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*A Measurement of $\Delta\sigma_L(np)$, the
Difference between Neutron-Proton
Total Cross Sections in Pure
Longitudinal Spin States*

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A MEASUREMENT OF $\Delta\sigma_L(np)$, THE DIFFERENCE BETWEEN
NEUTRON-PROTON TOTAL CROSS SECTIONS IN PURE LONGITUDINAL
SPIN STATES

by

MICHAEL ERVIN BEDDO

ABSTRACT

A measurement of $\Delta\sigma_L(np)$, the difference between neutron-proton total cross sections in pure longitudinal spin states, is described. The results will help determine the isospin-zero ($I = 0$) scattering amplitudes, which are not well known above laboratory energies of 500 MeV, whereas the isospin-one ($I = 1$) amplitudes are fairly well-determined to 1 GeV. Data points were taken at the Los Alamos Meson Physics Facility (LAMPF) at Los Alamos, New Mexico, for five neutron beam energies: 484, 568, 634, 720 and

788 MeV; they are the first in this energy range. Polarized neutrons were produced by charge-exchange of polarized protons on a liquid deuterium target (LD_2). Large-volume neutron counters detected the neutrons that passed through a polarized proton target. The counters subtended a range of solid angles large enough to allow extrapolation of the scattered neutrons to 0° . Two modifications to the LAMPF accelerator system which were made for this work are described. They included a "beam buncher," which modified the normal rf-time structure of the proton beam and allowed for the selection of peak-energy neutrons by time-of-flight means, and a computerized beam steering program, which reduced systematic effects due to beam motion at the LD_2 target. The experimental values of $\Delta\sigma_L(np)$ are found to be consistent with other np data, including preliminary data from SIN and Saclay, but not with some results from Argonne which used a polarized proton beam and a polarized deuteron target. The $I = 0$ component was extracted from $\Delta\sigma_L(np)$ using existing pp data ($I = 1$), with the unexpected result that $\Delta\sigma_L(I = 0)$ was found to be essentially identical in shape to $\Delta\sigma_L(I = 1)$. The significance of this is not yet understood.

CHAPTER 1 INTRODUCTION

1.1 Spin-Dependent Cross Section Experiments

This thesis will describe a recent measurement of $\Delta\sigma_L^{\text{tot}}(np) = \sigma(\vec{\uparrow}) - \sigma(\vec{\downarrow})$ which is the difference between total cross sections for the scattering of longitudinally polarized neutrons and protons with spins opposed and parallel. The longitudinal direction is defined to be the direction along the beam axis. The experiment was done at LAMPF at Los Alamos, New Mexico during the years 1987-1988, with the participation of Argonne National Laboratories, New Mexico State University, Texas A&M University, Los Alamos National Laboratories, University of Montana, Washington State University and Earlham College.

A cross section measures how much a beam is attenuated by a target and it is one of the experimentally measured observables that is needed to determine the nucleon-nucleon (NN) scattering matrix. A cross section has the units of area — all of the cross sections mentioned in this thesis are in units of millibarns ($1 \text{ barn} = 10^{-24} \text{ cm}^2$). There are a total of five complex parameters,^[1] or amplitudes, that have to be determined to

specify the elastic scattering matrix because nucleons have spin $1/2$. There are as many degrees of freedom involved in the scattering process as there are independent bilinear scalars that can be formed from the Pauli spin matrices that represent spin $1/2$ objects. For this reason, numerous spin dependent experiments over a range of energies and angles are required^[2,3] to specify the NN scattering matrix. The NN experimental program at the Los Alamos Meson Physics Facility (LAMPF) is designed to help provide the necessary data to determine these amplitudes for medium energies, up to 800 MeV.

The proton-proton (pp) amplitudes have been well determined up to 1 GeV, but the neutron-proton (np or pn) amplitudes have not.^[4] This is due to the difficulties associated with np experiments. Until quite recently, neutron beams, especially polarized ones, were not at all common, so neutrons from deuterons were used instead in quasi-free scattering. Usually, proton beams are used on deuteron targets, and the pn data is sorted out from the proton-deuteron (pd) data which involves complicated theoretical corrections. Since the neutron is bound in the deuteron nucleus, Fermi motion and screening are significant. Depolarization effects are present also because the deuteron exists in an admixture of S and D angular momentum states. It has not been until recently that highly polarized neutron beams became practical.

Both the total cross sections, σ^{tot} and $\Delta\sigma_L$, have been well measured over medium energies for pp scattering. The striking energy dependence^[5,6,7] of $\Delta\sigma_L(pp)$ between lab momenta 1 and 2 GeV/c came as a surprise (see Fig. 1.1) because $\sigma^{tot}(pp)$ shows

Figure 1.1

Experimental values of $\Delta\sigma_L^{tot}(pp)$, compared with the total cross section, $\sigma^{tot}(pp)$. The L subscript refers to beam and target spins aligned in the beam direction. Note the large energy variation in $\Delta\sigma_L^{tot}(pp)$, while the total pp cross section is without structure.

little structure in this momentum range. The spin dependent cross section $\Delta\sigma_T(pp)$, for transverse beam and target spins, also shows structure^[8,9] as seen in Fig. 1.2. A great deal of interest arose concerning the interpretation of these data because the energy dependence is quite large in comparison with the total cross section. It was postulated that these data were supportive evidence for "dibaryon" resonances^[10] which are bound states for six-quark objects. In fact, the "Bag Model" predicts many dinucleon resonances^[11] within the mass range 2.1 – 2.5 GeV. Others pointed out that pp scattering at medium energies is "clouded" by numerous inelastic thresholds (see Ref. [12], for instance), especially those thresholds associated with the $N\Delta$ and πd channels which open up near laboratory energies 630 MeV and 425 MeV, respectively. For a complete overview of experimental searches for dibaryons, see Ref. [13].

A complete picture of NN scattering requires a complete set of np experiments as well. This is because nucleons scatter in two isospin states. Isospin (I) is a quantum number which is conserved in strong interactions and has commutation relations similar in form to those of angular momentum.^[14] The isospin of a nucleon is $\frac{1}{2}$, and the proton is assigned $+\frac{1}{2}$ for the third component of isospin, I_3 . Similarly, the neutron has isospin $-\frac{1}{2}$ for I_3 . The isospin components I_1 and I_2 correspond to no measurable observable. Two protons always scatter in the isospin $I = 1$ state, while np scattering occurs from both isospin $I = 0$ and $I = 1$ states. The isospin $I = 0$ channel is a fundamental channel, and np scattering becomes important because at medium energies many of

Figure 1.2

Experimental values of $\Delta\sigma_T(pp)$, compared with the total cross section, $\sigma^{tot}(pp)$. The T subscript refers to beam and target spins aligned transversely to the beam direction. Note the large energy variation in $\Delta\sigma_T^{tot}(pp)$, while the total pp cross section is without structure.

the inelastic channels in pp scattering cannot occur from the isospin $I = 0$ state, for example, the $N\Delta$ and πd channels occur only from the isospin state $I = 1$ in NN scattering. The isospin $I = 0$ and $I = 1$ total cross sections (σ^{tot}) and their inelastic components (σ^{inel}) are plotted in Fig. 1.3. The isospin $I = 0$ cross sections σ^{tot} and σ^{inel} are markedly different from the isospin $I = 1$ cross sections above the threshold for pion production; the inelastic cross section $\sigma^{inel}(I = 1)$ rises rapidly and levels off at 25 mb, whereas the inelastic cross section $\sigma^{inel}(I = 0)$ is rising only slowly and is small. Consequently, np scattering in the isospin $I = 0$ state is mostly elastic,^[15,16,17] while the proposed dibaryon resonances are thought to be highly inelastic. Total cross section measurements are more sensitive to inelasticities than other scattering observables, and therefore the existence of highly inelastic $I = 0$ dibaryons should be easily detectable in a $\Delta\sigma_L$ experiment.

An attempt was made at the Argonne Zero-Gradient Synchro-cyclotron (ZGS) to make the first measurement of $\Delta\sigma_L(np)$ using a polarized proton beam and a polarized deuteron target.^[18] The spin-dependent cross section difference $\Delta\sigma_L(pd)$ was measured, and extensive theoretical corrections were done to find $\Delta\sigma_L(pn)$ from the deuteron data (see Fig. 1.4), since $\Delta\sigma_L(pn)$ is only approximately equal to $\Delta\sigma_L(pd) - \Delta\sigma_L(pp)$. As a systematic check, they also extracted $\Delta\sigma_L(pp)$ and found it consistent with $\Delta\sigma_L(pp)$ for free pp . However, the calculations were later found to be in error and were repeated. The revised ZGS $\Delta\sigma_L(pn)$ data^[19] are essentially flat (Fig. 1.5), and since pn data contains

*NN cross sections
for $I=0$ and $I=1$*

Figure 1.3

Comparison between the isospins $I = 0$ and $I = 1$ total and inelastic cross sections for NN scattering. The isospin $I = 0$ cross sections are markedly different from the isospin $I = 1$ cross sections above the threshold for pion production; the inelastic cross section for isospin $I = 1$ rises rapidly and levels off at 25 mb, whereas the inelastic cross section for isospin $I = 0$ is rising only slowly and is small.

Figure 1.4

The first $\Delta\sigma_L(pn)$ "data^[18]" (b), which were extracted from $\Delta\sigma_L(pd)$ (a). The isospin $I = 0$ component (c) was found using $\Delta\sigma_L(pp)$ data, which is isospin $I = 1$. The theory used to extract the np data (b) was later found to be in error and the calculations repeated. Figure taken from Ref. [18].

contributions from isospin $I = 0$ and $I = 1$ states, significant structure is implied in the isospin $I = 0$ channel as well. Because the theory involved is extremely complicated and not at all certain, further experiments which measure $\Delta\sigma_L(np)$ for free np scattering are essential. The free experiment has the added feature of helping to refine the theory for application to other data involving deuterons. Recently, $\Delta\sigma_L(np)$ (free) has been measured at various laboratories including SIN,^[20] Saclay,^[21,22] and LAMPF. The rest of this thesis will describe in detail the experiment that was done at LAMPF.

1.2 An Overview of the Experiment

In this experiment, accelerated polarized protons bombarded a liquid deuterium target to produce polarized neutrons which were then collimated to deliver a neutron beam to the experiment. The spins of the neutrons were aligned either parallel or antiparallel to the direction of the beam, using spin precession magnets. The longitudinally polarized neutron beam then passed through a proton target, which was also polarized either parallel or antiparallel to the beam direction. Typical values for the polarizations of the neutron beam and the proton target were about ± 0.6 and ± 0.7 , respectively. During the experiment the spins of the neutron beam and the proton target were flipped periodically to reduce systematic effects in the data.

Transmitted neutrons and scattered neutrons which passed through the proton target were detected downstream in a large volume segmented neutron hodoscope. From

Figure 1.5

The revised $\Delta\sigma_L(pn)$ data, taken from Ref. [19]. Compare with Fig. 1.4b; the revised data are smaller in magnitude by 4 mb, but remain essentially featureless in this momentum region.

the positions of the detected neutrons in the hodoscope the square of the momentum transfer (t) was calculated. A spin asymmetry was then calculated for each bin in the hodoscope:

$$\epsilon(t) = \frac{T^+(t) - T^-(t)}{T^+(t) + T^-(t)},$$

where $T^+(t)$ and $T^-(t)$ represented the number of beam normalized neutrons detected for a particular momentum transfer squared, with beam and target spins parallel (+) and opposed (-). The value of $\epsilon(t)$ at $t = 0$ was the final objective because $\epsilon(0)$ represents the spin asymmetry from the amount of beam transmitted through the target at 0° , which is related to the total cross section difference $\Delta\sigma_L$. Spin asymmetries at values of t other than zero were contaminated by neutrons scattered out of the beam through elastic and inelastic processes, and this contamination was removed by determining the functional form for the spin asymmetry $\epsilon(t)$ and applying the limiting process:[23]

$$\Delta\sigma_L \propto \lim_{t \rightarrow 0} \epsilon(t).$$

The spin asymmetries were fitted with a straight line and extrapolated to $t = 0$ to obtain $\epsilon(0)$.

The extrapolated asymmetry $\epsilon(0)$ is related to the total cross section difference $\Delta\sigma_L$ in a simple way. At 0° the number of beam particles transmitted through the target (N)

is expressed in terms of the number of incident beam particles (N_0) using the relation:

$$N^{\pm} = N_0^{\pm} \exp \left[-\alpha - \frac{1}{A} \left(\sigma^{\text{tot}} \mp P_n P_t \frac{\Delta \sigma_L}{2} \right) \right],$$

for beam and target spins parallel (+) and (-) opposed. The parameter α represents an attenuation constant for all material in the target except for the free protons, σ^{tot} is the spin-averaged np total cross section and A is the "target constant" for the proton target. The enhancement on N due to spin effects depends on the product of the neutron beam and proton target polarizations, $P_n P_t$.

The extrapolated asymmetry $\epsilon(0)$ is related to the number of transmitted beam particles by

$$\begin{aligned} \epsilon(0) &= \frac{N^+/N_0^+ - N^-/N_0^-}{N^+/N_0^+ + N^-/N_0^-} \\ &= \frac{\exp \left[\frac{P_n P_t}{2A} \Delta \sigma_L \right] - \exp \left[-\frac{P_n P_t}{2A} \Delta \sigma_L \right]}{\exp \left[\frac{P_n P_t}{2A} \Delta \sigma_L \right] + \exp \left[-\frac{P_n P_t}{2A} \Delta \sigma_L \right]}. \end{aligned}$$

The quantity $P_n P_t / 2A \cdot \Delta \sigma_L$ was sufficiently small that the approximation $\exp(x) \approx 1+x$ was used to obtain

$$\epsilon(0) = \frac{P_n P_t}{2A} \cdot \Delta \sigma_L.$$

The sensitivity of the experiment depended on the target constant A , which was defined as

$$A = (\rho_F L)^{-1},$$

ρ_F is the average number of free protons per unit volume in the target, and L is the length of the target. The target constant A has units of area.

Two new beam-control systems were set up at LAMPF for this experiment. One was a beam buncher,^[24] which compressed beam current into narrow bunches separated in time by 100 nsecs. This allowed the experiment to use time-of-flight information to select fast neutrons uniquely at the hodoscope. The second was a beam steering system. It had been known from earlier experiments that the beam position at the neutron production target was not as stable as required, so a computerized beam steering program was developed,^[25] which controlled the proton beam at the liquid deuterium target. A separate chapter has been set aside to describe these systems.

Various systematic effects associated with the LAMPF accelerator were studied in the off-line analysis of the data, and some pertinent correlations were found. These are discussed in later chapters.

1.3 Nucleon-Nucleon Scattering Formalism for $\Delta\sigma_L(np)$

Nucleon-Nucleon (NN) scattering experiments have indicated that the strong interaction is invariant under the symmetries of parity, spatial rotation, charge conjugation, and time reversal.^[26] This information shows that the NN scattering matrix, which contains all of the physics of the interaction, is a scalar quantity under each of the above transformations. With the above mentioned symmetries, five amplitudes are necessary for each isospin state to parameterize the NN scattering matrix. This follows from the consideration of the properties of spin 1/2 nucleons. Nucleons can scatter in

either a singlet state (with total spin $S = 0$) or one of three triplet states (with total spin $S = 1$). These spin states (χ) are listed below where the orbital angular momentum L is given for arbitrary total angular momentum J :

- singlet state ($S = 0; L = J$)

$$\chi(S = 0; S_z = 0) = \frac{1}{\sqrt{2}} (\chi(+-) - \chi(-+))$$

- triplet states ($S = 1; L = J - 1, J, J + 1$)

$$\chi(S = 1; S_z = 1) = \chi(++)$$

$$\chi(S = 1; S_z = 0) = \frac{1}{\sqrt{2}} (\chi(+-) + \chi(-+))$$

$$\chi(S = 1; S_z = -1) = \chi(--).$$

The two-nucleon wave function

$$\Psi(1, 2) = \chi_{spin} \phi_{isospin} \phi_{space}$$

must be antisymmetric. Therefore, transitions between singlet and triplet states do not occur (because of conservation of parity, isospin and the Pauli exclusion principle), but transitions among triplet states are allowed.

A physics argument explains why only five amplitudes are required for each isospin state.^[1] Two amplitudes are required to describe both singlet-singlet and triplet-triplet

transitions where $L = J$. To describe the other possible triplet-triplet transitions

$$L = J - 1 \rightarrow L = J + 1$$

$$L = J - 1 \rightarrow L = J - 1$$

$$L = J + 1 \rightarrow L = J + 1$$

three amplitudes are needed. The transition where $L = J + 1 \rightarrow L = J - 1$ is a time reversed process to the first transition shown above and therefore an additional independent amplitude is not required.

NN scattering formalism has been discussed in quite readable form by several authors, including Goldberger et al.,^[1] Hoshizaki^[2] and Bystricky et al.^[3] The helicity formalism was introduced by Jacob and Wick^[27] as a way of simplifying the appearance of spin in certain relativistic equations. The helicity of a particle with spin is defined as the projection of the spin of that particle onto the direction of its momentum, or $h = \vec{J} \cdot \vec{p}/p$ (note that $\vec{L} \cdot \vec{p} = 0$, since $\vec{L} = \vec{p} \times \vec{r}$). In the elastic scattering of two particles, $a + b \rightarrow c + d$, the total helicities for the initial state are $\lambda = h_a - h_b$ and $\mu = h_c - h_d$ for the final state.

The five helicity amplitudes $\phi = \langle \mu | \phi | \lambda \rangle$ are conventionally taken as

$$\phi_1 = \langle ++ | \phi | ++ \rangle, \text{ non-flip}$$

$$\phi_2 = \langle ++ | \phi | -- \rangle, \text{ double-flip, no net flip}$$

$$\phi_3 = \langle +- | \phi | +- \rangle, \text{ non-flip}$$

$$\phi_4 = \langle +- | \phi | -+ \rangle, \text{ double-flip}$$

$$\phi_5 = \langle ++ | \phi | +- \rangle, \text{ single-flip}$$

ϕ represents the NN spin scattering matrix operator. The symmetry of the amplitudes depends upon the isospin state of the interacting nucleons:^[1]

$$\phi_1(\pi - \theta) = (-1)^{I+1} \phi_1(\theta)$$

$$\phi_2(\pi - \theta) = (-1)^{I+1} \phi_2(\theta)$$

$$\phi_3(\pi - \theta) = (-1)^I \phi_4(\theta)$$

$$\phi_5(\pi - \theta) = (-1)^I \phi_5(\theta).$$

The np amplitudes are a linear combination of isospin $I = 0$ and $I = 1$ amplitudes, given by

$$\phi_{i,np} = (\phi_{i,I=0} + \phi_{i,I=1})/2.$$

If the isospin $I = 1$ amplitudes are known from pp scattering experiments, the isospin $I = 0$ amplitudes can be determined up to an arbitrary phase (ie., $\phi_{I=0} = 2\phi_{np} - \phi_{I=1}$).

At a scattering angle of 0° the helicity amplitudes ϕ_4 and ϕ_5 vanish identically because total helicity is conserved at 0° . This leaves the three helicity amplitudes ϕ_1 ,

ϕ_2 and ϕ_3 which can be non-zero. Thus when spin is considered, there are three cross sections which can be measured for NN scattering that relate to these 0° amplitudes. These are: the spin-averaged total cross section σ^{tot} , and the spin-dependent total cross sections $\Delta\sigma_L$ and $\Delta\sigma_T$, defined as^[28]

$$\begin{aligned}\sigma^{tot} &= \frac{1}{2} (\sigma(\uparrow\downarrow) + \sigma(\uparrow\uparrow)) = \frac{1}{2} (\sigma(\overleftarrow{}) + \sigma(\overrightarrow{})) \\ \Delta\sigma_L^{tot} &= \sigma(\overleftarrow{}) - \sigma(\overrightarrow{}) \\ \Delta\sigma_T^{tot} &= \sigma(\uparrow\downarrow) - \sigma(\uparrow\uparrow)\end{aligned}$$

where the notations $\uparrow\downarrow, \uparrow\uparrow$ ($\overleftarrow{}, \overrightarrow{}$) represent beam and target spins opposed and aligned transversely and longitudinally to the beam direction. Spin directions are defined in the laboratory frame.

These cross sections can also be expressed in terms of the helicity scattering amplitudes through the use of the generalized optical theorem:

$$\begin{aligned}\sigma^{tot} &= \frac{2\pi}{k} \Im m (\phi_1(0) + \phi_3(0)) \\ \Delta\sigma_L^{tot} &= \frac{4\pi}{k} \Im m (\phi_1(0) - \phi_3(0)) \\ \Delta\sigma_T^{tot} &= -\frac{4\pi}{k} \Im m \phi_2(0)\end{aligned}$$

where k is the relative wave number for the nucleons in the center of mass frame. The imaginary parts of the amplitudes ϕ_1 , ϕ_2 and ϕ_3 are determined at 0° by measuring the three cross sections. Also, since the scattering matrix is unitary, the real parts of

the forward scattering amplitudes can be found using dispersion relations, and having the real and imaginary parts of these amplitudes at 0° determines the phase for $\phi_1(0)$, $\phi_2(0)$ and $\phi_3(0)$. Having these amplitudes completely specified is especially useful in np scattering, because there is a general lack of data for the forward angles.

There are three elastic scattering spin observables that are intimately connected to the elastic parts of $\Delta\sigma_L^{tot}$ and $\Delta\sigma_T^{tot}$; these are the spin correlation parameters C_{LL} , C_{SS} and C_{NN} , measured for spins longitudinal (L), sideways (S) and normal (N) to the scattering plane. Measurements of C_{LL} for elastic np scattering at center of mass angles greater than about 70° can be found in References [29] and [30], for example. These spin correlations measure the spin asymmetry in the elastic differential cross sections for nucleon-nucleon scattering for spins parallel and anti-parallel. The total cross sections are a sum of elastic and inelastic contributions,

$$\sigma^{tot} = \sigma^{elastic} + \sigma^{inelastic}$$

and similarly for $\Delta\sigma_L$ and $\Delta\sigma_T$. The elastic part of $\Delta\sigma_L$ is related to C_{LL} , but it has not been measured for np scattering at small forward angles. These total cross section differences can be expressed in terms of the spin correlation parameters as

$$\begin{aligned}\Delta\sigma_L^{elastic} &= \int_{4\pi} -2C_{LL} \cdot \frac{d\sigma}{d\Omega} d\Omega \\ \Delta\sigma_T^{elastic} &= \int_{4\pi} -(C_{SS} + C_{NN}) \cdot \frac{d\sigma}{d\Omega} d\Omega\end{aligned}$$

where $d\sigma/d\Omega$ is the spin-averaged elastic differential cross section. The integrals represent the amount of beam removed by elastic scattering. It may be some time in the future before these integrals can be evaluated for the np system because of the lack of data at forward angles, but once they are obtained the inelastic parts of $\Delta\sigma_L(np)$ and $\Delta\sigma_T(np)$ will become known.

