Coherent η photoproduction from nuclei in a relativistic impulse approximation approach

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We study coherent η photoproduction from nuclei in a relativistic impulse approximation approach. For the elementary production amplitude we use a standard relativistic parametrization based on a set of four Lorentz- and gauge-invariant amplitudes. The photonuclear amplitude is evaluated without recourse to a nonrelativistic reduction; the full relativistic structure of the amplitude is maintained. On general arguments we show that the coherent process is sensitive to only one of the elementary amplitudes. Moreover, we show that the nuclear structure information is fully contained in the ground-state tensor density. The tensor density is evaluated in a mean-field approximation to the Walecka model and it is shown to be sensitive to relativistic effects. Distortion effects are incorporated through an η-nucleus optical potential that is computed in a simple ''$t\rho$'' approximation. [S0556-2813(97)02105-5]

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I. INTRODUCTION

The coherent photoproduction of pseudoscalar mesons (such as π^0 and η) offers a unique opportunity for the investigation of nucleon-resonance formation and propagation through the nuclear medium. The advent of more powerful and sophisticated machines—such as TJNAF and MAMI—will challenge, now more than ever, our theoretical understanding of this fundamental process.

Although experimentally challenging, the coherent photoproduction reaction offers numerous advantages. Because the nucleus remains in its ground state, all nucleons participate coherently in the reaction leading to enhanced cross sections relative to incoherent processes to discrete nuclear states. Theoretically, the coherent process acts as a spin-isospin filter by selecting a particular (scalar-isoscalar) component of the elementary photoproduction amplitude. In this way the reaction can be used to discriminate between various theoretical models that provide an equally good description of the elementary process. Moreover, all nuclear structure information is contained in, at most, a few ground-state densities. Indeed, nonrelativistic plane-wave-impulse-approximation analyses suggest that the photonuclear amplitude is directly proportional to the isoscalar (or matter) density [1–5]. Since the ground-state charge density is very accurately determined from electron-scattering experiments, the coherent reaction becomes—after making some plausible assumptions about the neutron density—an ideal tool in the determination of the scalar-isoscalar component of the elementary amplitude, and its possible modification in the many-body environment. Finally, because the coherent process is sensitive to the whole nuclear volume, one can place stringent limits on the form of the meson-nucleus optical potential. This unique character of the coherent reaction has already led to an experimental effort to measure coherent pion photoproduction from ^{4}He [6] and ^{12}C [7], and coherent η photoproduction from ^{4}He [8] up to photon energies of about 800 MeV; possibilities for extensions to higher energies and other nuclei, exist both at the Bonn ELSA facility and at TJNAF.

Most theoretical analyses of the elementary process start with a model-independent parametrization of the photoproduction amplitude in terms of four Lorentz- and gauge-invariant amplitudes [9]. It is then customary to evaluate this amplitude between on-shell nucleon spinors, thereby leading to the well-known Chew-Goldberger-Low-Nambu form for the photoproduction operator in terms of Pauli—rather than Dirac—spinors [9]. For the calculation of the photonuclear reaction one usually adopts the impulse approximation [1–5]; one assumes that the elementary (on-shell) amplitude is not modified in the many-body environment. For closed-shell (spin-saturated) nuclei the photonuclear process then becomes, in the plane-wave limit, a simple product of the elementary scalar-isoscalar amplitude times the Fourier transform of the ground-state matter density. One then improves on the plane-wave description by incorporating distortions into the propagation of the outgoing meson.

In this paper we propose to carry out the above theoretical program—without recourse to a nonrelativistic reduction of the elementary amplitude. The main difference relative to the standard nonrelativistic approach stems from the fact that in the present relativistic framework the lower components of the nucleon spinors will be determined dynamically, rather than from the free-space relation. The nuclear structure information will be contained in a few ground-state densities that will be computed using a mean-field approximation to the Walecka model [10]. It is worth mentioning that the Walecka model has enjoyed considerable success in describing ground-state properties of many nuclei and, indeed, rivals some of the best available nonrelativistic calculations.

