CHAPTER V
ANOMALY IN LOW-ENERGY ELECTRON-NUCLEON INELASTIC SCATTERING

(5.1) Introduction

The existence of weak neutral current couplings in the electron-nucleon system both in atomic physics and in electron-nucleon scattering processes as predicted by spontaneously broken gauge theories of Weinberg and Salam type has been studied in the recent past.\textsuperscript{101-113} The experimental results now available in atomic physics\textsuperscript{45} seem to provide positive evidence for the neutral current couplings even though first experimental results have been negative\textsuperscript{44,114}. The latest results from Novosibirsk\textsuperscript{45,115} are consistent with the predictions of the Weinberg-Salam model with $\sin^2 \theta_W = 0.25$ where $\theta_W$ is a parameter in the model called Weinberg angle.

Inelastic electron-nucleon scattering processes investigated experimentally\textsuperscript{116,117} as well as theoretically\textsuperscript{118}, occur predominantly via the parity conserving electromagnetic interaction and the parity violating non-electromagnetic interaction presumably weak interaction. Parity violation observed as the helicity dependence in the cross-section arises from interference between weak and electromagnetic amplitudes. The parity violating asymmetry between the right handed and left handed polarized electron-nucleon scattering has been theoretically analyzed mainly in parton model\textsuperscript{119-122}.

Parity violating asymmetry has been observed experimentally in the inelastic scattering of longitudinally polarized electron
beam from a unpolarized deuterium and hydrogen targets at 31A9. This process which was quite well understood phenomeno-
logically has a natural place to look for parity nonconservation in unified theories of weak and electromagnetic interaction
of Weinberg and Salam type model. The asymmetries have been
reported experimentally for electron energies between the range
of 16.2 GeV to 22.2 GeV. The experiment was performed by
Prescott et al.125 at SLAC at an electron energy of 19.4 GeV
and the experimental results are reported as124

\[ a = (-9.5 \pm 1.5) \times 10^{-6} \text{ e}^2/\text{GeV}^2 \text{ (deuterium)} \]
with \( \langle e^2 \rangle = 1.6 \text{ eV}^2 \) and the incoming electron energy
\( e = 19.4 \text{ GeV} \). The \( \langle e^2 \rangle \) is the average value of \( e^2 \). The
statistical and systematic errors are each about 9 percent of
the measured asymmetry.

The experimental results for asymmetry using the liquid
hydrogen (proton) target are reported as

\[ A = (-9.7 \pm 2.7) \times 10^{-3} \text{ e}^2/\text{GeV}^2 \text{ (proton target)} \]
with \( \langle e^2 \rangle \), the average value of \( e^2 \) for the reaction as
\( \langle e^2 \rangle = 1.46 \text{ eV}^2 \).

The theoretically calculated polarized asymmetry is deep
inelastic electron-deuteron scattering is in close agreement
with the experimentally observed values at SLAC and is consistent
with the Weinberg-Salam theory for Weinberg angle \( \sin^2 \theta_W = 0.20 \). Thus the inelastic \( e - d \) scattering experiment seems to give
good support for Weinberg-Salam model.
Experiments are presently underway to look for various 
parity violating effects in the electron-nucleon scattering 
at intermediate energies. Some theoretical calcula-
tions have also been made to estimate these effects. In the 
case of elastic electron-deuteron scattering at intermediate 
energies, calculations have been made for various parity 
violating effects. It has been shown in elastic 
electron-deuteron scattering that the parity violating neutral 
currents give rise to the vector polarization of the recoil 
deuteron and is of the same order as the other parity violating 
effects in these systems. These studies have been here exten-
ded to explain the asymmetry observed in deep inelastic electron-
deuteron scattering at intermediate energies.

In section (V-2.1) we evaluate the matrix element for 
the processes \( e^+ + d \rightarrow e^+ + n + p \) and discuss the weak and 
electromagnetic form factors for nucleons used to derive the 
expressions. Section (V-3) deals with the calculation of 
any asymmetry in the closure approximation for the outgoing 
nucleons. The results are applicable in the intermediate 
energy region of the incident electron where asymmetries are 
found to be in the range accessible to the present experiments.