The elastic parts of $\Delta\sigma_L$ and $\Delta\sigma_T$ have already been determined for pp scattering. There is more data at small angles for pp than for np scattering, and it is seen that the energy dependent structures in $\Delta\sigma_L(pp)$ and $\Delta\sigma_T(pp)$ are primarily caused by inelasticities,^[13] but the magnitude of $\Delta\sigma_L(pp)$ is primarily due to elastic scattering (see Fig. 1.6).

1.4 Nucleon-Nucleon Partial Waves

Data from spin scattering experiments for different energies and angles are combined in partial wave analyses to determine the NN amplitudes. Since nucleons scatter in singlet and triplet spin states, the amplitudes can be constructed from linear combinations of spin singlet and triplet partial waves. There are several groups active in partial wave and phase shift analysis (PSA), and each treats inelasticities differently. The solutions used here are those of Arndt et al.^[31] The ultimate goal is to completely determine the NN scattering matrix amplitudes in a model independent fashion.

In NN elastic scattering, total angular momentum, total parity and total isospin are

(a)

(b)

Figure 1.6

Elastic and inelastic contributions to $\Delta\sigma_L(pp)$ and $\Delta\sigma_T(pp)$. The energy-dependent structures in $\Delta\sigma_L(pp)$ and $\Delta\sigma_T(pp)$ are caused by inelasticities, while the size of $\Delta\sigma_L(pp)$ is due principally to the elastic component.

conserved. The Pauli exclusion principle holds that the total wave function describing the scattering of nucleons (fermions) is to be anti-symmetric, or

$$(-1)^{S+L+I} = -1.$$

In the case of pp scattering ($I = 1$), the total wave function is anti-symmetric only if the protons scatter in a singlet state with L even, or in a triplet state with L odd. The partial waves are denoted by $^{2S+1}L_J$, and the partial waves involved in pp scattering are, with the subscript J left off for clarity,

$$^1S, ^3P, ^1D, ^3F, ^1G, ^3H, \dots$$

Values of J range from $L - S$ to $L + S$ in integer steps.

For the case of np scattering, which occurs in a mixture of isospin $I = 0$ and $I = 1$ states, there are no anti-symmetry requirements. Therefore, one can describe np scattering with the extended set of partial waves

$$^1S, ^3S, ^1P, ^3P, ^1D, ^3D, ^1F, ^3F, ^1G, ^3G, \dots$$

The greater number of partial waves in np scattering reflects the extra degree of freedom in the np system because the nucleons can scatter in one of two isospin states.

The helicity amplitudes can be decomposed into spin partial waves^[1,32]

$$\begin{aligned}
\phi_1 &= \frac{1}{2ik} \sum_J \left\{ (2J+1)R_J + JR_{J-1,J} + (J+1)R_{J+1,J} + 2[J(J+1)]^{1/2}R^J \right\} d_{00}^J(\theta) \\
\phi_2 &= \frac{1}{2ik} \sum_J \left\{ -(2J+1)R_J + JR_{J-1,J} + (J+1)R_{J+1,J} + 2[J(J+1)]^{1/2}R^J \right\} d_{00}^J(\theta) \\
\phi_3 &= \frac{1}{2ik} \sum_J \left\{ (J+1)R_{J-1,J} + JR_{J+1,J} - 2[J(J+1)]^{1/2}R^J + (2J+1)R_{JJ} \right\} d_{11}^J(\theta) \\
\phi_4 &= \frac{1}{2ik} \sum_J \left\{ (J+1)R_{J-1,J} + JR_{J+1,J} - 2[J(J+1)]^{1/2}R^J - (2J+1)R_{JJ} \right\} d_{-11}^J(\theta) \\
\phi_5 &= \frac{1}{2ik} \sum_J \left\{ [J(J+1)]^{1/2}(R_{J-1,J} - R_{J+1,J}) + R^J \right\} d_{10}^J(\theta).
\end{aligned}$$

The $d_{\lambda\mu}^J$'s are reduced spin rotation matrices, parameterized by the total angular momentum J , and the total helicity for the initial (λ) and final (μ) states. At a scattering angle of 0° , the $d_{\lambda\mu}^J$'s are as follows:

$$d_{00}^J(0) = d_{11}^J(0) = 1$$

$$d_{-11}^J(0) = d_{10}^J(0) = 0.$$

The R_J amplitudes are the spin singlet partial waves (e.g., $^1S_0, ^1P_1, ^1D_2, \dots$) for $J = L$, and the R_{JJ} amplitudes are uncoupled spin triplet partial waves (eg. $^3P_1, ^3D_2, ^3F_3, \dots$) for $J = L$. Because the orbital angular momentum L is not a good quantum number in NN scattering, $L = J \pm 1$ triplet transitions also occur. The $R_{J\pm 1,J}$ and R^J amplitudes are coupled triplet partial waves (e.g., $^3S_1, \epsilon_1, ^3D_1, \dots$). The mixing terms ϵ describe mixed triplet transitions, such as $\epsilon_1 = ^3S_1 - ^3D_1$ for isospin $I = 0$ and $\epsilon_2 = ^3P_2 - ^3F_2$ for isospin $I = 1$.

One possible representation for the singlet and triplet partial wave amplitudes is

$$R_J = \cos \rho_J \exp(2i\delta_J) - 1$$

$$R_{JJ} = \cos \rho_{JJ} \exp(2i\delta_{JJ}) - 1$$

$$R_{J\pm 1,J} = \cos \rho_{J\pm 1,J} \cos 2\epsilon_J \exp(2i\delta_{J\pm 1,J}) - 1$$

$$R^J = i \sin 2\epsilon_J \exp(i(\delta_{J-1,J} + \delta_{J+1,J} + \alpha_J)).$$

The ρ 's are the inelasticities and the ϵ 's and α are mixing parameters for triplet states.

The δ 's are the phase shifts for the partial waves in question.

The cross sections σ^{tot} and $\Delta\sigma_L^{\text{tot}}$ can be expressed in terms of partial waves using the partial wave decompositions for the helicity amplitudes $\phi_1(0)$ and $\phi_3(0)$:

$$\sigma^{\text{tot}} = \frac{\pi}{k^2} \Im m \left(\frac{1}{i} (2J+1) [R_J + R_{J-1,J} + R_{J+1,J} + R_{JJ}] \right)$$

$$\sigma_L^{\text{tot}} = \frac{2\pi}{k^2} \Im m \left(\frac{1}{i} \left[(2J+1)(R_J - R_{JJ}) + (R_{J+1,J} - R_{J-1,J}) + 4[J(J+1)]^{1/2} R^J \right] \right).$$

No part of the partial wave decomposition for σ^{tot} can make a negative contribution because $\sigma^{\text{tot}} > 0$ always ($\Im m(R/i) > 0$). In $\Delta\sigma_L^{\text{tot}}$ the singlet and coupled triplet partial waves R_J , $R_{J+1,J}$ and R^J make positive contributions (can give a peak), and the uncoupled and coupled triplet partial waves R_{JJ} and $R_{J-1,J}$ make negative contributions (can give a dip). This is true for both isospin states $I = 0$ and $I = 1$.

If the imaginary and real parts of each partial wave are plotted for each energy in a complex plane, an Argand plot results. Any partial wave showing rapid change in its

phase shift δ by $\Delta\delta \approx \pi$ on a smooth background is considered to show resonance-like behavior.⁽³³⁾ Two candidates for dibaryon resonances in the pp system seem to occur in the partial waves 1D_2 and 3F_3 (Fig. 1.7).

Figure 1.7

Argand plots for the 1D_2 and 3F_3 partial waves. The axes $\text{Re}T$ and $\text{Im}T$ refer to the real and imaginary parts of the partial wave amplitude (1D_2 (a) and 3F_3 (b)), the ellipses indicate the error for each point plotted and the energy solutions used. Note that both these partial waves are highly inelastic ($\eta \approx 0.2$), with counter-clockwise loops indicating possible resonances. The 3F_3 partial wave exhibits a possible resonance on a background (loop located left of the $\text{Re}T$ axis). Taken from Ref. [4]

CHAPTER 2 THE ACCELERATOR

In this chapter, those portions of the experiment which are related to the LAMPF accelerator, or are generally permanent features of the experimental area used, are described.

2.1 The Lamb-Shift Polarized Ion Source

The LAMPF accelerator has the unique ability to accelerate two different ion beams, of opposite charge, simultaneously along the same linac. The beams are then split apart at the switch yard and delivered to various experimental areas. This experiment used the P^- beam, which is a polarized H^- atomic beam. Other beams available at LAMPF include unpolarized H^- and H^+ . P^- and H^+ are run with the same tune, which is usually optimized for H^+ .

Within the injector building at LAMPF there are three Cockcroft-Walton preaccelerators which contain ion sources for the beams. The preaccelerators operate at 750 KeV using voltage-multiplier circuitry. The voltage-multiplier circuits use solid-state rectifiers energized by 5 Khz transformers. A stack of capacitors are charged in

parallel and discharged in series, achieving a voltage multiplication. The preaccelerators are enclosed in smooth aluminum housings and are supported by insulating columns to minimize sparking and corona. Ions from these sources are extracted by potential differences and are pre-accelerated into the drift-tube linac, the second stage in the beams' journey to 800 MeV.

The polarized proton beam for the np experiment was generated using the LAMPF Lamb-shift H^- ion source,^[34] where the physics of hydrogen in metastable states is used to select ions in a particular spin state. This results in a beam with a high degree of polarization. The several processes which condition the proton beam are as follows:

- produce a proton beam.
- convert the proton beam into a neutral metastable atomic beam by adding an electron.
- apply some selective process which polarizes the metastable atomic beam by bringing those atoms with undesirable spin states to their ground state (quenching).
- convert the beam to negative ions by adding another electron (H^-) so that this polarized ion beam can be accelerated simultaneously with H^+ .

A block diagram of these processes is shown in Fig. 2.1. Each of these steps is discussed in this section.

The protons produced in the source have a kinetic energy of 500 eV and are di-

**LAMPF
LAMB-SHIFT
POLARIZED ION SOURCE
SCHEME**

Figure 2.1

Sequence of steps used to polarize a proton beam, for the LAMPF Lamb-shift source. The protons passed through a cesium-vapor cell, where an electron was captured, making an atomic beam of neutral hydrogen in an admixture of 2s and 1s states (the ground state, 1s, outnumbered the 2s state by 40 to 1). The spin filter polarized most of the 2s atoms and quenched the rest to the ground state. The beam next passed through an argon gas cell where negative ionization favored the 2s state over the 1s state by 400 to 1, and so most of the beam entering the accelerator consisted of polarized H^- .

rected through a cesium vapor cell where the protons capture an electron through the process of charge exchange. An atomic beam of neutral hydrogen, H^0 , then exits the cesium cell in a mixture of 1s and 2s states, where the population of 1s states, being of lower energy, is naturally greater than that of the 2s states. In the presence of an external magnetic field, the 2s state splits into four separate states going with various combinations of orientations of the electron spin and the nuclear spin. The 2s states are metastable because the transition $2s \rightarrow 1s$ is not allowed via one photon exchange, as there is no orbital angular momentum change. Notice in Fig. 2.2, however, that the lower 2s states crosses the higher 2p states in a magnetic field strength of around 575 Gauss. The 2p states have a very short lifetime compared with the metastable 2s state (1.6 nsec compared with 0.14 sec), and if we can connect the lower $2s_{1/2}$ states with the upper $2p_{1/2}$ states through some process, we can quench the lower $2s_{1/2}$ states without disturbing the upper $2s_{1/2}$ states. In practice this could be done by applying a transverse electric field of about 10 V/cm in conjunction with a longitudinal magnetic field of 575 Gauss, but this action leaves behind the upper two spin states. Therefore, a different approach is taken at LAMPF.

In this procedure, a nuclear spin filter is used which involves the three-level interaction discovered by Lamb and Retherford,^[35] in which all but one spin state is quenched. The filter uses a longitudinal rf field together with a transverse dc field and longitudinal magnetic field to accomplish this. By convention, the upper $2s_{1/2}$ states are denoted

HYDROGEN BREIT-RABI DIAGRAM.

Figure 2.2

Splitting of the $n=2$ states of the neutral hydrogen atom in a magnetic field.

as ' α ' states, the lower $2s_{1/2}$ states are ' β ' states, the upper $2p_{1/2}$ states are ' e ' states and the lower $2p_{1/2}$ states are ' f ' states. Each of these states has a degeneracy of two, which is removed in the presence of magnetic fields due to the coupling of electronic spins to nuclear spins (hyperfine structure). In the presence of a 540 Gauss longitudinal magnetic field, the metastable $2s$ state splits so that the β state ($2s_{1/2}, m_j = -1/2$) is nearly degenerate with the e state ($2p_{1/2}, m_j = +1/2$). A longitudinal rf electric field of strength 0.07 V/m and frequency 1608 MHz couples the α ($2s_{1/2}, m_j = 1/2$) and the e state, while at the same time a transverse DC electric field of strength 0.20 V/m couples the e state to the β state. The rf frequency is tuned to the $\alpha - \beta$ energy difference set up by the magnetic field, and so an equilibrium between the α and β mixtures is established.

Selection rules require that $\Delta m_j = 0$, and metastable atoms with other m_j values are quickly quenched by the fields set up in the spin filter. Most of the β state atoms are quenched through collisional processes or by the transverse electric field. The remaining long-lived β component is transformed into the α state through an adiabatic process; the rf field is adiabatically reduced along the z -axis, and then the magnetic field is reduced. The atomic hydrogen atoms exiting the spin filter are in a mixture of α states with spin $m_j = 1/2$ and ground state atoms.

The atomic beam then passes through an argon gas cell with a magnetic field of 6 Gauss, used to maintain the orientation of the spins. Here an extra electron is added

to the hydrogen atoms by charge exchange. The ground state atoms outnumber the 2s state atoms by 40 to 1, but the cross section for negative ionization in the argon cell favors the 2s state by approximately 400 to 1. The net result is that nearly all the ions exiting the argon cell are in the 2s state and are highly polarized (about 80%). These ions are then injected into the drift-tube linac for acceleration.

The LAMPF Lamb-shift source possesses the ability to measure the beam polarization to a high degree of accuracy because of the atomic physics of the source. Normally, the P^- beam exits the argon cell with nearly all ions in the 2s state with a high degree of polarization, but there is also some unpolarized background from the ionization of the ground-state atoms present. If somehow the 2s ions could be blocked, or quenched, one could obtain a measurement of this unpolarized background and calculate the beam polarization. This is actually done by periodically placing the source into a "Quench mode," in which the spin-filter is de-tuned so that the 2s ions are quenched, and then the cross section is down by a factor of 400 for ionization for all the atoms.

At LAMPF, the polarized/quench cycle operates with a 60-second period, asynchronously with the Normal/Reverse spin cycle of 130-second duration.^[36] The P^- beam would be completely polarized (100%) if not for the unpolarized background (I_q) which contributes to the total beam current. If the total beam current is I , then the polarized ions are represented by a beam current $I - I_q$. The ratio of beam intensities during which the spin-filter is tuned ($I - I_q$) and de-tuned (I_q) is then the polarization

of the beam,

$$P = \frac{(I - I_q)}{I}.$$

A correction term is added to this relation because the quenched beam typically has a small polarization ($P_q = -0.03$) opposite in direction to the P^- beam. A weighted average then allows one to obtain the true beam polarization

$$P = \frac{(I - I_q)}{I} + P_q \frac{I_q}{I}.$$

The quench period is normally 10 seconds in duration and run after every spin-flip and it offers a redundant check of the polarization measured by other means (beamline polarimeters, for example). The technique is straightforward. One only needs some arrangement of detectors and a clock to measure the beam intensity. The spin-filter cycles, "Normal," "Reverse," and "Quench", are available as logic levels to experimenters for use in data acquisition.

2.2 The Drift-Tube Linac

The first stage of the accelerator is the drift-tube linac, which accelerates the beams from 750 KeV to 100 Mev.^[37] Based upon the geometry of the Alvarez design,^[38] it uses a series of four cylindrical vacuum tanks which are one meter in diameter with the drift tubes suspended inside. Several resonant stems with eccentric nose pieces project from the sides of the tank towards the suspended drift tubes. These are known as

"post couplers" and are used for tuning and stabilizing the electric fields around the drift tubes. By inserting these post couplers to a chosen depth and rotating them, the capacitance and inductance of individual cavities may be adjusted and the system can be tuned to produce a resonant effect where little power is dissipated. This improves the stability of the drift-tube linac considerably.

Beam particles are accelerated by establishing a longitudinal alternating electric field within the tanks, which operate at a frequency of 201.25 Mhz. In order for H^+ and H^- (or P^-) particles to be accelerated simultaneously, the H^- particles must be injected into the drift-tube linac one-half rf cycle later than the H^+ particles. The lengths of successive drift tubes are increasingly greater, as required by the steadily increasing velocities of the beam particles. As the protons accelerate across the gaps of the drift tubes and coast through, the electric field changes direction just as the H^- particles coast into the gaps where they too are accelerated. By this time the protons are well on their way towards the next gap, but since they are within the drift tubes at this time they are not affected by the external electric fields which accelerate the H^- particles. This process continues until the beam particles emerge from the drift-tube linac, at which times they have final kinetic energies of 100 MeV. The drift-tube linac is about 62 meters long.

2.3 The Side-Coupled Linac

The beam particles which exit the drift-tube linac are injected into the side-coupled linac, which is the third and last stage of the LAMPF accelerator. Here the beam is accelerated to a final energy of 800 MeV and delivered to various experimental areas through the beam switchyard.

The unique feature of the LAMPF linac is the use of a series of external resonant cavities which are coupled to adjacent pairs of resonant cavities between drift tubes, a scheme invented by D. Nagel. The system operates at an rf frequency of 805 Mhz and accelerates beam particles in a similar manner to the drift-tube linac. A different design was used for this stage of the accelerator because the drift-tube design becomes less efficient at high energies and would dissipate a great deal of power in continuous operation at 800 MeV. This is because at high energies the drift tubes would be unreasonably long, resulting in a low average longitudinal accelerating gradient. While the drift-tube lengths could be shortened if the rf frequency were increased, this leads to the following problems:^[39]

- At high rf frequencies the acceleration gaps conduct large displacement currents, which load the rf generator.
- Adjacent tubes would act like dipole antennae at high rf frequencies and, therefore, radiate part of the rf energy instead of using that energy for accelerating.

A cut-away view of a segment of the side coupled linac is shown in Fig. 2.3, where an illustration of the location of the side-coupled resonant cavities with respect to the beam axis is also shown. The linac operates in the $\pi/2$ mode because energy transport between cavities is at a maximum for this mode, which means that it is the best choice for rf power coupling. However, $\pi/2$ modes have relatively low electric field gradients since half of the cavities are un-excited. The design at LAMPF overcame this limitation by displacing the unexcited cavities to the side and passing the particle beam through the excited cavities. Low-level electromagnetic oscillations in the side cavities act to transfer energy between excited cavities along the system. In this design, very little power is dissipated by the side cavities. The coupling system is designed to provide the necessary phase shifts and small power losses required in accelerating a particle beam to energies of 800 MeV in a relatively short distance (the linac is half a mile long). The side-couple cavities were then tuned by applying small physical distortions.

As in the case of the drift-tube linac, the physical spacing between the centers of the resonant cavities increases with increasing beam velocity, so each of the nearly five-thousand cavities which make up the linac have different lengths. The linac is divided into 352 sections, with various beam-handling devices mounted between sections, such as focusing magnets and sensors. When energies lower than 800 MeV are desired, some of the sections are turned off, lowering the beam energy by steps of 17 MeV. In practice the actual energies available to users are not so finely stepped.

Figure 2.3

Cut-away view of a segment of the side-coupled cavity structure.

2.4 The Line-B Beamline

After the accelerated protons have reached their final energies, they enter the beam switchyard where the P^- beam is directed into Line B. A diagram of the beam transport system used for this experiment is shown in Fig. 2.4. All beamline transport bends are in the horizontal plane. During the course of the experiment, which ran in 1987 and 1988, there were some changes made in the beamline. These will be pointed out in the following discussion. The primary objective of the beamline is to transport the polarized beam from the switchyard to the liquid deuterium target (LD_2). Polarized neutrons are produced by the charge-exchange reaction $\bar{p} + d = \bar{n} + X$ at the LD_2 target, collimated, and directed into the experimental area.

A new feedback system for steering the beam was designed for this experiment, which is discussed in the next chapter. Using components in the beamline, the beam position on the LD_2 target was continuously monitored and the beam re-steered, if necessary, to maintain a set position. This servo-steering was used with great success. In this section, several components of the beamline are discussed according to their importance in the the experiment.

2.4.1 The Spin Precession Solenoid

Upon entry into Line B, the beam passes through a stripper which removes the electrons from the atoms forming the polarized P^- beam, resulting in a polarized proton

Figure 2.4

The beamline used for the experiment. The polarized proton beam entered the beamline from the lower left in the figure, and the neutron beam was produced at the liquid deuterium target (at the top of the figure). The non-interacting proton beam was routed to a beam dump. Figure is taken from the LAMPF Users Handbook.

beam. Following this, the protons then pass through a set of bending magnets, LBBM2 and LBBM3, and pass through the Line B solenoid LBSOL. This is used to precess the spin of the protons in the beam so that the spin of the protons at the LD₂ target is nearly longitudinal in direction, or L-type, which is essential for the production of polarized neutrons. The spin direction of the beam upon entry to Line B was mostly N-type. Since all transport bends occur in the horizontal plane, the beam cannot be made L-type unless the spin is rotated from the vertical direction (N) to a sideways orientation (S). The N-type spin is rotated about the beam axis by a solenoid magnet LBSOL to give S-type spin, after which all subsequent bends in the beamline precess the spin from S to L at the LD₂ target. Occasionally, S-type beam was introduced into the beamline, during which times LBSOL was turned off. The magnetic field, and hence the current, required by LBSOL to rotate from N-type to S-type is directly proportional to the momentum of the beam. This follows from the consideration of spin precession for moving particles passing through a magnetic field.

