CHAPTER 1
HEAVY-NUCLEUS DISINTEGRATION PROCESSES

The existence of neutral current in strangeness conserving ($\Delta S = 0$) processes has now been established experimentally in high energy neutron reactions $^{23-30}$, but its space-time and isospin structure is still a subject of considerable theoretical study. Until now most of the theoretical analysis of the experimental results has been done assuming $V-A$ structure of neutral current $^{27-31}$ as predicted by spontaneously broken gauge theories. Recently, the question of the outgoing neutrino identity and its properties while interesting through neutral current interactions has been raised $^{32-34}$ and the need for studying the general space-time structure of this interaction has been emphasised $^{35}$. In particular, the possible presence of helicity flipping $S, P, T$ couplings in neutral current interaction has been studied in leptonically and semileptonic processes occurring in various branches of physics and attempts are made to analyse various experiments in these interactions $^{36-42}$. These theories are not ruled out from the data in particle physics processes and have acquired a new importance in the light of recent absence of parity violation in atomic physics experiments predicted by $V-A$ theories $^{43-45}$. It has been stressed by many people $^{46}$ that the nuclear physics processes can play a decisive role in elucidating the space-time and isospin structure of
these fundamental interactions. It is in this connection that much emphasis has been recently put on the process $\nu(\beta) + d \to \nu(\beta) + n + p$ both theoretically and experimentally. This process is presently being studied at "Savannah River Plant Fission Reactor neutrino facility" by Burr et al.\textsuperscript{24} at low energies with electron type of neutrino ($\nu_e$) and would soon be studied at LAMPF and ARL at intermediate energies\textsuperscript{49-54}. An older analysis\textsuperscript{55} of this process using the inverse antineutrino spectrum had pushed down the earlier results for $\frac{\sigma_{\text{expt.}}}{\sigma_{\text{theory}}} \approx 6$ to $\frac{\sigma_{\text{expt.}}}{\sigma_{\text{theory}}} \approx 3.6$. The improved experiment by vanderB et al.\textsuperscript{56} has observed this process and found the total cross-section to be

$$\sigma_{\text{expt.}} = (3.8 \pm 0.9) \times 10^{-45} \text{ cm}^2/\nu. $$

The experiments with $\bar{\nu}_e$ ($\nu_{\bar{e}}$) when analysed should also see this event at some level specially when the existence of the neutral current interaction is so well established in $\nu_e$ interactions. The reactor energies (45 MeV) are too small to include higher partial wave disintegrations of the deuteron but at the LAMPF and ARL energies there is considerable disintegration into P wave and higher partial waves. As shown in the case of charged reactions the P wave and higher partial waves become more important at relatively low incident neutrino energies\textsuperscript{57,58}. Keeping these experiments in mind here we develop the theoretical formulations in $V$, $A$ and $P$, $T$ theories. The wave functions used for initial deuteron state and final dinucleon states are described in section (1-3),
and the relative four-momentum squared of the final nucleons is calculated to be

\[ r^2 = r_0^2 = -\frac{q^2}{4} - y (p_y - p'_y - n) \] (1.6)

where \( y \) is the binding energy of the neutron and \( n \) is the nucleon mass. If the relative four momentum of the outgoing nucleons be \( r_0^2 \) in the centre of mass frame then \( n^2 = n_0^2 = -\frac{q^2}{4} \) and we get

\[ r^2 = \frac{q^2}{4} + y (p_y - p'_y - n) \] (1.6)

(c) The threshold energy of the final neutrino is given by the condition \( n^2 = 0 \) where the nucleons come out with zero kinetic energy. In this case the neutrino has the maximum energy available to it for a given scattering angle \( \theta \).