II. FORMALISM

The differential cross section in the center-of-momentum frame (c.m.) for the coherent photoproduction of η mesons is given by [3]

$$\left(\frac{d\sigma}{d\Omega}\right)_{\rm c.m.}=\left(\frac{M_T}{4\pi W}\right)^2\left(\frac{q_{\rm c.m.}}{k_{\rm c.m.}}\right)\frac{1}{2}\sum_{\lambda}|T_\lambda|^2, \qquad (1)$$

where M_T is the mass of the target nucleus, W is the total energy in the c.m. frame, while $k_{\rm c.m.}$ and $q_{\rm c.m.}$ are the three-momentum of the photon and η meson in the c.m. frame, respectively. Note that the formalism is independent of the mass of the produced meson so it can be applied without modification to the coherent photoproduction of any pseudoscalar meson. The relativistic-invariant amplitude for the coherent photoproduction from ($J^\pi=0^+$; $T=0$) nuclei can be written in the following way:

$$T_\lambda=\epsilon_\mu(\hat{\mathbf{k}},\lambda)\langle A(p');\eta(q)|J^\mu|A(p)\rangle, \qquad (2)$$

where $\epsilon_\mu(\hat{\mathbf{k}},\lambda)$ is the polarization vector of the photon which couples to a conserved electromagnetic current given by

$$\langle A(p');\eta(q)|J^\mu|A(p)\rangle\equiv\varepsilon^{\mu\nu\alpha\beta}k_\nu q_\alpha p_\beta\frac{1}{W}F_0(s,t). \qquad (3)$$

Here $p(p'=p+k-q)$ is the four momentum of the initial-(final) nucleus and $\varepsilon^{\mu\nu\alpha\beta}$ is the relativistic Levi-Civita symbol ($\varepsilon^{0123}\equiv-1$). Note that all the dynamical information about the coherent process is contained in a single Lorentz-invariant form factor $F_0(s,t)$, which depends on the Mandelstam variables $s=(k+p)^2$ and $t=(k-q)^2$. One can now carry out the appropriate algebraic manipulations to obtain the following—model-independent—form for the coherent photoproduction cross section:

$$\left(\frac{d\sigma}{d\Omega}\right)_{\rm c.m.}=\left(\frac{M_T}{4\pi W}\right)^2\left(\frac{q_{\rm c.m.}}{k_{\rm c.m.}}\right)\left(\frac{1}{2}k_{\rm c.m.}^2q_{\rm c.m.}^2\sin^2\theta_{\rm c.m.}\right)|F_0(s,t)|^2, \qquad (4)$$

where $\theta_{\rm c.m.}$ is the scattering angle in the c.m. frame.

We now proceed to compute the Lorentz invariant form factor in a relativistic impulse approximation. For the elementary $\gamma N\to\eta N$ amplitude we use a model-independent parametrization given in terms of four Lorentz- and gauge-invariant amplitudes:

$$T(\gamma N\to\eta N)=\sum_{i=1}^{4}A_i(s,t)M_i\,. \qquad (5)$$

For the invariant matrices we use the following standard set [3]:

$$M_1=-\gamma^5\not\epsilon\not k, \qquad (6a)$$

$$M_2=2\gamma^5[(\epsilon\cdot p)(k\cdot p')-(\epsilon\cdot p')(k\cdot p)], \qquad (6b)$$

$$M_3=\gamma^5[\not\epsilon(k\cdot p)-\not k(\epsilon\cdot p)], \qquad (6c)$$

$$M_4=\gamma^5[\not\epsilon(k\cdot p')-\not k(\epsilon\cdot p')]. \qquad (6d)$$

Note that, although standard, this is only one particular form for the elementary amplitude. Many other choices—all of them equivalent on shell—are possible [4,9]. However, in the absence of a detailed microscopic model of the elementary process it becomes ambiguous on how to take the elementary amplitude off shell. The above form of the elementary amplitude, or one very close to it, is ubiquitous in theoretical studies, so we adopt it here as well.