2.4.2 The Steering Magnets

Upon exiting LBSOL, the beam was collimated and passed through the bending magnet LBBM4, where it was directed through a set of three quadrupole magnets. These formed a quad triplet, in which the first and third quadrupole magnets had the same polarity and field strength and were opposite in polarity to the middle quadrupole

magnet, which had twice this field strength. The quadrupoles, LBQM1, LBQM2 and LBQM3, had been set further apart during the first year of data taking, where they operated with roughly equal currents, and they were moved closer together during the second year of running. The change was made because the beam could be more efficiently tuned in the new configuration. The beam next passed through a set of horizontal and vertical steering magnets, LBSM4X and LBSM4Y, which were used in the beam servo-steering system discussed in more detail in the next chapter.

2.4.3 The Polarimeters

The beam was next directed through the Line B polarimeter, LBPOL1, which monitored the transverse components of the proton beam polarization. A new polarimeter, LBPOL2, was added to the beamline for the experiment and was located further downstream from LBPOL1 and just upstream of the LD₂ target. LBPOL2 monitors the transverse components of the beam polarization at the LD₂ target, as there were no further bends until after the LD₂ target. LBPOL2 was functionally identical to LBPOL1.

A layout of the beamline polarimeters is given in Fig. 2.5, along with a plot of the *pp* analyzing power over the range of LAMPF energies (taken from Ref. [40]). The average analyzing power is about 0.5 in this energy range. In these polarimeters, a group of left/right (L/R) and up/down (U/D) counters are used to detect the scattered

Figure 2.5
Layout of the Line B polarimeters LBPOL1 and LBPOL2, and the CH₂ analyzing power for a laboratory scattering angle of 17° as a function of laboratory energy.

and recoil protons from a CH₂ target block. Each detector (L, R, U, D) is actually a coincidence between counters set at angles which are conjugate to each other (set to detect scattered and recoil protons simultaneously). These angles were chosen to be 17° and 67° lab, corresponding to *pp* elastic scattering, and were chosen for the following reasons:^[41]

- Near 800 MeV, the analyzing power for elastic *pp* scattering reaches a maximum.
- As a function of scattering angle, the analyzing power exhibits a broad maximum near 17° lab, making the polarimeter insensitive to errors in angle settings.
- The position of this maximum remains essentially unchanged over the range of LAMPF energies, minimizing the need for changes in detector positions with each energy change.
- The detection of both scattered and recoil protons in the final state gives a clean signal for elastic scattering. There are very few inelastic events detected in coincidence at these angles.

The protons elastically scattered from the CH₂ target are detected in each of four directions; left, right, up and down. By convention, the coincidence of the detectors at 17° left and 67° right is called Left, with the other directions defined in a similar manner. The left/right scattering plane analyzes the vertical component of beam polarization,

$$P_N = \frac{1}{A} \frac{(L - R)}{(L + R)}$$

while the up/down scattering plane analyzes the horizontal component of beam polarization,

$$P_S = \frac{1}{A} \frac{(D - U)}{(D + U)}.$$

It is impossible to measure the longitudinal component of beam polarization directly using this configuration, which is one of the problems of using an L-type beam.

In practice, the signals L, R, U and D are defined in terms of the geometrical means

$$L = (L_+ R_-)^{\frac{1}{2}}$$

$$R = (R_+ L_-)^{\frac{1}{2}}$$

$$U = (U_+ D_-)^{\frac{1}{2}}$$

$$D = (D_+ U_-)^{\frac{1}{2}}$$

because nearly all experimental uncertainties and instrumental asymmetries cancel to high order in such products.^[42] Physically, a beam proton polarized spin up and scattered left is equivalent to one polarized down and scattered right (and similarly for the other directions). The asymmetries are typically measured to a precision of ± 0.002 , with the error determined by counting statistics.^[43] It should be noted that events are not only detected from the elastic pp scattering, but also from quasi-elastic scattering of the beam protons from the carbon nucleus. These quasi-elastic scatters are about equal in number to the elastic scatters but have poor angular correlation. As a result, only about 5% of the total events are due to scatters off of the carbon. The effective

analyzing power for the polarimeters is then a few percent lower than for proton-proton scattering.

The beamline polarimeters were used to monitor the orientation of beam spin to insure that LBSOL was set correctly. If it was set correctly, the left/right asymmetries would be zero at both polarimeters, and the up/down asymmetry at LBPOL2 would also be zero, demonstrating L-type beam at the LD₂ target. The polarization of the proton beam was obtained using the quench method which was discussed at the end of Section 2.1 in this chapter.

2.4.4 Beam Monitors

In the beamline between the Line B polarimeters LBPOL1 and LBPOL2 there were an ion chamber, a wire scanner and a bending magnet. As a rule, the wire scanner was put into the beam only when beam tuning was in progress or when the quality of the beam tune was to be monitored which was done at regular intervals.

The bending magnet LBBM5 was the last bending magnet before the LD₂ target and was located just upstream of LBPOL2. With the Line B solenoid set correctly, the beam spin is very nearly L-type (between 2°–16° over the energy region 500–800 MeV) at LBPOL2. After passing through LBPOL2, the beam goes through a collimator in a steel wall and passes through the proton beam monitor used in this experiment. This was a secondary emission monitor (SEM).

The SEM monitor, or LBER, was chosen as the data normalizer in this experiment because neutrons are difficult to detect without seriously affecting the quality of the neutron beam. It was enclosed within a box located just upstream of the LD₂ target. Before this experiment the SEM was located further upstream between LBPOL1 and LBBM5. The SEM was moved to this location for this experiment in order to monitor the beam current at the LD₂ because some beam could be lost in transport to the LD₂ if the beam scraped against any obstacles in the beamline.

The SEM monitors the current of the proton beam by using the phenomenon of secondary electron emission. The passage of energetic charged particles through certain materials imparts sufficient energies to bound electrons to liberate them from their atomic bonds so that they escape from the surface of the material (secondary emission). The freed electrons are low in energy and originate from the thin region within 0.01 micrometers of the surface. The loss of these electrons induces a net charge on the surface, and a measurement of this charge provides an accurate relative determination of the total number of charged particles which passed through the material.

The SEM was constructed with one thin emitter foil located between two thin collector foils, enclosed within a box and operated within a vacuum. The emitter foil was maintained at a constant potential, typically -300 volts relative to the collector foils, and in this bias voltage those electrons liberated from the emitter are attracted to the collector foils. The charge on the collectors was integrated by a "RED NUN" RN-8111

charge digitizer, the output of which was a frequency proportional to beam current. By simultaneously scaling both the digitizer counts and pulses of a suitable reference clock, a measurement of the beam current was obtained.

The beam profile monitor LBII6 was used in the computerized beam feedback system, or servo-steering, to constantly detect and correct steering mis-alignments. This monitor was located in a box placed between the SEM and the LD₂ target. This box also contained the wire scanner LBWS6, which was enabled periodically in order to monitor the beam tune at the LD₂ target. Only one of LBII6 and LBWS6 could be in the beam at a time.

LBII6 was similar to a multi-wire proportional chamber. In the simplest proportional chamber, the detector is constructed of a plane of anode sense wires, with each wire parallel to neighboring wires, and with a cathode plate on each side of the anode plane. In the space between the two cathode plates the electric field lines are parallel and constant, except in the regions near the anode wires, where the electric field takes on a $1/r$ dependence. Any ions produced by the passage of energetic particles drifts along the electric field lines where they are quickly accelerated in the high field region near the wires and produce an avalanche. The ions produced by this multiplication process induce a signal on an anode sense wire, with the size depending on the gas mixture and anode voltage. By this means the profile of the beam can be obtained.

Perpendicular chambers allow for both x and y beam profiles. The resulting signals

were read out using a multiplexer, allowing the x and y beam profiles to be monitored by a dual trace oscilloscope as well as by computer. A typical beam profile from LBII6 is shown in Fig. 2.6, taken from an arbitrary run. The operation of the beam-steering feedback system which uses LBII6 is discussed in detail in the next chapter.

2.5 The LD₂ Target System

A good description of the LD₂ target can be found in the thesis of C. W. Bjork.^[44] The polarized neutrons are produced by the charge-exchange reaction $d(\vec{p}, \vec{n})X$, which produces a neutron spectrum at 0° characterized by a sharp peak located just below the incident proton momentum and a well-separated inelastic background. Fig. 2.7 shows the momentum spectra for neutrons produced at the LD₂ target by incident protons having kinetic energies of 647 and 800 MeV. Roughly, at all incident proton momenta above pion production, about 40% of the neutron beam lies in the peak region. The polarization of the neutron beam is found by measuring the observable called the longitudinal polarization spin transfer, or K_{LL} , which represents the amount of transfer of incident-beam polarization to the final-state beam. Experiments demonstrate that the polarization transfer in the above charge exchange reaction is highest for longitudinal spins.^[45,46] At zero degrees, there are only two spatial directions with any physical meaning; the longitudinal direction L and the normal direction N (in the absence of any scattering plane, N and S are equivalent). The spin-transfer function K_{LL} is larger than

Figure 2.6
A typical beam profile taken from LBII6 for an arbitrary run. The servo-steering system calculated x and y centroids for steering corrections.

Figure 2.7
Momentum spectra of neutrons exiting the LD₂ target for two incident proton beam energies. The peak neutrons resulted both from elastic and quasi-elastic scattering, while the broader peak neutrons resulted from inelastic scattering. Approximately 40% of the neutron beam lies within the elastic/quasi-elastic peak.

the quantity K_{NN} , as shown in Fig. 2.8. For longitudinal orientation, the polarization of the neutron beam is about 60% that of the proton beam, and it is polarized oppositely to the protons.

The liquid deuterium is located within a target cell situated within a vacuum. The entrance and exit windows for the LD_2 target are 3.2 cm wide and 15.24 cm long, and they are separated by a distance of 25.4 cm from each other. The windows are bowed outwards because there is vacuum outside the target cell. Therefore, depending on where the proton beam hits the entrance and exit windows, the number of neutrons produced will vary because the proton beam passes through various thicknesses of target material. Since the neutron beam was not used for the normalization of the data, a stable proton beam was essential for this experiment.

A Monte-Carlo study⁽⁴⁷⁾ of neutron production as a function of the path length of the proton beam through the LD_2 target, as well as of the beam spot size, illustrates the impact on the experiment of the current target configuration (the surface of the entrance and exit windows were assumed to be spherical and bowed outwards by 0.25 in). Fig. 2.9 shows some results of this study. Inputs for the study included known variations in the beam spot size and position at the LD_2 target windows. The model was also extended to include neutron transport through the collimator and on to the polarized target, which are also geometrical effects. It was concluded that typical systematic changes in the beam spot at the LD_2 target would effectively "kill" the experiment unless the

Figure 2.8

Plot of the spin-transfer functions K_{LL} and K_{NN} at LAMPF energies. Note that K_{NN} is smaller than K_{LL} , so to obtain a highly polarized neutron beam, it is desirable that the proton beam be nearly L-type at the liquid deuterium target. Data from Refs. [45,46].

Figure 2.9

Results of a beam-spot Monte-Carlo calculation at the liquid deuterium target. The number of neutrons produced at the liquid deuterium target (LD₂), as compared to the number of neutrons produced by a beam of zero width, is shown in (a) as a function of beam spot position and width at the LD₂ target. The fraction of these neutrons which then travel through the collimator and pass through the proton target downstream is shown in (b), which shows the sensitivity of the experiment to beam spot motion and width changes at the LD₂ target.

variations in the beam spot size and position could be fixed to less than 1 mm. Short term changes that average out in the long term should not affect the experiment.

These studies, together with experience with previous experiments in this beamline, led to the conclusion that control over the beam spot size and beam motion had to be instituted to prevent the experiment from being sensitive to the "false" asymmetries that would be generated. There was no way to control the size of the beam at the LD₂ target, but beam motion was controlled by a servo-steering beam feedback system. This is discussed in the next chapter.

A bending magnet (LBBM6) located downstream of the LD₂ target deflected the non-interacting proton beam into the beam dump while the neutrons continued on and passed through a collimator leading into the experimental area. Since the neutrons also passed through the magnetic field of LBBM6 (see Fig. 2.4), the neutron spins were rotated away from L-type.

CHAPTER 3 THE 10 MHZ BUNCHER AND THE SERVO-STEERING SYSTEM

3.1 The 10 Mhz Buncher

Normally, the LAMPF P^- beam travels along the accelerator in beam pulses which have a microstructure of 5 ns. That is, a detector set up to monitor the beam could expect to detect protons at times which are multiples of 5 ns. The flight path of the neutron beam in the experimental area is on the order of only 10 ns. The peak-energy neutrons would also have a microstructure of 5 ns, but one can easily envision a case where slow neutrons from previous beam pulses are mistaken for fast neutrons. As discussed in this section, this creates difficulties for one-armed experiments in extracting time-of-flight information.

3.1.1 Why a Beam Buncher was Required

A beam buncher was needed for this experiment because of the nature of the neutron beam used. As discussed in the last chapter, the neutron beam is produced by the

charge-exchange reaction $d(\vec{p}, \vec{n})X$, and at zero degrees this reaction produces a neutron spectrum with a sharp energy peak just under the incident proton energy, with a well-separated inelastic background. The intensity of the background is comparable to that of the peak. These low-energy neutrons make the experiment difficult. One must somehow discriminate against them in order to obtain a clean measure of the total cross section $\Delta\sigma_L$ at the peak energy. This required a "beam buncher."

The LAMPF accelerator uses a master timing frequency of 201.25 Mhz which is available to all experimenters for obtaining time-of-flight information during experiments. This signal is the basis of a "STOP" signal for a time-to-digital converter (TDC), with the detection of a neutron as a "START" signal. This time-of-flight information becomes ambiguous, however, since the neutron beam is made up of "fast" neutrons (peak energy) and "slow" neutrons. Consider a detector downstream of the LD₂ target which is used to count neutrons, but has no way of discriminating between fast and slow neutrons. If the distance between the detector and the LD₂ target is large, one has the case where slow neutrons originating at an earlier time overlap with fast neutrons originating at a later time and are detected at the same time.

The proton beam pulses which are incident upon the LD₂ target come every 5 ns, and the fast neutrons produced also have a 5-ns structure. The fast neutrons are detected at discrete times given by

$$t_{fast} = \frac{d}{\beta_0 c} + n\tau$$

where d is the distance between the LD₂ target and the detector, β_0 is the relative velocity of the fast neutrons and τ is the rf period or microscopic beam period. Similarly, the times at which slow neutrons are detected are

$$t_{slow} = \frac{d}{\beta c} + m\tau.$$

The integers n and m point out that we must consider which beam pulse the neutrons are to be associated with, and generally n and m are not equal. The allowed values of β , for which slow neutrons are indistinguishable from fast neutrons, is derived by equating t_{fast} and t_{slow} ,

$$\begin{aligned} t_{fast} &= t_{slow} \\ \frac{d}{\beta_0 c} + n\tau &= \frac{d}{\beta c} + m\tau \\ \text{or, } \frac{qc\tau}{d} &= \left\{ \frac{1}{\beta} - \frac{1}{\beta_0} \right\} \text{ where } q = n - m > 0. \end{aligned}$$

This can be inverted to solve for β , which is discreet and depends upon the integer q ;

$$\beta_q = \beta_0 \left\{ \frac{1}{1 + qc\tau\beta_0/d} \right\}.$$

Since we restrict q to non-negative integers, β_q is always less than β_0 for the fast neutrons.

If we now model the detector as having infinite timing resolution, or equivalently, infinite momentum resolution, we can now solve the problem of how many slow neutrons are mistaken for fast neutrons. Let $\mathcal{P}(\beta)$ represent the momentum spectrum of the

neutrons produced by an incident proton beam on the LD₂ target and let us use the delta function to model the response of the detector. Then the total number of neutrons is proportional to

$$\begin{aligned} N &\propto \sum_{q=0} \int d\beta \mathcal{P}(\beta) \delta(\beta - \beta_q) \\ &\propto \mathcal{P}(\beta_0) + \sum_{q=1} \mathcal{P}(\beta_q) \\ &\propto \mathcal{F} + \mathcal{S} \end{aligned}$$

where $\mathcal{F}$ and $\mathcal{S}$ represent the total amount of fast and slow neutrons, respectively. Using a proton beam of 800 MeV and a distance between the LD₂ target and detector of about 20 meter flight path, we find that the parameter $c\tau\beta_0/d$ has the value 0.063 (using $\tau = 5$ ns). From Fig. 3.1 both $\mathcal{F}$ and $\mathcal{S}$ may be estimated, and we find that the fraction of the total neutrons detected which are slow is approximately given by

$$f = \frac{\mathcal{S}}{\mathcal{F} + \mathcal{S}} \approx 0.22.$$

The relation $d^2\sigma/d\Omega d\beta = mc(1 - \beta^2)^{-3/2} d^2\sigma/d\Omega dp$ was used to transform from a momentum representation to one in β . This situation is known as “wrap-around”, whereby the detection of slow neutrons can coincide with the detection of fast neutrons. The buncher, which is discussed in the following section, extends the parameter τ from 5 ns to 100 ns. With beam pulses separated in time by 100 ns instead of 5 ns, the parameter $c\tau\beta_0/d$ becomes 1.26, and as a result the number of slow neutrons detected falls to

Figure 3.1

Plot showing how slow neutrons are mistaken for fast neutrons in the "wrap-around" effect. The normal rf time structure of the P^- beam at LAMPF was 5 ns/beam pulse; at 800 MeV those neutrons having momenta indicated in the figure would have time of flights consistent with fast neutrons (flight path of 20 meters). The detection of slow neutrons was essentially eliminated by stretching the rf time structure of the P^- from 5 ns to 100 ns, by using a beam buncher.

nearly 0 (from Fig. 3.1, the next wrap-around energy is less than 400 MeV/c), a very desirable result.

3.1.2 How the Buncher Works

The 10-MHz bunching system consists of two separate parts and is driven at a frequency of 10.0625 MHz. In this section we discuss how each of the components works with the current Lamb-shift source. Fig 3.2 shows where each of the buncher components are located inside and downstream from the P^- pre-accelerator. These components are the "first-stage" buncher inside the dome and the "low-frequency" buncher just upstream of the drift-tube linac. The purpose of the buncher is to increase the intensity of the individual P^- beam pulses and also to increase the time between the pulses so time-of-flight experiments can be done. It concentrates most of the beam current that is spread over 20 rf cycles into one rf cycle and chops that small fraction of current that is not bunched.

The first-stage buncher is basically an added section of insulated beam tube with a diameter matching nearby electro-optical elements (2.54 cm). The beam tube is separated from neighboring elements with about 1-cm gaps. The buncher was designed to operate at a frequency of 10.0625 MHz which is 1/20 of the normal 201.25-MHz operating frequency of the drift-tube linac. With respect to the P^- dome, the DC and AC voltages acting on the buncher tube are +4.00 KVDC and 400 – 800 VAC at

Figure 3.2
 Location of buncher components within the P- dome (not to scale). For maximum beam bunching, the lengths of both buncher tubes were matched to the relative velocity of the beam at each tube: 3.81 cm and 58.58 cm, respectively.

10.0625 Mhz, respectively. The buncher is "tuned" or optimized by adjustment of the AC component. The length of the first-stage buncher is a critical parameter if we are to bunch the current spread over 20 rf cycles into one rf cycle. For maximum efficiency, the length should be an integer multiple of the length parameter $\beta\lambda/2$, where β is the relative velocity of the beam and λ is the free-space wavelength (c/ν) of the operating frequency ν of the buncher. The product $\beta\lambda$ would then be the total distance traveled by the beam particles in one buncher rf cycle. This we can readily see from

$$\beta\lambda = \frac{v}{c} \frac{c}{\nu} = v\tau$$

where τ is the period of the buncher. With a sinusoidal drive, the second buncher gap will negate the bunching action of the first gap if the length is an even multiple of $\beta\lambda/2$ and will enhance the bunching action if the length is an odd multiple. Using the length $\beta\lambda/2$, we see that if in a particular rf cycle a beam particle is accelerated across the first gap, it will be further accelerated across the second gap because the phase of the sine-wave driver has changed by an odd multiple of π .

The change in kinetic energy imposed upon the beam particles after passing through the first-stage buncher (with voltage V) is given approximately by

$$KE \approx KE_0 \left(1 + \frac{2eV}{KE_0} \sin(\omega t + \phi) \right)$$

with $2eV$ being the change in kinetic energy, KE_0 the average kinetic energy and ϕ a phase. Immediately following the first-stage buncher, the beam is accelerated to about

10 KeV, and the net action over 20 rf cycles by the buncher is to slow down the beam particles in the first ten rf cycles and to speed up those in the second ten rf cycles. After a drift distance of 44 cm, the beam particles are bunched significantly, because the faster beam particles have caught up with the slower-moving ones. This can be illustrated by considering the change in expected drift times, t_0 , for the beam particles, given by

$$\frac{\Delta t}{t_0} = -\frac{1}{2} \frac{\Delta K E}{K E_0} \quad (3.1)$$

since t is proportional to $K E^{-1/2}$. At a beam energy of 10 KeV after the first-stage buncher, the drift time through a distance of 44 cm is about 320 ns, and if the buncher is operating at a voltage of 600 VAC the beam particles will be slowed down or speeded up by a maximum of 20 ns as seen in equation 3.1. The length of the first-stage buncher tube was 3.81 cm with gap sizes of 0.80 cm. It had a matching beam energy of about 4.5 KeV at this point.

The beam particles can be further bunched using the low-frequency buncher (LFB) located 10 m upstream of the drift-tube linac. The kinetic energy of the beam at this point has increased to about 750 KeV, where $\beta\lambda/2$ sets the length of this second buncher tube to about 58.58 cm. The voltage used on the LFB is 25 KVAC at an operating frequency of 10.0625 Mhz. The phase difference between the two bunchers is a critical parameter and must be set such that the LFB does not undo the bunching

accomplished by the first stage. When operating at nominal conditions, more than 90% of the P^- beam current is concentrated into one rf-cycle with a period between beam pulses of 100 nanoseconds (10.0625 Mhz), thus making time-of-flight determinations in experiments easier. A typical time history of the P^- beam current is displayed in Fig. 3.3, which was taken from a typical data run and is representative of most of the data-taking conditions.