The threshold energy is calculated to be given by

\[ E_{\nu} = n \cos \theta + \sqrt{(2m - E_{\nu} \cos \theta)^2 + \nu^2 (4m^2 - E_{\nu}))} - E_{\nu} \] (1.7)

which gives the limiting case i.e., \( n \to 0 \)

\[ \frac{\nu}{E_{\nu}} = \frac{p_{\nu}}{E_{\nu}} \left[ 1 + \left(\frac{\nu}{E_{\nu}}\right) \sin \theta / 2 \right] \] (1.8)

(d) When the target nucleon and the neutron are considered to be at rest (i.e., neuton-nucleon), then
In labra:

\begin{align}
1 &= \frac{1}{10} \sum_{i=1}^{10} n_i \text{ for the}\n
\begin{align}
2 &= \frac{1}{3} \sum_{i=1}^{3} n_i \\
3 &= \frac{2}{2} \sum_{i=2}^{2} n_i \\
4 &= \frac{1}{4} \sum_{i=1}^{4} n_i \\
5 &= \frac{1}{5} \sum_{i=1}^{5} n_i
\end{align}

\text{spectator and the}\n
\begin{align}
1 &= \frac{1}{1} \sum_{i=1}^{1} n_i \\
2 &= \frac{1}{2} \sum_{i=2}^{2} n_i \\
3 &= \frac{1}{3} \sum_{i=3}^{3} n_i \\
4 &= \frac{1}{4} \sum_{i=4}^{4} n_i \\
5 &= \frac{1}{5} \sum_{i=5}^{5} n_i
\end{align}

\text{atom}\n
\begin{align}
1 &= \frac{1}{1} \sum_{i=1}^{1} n_i \\
2 &= \frac{1}{2} \sum_{i=2}^{2} n_i \\
3 &= \frac{1}{3} \sum_{i=3}^{3} n_i \\
4 &= \frac{1}{4} \sum_{i=4}^{4} n_i \\
5 &= \frac{1}{5} \sum_{i=5}^{5} n_i
\end{align}

\text{interaction}\n
\begin{align}
1 &= \frac{1}{1} \sum_{i=1}^{1} n_i \\
2 &= \frac{1}{2} \sum_{i=2}^{2} n_i \\
3 &= \frac{1}{3} \sum_{i=3}^{3} n_i \\
4 &= \frac{1}{4} \sum_{i=4}^{4} n_i \\
5 &= \frac{1}{5} \sum_{i=5}^{5} n_i
\end{align}

\text{interaction with a neutral current}\n
\begin{align}
1 &= \frac{1}{1} \sum_{i=1}^{1} n_i \\
2 &= \frac{1}{2} \sum_{i=2}^{2} n_i \\
3 &= \frac{1}{3} \sum_{i=3}^{3} n_i \\
4 &= \frac{1}{4} \sum_{i=4}^{4} n_i \\
5 &= \frac{1}{5} \sum_{i=5}^{5} n_i
\end{align}

\text{interaction with a neutral current with an external}\n
\begin{align}
1 &= \frac{1}{1} \sum_{i=1}^{1} n_i \\
2 &= \frac{1}{2} \sum_{i=2}^{2} n_i \\
3 &= \frac{1}{3} \sum_{i=3}^{3} n_i \\
4 &= \frac{1}{4} \sum_{i=4}^{4} n_i \\
5 &= \frac{1}{5} \sum_{i=5}^{5} n_i
\end{align}

\text{interaction with an external field}\n
\begin{align}
1 &= \frac{1}{1} \sum_{i=1}^{1} n_i \\
2 &= \frac{1}{2} \sum_{i=2}^{2} n_i \\
3 &= \frac{1}{3} \sum_{i=3}^{3} n_i \\
4 &= \frac{1}{4} \sum_{i=4}^{4} n_i \\
5 &= \frac{1}{5} \sum_{i=5}^{5} n_i
\end{align}
where $z_3$ is the third component of the isospin and $a_y$ ($z = 0,1$) are the amplitudes for the isoscalar ($z = 0$) and isovector ($z = 1$) vector currents while $a_A$ ($z = 0,1$) are the corresponding quantities for axial vector currents.

