ABSTRACT

A SEARCH FOR LOW-MASS SHORT-LIVED STATES COUPLING TO e+ e-

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Following the discovery of correlated positron-electron peaks in heavy ion experiments, the existence of a previously undetected particle coupling to e+ e- was postulated. If such a particle were to exist, it should also be created in the time reverse process of the postulated decay, resonant e+ e- scattering. Several experiments had been carried out to search for such resonances, producing contradictory results.

A new experiment with improved sensitivity was therefore suggested. Monoenergetic positrons would be accelerated onto electrons in a lithium target and the kinematics of the e+ e- scattering events measured by energy and position sensitive detectors. The increase in sensitivity over other experiments was obtained by a large increase in beam intensity (5 x 10^4 positrons per second) and beam resolution (1 keV), the use of lithium targets and the better kinematic reconstruction of events possible with the superior position resolution of the detectors.

The experiment was carried out at Brookhaven National Laboratory where a 3 MV Dynamitron electrostatic accelerator was converted to a positron accelerator. The monoenergetic positrons were provided by a ^22Na source in conjunction with a thin single crystal tungsten moderator, and accelerated by the Dynamitron to a tunable energy between 1.0 and 2.5 MeV with keV resolution. The beam was transported through a new beam line to the target chamber where it could be focused down to 1 mm spot size. The energy of the interactions of the positrons with the lithium target was measured by plastic scintillators,
and the position of the scattered particles was recorded by a set of four multi-wire proportional chambers capable of sub millimeter resolution.

In a first experiment, the positron beam was scanned between 2150 and 2350 keV in 5 keV steps onto a 1.5 mg/cm² lithium target with at least 60 000 coincidences recorded per energy point. The full kinematics of each positron-electron scattering event was reconstructed. The positron-nucleus scattering events were also detected and served as a normalization to eliminate all beam and target effects. After appropriate cuts and corrections, a normalized $e^+ - e^-$ scattering excitation function was derived, leading to an upper limit (90% CL) to the cross section for pointlike particles of $\sigma = 1 - 5 \text{ mb}$, depending on the type of coupling assumed. This translates to lower limits on the lifetime of $\tau \geq 0.5$ to $7 \times 10^{-13}$ sec, depending on the energy.

A Monte Carlo simulation of the experiment was written, and the simulated event rates, energy and position distributions closely reproduced the experimental data.

In an extended run, data were acquired on a 2.5 mg/cm² lithium target from 1350 keV to 2350 keV in 3.5 keV steps with ~ 300 000 coincidences per point. These data were analyzed by X. Wu in an analogous way to the first run. These data yield lower limits on the lifetime of $3.3 \times 10^{-13}$ sec and $8.2 \times 10^{-13}$ sec for $J=0$ and $J=1$ resonances, respectively, over the invariant mass range $1560 \text{ keV}/c^2 < M_{X^*} < 1860 \text{ keV}/c^2$. 
A SEARCH FOR LOW-MASS SHORT-LIVED STATES COUPLING TO $e^+ e^-$

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# TABLE OF CONTENTS

Acknowledgements ii  
Table of Contents iii  
List of Figures viii  
List of Tables xvi

1. INTRODUCTION 1  

A. Motivation 1  
B. Present work 2  
C. Layout of this thesis 3

2. THEORETICAL AND EXPERIMENTAL BACKGROUND 5  

A. The e⁺ - e⁻ peaks in heavy ion scattering 6  
1. Motivation 6  
2. The peaks at EPOS 6  
2.1 Angular Information 7  
2.2 Possible Lifetime 8  
3. The ORANGE and TORI results 8  
4. Properties of the e⁺ - e⁻ peaks and new particles 9  

B. Possible Explanations 9  
1. Atomic Phenomena 10  
2. Nuclear Phenomena 10  

C. The Particle Scenario 11  
1. The Anomalous Magnetic Moment of the Electron 12  
2. Atomic Transitions 14  
3. Positronium Hyperfine Splitting 15  
4. Delbrück Scattering 16  
5. Nuclear Decay Experiments 16  

D. Particle Searches 17  
1. The Axion 17  
2. Particle Decays 19  
2.1 Heavy Meson Decay 19  
2.2 Rare Kaon Decay 20
2.3 Pion and Muon Decay
2.4 Sigma Decay
3. Beam Dump Experiments
4. Positive Result
E. Low Energy $e^+ - e^-$ Scattering
1. Mott Scattering
2. Positron Annihilation
3. $e^+ - e^-$ Scattering Kinematics
4. Bhabha Scattering
5. Resonant Scattering
6. Resonant Positron Scattering from Electrons in a Target
   6.1 The Importance of the Electron Momentum
   6.2 The Compton Profile
   6.3 Derivation of the Resonance Shape
   6.4 Choice of Target
   7.1 Source Experiments
   7.2 Positron Beam Experiments
   7.3 Search for Photon Sinal States

3. EXPERIMENTAL APPARATUS

A. The Beam Line
   1. The Source Deck
      1.1 The Source and Moderator
      1.2 The Positron Filter
   2. The Dynamitron
   3. The Beam Line
   4. The Target Chamber
   5. The Beam Dump

B. The Detectors
   1. The Scintillators
      1.1 The Energy Measurement
      1.2 The Bismuth Calibration
      1.3 The Corrections
      1.4 The Time Measurement
   2. The Wire Chambers
2.1 The Anode Plane 85
2.2 The Anode Calibration and the Linearity 87
2.3 The Cathode Plane 88
2.4 The Cathode Calibration and the Linearity 89
2.5 The Chamber Gain 90
C. The Data-Acquisition System 91
   1. Signal Processing 91
   2. The Trigger Logic 93
   3. The CAMAC Electronics 94
   3.1 Signals Recorded 94
   3.2 Digitizing Hardware 95
   4. The Data Analysis Computer 96
D. Experimental Procedure 97
   1. Procedure at each Energy 98
   2. The "Halo" Run 98

4. DATA ANALYSIS 132
   A. Input Register Check 132
   B. The Event Reconstruction 133
   C. The Kinematic Reconstruction 134
      1. Singles 134
      2. Coincidences 136
   D. Cuts 138
      1. Raw Cuts 138
      2. Physical Cuts 138
   E. Corrections 140
      1. Target Position 140
      2. Chamber Misalignment 141
      3. Chamber Gain 142
      4. Chamber Efficiency 143
      5. Effect on the Bhabha to Mott Ratio 143
   F. Experiment Monitoring 144
      1. Angular Parameters 144
      2. Detector Performance 145
5. MONTE CARLO SIMULATION

A. Overview 174
B. The Beam 175
C. The Target 175
D. The Wire Chambers and Position 176
   1. Geometric Effects 176
   2. Resolution 177
   3. Efficiency 178
E. The Scintillators and Energy 178
   1. Resolution 179
   2. Compton Scattered Annihilation Radiation 179
   3. Low Energy Tail 180
   4. Lower Level Discriminators 180
F. Interactions 180
   1. Mott Scattering 181
   2. Back-Scattered Positrons 182
   3. Bhabha Scattering 183
   4. Resonant Scattering 184
   5. Multiple Scattering 185
   6. Coincidence as Single 186
G. Cuts 186
H. Comparison of Simulated and Experimental Histograms 187
   1. Singles Histograms 187
   2. Coincidences Histograms 188
   3. Detector Histograms 189
I. Studies of Individual Effects 190
   1. Peak Broadening 190
   2. Low Energy Tail 192
   3. Resonance without Bhabha Scattering 192
J. Summary 192

6. RESULTS 221

A. The Bhabha/Mott Ratio 221
   1. General Fit 221
   2. Fit to the Resonance 222
7. CONCLUSIONS AND OUTLOOK

A. General Conclusions
B. Limits from this Run
C. Limits from the Extended Run
D. The Beam Dump Experiment
E. Future Experiments