To proceed, we perform some simple algebraic manipulations on the elementary amplitude so that the parity and Lorentz transformation properties of the various bilinear covariants become manifest. That is,

$$T(\gamma N\to\eta N)=(F_T^{\alpha\beta}\sigma_{\alpha\beta}+F_P\gamma_5+F_A^\alpha\gamma_\alpha\gamma_5), \qquad (7)$$

where tensor, pseudoscalar, and axial-vector amplitudes have been introduced:

$$F_T^{\alpha\beta}=\frac{1}{2}\varepsilon^{\mu\nu\alpha\beta}\epsilon_\mu k_\nu A_1(s,t), \qquad (8a)$$

$$F_P=2[(\epsilon\cdot p)(k\cdot p')-(\epsilon\cdot p')(k\cdot p)]A_2(s,t), \qquad (8b)$$

$$F_A^\alpha=[(\epsilon\cdot p)k^\alpha-(k\cdot p)\epsilon^\alpha]A_3(s,t)+[(\epsilon\cdot p')k^\alpha-(k\cdot p')\epsilon^\alpha]A_4(s,t). \qquad (8c)$$

Note that for this particular form of the elementary amplitude no scalar nor vector invariants appear. For closed-shell nuclei an enormous simplification ensues, as a result of the pseudoscalar and axial-vector ground-state densities being identically zero. This implies that the coherent reaction is sensitive to only the A_1 component of the elementary amplitude. Moreover, all the nuclear-structure information is contained in the ground-state tensor density [10]. This is in contrast to nonrelativistic approaches in which the coherent amplitude is proportional to the conserved vector density [1–5]. Thus, in a relativistic plane-wave impulse approximation, the Lorentz-invariant form factor acquires a remarkable simple form:

$$F_0^{\rm PW}(s,t)=iA_1(\tilde{s},t)\rho_T(Q)/Q. \qquad (9)$$

Note that $\tilde{s}$ represents the effective (or optimal) value of the Mandelstam variable s at which the elementary amplitude should be evaluated [2] and $Q\equiv|\mathbf{k}_{\rm c.m.}-\mathbf{q}_{\rm c.m.}|\simeq\sqrt{-t}$.

The elementary η-production amplitude used is constructed in an effective Lagrangian approach as detailed in Refs. [11,12]. The dynamical content of the amplitude consists of (i) s- and u-channel nucleon Born terms, (ii) t-channel vector-meson exchange, and (iii) intermediate excitation of spin-1/2 and spin-3/2 resonances in the s and u channels (resonating only in the s channel). The free parameters in the elementary model were fixed by achieving good fits to the latest $p(\gamma,\eta)p$ [13] and $d(\gamma,\eta)pn$ [14] data. The inclusion of two intermediate resonances, the $S_{11}(1535)$ and $D_{13}(1520)$, generates good fits. As is well known, the $S_{11}(1535)$ clearly dominates the elementary reaction, but the $D_{13}(1520)$ contribution is needed to reproduce the measured angular distributions. The model parameters we use here are identical to those determined in Ref. [12]. Figure 1 shows a sample comparison of this elementary model to some of the

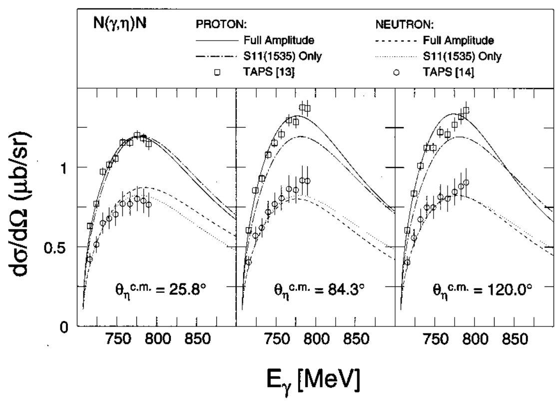

FIG. 1. A comparison of the $p(\gamma,\eta)p$ and $n(\gamma,\eta)n$ elementary amplitudes to a sample of the data from Refs. [13] (proton) and [14] (extracted neutron). Calculations are shown for the full amplitude [Born terms, vector-meson exchange, and intermediate $S_{11}(1535)$ and $D_{13}(1520)$ resonance excitation], as well as for the intermediate $S_{11}(1535)$ alone.

