$n\bar{n}$ transitions in nuclei

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It is shown that there is a double counting in the standard model of $n\bar{n}$ mixing in the medium, resulting in full cancellation of leading terms. The direct calculation of $n\bar{n}$ transition followed by annihilation is performed. The lower limit for the free-space $n\bar{n}$ oscillation time is $\tau_{n\bar{n}} \sim T_{n\bar{n}} > 10^{16}$ yr, where $T_{n\bar{n}}$ is the lifetime of neutron bound in a nucleus. This limit exceeds the previous one by 16 orders of magnitude.
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Any information on the occurrence of $n\bar{n}$ oscillation [1,2] is important in order to discriminate among various grand unified theories. The most direct limit on the free-space $n\bar{n}$ oscillation time $\tau_{n\bar{n}}$ is obtained using free neutrons: $\tau_{n\bar{n}} > 10^7$ s [3]. Alternatively, a limit can be extracted from the nuclear annihilation lifetime measured in proton-decay type experiments: $\tau_{n\bar{n}} > 10^8$ s$\sim$1 yr (see, for example, Refs. [4,5]). The process under consideration is (nucleus)$\rightarrow$($\bar{n}$ $-$nucleus)$\rightarrow$(annihilation products). The calculations were based on the potential model of $n\bar{n}$ mixing in the medium [4] or on the nonrelativistic diagram technique [5]. They predict the drastical suppression of $n\bar{n}$ transition in nuclei. However, $n\bar{n}$ conversion comes from the exchange of Higgs bosons with $m_H > 10^5$ GeV, so from the point of view of the microscopic theory (dynamic $n\bar{n}$ conversion [2], annihilation) the reasons for the suppression are not clear.

In this paper it is shown that the models used previously [4,5] are too crude or, more categorically, inapplicable to the problem under study. In particular, the potential model does not correspond to annihilation products in the final state. We perform the direct calculation of the process (nucleus)$\rightarrow$($\bar{n}$ $-$nucleus)$\rightarrow$(annihilation products). The interaction Hamiltonian is taken in the general form. Then the basic part of calculation is model independent. However, the amplitude in this case is singular. For solving the problem the approach with finite time interval [6] is used.

In the standard approach (later on referred to as potential model) the $n\bar{n}$ transitions in the medium are described by Schrödinger equations

$$(i\partial_t + \nabla^2/2m - U_n)n(x) = \epsilon\bar{n}(x),$$

$$(i\partial_t + \nabla^2/2m - U_{\bar{n}})\bar{n}(x) = \epsilon n(x). \qquad (1)$$

Here $\epsilon = 1/\tau_{n\bar{n}}$ is a small parameter [4]; U_n and $U_{\bar{n}}$ are the self-consistent neutron potential and the $\bar{n}$-nucleus optical potential, respectively. For U_n=const and $U_{\bar{n}}$=const in the lowest order in ϵ the probability of the process is

$$W_{\rm pot}(t) = 1 - |U_{ii}(t)|^2 = 2\,{\rm Im}\,T_{ii}(t),$$

$$T_{ii}(t) = i(\epsilon/\delta U)^2[1 - i\delta Ut - \exp(-i\delta Ut)], \qquad (2)$$

where U and T are the evolution operator and T operator, respectively; $U = 1 + iT$ and

$$\delta U = U_{\bar{n}} - U_n = {\rm Re}\,U_{\bar{n}} - i\Gamma/2 - U_n. \qquad (3)$$

Here $\Gamma \sim 100$ MeV is the annihilation width of $\bar{n}$-nucleus state.

What is described by $W_{\rm pot}(t)$? Let us take the imaginary part of Eq. (16) of Ref. [6]:

$$2\,{\rm Im}\,T_{ii}(t) = \epsilon^2 t^2 - \epsilon^2 \int_0^t dt_\alpha \int_0^{t_\alpha} dt_\beta 2\,{\rm Im}\,T^{\bar{n}}_{ii}(\tau), \qquad (4)$$

where $\tau = t_\alpha - t_\beta$. Here $T(t)$ and $T^{\bar{n}}(\tau)$ are the T operators of the whole process and $\bar{n}$-nucleus interaction, respectively. For the $T^{\bar{n}}$ operator $|i\rangle = |0\bar{n}_p\rangle$ and $\langle f|$ are the $\bar{n}$-nucleus and annihilation products, respectively. For the whole T operator $|i\rangle = |0n_p\rangle$ is the nucleus; the physical meaning of the final states will be cleared up later. The probability conservation $\Sigma_f |U_{fi}|^2 = 1$ gives us