REFERENCES
## LIST OF FIGURES

<table>
<thead>
<tr>
<th>Figure</th>
<th>Title</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>2.1</td>
<td>Atomic binding energies as a function of nuclear charge</td>
<td>46</td>
</tr>
<tr>
<td>2.2</td>
<td>The EPOS positron singles line family</td>
<td>46</td>
</tr>
<tr>
<td>2.3</td>
<td>Schematic view of the EPOS spectrometer</td>
<td>47</td>
</tr>
<tr>
<td>2.4</td>
<td>The EPOS sum-energy line family</td>
<td>47</td>
</tr>
<tr>
<td>2.5</td>
<td>The sum-energy spectra with different gates</td>
<td>48</td>
</tr>
<tr>
<td>2.6</td>
<td>Hemisphere correlations for the sum-energy lines</td>
<td>48</td>
</tr>
<tr>
<td>2.7</td>
<td>Schematic view of the ORANGE spectrometer</td>
<td>49</td>
</tr>
<tr>
<td>2.8</td>
<td>Schematic view of the TORI spectrometer</td>
<td>49</td>
</tr>
<tr>
<td>2.9</td>
<td>Different projections of the real and simulated coincidence data</td>
<td>50</td>
</tr>
<tr>
<td>2.10</td>
<td>Feynman diagrams of $e^+ - e^-$ interaction by photon exchange</td>
<td>50</td>
</tr>
<tr>
<td>2.11</td>
<td>Feynman diagrams of $e^+ - e^-$ interaction by particle exchange</td>
<td>51</td>
</tr>
<tr>
<td>2.12</td>
<td>Feynman diagrams for corrections to the electron magnetic moment</td>
<td>51</td>
</tr>
<tr>
<td>2.13</td>
<td>Mass dependence of the correction to the electron magnetic moment</td>
<td>52</td>
</tr>
<tr>
<td>2.14</td>
<td>Lifetime limits imposed by the electron g-factor</td>
<td>52</td>
</tr>
<tr>
<td>2.15</td>
<td>Feynman diagrams for Delbrück scattering</td>
<td>53</td>
</tr>
<tr>
<td>2.16</td>
<td>Composite of experimental limits on low-mass particles</td>
<td>53</td>
</tr>
<tr>
<td>2.17</td>
<td>Upper limit on $B(T'(1S) \rightarrow \gamma a)B(a \rightarrow e^+e^-)$</td>
<td>54</td>
</tr>
<tr>
<td>2.18</td>
<td>Limit to the branching ratio of Kaons decaying to a pion and neutral particle</td>
<td>54</td>
</tr>
<tr>
<td>2.19</td>
<td>Upper limits to branching ratios to a neutral particle of pions and muons</td>
<td>55</td>
</tr>
<tr>
<td>2.20</td>
<td>Feynman diagrams of the possible axion production mechanisms</td>
<td>55</td>
</tr>
<tr>
<td>2.21</td>
<td>Layout of beam dump experiment apparatus</td>
<td>56</td>
</tr>
<tr>
<td>2.22</td>
<td>Upper limit to pseudoscalar particle production from beam dump experiments</td>
<td>56</td>
</tr>
<tr>
<td>2.23</td>
<td>Mott scattering total cross section vs energy</td>
<td>57</td>
</tr>
<tr>
<td>2.24</td>
<td>Mott scattering differential cross section</td>
<td>57</td>
</tr>
<tr>
<td>2.25</td>
<td>Annihilation in flight total cross section vs positron energy</td>
<td>58</td>
</tr>
<tr>
<td>2.26</td>
<td>Annihilation in flight differential cross section</td>
<td>58</td>
</tr>
<tr>
<td>2.27</td>
<td>Coordinate system notation used in positron-electron scattering</td>
<td>59</td>
</tr>
<tr>
<td>2.28</td>
<td>Kinematic relations for $e^+ - e^-$ scattering</td>
<td>59</td>
</tr>
<tr>
<td>2.29</td>
<td>Opening angle sum distribution for $e^+ - e^-$ scattering</td>
<td>60</td>
</tr>
<tr>
<td>2.30</td>
<td>Minimum sum opening angle versus beam energy</td>
<td>60</td>
</tr>
<tr>
<td>2.31</td>
<td>Opening angle distribution of Bhabha scattered positrons</td>
<td>61</td>
</tr>
<tr>
<td>2.32</td>
<td>Bhabha scattering cross section in the center of mass</td>
<td>61</td>
</tr>
<tr>
<td>2.33</td>
<td>Cross section for particle decays in the center of mass reference frame</td>
<td>62</td>
</tr>
</tbody>
</table>
2.34 Cross section limits for new pointlike particles derived from the g-2 results 62
2.35 Resonance cross section for beryllium and carbon 63
2.36 Derived Lorentzian-like resonance shape 63
2.37 Experimental Compton Profile of lithium 64
2.38 Domains of integration for resonance derivation 64
2.39 Calculated resonance curves for Li, Be, Al and Ti 65
2.40 Compton profile of lithium and derived resonance cross section 65
2.41 Pair spectrometer set up used by Mills et al. 66
2.42 Positive result reported by Maier et al. 66
2.43 Difference of the azimuthal angles for a coincidence obtained by Maier et al. 67

3.1 Schematic drawing of the experiment 100
3.2 Detail of the $^{22}\text{Na}$ source 100
3.3 Beta decay spectrum from the decay of $^{22}\text{Na}$ 101
3.4 Energy distribution of positrons emitted from a Ni (100) surface 101
3.5 Angular distribution of positrons emitted from a W (110) moderator 102
3.6 Comparison of yields for moderated and unmoderated positrons 102
3.7 Drawing of the positron filter 103
3.8 Calculated trajectories of positrons through the filter 103
3.9 Calculated trajectories through the extractor and lens assembly 104
3.10 Calculated trajectories into the top dynodes of the accelerator 104
3.11 Schematic diagram of the operation of the Dynamitron 105
3.12 Beam energy width at different accelerator energies 105
3.13 Detail drawing of the beam line 106
3.14 Particle trajectories in a standard and a double focusing bending magnet 106
3.15 Schematic diagram of the object or image aperture 107
3.16 Envelope of the positron beam through the accelerator at 2.2 MeV 107
3.17 Schematic diagram of the target chamber 108
3.18 Diagram of one of the large aperture beam dump scintillators 108
3.19 Typical calibration spectrum obtained by the HPGe detector 109
3.20 Plot of the energy calibration of the germanium detector 109
3.21 Schematic diagram of the detectors 110
3.22 Blown up view of the detection system 110
3.23 The scintillator to light pipe to photomultiplier tube assembly 111
3.24 Raw calibration spectrum of a scintillator 111
3.25 Typical pedestal 112
3.26 Position of the bismuth source during calibrations
3.27 Fit of the nine parameter function to the bismuth spectrum
3.28 Two dimensional scintillator pulse height correction map
3.29 Scintillator ADC spectrum after corrections are applied
3.30 Energy measured by the scintillator versus beam energy
3.31 Time difference between the scintillators signals of a coincidence event
3.32 Frontal view and cross section of the wire chamber
3.33 Charge on the anode of the wire chamber versus applied voltage
3.34 Resistive division readout for the anode position measurement
3.35 Drawing of the printed circuit board tracks for the anode wires and the resistors
3.36 Diagram of the resistance to ground for an event on a given wire
3.37 Distribution of signals from the front cathode and sum of the anode signals
3.38 Distribution of anode signals measured with small range ADC
3.39 Distribution of anode signals measured with large range ADC
3.40 Distribution of pulse height as measured by the two different ADCs
3.41 Uncalibrated anode position histogram for uniform X-ray illumination
3.42 Channel of the centroid of each wire peak versus the wire number
3.43 Delay line readout diagram for the cathode position measurement
3.44 Drawing of the printed circuit board tracks for the cathode wires
3.45 Diagram outlining the generation of the cathode prepulses
3.46 Distribution of first stops from one side of the delay line readout
3.47 Distribution of second stops from one side of the delay line readout
3.48 Diagram showing the time taken by each signal to travel through the delay line
3.49 Distribution of sum of first stops
3.50 Distribution of the four sum combinations of the two stops from each side
3.51 Distribution of sum of correct stops
3.52 Uncalibrated cathode position histogram for uniform X-ray illumination
3.53 Linearity test of the cathode position
3.54 Cathode position showing the individual wires with dips for the steel wires
3.55 Channel number of the center of the cathode wires versus the wire number
3.56 Dependence of the chamber gain on the gas pressure
3.57 Gain variation as a function of the anode wire
3.58 Histogram of the sum of the anode signals
3.59 Diagram of the electronics for the scintillator detector
3.60 Diagram of the electronics for the wire chamber detector
3.61 Diagram of the trigger electronics
<table>
<thead>
<tr>
<th>Section</th>
<th>Description</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>3.62</td>
<td>Prescale timing diagram</td>
<td>130</td>
</tr>
<tr>
<td>3.63</td>
<td>Data from a halo run</td>
<td>131</td>
</tr>
<tr>
<td>3.64</td>
<td>Singles energy spectrum of a halo run with beam contamination</td>
<td>131</td>
</tr>
<tr>
<td>4.1</td>
<td>Corrected ADC channel number for all events in one scintillator</td>
<td>146</td>
</tr>
<tr>
<td>4.2</td>
<td>Reconstructed anode distribution for all events in chamber zero</td>
<td>146</td>
</tr>
<tr>
<td>4.3</td>
<td>Reconstructed cathode distribution for all events in chamber zero</td>
<td>147</td>
</tr>
<tr>
<td>4.4</td>
<td>Kinematics of a Mott scattered positron</td>
<td>147</td>
</tr>
<tr>
<td>4.5</td>
<td>Energy spectrum of single events in scintillator number zero</td>
<td>148</td>
</tr>
<tr>
<td>4.6</td>
<td>Energy spectrum of single events from all scintillators</td>
<td>148</td>
</tr>
<tr>
<td>4.7</td>
<td>Diagram of a possible path for positrons back-scattered from the steel pyramid</td>
<td>149</td>
</tr>
<tr>
<td>4.8</td>
<td>Energy spectrum for direct beam in a scintillator</td>
<td>149</td>
</tr>
<tr>
<td>4.9</td>
<td>Azimuthal distribution of single events</td>
<td>150</td>
</tr>
<tr>
<td>4.10</td>
<td>Scattering angle distribution of single events</td>
<td>150</td>
</tr>
<tr>
<td>4.11</td>
<td>Angular acceptance of wire chambers</td>
<td>151</td>
</tr>
<tr>
<td>4.12</td>
<td>Kinematics for Bhabha scattered positrons</td>
<td>151</td>
</tr>
<tr>
<td>4.13</td>
<td>Energy spectrum from individual particles in a coincidence</td>
<td>152</td>
</tr>
<tr>
<td>4.14</td>
<td>Total energy in a coincidence</td>
<td>152</td>
</tr>
<tr>
<td>4.15</td>
<td>Azimuthal scattering angle of individual particles in a coincidence</td>
<td>153</td>
</tr>
<tr>
<td>4.16</td>
<td>Azimuthal angular difference between the two particles in a coincidence</td>
<td>153</td>
</tr>
<tr>
<td>4.17</td>
<td>Azimuthal angular difference for Bhabha scattering from 100 μm beryllium foil</td>
<td>154</td>
</tr>
<tr>
<td>4.18</td>
<td>Scattering angle for individual particles in a coincidence</td>
<td>154</td>
</tr>
<tr>
<td>4.19</td>
<td>Center of mass scattering angle for individual particles in a coincidence</td>
<td>155</td>
</tr>
<tr>
<td>4.20</td>
<td>Total opening angle between the two particles in a coincidence</td>
<td>155</td>
</tr>
<tr>
<td>4.21</td>
<td>Scattering angle difference vs opening angle sum for a coincidence</td>
<td>156</td>
</tr>
<tr>
<td>4.22</td>
<td>Difference of the center of mass scattering angles</td>
<td>156</td>
</tr>
<tr>
<td>4.23</td>
<td>Raw cuts applied to the data</td>
<td>157</td>
</tr>
<tr>
<td>4.24</td>
<td>Singles energy spectrum with dashed line showing applied cut</td>
<td>158</td>
</tr>
<tr>
<td>4.25</td>
<td>Azimuthal angle distribution for singles with dashed line showing applied cut</td>
<td>158</td>
</tr>
<tr>
<td>4.26</td>
<td>Scattering angle distribution for singles with dashed line showing applied cut</td>
<td>159</td>
</tr>
<tr>
<td>4.27</td>
<td>Two dimensional distribution of selected singles hits</td>
<td>159</td>
</tr>
<tr>
<td>4.28</td>
<td>Two dimensional distribution of selected coincidence hits</td>
<td>160</td>
</tr>
<tr>
<td>4.29</td>
<td>Azimuthal angle distribution for coincidences with applied cut</td>
<td>160</td>
</tr>
<tr>
<td>4.30</td>
<td>Scattering angle distribution of coincidences in center of mass with applied cut</td>
<td>161</td>
</tr>
<tr>
<td>4.31</td>
<td>Reconstructed angles with target position uncorrected</td>
<td>161</td>
</tr>
<tr>
<td>4.32</td>
<td>Azimuthal angular distributions with the origin offset horizontally in software</td>
<td>162</td>
</tr>
</tbody>
</table>
4.33 Azimuthal angular distributions with the origin offset vertically in software
4.34 Center of mass scattering angle with the origin offset horizontally in software
4.35 Pulse height distribution of the front cathode signals with fitted Landau
4.36 Front cathode pulse height for all events with dashed line showing applied cut
4.37 Chamber efficiency cuts on coincidences
4.38 Coincidences to singles ratio and kinematics for uncorrected data
4.39 Coincidences to singles ratio and kinematics for gain corrected data
4.40 Coincidences to singles ratio and kinematics for position corrected data
4.41 Mean of the azimuthal angle difference for coincidences versus beam energy
4.42 Standard deviation of the mean of the azimuthal difference versus beam energy
4.43 Mean of the opening angle sum for coincidences versus beam energy
4.44 Standard deviation of the mean of the opening angle sum versus beam energy
4.45 Mean of the scattering angle center of mass difference versus beam energy
4.46 Standard deviation of the mean of the center of mass difference vs beam energy
4.47 Anode distribution from lowest and highest energy points
4.48 Nonlinearity of the anode
4.49 Difference between the real and reconstructed position of the anode
4.50 Fraction of events failing the cuts on ADC5 as a function of energy