3.1.3 A Computer Model of the Buncher System

A simple computer program was written in order to understand the action of the buncher in a qualitative manner. It involved a simple treatment of transport of the beam through electrostatic lenses. The input parameters required were the operating voltages and the relative phases between the first-stage buncher and the LFB. Technical information was obtained from Refs. [24] and [48] and incorporated in the program. These data included the physical dimensions of the buncher tubes and gap sizes, as well as drift distances. The gap sizes were on the order of 1 cm, and the lengths of the first-stage buncher and LFB tubes were 3.81 cm and 58.58 cm, respectively. The kinetic energy of the beam at the first buncher tube is about 4 KeV and 750 KeV at the LFB. All of the bunching action occurs when the beam particles are in the gaps, but concentration of the beam current does not occur until the beam drifts some distance. The drift distance between the two buncher tubes was taken to be 44 cm, and

Figure 3.3

Concentration of the P^- beam current with buncher on. Greater than 90% of the P^- beam current over 20 rf cycles was bunched within one rf cycle. This figure represents the typical quality of beam bunching for data acquisition in the experiment. Note the satellite peaks spaced 50 ns from the main peak; these smaller peaks reflect the underlying 5-ns time structure of the beam.

the drift distance between the LFB and the injector was taken to be 10 m which was representative of the actual setup.

When the beam travels through the first tube, the kinetic energies of beam particles are such that after a drift distance of 44 cm, about 40-50% of the beam current is bunched. In the simulation, beam particles were chosen uniformly throughout one buncher cycle (about 100 ns). After passing through the first-stage buncher tube, the beam was then instantaneously accelerated to 10 KeV. A density plot of the change in kinetic energies for each beam particle is shown in Fig. 3.4a. Before the beam particles drift there is no indication of any bunching action occurring. Fig. 3.4b is a density plot of the beam particles after drifting 44 cm, at which point about 40% of the beam lies within a 2 ns region (Fig. 3.4c). The average kinetic energy of the beam has been unaltered by the buncher while the beam current has been bunched together. Using equation 3.1, those particles which have a maximum change in kinetic energy have been displaced in time by about 20 ns.

The computer model next assumed that the beam was instantaneously accelerated to 750 KeV before it traveled through the LFB tube, therefore Fig. 3.4b is a representation of the initial conditions of the beam particles at this point. The LFB was then allowed to act upon these beam particles (Fig. 3.5a), and after exiting they drifted for about 10 meters before entering the linac (Fig. 3.5b). Just at the linac, the beam current was 60% bunched within 2 ns (see Fig. 3.5c). One detail included in the model is the fact

Figure 3.4

The bunching action by the first-stage buncher as modeled with a computer. The kinetic energies for the beam particles exiting the first-stage buncher tube (a) is dependent on the phase of the buncher at the moment in time that each particle passes through. Actual beam bunching does not begin until the particles drift a distance 44 cm; those particles with higher kinetic energies travel this distance in a shorter time than those particles with lower kinetic energies, which take longer (b). The density plots indicate the concentration of beam particles over 100 nsec periods. The x-axes represent one buncher period (100 nsec). About 40% of the beam current lies within 2 ns after the first-stage buncher (c).

Figure 3.5

Plots showing the kinetic energies for the beam particles exiting the LFB buncher tube (a) are also dependent on the phase of the LFB buncher at the moment in time that each particle passes through. The initial conditions for the beam at the LFB tube are those in Fig. 3.4. The density plots indicate the concentration of beam particles over 100 nsec periods; after drifting 10 meters there are very few beam particles that are not bunched (b). The x-axes represent one buncher period (100 nsec). About 60% of the beam current lies within 2 ns after the LFB buncher (c), showing the amount of bunching obtained with both buncher tubes operating.

that the linac has an acceptance window of 25 KeV; that is, beam particles which had kinetic energies outside the range

$$(KE_0 - 12.5 \text{ KeV}) \leq KE \leq (KE_0 + 12.5 \text{ KeV})$$

were not accepted by the linac, where KE_0 is the central energy for the linac acceptance window. This acceptance window is indicated in Fig. 3.5b, and those beam particles with energies outside this window are not accelerated by the drift-tube linac. The bunched beam which was accepted is shown in Fig. 3.6, where 90% of the P^- beam current that is accelerated was bunched within one rf cycle. When the P^- beam is injected into the linac, the beam current is almost completely bunched within one rf cycle, with less than 10% background lying 50 ns away from the P^- main peak. This is most representative of the actual data (refer to Fig. 3.3). The simulation also shows that only about half of the beam exiting the argon gas cell can be bunched within one rf cycle, so while the average beam current is lower with the buncher system operating, the instantaneous beam current is much larger.

3.2 The Beam Servo-Steering System

In an attempt to minimize systematic errors in this experiment due to beam effects, a feedback system was designed and constructed to automatically control beam steering. In previous experiments using the P^- beam, it had been observed that there

Figure 3.6

Plot showing the effect of the beam buncher as modeled with a computer. About 90% of the beam current accepted by the drift-tube linac is bunched within 2 ns. Because of the window on the beam particles kinetic energies imposed by the linac, nearly 60% of the total beam current exiting the argon gas cell is accelerated by the linac.

existed systematic beam motions between spin reversals at a level which would limit the sensitivity of this experiment.

3.2.1 Why a Beam Feedback System was Required

It had been known for some time that the position of the P^- beam in Line B varied according to the Normal and Reverse spin directions. Such variations at the LD₂ could be as large as 5 mm. Previous experiments were not very sensitive to this systematic effect. Because of the design of this experiment and the level of sensitivity required, it was clear that some control over systematic beam motion was required.

The beam current monitor used to normalize the data in this experiment was the SEM which was located just upstream of the LD₂ target in the proton beam. It is necessary that the number of neutrons produced per incident proton at the LD₂ target be insensitive to beam position at the LD₂. Unfortunately, the LD₂ target did not have flat windows, but rather the windows were bowed outwards. Since the number of neutrons produced for each incident proton depends upon the path length through the liquid target, one can easily see that the number of neutrons produced can depend relatively strongly upon the steering of the proton beam spot at the upstream LD₂ target window.

A good measure of $\Delta\sigma_L$ requires an asymmetry measurement on the order of ± 0.0001 . Experimentally, it was found that the beam spot position would have to

be fixed to better than ± 0.1 mm in order to meet this requirement. As a result of this, a beam steering feedback system was implemented in Line B at LAMPF by M. McNaughton (LANL).^[25] J. Faucett (LANL) wrote much of the software that interfaced the experiment with the servo-steering system. The purpose of this section is to discuss the operation and performance of this system.

3.2.2 How the Servo-Steering System Works

The servo-steering system was controlled by the Line B microVax computer. The program was written in Pascal and contains 2000 lines of code. The system reads the beam x and y profiles from LBII6 continuously, and computes means and standard deviations of beam centroids. Based upon these numbers, the currents to the beam steering magnets LBSM4X and LBSM4Y are modified by signals from the computer in order to maintain a desired spot position, which was set by the experimenter.

Data are accumulated in a total of twenty separate buffers: ten for spin Normal and ten for spin Reverse. Each of the buffers accumulates data over different time intervals for various levels of steering corrections; for example, buffer#1 might read the beam profile data four times before calculating a mean and standard deviation, while buffer#5 might read the data 1024 times before calculating a mean and standard deviation. This way, buffer#1 would find and correct large steering errors very quickly. Buffer#5, taking longer, would find and correct smaller steering errors.

The steering corrections are accomplished by writing to a CAMAC digital-to-analog (D/A) convertor, which in turn makes the required adjustments to the steering magnets in the proton beam. In the next section the performance of the servo-steering system is discussed.

3.2.3 Performance of the Servo-Steering System

The servo-steering system is capable of reading a beam profile from LBII6 every beam macropulse (40 Hz), calculating means and standard deviations from beam profiles and correcting steering errors, in about 0.1 second. The success of the servo-steering system is demonstrated in Fig. 3.7 which was taken from a typical run. In Fig. 3.7, note that the proton beam spot position at the LD₂ target was kept to within 0.1 mm during the run. Also note that the currents for the beam steering magnets in the proton beam were being constantly adjusted to maintain the spot position.

A few problems with the system were often encountered during data taking. The steering corrections were communicated to the steering magnets by CAMAC, where digital counts were converted to current adjustments for the steering magnets. Occasionally the magnets reached the limit of their current range, when the conversion of counts to current became meaningless and steering would cease. In this case, the proton beam would have to be re-steered and the servo-steering system re-started. Fortunately, this was not a common occurrence.

Figure 3.7
 Demonstration of the success of the servo-steering system (plots taken from a typical data run). Note that maintaining the beam spot position (a) to better than 0.1 mm at the liquid deuterium target required continuous adjustments of the steering magnets (b).

A second problem which often occurred was related to the beam tuning. The beam profile would often exhibit a long tail on one side of the centroid. The beam profile monitor was nonlinear, however, and such tails might not be real. Depending on a chosen background threshold, which was a user-set parameter, this could cause the centroid to be biased towards the side with the tail. The resulting steering corrections would, therefore, be incorrect. If the tail remained constant over some period in time, the beam spot coordinates x and y could be adjusted to compensate. Alternatively, the beam threshold could be set above the tail at the expense of statistics.

A related problem resulted from beam widths and centroids changing while the steering was on. The centroid, as determined by the computer, may have been different from the actual beam centroid in such a case because of the simple algorithm used.

Earlier servo-steering systems used analog electronics with exponential time constants rather than digital systems with intrinsically discrete changes. It was always thought that such digital systems would enter some oscillating mode which could never be controlled. However, a more stable performance was found with the digital computer approach by basing decisions on statistics before making any discrete steering corrections. The digital computer servo-steering system has shown several advantages over analog systems (from Ref. [25]):

- The signal from LBH6 could be conditioned by digital means. For example, the various wires could be assigned various gains to account for nonlinearities.

- The system used several data buffers, each with its own time constant. Major steering problems were found and corrected quickly, while smaller errors required more statistics before making corrections.
- During each spin reversal the computer switched buffers, anticipating systematic changes between N and R.
- The performance of the system was readily available to the experiment over Ethernet.

CHAPTER 4 THE EXPERIMENTAL SETUP

When the neutron beam exits the collimator and enters the experimental area named BR, it passes through various neutron monitors and magnets before it passes through a polarized proton target and is detected by an array of neutron counters located downstream. The neutron monitors, the magnets, the proton target and the neutron counters were optically aligned along the axis of the collimator. A diagram of the experimental setup used is illustrated in Fig. 4.1. In this chapter the experiment setup is discussed in detail.

4.1 Neutron Beam Collimator

After the neutrons are produced by the liquid deuterium target, they are collimated into a narrow beam before entering the experimental area. There is a 3.7-m-thick steel wall separating the LD_2 target and the experimental area which amounts to about 20 attenuation lengths for neutrons at 800 MeV. This steel wall is shielded further at the top with 0.91-m-thick concrete blocks. This amount of total shielding helped to reduce the background radiation in the area. Embedded within the steel shielding wall are six

Figure 4.1

Experimental setup used to measure $\Delta\sigma_L$. The neutrons produced at the LD_2 target passed through a collimator before entering the experimental area. The collimated neutron beam next passed through the magnetic fields of Lorraine and Castor, which rotated the spins of the neutrons to $\pm L$ -type. The proton target was polarized to $\pm L$ -type as well, and the neutrons passing through the polarized proton target were detected by a stack of neutron counters, located 4 meters downstream of the target.

gun barrels with diameters of 155 mm oriented at six angles to the proton beam. These serve as holders for collimator inserts. In this experiment, we used the gun barrel at 0° to the beam which is 365.78 cm long. It was filled with 22 steel slugs, each with a 2.54 cm hole centered on the gun-barrel axis. On the extreme downstream end of the collimator, lead plugs of various thicknesses (>10 radiation lengths) were inserted into the collimator to attenuate any gamma rays coming down the collimator.

4.2 The Front Neutron Monitors

When the neutron beam exits the collimator and enters the experimental area, it is monitored by the neutron front monitors (FMON). The neutrons detected by FMON serve to help monitor the neutron beam and are also involved in the normalization of our data.

A layout of the front monitor is displayed in Fig. 4.2. As the neutron beam travels through FMON, it first passes through the A counter, a veto counter used throughout the front monitor logic to reject charged particles coming through the collimator. Immediately following the A counter is a block of CH material in which a small fraction of neutrons interact to give charged particles to be detected in FMON. The neutron beam then passes through the M1 counter and the center of the M4 counter. Following the M4 counter are two scintillators in a left/right configuration making up the M2 counter. Between M2 and M3 there is a brass plate with a hole at its center. The last

Figure 4.2

Diagram of the neutron front monitor, which consisted of five independent monitors: M1; M2L and M2R; M3; M4U, M4D, M4L and M4R; and M5. Charged particles exiting the collimator were rejected using the charged-particle veto counter (A). A small CH₂ target converted some of the neutron beam to charged secondary particles; these were then detected by the set of M counters.

counter in the front monitor is M3 which also has a hole at its center. This is basically the same monitor used in two previous experiments at LAMPF, namely, experiment numbers 665 and 770 (see Refs. [29] and [30]). For our experiment, we added four more scintillators, since a more stable monitor and higher counting rates were needed. These new counters as a unit are called "M4," and as configured they offered four new signals to be monitored; M4 left/down, M4 left/up, M4 right/up and M4 right/down. Various combinations of these signals were used to monitor the steering of the neutrons coming out of the collimator into the area. For example, by a logical OR between the signals M4 left/down and M4 left/up, we obtained a signal called M4 left. Comparing this signal to one called M4 right, obtained in a similar manner, we could monitor the ratio of left/right counts (as well as up/down counts) from run to run and get a measure of the stability of the neutron beam position. A logical OR of the four M4 counters formed the neutron monitor M4Tot. M4Tot was used to correct the data for "dead time" effects (shown in Chapter 6).

4.3 The Spin Precession Magnets and the Neutron Analyzer

Downstream of the front monitor were two spin precession magnets called "Lorraine" and "Castor." These were used to re-orient the direction of the spin of the exiting neutron beam into the direction desired. Immediately downstream of these magnets was a neutron polarization analyzer called JPAN. The main purpose of JPAN was to

help determine the proper Lorraine and Castor currents for obtaining the proper spin direction of the neutron beam. During the data taking, JPAN was moved out of the beam line as its presence would only serve to degrade further the quality of the neutron beam.

After leaving the liquid deuterium target, the neutrons passed through the magnetic field of LBBM6. This deflected the non-interacting proton beam into the beam dump, but at the same time it precessed the neutron spin direction away from its original orientation. Lorraine and Castor corrected for this by re-rotating the neutron spins so that we could obtain the longitudinal spin direction needed for the experiment. Extensive analysis of the magnets, Lorraine and Castor, has been carried out by H. Spinka^[49] (ANL).

JPAN was designed by H. Shimizu, based on an earlier neutron polarimeter,^[50] and is diagrammed in Fig. 4.3. It consisted of a CH₂ target, a scintillation counter (N1) used to detect charged particles produced in the CH₂ block, brass absorbers, and four identical scintillation counters set up in left, right, up, down configuration located at a 30 degree angle to the beam. The brass absorbers were used to range out the low energy production and so gave greater weight to the peak energy spin asymmetries in the adjustment of Lorraine and Castor. Coincidences between N1 and the other four counters were recorded by the online computer for both spin states of the proton beam (Normal and Reverse), and from these data left/right and up/down spin asymmetries

Figure 4.3

The neutron polarimeter used in the experiment. It was similar in design to the proton beam polarimeters discussed in Chapter 2, Section 2.4.3. Not shown in the figure are two additional counters, NL (left) and NR (right).

were computed. The following formulae were used:

$$\epsilon_{LR} = \frac{\sqrt{L_1 R_1} - \sqrt{L_1 R_1}}{\sqrt{L_1 R_1} + \sqrt{L_1 R_1}}$$

$$\epsilon_{UD} = \frac{\sqrt{U_1 D_1} - \sqrt{U_1 D_1}}{\sqrt{U_1 D_1} + \sqrt{U_1 D_1}}$$

which are geometrical asymmetries. In order to determine the proper settings of Lorraine and Castor, several data points had to be taken by varying the current settings and calculating the resulting asymmetries. For L type beam, those current settings which resulted in no left/right or up/down asymmetries were found. Lorraine was set to rotate the neutron spin to a fully transverse horizontal direction (90 degrees to beam line), and Castor was set to rotate the spin back to a direction parallel (NOR) or anti-parallel (REV) to the beam direction. This operating procedure had the virtue that any charged particles exiting the collimator would be completely swept away from the beam line, and a neutral beam would continue to the polarized target downstream.

4.4 The Polarized Proton Target System

The polarized proton target system was, in a sense, the focal point of the experiment. A good measure of $\Delta\sigma_L$ would have not been possible if not for the large volume of the target which was the largest target ever used at LAMPF at the time the experiment ran. Such a large target has to be maintained by numerous external components (such as pumps, microwave power supplies, cryogenics, etc.) operating in concert with each

other. Some of these components are discussed in this section. The technology for polarizing targets has been around for some time, and much spin-spin physics has been done using such targets.

4.4.1 Dynamic Nuclear Polarization

The measurement of $\Delta\sigma_L$ requires both a polarized beam and polarized target. The only practical way to control the spin of a particle is through its magnetic moment which is fundamentally related to its spin and charge. In an external magnetic field, the magnetic moment responds to an external force which tends to align the spin vector along the direction of the field. However, the temperatures and field strengths required are not practical for polarizing protons by this brute-force method.

Consider an ensemble of protons in an external magnetic field H_0 at equilibrium at an absolute temperature T . Each proton can be in either one of two states; the state (+) where the proton is aligned parallel to the field, and the state (-) where the proton is aligned antiparallel to the field. The energies of these two states are

$$\epsilon_+ = -\mu H$$

$$\epsilon_- = +\mu H$$

where μ is the magnetic moment of the proton. Here, H should be considered as being the local magnetic field at the proton. If the density of protons is low, the local field acting on each proton is nearly that of the external field H_0 .

The probability of finding a proton in a state with energy ϵ is proportional to the Boltzmann factor $e^{-\epsilon/kT}$. The polarization, P , is given in terms of the average magnetic moment as follows,

$$\bar{\mu} = \mu \frac{e^{+\mu H/kT} - e^{-\mu H/kT}}{e^{+\mu H/kT} + e^{-\mu H/kT}}$$

from which $P = \tanh \mu H/kT$.

With the typical numbers in this experiment, $H_0 = 2.5$ Tesla and $T = 0.5$ K, the average proton polarizability is only 0.5%. Even at such low temperatures and high magnetic field strengths, the protons cannot be polarized beyond 1% which is hardly suitable for experiments of this type. On the other hand, consider the polarizability of electrons. With a magnetic moment nearly 660 times larger than that of the proton, electrons can be polarized to -99.8% (the electron's magnetic moment is negative and therefore the electron spin is antiparallel to the field). This provides the key to polarizing protons to a much higher degree, by a non-equilibrium process called "Dynamic Nuclear Polarization."^[51]

Consider a target sample that contains the nuclei we wish to polarize, along with other spins having large magnetic moments (such as free radicals, or ions with unpaired electrons). Because of the interactions between the various spin types, the high degree of ordering of the electron spins can be transferred to the nuclei, thereby enhancing their polarization. In this experiment, a polarized proton target was used, and we shall

restrict ourselves to this case where the polarizable nuclei are protons, which are located within a "lattice" of non-polarizable material that acts as a thermodynamic reservoir. We next place this proton-electron system in an external magnetic field, H_o , pointing in the z-direction. Let I represent the spin vector for the proton, and S the spin vector for the electron. The Hamiltonian for this system is then well represented by

$$\mathcal{H} = \mathcal{H}_o + \gamma_e \hbar H_o S_z - \gamma_p \hbar H_o I_z + AS \cdot I \quad (4.1)$$

where γ_e and γ_p are the electron and proton gyromagnetic ratios, respectively, and A is an electron-proton spin coupling constant, the principle source of the hyperfine structure. The coupling A is assumed to be very small in comparison with the electron and proton Zeeman interactions, and it is this hyperfine interaction which "transfers" the polarization of the electrons to the protons. The term $\mathcal{H}_o$ included in equation 4.1 is the Hamiltonian describing the interactions between the electron/proton spins and the surrounding target lattice. The electrons are strongly coupled to the lattice, and if perturbed they will quickly return to equilibrium with relaxation times of ≈ 40 ns at 0.5 K. The proton spins are weakly coupled and have relaxation times of several hours. As a result of this, the spin-interaction Hamiltonian may be extracted from equation 4.1 and is effectively

$$\mathcal{H}_{spin} = \gamma_e \hbar H_o S_z - \gamma_p \hbar H_o I_z + AS_z I_z \quad (4.2)$$

with energy eigenvalues

$$E_{spin} = \gamma_e \hbar H_0 m_S - \gamma_p \hbar H_0 m_I + A m_S m_I \quad (4.3)$$

where $m_S = \pm 1/2$ and $m_I = \pm 1/2$ for the electron and proton, respectively. There are four combinations of spins in which the electrons and protons may orient themselves in the following manner:

- electron spin "down"/proton spin "down" (--)
- electron spin "down"/proton spin "up" (-+)
- electron spin "up"/proton spin "down" (+-)
- electron spin "up"/proton spin "up" (++)

The energies for these four states are

$$E(--)= -\mu_e H_0 + \mu_p H_0 + A$$

$$E(-+)= -\mu_e H_0 - \mu_p H_0 - A$$

$$E(+ -)= +\mu_e H_0 + \mu_p H_0 - A$$

$$E(++)= +\mu_e H_0 - \mu_p H_0 + A.$$

The quantities m_S and m_I are the projections of the electron spin and proton spin onto the axis of quantization, and the selection rules for magnetic dipole transitions follow from equation 4.2 after one finds the non-zero matrix elements of the spin Hamiltonian

(equation 4.2) for magnetic dipole transitions. The allowed transitions are $\Delta m_S = 0$, $\Delta m_I = \pm 1$ and $\Delta m_S = \pm 1, \Delta m_I = 0$. A schematic of the energy levels with allowed transitions indicated is diagrammed in Fig. 4.4 for a magnetic field pointing "up". The so-called "forbidden" transitions are also shown, which occur with much smaller probability because they involve dipole-dipole transitions. In Fig. 4.4, two types of allowed transitions are clearly evident. They are electron single flip (I, II) and proton single flip (III, IV). Note that the electron single-flip transitions have lifetimes on the order of milliseconds while proton single-flip transition lifetimes are on the order of hours. The energies required to induce these single-flip transitions are

$$h\nu_p = 2\mu_p H_0 \pm A \approx 106 \text{ Mhz}$$

$$h\nu_e = 2\mu_e H_0 \pm A \approx 70 \text{ Ghz.}$$

The proton target may be polarized by using the forbidden transitions in which both the electron and proton spins flip (double flip transitions). As seen in Fig. 4.4, the electron/proton state $(-+)$ may be excited to the $(+-)$ state by irradiating the system with microwaves of frequency $\nu_e + \nu_p$. The state $(+-)$ then decays in either of one of two paths, the most likely of which is to the $(--)$ state where the population of protons in the $(-)$ state then increases. On the other hand, if the system is irradiated with microwaves of frequency $\nu_e - \nu_p$, the state $(--)$ is excited to the state $(++)$, where the most probable decay path is to the state $(-+)$, and the population of protons in the

Figure 4.4

Energy-level diagram for an electron-proton system in a strong external magnetic field H_0 . Allowed transitions are indicated by solid lines, "forbidden" transitions by dotted lines. The notation used is the spin orientation for the (electron, proton). The energy required to flip a proton corresponds to a frequency of 106 Mhz, while a frequency of 70 Ghz is required to flip an electron spin. Protons in the state $(--)$ can be aligned with the magnetic field by stimulating the transition $(--)\rightarrow(++)$, after which the $(++)$ state decays to $(-+)$. Protons in the state $(-+)$ are aligned against the magnetic field by stimulating the transition $(-+)\rightarrow(+)$, after which the $(+)$ state decays to $(--)$.