We express the matrix elements of the hadronic neutral current between the two nucleon states as

$$
\langle n (n') | J_\lambda^{\pi 0} | n (n') \rangle = \bar{u}(n') \left[ F_1^{\pi 0}(q^2) Y_\lambda + \frac{2}{Q M} \delta_{\lambda 0} F_2^{\pi 0}(q^2) \gamma_\lambda \right. \\
- \left. g_\lambda(q^2) Y_\lambda Y_5 - h_\lambda(q^2) \gamma_\lambda Y_5 \right] \frac{1}{M} \langle n | J^{\pi 0} | n \rangle (1.19)
$$

and assume that in analogy with the charged current, there is no definite evidence for the presence of the second class currents. The effect of the pseudoscalar term is very small and is neglected.

The transition matrix element for the process can be written in the standard notation as

$$
\mathcal{M} = \frac{g^2}{2} \bar{u}(n') \left\{ (1 - \gamma_5) u(n) \right\} \langle m | J_\lambda^{\pi 0} | d \rangle (1.19')
$$

where $k$ and $k'$ are the initial and final lepton (neutrino) momenta measured in the deuteron rest frame (i.e., lab frame) and $\langle m | J_\lambda^{\pi 0} | d \rangle$ is the matrix element of the hadronic current, arrived in the impulse approximation \[^{50}\text{to be}^\] $\langle m | J_\lambda^{\pi 0} | d \rangle = \int \Phi^*_\nu(z) \left[ \lambda_\lambda e^{i \frac{\mathbf{z}}{2}} + \lambda_\lambda e^{i \frac{\mathbf{z}}{2}} \right] \Phi^*_\nu(d) d^3 \mathbf{r} (1.19)$
\[ \bar{q} = \bar{p} - \frac{m}{E_{\bar{p}} c^2} \] is the three momentum transfer. \( \varphi_i(\vec{r}) \) is the initial wave function, \( \varphi_f(\vec{r}) \) is the final wave function and \( A_\lambda^{(\tau)}(\vec{r} = r, n) \) is the normal correlation of the single nucleon monomer.

The matrix element is derived between the two nucleon states by equation \((1.12)\) and written as follows:

\[ \langle \nu(\vec{r}') | \Gamma_A \lambda | \nu(\vec{r}) \rangle = \nu(\vec{r}) \left[ \frac{\bar{q} \cdot \vec{n}}{2M} \delta_{\lambda^* \lambda} \bar{\eta}_n F_2^{(\lambda)}(\vec{r}) \right] \]

\[ = \delta^{\lambda^* \lambda} \nu(\vec{r}) \nu(\vec{r}') \int \frac{d^4k}{(2\pi)^4} \nu(k) \]

The total form factors \( F_1^*(q^2) \), \( F_2^*(q^2) \), and \( F_3^*(q^2) \), refer to the charge correlation process and refers to the proton or neutron. These form factors for usual vector and axial vector currents are related with the form factors (isovector and isoscalar components) occur in the process of weak and electromagnetic process. In terms of isovector and isoscalar components, we get

\[ F_1^V = F_1^S + \frac{1}{2} \gamma_3 F_1^I \]
\[ F_2^V = F_2^S + \frac{1}{2} \gamma_3 F_2^I \]
\[ F_3^V = F_3^S + \frac{1}{2} \gamma_3 F_3^I \]

where \( \gamma_3 = +1 \) describes the process \( \nu + p \rightarrow \nu + p \) and \( \gamma_3 = -1 \) describes the process \( \nu + n \rightarrow \nu + n \). The quantities \( F_1^S, F_2^S \),
\( f_{1}^{0} \) denotes the isoscalar-vector and isoscalar-axial vector components of the currents and \( f_{1}^{1}, f_{2}^{1} \) are those for isovector components of vector and axial vector currents respectively. The quantities \( f_{1}^{2}, f_{2}^{2} \) for the neutral currents the role played by Fermi, magnetic and Coloumb-faller coupling constants for charged currents. At small \( a^{2} \) these quantities are related to the vector and axial vector amplitudes \( a_{U, A}^{0} ( \approx a_{0}^{2}) \) by the relations

