Delta electroproduction in a chiral bag model approach

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We study the $\gamma^* N \to \Delta$ transition amplitudes in a recoil corrected cloudy bag model approach. A modified Peierls-Thouless projection method is used to construct the Galilean invariant baryon states. The pionic contribution is found to be significant. The effect of the recoil correction is to reduce the magnitude of the transition amplitudes at small momentum transfer and to enhance them at modest momentum transfers.

1 Introduction

The nucleon-delta electromagnetic transition amplitude is an outstanding example of the success of the quark model. There have been many theoretical and experimental explorations of this transition process. In a naive quark model the $\gamma^* N \to \Delta$ transition occurs only by an M1 transition, while the E2 process is fully suppressed. In more sophisticated models, quarks can interact through, for example, one-gluon exchange in addition to the confinement potential between them. Then it is possible for configuration mixing, involving the excitation of one quark to a $d$-state, to generate a small, but nonvanishing, E2 amplitude. To extract the $\gamma^* N \to \Delta$ amplitude from experimental data is not an easy task. There are some uncertainties in the subtraction of background, and the results are somewhat model dependent. With the advent of the new generation of accelerators, much more accurate measurements will be made. The anticipated high quality data should test various hadron models and help to build more realistic ones.

The cloudy bag model (CBM) improves the MIT bag model by introducing an elementary, perturbative pion field which couples to quarks in the bag in such a way that chiral symmetry is restored. The pion field significantly improves the predictions of the static properties of baryons. Previous calculations of delta photoproduction in the cloudy bag model differ from the results presented here and neglected the recoil correction. The baryon wave function is simply a direct product of individual quark wave functions, similar to the nuclear shell-model wave function (independent particle motion). This type of wave function is not a momentum eigenstate although the Hamiltonian commutes with the total momentum operator. The matrix elements evaluated between such static states contain spurious center of mass motion which ought to be removed. Early studies indicated that the correction for spurious center of mass motion is significant. It is expected to be most important in calculations...
where relatively large momentum transfers are involved. There are several intuitively motivated prescriptions\textsuperscript{8} for the correction of center of mass motion, however, none of them are fully satisfactory. In this work, we have chosen to use the Peierls-Thouless (PT)\textsuperscript{9} method to eliminate the center of mass motion, since it is the most convenient for our purposes.

In this paper, we calculate the nucleon-delta electromagnetic transition amplitudes with respect to the virtual photon. The spurious center of mass motion is corrected by using the PT projection method. As a first step, we assume exact $SU(6)$ symmetry for the quark structure of the baryons, so that all quarks in the ground state of the $N$ and $\Delta$ are in the $s$ state. The notation of references\textsuperscript{10,11} is followed. We briefly review the method to construct the PT wave function in Sec. II. The calculation of helicity amplitudes for virtual photoproduction of the delta is performed in Sec. III, and in Sec. IV we present the numerical results. Finally in Sec. V we summarize our results.

2 Galilean invariant baryon states

We start with the chirally invariant Lagrangian density of the cloudy bag model\textsuperscript{4}

$$L = (i\overline{q}\gamma^\mu \partial_\mu q - B)\theta_V - \frac{1}{2} \gamma q \Delta_S$$

$$+ \frac{1}{2}(\partial_\mu \pi)^2 - \frac{1}{2}m_\pi^2 \pi^2 - \frac{i}{2f} \gamma_5 \tau \cdot \pi q \Delta_S,$$

where $\theta_V$ is a step function which is one inside the bag volume $V$ and vanishes outside, and $\Delta_S$ is a surface delta function. In a perturbative treatment of the pion field, the quark wave function is not affected by the pion field and is given by the MIT bag solution\textsuperscript{5}

$$q(r) = \begin{pmatrix} g(r) \\ i\sigma \cdot \hat{r} f(r) \end{pmatrix} \theta(R - r),$$

where $R$ is the spherical bag radius. For the ground state of a massless quark $g(r) = N_s j_0(\omega_s r/R)$, $f(r) = N_s j_1(\omega_s r/R)$, where $\omega_s = 2.04$ and $N_s^2 = \omega_s/8\pi R^3 j_0^2(\omega_s)/(\omega_s - 1)$.