(+) state is increased. Note that, in both schemes, the spins of the electrons return to their initial state, but in both cases a proton spin has been flipped. In either case the desired proton state is "pumped", much as with a laser or maser, and the protons can be polarized to a high degree. Note also that the bandwidth of the microwave source must be sufficiently narrow ($\Delta\nu \ll \nu_p$) for the target to polarize. Otherwise the two processes compete with no net polarization. An additional requirement on the microwave source is that enough power must be delivered to the target so that the probability of double flip transitions is greater than the thermal transition rate III (see Fig. 4.4) which tends to depolarize the protons. This puts the system in a nonequilibrium state, and microwave power must be continually delivered. However, too much power will destroy the high ordering of the electron spins because of thermal heating, and in the extreme limit there will not be enough electrons in the (-) state for this dynamic polarization process to continue. In the experiment, target polarizations close to $\pm 70\%$ were achieved. One of the remarkable features of the dynamic polarization approach is that the target can be polarized either aligned or anti-aligned with the same external magnetic field by simply re-tuning the frequency of the microwave power source.

4.4.2 The Polarized Proton Target Magnet

One of the factors which directly affects the quality of the measured $\Delta\sigma_L$ is the so-called target constant, explicitly

$$A = (\rho_F L)^{-1}.$$

Here ρ_F is the number density for free protons in the target and L is the target length. The free proton density represents those protons which are available for polarization in the region spanned by the beam. All things being equal, the quality of the measurement scales with the target length and the extent to which the beam is subtended by the target. With a longer target, the raw signal is larger resulting in a smaller error overall. It was also desirable that the diameter of the target be large enough so that all of the neutron beam passes through the target.

To help satisfy these requirements for the experiment, a new superconducting magnet was used at LAMPF. Dubbed "Cyclops", this magnet generated a solenoid field with a strength of 2.5 Tesla, uniform to 10^{-4} , within a cylindrical volume of 5 cm diameter by 25 cm long along the axis of the magnet. This uniformity is important in polarizing the target material, as different portions of the target would require microwaves of different frequencies if the magnetic field were not uniform. With the target material constrained within this volume, the target cell used in the experiment measured 4.7 cm in diameter and 12.5 cm long (217 cm^3), the largest target used to date at LAMPF.

The large size of this target posed several problems, one of which was the question as to whether the existing refrigerator would be adequate in maintaining the target material at a low temperature. When the experiment was begun, a target cell 20-cm long was tested. It soon became apparent that a smaller target would have to be used since the refrigerator could not keep the target adequately cold. As a general rule of thumb, about 1 mW of microwave power has to be delivered for each gram of target material in order to polarize it.^[52] Still more power is required when reversing the polarization. With the target 20 cm long, the refrigerator would have to provide about 400 mW of cooling. The refrigerator could provide only 250 mW of cooling, so the target length was shortened to 12.5 cm.

Another problem using a large target was related to the delivery of the microwaves into the target cell. The microwaves did not uniformly illuminate the target material at first, and as a result there was a substantial deviation in the polarization between the upstream side of the target to the downstream side. Initially, the microwaves were delivered from the rear of the target cavity which was incased in copper. The microwaves suffered attenuation in dissipating within the target material. The target cavity was therefore modified by removing the copper casing surrounding the target cell and replacing it with teflon which is transparent to microwaves. The microwave wave guide was then moved from the rear of the target cavity and modified to deliver microwaves along the side of the target cavity, through the teflon. This illuminated the target ma-

terial uniformly, and the polarization was then found (using NMR sense coils) to be satisfactorily uniform throughout the target.

A layout of the polarized target system used in the experiment is shown in Fig. 4.5. A μ PDP computer periodically measured the polarization of the target using NMR techniques. The polarization was then available to the experiment through CAMAC. A unique addition to the generic target systems previously used was an IBM computer which monitored the target system, reporting any deviations from standard normal operating conditions. The computer also had some limited control over the target system, such as the microwave frequency used to polarize the target, and the system was "tweaked" occasionally in order to optimize the target polarization. This was developed by J. Faucett (LANL). The IBM computer was interfaced to the target system via CAMAC.

4.5 The Neutron Counter Array

Before the experiment ran, a prototype detector was built and beam tests were conducted to determine the feasibility of the proposed experiment. The detector support frame was built of "unistrut" materials and supported two planes of neutron counters, six counters/plane. The tests were conducted during the running of E770,^[29] which was also a Ph.D. thesis experiment. The preliminary studies were useful in characterizing the counters and learning how to use special fast timing electronics built by Argonne

Figure 4.5

Layout of the PPT system. The target material was contained within a target cell located inside a cryostat (Gandolf), at the center of a superconducting magnet (Cyclops). The microwaves used to polarize the proton target were delivered from the rear of the cryostat.

National Laboratory for these counters.

4.5.1 Construction Details

The neutron counters downstream of the PPT were used to detect the neutrons passing through the polarized target. There were a total of 24 counters arranged in planes of six counters each, with a charged-particle veto counter in front. Each neutron counter was made up of two blocks of NE-102 plastic scintillator glued together to form a block of scintillator with an active volume of $11.4 \times 25.4 \times 61.0 \text{ cm}^3$. A lucite light guide, 65 cm long, was glued onto each end of a counter, and Amperex 58 AVP photomultiplier tubes were optically coupled to the light guides to view the scintillation light (a tube base schematic is shown in Fig. 4.6). In order to capture scintillator light entering at oblique angles, Dow-Corning optical grease was used to join the photomultiplier tubes on both ends of the counters. The photomultipliers were offset 5 cm to allow the counters to nest together snugly without sizable gaps between the stacked counters.

Each scintillator and its light guides were wrapped in a layer of aluminized mylar. The counter was then cushioned on four sides by 3.2-mm-thick synthetic rubber, and the entire assembly was inserted into a rectangular steel can made of 1.6-mm-thick steel sheet metal. The photomultiplier tubes were surrounded by mu-metal shields and cylindrical cans made of 1.6-mm-thick steel. The cylindrical cans were bolted on each end of the rectangular can. Large springs pressed the photomultiplier tubes against the

EXP 980 NEUTRON COUNTER
TUBE BASE NOV-15-1988

Figure 4.6

Photomultiplier tube-base electronic design used in the experiment.

light guides.

The neutron counters were then stacked with the 25.4-cm dimension along the beam direction. Adjacent planes of counters were oriented perpendicular to each other. A heavy-duty frame was constructed at LAMPF to hold the counters. The neutron counters could be rotated 360°, raised or lowered, or moved transversely to the beam. The frame was also equipped with wheels so that the distance to the proton target could be varied. This arrangement allowed systematic studies to be carried out during data taking and made it possible to calibrate the neutron counters using cosmic rays.

Two techniques were used to determine the interaction points for detected neutrons in the hodoscope. The first method used fast time-to-digital convertors (FTDCs) which were constructed at Argonne.^[53] If a neutron counter detected a hit, a FTDC calculated the time difference between the photomultiplier signals from that counter. This determined the interaction point. The second method used to determine the interaction point for a detected neutron used the feature that adjacent planes were perpendicular to each other. A coincidence between counters in adjacent planes fixes the position of the neutron to within the overlap of the orthogonal counters.

4.5.2 Cosmic Ray Calibration

When the anode signals from the photomultiplier tubes on each side of a counter were put into a logical coincidence, the coincidence signal was most likely to be due to

a real event such as a neutron, rather than a random coincidence between noise pulses. When the neutron beam was off, we still found signals due to particles detected in the neutron counters. It happens that this neutron counter stack, like all scintillation counters, was also a cosmic-ray detector, although the rates were small compared with the neutron detection rate with the beam on. We made use of this "built in" feature, and of the macro structure of the neutron beam (40 Hz), to monitor the gains of the counters at the same time we took data. Neutrons do not interact in scintillators in the same manner as charged particles, and as a result the neutron counters could not be plateaued as most scintillators are. This is because neutrons interact anywhere within the counters and give a continuum of pulse heights, whereas charged particles produce a narrow peak in the pulse height distribution since they lose energy fairly uniformly as they pass through the counters. The neutron beam has a duty cycle of about 9% ($\leq 3\%$ for energies below 800 MeV), and so most of the time the neutron beam was off. By creating a logic gate representing the time when the neutron beam was off, we could open a time window in which cosmic rays could be detected and sampled by the online computer. The particles of interest were the cosmic ray muons, which were minimum ionizing. By accumulating a histogram of the anode pulse heights from each photomultiplier tube that detected a muon, we could monitor the gains of the counters. Tube base voltages would then be adjusted in order to maintain these gains.

The online computer sampled the cosmic ray events and histogrammed the cosmic

ray pulse heights for the neutron counters only if at least three adjacent counters detected a charged particle at the same instant. This was found to yield clean muon spectra which were used to set the gains of the counters. The pulse-height distributions for counters which were in a horizontal orientation differed somewhat from those counters oriented vertically, as seen in Fig. 4.7, which was taken from a typical data-taking run. The histograms for those counters in a vertical orientation had considerably fewer counts than the histograms of horizontal counters. This is primarily due to the fact that more of the cosmic rays come from a vertical direction, so that many of these cosmic rays cannot satisfy the requirement of traveling through three adjacent vertical counters. The shapes and peak positions of the pulse-height distributions for the two orientations are also different in that the peaks of the distributions from vertical counters are wider and shifted towards larger pulse heights than those of horizontal counters. Since the pulse height is related to the amount of energy deposited into the counter from the passage of a charged particle, we expect the pulse height distributions for vertical counters to be shifted toward higher energies. This is because any charged particles which go through at least three adjacent counters must traverse longer path lengths than for horizontal counters, depositing more energy in the process. Pulse-height distributions were accumulated for the counters arranged horizontally, and these were compared with the pulse-height distributions accumulated for the same counters when rotated into a vertical orientation.

Figure 4.7

Pulse height distributions for muons in the neutron counters, taken from a typical run. The angular distribution for cosmic ray muons is peaked towards the zenith direction, indicating that the difference between the horizontal and vertical pulse height spectra was primarily due to geometry.

It was found that the peak position shifted, for a given counter, by a factor of nearly 1.4 between horizontal and vertical orientations. This was later modeled by a simple computer program code where it was assumed that:

- The muons were monoenergetic.
- The amount of energy deposited within a counter is proportional to the path length within the counter.
- The muon flux distribution was concentrated towards the zenith, with a distribution given by (see Ref. [54])

$$\frac{d^2 \dot{N}}{d\Omega d\sigma} \propto \cos^2 \theta,$$

where $\dot{N}$ is the number of muons per unit time, $d\Omega$ is a unit of solid angle and $d\sigma$ is unit surface area.

- The muons had to go through three adjacent counters.

Very roughly, 75% of the cosmic rays at sea level are muons having an average energy of 2 GeV. The energy distribution for these muons varies as E^{-2} , but this was ignored in the model as such muons are minimum-ionizing. The energy loss rate for minimum-ionizing particles, dE/dx , is essentially constant for all energies above the minimum-ionizing energy. The resolution of the photomultiplier tubes were not taken into consideration, but the results of the model show that the differences between the pulse height distributions between horizontal and vertical counters is mainly due to ge-

ometrical considerations (see Fig. 4.8). The ratio of the centroids for the distributions given by the model is 1.4, matching the data nicely.

4.5.3 Efficiency of the Counters

A measurement of the detection efficiency,^[55] for neutrons of energies 289 and 435 MeV, was done by replacing the polarized proton target with a liquid hydrogen (LH₂) target which was a flask measuring 6.81 cm in diameter and with rounded ends that were separated by a length of 6.76 cm (some background data were collected with the target filled with gaseous hydrogen). The neutrons, which were elastically scattered from the LH₂ target, were "tagged" by the recoil protons. The energy of the incident neutron beam was 484 MeV at the LH₂ target. A layout of the experiment for the efficiency measurements is shown in Fig. 4.9.

The recoil protons were detected by a set of scintillation counters which were mounted on a frame constructed of unistrut material. This frame was attached at a pivot point located beneath the LH₂ target center which allowed the frame to be rotated about this pivot. The readout trigger (GT) was a coincidence of the scintillation counters P₁, P₂ and R₁. The trigger rate was less than 1 Hz, and accidentals were found to be negligible. The average proton beam intensity was 5-6 nA which corresponded to an estimated average of about 12,000 neutrons/sec (incident peak energy neutrons).

The P₁ and P₂ counters had dimensions 10.2×10.2 cm² and 20.3×20.3 cm², and

Figure 4.8
Results of a model simulating muon detection in the counters, using the geometry of the neutron counters as inputs. Compare with Fig. 4.7.

Figure 4.9
Experimental layout for the measurement of neutron counter efficiency. The liquid hydrogen target replaced the polarized proton target used in the experiment.

were both .32-cm thick. The acceptance solid angles for elastic events were defined by the P_2 counter which was located 257 cm from the center of the LH_2 target. The corresponding spot size of the neutrons at the neutron hodoscope was estimated at 12 cm (FWHM). The neutron detectors were located 236 (248) cm from the LH_2 target for 289 (425) MeV. The R_1 and R_2 counters were 0.64 cm thick and measured 27.9 cm high and 30.5 cm wide. The distances from the pivot to the P_1 , R_1 and R_2 counters were 50, 275 and 315 cm, respectively. The proton and neutron laboratory angles were 47.3° (69.4°) and 36.0° (16.6°) for 289 (435) MeV.

Background detection was reduced using various requirements in the recoil proton arm. Elastic recoil protons were selected on the basis of their time-of-flight (TOF) through the proton arm which was compared to the 10 Mhz timing of the buncher. Most events associated with low energy incident neutrons in the beam were then rejected. Other charged particles (mostly pions) were distinguishable by their TOF and absorption range. Aluminum plates of various thicknesses were placed in the proton arm and used to range out the pions and any inelastic or slow recoil particles. The thicknesses were chosen to permit elastic protons associated with peak energy neutrons to reach R_1 but to be stopped before reaching R_2 . High energy pions that managed to reach R_2 were rejected by their TOF in the offline analysis. The plates were located between P_2 and R_1 and between R_1 and R_2 . For scattered neutron energies of 289 (435) MeV and proton laboratory angles of 47.3° (69.4°), the aluminum plate

thicknesses used were 10.2 (0.32) cm and 3.8 (1.27) cm.

The number of good triggers (GT) at 289 MeV was 677 and 676 at 435 MeV. The proton arm timings that were gated by GT, as well as the timings and pulse heights for the counters in the neutron hodoscope, were recorded on magnetic tape and analyzed offline by M. Devereux.^[56] For each GT, the pulse heights and timing data for the neutron counters in the hodoscope were analyzed. It was assumed that either neutrons or cosmic rays were detected by the neutron counters. The neutron TOF and pulse height data were used to establish the identity of the particles detected. In this way, the cosmic ray background (BKG) was estimated to be 4 ± 4 events per plane at either neutron energy. Cuts were made on the histograms for the proton TOF data, to reject pions and other inelastics which made a GT at the time of data taking.

The efficiencies ϵ for the first three planes (P1, P2 and P3) were given by the fraction of neutrons to protons. Plane 4 was not operational at this time. The separate efficiencies that were calculated for each plane in the neutron hodoscope were obtained using the relations

$$\begin{aligned}\epsilon(P1) &= \frac{GT \cdot P1 - BKG}{GT}, \\ \epsilon(P2) &= \frac{GT \cdot P2 \cdot \overline{P1} - BKG}{GT}, \\ \epsilon(P3) &= \frac{GT \cdot P3 \cdot \overline{P2} \cdot \overline{P1} - BKG}{GT}.\end{aligned}$$

The errors on the number of neutrons detected were estimated from the binomial dis-

tribution and were added in quadrature with the uncertainty in BKG to get the total errors for the efficiencies $\epsilon(P1)$, $\epsilon(P2)$ and $\epsilon(P3)$. Table. 4.1 summarizes the efficiencies found from the analysis.

The scintillator thresholds were determined from the muon pulse height spectra collected in cosmic ray runs. Assuming that the muon peak in the ADC spectra corresponded to muons depositing 19.4 MeV in the neutron detectors (minimum-ionizing energy loss for muons is -1.7 MeV/cm), we found the discriminator thresholds for planes P1, P2 and P3 to be equivalent to 2.3 ± 0.6 , 2.4 ± 0.2 and $3.4 \pm 0.6 \text{ MeV}$. The experimental results were compared with a simulation using the Monte-Carlo of Sailor, et al.^[57] The steel rectangular cans, mylar, and synthetic rubber cushioning the detectors were not included in the model. The experimental values collected in Table. 4.1 are within two standard deviations of the efficiencies predicted.

Table 4.1

Neutron counter efficiencies for 289 and 435 MeV neutrons.

	289 MeV	435 MeV
$\epsilon(P1)$	0.207 ± 0.017	0.237 ± 0.017
$\epsilon(P2)$	0.150 ± 0.015	0.153 ± 0.015
$\epsilon(P3)$	0.126 ± 0.014	0.133 ± 0.014
ϵ_{TOT}	0.483 ± 0.022	0.523 ± 0.022

CHAPTER 5 DATA ACQUISITION

Because this experiment was basically a scaler experiment, the incoming data had to be pre-processed by some hardware arrangement before being processed by a computer. The hardware logic used in this experiment was extremely complex (the number of scalers used probably set a new record at LAMPF). This chapter is, therefore, divided into several sections in an attempt to coherently discuss how the data were acquired in the experiment.

5.1 Basic Philosophy of the Logic

This experiment was strictly a scaler experiment; that is, scattering events and monitor counts were accumulated in hardware scalers and read out after some period of time by a computer, to be written on magnetic tape. However, in order to monitor the state of the experiment, real-time sampling of various quantities was necessary. We could monitor such things as the quality of the beam bunching, the stability of the beam steering at the liquid deuterium target and the gains of the neutron counters. The number of events read by the computer were then classified by types as follows:

- EVENT5 - The beam servo-steering is monitored.
- EVENT6 - The PPT target polarization is recorded.
- EVENT7 - All the scalers are read. This usually occurs during spin changes.
- EVENT8 - Cosmic ray events detected in neutron bar counters are read.
- EVENT9 - Beam related neutrons are detected in counters.

The reason for the event nomenclature above is related to the standard computer hardware and software used at LAMPF for data acquisition.

With a scaler experiment, nearly all of the logical decision making had to be done with hardware. At the event rates we experienced, the online computer could not examine each event or make any software decisions concerning each event, so the hardware had to decide whether an event was good (i.e., should be scaled) or whether it should be rejected. While the basic logic used in the experiment is relatively simple to understand, with 24 neutron counters it became somewhat intricate to handle in a practical sense.

The data acquisition hardware used in the experiment was standard at LAMPF and included a CAMAC system, a Micro-programmable Branch Driver (MBD), two LAMPF gate generators (LGG) and a LAMPF trigger module. The online computer used in the data acquisition was a MicroVAX II, which also performed some analysis of the data. With the system at LAMPF, it is the experimenter's responsibility to determine which logic events to set up for communication between the experiment and

the computer. Once this is done, signals representing these events are patched to a LAMPF trigger module. These modules have several inputs accommodating a wide variety of possible events, and when a particular input is driven by some event signal, the MBD is triggered. The MBD is a programmable device with a large data buffer, which reads data from the CAMAC system and passes this data to the on-line computer through a data bus. The MBD and CAMAC system serves as the interface between the experiment and the computer.

5.2 The Neutron and Cosmic Ray Triggers

The trigger logic is diagrammed in Fig. 5.1 where, for simplicity, only one counter of each plane is considered. The notation "Sp_c" is used throughout the rest of this chapter, where "p" is the plane number, "c" is the counter number and "S" stands for scintillator. Raw signals from each of the neutron counters were brought into the counting house from the experimental area using BNC signal cables. Each of the 24 neutron counters had one phototube on each end which, for convenience of notation, are called sides "A" and "B."

In plane 1, phototube signals S13A and S13B were patched to a signal splitter having two outputs, as were the other counters. One output of each splitter was an input to CAMAC Analog to Digital Convertors (ADCs - LRS 2249) these were used to monitor the gains of the counters throughout the experiment. The other output of the splitters was

Figure 5.1

The trigger logic used in the experiment. The numbers inside the logic symbols refer to LeCroy Research Systems (LRS) logic units.

an input to Constant Fraction Discriminators (CFDs - ORTEC 934 or PHILLIPS 715) used to convert the analog raw signal into a NIM logic pulse. The NIM logic standard, ("FALSE" = 0 V and "TRUE" = -0.7 V) is used throughout experimental nuclear and high-energy physics. Most of the NIM logic modules used were manufactured by LeCroy Research Systems and can operate at rates approaching 100 Mhz.

The thresholds of the CFDs were set to -30 mV; raw signals with pulse heights $PH < -30$ mV passed the threshold and were converted to logic pulses. The output signals of the CFDs were called S13A, S13B for plane 1/counter 3 and similarly for the other counters and planes. In Table. 5.1 some Boolean algebra summarizes the forming of all important signals making up the trigger logic, and it should be compared with Fig. 5.1.

One output from each "A" CFD was combined with the corresponding "B" CFD as inputs to a Time Averager (TAV) device.^[58] This device essentially behaves as an AND gate but differs in that the timing of the output signal is not dependent upon the overlap of A/B timings. A coincidence output signal was produced which has an output timing determined by the average of the A and B times. An event which occurs at the center of neutron counter S13, for example, would have equal A and B timings, while an event near one end of the counter would have different relative timings. But, the TAV output would indicate that the event occurred at the center of the counter in both cases. This allowed very precise timing between neutron counters, and various cable

Table 5.1

Boolean algebra for the trigger logic. The column "point in logic" refers to points in Fig. 5.1.