\[
\begin{align*}
  f_{1}^{0} &= 1.5a_{V}^{0}, \quad f_{2}^{0} = 1.5a_{V}^{0}(\mu_{p} + \mu_{n}) \quad a_{V}^{0} = 0.3a_{A}^{0}(r_{A}/r_{V}) \\
  f_{1}^{1} &= a_{V}^{1}, \quad f_{2}^{1} = a_{V}^{1}, \quad a_{A}^{0} = a_{A}^{0} (7.17)
\end{align*}
\]

where \( f_{1}^{00}, f_{2}^{00} \) and \( f_{1}^{00} \) are the quantities from the charged currents and are given as

\[
\begin{align*}
  f_{1}^{00} &= 1.0, \quad f_{2}^{00} = \mu_{p} - \mu_{n}, \quad f_{1}^{00} = (r_{A}/r_{V}) = 1.04 \quad (7.18)
\end{align*}
\]

with

\[
\mu_{p} = 1.79 \quad \text{and} \quad \mu_{n} = 1.91
\]

In the light of the unified gauge theories (\(^{[12\text{a}-12\text{c}]}\)) for weak neutral currents, there exist various neutral currents models in the literature proposed by many people\(^{[4-7]}\) which predict the value of isoscalar and isovector parameters \( a^{0} \)
and $\alpha_A^2$. The well known models usually satisfy an $SU(2) \times U(1)$ group algebra for the weak and electromagnetic interaction and predict the essential parameters in it. The essential parameters are $a$, $b$ which are internally fixed with each model, as well as the Weinberg angle $\theta_W$ (a parameter in "weinberg - salam theory"), and the ratio $\sin^2 \theta_W = \left( \frac{m_W}{m_\pi} \right)^2$, where $m_W$ and $m_\pi$ are masses of intermediate vector bosons. These parameters are determined by fitting high energy experimental data as given in Table I.

The isoscalar and isovector constants appearing in the equation (I.11) are also given in this table which are calculated by relations

$$ a_V^0 = \frac{1}{2} \int (a-b - \frac{1}{2} \sin^2 \theta_W) $$

$$ a_A^0 = \frac{1}{2} \int (b-a) $$

$$ a_V^1 = \frac{1}{2} \int (2a+b-4\sin^2 \theta_W) $$

$$ a_A^1 = \frac{1}{2} \int (2a-b) $$

The form factor $P_{1,2}^\pi(n)$, $P_A(n)$ are calculated for proton and neutron separately by equations (I.16) - (I.19) for various models and given in Table II.

Now we can calculate the matrix element $M$ as defined by equation (I.13) using the transition operator's matrix
element given by equation (I.14), if we know how to describe initial state involving deuteron and final state involving the two nucleons in a relativistically covariant way. Since the nucleons inside deuteron are at rest we can take a non-relativistic wave function for the deuteron using small spinors. We calculate the non-relativistic reduction \( \Lambda_2^N \) of the single nucleon operator \( \bar{t} \Gamma_2^N \) between the two nucleon states retaining the terms up to order \( \bar{n}^2/m^2 \). The terms of the order \( \bar{n}^2/m^2 \) are related through the energy momentum conservation to \( \bar{p}^2/m^2 \) where \( \bar{p} \) is the relative momentum of the two nucleons inside deuteron. The relations between \( \bar{n} \) and \( \bar{n}' \) have been obtained in earlier section (II-1) as given by equation (I.9). As the magnitude of \( \bar{n} \) becomes 1 more the nucleons inside deuteron are off the mass shell and therefore the large \( \bar{p}^2/m^2 \) terms are related to the off mass shell effects which we should neglect. The effect of neglecting the off mass shell effects, retaining the terms of \( \bar{n}^2/m^2 \) order and neglecting the terms higher than \( \bar{n}^2/m^2 \) are briefly discussed in the chapter III.