The bare baryon is taken to be composed of three quarks with the spin-flavor wave function given by $SU(6)$ symmetry. Naively the space component is the direct product of three quark wave functions in coordinate space

$$\Psi(x_1, x_2, x_3; x) = q(x_1 - x)q(x_2 - x)q(x_3 - x).$$

Here $x$ indicates the location of the bag center, while $x_1$, $x_2$, and $x_3$ specify the positions of the three quarks. Clearly this wavefunction does not have definite momentum and is not a momentum eigenstate. A momentum eigenstate of the baryon can be constructed by making a linear superposition of the localized states, namely,

$$\Psi_{PV}(x_1, x_2, x_3; p) = N'(p) \int d^3xe^{ip\cdot x}\Psi(x_1, x_2, x_3; x),$$
where the subscript PY stands for Peierls-Yoccoz projection, and \( N'(p) \) is a momentum-dependent normalization constant. It can be shown that \( \Psi_{\text{PY}}(p) = e^{ip\cdot x_{\text{cm}}} \Psi_{\text{in}}(p) \), where \( x_{\text{cm}} = (x_1 + x_2 + x_3)/3 \) is the center of mass of the baryon and \( \Psi_{\text{in}}(p) \) is the appropriately defined intrinsic part of the wave function. Since \( \Psi_{\text{in}}(p) \) still depends on the c.m. momentum, it violates translational invariance. To overcome this problem, Peierls and Thouless (PT) proposed to make another superposition of these momentum eigenstates, i.e.,

\[
\Psi_{\text{PT}}(x_1, x_2, x_3; p) = N(p) \int d^3p' w(p') e^{i(p-p') \cdot x_{\text{cm}}} \Psi_{\text{PY}}(x_1, x_2, x_3; p').
\]  

(5)

The weight function, \( w(p') \), should in fact be chosen to minimize the total energy, but this would be quite complicated to implement here. Instead, we choose \( w(p') = 1 \) for simplicity and convenience. Then integrations over \( x \) and \( p' \) can be carried out easily. This leads to a much simplified PT wave function,

\[
\Psi_{\text{PT}}(x_1, x_2, x_3; p) = N_{\text{PT}} e^{ip \cdot x_{\text{cm}}} q(x_1 - x_{\text{cm}}) q(x_2 - x_{\text{cm}}) q(x_3 - x_{\text{cm}}),
\]  

(6)

where the normalization factor, \( N_{\text{PT}} \), is given by the condition

\[
\int d^3x_1 d^3x_2 d^3x_3 \Psi_{\text{PT}}^\dagger(x_1, x_2, x_3; p') \Psi_{\text{PT}}(x_1, x_2, x_3; p) = (2\pi)^3 \delta^{(3)}(p' - p).
\]  

(7)

This leads to

\[
N_{\text{PT}} = \left[ 3 \int d^3r_1 d^3r_2 \rho(r_1) \rho(-r_1 + r_2) \right]^{-1/2},
\]  

(8)

where \( \rho(r) \equiv q^\dagger(r) q(r) = [g^2(r) + f^2(r)] \theta(R - r) \). Notice that, for this simple version of the PT projection, \( N_{\text{PT}} \) is a momentum independent constant and the wave function in Eq. (3) is manifestly Galilean invariant.