5.1 Electron momentum distribution in lithium
5.2 Simulated cathode distribution
5.3 Simulated excitation functions with different chamber efficiencies
5.4 Energy distribution of direct beam sent into the scintillator
5.5 Azimuthal angular distribution of Mott scattered positrons
5.6 Schematic path of back-scattered positrons
5.7 Energy dependence of the back-scattering fraction of electrons
5.8 Back-scattering energy dependence of electrons incident on iron
5.9 Fit to the experimental singles spectrum to parametrize the tail
5.10 Energy dependence of the singles energy tail parameter
5.11 Center of mass scattering angle distribution of both positrons and electrons
5.12 Difference in center of mass scattering angle for a coincidence
5.13 Simulated laboratory scattering angle distribution for positrons in a coincidence
5.14 Simulated laboratory scattering angle distribution for electrons in a coincidence
5.15 Combined laboratory scattering angle distribution for positrons and electrons
5.16 Sum opening angle for a coincidence in the laboratory frame of reference
5.17 Azimuthal angle distribution for positrons in a coincidence
5.18 Azimuthal angle difference for a coincidence
5.19 Center of mass scattering angle difference after various effects are included
5.20 Azimuthal angle difference after various effects are included
5.21 Laboratory opening angle sum after various effects are included
5.22 Center of mass scattering angle for positrons and electrons
5.23 Singles energy spectrum with low energy tail and misidentified coincidences
5.24 Singles energy distribution with cuts
5.25 Singles energy distribution without cuts
5.26 Singles scattering angle distribution with cuts
5.27 Singles scattering angle distribution without cuts
5.28 Singles azimuthal angle distribution with cuts
5.29 Singles azimuthal angle distribution without cuts
5.30 Coincidence sum energy distribution
5.31 Energy distribution of individual particles in a coincidence
5.32 Opening angle sum distribution for coincidences
5.33 Center of mass scattering angle difference distribution for coincidences
5.34 Center of mass scattering angle for positrons and electrons in coincidences
5.35 Laboratory scattering angle for positrons and electrons in coincidences
5.36 Azimuthal angle difference distribution for coincidences
5.37 Azimuthal angle for both positrons and electrons in coincidences, with cuts
5.38 Azimuthal angle for both positrons and electrons in coincidences, without cuts
5.39 Distribution of all hits along the cathode readout of the wire chambers
5.40 Distribution of all hits along the anode readout of the wire chambers
5.41 Histograms of run without corrections
5.42 Histograms of run with the wire chamber position corrections
5.43 Histograms of run with the beam spot size added
5.44 Histograms of run with the detector resolution added
5.45 Histograms of run with multiple scattering added
5.46 Histograms of run with beam and target misalignment added
5.47 Singles energy spectrum without back-scattered positrons
5.48 Excitation function with Bhabha scattering turned off