(V-2) Inelastic electron-deuteron scattering

(V-2.1) Matrix element

In this section we shall derive the matrix element for 
the process,

\[
\tag{V.1}
\end{equation}
where \((\vec{k}_1, \vec{p}_1)\) are the initial electron energy and momentum
and \((\vec{k}_2, \vec{p}_2'), (\vec{q}_1, \vec{q}_1), \text{ and } (\vec{q}_2, \vec{p}_2)\) are the momentum and
energy for electron, proton and neutron in the final state.

The matrix element for the process \(e^\pm + d \rightarrow e^\mp + n + p\)
is denoted as

\[
M^\pm = i S_{\lambda\gamma}^{EM} \gamma^\gamma + i S_{\lambda\gamma}^{W} \gamma^\gamma
\]

where \(\gamma\)'s in the above equation (V.2) refer to the particular helicity state of the electron i.e., for the left handed
and right handed polarized electrons respectively. \(S_{\lambda\gamma}^{EM}\) and
\(S_{\lambda\gamma}^{W}\) denote the leptonic and hadronic matrix elements for
electromagnetic interaction and \(S_{\lambda\gamma}^{W}\) denotes the leptonic
matrix element and hadronic matrix element for weak neutral
current interaction.

The leptonic matrix elements \(S_{\lambda\gamma}^{EM}\) and \(S_{\lambda\gamma}^{W}\) are given by

\[
S_{\lambda\gamma}^{EM} = \bar{u}(k') \gamma^\gamma \frac{(1 + \gamma_5)}{2} u(k)
\]

and

\[
S_{\lambda\gamma}^{W} = \bar{u}(k') \gamma^\gamma i \gamma^\gamma \gamma^\gamma \frac{(1 + \gamma_5)}{2} u(k)
\]

where \(u, \bar{u}\) denote the spinors of the electrons, \(k\) and \(k'\)
are the initial and final leptons (electrons) momenta measured
in the deuteron rest frame. The factor \(\frac{(1 + \gamma_5)}{2}\) projects out the
particular helicity for left handed and right handed polarized
electrons which we assume to be moving with relativistic
velocities.

The hadronic matrix elements for electromagnetic and
weak interaction processes, \(S_{\lambda\gamma}^{EM}\) and \(S_{\lambda\gamma}^{W}\) in the impulse
approximation are given by
\[ J_{\lambda}^{\mu} = -\frac{e^2}{q^2} \int \phi_2^*(\vec{r}) \left[ \mathcal{F}_{\lambda}^b e^{i\vec{q}\cdot\vec{r}/2} + \mathcal{F}_{\lambda}^n e^{-i\vec{q}\cdot\vec{r}/2} \right] \phi_1(\vec{r}) d^3r \]  
(7.3)
and
\[ J_{\lambda}^{\nu} = -\frac{G}{q^2} \int \phi_2^*(\vec{r}) \left[ \mathcal{\Lambda}_{\lambda}^b e^{i\vec{q}\cdot\vec{r}/2} + \mathcal{\Lambda}_{\lambda}^n e^{-i\vec{q}\cdot\vec{r}/2} \right] \phi_1(\vec{r}) d^3r \]  
(7.6)
where \( e^2 = 4\pi\alpha \) is substituted in equation (7.5) for the electromagnetic coupling constant with \( \alpha = 1/137 \) as the fine structure constant. \( \vec{q} = \vec{k} - \vec{k}' \) is the three-momentum transfer for the process. \( \phi_1(\vec{r}) \) and \( \phi_2(\vec{r}) \) are the initial deuteron wave function and final dimuon wave function respectively.