data from Refs. [13] and [14]; the data for the elementary neutron amplitudes are from Ref. [14], as extracted from the measured $d(\gamma,\eta)np$ data. Good agreement is found between the elementary model and the data for all energies and angles measured, and for both the proton and extracted neutron amplitudes. Also shown in Fig. 1 is the contribution from the $S_{11}(1535)$ excitation term only. Clearly, intermediate $S_{11}(1535)$ resonance excitation dominates the production amplitude, but is not sufficient on its own to adequately describe the data. Note that, when embedded into the impulse-approximation calculation for coherent production from closed-shell nuclei, intermediate $S_{11}(1535)$ excitation is suppressed due to spin-isospin considerations, and the role of the $D_{13}(1520)$ is enhanced. Indeed, it is just this suppression of the dominant s-wave term—allowing enhancement of the nondominant contributions—which initially provoked interest in this coherent reaction [3].

The ground-state tensor density is evaluated in a mean-field approximation to the Walecka model. For closed-shell nuclei the relativistic spinors can be classified according to a generalized angular momentum κ and can be written in a two-component representation; i.e.,

$$\mathcal{U}_\alpha(\mathbf{x})=\frac{1}{r}\begin{bmatrix} g_a(r)\mathcal{Y}_{+\kappa m}(\hat{\mathbf{r}}) \\ if_a(r)\mathcal{Y}_{-\kappa m}(\hat{\mathbf{r}})\end{bmatrix};\quad [\alpha\equiv(a;m)=(n,\kappa;m)], \tag{10}$$

where the upper and lower components are expressed in terms of spin-spherical harmonics defined by

$$\mathcal{Y}_{\kappa m}(\hat{\mathbf{r}})\equiv\langle\hat{\mathbf{r}}|l_{1/2}jm\rangle;\quad j=|\kappa|-\frac{1}{2};$$

$$l=\begin{cases}\kappa, & \text{if } \kappa>0;\\ -1-\kappa, & \text{if } \kappa<0.\end{cases} \tag{11}$$

For spin-saturated nuclei only three ground-state densities do not vanish [10]; these are the scalar and timelike-vector densities—used to compute the mean-field ground state—and the tensor density defined by

$$[\rho_T(r)\hat{r}]^i=\sum_\alpha^{\text{occ}}\bar{\mathcal{U}}_\alpha(\mathbf{x})\sigma^{0i}\mathcal{U}_\alpha(\mathbf{x}), \tag{12a}$$

$$\rho_T(r)=\sum_a^{\text{occ}}\left(\frac{2j_a+1}{4\pi r^2}\right)2g_a(r)f_a(r). \tag{12b}$$

For the coherent reaction it is only the Fourier transform of the latter that is needed:

$$\rho_T(Q)=4\pi\int_0^\infty dr r^2 j_1(Qr)\rho_T(r). \tag{13}$$

Note that the tensor density is linear in the lower (or small) component of the single-particle wave function; this is in contrast to the scalar and vector densities where the lower component enters as an $(f/g)^2$ correction. Thus, the tensor density is small and not well constrained by experiment. Yet, it is as fundamental as the vector density measured in electron scattering. Moreover, the tensor density is interesting because it is sensitive to the relativistic components of the wave function. Indeed, the mean-field approximation to the Walecka model is characterized by the existence of large Lorentz scalar and vector potentials that are responsible for a substantial enhancement of the lower components of the single-particle wave functions. This enhancement is at the heart of the phenomenological success enjoyed by the Walecka model. Thus, the enhancement of the tensor density (to be shown later) represents an inescapable prediction of the model.

Although simple and illuminating, the plane-wave formalism must be modified so that final-state interactions between the outgoing meson and the nucleus can be incorporated. This is done via an η-nucleus optical potential of the $t\rho$ form

$$2\omega_\eta V(r)=-b\rho_V(r). \tag{14}$$

Here $\rho_V(r)$ is the conserved vector density—also computed in the mean-field approximation to the Walecka model—and b is a two-body (ηN) parameter taken from Ref. [3] and parametrized in the following way:

$$b(p_{\text{lab}})\equiv(\alpha+\beta p_{\text{lab}}+\gamma p_{\text{lab}}^2)^{-1}, \tag{15a}$$