3 The helicity amplitudes in the cloudy bag model

From the CBM Lagrangian given in Eq. (3), the conserved local electromagnetic current can be derived using the principle of minimal coupling \( \partial_{\mu} \rightarrow \partial_{\mu} + iq A_{\mu} \), where \( q \) is the charge carried by the field upon which the derivative operator acts. The total electromagnetic current is then

\[
J^\mu(x) = j^\mu_q(x) + j^\mu_u(x),
\]  

(9)

\[
j^\mu_q(x) = \sum_a Q_a \overline{q}_a(x) \gamma^\mu q_a(x),
\]  

(10)

\[
j^\mu_u(x) = -ie [\pi^\dagger(x) \partial^\mu \pi(x) - \pi(x) \partial^\mu \pi^\dagger(x)],
\]  

(11)

where \( q_a(x) \) is the quark field operator for flavor \( a \), \( Q_a \) is its charge in units of \( e \), and \( e \equiv |e| \) is the magnitude of the electron charge. The charged pion field operator,
\( \pi(x) = \frac{1}{\sqrt{2}} [\pi_1(x) + i\pi_2(x)] \), either destroys a negatively charged pion or creates a positively charged one.

It is customary to define the helicity amplitudes for the electroproduction of the delta as

\[
A_{3/2} = \frac{1}{\sqrt{2\omega_\gamma}} \langle \Delta; s_\Delta = 3/2 | \vec{J}(0) \cdot \vec{e} | N; s_N = 1/2 \rangle, \tag{12}
\]

\[
A_{1/2} = \frac{1}{\sqrt{2\omega_\gamma}} \langle \Delta; s_\Delta = 1/2 | \vec{J}(0) \cdot \vec{e} | N; s_N = -1/2 \rangle, \tag{13}
\]

where the \( \Delta \) is at rest and the photon is travelling along the z-axis with right-handed polarization, \( \vec{e} = -\frac{1}{\sqrt{2}} (1, i, 0) \). The spin projections of \( \Delta \) and \( N \) along the z-axis are denoted as \( s_\Delta \) and \( s_N \) respectively. For a virtual photon, the three-momentum in the \( \Delta \) rest frame is given by

\[
|\vec{q}|^2 = Q^2 + \frac{(M_\Delta^2 - M_N^2 - Q^2)^2}{4M_\Delta^2}, \tag{14}
\]

with \( Q = \sqrt{-q^2} \) the magnitude of the four momentum transfer. The photon energy is related to this by \( q_0^2 = |\vec{q}|^2 - Q^2 \), where for a real photon we have \( Q^2 = 0 \), so that \( \omega_\gamma = |q_0| = (M_\Delta^2 - M_N^2)/2M_\Delta \). The experimentally extracted, resonant, helicity amplitudes are to be associated with the fully dressed initial and final baryons. In the cloudy bag model, due to the \( \piBB' \) coupling, a physical baryon state is described as a mixture of a bare bag and its surrounding pion cloud,

\[
|A\rangle = \sqrt{Z_2^A} [1 + (E_A - H_0 - \Lambda H_I \Lambda)^{-1} H_I] |A_0\rangle, \tag{15}
\]

where \( Z_2^A \) is the bare baryon probability in the physical baryon states, \( \Lambda \) is a projection operator which projects out all the components of \( |A\rangle \) with at least one pion, and \( H_I \) is the interaction Hamiltonian which describes the process of emission and absorption of pions. The matrix element of \( H_I \) between the bare baryon states is given by

\[
u_0^{AB}(\vec{k}) \equiv \langle A_0 | H_I | \pi_j(\vec{k})B_0 \rangle \]

\[
= \frac{if_0^{AB}}{m_\pi} \frac{u(kR)}{[2\omega_k(2\pi)^3]^{1/2}} \sum_{m,n} C_{S\mu_1S_A}^{\mu S_A,\mu S_A}(\hat{s}_m \cdot \vec{k}) C_{T\nu_1T_A}^{\nu T_A,\nu T_A}(\hat{t}_n \cdot \vec{e}_j), \tag{16}
\]

where the pion has momentum \( \vec{k} \) and isospin projection \( j \), \( f_0^{AB} \) is the reduced matrix element for the \( \pi B_0 \to A_0 \) transition vertex, \( u(kR) = 3j_1(kR)/kR \), \( \omega_k = \sqrt{k^2 + m_\pi^2} \), and \( \hat{s}_m \) and \( \hat{t}_n \) are spherical unit vectors for spin and isospin, respectively.