Since we are interested to calculate the hadronic matrix elements between the initial deuteron state and final dimuon state we shall take the nucleons in the final state as non-relativistic. \( \mathcal{F}_{\lambda}^N \) and \( \mathcal{\Lambda}_{\lambda}^N (N = p, n) \), both for proton as well as neutron, are the non-relativistic reductions of the single nucleon covariant operators \( T_{\lambda}^{YN} \) and \( T_{\lambda}^{WN} \) for \( \gamma \)-exchange (electromagnetic) process and for \( W \)-exchange (weak interaction) process involving neutral currents. These single nucleon operators, \( T_{\lambda}^{YN} \) and \( T_{\lambda}^{WN} \), are defined as
\[ T_{\lambda}^{YN} = \left[ G_{1}^N(q^2) Y_{\lambda} + 2 \left( \frac{\not{\alpha}_1}{2M} \right) G_{2}^N(q^2) \right] \]  
(7.7)
and
\[ T_{\lambda}^{WN} = \left[ F_{1}^N(q^2) Y_{\lambda} + 2 \left( \frac{\not{\alpha}_1}{2M} \right) F_{2}^N(q^2) - G_{1}^N(q^2) Y_{\lambda} Y_{5} - \frac{\alpha N(q^2)}{\alpha A(q^2)} \right] \]  
(7.8)
where $G_1^N(q^2)$ and $G_2^N(q^2)$ in equation (V.7) are the vector form factors for electromagnetic processes having the four momentum transferred square i.e. $q^2$, dependence and defined for proton and neutron separately. In equation (V.8), $F_1^N(q^2)$ and $F_2^N(q^2)$ are the vector form factors for nucleons in weak neutral current interaction, $g_A^N(q^2)$ is the axial vector form factor for nucleons as defined for proton and neutron separately.

Following the standard method as discussed in chapter I, we obtained the nonrelativistic reductions of single nucleon operators between the initial deuteron state and final dimuon state i.e. $\langle np|\Gamma^{^1S_0}_N|d\rangle$ and $\langle np|\Gamma^{^3P_1}_N|d\rangle$. We get the following results for $\mathcal{F}_\lambda^N$ and $\mathcal{A}_\lambda^N$.

$$\mathcal{F}_\lambda^N = G_3^N(q^2) \frac{P^\lambda}{M} + 1 G_M^N(q^2) \varepsilon_{ijk} \frac{q_i q_k}{2M^2} \tag{V.9a}$$

$$\mathcal{A}_\lambda^N = G_3^N(q^2) \left(1 + \frac{P^\lambda}{2M^2}\right) + 1 G_M^N(q^2) \frac{q_i P^\lambda}{4M^2} \tag{V.9b}$$

where

$$G_3^N(q^2) = G_1^N(q^2) - \frac{2M^2}{M^2} G_2^N(q^2)$$

$$G_3^N(q^2) = G_1^N(q^2) + G_2^N(q^2) + 1 G_M^N(q^2)$$

$$G_3^N(q^2) = G_1^N(q^2) + 2G_2^N(q^2)$$

The results for $\mathcal{A}_i^N$ and $\mathcal{A}_o^N$ are same as defined in equation (I.20) and quoted as

$$\mathcal{A}_i^N = \left[ F_i^N(q^2) \frac{P_i^i}{M} + 2 F_m^N(q^2) \varepsilon_{ijk} \frac{q_i q_k}{2M} - 9A(q^2) \left(1 + \frac{q^2}{8M^2}\right) \sigma_i^p \right.\left. + \frac{h^2}{2M^2} \sigma_i - \frac{q_i q_j}{8M^2} \sigma_j^p \sigma_j \right] \tag{V.11a}$$
\[ \chi^N_0 = g^{-1} \left( \frac{q^2}{2M^2} \right) - g^N_A(q^2) \frac{\sigma^2 P^2}{M} + i \varepsilon_{ijk} g^N_R(q^2) \frac{g^j P^k \sigma^l}{4M^2} \] (V.11b)

where \( g^N_A(q^2) \), \( g^N_D(q^2) \) and \( g^N_R(q^2) \) are some as given in equation (I.21). \( g^N_A(q^2) \) and \( g^N_R(q^2) \) are the usual form and magnetic form factors for nucleons. It should be noted here that in deriving the equations (V.11) for the non-relativistic reduction \( \chi^N_0 \), we have neglected the term containing the pseudoscalar form factor which is proportional to electron mass.

