Delta Decay in the Nuclear Medium

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Abstract

The Δ decay in the nuclear medium is calculated in the relativistic meson-nucleon model. The delta spreading width is calculated and compared with the Pauli-blocked πN decay width. The influence of relativistic mean fields is also studied. We stress the importance of understanding the delta spreading width in interpreting experiments involving delta resonances.

I. INTRODUCTION

The understanding of the properties of the Δ resonance in the nuclear medium, especially the branching ratios of its various decay modes, is important for the analysis of scattering processes involving Δ-hole excitations.

A Δ, which mainly decays into πN in free space, acquires additional decay channels in the medium. For example, it can decay through the mechanism, Δ + N → N + N, commonly known as “delta spreading width” which can be seen as reabsorption of the virtual pion through a particle-hole excitation of the nucleus. Indeed, a number of experiments involving (p, n) reactions [1] show that the additional channels of delta decay seem to dominate over the usual πN decay.

Recently, we have proposed that the spreading width could be the source of errors in measuring the nucleon axial form factor from charged-current neutrino-nucleus scattering [2]. Namely, in the BNL experiment of neutrino-nucleus scattering [3] where neutrino beams with an average energy of 1.3 GeV are used, the data are usually fitted to the theoretical curve calculated in a Fermi gas model of the nucleus in order to extract the axial form factor. In this approach Δ-hole excitations are assumed to be completely excluded by experimentally rejecting events involving pions in the final state. However, when the Δ obtains a non-zero

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spreading width, some events which generate $\Delta$-hole excitations lead to pionless decays of the $\Delta$ which cannot be excluded by merely detecting the pions. To estimate these events quantitatively, it is important to know the branching ratio of the spreading width to the $\pi N$ decay width before one can deduce the nucleon axial form factor accurately.

Further, the knowledge of the branching ratio can be used to improve Monte Carlo simulations of neutrino detectors, e.g. in the Kamiokande experiment, where the flavor ratio of atmospheric neutrinos are measured [1]. It was suggested that pions produced from the charged-current neutrino-nucleus reactions cause uncertainties in identifying the flavor of the neutrinos created in the earth’s atmosphere [2]. In these experiments the main source of pions is believed to come from delta decays. Therefore, a thorough knowledge of the branching ratio between pionic and non-pionic decays of deltas is crucial for estimating the pion events in neutrino-nucleus scattering of atmospheric neutrinos.

In this work, we present a calculation of the delta decay width and its modification in nuclear matter within the framework of the Walecka model [3]. The model is a relativistic meson-nucleon model which has been quite successful in describing nuclear properties, including response functions [4]. Crucial ingredients of this model are the presence of large vector and scalar self-energies which modify the nucleon properties in the medium enhancing relativistic effects. Our concern is to study the modification of the delta decay width driven by these relativistic effects. Nonrelativistic calculations of the delta decay in the medium can be found in Refs. [5,6].

II. FORMALISM

In this section, we present the relativistic calculation of the $\Delta$ decay in the nuclear medium. First we review the calculation of the $\pi N$ decay which has been studied in Ref. [7]. Then we will present our calculation of the spreading width and its modification due to nuclear environment.

In free space, a delta decays to $\pi N$. The decay width for this mode can be obtained from the imaginary part of the delta self-energy as shown in Fig. 1 (a). Formally, the decay width is given in terms of the self-energy $\Sigma_{\mu\nu}$ and the Rarita-Schwinger spinor of the delta, $\Delta^\mu$,

$$\Gamma_f = -2 \Im(\Delta^\mu \Sigma_{\mu\nu} \Delta^\nu)$$

(1)

To calculate the self-energy, we introduce the $\pi N \Delta$ interaction Lagrangian,

$$\mathcal{L}_{\pi N \Delta} = \frac{f_\Delta}{m_\pi} \bar{\Delta}^\mu \mathbf{T} N \partial_\mu \pi + \text{h.c.}$$

(2)