$$\alpha=(+0.136,-0.052)\ \text{fm}^{-1}, \tag{15b}$$

$$\beta=(+0.035,-0.072), \tag{15c}$$

$$\gamma=(-0.061,+0.009)\ \text{fm}. \tag{15d}$$

The Klein-Gordon equation for the outgoing η meson can now be solved in each angular-momentum channel resulting in a coherent form factor expressed in a partial-wave series. That is,

$$F_0^{DW}(s,t)=iA_1(\tilde{s},t)\rho_T^{DW}(\mathbf{q}_{\text{c.m.}},\mathbf{k}_{\text{c.m.}})/q_{\text{c.m.}}, \tag{16}$$

where the distorted-wave tensor density has been introduced:

$$\rho_T^{DW}(\mathbf{q},\mathbf{k})=\sum_{l=1}^{\infty}\varrho_l(q,k)P_l'(\hat{\mathbf{q}}\cdot\hat{\mathbf{k}}), \tag{17a}$$

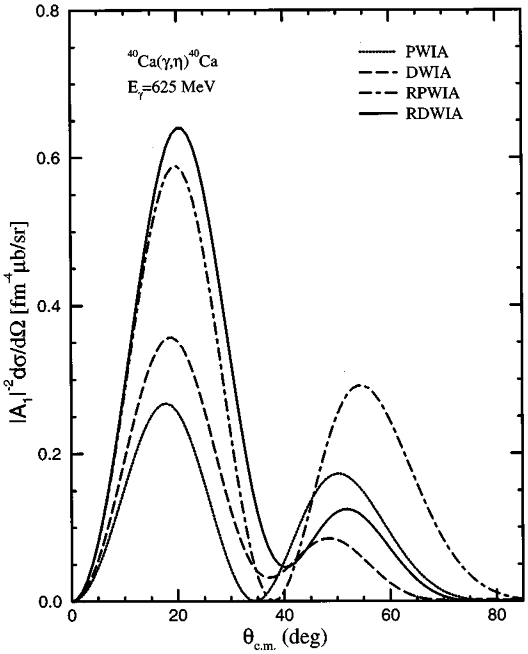

FIG. 2. The coherent η-photoproduction cross section from ^{40}Ca at a photon laboratory energy of $E_\gamma = 625$ MeV; note that the elementary amplitude A_1 has been factored out. The dotted (dashed) line represents a plane-wave (distorted-wave) calculation in which the lower components of the wave functions were determined from the free-space relation. The dash-dotted (solid) line represents a plane-wave (distorted-wave) calculation in which the lower components were generated dynamically from the Dirac equation.

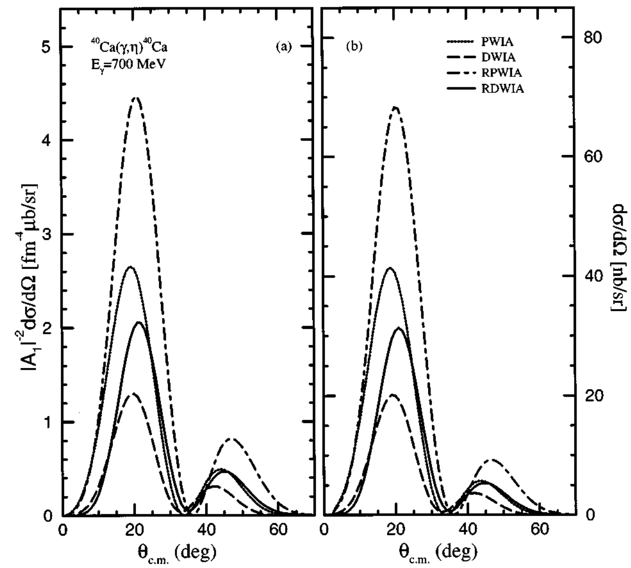

FIG. 3. The coherent η-photoproduction cross section from ^{40}Ca at a photon laboratory energy of $E_\gamma = 700$ MeV; with (a) and without (b) the elementary amplitude A_1 factored out. The dotted (dashed) line represents a plane-wave (distorted-wave) calculation in which the lower components of the wave functions were determined from the free-space relation. The dash-dotted (solid) line represents a plane-wave (distorted-wave) calculation in which the lower components were generated dynamically from the Dirac equation.

$$\varrho_l(q,k) = 4\pi \int_0^\infty dr r^2 \phi_{lq}^{(+)}(r) \rho_T(r) [j_{l-1}(kr) + j_{l+1}(kr)]. \tag{17b}$$

Note that $\phi_{lq}^{(+)}(r)$ is the η-meson distorted wave with angular momentum l and P_l' is the derivative of the Legendre polynomial of order l.