Using the above mentioned techniques we evaluate the nonrelativistic reduction \( \Lambda_2^N \) as outlined in Appendix A. We get the following result for the operators \( \Lambda_2^N \) and \( \Lambda_0^N \):

\[
\Lambda_2^N = \left[ \Gamma_2^N (\gamma^0) \frac{\hat{P}_j}{M} + \left( \Gamma_2^N (\gamma^0) \gamma^j \frac{\hat{P}_j}{2N} \right) \gamma_R - \gamma_0^N \gamma^j \left( (1 - \gamma_{3Nj}) \gamma^0 \right) \gamma_R + \frac{1}{2 M^2} \gamma^j \gamma_R \right] \\
+ \frac{1 - \gamma_{3Nj}}{8 N^2} \gamma^j \gamma_R.
\]
and

\[ \Lambda^N_c = \left[ F^u_R(y, 1 + \frac{L}{M^2}) - \frac{\Gamma^u}{M^2} \right] \left[ \frac{\phi^u(q^2)}{\phi^u(q^2)} \right] + \frac{2 \epsilon}{3} \frac{F^u_N(y)}{4 M^2} \phi^u(q^2) \]

(I.39)

where

\[ F^u_{\mu}(q^2) = F^u_R(q^2) + \frac{\phi^u}{\phi^u(q^2)} F^u(q^2) \]

\[ F_{\mu N}(q^2) = F^u_R(q^2) + 2 F^u(q^2) \]

(I.40)

(I.2.2) Helicity Limiting Theories

We shall consider here the most general \( n, p, T \) (helicity limiting) neutral current which can be formed from members of the usual quark model scalar, pseudoscalar and tensor currents.

The 3 \( \overline{P} \) conserving neutral current effective interaction is given by \( 69-71 \)

\[ \mathcal{L}^{\mu} = \frac{1}{\sqrt{2}} \left\{ \frac{\phi^u(q^2)}{\phi^u(q^2)} \right\} \left( \overline{f} \gamma^\mu \gamma^5 f + \overline{f} \gamma^\mu \gamma^\lambda \gamma^5 \right) \]

(I.42)

where \( \overline{f} \) is the neutrino field and \( \overline{f}, \gamma^5, \gamma^\lambda \) are the hadronic scalar, pseudoscalar and tensor currents respectively.

Since the incoming neutrino (anti-neutrino) is left (right) handed, the effective matrix element is obtained from equation (I.42) by making the substitution \( \overline{f} \rightarrow \frac{1}{2} (\gamma_5 \overline{f}) \), ensuing

\[ \mathcal{L}^{\mu} = \frac{1}{\sqrt{2}} \left\{ \frac{\phi^u(q^2)}{\phi^u(q^2)} \right\} \left( \overline{f} \gamma^\mu \gamma^5 f \overline{f} \gamma^\lambda \gamma^5 \right) \]

(I.43)
The most general-quark-model structure for the scalar,
pseudo-scalar and tensor currents is

\[ J = g_0 J_0 + g_\gamma J_\gamma + g_\rho J_\rho \]
\[ J^0 = g_0 J_0^0 + g_\gamma J_\gamma^0 + g_\rho J_\rho^0 \]  
\[ J^{\lambda\gamma} = g_0 J_0^{\lambda\gamma} + g_\gamma J_\gamma^{\lambda\gamma} + g_\rho J_\rho^{\lambda\gamma} \]  

(1.24)

where in this fig: \( i = 1, \rho, \gamma; J = 0, 3, 8 \) are real numbers.

The \( J_0, J_0^0 \) and \( J_0^{\lambda\gamma} \) \( \gamma \)-mixing currents are written
in quark model (with quark "field map") as

\[ J_0 = \bar{u} \gamma_\tau v \]
\[ J_0^0 = \bar{u} \gamma_\tau v \gamma_5 \]
\[ J_0^{\lambda\gamma} = \bar{u} \gamma_\lambda \gamma_\gamma v \]  

(1.25)