Here $N$, $\pi$, and $\Delta$ are the nucleon, pion and delta field, respectively. The $2 \times 4$ isospin matrices $\mathbf{T}$ satisfy [8]

$$T^i(T^i)^j = \delta^{ij} - \frac{1}{3} \tau^i \tau^j.$$

(3)

There is some ambiguity in the form of the Lagrangian of a spin-3/2 field, which affects the off-shell behavior of the particle. We will not discuss this complication in the following
because the delta is on-shell in our calculation. Using this Lagrangian and standard Feynman rules, one writes the self-energy in free space,

$$ \Sigma_{\mu\nu}(q) = \left( \frac{f_\Delta}{m_\pi} \right)^2 i \int \frac{d^4k}{(2\pi)^4} g_\mu q_\nu D_0(q) G_f(k) , $$  

where the pion \([D_0(q)]\) and nucleon \([G_f(k)]\) propagators in free space are given as

$$ D_0(q) = \frac{1}{q^2 - m_\pi^2 + i\epsilon} , $$  

$$ G_f(k) = \frac{k^2 + m}{k^2 - m^2 + i\epsilon} . $$  

Putting these into Eq. (1), one can easily calculate the free decay width,

$$ \Gamma_f = \left( \frac{f_\Delta}{m_\pi} \right)^2 \frac{(E_k + m)}{12\pi m_\Delta} |k|^3 , $$  

with

$$ k^2 = \left( \frac{m_\Delta^2 + m_\pi^2 - m^2}{2m_\Delta} \right)^2 - m_\pi^2 ; \quad E_k = \sqrt{k^2 + m^2} . $$

Here \(m, m_\Delta, \) and \(m_\pi\) represent the mass of nucleon, delta and pion respectively. By comparing Eq. (7) with its experimental value of 115 MeV, one can fix the coupling constant \(f_\Delta = 2.15.\)

In the nuclear medium viewed as a free Fermi gas of nucleons, the \(\pi N\) decay is suppressed by Pauli blocking, since the phase space available to the decaying nucleon is restricted by the occupied Fermi sea. Indeed, at nuclear saturation density the \(\pi N\) decay is completely blocked for a delta at rest in the medium. When the delta is moving, however, the decay sphere of the nucleon in momentum space becomes an ellipsoid, which only partially overlaps with the Fermi sphere \([11]\). Therefore, the delta obtains a non-zero energy-dependent width.

The Pauli blocking can be incorporated by replacing \(G_f(k)\) with

$$ G_0(k) = G_f(k) + 2\pi\delta(k^2 - m^2)\Theta(k_0)\Theta(k_F - |k|)(k + m) . $$  

Here \(k_F\) is the Fermi momentum whose value at the nuclear saturation is 1.3 fm\(^{-1}\). Note that in this form the nucleon propagator is given in the rest frame of the nuclear matter where the delta is moving with a certain relative velocity. From the following identity \([11]\) for the \(\Delta\) spinor with momentum \(p^\mu = (E_\Delta, \vec{p}_\Delta)\)

$$ \sum_\sigma \Delta^\mu(\sigma)\Delta^\nu(\sigma) = -\frac{\not{p}_\Delta + m_\Delta}{2m_\Delta} \left[ g^{\mu\nu} - \frac{2p_\mu^\Delta p_\nu^\Delta}{3m_\Delta^2} + \frac{1}{3m_\Delta} (p_\mu^\Delta\gamma^\nu - p_\nu^\Delta\gamma^\mu) - \frac{1}{3} \gamma^{\mu\nu} \right] \equiv -\frac{\not{p}_\Delta + m_\Delta}{2m_\Delta} P_{3/2}^{\mu\nu} , $$  

the decay width averaged over the delta spin becomes

$$ \bar{\Gamma}_{\pi N}(|\vec{p}_\Delta|) = -\frac{1}{8m_\Delta} \frac{f_\Delta}{m_\pi} \left( \frac{f_\Delta}{m_\pi} \right)^2 i \int \frac{d^4k}{(2\pi)^4} \Theta(q_0)\delta(q^2 - m_\pi^2)\Theta(k_0 - E_F)\delta(k^2 - m^2) \times \text{Tr} \left[ g_\mu q_\nu (k + m)(\not{p}_\Delta + m_\Delta)P_{3/2}^{\mu\nu} \right] . $$  


