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Nuclear Physics A699 (2002) 165c–168c

NUCLEAR PHYSICS A

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Eta photoproduction off three body nuclei

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The process of coherent photoproduction of η-mesons on light nuclei is investigated on the bases of microscopic Finite Rank Approximation method. Numerical calculations are performed for the ^{3}He and ^{3}H targets and different inputs for η-nucleon interaction. Comparison with the modified optical model approach is given.

1. INTRODUCTION

The study of the η-nucleus interaction is motivated by various reasons. The first one is purely of nuclear nature and is addressing the possibility of forming quasi–bound states or resonances [1] in the η-nucleus system. Another reason is related to the study of the $S_{11}(1535)$ resonance structure and charge symmetry breaking effects, which may appear due to the $\eta - \pi_0$ mesons admixture.

The η-nucleus interaction can manifest itself in the final states of certain nuclear reactions, such as photoproduction, for example. Theoretical analysis of (γ,η)-reactions on nuclei is hampered by the three major problems: the unknown off-shell behavior of the two-body $\gamma N \to \eta N$ amplitude, inaccuracies in the description of the nuclear target as a many-body system, and rescattering effects in the final state. To minimize these uncertainties we employ fully microscopic Finite Rank Approximation method modified to include the electromagnetic interaction.

2. FORMALISM

Let us briefly discuss first the FRA method [2] describing the elastic scattering of η-mesons off A nucleons. To derive the FRA equations, we start from the many–body Lippmann–Schwinger equation for the total T–operator, and after certain algebraic manipulations we obtain the following equivalent system of equations

$$\begin{aligned} T(z) &= \sum_{i=1}^{A} T_i^0(z) + \sum_{i=1}^{A} T_i^0(z)\,\frac{1}{z-H_0}\,H_A\,\frac{1}{z-H_0-H_A}\,T(z) \\ T_i^0(z) &= t_i(z) + t_i(z)\,\frac{1}{z-H_0}\,\sum_{j\neq i}^{A} T_j^0(z)\,, \end{aligned} \tag{1}$$

PII S0375-9474(01)01488-9

where A is the number of nucleons, z the total ηA–energy, H_0 the Hamiltonian of the free motion of η with respect to the *c.m.* of the nucleus, T_i^0 are the auxiliary operators, and t_i the two–body η-nucleon T–operators in the many–body space.

Having in mind, that for the lightest nuclei we have only one bound state and that we are interested in the low energy η-nucleus scattering processes, we may retain in the spectral representation of H_A the first term: $H_A \approx \mathcal{E}_0|\psi_0\rangle\langle\psi_0|$. In the FRA this approximation is the only one used while the rest of the calculations can be done exactly.

To introduce the electromagnetic interaction in equations (1) we follow the same procedure as the one used in [3]. Instead of the scalar ηN t-matrix we introduce the two-channel generalization of it, and this transformation will produce the similar matrix form for the solutions of equations (1):

$$t_i \rightarrow \begin{pmatrix} t_{\gamma\gamma} & t_{\gamma\eta} \\ t_{\eta\gamma} & t_{\eta\eta} \end{pmatrix}_i, \qquad T_i^0 \rightarrow \begin{pmatrix} T^0_{\gamma\gamma} & T^0_{\gamma\eta} \\ T^0_{\eta\gamma} & T^0_{\eta\eta} \end{pmatrix}_i. \tag{2}$$

Here $(t_{\gamma\gamma})_i$ describes the Compton scattering, $(t_{\eta\gamma})_i$ the photoproduction process, and $(t_{\eta\eta})_i$ the elastic ηN scattering on the nucleon number i.

It is technically more convenient to consider the reaction of η-photoabsorption, and then the photoproduction cross-section can be obtained by applying the detailed balance principle. In the first order of electromagnetic interaction we get from (1) and (2):

$$\begin{aligned} T^{\gamma\eta}(z) &= \sum_{i=1}^{A} T^0_{\gamma\eta}(z)_i + \sum_{i=1}^{A} T^0_{\gamma\eta}(z)_i|\psi_0\rangle \frac{\mathcal{E}_0}{(z-H_0)(z-H_0-\mathcal{E}_0)} \langle\psi_0|T^{\eta\eta}(z)\,, \\ T^0_{\gamma\eta}(z)_i &= t_{\gamma\eta}(z)_i + t_{\gamma\eta}(z)_i \frac{1}{z-H_0} \sum_{j\neq i}^{A} T^0_{\eta\eta}(z)_j\,. \end{aligned} \tag{3}$$

Here $T^{\gamma\eta}$ is the photoabsorption full T-matrix, $(T^0_{\gamma\eta})_i$ is the corresponding auxiliary operator, whereas $T^{\eta\eta}$ describes elastic scattering of η-meson off a nucleus, satisfying (1).

3. INPUT DATA

To obtain the necessary nuclear wave functions ψ_0 we employed the Malfliet–Tjon NN–potential [4] and the integro–differential equation approach (IDEA) [5]. This approach is based on the Hyperspherical Harmonic expansion method applied to Faddeev-type equations, and in fact it is fully equivalent to the exact Faddeev equations for S–wave projected potentials.