$$2\,{\rm Im}\,T_{ii} = \sum_{f\neq i} |T_{fi}|^2 + |T_{ii}|^2,$$

$$\sum_{f\neq i} |T_{fi}(t)|^2 = W(t). \qquad (5)$$

The probability of the process $W(t)$ will be defined below; $|T_{fi}|^2 \sim \epsilon^2$, whereas $|T_{ii}|^2 \sim \epsilon^4$ [see Eqs. (2) and (15)]. So for the left-hand side (LHS) of Eq. (4) $2\,{\rm Im}\,T_{ii}(t) = W(t)$, which was taken into account in Eq. (2). For the T matrix of the $T^{\bar{n}}(\tau)$ operator Eq. (5) has the form

$$2\,{\rm Im}\,T^{\bar{n}}_{ii}(\tau) = \sum_{f\neq i} |T^{\bar{n}}_{fi}(\tau)|^2 + |T^{\bar{n}}_{ii}(\tau)|^2, \qquad (6)$$

$T^{\bar{n}}_{ii} = \langle 0\bar{n}_p|T^{\bar{n}}|0\bar{n}_p\rangle$, $T^{\bar{n}}_{fi} = \langle f|T^{\bar{n}}|0\bar{n}_p\rangle$. The $\bar{n}$-nucleus interaction is a nonperturbative process and $|T^{\bar{n}}_{ii}|^2 \sim \Sigma_{f\neq i} |T^{\bar{n}}_{fi}|^2$. Now Eq. (4) is

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$$W(t)=\epsilon^2t^2-\epsilon^2\int_0^t dt_\alpha\int_0^{t_\alpha}dt_\beta|\langle 0\bar{n}_p|T^{\bar{n}}(\tau)|0\bar{n}_p\rangle|^2$$

$$-\epsilon^2\int_0^t dt_\alpha\int_0^{t_\alpha}dt_\beta\sum_{f\neq i}|\langle f|T^{\bar{n}}(\tau)|0\bar{n}_p\rangle|^2. \quad (7)$$

Let us calculate $T_{ii}^{\bar{n}}$ and $T_{fi}^{\bar{n}}$ in the framework of the potential model. The wave function of initial state is described by the equation

$$i\frac{\partial\Phi}{\partial t}=H_0\Phi, \quad (8)$$

$H_0=-\nabla^2/2m+U_n$. At the moment $t=0$ the interaction δU is turned on. We have

$$i\frac{\partial\Psi}{\partial t}=(H_0+\delta U)\Psi, \quad (9)$$

$\Psi(0)=\Phi(0)$. The projection to the initial state and T matrix at $t=\tau$ are

$$\langle\Phi|\Psi\rangle=U_{ii}^{\bar{n}}(\tau)=\exp(-i\delta U\tau), \quad (10)$$

$$T_{ii}^{\bar{n}}(\tau)=i[1-\exp(-i\delta U\tau)],$$

$$\sum_{f\neq i}|T_{fi}^{\bar{n}}(\tau)|^2=1-|U_{ii}^{\bar{n}}(\tau)|^2=1-e^{-\Gamma\tau}=W_{\bar{n}}(\tau), \quad (11)$$

where $W_{\bar{n}}(\tau)$ is the $\bar{n}$-nucleus decay probability. Note that Γ corresponds to all $\bar{n}$-nucleus interactions followed by annihilation. However, the main contribution comes from the annihilation without rescattering of $\bar{n}$ [5], because $\sigma_{ann}>2\sigma_{sc}$. Substituting these expressions in Eq. (7), one obtains the potential model result (2).

Therefore, the approach with the finite time interval was verified by the example of the exactly solvable potential model. It is involved in Eq. (4) as a special case. Solving Eqs. (1) by the method of Green's functions we will obtain the same results. We have started from Eq. (4) only for verification of the finite time approach.

