right now, the integral

\[^4\text{It has been considered in fact in [19] within reasonable framework of diagram technique, however, the author has drawn later wrong conclusions from this consideration.}\]
over $d^3p$ converges at small $p \sim \kappa$ which corresponds to large distances, $r \sim 1/\kappa$. By this reason the annihilation amplitude can be taken out of the integration in some average point, and we obtain the approximate equality

$$T(d \rightarrow mesons) = g_{dnp}m_N\mu_{\bar{n}n}I_{dNN}T(\bar{n}p \rightarrow mesons)$$ (11a)

with

$$I_{dNN} = \frac{i}{(2\pi)^4} \int \frac{d^4p}{(p^2 - m_N^2)[(d - p)^2 - m_N^2]^2} \approx \frac{i}{(2\pi)^4(2m)^3} \int \frac{d^4p}{(p_0 - m_N - \bar{p}^2/(2m_N) + i\delta)(m_d - m_N - p_0 - \bar{p}^2/(2m_N) - i\delta)^2} =$$

$$= \int \frac{d^3p}{(2\pi)^38m_N[\kappa^2 + \bar{p}^2]^2} = \frac{1}{64\pi m_N\kappa},$$ (13)

This integral converges at small $|\bar{p}| \sim \kappa$, more details can be found in the next section. The decay width (probability) is, by standard technique,

$$\Gamma(d \rightarrow mesons) \approx \mu_{\bar{n}n}^2g_{dnp}^2I_{dNN}^2m_N \int |T(\bar{n}p \rightarrow mesons)|^2d\Phi(mesons),$$ (14)

$\Phi(mesons)$ is the final states phase space. Our final result for the width of the deuteron decay into mesons is

$$\Gamma_{d \rightarrow mesons} \approx \frac{\mu_{\bar{n}n}^2}{16\pi\kappa}m_N^2[v_0\sigma^{ann}(\bar{n}p)]_{v_0 \rightarrow 0} \approx \frac{\mu_{\bar{n}n}^2}{8\pi\kappa}m_N[p_{c.m.}\sigma^{ann}_{\bar{n}p}]_{p_{c.m. \rightarrow 0}},$$ (15)

where $p_{c.m.}$ is the (anti)nucleon momentum in the center of mass system. This result is very close to that obtained by L.Kondratyuk (Eq. (17) in [14]) in somewhat different way, using the induced $\bar{n}p$ wave function.

The annihilation cross section of antineutron with velocity $v_0$ on the proton at rest equals

$$\sigma(\bar{n}p \rightarrow mesons) = \frac{1}{4M_N^2v_0} \int |T(\bar{n}p \rightarrow mesons)|^2d\Phi(mesons).$$ (16)

According to PDG at small $v_0$, roughly, $[v_0\sigma^{ann}_{\bar{n}p}]_{v_0 \rightarrow 0} \approx (50 - 55)mb \approx (130 - 140)GeV^{-2}$. So, we obtain $\mu_{\bar{n}n} \leq 2.5 \times 10^{-24}eV$, or $\tau_{\bar{n}n} > 5 \times 10^8$ sec if we take

\footnote{The result Eq. (17) in [14] can be rewritten in our notations as

$$\Gamma_{d \rightarrow mesons} \approx 0.01\mu_{\bar{n}n}^2\frac{m_N^2}{\kappa}[v_0\sigma^{ann}_{\bar{n}p}]_{v_0 \rightarrow 0},$$ (17')

which differs from our result by some numerical factor, close to 1 and not essential for our conclusions.}
optimistically same restriction for the deuteron stability as it was obtained for the Fe nucleus, $T_d \simeq T_{Fe} > 6.5 \times 10^{31} yr$ [9]. Our result (15) is valid up to numerical factor of the order $\sim 1$, since we did not consider explicitly the spin dependence of the annihilation cross section and the spin structure of the incident nucleus. Same holds in fact for the results obtained in preceding papers, see [2, 14] e.g.

Additional suppression factor in comparison with the case of a free neutron is of the order of

$$\mu_{\bar{n}n}/\kappa \sim 10^{-31}$$

in agreement with our former rough estimate (9), and disappears, indeed, when the binding energy becomes zero. The binding energy of the deuteron should be very small, to give $\kappa \sim \mu_{\bar{n}n}$, to avoid such suppression. At such vanishing binding energy the nucleons inside the deuteron are mostly outside of the range of nuclear forces, similar to the vacuum case.

Results similar to (15) can be obtained for heavier nuclei, see [5, 6, 14, 16, 17]. The physical reason of such suppression is quite transparent and has been discussed in the literature long ago (see [2, 13, 15] e.g.): it is the localization of the neutron inside the nucleus, whereas no localization takes place in the vacuum case. In the case of the deuteron or heavier nucleus the annihilation of antineutron takes place, and final state is some continuum state containing mesons. By this reason the transition of the final state back to the incident nucleus is not possible in principle, and there cannot be oscillation of the type, e.g. $d \rightarrow \text{mesons} \rightarrow d$. This is important difference from the case of the free neutron.