6.1 Ratio of coincidences to singles
6.2 Linear fit to the data
6.3 Quadratic fit to the data
6.4 Residuals of data with linear background subtracted
<table>
<thead>
<tr>
<th>Section</th>
<th>Title</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>6.5</td>
<td>Fit of a line and a gaussian with a fixed mean and standard deviation</td>
<td>231</td>
</tr>
<tr>
<td>6.6</td>
<td>Amplitude and error in the amplitude of the fit gaussian at each energy</td>
<td>231</td>
</tr>
<tr>
<td>6.7</td>
<td>Fits to the resonance</td>
<td>232</td>
</tr>
<tr>
<td>6.8</td>
<td>Amplitude and error in the amplitude of the fit gaussians at each energy</td>
<td>233</td>
</tr>
<tr>
<td>6.9</td>
<td>Two standard deviation upper limit to the data</td>
<td>233</td>
</tr>
<tr>
<td>6.10</td>
<td>Three standard deviation upper limit to the data</td>
<td>234</td>
</tr>
<tr>
<td>6.11</td>
<td>Diagram of renormalized gaussian with lower bound</td>
<td>234</td>
</tr>
<tr>
<td>6.12</td>
<td>Normalized upper limit (90% CL) to the amplitude of a fitted resonance</td>
<td>235</td>
</tr>
<tr>
<td>6.13</td>
<td>Normalized upper limit (95% CL) to the amplitude of a fitted resonance</td>
<td>235</td>
</tr>
<tr>
<td>6.14</td>
<td>Upper limit to the cross section for scalar and pseudo-scalar coupling</td>
<td>236</td>
</tr>
<tr>
<td>6.15</td>
<td>Upper limit to the cross section for vector coupling</td>
<td>236</td>
</tr>
<tr>
<td>6.16</td>
<td>Upper limit to the cross section for axial vector coupling</td>
<td>237</td>
</tr>
<tr>
<td>6.17</td>
<td>Upper limit to the cross section in barns-electronVolts/sterradian</td>
<td>237</td>
</tr>
<tr>
<td>6.18</td>
<td>Energy dependence of the peak resonance cross section</td>
<td>238</td>
</tr>
<tr>
<td>6.19</td>
<td>Lifetime upper limit to the data for the different resonance couplings</td>
<td>239</td>
</tr>
<tr>
<td>6.20</td>
<td>Simulated excitation function</td>
<td>240</td>
</tr>
<tr>
<td>6.21</td>
<td>Experimental excitation function</td>
<td>240</td>
</tr>
<tr>
<td>6.22</td>
<td>Ratio of Bhabha cross section to Mott cross section</td>
<td>241</td>
</tr>
<tr>
<td>6.23</td>
<td>Fit of the resonance shape and a linear background to the simulated data</td>
<td>241</td>
</tr>
<tr>
<td>6.24</td>
<td>Fitted amplitudes of a resonance versus energy for the simulated data</td>
<td>242</td>
</tr>
<tr>
<td>6.25</td>
<td>Upper Limit to a resonance for the simulated data</td>
<td>242</td>
</tr>
<tr>
<td>6.26</td>
<td>Simulation with .25% errors in the good coincidence rate</td>
<td>243</td>
</tr>
<tr>
<td>6.27</td>
<td>Simulation with a $10^{-14}$ second scalar or pseudo scalar resonance</td>
<td>243</td>
</tr>
<tr>
<td>6.28</td>
<td>Simulation with a $10^{-14}$ second vector resonance</td>
<td>244</td>
</tr>
<tr>
<td>6.29</td>
<td>Simulation with a $10^{-14}$ second axial vector resonance</td>
<td>244</td>
</tr>
<tr>
<td>6.30</td>
<td>Simulation with a $5 \times 10^{-14}$ second scalar or pseudo scalar resonance</td>
<td>245</td>
</tr>
<tr>
<td>6.31</td>
<td>Simulation with a $10^{-13}$ second scalar or pseudo scalar resonance</td>
<td>245</td>
</tr>
<tr>
<td>7.1</td>
<td>90% confidence limit on the lifetime for a $J = 0$ particle vs beam energy</td>
<td>252</td>
</tr>
<tr>
<td>7.2</td>
<td>90% confidence limit on the lifetime for a $J = 0$ particle vs invariant mass</td>
<td>252</td>
</tr>
<tr>
<td>7.3</td>
<td>Composite of this result with those of other experiments</td>
<td>253</td>
</tr>
<tr>
<td>7.4</td>
<td>Coincidence to singles ratio versus invariant mass for the extended run</td>
<td>253</td>
</tr>
<tr>
<td>7.5</td>
<td>Normalized residuals versus invariant mass</td>
<td>254</td>
</tr>
<tr>
<td>7.6</td>
<td>Composite of the extended thin target run and the other limitations</td>
<td>254</td>
</tr>
<tr>
<td>7.7</td>
<td>Diagram of the target-beam dump assembly</td>
<td>255</td>
</tr>
<tr>
<td>7.8</td>
<td>Expected sensitivity to $J=0$ particle versus lifetime of the particle</td>
<td>255</td>
</tr>
<tr>
<td>Section</td>
<td>Title</td>
<td>Page</td>
</tr>
<tr>
<td>---------</td>
<td>----------------------------------------------------------------------</td>
<td>------</td>
</tr>
<tr>
<td>7.9</td>
<td>Unvetoed coincidence rate versus beam energy without cuts</td>
<td>256</td>
</tr>
<tr>
<td>7.10</td>
<td>Unvetoed coincidence rate versus beam energy with cuts</td>
<td>256</td>
</tr>
<tr>
<td>7.11</td>
<td>90% confidence limit on the excluded lifetimes vs beam energy</td>
<td>257</td>
</tr>
<tr>
<td>7.12</td>
<td>Composite of the limitations on a particle</td>
<td>257</td>
</tr>
</tbody>
</table>
# LIST OF TABLES

<table>
<thead>
<tr>
<th>Table</th>
<th>Title</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>2.1</td>
<td>Upper limits to branching ratios from various particle decays</td>
<td>19</td>
</tr>
<tr>
<td>2.2</td>
<td>Summary of beam dump experiments</td>
<td>22</td>
</tr>
<tr>
<td>2.3</td>
<td>Summary of analysis of positive result experiment</td>
<td>23</td>
</tr>
<tr>
<td>2.4</td>
<td>Elementary vertices for different couplings of resonant states to electrons</td>
<td>28</td>
</tr>
<tr>
<td>2.5</td>
<td>Comparison of different target materials</td>
<td>39</td>
</tr>
<tr>
<td>2.6</td>
<td>Searches for $e^+ - e^-$ resonances with radioactive source</td>
<td>41</td>
</tr>
<tr>
<td>2.7</td>
<td>Searches for $e^+ - e^-$ resonances with positron beams</td>
<td>43</td>
</tr>
<tr>
<td>2.8</td>
<td>Searches for long-lived $e^+ - e^-$ resonances</td>
<td>44</td>
</tr>
<tr>
<td>3.1</td>
<td>Parameters of the Dynamitron positron beam</td>
<td>74</td>
</tr>
<tr>
<td>3.2</td>
<td>Sources used for germanium detector calibration</td>
<td>78</td>
</tr>
<tr>
<td>3.3</td>
<td>Bismuth 207 electron energies and intensities</td>
<td>81</td>
</tr>
<tr>
<td>4.1</td>
<td>Final corrections on the position of the target and wire chambers</td>
<td>142</td>
</tr>
<tr>
<td>5.1</td>
<td>Simulated and experimental widths of angular distributions</td>
<td>189</td>
</tr>
<tr>
<td>6.1</td>
<td>Simulated and measured cross-sections</td>
<td>228</td>
</tr>
</tbody>
</table>
1. INTRODUCTION

A. MOTIVATION

The discovery of sharp positron peaks in heavy ion collisions at energies close to the Coulomb barrier [Sch83], followed by the discovery of sharp correlated $e^+ - e^-$ peaks [Cow86], has led to speculations about an as yet undiscovered neutral particle coupling to $e^+ - e^-$. The seemingly back-to-back nature of the positrons and electrons, in conjunction with the sharp sum energy has led naturally to the postulate of a particle of mass equal to the sum energy peak decaying into a positron electron pair [Sch85]. The discovery of a new particle in this energy range would be in conflict with the standard model of particle physics and hence of enormous significance. It is therefore imperative to further investigate and determine the existence of such a new particle.

The heavy ion measurements are subject to significant lepton backgrounds and there are speculations [Gre87] that the peaks could be due to a nuclear effect from the scattered heavy ions. It would therefore be most desirable to detect the particle in an unrelated physical system.

If a particle were to exist, it should also be produced by resonant $e^+ - e^-$ scattering [Rei86], the time reverse process of the postulated decay measured in heavy ion experiments. The search for low energy resonances in $e^+ - e^-$ scattering was therefore suggested [Sch85] as an independent, unambiguous identification of a new particle. Several experiments ([Erb86], [Bar87] and [Bar89], [Mil87], [Pec87], [von87], [Wan87] and [Sak88]) performed with positron sources resulted in contradictory evidence for a new particle. The radiation source experiments utilize positrons from $b^+$ decaying isotopes. Since the positrons are the result of $b^+$ decay, they have the typical broad energy spectrum. The experiments therefore do not probe one energy at a time, but acquire data over a range of energies simultaneously. The narrow resonances therefore have to be extracted from broad backgrounds that are often not well understood.

An improved set of experiments was therefore needed, preferably with monoenergetic positron beams. The first such experiment provided a positive result, although a more complete execution of that experiment yielded a negative result ([Mai87], [Mai88], [Mai89]). The experiment utilized an existing set of detectors with limited position resolution and a relatively low intensity positron beam from the pelletron at
Stuttgart, limiting the quality of the data. Another experiment, by Tsertos et al. ([Tse88a], [Tse88b], [Tse89a], [Tse89b], [Tse90], [Tse91a], [Tse91b]), generated a high intensity of positrons with the ILL reactor at Grenoble and selected the positron energy with a spectrometer. Although the experiment achieved very high energy resolution in detecting scattered positrons and electrons by the use of silicon detectors, the position resolution was limited to the size of the detectors. The energy resolution of the spectrometer was limited, but the experiment used the dispersion of the spectrometer to scan different energies at different positions. Other competitive experiments were carried out ([Kie87], [Lor88], [Jud90], [Hal91]).

The experiment that is the subject of this thesis aims to search for resonances in e\textsuperscript{+}-e\textsuperscript{-} scattering with a high intensity monoenergetic beam and high position and energy resolution detectors.

B. PRESENT WORK

To execute an unambiguous low energy scan of e\textsuperscript{+} - e\textsuperscript{-} scattering, a tunable monoenergetic positron beam was required. It was generated by a 22Na source in conjunction with a positron "moderator", a low energy positron filter and the 3 MV Dynamitron electrostatic accelerator at Brookhaven National Laboratory.

A collaboration of Brookhaven National Laboratory, the City College of New York and Yale University refurbished the accelerator, converted it to a positron accelerator and designed a filter to eliminated all positrons that had not been thermalized. A new beam line was built that could focus the beam down to a 1 mm spot and monitor the beam energy as data was being acquired. A set of multiwire proportional chambers were built for the position reconstruction of the events with a resolution of ~1 mm. The energy of events was measured with plastic scintillation detectors. A precise kinematic reconstruction of the events was thereby made possible.

In a first experiment, the monoenergetic positron beam was scanned between 2150 keV and 2350 keV in 5 keV steps on a 1.5 mg/cm\textsuperscript{2} lithium target, the source of electrons for the positron electron scattering. The scattering events (typically 60 000 per energy point) were detected by the scintillators and wire chambers so as to allow full kinematic reconstruction. The positron-nucleus scattering (Mott scattering) events were also detected and served as a normalization to eliminate all beam and target effects. Data were acquired on a CAMAC based data acquisition system and recorded for off-line analysis.
The data were analyzed off-line and after calibrations, appropriate cuts and corrections, a normalized $e^+ - e^-$ scattering excitation function was derived. From the normalized $e^+ - e^-$ scattering rates, upper limits to the cross section for pointlike resonances were derived. These were then translated into lower limits to the lifetime for the different types of coupling.