(V-3) Calculation of Asymmetry

The asymmetry in the polarized electron-deuteron inelastic scattering is defined as

\[ A = \left( \frac{d\sigma^-}{d\Omega} \right)_+ - \left( \frac{d\sigma^-}{d\Omega} \right)_- \] (V.12)

where \( \left( \frac{d\sigma^-}{d\Omega} \right)_+ \) and \( \left( \frac{d\sigma^-}{d\Omega} \right)_- \) are the differential cross-sections for the inelastic scattering of left handed and right handed electrons on unpolarized deuteron targets. That is, we separately calculate the cross-sections for the scattering of left handed and right handed electrons on an unpolarized deuteron target, then the asymmetry \( A \) is obtained in equation (V.12) by dividing the difference by the sum of these cross-sections. The asymmetry is parity violating and, therefore,
it is a measure to determine that the electron is the parity violating weak neutral interaction with the deuteron target. Since the asymmetry arises due to interference between weak and electromagnetic amplitudes, thus it is proportional to a factor $G^2$, where $G$ is usual fermi coupling constant and $\alpha$ as the fine structure constant, which is larger than the usual weak interaction effects which are of the order of $\alpha^2$. Since asymmetry is ratio, many theoretical as well as experimental uncertainties common to both numerator and denominator in equation (V.12) cancel. Secondly, it is useful to predict the relative sign between weak and electromagnetic interactions. The determination of the algebraic sign between weak and electromagnetic interference term is the unique prediction of unified gauge theory of weak and electromagnetic interactions. Thus the calculation of asymmetry provides an additional check on the various neutral current $SU(2) \times U(1)$ type models.

Using the matrix element defined through equations (V.2)-(V.7), the differential cross sections $\left(\frac{d\sigma}{d\omega}\right)_\pm$ are written as

$$\left(\frac{d\sigma}{d\omega}\right)_\pm = \frac{G^2}{192\pi^5} \frac{m_e m_e' e'}{E_e} \int \left|M^{\pm}_2\right|^2 d^3p',$$

(V.13)

where $\left|M^{\pm}_2\right|^2$ is the absolute square of the matrix element defined in equation (V.2), $p'$ is relative momentum of the two nucleons in their centre of mass frame. $E_e, m_e$ are the initial
electron energy and mass and \( E_e \), \( m_e \) are the final electron energy and mass respectively.

The matrix element square, \( |m|^{\pm}|^{2} \) is defined as

\[
|m|^{\pm}|^{2} = \sum_{\nu j} L_{\lambda j}^{\pm \lambda} \cdot S_{\nu j}^{\pm}
\]  

(7.14)

where

\[
L_{\lambda j}^{\pm \lambda} = \sum_{\nu j} \sum_{\text{spins}} \left( \ell_{\lambda j}^{\pm \lambda} \right)^{\dagger} \left( \ell_{j}^{\pm \lambda} \right)
\]  

(7.15)

and

\[
S_{\nu j}^{\pm} = \sum_{\text{spins}} \sum_{\nu j} \left( J_{\lambda}^{\dagger} \right)^{\dagger} \left( J_{j}^{\dagger} \right)
\]  

(7.16)

\( i \) and \( j \) take values 1 and 2. \( \ell_{\lambda j}^{\pm \lambda} \) and \( J_{\lambda}^{\dagger} \) refer to the leptonic matrix elements corresponding to \( y_{\ldots} \) (electromagnetic) and \( \nu_{\ldots} \) (weak neutral current) couplings defined in equations (7.3) and (7.4). \( J_{\lambda}^{\dagger} \) and \( J_{\lambda}^{\dagger} \) refer to the hadronic matrix elements corresponding to \( Y_{\ldots} ^{\ldots} \) and \( \bar{\nu}_{\ldots} ^{\ldots} \) couplings defined in equations (7.5) and (7.6). The integration in equation (7.15) is performed over \( \nu j \) and the differential cross sections \( \left( d\sigma/d\omega \right) \) can be calculated.

We calculate the asymmetry in equation (7.12) by performing the \( d\phi \) integration in equation (7.13) in the closure approximation following the standard method used in chapter IX. This involves replacing the leptonic factor \( L_{\lambda j}^{\pm \lambda} \) by their average value as \( \langle L_{\lambda j}^{\pm \lambda} \rangle \) which is then assumed to be a slowly varying function over the range of hadronic variables integrated.
and can be taken outside the integration. Care must be taken to choose \( \langle \mathcal{L}_{xy} \rangle \). \( \langle \mathcal{L}_{xy} \rangle \) is chosen by assuming

\[ \langle \hat{E}^2 \rangle = \frac{q^2}{2m} \]

corresponding to the quasi elastic peak in the double differential cross section \( d^2\sigma/d\Omega dE \). The \( d^3p \) integration is then performed using the closure approximation over the final dimuon wave function as given in equation (II.9) where \( \psi_\alpha(x) \) is the same final dimuon wave function in this process.