5. As we noted previously, the presence of the second order pole in intermediate energy variable is characteristic for the processes with the neutron - antineutron transition, but it is in fact not a new peculiarity, it takes place also for the case of the nucleus formfactor with zero momentum transfer, $F_A(q = 0)$. Let us consider as an example the deuteron charge formfactor. In the zero range approximation it can be written as

$$F_d(q) = \frac{i(2mg_{dnp})^2}{(2\pi)^4} \int \frac{d^4p}{(p^2 - m_N^2)[(d - p)^2 - m_N^2][(d - p + q)^2 - m_N^2]}.$$ (17)

\[ \text{There is no final formula for } \Gamma_{d\rightarrow \text{mesons}} \text{ in [19] to be compared with our result (14), (15). Numerically, however, the result of [19] is in rough agreement with our and [14] estimate.} \]
Behind the zero range approximation $g_{dnp}$ should be considered as a function of the relative $n - p$ momentum, not as a constant. For $q = 0$ second order pole appears, and we come to the expression for $F(q = 0)$ containing the integral $I_{dNN}$ introduced above in Eq. (13):

$$F_d(0) = (2m_N g_{dnp})^2 I_{dNN}. \quad (18)$$

![Figure 3: The diagram describing the deuteron charge formfactor.](image)

In the nonrelativistic approximation, when only the nearest in energy $E = p_0$ singularities are taken into account, the integral over the energy has the structure

$$I_{dNN} \sim \int \frac{dE}{(E - a + i\delta)(E - b - i\delta)^2} = \frac{-2\pi i}{(a - b)^2}.$$ 

(19)

$a = m_N + \vec{p}^2/2m_N$, $b = m_d - m_N - \vec{p}^2/2m_N$, $a - b = \epsilon_d + \vec{p}^2/m_N$, and can be calculated using the lower contour, or the upper contour, with the help of formulas known from the theory of functions of complex variables. After this we obtain

$$F_d(q = 0) = \frac{g_{dnp}^2 m_N}{16\pi^3} \int \frac{d^3 p}{(\kappa^2 + \vec{p}^2)^2}. \quad (20)$$

Since $F_d(0) = 1$, this leads to the above mentioned normalization condition

$$g_{dnp}^2/(16\pi) = \sqrt{\epsilon_d/m_N}. \quad (7)$$

This relation between the constant $g_{dnp}$ and the binding energy of the weakly bound system (deuteron in our case) is known for a long time [20, 21, 22].

7As it is known from nonrelativistic diagram technique, the wave function of the deuteron in momentum representation is $\Psi_d(\vec{p}) = g_{dnp}/[4\pi^{3/2}(\kappa^2 + \vec{p}^2)]$, therefore, the normalization of the charge formfactor $F_d(0) = 1$ follows from the normalization of the deuteron wave function, which is also well known from quantum mechanics.
It was obtained in [20, 21, 22] using different methods, dispersion relation, for example.

If the infrared divergence discussed in [10, 11] takes place for the process of $n - \bar{n}$ transition in nucleus, it should take place also for the nucleus form-factor at zero momentum transfer. But it is well known not to be the case. There is no ”new limit on neutron - antineutron transition” [10]; instead, one should treat correctly singularities of the transition amplitudes in the complex energy plane, in particular, the second order pole contribution to the transition amplitudes. The author of [10, 11] tries to reconstruct the space-time picture of the process, but the correspondence of this picture to the well justified amplitude, as it appears from Feynman diagrams, is questionable. The infrared divergence discussed in [10, 11] is an artefact of this inadequate space-time picture of the whole process of $n - \bar{n}$ transition with subsequent antineutron annihilation. Another quite unrealistic consequence of this space-time picture is the nonexponential law of the nucleus decay.

Field-theoretical description of nuclear reactions and processes is potentially useful, it allows to study some effects which is not possible in principle to study in other way, e.g. relativistic corrections to different observables. One should be, however, very careful to treat adequately analytical properties of contributing amplitudes. E.g., in the case of the parity violating amplitude of $np \rightarrow d\gamma$ capture it was necessary to take into account contributions of all singularities (poles) of the amplitude in the complex energy plane, not only contributions of the nearest poles in energy variable, as it is made usually in nonrelativistic calculations. The nonrelativistic diagram technique developed up to that time turned out to be misleading for the case of physics problem considered in [23]. Besides, and it is the specifics of the processes with photon emission, the contact terms should be reconstructed to ensure the gauge invariance of the whole amplitude of photon radiation [23].

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8In his latest comment [24] Nazaruk makes the statement "In the correct model the neutron line entering the \(n\bar{n}\) transition vertex should be the wave function." It means that new rules are proposed in [10, 11] instead of well known Feynman rules. These "new rules" should be, at least, clearly formulated. and, second, these rules should allow to reproduce all well known results of nuclear theory. There is no need of further discussion.
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