To prove the experiment is well understood, a Monte Carlo simulation was written. It includes the relevant cross sections, the beam parameters, the detector geometry, the detector resolution, multiple scattering, etc... The simulated distributions can then be compared to the raw data histograms, and the normalized coincidence rate compared to the experimental excitation function. The simulation results closely reproduced the experimental data.

In an extended run, data were acquired on a 2.5 mg/cm$^2$ lithium target from 1350 keV to 2350 keV in 3.5 keV steps with ~ 300 000 coincidences per point. These data were analyzed by X. Wu [Wu92], in an analogous way to the first run, and his results will be shown in the final chapter. From the higher statistics and extended range of the data, a more sensitive search for resonances was possible and more restrictive upper limits were obtained.

C. THESIS LAYOUT

In the next chapter the theoretical and experimental bounds on a new particle are explored. The heavy ion results will be reviewed and the evidence for a new particle outlined. Limits on the existence of a new particle are set from precision QED experiments and from particle searches. The low energy $e^+ - e^-$ particle searches are then reviewed, pointing to possible improvements in new experiments.

In Chapter 3 the experimental apparatus and its operations are described, from the production of the beam to the design and operation of the detectors. The experimental procedures followed for the data acquisition are also explained.

Chapter 4 covers the data analysis. It explains how events were reconstructed from chamber signals and how they were processed. The cuts and corrections that were applied to the data are shown and the final excitation function: a normalized $e^+ - e^-$ count versus energy is then obtained.
The contents and results of the Monte Carlo computer simulation of the experiment are detailed in Chapter 5. The simulation results are then compared to the experimental data.

The results from the analysis revealed no resonance, and the physical limits on particle resonances are derived in Chapter 6. The limits are also compared to the simulation results.

The need for an expanded run of the experiment becomes clear when the limits from this experiment are combined with the other experimental and theoretical limits on a particle. In Chapter 7, the results of a dual approach to improving the lifetime limits are shown. The current experiment was expanded in energy range and in statistical accuracy, and an active target-beam dump experiment was carried out to exclude longer lifetimes. Finally, we discuss the future prospects of low energy particle searches and a new generation of heavy ion experiments.
2. THEORETICAL AND EXPERIMENTAL BACKGROUND

Although the presence of correlated peaks in heavy ion interactions is by now well established, no satisfactory explanation for their origin has yet been produced. The evidence pointing to the possibility of a neutral object being at the origin of the peaks will be presented, as will the problems of this interpretation. The consequences of the data for a neutral object are outlined and other explanations for the peaks will be explored.

If a neutral object is accepted as the origin of the peaks, there are far-reaching implications for the current standard model of particle physics. The current standard model does allow for a new neutral particle, the Higgs boson, but at vastly different energies. If a new particle were to exist, the standard model would have to be expanded to include this particle, or more likely, a new family of particles. For each new particle, there would be a wide range of new possible Feynman diagrams from new possible reactions. Since every Feynman diagram has an associated cross section that can be calculated, the introduction of a new particle would affect the theoretical predictions of a wide range of cross sections.

The standard model has been successfully tested by a variety of experiments, and from the agreement between these experiments and the theoretical predictions, tight limits on the presence of a new particle can be set. Strong constraints exist for both a pointlike object and a composite object. These constraints, obtained mostly from precision tests of QED, high energy beam dump experiments and heavy meson decay measurements have all but excluded the possibility of a new short-lived neutral particle. However, most of the constraints assume the creation of just one pointlike particle. If more than one particle was created, or if it was actually a composite object, most calculated limits would be invalid. The beam dump limits have the additional problem that they assume a created particle would interact weakly enough to escape the beam dump.

By performing the time reverse reaction of the neutral object decay, i.e. e^+ - e^- scattering, a direct search for new states can be executed. This experiment is a much more direct way of probing for new particles, and the interpretation is therefore less model dependent than for other particle searches. The existence of a low mass object coupling to e^+ - e^- would produce a resonance in positron-electron scattering, i.e. Bhabha scattering. The properties of such a resonance are derived in sections 5 and 6, as are the backgrounds encountered in the implementation of the experiment. Finally, the current status of low mass e^+ - e^- searches is reviewed in section 7.
A. **THE e+ - e- PEAKS IN HEAVY ION SCATTERING**

The production of correlated e+ - e- pairs at certain sharp energies, first reported by the EPOS collaboration [Cow86] at GSI in Germany, was discovered in the scattering of U and Th near the Coulomb barrier, and similar observations were then made in Th+Th and Th+Cm collision systems [Cow87, Cow88].

1. **Motivation**

The presence of the positron-electron pairs was postulated as an explanation of sharp positron peaks observed earlier by the EPOS [Sch83] and ORANGE collaborations. The positron peaks were first believed to be evidence for the spontaneous decay of the QED vacuum, or "sparking of the vacuum" (see Figure 2.1). The sparking of the vacuum was the original motivation for the experimental program, and its predictions were successful in accounting for the properties of the first peak observed in U+Cm scattering. However, later experimental data showed the energy of the positron peak to be almost independent of the total charge involved [Cow87], as opposed to the $Z^{20}$ dependence predicted by the spontaneous emission model. A current set of positron spectra is shown in Figure 2.2. Three distinct lines have now been observed: at ~ 320 keV, ~ 360 keV and ~ 430 keV [Bok91]. The width of the lines is typically ~ 80 keV, corresponding to a Doppler broadening of $v \sim 0.05 \, c$, the speed of the heavy ion center of mass. The lines are essentially constant in energy, and are observed for systems with a total charge as low as $Z = 163$, below the predicted $Z = 173$ required for onset of spontaneous positron emission.

Since the peaks in positron energy were constant and narrow in energy for different beam-target combinations, and seemingly did not originate from nuclear internal pair conversion or other known processes, the possibility of them being produced by the two body decay of a previously undetected neutral particle was suggested. If such was the case, a correlated monoenergetic electron peak should be visible in electron spectra.

2. **The Peaks at EPOS**

The EPOS experiment was modified to measure electrons and positrons in coincidence, and correlated peaks were discovered [Cow86]. The EPOS experiment consists of a pair of solenoids (see Figure 2.3), perpendicular to the heavy ion beam direction, to transport leptons away from the high background beam-target interaction point. Each solenoid contains Si(Li) detectors to measure the energy of leptons. There is a
cylindrical NaI array to detect the positron annihilation radiation, and a set of position sensitive parallel plate avalanche detectors for the coincident measurement of heavy ions.

There currently is an ensemble of three lines [Bok91] identified in various collision systems, as shown in Figure 2.4. The lines have a mean total kinetic energy of ~ 620 keV, ~750 keV and ~810 keV and share a set of unique features. They have similar cross sections \(-\frac{d\sigma}{d\Omega} = 5 \mu b / sr\), and most have an electron peak of the same energy and width as the positron peak. The line width of the sum energy peak is \(\leq 40 \text{ keV}\), much less than expected from the sum of the individual peaks. This can be explained when one considers the Doppler broadening of an e\(^+\) - e\(^-\) pair emitted back-to-back in the heavy ion center of mass reference frame. The Doppler broadening is proportional to \(\cos \theta_{e,HI}\), where \(\theta_{e,HI}\) is the angle between the emitted lepton and the emitting system. For a pair emitted back-to-back, the effect from both leptons cancel and a narrow sum energy peak is obtained. Conversely, when the difference in energy between the positron and electron energy is calculated, the broadening is enhanced and a very wide peak at zero energy difference is produced, as in Figure 2.4. A slightly higher average positron energy than electron energy could be caused by the Coulomb field of the heavy ions.

The narrow sum energy and broad energy difference exclude conventional internal pair conversion (IPC) ([Cow86], [Cow87], [Cow88]) because in order to obtain a narrow sum energy line from IPC, the emitter has to move at \(v < 0.01c\) in the laboratory frame of reference. However, the observed pairs are detected in coincidence with the heavy ions traveling at \(\sim 0.05c\) and no peaks were detected in delayed coincidences with the heavy ions.

2.1 Angular Information

The most direct way to determine whether the lepton pairs are caused by a two body decay is to measure the momentum correlation between them. This seemingly simple task is complicated by experimental considerations and the scattering of the leptons in the target (\(\sim 300 \mu g/cm^2\)). Some angular information can be retrieved by indirect measurements though.

Using the time of flight of the leptons, some rough emission angle cuts can be made. The coarse time of flight information is sufficient to distinguish [Bok91] e\(^+\) - e\(^-\) pairs with the positron emitted opposite from the those emitted perpendicular to the electrons (see Figure 2.5). Although the 810 keV line is enhanced to a 6 \(\sigma\) level with a cut on opposite events, region A in Figure 2.5, the 620 keV line needs a cut on regions B and
8

C to enhance the peak. The need for this cut cannot be explained by a classic two body decay scenario.