We get the following results for the various components of the hadronic tensors \( J_{\alpha\beta} \), i.e.,

\[
J_{12}^{12} = \int \int \frac{d^3p}{c^3} \left[ \left( \frac{\partial^2}{\partial M^2} - \frac{1}{m} \right) \left( \frac{1}{1 + \frac{q^2}{2m}} \right) \right] (V.17a)
\]

\[
J_{13}^{13} = \int \int \frac{d^3p}{c^3} \left[ \left( \frac{\partial^4}{\partial M^4} - \frac{1}{m} \right) \left( \frac{1}{1 + \frac{q^2}{2m}} \right) \right] \frac{M}{p} (V.17b)
\]

\[
J_{13}^{13} = \int \int \frac{d^3p}{c^3} \left[ \left( \frac{\partial^2}{\partial M^2} - \frac{1}{m} \right) \left( \frac{1}{1 + \frac{q^2}{2m}} \right) \right] \frac{M}{p} (V.17c)
\]
\[
\begin{aligned}
\mathbf{J}_{00}^{22} &= \int_{-\infty}^{\infty} d^2 \mathbf{p}' \cdot 24\pi^2 \left[ \left( \rho_R^n(a^2) + \rho_R^0(a^2) \right) \left( 1 + \frac{a^2}{2m_0} - \frac{M_a}{2m_0} \right) \\
&\quad + 2 \rho_R^0(a^2) \rho_R^0(a^2) \left\{ \frac{1}{m_0} \left( 1 + \frac{a^2}{2m_0} \right) - \frac{\eta(a)}{m_0} \right\} \\
&\quad + \left( \rho_R^0(a^2) + \rho_R^0(a^2) \right) \left( \frac{a^2}{2m_0} - \frac{M_a}{2m_0} \right) + 2 \rho_R^0(a^2) \rho_R^0(a^2) \left( \frac{a^2}{2m_0} - \frac{M_a}{2m_0} \right) \left( m_0 + \frac{M_a}{2m_0} \right) \right] \\
\mathbf{J}_{\alpha 1}^{22} &= \int_{-\infty}^{\infty} d^2 \mathbf{p}' \cdot 24\pi^2 \left[ \left( \rho_R^n(a^2) + \rho_R^0(a^2) \right) \left( 1 + \frac{a^2}{2m_0} - \frac{M_a}{2m_0} \right) \right] \\
&\quad + \frac{\eta(a)}{m_0} \left( \rho_R^0(a^2) + \rho_R^0(a^2) \right) \rho_R^0(a^2) \left( \frac{a^2}{2m_0} - \frac{M_a}{2m_0} \right) \left( m_0 + \frac{M_a}{2m_0} \right) \right] \\
\mathbf{J}_{1j}^{22} &= \int_{-\infty}^{\infty} d^2 \mathbf{p}' \cdot 24\pi^2 \left[ \left( \rho_R^n(a^2) + \rho_R^0(a^2) \right) \left( 1 + \frac{a^2}{2m_0} - \frac{M_a}{2m_0} \right) \right] \\
&\quad + \frac{\eta(a)}{m_0} \left( \rho_R^0(a^2) + \rho_R^0(a^2) \right) \rho_R^0(a^2) \left( \frac{a^2}{2m_0} - \frac{M_a}{2m_0} \right) \left( m_0 + \frac{M_a}{2m_0} \right) \right] \\
\mathbf{J}_{\lambda \eta}^{22} &= \mathbf{J}_{\lambda \eta}^{22} \left( \rho_R^n(a^2) = 0 \right) \left( \rho_R^0(a^2) \rightarrow \rho_R^0(a^2) \right) \left( \rho_R^0(a^2) \rightarrow \rho_R^0(a^2) \right) \\
\end{aligned}
\]
The functions $I(q)$, $II(q)$ and $K$ are defined by the expressions (11.12). These functions contain the lepton wave function for initial state $\psi_1(x)$ and thus the lepton effect enters through these functions.