<i>Point in logic</i>	<i>Signal</i>
Charged Veto	$A = AL + AR$
CFDs	$SpcA, SpcB$ ($p=1,2,3,4$; $c=1,2,3,4,5,6$)
TAVs	$Spc = SpcA \cdot SpcB$
Fanins/outs	$P1 = (S11+S12+S13+S14+S15+S16)$ $P2 = (S21+S22+S23+S24+S25+S26)$ $P3 = (S31+S32+S33+S34+S35+S36)$ $P4 = (S41+S42+S43+S44+S45+S46)$
Vetos	$V1 = (A + P1 \geq 2)$ $V2 = (A + P1 \geq 1 + P2 \geq 2)$ $V3 = (A + P1 \geq 1 + P2 \geq 1 + P3 \geq 2)$ $V4 = (A + P1 \geq 1 + P2 \geq 1 + P3 \geq 1 + P4 \geq 2)$
AND gates	$(P1=1) = P1 \cdot \overline{V1}$ $(P2=1) = P2 \cdot \overline{V2}$ $(P3=1) = P3 \cdot \overline{V3}$ $(P4=1) = P4 \cdot \overline{V4}$
STR 8	$STR\ 8 = (P1 \geq 2) + (P2 \geq 2) + (P3 \geq 2) + (P4 \geq 2)$
STR 9	$STR\ 9 = (P1=1) + (P2=1) + (P3=1) + (P4=1)$

lengths and delay boxes were used to insure that all counters were timed together to better than a fraction of a nanosecond. This was accomplished using a fast oscilloscope to adjust the timings to ± 1 ns. Time-of-flight was then used for timing adjustments to fractions of ns (Section 5.4). The noise level of each counter was reduced substantially at this point because of the coincidence that was formed between both ends of each counter.

Four linear fanins/outs were used to fan in all counters in each plane. For example, the linear fanin of the TAV outputs S11, S12, S13, S14, S15 and S16 formed the signal called "P1." Linear fanins/outs differ from logic fanins/outs in that the input/output signals are "OR-ed" in an analog sense. If two or more inputs to the fanins/outs are logical TRUE, the output will be a logical TRUE but also be an analog sum of the inputs.

Each plane had its own unique veto signal. For example, the plane-1 veto was a logical OR between the charged-particle veto counter A and a signal called "P1 $\geq$ 2" which was formed by using a discriminator with a threshold set above the level for logical TRUE. Because of the linear fanins/outs, when two or more counters in a plane fired, the corresponding ≥ 2 discriminators generated a veto signal and caused that plane to be vetoed. Plane 1 was vetoed if a charged particle was detected, or if two or more counters in plane 1 were hit. The vetos for planes 2, 3 and 4 were similar. Any plane which was hit vetoed all following planes. The trigger required a single counter in a

plane to be hit, and all following planes to be vetoed.

In the end, two triggers were formed which were called "STR 8" and "STR 9." The signal STR 8 was a "loose" trigger requiring two or more hits per plane and was used for cosmic ray data. STR 9 was intended to be a good event guaranteed to be due to one and only one counter firing in the first plane hit.

During the experimental setup, the trigger logic was checked out and timed using both a pulser and the production-type beam in the experimental area. STR 8 and STR 9 were inputs to the computer latch logic which conditioned the triggers STR 8 and STR 9 for the sampling of various events by the on-line computer.

5.3 Computer Sampling of Events

The logic was set up to allow the computer to sample two classes of events in real time. These were cosmic ray events, governed by EVT 8, and neutron events governed by EVT 9 as discussed in the last section.

Raw signals from the anode of each phototube were divided into two halves, using a passive 50 Ohm splitter. One output of each splitter was routed to CAMAC analog-to-digital converters (ADCs - LRS 2249) through 230 nanoseconds of delay cable. The second output of each splitter was an input to a constant-fraction discriminator (CFDs - ORTEC 934 or PHILLIPS 715). One discriminator output went to CAMAC time-to-digital converters (TDCs - LRS 2228A) after about 180 nanoseconds of delay. A second

discriminator output was sent to a TAV to eventually form the triggers as shown in the last section on the trigger logic. One output of each TAV was also sent to a CAMAC coincidence register (LRS 2341) after about 180 nanoseconds delay. The delay cables were necessary because the generation of the triggers EVT 8 and EVT 9 required some time. The logical OR of these triggers was used as gates for the ADCs, TDCs and the coincidence register.

The ADCs and TDCs required a certain amount of time ($\approx 60 \mu\text{sec}$) to complete their conversions. During this time, subsequent triggers had to be blocked. In addition, the ADCs, TDCs and coincidence register were read and cleared by the microVax computer via a BiRa Micro-programmable Branch Driver (MBD), an operation which took several microseconds. For these reasons, additional logic was set up to prevent any triggers while conversions or read operations were underway. A diagram of the computer gating logic is shown in Fig. 5.2.

In Fig. 5.2, two important signals are used as discriminator (LRS 365) inhibits. The STR 9 signal was vetoed if there was no polarized beam, data taking was not in progress or a spin change was under way (Normal beam to Reverse beam or Reverse to Normal). The veto signal for STR 9, $P \cdot BG \cdot \overline{SC} \cdot RUN$, was also used as scaler gates elsewhere in the logic. The signal $P \cdot BG$ denoted that the polarized beam was on, with BG representing a "beam gate" signal, and SC denoting a spin change in progress. All signals used in gating throughout the experimental hardware were generated by the

LAMPF gate generators using signals provided by the accelerator. Because the beam is pulsed with a duty cycle of at most 9% (for 800 MeV beam); the beam was "off" most of the time. The cosmic ray data was gathered during this beam-off period. The STR 8 signal was vetoed by $\overline{BG \cdot RUN}$ which specified that beam was on and data taking was not in progress.

Any STR 8 or STR 9 signals not vetoed then formed inputs to latched LRS 304 majority logic units to form the signals EVT 8 or EVT 9. EVT 8 or EVT 9 would then remain latched until vetoed by a special gate called "ALL BUSY." ALL BUSY, in turn, was made up by a logical OR of a 100-microsecond gate generated by a LRS 222 gate generator (started by either STR 8 or STR 9) and the signals EVT 4 and "Computer Busy." The 100-microsecond gate was necessary to allow time for the conversion processes in the CAMAC ADCs and TDCs. The signal EVT 4 specifies a special signal used in the experiment to flush the MBD buffer and reset various electronics. The Computer Busy signal was active whenever the computer was otherwise too busy to sample any events.

The signal "ANY EVENT," which was a logical OR between EVT 8 and EVT 9, was used for the gating of the ADCs, TDCs and the TAV coincidence registers. The computer-sampled events provided an invaluable monitor on the experiment because neutron counter gains and timings could be monitored in real time.

Figure 5.2
The computer latch logic used to hold off triggers while the computer was sampling various neutron and cosmic ray events. A latch 100 μ secs long allowed ample time for CAMAC ADCs, TDCs and latches to convert and be read and cleared by the computer.

5.4 The Time-of-Flight Logic

As discussed earlier, the accelerator timing is based upon a 201.25 Mhz reference frequency. With the buncher in operation, a 10 Mhz reference is available to experimenters. Early studies indicated that the 10 Mhz timing signal was not an adequately stable reference for obtaining time-of-flight information. Various factors, such as daily temperature changes, shifted the relative timings between the 10 Mhz and the 201 Mhz in an unusable manner. The 201 Mhz was a very stable timing reference, but a stable 10 Mhz reference was also needed for timing when the buncher was operating, to eliminate problems in discriminating against the low-energy component of the neutron beam (see Chapter 3). A stable 10-Mhz reference was obtained by using the MP-13 Synchronizer. This unit acted to "lock-on" the wandering 10-Mhz signal to the stable 201-Mhz and generate a stable 10-Mhz timing reference used for timing in the experiment.

The MP-13 synchronizer required two inputs and had one output: "pulse in," "synchronize in" and "synchronized out." Via a comparison procedure, the 201 rf pulses closest in time to the 10 Mhz pulses were output, and this stable 10-Mhz was used for timing in the experiment. With a QVT (charge/voltage/time analyzer) module, the buncher was monitored during data taking by using signals from the beam polarimeter (LBPOL) as a START and the synchronized 10 Mhz as a STOP. All of the timing logic is shown in Fig. 5.3a. The buncher was not monitored during the first part of

the experiment. This was due to lack of experience with the operating characteristics of the newly-constructed buncher. Studies showed that the amount of beam bunching systematically changed between the P^- source cycles Normal and Reverse, a problem which led to this logic being set up to monitor this effect.

The synchronized 10-Mhz signal was also used in conjunction with the WIDESTART signal of the FTDC logic for forming the time-of-flight (TOF) scaler gates. This is discussed in the next section. With a second MP-13 synchronizer, logic similar to that of the QVT logic was set up. The WIDESTART was used as a FTDC START, while a "synchronized" WIDESTART served as the FTDC STOP. Bin widths on the TOF FTDC were calibrated at 0.75 ns per bin. A typical histogram of the TOF FTDC output bins is shown in Fig. 5.3b. The quasi-elastic peak was contained within TOF FTDC bins 3, 4 and 5, which were fanned into a logic fanin/out. The resulting gate was called TOF, and the coincidence and FTDC scalers were gated on this.

The counts in each of the TOF FTDC bins were monitored throughout the experiment, and timing adjustments were made to any counters not properly "timed-in." These adjustments usually consisted of adding or subtracting 0.25 ns delay (≈ 5 cm of 50 Ω cable) before the CFDs, followed by an inspection of the TOF FTDC histogram.

Occasionally, timing shifts on the order of several ns occurred. In such cases the 10 Mhz reference had shifted in phase, relative to the more stable 201 Mhz reference. Adjustments to timing were then made at the 10 Mhz/201 Mhz MP-13 synchronizer.

Figure 5.3

Logic used to determine time-of-flight (TOF) in the experiment. The first section of the logic (a) was used to establish a stable 10-Mhz reference frequency. The beam buncher was also monitored at this point. TOF information was obtained by using the WIDESTART signal from the position FTDCs (next section) as a START for the TOF FTDC and the 10-Mhz reference as a STOP. A typical histogram of the TOF FTDC outputs is shown in (b); a fanin of output bins 3, 4 and 5 was used to define "good TOF" for detected neutrons.

5.5 The Fast Time-to-Digital Converter Logic

The relative timings between sides A and B of each neutron counter were used to locate the spatial position of an event in the xy plane. Fast Time-to-Digital Converters (FTDCs) were constructed at Argonne^[53] for the neutron counters because CAMAC TDCs require several microseconds to convert and could not handle the rates experienced during data taking. The FTDC's have a total of eight output channels, two of which represent underflows and overflows, leaving a total of six channels for extracting positional information. The bin widths for the position FTDC's were calibrated to 1.25 ns per bin. This corresponded to subdividing the length of each counter into ≈ 10 cm bins. This was checked experimentally by moving the hodoscope by small amounts while in a beam. FTDC histograms were accumulated, and the centroid positions in these histograms were related to the physical positions for each counter.

A diagram of the FTDC logic is found in Fig. 5.4. The signal from side A was used as a START, and the signal from side B was used as a STOP. The FTDC required 250 ns to convert relative timing information into positional information, and for this reason, a latch was set up to kill all successive inputs to allow for conversion of a previous FTDC START/STOP. The signals $P1=1$, $P2=1$ and $P3=1$ were logically ORed to provide a latch called "WIDESTART." This set off a 250-ns wide latch which, through feedback, vetoed any further WIDESTARTs for a period of 250 ns. This dead-time ("DT") was

used in several places in the logic and was important for the normalization of the data.

Each of the inputs to the FTDCs was taken from the CFD's outputs which were gated by DT and the particular plane corresponding to that counter. For example, the FTDC13 inputs were S13A and S13B, gated on $P1 \cdot \overline{DT}$. With this scheme, those FTDCs belonging to vetoed planes were not gated, and only the first plane hit was analyzed. Underflows and overflows occurred when there was no FTDC START or STOP, respectively.

The FTDC outputs correspond to positional information. These were scaled by CAMAC scalars (LRS 2551) which were read into the online computer during P^- spin changes. Two sets of scalars were used for the FTDCs: one set gated by DT only, and a second set gated by TOF. The TOF scalars were a subset of the DT scalars. The experiment used FTDCs for three planes, and with the two scalar gates DT and TOF, a total of 216 scalar channels were required.

5.6 The Coincidence Logic

The coincidence between perpendicular counters in adjacent planes would be used, in addition to the FTDC technique, for an additional measurement of $\Delta\sigma_L$. In the experiment, there were four planes of counters yielding three sets of coincidence planes. With six counters per plane, the intersection of counters in adjacent planes divided the xy plane into 36 bins with each bin having an area of 12.7×12.7 cm². It was realized

later in the experiment that only 24 bins could be used because the physical size of the scintillator limited the spatial extent of detection to exclude the outermost counters in each plane. In any case, the number of scalers procured were sufficient to scale a total of 36 bins/plane, including accidentals, and numbered a total of 216 scaler channels.

The experimental logic required to scale coincidence events was complex indeed. The required number of various logic units was greater by far than the number of units available at LAMPF. One type of logic unit required was the strobed-coincidence module. Having six inputs and a strobe input, the strobed-coincidence module outputs up to six separate coincidence signals (made between the inputs and the strobe). A total of six strobed-coincidence units were required for each coincidence plane (not including accidentals). Only six units were obtained at LAMPF, and six more came from Argonne National Laboratory. The total number of units available, therefore, was sufficient to scale only two sets of coincidence planes, or one coincidence plane with its corresponding accidentals.

This problem had a simple solution, once one understood the nature of the trigger logic which was discussed in section 5.2 of this chapter. The main trigger contains the information necessary for determining which plane first detected an event with all following planes vetoed. Using this feature, the logic in Fig. 5.5 was constructed.

In the trigger logic discussed earlier, the TAVs formed coincidences between both ends of each detector. These TAV signals were also used in the coincidence logic. In

Figure 5.5
The coincidence logic used in the experiment to determine the position of a hit in a plane. Inputs for the coincidence logic were taken from the TAVs used in the trigger logic (Fig. 5.1).

Fig. 5.5, note that a series of logic fanins/outs are used to combine signals from each of the detector planes. Six of the fanins/outs form the signals STR1...STR6 and are used to strobe the strobed-coincidence units located later in the logic. Each of the STR signals was designed to represent a fanin of a particular counter from each plane. For example, STR3 was a fanin of counters S13, S23 and S33, and similarly for all the other STR signals. An additional six fanins/outs were used in a similar manner for the formation of the signals CTR1...CTR6 which were the inputs for the strobed-coincidence units. As an example, CTR3 represents the fanin of counters S23, S33 and S43. Any coincidence made between the STR and CTR signals localizes the event. Lastly, six fanins/fanouts were used to form accidental CTR signals denoted by CTR1_{acc}...CTR6_{acc}.

Three banks of gated discriminators were used to properly form the STR and CTR signals so that no information about the event was lost. Although the following planes were vetoed by the trigger logic, the raw TAV signals in vetoed planes were not blocked. This had to be handled properly so that the STR and CTR signals did not represent a combination of counters from various planes. Because the trigger logic was set up to determine the first plane hit, the signals P1, P2 and P3 were used for gating the discriminators. The discriminators that were used in forming the STR signals were set up in such a way that signals from counters S11...S16 were gated on P1, counters S21...S26 gated on P2, and counters S31...S36 gated on P3. The logic operated in this fashion: if an event was detected in counters S13 and S23, the first bank of discriminators

insured that STR3 was due exclusively to S13 because the first plane hit contained S13 (S23 did not get through because the gate P2 is FALSE). In this manner the STR signals represented the first plane hit in the detector stack.

The second bank of discriminators was used to form the CTR signals and was set up in a similar manner; counters S21...S26 were gated on P1, counters S31...S36 were gated on P2, and counters S41...S46 gated on P3. Naturally, for coincidences between adjacent planes, the TAV signals for plane 2 were gated on P1 and similarly for the other planes. In this fashion the CTR signals represented the second plane hit. Using the previous example, with counters S13 and S23 hit, the signal CTR3 would be $CTR3 = S23 \cdot P1$. In this scheme, any coincidences between STR3 and CTR3 would represent the possible coincidences S13-S23, S23-S33 or S33-S43, but only one of these at any given time.

The third bank of discriminators was used for monitoring accidentals. The signals $CTR1_{acc} \dots CTR6_{acc}$ were formed in the same manner as the CTR signals discussed above with the exception that the inputs to the gated discriminators were delayed by 100 ns. With the buncher on, the beam had a RF period of 100 ns. Accidentals would, as an example, be represented by those events in which one plane might detect a real event while a different plane may detect a straggler from some previous beam pulse. In the coincidence logic, this would be misinterpreted as a real event. A measure of the accidental rate was made by simply scaling the coincidences between the STR and CTR_{acc} signals. Typically the accidental rate was found to be on the order of half a

percent.

Each of the strobed-coincidence units was gated on the STR signals. The 36 outputs, STR·CTR and their corresponding accidentals, STR·CTR_{acc}, were scaled by six banks of CAMAC scaler units (108 scalers for TRUES and 108 scalers for ACCIDENTALS). In the earlier example of a hit in counters S13 and S23, the strobed-coincidence outputs, STR3CTR3, STR3CTR3_{acc} were scaled by three banks of scalers each. Those scalers corresponding to coincidences in planes 1 and 2 had to be properly gated while all other scalers (i.e., planes 2 and 3 and planes 3 and 4) were inhibited. A total of 72 scalers were used for scaling P1P2 coincidences and accidentals and the other coincidence planes. The problem of how to gate the scalers, i.e. decoding the strobed-coincidence outputs to properly assign a given coincidence with the actual planes hit, was easily solved by special scaler gates. Those scalers representing P1P2 coincidences and accidentals were gated with the scaler gate P1·TOF, which not only decoded the strobed-coincidences outputs, but also required the events to be in good TOF (see section 5.4). The scaler gates P2·TOF and P3·TOF were used for scaling the other coincidence planes.

The coincidence logic, despite its complexity, had the desirable advantage that the number of electronic units became manageable and hence required a minimum of space. As the logic was put together, a pulser was used to test the reliability of the logic in order to determine whether it functioned as desired. For example, the inputs S13 and S23 were pulsed, and the scalers monitored. There was only one unexpected pulse out

of several thousand input pulses which ended up in the wrong place, i.e., scaler channel.

This is consistent with the nominal error rate for the logic modules alone.

CHAPTER 6 DATA ANALYSIS

This chapter will describe how the data were analyzed and how $\Delta\sigma_L(np)$ was finally obtained. The reduction of data is a major part of any experiment and it includes a search for any systematic effects in the data, an estimation of their effect, and sometimes a correction for them (e.g., error in normalization is a systematic effect not usually corrected for). In the following sections the steps used in the analysis are described.

6.1 Filtering the Data

During data taking, scaler ratios and asymmetries were plotted on a run-by-run basis to monitor the general state of the experiment. In addition, $\Delta\sigma_L$ was plotted every two minutes using accumulating scalers. The cumulative $\Delta\sigma_L$ generally showed the expected behavior (with decreasing error bars and asymptotically approaching a final value), but occasionally there were large variations outside the expected statistical errors. Since $\Delta\sigma_L$ is a small effect ($\epsilon(0) \approx 10^{-4}$), it was clear that in such cases large systematic effects had occurred in the data. For this reason, the data were analyzed offline on a spill-by-spill basis.

The scaler data were analyzed using the Q-system and a modified analyzer which calculated scaler ratios and various asymmetries for each spill. The analyzer began each run by reading in a list of "formulae" from an input file and creating "plot files" for that run on disk. As each spill was read back from magnetic tape, the various scaler ratios and asymmetries that were calculated by the analyzer for that spill were written to the plot files. These were then displayed on a graphics terminal to identify "bad spills" interactively, and additional statistical analysis could be performed.

Each scaler was identified by an index, and the formulae were written using these indices and the arithmetic operators +, -, /, and *. For example, if the ratio of scalers #312 and #423 was desired, the formula used was simply "312/423". This feature of the analyzer had the advantage that the scaler calculations were not "hard-wired" in to the analyzer. This allowed for different ratios to be calculated by simply modifying the formulae. The bad spills which were identified on a statistical basis were rejected in the final pass through the data.

The scaler quantities which proved most important for this spill-by-spill analysis were

- SEM/CLK (proton monitor)
- M4/SEM (front monitor)
- QVT (buncher) fraction

- TOF (time-of-flight) fraction.

Fig. 6.1 was taken from a typical run and shows each of these ratios on a spill-by-spill basis. The TOF fraction was the most useful in practice. This was the fraction of detected neutrons having good time-of-flight. Those runs (or spills) with a small TOF fraction were rejected because, in this case, $\Delta\sigma_L$ might be contaminated with low energy neutrons. The QVT fraction was used in the same manner but it was more sensitive to the quality of the proton beam bunching. Usually the TOF fraction was an indicator of the buncher performance as well, but more often it pointed out timing shifts in the TOF logic.

Because the proton beam was monitored rather than the neutron beam, it was important that the ratio of neutrons to protons produced in the beam was a constant within each run. The scaler ratio M4/SEM was an indicator of this because M4 sampled the neutron beam, and the SEM was the proton beam monitor.

The scaler ratio SEM/CLK monitored the proton beam intensity. One of the motivations for monitoring the buncher in an active way came about when it was observed that large asymmetries in beam intensity affected the quality of beam bunching. As noted in an earlier chapter, some data were taken with the buncher not monitored (720 MeV and 788 MeV).

Figure 6.1

Example of the plot ratios used to filter the data, taken from a typical run. The ratio SEM/CLK was proportional to the proton beam intensity. The second ratio, M4/SEM, was proportional to the number of neutrons produced per incident beam proton. The buncher performance (QVT) was monitored with the fraction of beam protons bunched within 100 ns-spaced pulses. Timing shifts in the experimental logic were found using the time-of-flight (TOF) fraction. This was a measure of the fraction of neutrons with good TOF.

Some statistical criterion was needed for filtering the data. Eventually, Chauvenet's criteria was adopted, [59, pp. 170-1] It states

... any reading of a series of n readings shall be rejected when the magnitude of its deviation from the mean of the series is such that the probability of occurrence of all deviations that large or larger does not exceed $1/2n$.