Another method for gaining angular information is to monitor the position where the lepton trajectory hits the detector. A lepton emitted in the forward direction will always hit the detectors from a backward direction, and vice versa, assuming the leptons originate at the target along the symmetry axis of the solenoids. By separating the lepton pairs by equal and opposite side detection, the 620 and 810 keV peaks can be associated [Bok90] with back-to-back emission (see Figure 2.6). However, the 750 keV peak seems to be associated with the emission of the electron in the forward direction irrespective of the positron angle of emission.

2.2 Possible Lifetime

Some information is available on the possible lifetime of a neutral object. From the width of the lines, using the Heisenberg uncertainty principle, the object would have a lifetime $\tau \geq 10^{-19}$ s. Since the fiducial volume of the detectors is rather small (1-2 cm), an upper limit of $10^{-9}$ s can be derived. Therefore, the lifetime constraints of a possible particle are $10^{-19} \leq \tau \leq 10^{-9}$ s.

3. The ORANGE and TORI Results

A separate experiment, the ORANGE collaboration [Koz87], also operated at GSI, has performed separate searches for positrons and $e^+ - e^-$ peaks in heavy ion collisions. The experiment (Figure 2.7) uses two dispersive toroidal magnetic spectrometers of the orange type to focus leptons onto silicon detector arrays. The experiment confirmed the existence of positron peaks and later of correlated positron-electron pairs. Although two peaks ($\sim 630 \pm 15$ keV and $\sim 805 \pm 15$ keV) agree with the EPOS results, there is no full agreement on the total number of lines and their energies. This may be due to the fact the groups have explored slightly different beam target combinations.

The TORI experiment [Krä87], a third collaboration also running at GSI, uses yet another magnetic configuration to transport leptons (Figure 2.8). Unlike the other two experiments, the group has not reported any sharp peaks in $e^+ - e^-$ coincidences, although a slight (20%) enhancement in the positron yield between 200 and 500 keV was seen under certain kinematic conditions for elastic nuclear collisions.
4. Properties of the $e^+ - e^-$ peaks and new particles

In summary, the particle scenario nicely explains the energies of the positron and electron in a peak coincidence, reproducing the sharp sum energy, broad difference in energy, and for some cases near identical energy for each lepton. There are some problems however with the kinematic enhancements of the peaks: the time of flight cut and the hemisphere correlations do not seem to apply uniformly to all three lines. Another problem is the requirement that the lepton emitter be produced nearly at rest in the heavy ion center of mass frame of reference. From the width of the energy difference peaks, the mean kinetic energy of the lepton emitter, $<T_x> \sim 9.5$ keV, is considerably lower than predicted in most models.

If the origin of the peaks is assumed to be a new particle ([Bal85], [Sch85]) or single object, one would have to be postulated for each observed peak. Postulating a composite object would allow different peaks to be resonances of the same object. The composite object would have to be made up of known particles, i.e. electrons and positrons, or else multiple new particles have to be assumed.

A recent claim of $180^\circ$ correlated equal energy photons with a sum energy of $1062\pm1$ keV from 5.9 MeV/nucleon U+Th collisions [Dan87] was thought to be related to the $e^+ - e^-$ peaks. It was suggested this could be the $\gamma\gamma$-decay channel of a neutral particle, similar to the one postulated with respect to the positron-electron peaks in the GSI experiments. These $\gamma$ peaks turned out to be caused by a detector effect, and hence no other decay modes are now claimed. The $\gamma\gamma$ decay channel of a neutral particle is not a priori excluded and the searches and possibilities for $\gamma\gamma$ decays of a particle produced in $e^+ - e^-$ scattering will be discussed later.

B. POSSIBLE EXPLANATIONS

Since the discovery of the peaks, a wide variety of explanations have been proposed. They can be separated into proposals associating the peaks with known atomic or nuclear physics phenomena that had so far been overlooked, or proposals that postulate the existence of new physical phenomena. The models based on atomic and nuclear phenomena will be discussed here, and the possibility and limits on a possible new particle will be explored in detail in the next sections.
1. Atomic Phenomena

Lichten and Robatino early on proposed an atomic mechanism to explain the positron production in heavy ion collisions ([Lic85a], [Lic85b], [Rei85]). The interacting heavy ions are considered to form a molecular system, with the electrons in molecular orbitals. By considering the effect of multiple vacancies being transferred from outer shells into inner shells, the authors were able to produce a good fit to the positron peaks in U+U scattering systems. This mechanism does not accommodate the important fact that the positrons are correlated to electrons, and it is therefore not considered an adequate model for the e⁺ - e⁻ peaks.

Another suggestion for positron-electron production was the possibility of vibrational and other excitation modes of the heavy ion quasi molecule ([Car86], [Car88], [Cho88]). By calculating the pair creation probability from the time varying electromagnetic field of the vibrating nucleus, multiple peaks were reproduced. Although the model generated peaks in a subcritical system and with a soft Z-dependence, it failed to predict the strict kinematic relations between the observed positrons and electrons. The required quasi molecule lifetimes are also very long, ~5 x 10⁻¹⁶ s, as compared to a Rutherford scattering time of ~ 10⁻²¹ s.

If the e⁺ - e⁻ peaks are considered to be decay products of a neutral object, the simplest compound objects would be e⁺ - e⁻ bound states, i.e. micro-positronium. A prediction [Won86] of a 1.58 MeV bound state was produced from calculations with a classic magnetic dipole interaction. It was later generally proven [Gra88] that no e⁺ - e⁻ bound states can exist with a total energy above 2 mₑ.

The next step in complication for a neutral object of MeV mass is the existence of bound states of many positron-electron pairs, i.e. poly-positronium [Mül86]. In order to bind such systems, a new force is needed [Ion88], with a coupling constant far too large to have been undetected. It therefore seems unlikely that bound states of e⁺ - e⁻ can account for the peaks detected in heavy ion collisions without the addition of new physical phenomena, although not all positron-electron configurations can yet be ruled out [Gri89].

2. Nuclear Phenomena

An obvious source of correlated positrons and electrons in a heavy ion nuclear system is internal pair conversion (IPC). The transition strengths required are so large that the internal pair conversion process can be excluded on grounds of intensity. The sharp
energy feature of the peaks is also not replicated ([Cow86], [Cow87], [Cow88]) in simulated IPC pairs, as shown in Figure 2.9. The narrow region of $E_{e^+} - E_{e^-}$ space where the experimental peaks are found (with $E_{e^+} \approx E_{e^-}$) is not properly reproduced in the IPC process and the broad peak in the positron-electron energy difference is completely absent in the IPC simulation.

One particular case of IPC has been discussed extensively [Kra86b] as a possible source of the observed lepton pairs: the 1.76 MeV E0 transition in $^{90}$Zr. Another isotope of Zr, $^{96}$Zr also has an E0 transition at 1.59 MeV, possibly explaining the multiple peaks. However, as illustrated in Figure 2.9, the features in the energy sum and difference are not reproduced, and the lifetimes of the $^{90}$Zr state is such that the heavy ion would probably escape the EPOS fiducial volume.

More complicated processes can be obtained by combining IPC with subsequent interactions of the electron: capture into empty inner electron shells, and electron internal conversion in cascade [Cow88]. These possibilities were excluded by comparison of Monte Carlo simulations with the experimental data. Internal pair conversion from nuclear fragments from the heavy ion collision was also investigated and excluded as a source of the peaks.

Other production mechanisms were similarly [Cow88] excluded. External pair creation, where a gamma ray emitted from the nucleus converts into a positron-electron pair away from the nucleus, e.g. in the wall of the spectrometer, is one such mechanism. Finally, models where the positron and electron are monoenergetically emitted sequentially in a cascade nuclear decay also fail to reproduce the energy features of the experimental data.

Since no conventional scenario has been produced to satisfactorily model the correlated $e^+ - e^-$ peaks in heavy ion scattering, and given the immense implications of a possible new particle, the particle models are investigated in detail in the next section.

C. THE PARTICLE SCENARIO

The existence of a low mass particle coupling to $e^+ - e^-$ would significantly alter certain Quantum Electrodynamics (QED) calculations. New Feynman diagrams with virtual processes would have to be included, causing a shift in precision QED predictions.
From the agreement between theoretical calculations and experimental measurements of certain QED constants, limits on the existence of a new particle can be obtained [Suz86].

The lowest order Feynman diagrams of $\ell^+ - \ell^-$ scattering, or Bhabha scattering, are shown in Figure 2.10, with the corresponding diagrams that would have to be added for a neutral object shown in Figure 2.11. The left hand diagrams show the exchange of a virtual photon or $X^0$ (a t-channel process), whereas the right hand diagrams show the annihilation of the lepton pair into a virtual $\gamma$ or $X^0$. The possible contribution to Bhabha scattering from the $X^0$ graph was experimentally measured in the experiment described in this thesis, but the existing constraints are significant, originating from a varied field of experiments in atomic, nuclear and particle physics. The limits from atomic and nuclear phenomena are covered in this section and the vast amount of data from particle physics experiments is the subject of the next section.

When considering resonant scattering, the annihilation graph contribution is enhanced by the $X^0$ propagator when the on-mass shell condition is met. Since Bhabha scattering is known to provide an adequate description of $\ell^+ - \ell^-$ scattering, the coupling constant $\alpha_{X^0}$ must be small compared to $\alpha$, and the exchange graph contribution must be small compared to the resonantly enhanced annihilation graph. Strict limits on this $\alpha_{X^0}$ can be now derived from QED measurement.