III. RESULTS

In Fig. 2 we display the coherent η-photoproduction cross section from ^{40}Ca at a photon laboratory energy of $E_\gamma = 625$ MeV. In order to identify those effects arising exclusively from relativity and final-state interactions we have factored out the elementary amplitude from the cross section. Thus, the only dynamical information that this calculation is sensitive to is nuclear structure and distortions. The dotted line represents a plane-wave-impulse-approximation (PWIA) calculation in which the lower components of the single-particle wave functions were determined from the free-space relation; the upper component remained unchanged, apart from a small normalization correction. This represents our best attempt at reproducing standard nonrelativistic calculations which employ free, on-shell spinors to effect the nonrelativistic reduction of the elementary amplitude. Indeed, in this ''nonrelativistic'' limit there is a simple relation between the tensor and vector densities of closed-shell nuclei:

$$\rho_T(Q) \approx -\frac{Q}{2M_N} \rho_V(Q), \tag{18}$$

where M_N is the free nucleon mass and $\rho_V(Q)$ is the Fourier transform of the ground-state matter density. In this way, the PWIA calculation establishes a baseline, against which possible relativistic effects may be inferred. The effect of distortions on this nonrelativisticlike calculation is depicted with the dashed line. At these low energies ($q_{\text{c.m.}} \lesssim 300$ MeV) the real part of the optical potential is attractive. This creates a competition with the imaginary part that results in a distorted-wave-impulse-approximation (DWIA) cross section relatively close to its plane-wave value. This picture, however, changes dramatically once the lower components of the wave functions are determined dynamically, rather than from the free-space relation. In the particular case of the Walecka model the lower components are enhanced substantially in the medium as a consequence of the large scalar and vector mean-field potentials. Since the coherent cross section becomes dominated by the tensor density—which is very sensitive to this enhancement—we obtain a relativistic-plane-wave-impulse-approximation (RPWIA) cross section that is, at least, twice as large as its nonrelativistic counterpart (dot-dashed line). Note that as in the nonrelativistic case, distortion effects do not seem to play an important role at this low energy (solid line).

In Fig. 3(a) we show the corresponding behavior of the coherent cross section at a photon energy of $E_\gamma = 700$ MeV. At these energies ($q_{\text{c.m.}} \gtrsim 400$ MeV) the real part of the optical potential has become repulsive. This, in combination with a relatively large imaginary component, results in a large quenching of the distorted-wave cross sections relative to their plane-wave values. Note that in spite of distortions—which make the interior of the nucleus, and thus the region of small effective nucleon mass, largely inaccessible—the

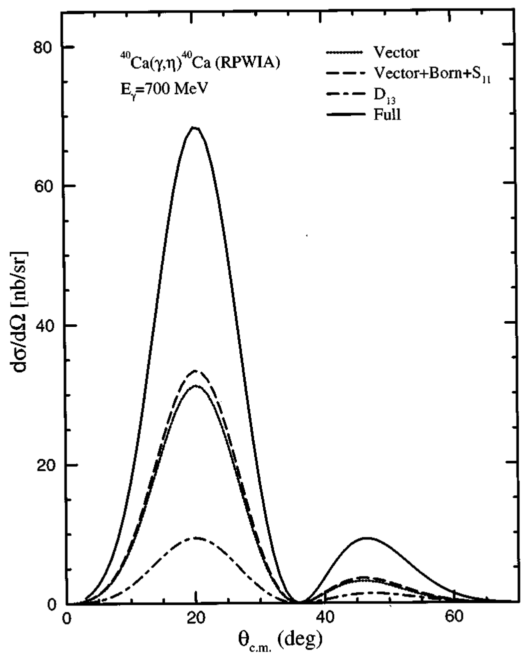

FIG. 4. Breakdown of the elementary contributions to the coherent η-photoproduction cross section from ^{40}Ca at a photon laboratory energy of $E_{\gamma}=700$ MeV. All curves were generated in a relativistic plane-wave-impulse approximation.