The criterion probabilities were interpolated from values in a table in Ref. [59, pg. 319]. This was found to be a robust recipe for filtering the scaler data because of its conservative nature in identifying "outliers." No more than 10% of the data was rejected in this manner. The criteria failed for runs in which many spills showed large variations from their mean. In such cases, the outliers were identified interactively, but occasionally the whole run was rejected.

A list of the rejected spills for each run was kept in a file, and in the final pass through the data, the analyzer used this file to reject the bad spills. The filtered scalers were then stored on disk and tape. Using these scalers $\Delta\sigma_L(np)$ was extracted and correlation studies (systematics) were conducted.

6.2 Data Normalization

From the final scalers, spin asymmetries were calculated for each bin in the hodoscope which corresponded to different momentum transfers (t). These asymmetries were then extrapolated to 0° to find $\Delta\sigma_L(np)$. The momentum transfers (t) for each bin were

taken to be the unweighted mean t for that bin:

$$\langle t \rangle = \frac{1}{A} \int t dA$$

with $t = (p_1 - p_3)^2$.

In the relations above, p_1 and p_3 are the four-momenta for the incident and scattered neutrons, respectively. The mean t was found to be a reasonable choice, as it was found that the extrapolations to $t = 0$ were insensitive to reasonable variations in how t was defined.

The neutron data (N) were first normalized with the proton beam monitor (SEM). Deadtime (DT) corrections were made with the front monitor ($M4$) gated by both the RUN gate and the DT gate. The deadtime for the experiment was due, primarily, to the FTDCs which was on the order of 250 ns. The normalizer for the neutron data was

$$\mathcal{N}_{beam} = SEM_{run} \cdot \frac{M4_{dt}}{M4_{run}}$$

where the subscripts denote the gating scheme used. The ratio $M4_{dt}/M4_{run}$ was an estimate of the experimental "livetime."

If $N^\pm$ represents the number of detected neutrons and $\delta N^\pm$ the statistical uncertainties for spins parallel and anti-parallel, then $T^\pm$ represents normalized neutrons

$$T^\pm = \frac{N^\pm}{\mathcal{N}_{beam}^\pm}$$

with the statistical uncertainty in $T^\pm$ given by

$$\delta T^\pm = T^\pm \left[\left(\frac{\delta N^\pm}{N^\pm} \right)^2 + \left(\frac{\delta \mathcal{N}_{\text{beam}}^\pm}{\mathcal{N}_{\text{beam}}^\pm} \right)^2 \right]^{1/2}.$$

The statistical errors for the neutrons and beam are added in quadrature because N and $\mathcal{N}_{\text{beam}}$ are independent measurements of different events and therefore were not correlated. In terms of the normalized neutrons T , the asymmetries for each bin in the hodoscope are calculated by

$$\epsilon = \frac{T^+ - T^-}{T^+ + T^-}$$

with a statistical uncertainty given by

$$\delta \epsilon = \frac{1}{T^+ + T^-} \left[(\delta T^+)^2 (1 - \epsilon)^2 + (\delta T^-)^2 (1 + \epsilon)^2 \right]^{1/2}.$$

Here, T^+ and T^- are independent measurements so their uncertainties are added in quadrature.

The number of detected neutrons N followed a Poisson distribution (counting statistics), and therefore the uncertainties were $(\delta N)^2 = N$. In the coincidence logic, accidentals were scaled, and the recipe changes somewhat:

$$\begin{aligned} N_{\text{true}} &= N - N_{\text{acc}} \\ (\delta N)_{\text{true}}^2 &= (N + N_{\text{acc}}). \end{aligned}$$

In the relations above, N_{true} was the number of detected neutrons corrected for accidentals, and δN_{true} a conservative estimate of the statistical uncertainty. The accidental rates in the coincidence logic were always less than a few percent during the experiment.

The uncertainty in the normalizer $N_{beam}^{\pm}$ was a different matter because $M4_{dt}$ was a subset of $M4_{run}$ and therefore these were not independent measurements. In this case, $M4_{dt}$ followed the binomial distribution

$$Prob(m) = \frac{M!}{m!(M-m)!} p^m (1-p)^{M-m}$$

where $p = \frac{E(M4_{dt})}{M4_{run}}$.

The probability of getting $m = M4_{dt}$ counts, given $M = M4_{run}$ counts, is given by $Prob(m)$. The probability p is the ratio of the expected number of counts $E(M4_{dt})$ to the total counts in $M4_{run}$. If we define

$$r = \frac{M4_{dt}}{M4_{run}},$$

then the statistical uncertainty on r is given by

$$\sigma_r^2 = \frac{1}{M4_{run}^2} E(m^2 - \langle m \rangle^2)$$

where $E(m^2 - \langle m \rangle^2)$ is just the variance in m , $Var(m)$. Since the binomial distribution has a statistical variance $Var(M4_{dt}) = M4_{run}p(1-p)$, we find that

$$\sigma_r^2 = r^2 \left(\frac{1}{M4_{dt}} - \frac{1}{M4_{run}} \right).$$

The uncertainty in the *SEM* was not included in the error analysis for the beam normalizer because it stemmed from a systematic effect. The *SEM* had a substantial "dark current" associated with it which introduced systematic effects that were much

larger than any statistical fluctuations. The scaled *SEM* counts can be thought of as the sum of two parts: those counts which are beam related (SEM_{beam}); and those counts which are generated spuriously by the *SEM* digitizer (SEM_{dark}). The dark current was present even in the absence of beam, and it was estimated from the RUN and Quench scalers to be about 1 count/sec. This is to be compared with $10^1 - 10^2$ counts/sec of signal.

6.3 The Target Constant

$\Delta\sigma_L(np)$ was related to the extrapolated asymmetry $\epsilon(0)$ through the equation

$$\epsilon(0) = \frac{P_n P_t}{2A} \Delta\sigma_L.$$

The density ρ_F that appears in the target constant $A = (\rho_F L)^{-1}$ can be rewritten as

$$\rho_F = \frac{N_A \rho n f}{M},$$

where ρ , n , f and M are defined below (N_A is Avogadro's number). The target material was propanediol and had a mass density ρ at low temperatures (0.5 K) of 1.12 gm/cm³. The number of free protons (polarizable) for each molecule of propanediol was represented by n and its value was approximately 8. The packing fraction f took into account that the actual density of protons was reduced because the target volume was packed with tiny spheres of propanediol (for random packing, one expects $f \approx 0.64$). The molecular weight of the target was M ; 76.09 gm/mole. One other factor that should

be included is the fraction of the beam actually hitting the target. Since this fraction is known only as a lower limit ($\approx 3\%$), it was treated as contributing to the systematic normalization error.

At the conclusion of each run cycle, the target was emptied and the target material was weighed on a laboratory balance. When each target was emptied, a liquid-nitrogen bath was used to maintain the beads of the target material at low temperatures because the propanediol absorbs moisture at warm temperatures. A small portion of the target material was then chemically analyzed for hydrogen content. This yielded the fraction of target material that was polarizable. The packing fraction was estimated from the mass and volume of the target at low temperatures using:

$$f = \frac{\text{mass/volume}}{\rho} \Big|_{0.5K}$$

The linear expansion of teflon at low temperatures^[60] was used to estimate the volume of the target at liquid helium temperatures. The linear dimensions of a teflon object shrinks by 1.1 % in going from room temperature to that of liquid nitrogen. An additional shrinkage of 0.1 % results in going from the temperature of liquid nitrogen to that of liquid helium. The corresponding change in the room temperature volume is then -3.6 % with a relative error of 10 %. The estimated uncertainty in the polarization of the proton target was estimated by D. Hill^[60] (ANL) to be $\pm 3.8\%$. He also calculated other target normalization corrections which were used in calculating $\Delta\sigma_L$.

A list of the target parameters for both years of running is given in Table 6.1. Note that in Table 6.1 the packing fraction differs by 12% between the two years of data taking. This is due to the difference in the target masses which were measured at the end of 1987 and 1988.

Table 6.1
Target parameters for the experiment.

	1987	1988
mass of target	154.3	137.2 gm
hydrogen content by weight	9.95%	10.63%
packing fraction	0.659 ± 0.029	0.586 ± 0.026
target constant	1823 ± 80	1920 ± 85 mb

- target material: propanediol
- molecular weight: 76.09 gm/mole
- low temperature density: 1.12 gm/cm^3
- length of target cell: 12.5 cm
- volume of target cell at low temp: $209 \pm 9 \text{ cm}^3$

0.4 The Sign of $\Delta\sigma_L$

The directions of the magnetic fields of the spin-precession magnets, Lorraine and Castor, were determined by a current-carrying wire. A current-carrying wire placed within a magnetic field experiences a force upon it which is proportional to the amount of current passing through the wire. The direction of the force depends on both the direction of current flow and the direction of the magnetic field. These magnet polarities were needed to determine the spin direction for the neutrons at the proton target. The spin directions for the neutrons entering the experimental area had been known from earlier work (see Ref. [49]), but the magnets had been connected afresh before this experiment was set up.

The proton beam polarization at the LD₂ target was mostly L-type, and the polarization of the neutron beam was proportional to that of the proton beam through the spin transfer functions K_{LL} and K_{NN} . In terms of the proton beam polarization components, the neutron beam polarization components are

$$\vec{P}_n \cdot \hat{L} = K_{LL} \vec{P}_p \cdot \hat{L} = K_{LL} P_p \cos \theta_p$$

$$\vec{P}_n \cdot \hat{N} = K_{NN} \vec{P}_p \cdot \hat{N} = K_{NN} P_p \sin \theta_p$$

where θ_p is the proton spin direction at the LD₂ target. It was shown earlier that K_{LL} was greater than K_{NN} , so it was advantageous to keep θ_p small. This led to a neutron beam exiting the LD₂ target that being primarily L-type. Because K_{LL} is negative, the

neutrons are polarized oppositely to the protons.

The magnitude of the beam polarization for the neutrons is needed to obtain $\Delta\sigma_L$ and is given by an effective spin transfer function called K_{eff} :

$$K_{eff} = (K_{LL}^2 \cos^2 \theta_p + K_{NN}^2 \sin^2 \theta_p)^{1/2}.$$

The magnitude of the neutron beam polarization is just $|P_n| = K_{eff}|P_p|$. Table 6.2 gives the various proton spin directions and K_{eff} for the data taken at different energies. The pairs of angles listed for each energy correspond to the proton spin directions during the Normal and Reverse spin states, respectively. The direction along the beam momentum is taken to be 0° . The values of K_{LL} and K_{NN} at the energies 568 and 720 MeV in the table above were interpolated because there is no data at those energies.

Table 6.2

Spin directions for the protons at the LD₂ target and corresponding effective spin transfer K_{eff} .

T_{lab}	θ_{p,LD_2} (degrees)		K_{LL}	K_{NN}	K_{eff}
484	182.0	2.0	-0.499 ± 0.010	-0.160 ± 0.013	0.499
568	177.0	-3.0	-0.568	-0.119	0.567
634	173.0	-7.0	-0.637 ± 0.010	-0.079 ± 0.015	0.632
720	167.8	-12.2	-0.620	-0.078	0.606
788	163.6	-16.4	-0.604 ± 0.016	-0.078 ± 0.021	0.580

After the LD_2 target, the neutrons passed through the magnetic field of LBBM6 which deflected the non-interacting proton beam into a beam dump. This magnet also precessed the spins of the neutrons away from pure L-type by about 45° in a clockwise fashion. The rate of change in the spin direction $\vec{\sigma}$ for a particle with magnetic moment μ moving perpendicularly to an external magnetic field $\vec{B}$ is

$$\frac{d\vec{\sigma}}{dt} \propto \vec{\mu} \times \vec{B}.$$

The direction the dipole rotates is easily found from the right-hand rule (neutrons have a negative magnetic moment). A diagram showing the spin directions of the neutron beam in the experimental area is shown in figure 6.2 for the NOR and REV orientations of Castor. The field direction for Lorraine was down. This rotated the neutron spins clockwise. The current was adjusted (with Castor turned off) so that there was a maximum Up/Down asymmetry at the neutron polarimeter NPAN. The spins of the neutrons leaving the field of Lorraine were then oriented at -90° ($+90^\circ$) when the proton spin state was Normal (Reverse). The field direction for Castor was pointed up (down) when Castor was NOR (REV). This rotated the neutron spin counter-clockwise (clockwise). The current for Castor was adjusted so that Up/Down and Left/Right asymmetries vanished at NPAN which indicated that the spins were L-type and oriented at either 0° or 180° . The actual spin direction depended on the field direction of Castor, and the spin state of the protons.

Figure 6.2

Spin directions for the neutrons in the experimental area. When the P^- beam was in the "Normal" polarization state, the neutrons were either +L-type (Castor NOR) or -L-type (Castor REV) at the proton target.

The magnetic field of the polarized proton target (PPT) was found to be pointing downstream along the direction of the neutron beam. When the proton polarization was positively (POS) enhanced, the proton spins pointed along the magnetic field direction. The proton spins were pointed against the field direction when the proton polarization was negatively (NEG) enhanced.

The spin asymmetries calculated from the scalars were the asymmetries in Normal and Reverse scalars. That is,

$$\epsilon = \frac{\text{NOR} - \text{REV}}{\text{NOR} + \text{REV}}$$

were based on the accelerator signals (proton beam polarization states). $\Delta\sigma_L$ was related to the extrapolated asymmetry at $t = 0$ by

$$\begin{aligned}\Delta\sigma_L &= \frac{2A}{P_n P_t} \cdot \epsilon(0) \\ &= \text{sign}(\text{Cas}, \text{Tar}) \times \frac{2A}{P_n P_t} \epsilon(0)\end{aligned}\tag{6.1}$$

where the superscript + (-) represent spins parallel (anti-parallel). From the arguments presented above, the correct signs for the different combinations of Castor and PPT combinations were

$$\text{sign}(\text{NOR}, \text{POS}) = +1$$

$$\text{sign}(\text{NOR}, \text{NEG}) = -1$$

$$\text{sign}(\text{REV}, \text{POS}) = -1$$

$$\text{sign}(\text{REV}, \text{NEG}) = +1.$$

In the correlation fits discussed in the next section, these signs were incorporated into the beam and target polarizations (P_n, P_t) for each run.

6.5 Correlation Studies

Correlation studies were conducted on the data to look for any systematic effects in the measurement of $\Delta\sigma_L$. One looks for systematic effects in the beam normalizer and in the number of detected neutrons. Several candidates showing correlations with

$\Delta\sigma_L$ presented themselves during the data taking. The results of the off-line studies are presented in this section.

The most important effect was due to an asymmetry in the buncher "fraction." This was realized in the early stages of data acquisition, but unfortunately data presented here for 720 MeV and 788 MeV were gathered without any of the necessary buncher information. This was due to our lack of experience with the (then) new buncher and its direct systematic effect on the data. The buncher fraction was an important quantity to monitor because of the following reason: not all of the monitors used in the experiment were sensitive to the time structure of the beam, while the number of detected neutrons depended on the fraction of protons bunched within a single rf cycle amongst twenty rf cycles.

It was therefore necessary to correct for the actual fraction of the proton beam responsible for the neutrons detected through time-of-flight (TOF). Instead of "hard-wiring" this correction into the analyzer, correlation studies were done so that the size of the effect could be compared with a model. In the model, f^+ and f^- were taken to be the fractions of beam bunched within the single rf cycle (f was typically 90–95%) for beam states Normal and Reverse. The corrected beam normalizer is then

$$\mathcal{N}_{beam}^{\pm} = SEM_{run}^{\pm} \cdot \frac{M4_{dt}^{\pm}}{M4_{run}^{\pm}} \cdot f^{\pm}.$$

The "corrected" spin asymmetry (in terms of the extrapolated asymmetry $\epsilon(0)$) is

$$\epsilon_{\text{corr}} = \frac{\epsilon(0) - \epsilon_f}{1 - \epsilon(0)\epsilon_f}$$

with $\epsilon_f = \frac{f^+ - f^-}{f^+ + f^-}.$

The product $\epsilon(0)\epsilon_f$ was typically 10^{-5} or smaller, so the corrected asymmetry was taken to be $\epsilon(0) - \epsilon_f$.

The systematic effects of the buncher were studied using a slightly modified expression involving $\Delta\sigma_L$ as an independently fitted parameter. From Eqn. 6.1, one can construct the function

$$\epsilon(0) = \frac{P_n P_t}{2A} \Delta\sigma_L + \alpha \epsilon_f \quad (6.2)$$

where $\Delta\sigma_L$ is proportional to ϵ_{corr} instead of $\epsilon(0)$. The parameter α was included to be determined from correlation fits to the data, as our definition of f may have been somewhat in error (e.g., the time-of-flight gate for the neutrons may have actually wanted a different fraction). The expected value for α was $\alpha = 1$ and was confirmed from fits to the data. Note that correlation fits to Eqn. 6.2 give a $\Delta\sigma_L$ corrected for buncher asymmetries.

There were three additional effects which might have biased the measurement of $\Delta\sigma_L$. Accidentals were a potential problem because the FTDC logic was not set up to sample beam-related accidentals. The coincidence logic was set up to sample accidentals, however, and the accidental rate in the coincidence scalers could be used as a measure

of the accidental rate for the FTDC scalers. A second systematic effect, which was related to beam accidentals, was the detection of cosmic rays within the TOF window. It was possible that a small number of cosmic rays having good TOF were detected by mistake and scaled. Naturally, the fractional number of cosmic rays detected decreases with increasing beam intensity. We can correct for these two effects by subtracting them off:

$$N_{corr} = N_{det} - ACC$$

$$\text{with } ACC = C \cdot T + acc_{beam}$$

where C is the cosmic ray rate (assumed constant) and T is the duration of a data run. The quantity acc_{beam} was taken to be beam related accidentals; a rate effect that increases with beam intensity. As mentioned before, the accidental rate for the coincidence logic was on the order of one percent, and this can be taken as an indication of the accidental rate for the FTDCs.

The dark current of the *SEM* biased the normalization of the neutron data and became important if there was an asymmetry in the beam intensity between spin states Normal and Reverse. The dark current rate was deduced to be about 1 count/sec, and if this rate is assumed constant with beam intensity the *SEM* counts can be corrected by subtracting the estimated amount of dark current:

$$SEM_{corr} = SEM - DC \cdot T.$$

One can then form the ratio of normalized neutrons for spin states Normal and Reverse, corrected for the systematic effects mentioned above:

$$r_{corr} = r \cdot \left(\frac{1 - (CT^+ + acc_{beam}^+)/N^+}{1 - (CT^- + acc_{beam}^-)/N^-} \right) \cdot \left(\frac{1 - DC \cdot T^-/SEM^-}{1 - DC \cdot T^+/SEM^+} \right).$$

The asymmetry in r_{corr} is then (to first order)

$$\epsilon = \epsilon_{corr} + \left(C \left(\frac{T}{N} \right) - DC \left(\frac{T}{SEM} \right) \right) \epsilon_{SEM/T} - \left(\frac{acc}{N} \right) \epsilon_{acc/N}.$$

The ratios N/T and SEM/T were taken to be average neutron and proton beam intensities (the asymmetries in the neutron and proton beam intensities are taken to be equal) and acc/N a measure of the accidental rate. Note that the effects due to the accidental detection of cosmic rays and the dark current of the SEM are not distinguishable in this form. Eqn. 6.2 is then generalized to

$$\epsilon(0) = \frac{P_n P_t}{2A} \Delta\sigma_L + \alpha\epsilon_f + \beta \left(\frac{T}{SEM} \right) \epsilon_{SEM/T} + \gamma \left(\frac{acc}{N} \right) \epsilon_{acc/N} \quad (6.3)$$

where the neutron intensity is taken to be proportional to the proton intensity. The parameter β was to be determined as was γ (which should be about 1). To complete the model hypothesized to correct $\Delta\sigma_L$ for systematic effects, a discussion of a strange effect that was found in the data is in order.

When correlation fits to Eqn. 6.3 were done, it was found that they were in general quite bad. The total chi-squared values (χ^2) were large with respect to the number of degrees of freedom for the fits. This indicated a bad model. The data were then

binned according to the orientations of Castor and the PPT, and each of the four resulting data sets were fit separately. It was found that the total χ^2 for each bin then became quite reasonable except that a strange property of the data became evident. It was found, that for those orientations in which Castor and the PPT were either (NOR,POS) or (REV,NEG), the values of $\Delta\sigma_L$ obtained differed by more than three standard deviations from those obtained for those orientations of Castor and the PPT that were either (NOR,NEG) or (REV,POS) which were consistent with zero.

This behavior in the data seemed consistent with the presence of an overall offset constant. Adding a constant term to the four extrapolated asymmetries, binned according to the orientation of Castor and the PPT, resulted in the four $\Delta\sigma_L$'s becoming more consistent with each other. The source of this constant has not yet been determined, but it was included in the correlation fits.

The presence of an overall constant term implies that the systematic effect may have originated in the proton beamline, as it was independent of the orientation of Castor and the PPT. Studies were conducted to determine the source of this offset term. The beam-steering data were investigated first because a likely candidate might have been a systematic change in the beam-spot width. Although the servo-steering maintained a spot position at the LD₂ target constant to better than 0.1 mm, it could not control the beam width. As mentioned in Section 2.5 of Chapter 2, Monte-Carlo studies indicated that the experiment was extremely sensitive to beam-width changes. Data from the

beam steering system were therefore included in the fits. However, it was found that the constant term was not correlated with changes in the width of the beam spot.

The final correlation fits were carried out using the following:

$$\epsilon(0) = \frac{P_n P_t}{2A} \Delta\sigma_L + \alpha \epsilon_f + \beta \left(\frac{T}{SEM} \right) \epsilon_{SEM/T} + \gamma \left(\frac{acc}{N} \right) \epsilon_{acc/N} + \delta, \quad (6.4)$$

The fit parameter δ represents the constant offset term. In the next section, the results of these fits are given for each energy.

6.6 Results of the Correlation Studies

The basic procedure followed in obtaining the final results was straightforward. The variables in Eqn. 6.4 (P_n , P_t , ϵ_f , T/SEM and $\epsilon_{SEM/T}$, and acc/N and $\epsilon_{acc/N}$) were found from the final scalars. The extrapolated asymmetries for each plane were then subdivided and fit according to the "scheme" used (FTDC or coincidence). The procedure began with the five fit parameters; $\Delta\sigma_L \cdot K_{eff}/2A$, α , β , γ and δ . A least-squares fit^[61] was done, and estimates of the parameters and their corresponding uncertainties were then obtained. The uncertainties for the estimates were taken from the error matrix. A total χ^2 for each fit was also calculated.