To find the $(T^0_{\eta\eta})_i$ and $T^{\eta\eta}$ from (1) it is necessary to know the η-nucleon t-matrix. We use separable form for it:

$$t_{\eta\eta}(k', k; z) = g(k')\,\tau_{\eta\eta}(z)\,g(k), \tag{4}$$

where the vertex function for the ηN $\leftrightarrow$ N* is chosen as $1/(k^2+\alpha^2)$, being is of Yukawa-type in configuration space. The value of range parameter used $\alpha = 3.316\,\mathrm{fm}^{-1}$ was determined in Ref. [6]. The first propagator $\tau^{BW}_{\eta\eta}(z)$ has a simple Breit-Wigner form

$$\tau^{BW}_{\eta\eta}(z) = \frac{\lambda}{z - E_0 + i\Gamma/2}, \tag{5}$$

which is motivated by the dominance of the S_{11}–resonance. The strength parameter λ is chosen to reproduce the η-nucleon scattering length $a_{\eta N}$, the imaginary part of which accounts for the flux losses into the πN channel. The value of the scattering length $a_{\eta N}$ is not accurately known, here we use $a_{\eta N} = 0.55 + i0.3\,\mathrm{fm}$ which is in the range of uncertainties. The E_0 and Γ are the parameters of the S_{11} resonance [7].

An alternative propagator $\tau^{LS}_{\eta\eta}(z)$ was found by solving the Lippmann-Shwinger equation for an energy-dependent separable potential with the same form-factors as it was proposed in [8]:

$$\tau^{LS}_{\eta\eta}(z) = -\frac{4\pi\,\alpha^3}{\mu_{\eta N}}\,\frac{\lambda_{\eta N}(E_{\mathrm{res}} - z) + C E_{\mathrm{res}}}{E_{\mathrm{res}} - z - [\lambda_{\eta N}(E_{\mathrm{res}} - z) + C E_{\mathrm{res}}]/(1 - ip/\alpha)^2}\,. \tag{6}$$

Here α is the same parameter as in the form-factors, E_{res}, $\lambda_{\eta N}$, and C are new ones, and p is the on-energy shell momentum.

It was experimentally shown [9] that, at low energies, the reaction $\gamma N \to \eta N$ proceeds mainly via the formation of the S_{11}-resonance. This implies that $t^{\gamma\eta}$ in the near–threshold region can be written in a separable form similarly to Eq. (4). To construct such a separable t–matrix, we used the results of Ref. [10] where the $t^{\eta\gamma}$ (which is equal to $t^{\gamma\eta}$) was obtained as an element of a multi-channel T-matrix on the energy shell. For our calculations, we extended this T-matrix off the energy shell, using the Yamaguchi form-factors which become unit on the energy shell. It is generally believed, that $t^{\gamma\eta}$ is different for neutron and proton. We assumed that they have the same functional form and differ by a constant factor $t^{\gamma\eta}_n = A\,t^{\gamma\eta}_p$. A multipole analysis gives for this factor the following estimate [11]: $A = -0.84 \pm 0.15$.

4. RESULTS

We calculated the total cross-sections for ^{3}H and ^{3}He using different $t_{\eta\eta}$-matrices: t_{BW} and t_{LS}. The results can be seen in Fig. 1. The photo-neutron to photo-proton ratio $t_{\mathrm{n}}/t_{\mathrm{p}}$ for the elementary $t^{\gamma\eta}$-matrix was fixed at $A = -0.75$. Below the threshold, calculations with the different forms of $t_{\eta\eta}$ give the same size for the cross-sections, especially in the case of ^{3}H target. The common feature of all curves (independently from the type of two-body inputs) is the cusp near the threshold of the full disintegration of nuclei.

The dependence of the cross-section on the parameter A for the ^{3}He is shown in Fig. 2, three curves correspond to $A = 1, -0.75$ and -1 (from up to down). We used here t_{LS}, but for the Breit-Wigner $t_{\eta\eta}$-matrix the dependence is the same. One can see that the cross-section has a rather weak dependence for negative values of the parameter A while the limiting values $A = \pm 1$ are well separated. The same calculations for the ^{3}H show that there are more possibilities to distinguish negative values of A in this case.

At the same figure we compare our results with those obtained with a modified optical model [12] (filled triangles). The difference in the results is drastic, especially in the energy region of the threshold for the photoproduction of η-mesons. The reason of this difference is probably due to the first order approximation of the optical potentials used. In contrast to the low-energy π-nucleon scattering, where the two-body πN-interaction is weak and thus the first order optical π-nuclear potential is sufficient, the ηN interaction is dominated by the near threshold S_{11} resonance and one should expect that higher order

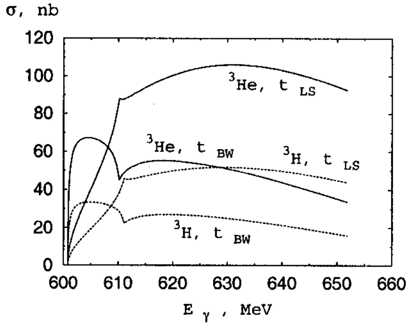

Figure 1. Total photoproduction cross-section for the ^{3}H and ^{3}He with different elementary $t_{\eta\eta}$-matrices.

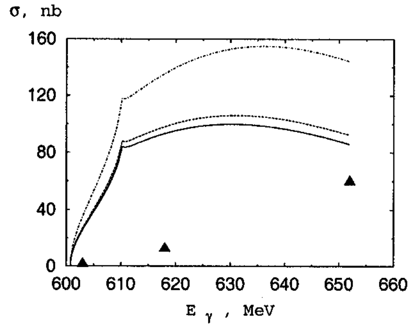

Figure 2. Total photoproduction cross-section on the ^{3}He for different values of parameter A.

terms are important. In the FRA approach these terms are naturally included in the formalism.

ACKNOWLEDGMENTS

The authors gratefully acknowledge financial support from the University of South Africa, the Division for Scientific Affair of NATO (grant CRG LG 970110), and the DFG-RFBR (grant 436 RUS 113/425/1). One of authors (V.B.B.) willing to thank Physikalishes Institut of Bonn University for hospitality.

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