We express the nucleon matrix element of the neutral
member of these currents notate (neglecting the second \( \gamma \)-currents as

\[ \langle N(h) | J | N(h) \rangle = \gamma_\lambda \bar{u}(h) \Gamma_8 \left( g_\lambda \right) t_\gamma u(h) \]
\[ \langle N(h) | J_0^0 | N(h) \rangle = \gamma_\lambda \bar{u}(h) \Gamma_5 \left( g_\lambda \right) t_\gamma u(h) \]
\[ \langle N(h) | J_0^{\lambda\gamma} | N(h) \rangle = \gamma_\lambda \bar{u}(h) \left[ \Gamma_1 \left( g_\lambda \right) \gamma_\lambda \gamma_\gamma + \frac{i T_3 \left( g_\lambda \right) \left( -\gamma_\tau \gamma_\gamma - \gamma_\gamma \gamma \right)}{M} \right] u(h) \]  

(1.26)
where
\[ a = \sqrt{a^2 - a^1}, \quad t^2 = t_1 + t_2 = \frac{1}{2} \sqrt{a^2} \]
\[ t^3 = \frac{1}{2} \xi_3, \quad t_0 = \frac{1}{2} (\xi_1 + \xi_2), \quad t_p = \frac{1}{2} (1/3)^{\frac{1}{2}}. \]  

Since the nucleons involved are nonrelativistic we may

a nonrelativistic reduction of these single nucleon center-

terms given in Appendix A, and write the matrix elements between

nucleons state (omitting from kinematic factors) as

\[ N = \sum_{i,j} \left[ \xi_j \left( \frac{1}{2} \left( A_i \right) \chi_i \left( \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \chi_i \right) + \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \chi_i \right) \right] 

\] 

\[ + \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \chi_i \left( \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \chi_i \right) \right] 

\] 

where the total nuclear form factors \( F_0(0^2), \quad F_0(0^2) \) and

\( F_1(0^2) \) are defined as

\[ F_0(0^2) = \sum_{i,j} \left( \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \chi_i \right) \] 

\[ + \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \chi_i \left( \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \sum_{\xi_i} \frac{1}{2} \chi_i \right) \]
\[ \chi_1 \text{ and } \chi_2 \text{ are multi spin operators for initial and final state nucleons and } \varepsilon = \varepsilon_1 \text{ described the process } \pi^0 \rightarrow \pi^+ \pi^-; \varepsilon = \varepsilon_2 \text{ described the process } \pi^+ \rightarrow \pi^0 \pi^- . \]

The numerical value of renormalization constant for scalar, pseudoscalar and tensor couplings is given in Table III and then the numerical value of those form factors for proton and neutron has been calculated from them in the quark model phenomenological theory by equation (1,20) and given in Table IV.

The matrix element for the process can be written in the inverse approximation from equation (1,28) as

\[ M = \left[ \frac{g}{2} \left( \chi_1 \chi_2 \right) - \left\langle n \eta \mid \Lambda \mid d \right\rangle + \varepsilon_1 \varepsilon_2 \left( \left\langle n \eta \mid \Lambda_1 \mid d \right\rangle \right) \left( \varepsilon_1 \varepsilon_2 \left\langle n \eta \mid \Lambda_2 \mid d \right\rangle \right) \right] (1.30) \]

where

\[ \left\langle n \eta \mid \Lambda \mid d \right\rangle = \int p^*(r) \left( \Lambda_1 \xi \xi \xi \xi + \Lambda_2 \xi \xi \xi \xi \right) f_1(r) \rho(r) d^3r \]

\[ \left\langle n \eta \mid \Lambda_1 \mid d \right\rangle = \int p^*(r) \left( \Lambda_1 \xi \xi \xi \xi \right) f_1(r) \rho(r) d^3r \]

\[ \left\langle n \eta \mid \Lambda_2 \mid d \right\rangle = \int p^*(r) \left( \Lambda_2 \xi \xi \xi \xi \right) f_1(r) \rho(r) d^3r \]

and \( \Lambda_1, \Lambda_2 \) and \( \Lambda_3 (\varepsilon = \varepsilon_1) \) are given in the lowest order of \( (1/M) \) by
\[ \psi^p = \pi^p(q^p)(\epsilon = -1) + a^p(q^p)(\epsilon = +1) \frac{\sigma^+}{2M} \]
\[ \psi^j = \pi^j(q^j)(\epsilon = -1) \xi_{jk} \frac{\sigma^y}{2M} \]
\[ \psi^l = \pi^l(q^l)(\epsilon = -1) \frac{\sigma^-}{2M} \]

\[ \psi^n = \pi^n(q^n)(\epsilon = -1) + a^n(q^n)(\epsilon = +1) \frac{\sigma^y}{2M} \]
\[ \psi^j = \pi^j(q^j)(\epsilon = -1) \xi_{jk} \frac{\sigma^y}{2M} \]
\[ \psi^l = \pi^l(q^l)(\epsilon = -1) \frac{\sigma^-}{2M} \]