relativistic calculation (RDWIA) still shows a substantial enhancement relative to the nonrelativistic result (DWIA). We also display the coherent cross section—including the effect from the elementary amplitude—in Fig. 3(b). We conclude, because the elementary amplitude varies relatively slow in this region, that most of the qualitative features identified in Fig. 3(a) remain. Finally, in Fig. 4 we show a breakdown of the elementary contributions to the RPWIA calculation of Fig. 3(b). This figure shows that a significant portion of the strength arises from the individual contributions from the $D_{13}(1520)$ excitation and the t-channel exchange of vector mesons, while as expected very little strength is contributed by $S_{11}(1535)$ excitation or the Born terms. Also note that the constructive interference between $D_{13}(1520)$ excitation and vector-meson exchange results in a cross section substantially stronger than their incoherent sum.

IV. CONCLUSIONS

We have computed the coherent η-photoproduction cross section in a relativistic-impulse-approximation approach. Our formalism differs from conventional approaches in that we do not perform a nonrelativistic reduction of the elementary amplitude. In this manner, we could explore possible relativistic corrections to the coherent amplitude arising from medium modifications to the nucleon spinors. We have adopted a relativistic parametrization of the elementary amplitude in terms of four Lorentz- and gauge-invariant amplitudes. Using general arguments we have shown that the coherent reaction is sensitive to only one of the components of the elementary amplitude. Moreover, we identified the ground-state tensor density—and not the vector density—as the essential component of the nuclear structure.

The very high accuracy proton data [13] and the inferred neutron data from the deuteron target experiment [14] are well reproduced within the effective Lagrangian model which in turn is used to construct the elementary transition operator; this is the input to our nuclear calculations. The near cancellation of the sum of the proton and neutron $S_{11}(1535)$ helicity amplitudes allow us to examine the background contributions in great detail. In particular, our calculations indicate that the vector mesons and the $D_{13}(1520)$ account for most of the strength in the differential cross section. This can provide further theoretical insights into the background mechanism of η photoproduction. For the nuclear tensor density we used a mean-field approximation to the Walecka model. The Walecka model is characterized by large scalar and vector potentials that induce a large enhancement in the lower (''small'') components of the single-particle wave functions. The tensor density depends linearly on these lower components and is, thus, sensitive to their in-medium enhancement. Indeed, we reported relativistic cross sections that were substantially larger than to those obtained from assuming the free lower-to-upper ratio, as is usually done in nonrelativistic calculations. Further, we demonstrated that in the low-energy region, where the real part of the optical potential is attractive, these large enhancements are preserved even in the presence of distortions.

To conclude, we address two possible complications to the simple picture presented here: (1) violations to the impulse approximation and (2) off-shell ambiguities in the elementary amplitude. In the impulse approximation one assumes that the elementary amplitude—which only contains on-shell information—can be used without modification in the nuclear medium. However, any microscopic model of the reaction is bound to predict some sort of modification to the elementary amplitude as the process is embedded in the medium. For example, the coupling to intermediate N^* resonances represents an important component of most realistic models of the photoproduction amplitude. However, a variety of processes, such as the interaction of the resonances with the nuclear mean fields as well as Pauli blocking, can affect the formation, propagation, and decay of these resonances in the nuclear medium. Thus, it is important to have a reliable microscopic model in which to test the validity of the impulse approximation. A microscopic model can also provide guidance on how to take the elementary amplitude off shell. The form of the elementary amplitude used here, although standard, is not unique. Many other choices—all of them equivalent on shell—are possible. While all these choices are guaranteed to give identical results for on-shell observables, they can yield vastly different predictions off shell. Without theoretical guidance, there is no hope of resolving the off-shell ambiguity. Without experimental support, there is no way of testing the validity of our models. Indeed, much work remains to be done—on both theoretical and experimental fronts—before a clear picture of the coherent process can emerge.

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