$p_\Delta$ is the delta momentum relative to the nuclear matter, corresponding to a relative velocity $v = |p_\Delta/E_\Delta|$. Note that we have made use of Cutkosky rule [12] in deriving this. The integration over the nucleon momentum $k$ leads to

$$\bar{\Gamma}_{\pi N}(|p_\Delta|) = \frac{1}{192\pi} \left( \frac{f_{\Delta}}{m_\pi} \right)^2 \frac{E_2}{|p_\Delta|m_\Delta} \left[ \frac{(m_\Delta + m_\pi)^2 - m^2}{(m_\Delta - m_\pi)^2 - m^2} \right] \left[ \frac{(m_\Delta + m)^2 - m^2}{(m_\Delta - m)^2 - m^2} \right] ,$$

where

$$E_2 = \left\{ \min(E_\Delta, E_\pi) - \max(E_F, E_-) \right\} \Theta(E_2) ,$$
$$E_\pm = \max \left\{ E_F, \frac{(m_\Delta^2 + m^2 - m_\pi^2)E_\Delta \pm |p_\Delta|\sqrt{[(m_\Delta - m)^2 - m_\pi^2][(m_\Delta + m)^2 - m_\pi^2]}}{2m_\Delta^2} \right\} .$$

A similar calculation has also been done in Ref. [10]. The authors of Ref. [10] also discussed the results for different magnetic quantum numbers of the delta, which is qualitatively similar to the averaged result.

The free Fermi gas calculation can be improved by incorporating mesonic mean fields. In a mean-field approximation to the Walecka model, the propagation of the delta and the nucleon is modified by the presence of constant scalar and vector mean fields which are determined from nuclear saturation properties.

The $\pi N$ decay width in this approximation can be obtained with the replacements

$$m_\Delta \to m_\Delta^* = m_\Delta - S_\Delta ,$$
$$E_\Delta \to E_\Delta^* = \sqrt{m_\Delta^2 + p_\Delta^2} ,$$
$$m \to m^* = m - S_N ,$$
$$k^\mu \to k'^\mu = k^\mu - V^\mu_N ,$$
$$p_\Delta^\mu \to p_\Delta'^\mu = p_\Delta^\mu - V^\mu_N .$$

The nucleon self-energies $(V^\mu_N, S_N)$ are assumed to be the same as the delta self-energies $(V^\mu_\Delta, S_\Delta)$ as is motivated from the universal coupling assumption. In symmetric nuclei, the space components of the vector self-energy are averaged to zero. Also its time component can be eliminated by appropriately shifting integration variables. However, the scalar self-energy $S_N$ does contribute. Its value at a given nuclear density can be determined by solving the self-consistency equation [8]

$$m^* = m - \frac{g_s^2}{m_s^2} \frac{4}{(2\pi)^3} \int_0^{k_F} d^3k \frac{m^*}{\sqrt{k^2 + m^2}} ,$$

where the ratio of the scalar meson coupling to its mass, $g_s^2/m_s^2$, is determined from the equilibrium properties of nuclear matter [8]. At nuclear saturation of $k_F = 1.3$ fm$^{-1}$, Eq. (13) yields $m^* = 510$ MeV.