Under the approximation that there is at most one pion in the air, there are three different processes contributing to the \( \gamma^*N \to \Delta \) vertex, as shown in Fig. 1. For
Figure 1. Diagrams illustrating the various contributions included in the calculation. The intermediate baryons $B$ and $B'$ are restricted to the $N$ and $\Delta$ here.

Figs. 1(a) and 1(b), we substitute Eqs. (10) and (15) into Eqs. (12) and (13), and obtain the helicity amplitudes,

$$A_{3/2}^{(a)}(Q^2) = \sqrt{3} A_{1/2}^{(a)}(Q^2) = A_{\text{bare}}(Q^2) \sqrt{Z_N^2 Z_\Delta^2},$$

$$A_{3/2}^{(b)}(Q^2) = \sqrt{3} A_{1/2}^{(b)}(Q^2) = A_{\text{bare}}(Q^2) \left( \frac{f_{NN}^2}{27 \pi^2 m_\pi^2} \right) \int \frac{dk k u^2(kR)}{k} \left[ \frac{5/4}{\omega_k (\omega_k + \delta - \omega_\gamma)} + \frac{1}{(\omega_k + \delta)(\omega_k + \delta - \omega_\gamma)} + \frac{2/25}{(\omega_k + \delta)(\omega_k - \omega_\gamma)} + \frac{1}{\omega_k (\omega_k - \omega_\gamma)} \right],$$

where $\delta = m_\Delta - m_N$, and $f_{NN}$ is the renormalized $\pi NN$ coupling constant. The four terms in the right-hand side of Eq. (18) correspond to four possible intermediate states, ($N\Delta$), ($\Delta\Delta$), ($\Delta N$), and ($NN$), respectively. The recoil corrected bare $\gamma N_0 \to \Delta_0$ transition amplitude is

$$A_{\text{bare}}(Q^2) = -\frac{e}{\pi \sqrt{6} \omega_\gamma} \int_0^R dr r^2 \rho(r) f(r) j_1(qr) K(r)$$

where $K(r) = \int d^3x \rho(x) \rho(-x - r)$ is the recoil function to account for the correlation of the two spectator quarks. The renormalization constants, $Z_A^4$, are determined by the normalization condition for the physical baryon state. In this work, we have adopted the usual philosophy for the renormalization in the CBM. Throughout this
work approximate relation, $f^{AB} \approx \left(\frac{f^{AB}}{f_0^{NN}}\right) f^{NN}$, is always used. There are uncertain corrections on the bare coupling constant $f_0^{NN}$, such as the nonzero quark mass and correction of center of mass motion. Therefore, we use the renormalized coupling constant in our calculation, $f^{NN} \approx 3.03$, which correspond to the usual $\pi NN$ coupling constant, $f_\pi^{3NN} \approx 0.081$. As a result, the factor $\sqrt{Z^N Z^A}$ is absorbed into the renormalized coupling constants in Fig. 1(b). This treatment is equivalent to the original CBM formalism up to order $(f^{NN})^2$ and consistent with current conservation.

To evaluate the contribution caused by the photon-pion-pion coupling vertex [see Fig. 1(c)], we use the usual plane wave expansion for the quantized pion field

$$\pi_j(\vec{x}, t = 0) = \int \frac{d^3k}{(2\pi)^32\omega_k |k|^{1/2}} \left[a_j(\vec{k}) e^{ik \cdot \vec{x}} + a_j^\dagger(\vec{k}) e^{-ik \cdot \vec{x}}\right],$$

where $a_j(\vec{k})$ ($a_j^\dagger(\vec{k})$) annihilates (creates) a pion with momentum $\vec{k}$ and isospin $j$. With the identity, $a_j(\vec{k}) |A\rangle = (E_A - \omega_k - H)^{-1} H_i^j(\vec{k}, j) |A\rangle$, we obtain the transition amplitude at position $\vec{x}$,

$$\langle \Delta, s_\Delta | j_\pi(\vec{x}) | N, s_N \rangle = -ie \sum_{j'j} \epsilon_{jj'3} \int d^3k d^3k' e^{i\vec{k} \cdot \vec{x}} (2\pi)^3 2(\omega_k \omega_{k'})^{1/2} \times \sum_B \left[\eta_{jj'}^B(\vec{k}, \vec{k}) G^B(\vec{k}, \vec{k}) + \eta_{jj'}^B(\vec{k}, \vec{k}) G^B(\vec{k}, \vec{k})\right].$$