Using these equations, the asymmetry is calculated to be lowest order in weak interaction coupling constant to be

$$A(q) = \frac{\alpha^2}{\beta_2 - \alpha^2} \frac{X}{Y} \quad (\text{11.20})$$

where

$$X = \left\{ -a(1-\cos q) \frac{(n^2 + l^2)}{l} \right\} (\frac{\alpha^2}{\beta_2 - \alpha^2} + \frac{\alpha^2}{\beta_2 - \alpha^2} - \frac{a^2}{\beta_2 - \alpha^2})$$

$$+ \alpha^2 (\frac{\alpha^2}{\beta_2 - \alpha^2} - \frac{a^2}{\beta_2 - \alpha^2}) \frac{I(q)}{l} + a \left\{ \frac{a^2}{\beta_2 - \alpha^2} \right\} \frac{I(q)}{l} \right\}$$

$$\cdot \frac{\alpha^2}{\beta_2 - \alpha^2} + \frac{\alpha^2}{\beta_2 - \alpha^2} \frac{I(q)}{l} \frac{\alpha^2}{\beta_2 - \alpha^2} \frac{I(q)}{l} \frac{\alpha^2}{\beta_2 - \alpha^2}$$

and

$$Y = \frac{\alpha^2}{\beta_2 - \alpha^2} \left\{ 1 + \frac{(n^2 + l^2) \frac{I(q)}{l} (1+\cos q)}{\beta_2 - \alpha^2} \right\} (1+\cos q)$$

$$- \frac{\alpha^2}{\beta_2 - \alpha^2} \frac{I(q)}{l} (1+\cos q) - \frac{\alpha^2}{\beta_2 - \alpha^2} \frac{I(q)}{l} (1+\cos q)$$

$$\cdot \frac{\alpha^2}{\beta_2 - \alpha^2} \frac{I(q)}{l} (1+\cos q) \quad (\text{11.21})$$
(4.4) Results and Discussion

In order to calculate the asymmetry from equations (4.20), (4.21) and (4.22) we need a model to describe the neutral weak interaction Hamiltonian which provides the value for various weak and electromagnetic form factors defined in equations (1.91) and (1.10). We have used various neutral current models based on C(2) algebra to calculate the weak form factors for the neutrino neutron disintegration process in chapter 11. In the following we give a discussion of a model similar to one used in chapter 11 and describe the form factors and wave function in calculating the asymmetry for the process \( e^\pm + n \rightarrow e^\pm + n + p \).

(a) Weak interaction model

The interaction Lagrangian for the inelastic e-n scattering is written in the form

\[
\mathcal{L}_{\text{int}} = \frac{1}{\sqrt{2}} e \bar{\psi}_e \gamma_\mu (\not{p} - b \not{d}_S) \not{Q}_e \gamma_\mu \not{S}_e
\]

(4.23)

where \( J_{\not{Q}_e}^{\not{S}_e} \) is given in quark model by

\[
J_{\not{Q}_e}^{\not{S}_e} = \frac{1}{2} \left[ \bar{d}_A (\gamma_\mu - \gamma_\mu \not{Q}_e) \not{S}_e + \bar{d}_A (\not{S}_e \gamma_\mu - \not{S}_e \gamma_\mu \not{Q}_e) \not{Q}_e \right]
\]

(4.24)

\( \chi_\not{Q}_e, \chi_\not{S}_e \) are the amplitudes for isoscalar and isovector vector currents while \( \chi_{A} (\not{Q} = 0, 1) \) are the amplitudes for axial vector current. The values of \( \chi_{\not{Q}} \) and \( \chi_A \) used in the calculation of form factors are evaluated from equation (1.10) as given in table I. In this model the form factors used for weak and electromagnetic interactions are given by
\[ F^p_1(c) = \frac{1}{2} (3a_y^c + a_1^c), \quad F^p_1(0) = \frac{1}{2} (3a_y^u - a_1^u) \]

\[ F^p_2(0) = \frac{1}{2} \left[ F^p_1(3a_y^u - a_1^u) + \mu_n (3a_y^c - a_1^c) \right] \]

\[ F^n_2(0) = \frac{1}{2} \left[ F^p_1(3a_y^c - a_1^c) + \mu_n (3a_y^u + a_1^u) \right] \]

\[ F^n_A(0) = \frac{1}{2} \left( \frac{3}{2} a_h^c + a_1^c \right) \]

\[ F^n_A(0) = \frac{1}{2} \left( \frac{3}{2} a_h^u - a_1^u \right) \]

\[ a' = 1 - 4 \ln^2 \eta_w, \quad b' = 1 \]  