In addition to the $\pi N$ decay, a delta in the medium can also decay into 2p-1h through the process, $\Delta + N \to N + N$. The lowest-order self-energy diagram for this process is shown in Fig. 1 (b). This spreading width can be calculated in the same way as the $\pi N$ decay by replacing the free pion propagator with the pion propagator $D(q)$ including particle-hole polarization.
\[ D(q) = \frac{1}{q^2 - m_\pi^2 - \Pi(q) + i\epsilon} . \]  

(14)

We proceed by introducing the interaction Lagrangian for nucleons and pion using the pseudovector coupling for the pion,

\[ \mathcal{L}_{\pi NN} = \frac{f_N}{m_\pi} N\gamma^\mu\gamma_5 N \partial_\mu \pi , \]

(15)

with \( f_N = 1.01 \). Using the interaction Lagrangian Eq. (15) we can easily determine the lowest order particle-hole polarization \( \Pi_{ph}(q) \):

\[ \text{Re} \Pi_{ph}(q) = \left( \frac{f_N}{m_\pi} \right)^2 \frac{q^2 m^2}{\pi^2|q|} \int_m^{E_F} dE_k \log \left| \frac{(q^2 - 2|k||q|)^2 - 4q_0^2 E_k^2}{(q^2 + 2|k||q|)^2 - 4q_0^2 E_k^2} \right| , \]

(16)

\[ \text{Im} \Pi_{ph}(q) = \left( \frac{f_N}{m_\pi} \right)^2 \frac{q^2 m^2 E_1}{\pi|q|} \]

(17)

with

\[ E_1 = E_F - E_\pi \quad ; \quad E_\pi = \min(E_F, E_{\text{max}}) \]

(18)

\[ E_{\text{max}} = \max \left[ m, E_F - q_0, -\frac{q_0}{2} + \frac{|q|}{2} \sqrt{1 - 4m^2/q^2} \right] . \]

(19)

The \( \Delta \) decay calculated with Eq. (14) contains two decay channels. First, the contribution from the usual \( \pi N \) decay which can be identified with the pole contribution of the pion propagator Eq. (5) with the shifted mass

\[ \tilde{m}_\pi^2 = m_\pi^2 + \Pi(q^2 = \tilde{m}_\pi^2) . \]

(20)

If the lowest-order nucleon-hole polarization is used for \( \Pi(q) \), then, in some kinematics, Eq. (20) is solved by negative \( \tilde{m}_\pi^2 \) indicating a tachyonic pole. This pole can be removed by introducing the Landau-Migdal parameter \( g' \) which is incorporated by writing

\[ \Pi(q) = \frac{q^2 \Pi_{ph}}{q^2 + g' \Pi_{ph}} . \]

(21)

In our calculation, we take a conventional value of the Landau-Migdal parameter, \( g' = 0.7 \). One has to be a little bit careful with the pole contribution since the residue of the propagator is now changed due to the real part of \( \Pi(q) \). However, the modification of \( \pi N \) decay due to the real part of Eq. (21) is very small.

The second contribution arises from the imaginary part of the polarization Eq. (21). This part can be identified with the particle-hole decay channel of the virtual pion generating the spreading width of the \( \Delta \). Note that the imaginary part of the particle-hole polarization is only non-zero at space-like momenta of the pion. Therefore, the spreading width can be separated from the \( \pi N \) decay in the medium due to their different kinematic region. Replacing \( D_0(q) \) by the expression Eq. (14) in Eq. (2) and taking the contribution arising from the imaginary part of the polarization, we obtain the spreading width in the nuclear rest frame,
\[ \tilde{\Gamma}_\Delta(|p_\Delta|) = \frac{1}{12m^3\Delta\pi^2} \left( \frac{f_\Delta}{m_\pi} \right)^2 \int_{E_F}^{E_k} dE_k \int_{-1}^{1} d\cos\theta \]

\[ \times \frac{\text{Im}\Pi(q)\Theta(-q^2)}{[q^2 - m^2_\pi - \text{Re}\Pi(q)]^2 + [\text{Im}\Pi(q)]^2} \left( k \cdot p_\Delta + m_\pi \Delta \right) [q^2m^2_\Delta - (p_\Delta \cdot q)^2] , \quad (22) \]

which can be solved numerically. The lowest order spreading width shown in Fig. (1) (b), which is usually referred to as the spreading width, can be obtained from Eq. (22) by taking the lowest-order contribution from the polarization, only, and by neglecting the Landau-Migdal correction. Note that, unlike in the case of the free \( \pi N \) decay, the pion momentum, \( q^2 \), is not restricted to its on-shell value in the course of the integration. Therefore, one should consider the momentum dependence of the \( \pi N\Delta \) and \( \pi NN \) vertices. In this regard, we employ the monopole form for the form factor,