1. The Anomalous Magnetic Moment of the Electron

The most stringent limit on the existence of an additional particle in the MeV range is obtained from the anomalous magnetic moment of the electron [Rei86].

The magnetic moment of a particle in term of its spin $\vec{s}$ is $\vec{\mu} = g \mu_B \vec{s}$, where $\mu_B$ is the Bohr magneton ($\mu_B = \frac{e\hbar}{2mc}$) and $g$ is the Lande g-factor. For a spin $\frac{1}{2}$ particle, like the electron, the Dirac theory predicts $g = 2$ [Per82]. However, the experimentally measured value differs from 2 by ~ 0.2%, and this difference is called the anomalous magnetic moment of the electron, or $g-2$. The difference is caused by radiative corrections, the additional Feynman diagrams from the emission and recapture of virtual photons. The lowest order contributions (to the order $\alpha^2$) to the anomalous magnetic moment are shown in Figure 2.12. By including radiative corrections to the order $\alpha^4$, the theoretical $g-2$ anomaly $a^\text{th}_e$ has been calculated [Kin87] to be
with (22) being the error in the calculation and (104) the propagated uncertainty in the value of $\alpha$, the fine structure constant. The value of $a_e^{\text{th}}$ above does not include the contributions from the $\mu^\pm$ and $\tau^\pm$ lepton vacuum polarization loops, hadronic effects and weak interactions.

The experimental value [Van84] for $a_e$ is

$$a_e^{\text{exp}} = 1159 652 193 (4) \times 10^{-12}. \quad (2.2)$$

The maximal difference between the experimental value and the theoretical calculation of $a_e$ can be considered as the largest possible contribution from a neutral particle $X^0$ (Figure 2.12) and from the amplitude of this contribution a limit on the $X^0$ lifetime can be obtained. As Reinhardt [Rei86] derived, $\Delta a$ can be expressed as

$$\Delta a_i = \left( \frac{\alpha_i^X}{2\pi} \right) K_i \left( \frac{m_X}{m} \right), \quad i = S, P, V, A, \quad (2.3)$$

where $i$ refers to the type of coupling: $S$ for scalar coupling ($\Gamma_S = 1$), $P$ for pseudoscalar coupling ($\Gamma_P = iy^5$), $V$ for vector coupling ($\Gamma_V = \gamma^\mu$) and $A$ for axial vector coupling ($\Gamma_A = \gamma^\mu\gamma^5$). The functions $K_i$ are shown in Figure 2.13 as a function of the mass of the neutral particle.

For a given type of coupling and a given mass, an upper limit on the $X^0$ coupling $\alpha_i^X$ can be calculated. The resonance width in terms of the coupling is then

$$\Gamma_i = \frac{1}{2} m_X \alpha_i^X F_i(\rho) \quad (2.4)$$

with $F_i$ slow functions of the mass ratio $\rho = \left( \frac{m_X}{m} \right)^2 \quad (2.5)$

given by:

$$F_S(\rho) = \left( 1 - \frac{4}{\rho} \right)^{3/2}$$

$$F_P(\rho) = \left( 1 - \frac{4}{\rho} \right)^{1/2}$$

$$a_e^{\text{th}} = \left( \frac{g-2}{2} \right) = 1159 652 263 (22) (104) \times 10^{-12} \quad (2.1)$$
\[ F_Y(\rho) = \frac{2}{3} \left( 1 - \frac{4}{\rho} \right)^{1/2} \left( 1 + \frac{2}{\rho} \right)^{1/2} \]
\[ F_A(\rho) = \frac{2}{3} \left( 1 - \frac{4}{\rho} \right)^{3/2} \]  

The lifetime lower limit on a particle, \( \tau_i \), is then simply given by \( \tau_i = \frac{\hbar}{\Gamma_i} \) and is shown in Figure 2.14 for the different types of coupling. Combining the g-2 results and the heavy ion data, the range of possible lifetimes for a pointlike particle would be narrowed from

\[ 10^{-19} \, \text{s} \leq \tau \leq 10^{-9} \, \text{s} \] to \( \sim 10^{-13} \, \text{s} \leq \tau \leq 10^{-9} \, \text{s} \).  

2. Atomic Transitions

A limit on the coupling of a particle to \( e^+ - e^- \) can also be obtained from the effect a particle would have on atomic binding energies. Unfortunately, this method sets bounds on the combined coupling constant \( g^e g^p \) and \( g^e g^n \), where \( g^e \) is the coupling to electrons, \( g^p \) the coupling to protons and \( g^n \) the coupling to neutrons (the coupling constants \( g \) and \( \alpha \) are related by

\[ \alpha = \frac{g^2}{4\pi} \]  

Using the Lamb shift in hydrogen [Lud81], with \( \Delta E = 1057\,845 \pm 9 \, \text{kHz} \), and its discrepancy from the theoretical prediction of 30 kHz [Moh75,Don91], a limit of \( g^e g^p < 2 \times 10^{-8} \) is obtained. The limit is valid for scalar, vector and axial vector coupling only and since the three cases produce almost identical values, only the scalar limit is quoted. For pseudoscalar coupling, there are no constraints from these results. From the \( K \alpha \) transition in heavy nuclei, the limit on \( g^e g^N \) [Rei86], where \( g^N = (Z g^p + N g^n) / (Z + N) \), is \( g^e g^N < 10^{-6} \).

A limit for \( g^e g^n \) can also be derived from electron-neutron scattering data, resulting in \( g^e g^n < 3 \times 10^{-8} \). From neutron-nucleus scattering data and other work, \( g^n \) and \( g^p \) can be estimated to be \( g^n, p < 10^{-5} \).

The final limits on \( g^e \) are therefore not as stringent as the g-2 limits \((\frac{g^2}{4\pi} < 1.6 \times 10^{-8})\), but they provide an independent check on the previously obtained limit.
3. Positronium Hyperfine Splitting

If an object couples to \( e^+ - e^- \), it will affect the bound states of \( e^+ - e^- \), i.e. the energy levels of positronium. As for Bhabha scattering, the Feynman diagrams of Figure 2.11 would contribute and would shift the hyperfine splitting of the energy levels of positronium. The t-channel diagram would change all energy levels equally and would therefore not affect the energy level differences. The s-channel diagram will contribute when the spin and parity of the exchanged particle is matched with the positronium levels (a vector coupling would contribute to the triple ground state, etc.).

The matrix element calculated by Schäfer [Sch86] for this process is:

\[
M = \frac{g^2}{m_X - 4m_e} A - \frac{g^2}{m_X} B, \tag{2.9}
\]

where \( A \) and \( B \) are given by, in the non-relativistic limit:

- for a scalar \( 0^+ \): \( A = 0 \) \hspace{1cm} \( B = 1 \)
- for a pseudoscalar \( 0^- \): \( A = \frac{1 - <\sigma_1\sigma_2>}{2} \) \hspace{1cm} \( B = 0 \)
- for a vector \( 1^- \): \( A = \frac{-3+<\sigma_1\sigma_2>}{2} \) \hspace{1cm} \( B = 1 \)
- for an axial vector \( 1^+ \): \( A = \frac{1 - <\sigma_1\sigma_2>}{2} \) \hspace{1cm} \( B = <\sigma_1\sigma_2> \) \tag{2.10}

with the spin factor \( <\sigma_1\sigma_2> \) equal to \(-3\) for the singlet state and \(1\) for the triplet state, respectively. Since the energy shifts produced by a scalar resonance are identical, there is no change in the hyperfine splitting and no bound on the existence of a neutral scalar particle. Using experimental results [Rit84] and the calculated value for the hyperfine splitting, the maximum discrepancy is:

\[
|\Delta E_{\text{exp}} - \Delta E_{\text{th}}| < 10 \text{ Mhz.} \tag{2.11}
\]

As for the anomalous magnetic moment of the electron, the discrepancy can be considered the maximum possible contribution of the neutral particle diagrams. The upper limits on the coupling for the P,V,A cases are then [Sch86], for a particle mass of 1.7 MeV/c²:

\[
\alpha_{X_0}^P < 10^{-6}, \alpha_{X_0}^V < 10^{-6} \text{ and } \alpha_{X_0}^A < 7 \times 10^{-7}, \tag{2.12}
\]

with corresponding lifetime limits of:
\[ \tau_{X_0}^P \geq 1 \times 10^{-15} \text{ s}, \tau_{X_0}^V \geq 1.2 \times 10^{-15} \text{ s} \text{ and } \tau_{X_0}^A \geq 3 \times 10^{-15} \text{ s}. \]  

(2.13)

Again, these limits are not as restrictive as the g-2 limits, but they provide an independent check of the possible domain for a new particle.

4. Delbrück Scattering

Schafer et al. [Sch86] derived limits on the coupling of a light neutral boson to e⁺ - e⁻ from a study of photon scattering in an external electric field, i.e. Delbrück scattering. They compared the experimental cross sections of the scattering of photons off uranium nuclei with the predicted [Mor73] cross sections obtained from the Feynman diagrams (Figure 2.15) contributed by a neutral particle. Upper limits on the scalar (g_S ≤ 50 GeV⁻¹) and pseudoscalar (g_P ≤ 2 GeV⁻¹) couplings were derived. However, the calculations do not match the data very well and they do not take into account the interference terms from the QED photon exchange graphs.