Let us return to Eq. (7). It is at least unclear. (1) The first term is the free-space $n\bar{n}$ transition probability. The matrix elements $T_{ii}^{\bar{n}}$ and $T_{f\neq i}^{\bar{n}}$ describe transitions ($\bar{n}$-nucleus) $\rightarrow$($\bar{n}$-nucleus) and ($\bar{n}$-nucleus)$\rightarrow$(annihilation products), respectively. So the first and the second terms correspond to the $\bar{n}$ nucleus in the final states. However, in the experiment only the annihilation products are detected (the $\bar{n}$ nucleus is unobservable) and the result should be expressed only in the terms of $T_{f\neq i}^{\bar{n}}$. Moreover, the $\bar{n}$-nucleus decays into final states identical with the states given by the third term. This suggests that the potential model contains the double counting. Expression $1-|U_{ii}|^2$ from Eq. (2) describes the inclusive decay of the initial state and so the $n\bar{n}$ transition with the $\bar{n}$ nucleus in the final state is also included in $W_{\rm pot}$, unless additional limits are imposed. To exclude the double counting the annihilation products in the final state should be fixed. (2) Let $|\delta Ut|\ll 1$. (This is the case in some other problems.) When $\Gamma=0$, the third term equals to zero. When $\Gamma\neq 0$, the contribution of the third term is negative and $dW/d\Gamma<0$, whereas the opening of the new channel (annihilation) should increase W.

How big is the probable error? The contributions of the second and third terms are $x_2=-\epsilon^2t^2/2+F_2$, $x_3=-\epsilon^2t^2/2+F_3$. The functions $F_{2,3}$ contain the terms proportional to t and $\exp(-i\delta Ut)$. So the ϵ^2t^2 term produced by the third term is fully canceled. This is a consequence of double counting. Therein lies the reason of the discrepancy between our result and the result of the potential model.

As noted in [6], Eqs. (11) and (2) can also be obtained by means of microscopic variant of the potential model (zero angle rescattering diagrams of $\bar{n}$). In this case the Hamiltonian of the $\bar{n}$-medium interaction is $H=\delta U$. The same calculation was repeated by Dover *et al.* [4]. They substitute $H=-i\Gamma/2$ in Eq. (4) and obtain Eq. (2). On the basis of this and only this they refute the result of Ref. [6]. In other words they refute our limit because it differs from the prediction of the potential model ($H=-i\Gamma/2$ [4]).

What is wanted is $\Sigma_{f\neq i}|T_{fi}|^2$, where $\langle f|$ is the annihilation product. It is connected with the diagonal matrix element by Eq. (5):

$$2\,\mathrm{Im}\,T_{ii}=\sum_{f\neq i}T_{fi}^*T_{fi}. \quad (12)$$

Calculation of T_{ii} is determined by the right-hand side (RHS) of Eq. (12): the cut corresponding to T_{ii} must contain only annihilation products, which is not in accordance with Eq. (7). It includes redundant states $f=(\bar{n}$-nucleus) forbidden by the unitarity condition. The relation (12) is not fulfilled. Also the eigenfunctions of $H_0+\delta U$ do not form the complete orthogonal set. Due to this fact the $\bar{n}$ nucleus (described by $U_{\bar{n}}$) also cannot appear in Eq. (12) as the intermediate state. So the model (1) is inapplicable in our case because it leads automatically to incorrect matrix element T_{ii}. Elimination of redundant trajectories from T_{ii} means the direct calculation of T_{fi}.

In Ref. [6] the first and the third terms were taken into account. The second one was omitted. The first term corresponds to low density limit and is meaningful for $n\bar{n}$ transitions in the gas. This scheme is not quite correct here. In this paper we perform the direct calculation of the process (nucleus)$\rightarrow$($\bar{n}$-nucleus)$\rightarrow$(annihilation products). We have

$$\langle f|U(t,0)-I|0n_p\rangle=iT_{fi}(t)=\sum_{k=1}^{\infty}(-i)^{k+1}\langle f|\int_0^t dt_1\ldots\int_0^{t_{k-1}}dt_k\int_0^{t_k}dt_\beta H(t_1)\ldots H(t_k)H_{n\bar{n}}(t_\beta)|0n_p\rangle, \quad (13)$$

where

$$H(t)=(\text{all } \bar{n}\text{-medium interactions})-U_n,$$

$$H_{n\bar{n}}(t)=\epsilon\int d^3x(\bar{\Psi}_{\bar{n}}\Psi_n+\text{H.c.}), \qquad (14)$$