\[ F(q^2) = \frac{\Lambda^2 - m^2_\pi}{\Lambda^2 - q^2} , \quad (23) \]

in our discussion. This is to take into account the finite size of the nucleons as well as the short-range effects of NN interactions. However, the value of the cut-off parameter \( \Lambda \) is in doubt. In principle, this should be determined by summing the vertex diagrams of the underlying field theory. Here we use the Bonn potential values [14]; for \( \pi NN \), \( \Lambda = 1.8 \text{ GeV} \) and for \( \pi N\Delta \), \( \Lambda = 0.85 \text{ GeV} \). Finally, in order to implement the effects due to the nuclear self-energies, the replacements in Eq. (12) are done for the calculation of the spreading width.

In addition to the particle-hole polarization, \( \Delta \)-hole excitations can contribute. This means that a delta rescatters in the medium by \( \Delta + N \rightarrow \Delta + N \), which can be obtained by taking the imaginary part of \( \Delta \)-h polarization. The difference to the particle-hole polarization is that the imaginary part has finite values even for positive \( q^2 \). Depending on the four-momentum the pion pole is shifted into the complex plane as the pion obtains a width from the decay channel to a \( \Delta \)-h excitation in the medium. The rescattered \( \Delta \) subsequently decays leading to more complex decay channels, for example \( 2N - 1h - 1\pi \), etc. The possible influence of this contribution will be studied in a forthcoming publication [15].

III. RESULT

We now present the result for the various delta widths. Our concern is to understand the momentum dependence of the decay width as well as the role of the nuclear mean fields.

In the following, we will show our results calculated at half, normal and twice the nuclear saturation density with the corresponding Fermi momenta \( k_F = 203 \text{ MeV} \), 256 MeV, and 323 MeV, respectively. The self-consistency equation of Eq. (13) yields an effective nucleon mass \( m^* = 510 \text{ MeV} \) at the saturation density, \( m^* = 715 \text{ MeV} \) at half saturation density, and \( m^* = 261 \text{ MeV} \) at twice nuclear density. The reduction of the nucleon mass in the medium arises due to the strong scalar self-energy. In principle, the self-energies are functions of the momentum of the propagating nucleon [16]. Specifically, the absolute values of the self-energies decrease as the momentum increases. This momentum dependence could be interesting especially in the high momentum region. However, in most of the momentum
range of our concern, the momentum dependence of the self-energies is not expected to be important. The result with half the saturation density is applicable to \((p, n)\) reactions where the interactions most likely occur on the nuclear surface. On the other hand, the result at saturation density could be used in lepton-nucleus scattering as the incoming lepton can probe the region deeply inside the nucleus. The result with higher density could be interesting in the heavy-ion collision even though the neglect of contributions from the lower continuum of the nucleons in our calculation is questionable.

Figures 2 show the results for the \(\pi N\) decay width as a function of the delta momentum \(|\mathbf{p}_\Delta|\). Fig. 2 (a) shows the result of free Fermi gas calculation while Fig. 2 (b) includes the nuclear mean fields. At saturation density or at twice the saturation density, the decay width increases as the delta momentum increases. This is because the overlap between the momentum ellipsoid of the decay nucleon and the Fermi sphere is getting smaller and the Pauli blocking becomes less important. Indeed, as is shown by the dashed curve, the ellipsoid is completely out of the Fermi sphere for \(|\mathbf{p}_\Delta| > 650\) MeV, and the \(\pi N\) decay width is the same as its value in free space. A similar trend can be observed for the mean field calculation in Fig. 2 (b). At half saturation density (the solid curve in both figures), however, Pauli blocking is less significant as could be expected. Our results indicate that the mean field combined with the Pauli blocking reduces the \(\pi N\) decay width substantially from its free value of 115 MeV which agrees with Wehrberger and Wittman [10].