The number of fit parameters would then be adjusted, based on the outcome of each fit. The "three-sigma" rule was used, which meant that if any of the parameter estimates were consistent with zero to within three standard deviations, that parameter

was dropped from the fit (with the exception of $\Delta\sigma_L$). In practice, the removal of non-significant parameters generally lowered the total value of χ^2 and improved the fit. The data were then fit again, and the fit was continually adjusted until only significant parameters were left (see Tables 6.3, 6.4, 6.5, 6.6 and 6.7 for which parameters were deemed significant).

Any run that had a contribution to the total χ^2 of more than 9 (dimensionless) was investigated for the possibility of excluding that run from the fits. The experiment logbooks were searched for valid reasons for excluding the run, such as "steering turned off" or "TOF shifted". In practice, only a small number of runs were excluded from the fits for these reasons.

The results of the fits are presented on the following pages in tabular form (a blank entry for a fit parameter means that parameter was not significant in the fit) for each energy. The square root of the χ^2 per degree of freedom, called the "reduced chi squared," are listed as indicators of the quality of the fits. The target constant for each energy is also listed, as well as is the effective spin transfer K_{eff} (from Table 6.2). The data analyzed were from 484, 720 and 788 MeV runs taken in 1987, and the 568 and 634 MeV runs taken in 1988. As more data were taken, additional planes of counters were set up and incorporated into the experimental hardware. This is reflected in the number of planes listed in each table.

Table 6.3 lists fit results at 484 MeV for two FTDC planes and two coincidence

planes. The bottom portion of Table 6.3 reflects a natural division in the data because the counters were moved further away from the proton target. The counters were also moved 2.5-in beam right and raised 2.5-in up. The most important systematic effects found were the buncher and beam intensity, which have been corrected for by the fits.

At 568 MeV, the most important effect is that of the buncher (see Table 6.4). The data taken at 634, 720 and 788 MeV have the offset term included in the fits. Note that for all of the energies in which the buncher was monitored, $\alpha \approx 1$ is the dominant systematic effect.

The final $\Delta\sigma_L(np)$ results were found by taking a weighted average of the estimates listed in Tables 6.3, 6.4, 6.5, 6.6 and 6.7. The recipe used was

$$\mu = \frac{1}{w} \sum_{i=1}^N w_i y_i$$

where $w_i = \frac{1}{\sigma_i^2}$ and $w = \sum_{i=1}^N w_i$

which is valid for data with a common mean μ . The quantity σ above is the uncertainty for each data point. In order to obtain a more realistic estimate of the overall statistical uncertainty in $\Delta\sigma_L$, the uncertainties for each of the $\Delta\sigma_L$'s were multiplied by the square root of the reduced χ^2 before averaging the data together. Table 6.8 summarizes the $\Delta\sigma_L$'s found for both the FTDC scheme and the coincidence scheme. While the agreement is good at four energies, the FTDC and coincidence results in Table 6.8 differ by more than three standard deviations at 788 MeV. It is not clear why there is such a

difference. At 788 MeV the coincidence scheme was set up for only one set of coincidence planes (P1-P2), with the coincidence scalers gated on time-of-flight and deadtime. The deadtime scalers gave an extrapolated asymmetry $\epsilon(0)$ which was more consistent with that from the FTDC scalers, and it may have been that the time-of-flight gate was not set up properly for the coincidence scalers.

In the next chapter, some additional comments are made concerning the results of this experiment. Also, a comparison is made to the revised Argonne data and previous measurements made with free polarized neutron beams.

Table 6.3
Results from the 484 MeV correlation studies

484 MeV

target constant (A) = 1823 mb
spin transfer (K_{eff}) = 0.499

scheme	$\Delta\sigma_L \left(\frac{K_{eff}}{2A} \right)$	α	β	γ	δ	$\sqrt{\chi^2/df}$
P1	-0.00142 (0.00040)					1.47
P2	-0.00096 (0.00031)	0.79 (0.07)				0.95
P1P2	-0.00175 (0.00052)					1.59
P2P3	-0.00136 (0.00061)	0.97 (0.13)				1.14
P1	-0.00143 (0.00029)		0.00177 (0.00023)			1.40
P2	-0.00055 (0.00028)		0.00116 (0.00016)			1.43
P1P2	-0.00187 (0.00031)		0.00145 (0.00020)			1.40
P2P3	-0.00109 (0.00044)		0.00144 (0.00029)			1.38

Table 6.4
Results from the 568 MeV correlation studies

568 MeV

target constant (A) = 1920 mb

spin transfer (K_{eff}) = 0.567

scheme	$\Delta\sigma_L \left(\frac{K_{eff}}{2A} \right)$	α	β	γ	δ	$\sqrt{\chi^2/df}$
P1	-0.00091 (0.00030)	0.806 (0.068)				1.20
P2	-0.00034 (0.00037)	0.999 (0.107)				1.29
P3	-0.00062 (0.00033)	1.269 (0.083)				1.36
P1P2	-0.00095 (0.00043)	0.842 (0.097)				1.06
P2P3	-0.00053 (0.00063)	1.264 (0.148)				1.09
P3P4	-0.00089 (0.00061)	1.520 (0.148)				1.21

Table 6.5
Results from the 634 MeV correlation studies

634 MeV

target constant (A) = 1920 mb

spin transfer (K_{eff}) = 0.632

scheme	$\Delta\sigma_L \left(\frac{K_{eff}}{2A} \right)$	α	β	γ	δ	$\sqrt{\chi^2/df}$
P1	-0.00020 (0.00025)	1.183 (0.080)			-0.00065 (0.00012)	1.30
P2	-0.00079 (0.00029)	1.246 (0.092)			-0.00084 (0.00014)	1.31
P3	-0.00038 (0.00034)	1.502 (0.107)			-0.00072 (0.00017)	1.24
P1P2	-0.00011 (0.00034)	1.107 (0.104)			-0.00068 (0.00016)	1.10
P2P3	-0.00097 (0.00046)	1.185 (0.145)			-0.00056 (0.00023)	1.27
P3P4	-0.00074 (0.00066)	1.134 (0.138)			-0.00091 (0.00022)	1.12

Table 6.6
Results from the 720 MeV correlation studies

720 MeV

target constant (A) = 1823 mb
spin transfer (K_{eff}) = 0.606

scheme	$\Delta\sigma_L \left(\frac{K_{eff}}{2A} \right)$	$\alpha^{(a)}$	β	γ	δ	$\sqrt{\chi^2/df}$
P1	-0.00131 (0.00023)			1.42 (0.48)	0.00039 (0.00011)	1.16
P2	-0.00141 (0.00033)					1.29
P1P2	-0.00134 (0.00029)			3.64 (0.53)	0.00076 (0.00014)	1.13

^(a)Buncher was not monitored.

Table 6.7
Results from the 788 MeV correlation studies

788 MeV

target constant (A) = 1823 mb
spin transfer (K_{eff}) = 0.580

scheme	$\Delta\sigma_L \left(\frac{K_{eff}}{2A} \right)$	$\alpha^{(a)}$	β	γ	δ	$\sqrt{\chi^2/df}$
P1	-0.00140 (0.00023)		0.00052 (0.00010)			1.36
P2	-0.00097 (0.00024)				-0.00055 (0.00012)	1.28
P1P2	+0.00026 (0.00031)				-0.00074 (0.00015)	1.38

^(a)Buncher was not monitored.

Table 6.8
 $\Delta\sigma_L(np)$ found for both the FTDC and coincidence scheme

$\Delta\sigma_L(np)$

Energy	scheme	$\Delta\sigma_L(\text{mbarns})$
484	FTDC	-7.46 ± 1.41
	COINC	-11.55 ± 2.15
568	FTDC	-4.60 ± 1.64
	COINC	-5.66 ± 2.29
634	FTDC	-2.62 ± 1.29
	COINC	-2.54 ± 1.76
720	FTDC	-8.05 ± 1.36
	COINC	-8.06 ± 1.97
788	FTDC	-7.42 ± 1.38
	COINC	$+1.63 \pm 2.69$

CHAPTER 7 CONCLUSIONS

7.1 Normalization Error

A plot of the results of this experiment is shown in Fig. 7.1 together with Arndt's phase-shift prediction.^[62] The phase-shift solution used was "SM89". The data are in agreement with the prediction which implies that probably there will be little change in any of the values of the phase shifts if this data were added to the data base. The purpose of this chapter is to make some comments about the data, including previously published measurements of $\Delta\sigma_L(np)$.

The uncertainties given for $\Delta\sigma_L(np)$ in Table 6.8 (last chapter) are statistical only; that is, the uncertainties were obtained from the error matrix for each fit and scaled by the square root of the reduced χ^2 per degree of freedom. In order to complete the estimates for the uncertainties in $\Delta\sigma_L(np)$, the overall normalization error was also estimated. Let

$$\Delta\sigma_L = \frac{2A}{K_{eff}} \cdot \frac{\epsilon}{P_b P_t},$$

where $2A/K_{eff}$ is the normalizer and $\epsilon/P_b P_t$ is the extrapolated asymmetry corrected

Figure 7.1

Values of $\Delta\sigma_L(np)$ found in this experiment. The error bars are statistical only.

for proton beam and target polarizations. The uncertainty in $\Delta\sigma_L(np)$ is found by

$$\begin{aligned}\delta^2(\Delta\sigma_L) &= \delta^2\left(\frac{2A}{K_{eff}} \cdot \frac{\epsilon}{P_b P_t}\right) \\ &= \epsilon^2 \cdot \delta^2\left(\frac{2A}{K_{eff} P_b P_t}\right) + \left[\frac{2A}{K_{eff} P_b P_t}\right]^2 \cdot \delta^2(\epsilon).\end{aligned}$$

The right term in the equation above is the purely statistical part while the term on the left is the systematic uncertainty of the normalizer. An estimate for the systematic error due to the neutron beam missing the target was also found. A sweep of the beam profile at the proton target at 720 MeV showed that the beam had a full width at

1/10 maximum ($\text{FW}_{10}M$) of 4.8 cm corresponding to a σ of 1.12 cm (the diameter of the target was 4.7 cm). A gaussian profile was assumed for the beam, and the effect of beam missing the target was taken to be one-half the fraction of beam missing the target, or

$$\begin{aligned} f_{\text{miss}}/2 &= \frac{1}{2} \frac{\int_{2.14}^{\infty} \exp[-t^2/2] dt}{\int_0^{\infty} \exp[-t^2/2] dt} \\ &\approx 0.01. \end{aligned}$$

By collecting terms, we obtain

$$\delta^2(\Delta\sigma_L) = (\Delta\sigma_L)^2 \cdot \left(\frac{\delta^2 A}{A^2} + \frac{\delta^2 K_{eff}}{K_{eff}^2} + \frac{\delta^2 P_t}{P_t^2} + \frac{\delta^2 P_n}{P_n^2} + \frac{f_{\text{miss}}^2}{4} \right) + \delta^2(\Delta\sigma_L)_{\text{statistical}}.$$

Note that the systematic normalization error scales with $\Delta\sigma_L$.

The fractional uncertainty on the target constant is estimated to be 0.044. The fractional uncertainty for the spin transfer K_{eff} is 0.028 (Table 6.2, last chapter), and the fractional uncertainty for the polarization of the proton target is 0.038. The polarization of the proton beam is known to about 1%, and the error due to beam missing the target was estimated at 1%. Therefore, these fractional errors were not included in the calculation of the normalization uncertainty. By adding the more important terms in quadrature, an estimate of the overall normalization uncertainty is $0.065|\Delta\sigma_L|$ at all energies. Table 7.1 itemizes the normalization error for the FTDC data. The sign of the necessary normalization correction may be different between the two years of running because different targets were used.

Table 7.1

$\Delta\sigma_L(np)$ from the FTDC scheme, with estimated normalization errors.

$\Delta\sigma_L(np)$			
Energy	scheme	$\Delta\sigma_L(\text{mbarns})$	normalization error(mbarns)
484 ^(a)	FTDC	-7.46 ± 1.41	0.48
568 ^(b)	FTDC	-4.60 ± 1.64	0.30
634 ^(b)	FTDC	-2.62 ± 1.29	0.17
720 ^(a)	FTDC	-8.05 ± 1.36	0.52
788 ^(a)	FTDC	-7.42 ± 1.38	0.48

^a First year of running (1987).

^b Second year of running (1988).

One must be cautious about combining the FTDC and coincidence estimates for $\Delta\sigma_L$ to obtain a final number. This is because the FTDC setup contained the coincidence setup as a subset, and in this case, one is not permitted to combine the data by weighted averages because the resulting error would be smaller than either of the FTDC and coincidence estimates for the error. One can adopt the view that we measured $\Delta\sigma_L$ using the FTDC method and used the coincidence scheme as a cross check of the overall measurement. For this reason, Table 7.1 will be taken to be the final values.

7.2 Comparison with Other Data

The data were also compared with other measurements of $\Delta\sigma_L(np)$ to check for consistency. It was mentioned in Chapter 1 that the first "measurement" of $\Delta\sigma_L(np)$ was extracted from deuteron data using extensive theoretical corrections,^[18] and therefore a cross check over the overlapping energy range was in order. In addition, preliminary data from SATURNE II^[22] (Saclay) and SIN^[20] are available. The revised Argonne data^[19] are plotted in Fig. 7.2 as well as the data from SATURNE,^[21] SATURNE II, SIN and this experiment (Table 7.1). Arndt's phase shift prediction is also indicated in the figure. The uncertainty on Arndt's prediction is about ± 2 mb at energies above 600 MeV and on the order of ± 1 mb at energies lower than 600 MeV. The agreement among the different data sets is reasonable (a preliminary estimate of the normalization error for the SIN data is $\pm 10\%$). Note that the Argonne data is essentially flat over a large energy range while the points from SATURNE, SATURNE II and this experiment suggest that there is a peak near 600 MeV and a valley near 900 MeV. This is illustrated in the figure by a dashed line. Note that there is about a 6-mb change in $\Delta\sigma_L(np)$ in going from the peak to the valley.

In order to quantify the agreement among the different data sets in some fashion, a test statistic was used. Using the t -test^[61] (t -distribution) with

$$t = \frac{x - \mu}{\sqrt{\sigma_x^2 + \sigma_\mu^2}},$$

Figure 7.2

Plot of the data from this experiment, compared with revised Argonne results, SATURNE, SATURNE II and SIN data. The dashed curve is intended to guide the eye. The solid curve is Arndt's phase shift prediction.^[62]

one can make a estimate of the probability that two experiments give the same result. Given the test statistic t for a particular experiment, the t -distribution gives the probability that other experiments have worse t -values.

If μ represents the Argonne data, some conclusions can be made about whether the Argonne data are consistent with other measurements using free polarized neutron beams. The t values for the $\Delta\sigma_L(np)$ points were calculated using numbers from References [18] and [21] and the data from Table 7.1, and the results are presented in

Table 7.2

Comparison of the Argonne data with results of this experiment, and with data from SATURNE, using the test statistic t .

Energy	t-test for $\Delta\sigma_L(np)$ data		
	FTDC	SATURNE	$P(t > r)$
484	2.4		0.016
568	0.6		0.5
633		0.3	0.7
634	1.4		0.2
720	2.0		0.04
788	1.7		0.09
882		1.9	0.05
977		1.6	0.1
1083		0.3	0.7

Table 7.2.

If the Argonne data are correct, and if the measurements of $\Delta\sigma_L(np)$ are unbiased, then the probability of an experiment obtaining a $\Delta\sigma_L(np)$ which differs as much as this one and SATURNE's from the Argonne data is given in the last column of Table 7.2. These probabilities were read from a graph in Ref. [61] (pg. 50). The preliminary data from SATURNE II differ from the Argonne data by $|t| \geq 3$, corresponding to a probability $P(|t| > r) \leq 0.002$ in the energy region 788 to 1100 MeV. These probabilities also make a statement about the validity of the Argonne data and indicate that the Argonne data are inconsistent with the data obtained using free neutron beams.

The dip in Fig. 7.2 located around 900 MeV (1.58 GeV/c) has a FWHM width of about 350 MeV (0.4 GeV/c). One detail included in the theory used to extract the pn data from the Argonne pd data was the smearing due to Fermi motion of the bound neutron (Ref. [18]) which is on the order of 0.1 GeV/c. However, a smearing of this size obscuring a 6 mb 0.4 GeV/c-wide dip to a background level is not likely due to only Fermi motion.

7.3 Extraction of the Isospin $I=0$ Component

The scattering of neutrons and protons occurs with equal probability in isospin states $I = 0$ and $I = 1$. Protons scatter only in the isospin state $I = 1$. The formula for extracting the isospin state $I = 0$ component from $\Delta\sigma_L(np)$ is

$$\Delta\sigma_L(I = 0) = 2\Delta\sigma_L(np) - \Delta\sigma_L(pp)_{corr}$$

where the pp data must be corrected for coulomb effects. Fig. 7.3 is a plot of $\Delta\sigma_L(pp)$ which have been measured by several experiments (references for the pp data can be found in Arndt's SAID program.^[62]) Using these pp data, the isospin $I = 0$ component was extracted and is shown in Fig. 7.4 along with Arndt's prediction. There is some evidence for structure in the isospin $I = 0$ channel. Argonne has a 5 mb enhancement at 788 MeV, while data from SATURNE, SATURNE II, in conjunction with this experiment, show a 5 mb dip near 1000 MeV.

Figure 7.3

Plot of $\Delta\sigma_L(pp)$ ($I = 1$) data.

One can make several comments about the isospin $I = 0$ data. First, the error bars increase by a factor of two (approximately) with respect to the $\Delta\sigma_L(np)$ data, thereby reducing the significance of any structures present. Second, Arndt's prediction is unsuitable at energies higher than 788 MeV as there are almost no data above 800 MeV in the np channel to constrain the partial wave analysis. Arndt's prediction for isospin $I = 0$ $\Delta\sigma_L$ is a curve with a slope of about 3 mb/100 MeV whereas the np data are relatively flat. This implies that when all the $\Delta\sigma_L(np)$ data are included in partial wave

Figure 7.4

Values of $\Delta\sigma_L$ for isospin $I = 0$, extracted from $\Delta\sigma_L(np)$ using pp data. Arndt's phase shift prediction^[62] is also shown.

analyses, the isospin $I = 0$ partial wave amplitudes should change markedly.

The isospin $I = 0$ and $I = 1$ data are plotted together in Fig. 7.5 for comparison (the dashed curve in the figure is for purpose of illustration). One can readily see the substantial difference in the size of the error bars in each curve. Clearly, more data are needed to confirm the shape of the structure in $\Delta\sigma_L$ for the isospin $I = 0$ channel. The structure in Fig. 7.5 for isospin $I = 0$ is probably not an $N^*(1440)$ threshold effect because the $N^*(1440)$ threshold is near 1140 MeV.

Figure 7.5
Experimental values of $\Delta\sigma_L$ for isospin states $I = 0$ and $I = 1$. The dashed curve is meant to guide the eye. It is identical in shape to the isospin $I = 1$ curve, but displaced along the horizontal and vertical axes.

It is interesting to note that the dashed curve for the isospin $I = 0$ data in Fig. 7.6 was drawn by translating the isospin $I = 1$ curve by 125 MeV and 13.2 mb. The fact that the isospin $I = 0$ channel is mostly elastic seems to suggest the inelastic channels that occur in the pp channel ($N\Delta$ and πd) may not be responsible for the structure in $\Delta\sigma_L(pp)$; that is, these structures may not be threshold effects and may, in fact, be indicators of dibaryon resonances. It should prove interesting to repeat this experiment for transverse beam and target spins—will $\Delta\sigma_T(I = 0)$ be similar to $\Delta\sigma_T(I = 1)$?

7.4 Impact on the Partial Waves

In order to assess the impact of the $\Delta\sigma_L(np)$ data on the np partial waves, the SAID program was used (Ref. [62]). This is an interactive computer program used to examine the NN database, and among other things, allows the input of experimental data so that changes in partial waves can be studied. The SIN, SATURNE, and SATURNE II data were input as well as the data from this experiment (the Argonne data was not used).

The SIN data were assigned a normalization of approximately 0.8 by the SAID program in order to make the data more consistent with the np data base. This is consistent with SIN's preliminary estimate of a $\pm 10\%$ normalization error. The SATURNE and SATURNE II data appeared to have the correct normalizations as did the data from this experiment.

The SAID program next calculated changes in phase shifts (δ) for all the partial waves, and as expected, the isospin $I = 1$ partial waves were unchanged. The isospin $I = 0$ partial waves, on the other hand, did change. These changes are tabulated in Table 7.3 for the data from SATURNE II and this experiment. The data from SIN and this experiment did not change any of the phase shifts by more than about 1° . The data from SATURNE II which are at high energies will probably impact future partial wave analyses for the isospin $I = 0$ partial waves to a greater degree than this experiment. It would be interesting to measure the high energy behavior of $\Delta\sigma_L(np)$ ($T_{lab} \geq 1200$ MeV) since this is the energy region which differs from phase shift predictions the most.

Table 7.3

Impact on isospin $I = 0$ partial waves, using data for $\Delta\sigma_L(np)$

partial wave	$\Delta\delta$ for this experiment					$\Delta\delta$ for SATURNE II		
	484	568	634	720	788	800	840	880
	MeV					MeV		
1P_1	0.32	-0.15	-0.27	0.05	1.38	3.30	9.02	8.12
3S_1	-0.24	0.10	0.15	-0.73	-0.75	-1.79	-1.81	-1.52
e_1	0.26	-0.14	-0.16	1.30	0.96	2.30	4.79	4.05
3D_1	0.11	-0.11	-0.10	0.60	0.18	0.43	-0.87	-1.08
3D_2	0.10	-0.01	0.08	-0.88	-0.49	-1.17	-14.01	-12.32
1F_3	-0.06	0.04	-0.03	-0.26	0.09	0.21	-3.84	-3.44
3D_3	0.01	-0.04	-0.04	0.06	-0.15	-0.37	3.93	3.46
e_3	-0.06	0.02	-0.06	-0.13	-0.30	-0.71	6.91	6.01
3G_3	-0.07	0.07	0.06	-0.26	-0.27	-0.66	5.13	4.54
3G_4	0.17	-0.07	-0.04	0.80	0.48	1.15	-0.35	-0.29
1H_5					-0.21	-0.50	5.47	4.78
3G_5					-0.04	-0.09	-3.32	-2.91
e_5					0.06	0.14	0.73	0.64
3I_5					0.00	0.00	-3.62	-3.12

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