$H+H_{n\bar{n}}=H_I$. Here $|0n_p\rangle$ is the state of the medium containing the neutron with the 4-momenta $p=(\mathbf{p}_n^2/2m+U_n,\mathbf{p}_n)$, $\langle f|$ represents the annihilation products; $H_{n\bar{n}}$ is the oscillation Hamiltonian [4]. In the case of the formulation of the S-matrix problem $(t,0)\to(\infty,-\infty)$ Eq. (13) in the momentum representation includes the singular propagator $G=1/(\epsilon_n-\mathbf{p}_n^2/2m-U_n)\sim 1/0$. Taking into account that $H_{n\bar{n}}|0n_p\rangle=\epsilon|0\bar{n}_p\rangle$, we change the order of integration and obtain

$$T_{fi}(t)=-\epsilon\int_0^t dt_\beta i T_{fi}^{\bar{n}}(t-t_\beta),$$

$$iT_{fi}^{\bar{n}}(\tau)=\sum_{k=1}^{\infty}(-i)^k\int_{t_\beta}^{t}dt_1\ldots\times\int_{t_\beta}^{t_{k-1}}dt_k\langle f|H(t_1)\ldots H(t_k)|0\bar{n}_p\rangle, \qquad (15)$$

where $|0\bar{n}_p\rangle$ is the state of the medium containing the $\bar{n}$ with the 4-momenta p; $\tau=t-t_\beta$. The 4-momenta of n and $\bar{n}$ are equal. $T_{fi}^{\bar{n}}$ is an exact amplitude of $\bar{n}$-nucleus decay. It includes all the $\bar{n}$-nucleus interactions followed by annihilation. The expression for $T_{fi}(t)$ was obtained in perfect analogy to Eq. (4) that can be considered as a test for Eq. (15).

The two-step process was reduced to the annihilation decay of the $\bar{n}$ nucleus. (The slightly different method is the separation of the antineutron Green's function [6].) It is seen from Eqs. (13) and (15) that both pre- and post-$n\bar{n}$ conversion spatial wave functions of the system coincide:

$$|0n_p\rangle_{sp}=|0\bar{n}_p\rangle_{sp}. \qquad (16)$$

$\bar{n}$ appears in the state with $\delta U=0$. We would like to stress that in the potential model (1) the picture of $\bar{n}$-nucleus formation is exactly the same: in Eq. (4) for $T_{ii}=\langle 0n_p|T|0n_p\rangle$ and $T_{ii}^{\bar{n}}=\langle 0\bar{n}_p|T^n|0\bar{n}_p\rangle$ condition (16) was fulfilled. Hereafter, the potential model of the $\bar{n}$-medium interaction (block T^n) was used and W_{pot} was reproduced, which confirms the picture of $\bar{n}$-nucleus formation given above. Solving Eqs. (1) by method of Green's functions we will obtain the same results, including Eq. (16). The equality of vectors of state (16) is also evident from the continuity of solution of Eqs. (1).

In both models the first stage of the process ($n\bar{n}$ conversion) is described identically. The basic difference centers on the next stage—annihilation. In the potential model $T_{ii}^{\bar{n}}$ is calculated (as a result the self-energy part $\Sigma=\delta U$ appears) and is used in Eq. (7), which is wrong. We calculate $T_{fi}^{\bar{n}}$ starting from the same point (16). The result will be expressed through Γ [see Eqs. (19), (11)], but not through $\delta U=\text{Re }\delta U-i\Gamma/2$, as is usually the case in decay calculations. The standard δU dependence is manifested in scattering problems, when the diagonal matrix element T_{ii} in the RHS of Eq. (12) should be taken into account. It corresponds to the observarble process—zero angle scattering of the incident particle.

The characteristic annihilation time of $\bar{n}$ is $\Delta=1/\Gamma\sim 10^{-23}$ s. When $\tau\gg\Delta$, $T_{fi}^{\bar{n}}(\tau)$ reaches its asymptotic value $T_{fi}^{\bar{n}}$:

$$T_{fi}^{\bar{n}}(\tau\gg\Delta)=T_{fi}^{\bar{n}}(\infty)=T_{fi}^{\bar{n}}=\text{const}. \qquad (17)$$

The expressions of this type are the basis for all S-matrix calculations. [Measurement of any process corresponds to some interval τ. So it is necessary to calculate $U(\tau)$. The replacement $U(\tau)\to S(\infty)$ is equivalent to Eq. (17).] Let us take $t\gg\Delta$. From Eqs. (15) and (17) we have

$$T_{fi}(t)=-i\epsilon\left[\int_0^{t-\Delta}dt_\beta T_{fi}^{\bar{n}}(t-t_\beta)+\int_{t-\Delta}^{t}dt_\beta T_{fi}^{\bar{n}}(t-t_\beta)\right]\sim -i\epsilon t T_{fi}^{\bar{n}}. \qquad (18)$$