Here $B$ denotes the intermediate baryon states (restricted to $N$ and $\Delta$ here), and

$$\eta_{jj'}^B(\vec{k}, \vec{k}) \equiv \frac{f^B f^{NB}}{m_n^2} \frac{u(kR)u(k'R)}{16\pi^3(\omega_k + \omega_{k'})^{1/2}} \sum_B C_{S_B1S_{\Delta}S_N}^{s_Bm's_{\Delta}} C_{S_B1S_{\Delta}S_N}^{s_Bm's_{\Delta}} (\vec{s}_m \cdot \vec{k})(\vec{s}_m \cdot \vec{k})$$

$$\times \sum_t C_{\bar{T}_B1T_{\Delta}}^{tBn1\bar{n}_{\Delta}} C_{\bar{T}_B1T_{\Delta}}^{tBn1\bar{n}_{\Delta}} (\vec{t}_n \cdot \vec{e}_j)(\vec{t}_n \cdot \vec{e}_j),$$

$$G^N(\vec{k}, \vec{k}) \equiv \frac{1}{(\omega_k + \omega_{k'}) + \frac{1}{(\omega_k' - \omega_\gamma)(\omega_k + \omega_{k'} - \omega_\gamma)}} \left(\frac{1}{(\omega_k' + \omega_{\gamma})(\omega_k + \omega_{k'} - \omega_{\gamma})}\right)$$

$$G^\Delta(\vec{k}, \vec{k}) \equiv \frac{1}{(\omega_k + \omega_{k'}) + \frac{1}{(\omega_k' + \omega_{\gamma})(\omega_k + \omega_{k'} - \omega_{\gamma})}} \left(\frac{1}{(\omega_k' + \omega_{\gamma})(\omega_k + \omega_{k'} - \omega_{\gamma})}\right).$$

$G^\Delta(\vec{k}, \vec{k})$ and $G^\Delta(\vec{k}, \vec{k})$ are obtained by the interchange of $\vec{k}$ and $\vec{k}'$ in the corresponding equation. The three terms in Eqs. (23) and (24) indicate the three different time orders in the time-ordered perturbation theory, as illustrated in Fig. 1(c). Using the translational invariance of the electromagnetic current operator, $j^\mu(x) = e^{ip \cdot x} j^\mu(0)e^{-ip \cdot x}$, then the $\gamma^* N \rightarrow \Delta$ helicity amplitudes due to the $\gamma^* N$ interaction are simply given by

$$A(Q^2) = \int d^3x e^{i q \cdot x} \langle \Delta, s_\Delta | j_\pi(\vec{x}) \cdot \vec{e} | N, s_N \rangle.$$
Figure 2. The effect of the center of mass correction on the helicity amplitude, $A_{3/2}$, for the bare bag. The number on each curve indicates the bag radius in fm for the calculation.

After performing some spin and isospin algebra, we obtain

$$A_{3/2}^{(c)}(Q^2) = -\frac{(f^{NN})^2|\vec{q}|}{240\sqrt{6}\omega_{\gamma\pi^3}m^2_\pi} \int \frac{d^3k k^4 \sin^2 \theta u(kR)u(k'R)}{\omega_k\omega_{k'}} \times \left[ G^N(\vec{k}, \vec{k'}) + 3G^\Delta(\vec{k'}, \vec{k}) + 2G^\Delta(\vec{k}, \vec{k'}) \right],$$