The numerical values for weak form factors in various neutral current models calculated from equation (V.25) are given in Table II. The electromagnetic form factors are evaluated as

\[ G^p_1(c) = 1, \quad G^p_n(c) = 0, \quad G^n_1(c) = 1 + \mu_p, \quad G^n_2(c) = \mu_n \]  

with

\[ \left( \frac{\mu}{\mu^p} \right) = 1.24; \quad \frac{\mu_n}{\mu^p} = 1.79 \text{ and } \frac{\mu}{\mu_n} = -1.99 \]

The \( q^2 \) dependence of the various weak and electromagnetic form factors is taken to be same i.e.,

\[ g(q^2) = \frac{g(c)}{(1 + c^2/m_n^2)^2}, \quad v(q^2) = \frac{v(c)}{(1 + c^2/m_n^2)^2} \]

with \( \sqrt{v} = 840 \text{ MeV} \).

(b) Wave function

The effect of neutron wave functions enters through the functions I(q), II(q) and K defined in equation (VII.12). We have used here the Hulthen wave function for the neutron.
(neglecting the n state) to calculate these functions. These functions evaluated from equation (II.12) are given in equation (II.23).

Using the form factors from equations (7.25) and (7.26) and the functions H(n), H(q) and k from equation (II.23), the asymmetry A(θ) is numerically calculated from equation (7.20) for various values of energies relevant to the presently proposed experiments in the intermediate-energy range. In figure 11 we have shown the asymmetry A(θ) versus scattering angle θ for various energies. We see that for a given energy the asymmetry is higher at backward angles and can be large enough to be experimentally accessible. The asymmetry increases with the incident electron energy. The maximum asymmetry which is expected at backward angles (high q²) however does not increase very much due to the damping induced by the weak and electromagnetic form factors at high q². For comparison we have also shown in figure 12, the asymmetry A(θ) versus the scattering angle θ for various weak interaction models discussed in text.

In the end we would like to make some comments about the closure approximation (see chapter II) in which these results have been calculated. This approximation gives fairly good results in the energy and q² range considered here. The major uncertainty in closure approximation comes from (a) extending the d³p from p_max to ∞ where p_max is the momentum allowed by the energy momentum conservation and (b) in making the proper choice of <q>. It has been shown elsewhere in the related
processes $\gamma + d \to \gamma + p + p$ and in $\gamma + d \to \gamma + n + p$ (chapter 11) that the additional dinucleon states which are included by extending the $d^3p$ integration from $p_{\text{max}}$ to $\infty$ contribute very little to the differential cross section. We therefore do not make much error in extending the momentum integration from $p_{\text{max}}$ to $\infty$ for $<e'\gamma>$. We have taken the energy corresponding to quasi-elastic peak, i.e., $<e'\gamma> = E_e - q^2/2M$. This may not be a very good choice for $<e'\gamma>$ at very small $q^2 \approx 0$ where the internal motion of the nucleons inside the deuteron becomes important. But even at small $q^2$ this choice for $<e'\gamma>$ is not bad at intermediate energies where $E_e >> q^2/2M$. It is thus clear that except for the low $q^2$, where $E_e$ is small, the closure approximation should be quite good. The effect of final state interactions in $q^2 + dq^2 + n + p$ can also shift the quasi-elastic peak thus changing $<e'\gamma>$ but it has been shown that the effect of final state interactions in $d\gamma/dq^2$ and $\gamma$ are very small even though they may be large for $d^2\gamma/dq^2dq'$. 

We can therefore conclude that the results presented in the closure approximation should be quite good at intermediate energies and in the backward direction where $q^2$ is relatively large. It is in this region, that the asymmetries become large enough to be experimentally observable.