(1.7) Wave functions

In this section we shall describe the initial state wave function for the deuteron \( \psi_1^p(r) \) and the final state wave function for dimuon, \( \psi_f^j(r) \). As we have discussed earlier that the matrix elements for helicity conserving and helicity flipping theories can only be determined if we know the suitable form of these wave functions \( \psi_1^p(r) \) and \( \psi_f^j(r) \).

Here, the final state wave function is described in the absence of the final state interactions by a plane wave.

The calculation of the final state interaction effect, however, involves a detailed knowledge of \( \psi_f^j(r) \), discussion of which we defer until chapter IV.

In order to describe the nonrelativistic wave function for the deuteron we are faced with certain difficulties on the
wave function is described by the usual spinor and the relative motion of the two nucleons inside deuteron in to be described by the bound state wave function. So it is proper assumption to choose the equal time wave function for initial deuteron and final dimuon states, i.e.

\[ \Phi (r) = \phi (r) \chi (r) \]  \hspace{1cm} (7.24)

Thus we write the deuteron wave function as

\[ \psi (r_p, r_n) = e^{i \Phi (r)} \phi (r) \]  \hspace{1cm} (7.24a)

with

\[ r = \frac{r_p + r_n}{2} \quad \text{and} \quad r = r_p - r_n \]  \hspace{1cm} (7.24b)

where \( e^{i \Phi (r)} \) factor correctly describes the centre of mass motion of the two nucleons inside deuteron. The space part of initial state wave function, \( \phi (r) \) is the usual triplet wave function for the deuteron where the time coordinates of the two nucleons are taken to be equal. \( \phi (r) \) is not very well known at very short distances and thus we choose the Bohmian wave function\(^{75}\) neglecting P states as

\[ \phi (r) = \frac{1}{\sqrt{4\pi}} \frac{U(r)}{r} \chi _{\text{in}} \]  \hspace{1cm} (7.24c)

where \( U(r) \) is the radial part of the wave function with \( r \) as the relative position of the proton with respect to
neutron in deuteron and given for Nelthen wave function by

\[ \psi(r) = \sqrt{\frac{\lambda R}{(a^2 + \lambda^2)^{\frac{3}{2}}}} \left( \frac{r + a}{a^2 + \lambda^2} \right) \times (e^{-\lambda r} - e^{-\lambda a}) \]  \hspace{1cm} (7.36)

with

\[ \lambda = 46.5 \text{ keV} \quad R = 2.57 \text{ fm}. \]

In a way similar to the equation (1.34) we write down the wave function for final state as

\[ \psi_i(r_1, r_2, \ldots, r_n) = e^{-i(p_i \cdot r_i)} \cdot \phi_i(r) \]  \hspace{1cm} (7.37)

where \( \phi_i(r) \) the space part of \( \varphi_i(r) \) in the radial wave function of the final two-nucleon state obtained by solving the equal time wave function given in equation (7.36). In first, neglect the final state interaction effect on the wave function, the \( \phi_i(r) \) can be expressed by

\[ \phi_i(r) = e^{i\vec{p}_i \cdot \vec{r}} \chi_i \]  \hspace{1cm} \text{for triplet state} \hspace{1cm} (7.38)

\[ \chi_i \]  \hspace{1cm} \text{for singlet state}

\[ \chi_i, \chi_0 \]  \hspace{1cm} \text{being the null spinors for the triplet and singlet states respectively, \( \vec{p}_i \) is the relative momentum of containing nucleon in their centre of mass frame which is suppressed by equation (1.5) in section (I-1).}