The contribution of the second term is negligible since $|T_{fi}^{\bar{n}}(\tau)|^2\leqslant 1$. The probability of the whole process is

$$W(t)=\sum_{f\neq i}|T_{fi}(t)|^2\sim\epsilon^2t^2\sum_{f\neq i}|T_{fi}^{\bar{n}}(t)|^2\sim\epsilon^2t^2, \qquad (19)$$

where Eq. (11) has been taken into account. The value $\epsilon^2t^2=t^2/\tau_{n\bar{n}}^2$ is the free-space $n\bar{n}$ transition probability. Due to the annihilation channel $n\bar{n}$ conversion is practically unaffected by the medium. So $\tau_{n\bar{n}}\sim T_{n\bar{n}}$, where $T_{n\bar{n}}$ is the oscillation time of neutron bound in a nucleus. In order to find the limit for $\tau_{n\bar{n}}$ from experimental data on nuclear stability, the distribution (19) should be used (but not the exponential decay law). Let us take N_n, T_0, ϵ_1, and θ as the total number of neutrons under observation, the observation time, the overall $n\to\bar{n}$ detection efficiency, and the average number of observable $n\to\bar{n}$ events, respectively. From the inequality

$$N_n(T_0/\tau_{n\bar{n}})^2(\epsilon_1/\theta)<1 \qquad (20)$$

one obtains $\tau_{n\bar{n}}>10^{16}$ yr, where the values $T_0=1.3$ yr, $N_n=2.4\times10^{32}$, $\epsilon_1=0.33$, and $\theta=2.3$ [7] were used.

Our previous result [6] is different from Eq. (19) only by a factor of 2. However, in Ref. [6] we used the limit $T_{n\bar{n}}>4.3\times10^{31}$ yr [7] deduced from the experimental data by means of exponential decay law which does not agree with Eq. (19).

Let us return to the reason for the enormous quantitative disagreement between our result and the potential model one. The strong sensitivity of the results should be expected. Really, in the momentum representation the S-matrix amplitude M_s, corresponding to the $n\bar{n}$ transition followed by annihilation [see Eq. (14)] diverges

$$M_s=\epsilon\frac{1}{\epsilon_n-\mathbf{p}_n^2/2m-U_n}M\sim\frac{1}{0}, \qquad (21)$$

where M is the annihilation amplitude. These are infrared singularities conditioned by zero momentum transfer in the ϵ

vertex. It can be seen that $M_s \sim 1/0$ for any bound state wave function of a neutron (i.e., for any nuclear model). On the other hand in the potential model the energy is not conserved and becomes complex: $M_A \to M_A + \delta U$ (M_A is the nuclear mass). In this case we have instead of Eq. (21)

$$M_{\rm pot} = \epsilon \frac{1}{\epsilon_n - \mathbf{p}_n^2/2m - U_n^-} M = \epsilon \frac{1}{\delta U} M. \tag{22}$$

This is a potential model amplitude. Really, the process width is $\Gamma_{\rm pot} = \int d\Phi |M_{\rm pot}|^2/2M_A = \epsilon^2 \Gamma / |\delta U|^2 = W_{\rm pot}/t$, which coincides with Eq. (2) when $|\delta U t| \gg 1$. It is seen that: (1) There is a double counting in M and G with respect to H. $M_{\rm pot}$ does not agree with Eq. (14) as well. (2) $\delta U = 0$ is the singular point and due to zero momentum transfer $q=0$ in the vertex corresponding to $H_{n\bar{n}}$ we are in this point. So the result is extremely sensitive to δU. (Usually, in the reactions and decays the momentum transferred is $q \neq 0$. In this case the δU dependence of G is masked by q: $G^{-1} = (\epsilon_n - q_0) - (\mathbf{p}_n - \mathbf{q})^2/2m - U_n - \delta U$. We deal with the 2-tail and $q = 0$.]

Comparing Eqs. (21) with (18) one sees that in principle the limit $\delta U \to 0$ corresponds to the replacement

$$1/\delta U \to t. \tag{23}$$

Certainly, we do not set $\delta U = 0$. U_n^- is not introduced at all. In the calculation of Eq. (13) the multiplier t [see Eq. (18)] arises automatically instead of $1/\delta U$ in the potential model, or $1/\Delta q$ in the case $q \neq 0$. When $q \neq 0$ in the ϵ-vertex, Eq. (13) leads to the usual S-matrix result (see below). The formal reason for the differences in the results is the full cancellation of the terms $\sim t^2$ in Eq. (7). The erroneous structure of Eq. (7) is caused by the nonperturbative and two-step nature of the process. The fact that $q=0$ increases the disagreement.