Unlike the \(\pi N\) decay, the spreading width, which is represented by the mechanism \(\Delta + N \rightarrow N + N\), is plotted in Figs. 3. Figure 3 (a) shows our calculated spreading width without the mean fields and Figure 3 (b) includes the mean fields. First note that the momentum dependence of the spreading width is not as significant as the \(\pi N\) decay. This means, when added to the \(\pi N\) decay width, the momentum dependence of the in-medium decay width is mainly driven by the Pauli blocking.

Conventionally the spreading width is scaled as the nuclear density [17]. This is reflected in the Fermi gas results in Fig. 3 (a) where the dashed curve is about a factor of 2 smaller than the dot-dashed curve and the solid curve is about a half of the dashed curve. But this kind of scaling is not applicable in the mean field case. Therefore, scaling the spreading width with density has to be applied with caution. Our mean field results at the saturation density which vary from 75 MeV to 57 MeV qualitatively agrees with the conventionally adopted value of 70 MeV in the nonrelativistic calculations [17,18]. Also at this density, the branching ratio of pionic decay to the spreading width, \(\Gamma_{\pi N}/\Gamma_s\), varies from 0 to 1.33 indicating that a significant portion of deltas leads to the pionless decays. However the role of the spreading width compared to the \(\pi N\) decay becomes smaller at half saturation density.

On the other hand, the mean field result for the spreading width at the double saturation density is rather abrupt. Its magnitude is only \(6 \sim 12\%\) of the Fermi gas result. This kind of huge reduction due to the mean fields might come from neglecting the vacuum contribution to the particle-hole loop in Hartree approximation. At this density in our simple calculation, the effective mass of nucleon is 261 MeV, much smaller than its free value. Therefore, antinucleons can be easily excited which could lead to the appreciable contribution from the vacuum. More work on the vacuum contribution should be done [15]. However, as indicated in Fig. 3 (a) and (b), the spreading width has a very strong dependence on the mean field given our approximations.

\(\Delta\)-hole contributions in neutrino-nucleus scatterings have been discussed in Ref. [21].
There it was stressed that the pionic versus non-pionic decay could be an important source of uncertainty in determining the nucleon axial form factor or estimating the pion events in atmospheric neutrinos. For example, in Ref. [2], we have found that the theoretical error of about 0.1 GeV in the extracted axial mass $M_A$ is expected from the events due to delta-hole excitations. This conclusion is drawn by assuming that the branch ratio between the spreading width and the $\pi N$ decay is roughly one. However, our calculation in this work show that the branching ratio has a clear momentum dependence which should be properly taken into account. In this regard, our results, especially the results at saturation density, could provide important constraints.

In summary, we have calculated the delta spreading width in the framework of a relativistic nucleon-meson model. Our concern is to present the branching ratio of the spreading width to $\pi N$ decay so that it can be used for the analysis of measurements of the nucleon axial form factor [3] or for the atmospheric neutrino experiments [4]. $\pi N$ decay is substantially reduced by Pauli blocking and mean fields. Also, mean fields play an important role in the spreading width. We have found that the spreading width is large even though its relative strength to the $\pi N$ decay depends on the nuclear density and the relative motion of the delta with respect to the nuclear medium.

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FIGURES

FIG. 1. Diagrams for $\Delta$ self-energy. Figure (a) is the self-energy diagram representing the $\pi N$ decay and figure (b) is the lowest diagram for the spreading width.

FIG. 2. Pionic decay of $\Delta$ with respect to the delta momentum. (a) is for the $\pi N$ decay without mean fields and (b) includes the mean fields. The solid lines show the results for the half saturation density, the dashed lines for the saturation density and the dot-dashed lines for twice the saturation density.

FIG. 3. Spreading width of $\Delta$. (a) shows the result without mean fields and (b) includes the mean fields. The solid lines show the results for the half saturation density, the dashed lines for the saturation density and the dot-dashed lines for twice the saturation density.
Figure 1
Figure 2
Figure 3