$$A_{1/2}^{(c)}(Q^2) = -\frac{(f^{NN})^2|\vec{q}|}{720\sqrt{2}\omega_{\gamma\pi^3}m^2_\pi} \int \frac{d^3k k^4 \sin^2 \theta u(kR)u(k'R)}{\omega_k\omega_{k'}} \times \left[ 2G^N(\vec{k'}, \vec{k}) - G^N(\vec{k}, \vec{k'}) + G^\Delta(\vec{k'}, \vec{k}) + 4G^\Delta(\vec{k}, \vec{k'}) \right],$$

where $\vec{k'} = \vec{k} + \vec{q}$, $\omega_{k'} = \sqrt{k'^2 + m^2_\pi}$, and $\theta$ denotes the angle between $\vec{k}$ and $\vec{q}$. It is worthwhile to mention that the form of our results for Fig. 1(c) are quite different from those of KE$^6$ and Bermuth et al.$^7$ where the integral variables are $k$ and $k'$ in their formulations. We believe that our expressions are more straightforward and manifestly respect the three momentum conservation at the $\gamma\pi\pi$ vertex.

4 Results

The overall effect of the PT recoil correction on the bare bag contribution to the typical $\gamma^*N \rightarrow \Delta$ helicity amplitude, $A_{3/2}$, is shown in Fig. 2. In the real photon limit ($Q^2 \rightarrow 0$), the magnitude of the $\gamma^*N \rightarrow \Delta$ transition amplitude increases with the bag radii in a fashion similar to that of the magnetic moment of bare baryons. The
correction of the center of mass motion usually reduces the bare transition amplitudes by 5 to 10 % for $Q^2 \lesssim 0.5 \text{ GeV}^2$ within a reasonable range of bag radii. However, this recoil correction would flip sign and increase the transition amplitude for larger momentum transfers.

The real parts of total helicity amplitudes, $A = A^{(a)} + A^{(b)} + A^{(c)}$, are presented in Fig. 3. With the contributions of the pion cloud, the bag radius dependence is quite different from that for the bare transition amplitudes shown in Fig. 2. This can be explained by the fact that the pionic contribution is competing with that of quark core, since a small bag radius means a strong pion cloud. In the small $Q^2$ region, the pion cloud compensates more than the loss in the bare amplitude when using a small bag radius. Generally, the smaller the bag radius, the larger the total transition amplitude. We list the helicity amplitudes corresponding to the real photon limit at the $\Delta$ resonance in Table 1. With the small bag radius $R = 0.7 \text{ fm}$, we are able to reproduce the experimental helicity amplitude in this model.

5 Summary

In conclusion, we have calculated the $\gamma^* N \to \Delta$ transition form factors in the cloudy bag model, including the center of mass correction via a simplified Peierls-Thouless projection method. The effect of this recoil correction is to slightly reduce the magnetic form factor at small momentum transfer and to make the form factor slightly harder. Generally, with the PT projection the transition moment is reduced
Table 1. Helicity amplitude of delta photoproduction, $A_{3/2}^*$, in units of $10^{-3}\text{GeV}^{-1/2}$. Here static denotes the static calculation and PT denotes the Peierls-Thouless projection. The indices a, b, and c correspond to the Figs. 1(a), 1(b), and 1(c) respectively. The latest estimate by Particle Data Group is $-258 \pm 6$.

| R(fm) | static | PT |
|-------|--------|----|
|       | a      | b  | c  | total | a  | b  | c  | total |
| 1.0   | -115   | -53 | -21i | -205 | -106 | -49 | -18i | -195 |
| 0.9   | -106   | -62 | -20i | -216 | -98  | -58 | -18i | -205 |
| 0.8   | -96    | -73 | -19i | -233 | -86  | -68 | -17i | -222 |
| 0.7   | -85    | -87 | -17i | -260 | -79  | -80 | -16i | -249 |

by about $5 \sim 8\%$.

The pion cloud contribution proved to be crucial to account for the measured helicity amplitudes using a reasonable bag radius in this model. In similar calculations using constituent quark models (the nonrelativistic and relativized quark models), the helicity amplitudes are usually significantly underpredicted. Further details and extensions of this work will be presented elsewhere.

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