An additional comment is necessary regarding the t dependence of the whole process probability $W(t)$. Equation (19) has been obtained in the lowest order in ϵ. The exact distribution $W_{pr}(t)$ which accounts for the all orders in ϵ is unknown. However, W is the first term of the expansion of W_{pr} and we can restrict ourselves to the lowest order $W_{pr} = W$, as it is usually the case for rare decays. W_{pot} is also calculated in the lowest order in ϵ.

The protons must be in a very early stage of the decay process. Thus the realistic possibility is considered [8–11] that the proton has not yet entered the exponential stage of its decay but is, instead, subject to non-exponential behavior which is rigorously demanded by quantum theory for sufficiently early times. At first sight, since $\tau_{n\bar{n}} > 10^{16}$ yr for $n\bar{n}$ mixing in a nuclei, the non-exponential behavior should be expected too. In fact, there is one more problem: we deal with the two-step process. When trying to calculate M_s and Γ_s in the framework of standard S-matrix theory we get $\Gamma_s \sim 1/0$. So the decay law $\exp(-\Gamma_s t)$ is irrelevant and it is necessary to deduce the distribution $W(t)$ as was done above.

We have to mention the main points made by Krivoruchenko's preprint [12]. (1) The $n\bar{n}$ transition followed by annihilation (two-step nuclear decay) and motion of particle in the classical field are two different problems. Describing the first one by Eqs. (1) we understand that this is an effective procedure. From a formal standpoint in the first and second cases the potentials are complex and real, respectively. Unfortunately, sometimes the literal analogy between these problems is drawn [12]. (2) The initial Eq. (11) of Ref. [12] must describe the $n\bar{n}$ transition followed by annihilation. However, the LHS of Eq. (11) is free of $\bar{n}$-nucleus interaction at all. The RHS contains annihilation width Γ (we stress this point) and coincides with the potential model result. We also would be glad to get the result without calculations, but some difficulties emerge in reaching this goal.

The interaction responsible for the $n\bar{n}$ conversion is ultraweak. Therefore, the n-nucleus interaction in the initial state should be taken into account exactly. The neutron line entering into the $n\bar{n}$ transition vertex should be the wave function of the bound state [see Eq. (8)], but not the propagator, as in the model based on the diagram technique [5,13]. As a result, in this model the $n\bar{n}$ transition is possible only between the acts of interactions of an oscillating particle and a nucleus. These interactions lead to total suppresion of $n\bar{n}$ conversion, which is incorrect. This can be understood using the analogy with β decay and taking into account that $m_H > 10^5$ GeV.

Some additional remarks on Ref. [13] are necessary. (1) The picture described in Sec. 1 is valid only for a simple interaction operator. We deal with products of operators [see Eq. (13) of this paper]. (2) In Sec. 2 it is claimed that the amplitude should be singular at $B \to 0$, where B is the binding energy. Accordingly, as $B \to 0$, the amplitude obtained is $|A_1|^2 \sim 1/0$. In fact, they are the usual infrared singularities mentioned above, which must be avoided. The correct model should reproduce the low density limit $W(t) = \epsilon^2 t^2$. (3) The cut corresponding to the diagonal matrix element (18) is completely free of annihilation products.

We try to calculate the process amplitude starting from Eq. (14). The S-matrix theory gives Eq. (21). The approach with a finite time interval is infrared-free. Its verification for the diagrams with $q=0$ was made above by the example of the potential model. For nonsingular diagrams the test is obvious. Let us have $q \neq 0$ in the ϵ vertex. The appropriate calculations with a finite time interval (adiabatic hypothesis should be used) give the S-matrix result (we stress this fact because it means the verification of the approach): $T_{fi} = i\epsilon'(1/\Delta q)\mathcal{T}^n_{fi}$, where $\mathcal{T}^n_{fi}$ is the S-matrix amplitude of annihilation of virtual $\bar{n}$ with the 4-momenta $k = p - q$. Comparing with Eq. (18) one sees that the limit $\Delta q \to 0$ corresponds to the replacement $1/\Delta q \to t, \mathcal{T}^n_{fi} \to T^n_{fi}$ [compare with Eq. (23)]. The similar problem for matrix element T_{ii} was solved in Ref. [14].

The main results of this paper are given in the abstract. In the next paper the following statements will be proved: (1) All the results are true for any nuclear model. (2) The contribution of the corrections is negligible. (3) Further investigation and verification of the approach will be presented as well. In our opinion, it makes sense to look at some other problems on oscillation of particles in a medium from the standpoint given above.

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