5. Nuclear Decay Experiments

A number of nuclear physics experiments have been performed to search for neutral particles, more specifically the axion discussed in the next section. These experiments vary from reactor based experiments to standard accelerator beam on target experiments to nuclear beam dump experiments. The reactor based experiments ([Zeh82], [Ana85], [Koc85], [Cav86]) executed searches for particles with a mass < 2 m_e and therefore provide no limits for the energy range considered here.

The first experiment designed to look for axions in a nuclear system, performed by Calaprice et. al. [Cal79], studied the decay of excited states of ¹²C. A series of experiments have since been executed in ⁴He [Fre84], ¹²C [Frö91], ¹³C [Bab86], ¹⁰B, ¹⁴N, ⁸Be ([Hal86], [deB86]) and ¹⁶O systems and previous experiments in ⁶Li, ¹⁰B and ¹⁴N were reanalyzed. The beam dump type experiments, where the created neutral particle has to cross an absorber before decaying, are not sensitive to lifetimes smaller than 10⁻¹¹ seconds [Muk86]. The other experiments, most recently those of Savage et al. ([Sav86], [Sav88]), put strict limits on the branching ratio to e⁺ - e⁻. For pseudoscalar particles in the lifetime range 10⁻¹⁹ s < \( \tau_X \) < 10⁻¹¹ s and for masses ~ 1.8 MeV, \( \frac{\Gamma_X}{\Gamma_Y} \) is 4 × 10⁻⁴. If the interaction between a 0⁺ particle X and the nucleons is parametrized by the following Lagrangian:
\[ L_{m} = \overline{\psi} \gamma_{5} (g^{(0)} + \tau, g^{(1)}) \psi X \]  
(2.14)

where \( X \) and \( \psi \) are the particle and nucleon fields, and \( \tau \) is the Pauli matrix, then the derived upper limits on the coupling constants \( g^{(0)} \) and \( g^{(1)} \) for isoscalar and isovector coupling, respectively, are:

\[ g^{(0)} < 1.6 \times 10^{-2} \quad \text{and} \quad g^{(1)} < 2.0 \times 10^{-2}. \]  
(2.15)

D. PARTICLE SEARCHES

Most searches for new particles were motivated by the possible existence of axions. After a brief discussion of axion theory, the particle searches in rare decay experiments and in beam dump experiments will be reviewed.

1. The Axion

In the original formulation of Quantum Chromodynamics (QCD), the gauge theory of the strong interaction, there were solutions predicted that would strongly violate CP (Charge conjugation and Parity) and P conservation. A mechanism proposed by Peccei and Quinn was introduced to suppress these conservation law violating processes. They postulated a global U(1) symmetry with a Higgs-field doublet \( \langle \phi_{1} \rangle \) and \( \langle \phi_{2} \rangle \) coupling to the up and down quarks. The breaking of the symmetry cancels the CP conservation violating terms and produces a Goldstone boson, the axion. A detailed review of axion physics was published by Kim in [Kim87].

In terms of the ratio of the vacuum expectation values of the Higgs fields

\[ x = \frac{\langle \phi_{1} \rangle}{\langle \phi_{2} \rangle}, \]  
(2.16)

the axion mass can be estimated to be

\[ m_{a} = 75 \left( x + \frac{1}{x} \right) \text{keV}, \]  
(2.17)

where three quark generations were assumed and the up to down quark mass ratio was put at 0.56. Given a certain mass particle, in this case 1.8 MeV, two possible values of \( x \) are obtained:

\[ x \sim 0.04 \quad \text{and} \quad x \sim 24.42. \]

The axion lifetime \( \tau_{a} \) can be calculated to be
\[ \tau_a = \frac{8\pi x^2 f^2}{m_e^2 (m_a^2 - 4m_e^2)^{1/2}}, \quad (2.18) \]

where \( m_e \) is the rest mass of the electron and \( f^2 \) is given by

\[ f^2 = 2 \left( \langle \phi_1 \rangle^2 + \langle \phi_2 \rangle^2 \right) = \frac{1}{\sqrt{2}G_F} \quad (2.19) \]

with \( G_F \) the Fermi coupling constant. For a 1.8 MeV axion, the lifetimes are:

\[ 4.4 \times 10^{-12} \text{s and } 1.6 \times 10^{-6} \text{s}. \quad (2.20) \]

The longer lifetime solution is not relevant to the heavy ion results, since the experimental fiducial volume limits the lifetime to \( \tau \sim 10^{-9} \text{s} \). This leaves the \( 4.4 \times 10^{-12} \text{s} \) solution with \( x \sim 0.04 \). It should be mentioned ([Kra86a], [Pec86]) that several variant axion models have been produced to allow the possibility of shorter lifetimes, but they are not included in this discussion.

Given a value of \( x \), the coupling strength to the electron is given by

\[ \frac{g^2}{4\pi} = \frac{m_e^2}{4\pi f^2 x^2}. \quad (2.21) \]

By measuring the \( e^+ - e^- \) production rates in beam dump and in nuclear experiments or by searching for heavy quark decays into axions [Kra86b], strong limits can be set on the axion coupling strength to electrons, and hence entire domains of \( x, \) mass and lifetime can be excluded.

The search for particles coupling to \( e^+ - e^- \) is summarized in Table 2.1 and Figure 2.16. Only the experiments with a limit relevant to the heavy ion peaks were included. Most of these experiments were performed as standard axion searches and are applicable to the mass-lifetime range of interest, while some experiments were specifically refined to explore the possibility of particles near 1.8 MeV/c^2 after publication of the EPOS results. The nuclear searches were described earlier. The heavy decay and beam dump experiments are now discussed in greater detail.
Table 2.1:
Limits on branching ratios of searches for particles coupling to $e^+ - e^-$.  

<table>
<thead>
<tr>
<th>Branching Ratio Measured</th>
<th>Upper Limit</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{16}O^*$ study, $\Gamma_X/\Gamma_\pi$</td>
<td>$5 \times 10^{-2}$</td>
<td>[Hal86]</td>
</tr>
<tr>
<td>$^{14}N^<em>$, $^8$Be</em> studies, $\Gamma_X/\Gamma_\gamma$</td>
<td>$4 \times 10^{-4}$</td>
<td>[Sav88]</td>
</tr>
<tr>
<td>$J/\Psi \rightarrow \gamma + a$</td>
<td>$1.4 \times 10^{-5}$</td>
<td>[Edw82]</td>
</tr>
<tr>
<td>$B(J/\Psi \rightarrow \gamma + a)B(\Upsilon \rightarrow \gamma + a)$</td>
<td>$1.6 \times 10^{-9}$</td>
<td>[Sie82]</td>
</tr>
<tr>
<td>$B(\Upsilon(1S) \rightarrow \gamma a)B(a \rightarrow e^+e^-)$</td>
<td>$2.2 \times 10^{-3}$</td>
<td>[Bow86]</td>
</tr>
<tr>
<td>$B(K^+ \rightarrow \pi^+e^+e^-)$</td>
<td>$4 \times 10^{-7} \rightarrow 4 \times 10^{-6}$</td>
<td>[Bak87]</td>
</tr>
<tr>
<td>$\Gamma(\mu^+ \rightarrow e^+a)/\Gamma(\mu^+ \rightarrow e^+\nu\nu)$</td>
<td>$10^{-12} \rightarrow 10^{-10}$</td>
<td>[Eic86]</td>
</tr>
<tr>
<td>$B(\Sigma^+ \rightarrow pa)B(a \rightarrow e^+e^-)$</td>
<td>$7 \times 10^{-6}$</td>
<td>[Ang69]</td>
</tr>
</tbody>
</table>

2. Particle Decays

The presence of axion would produce new branching ratios in many decays, and hence can be searched for in a wide variety of decays.

2.1 Heavy Meson Decay

An axion can be searched for by the decay $V \rightarrow \gamma + a$ where $V$ is a heavy vector meson ($J/\Psi$, $\Upsilon$) and $a$ is an axion.

The search in $J/\Psi$ decay was executed by Edwards et al. [Edw82]. The group put a limit on the branching ratio $B(J/\Psi \rightarrow \gamma + a)$ by measuring the ratio

$$\frac{B(J/\Psi \rightarrow \gamma + a)}{B(J/\Psi \rightarrow \mu^+ + \mu^-)}.$$  \hspace{1cm} \text{(2.22)}

Assuming a long lifetime ($\tau > 10^{-6}$s), they derived a limit of

$$B(J/\Psi \rightarrow \gamma + a) < 1.4 \times 10^{-5} \text{ (90\% CL).}$$  \hspace{1cm} \text{(2.23)}

This translates into an upper limit for the axion parameter $x$ of $x < 0.6$. To obtain a measurement that is independent of the relative axion coupling to the up and down quarks, the product of the upper limits on the branching ratios for the radiative decays of the $J/\Psi$ and the $\Upsilon$ was measured. The theoretical estimate ([Sie82], [Kra86b]) for this process is

$$P = B(J/\Psi \rightarrow \gamma + a)B(\Upsilon \rightarrow \gamma + a) = 1.6 \times 10^{-3},$$  \hspace{1cm} \text{(2.24)}
whereas the experimental value [Sie82] is $0.6 \times 10^{-9}$, excluding a standard axion.

The experiments were not sensitive to particles with $\tau < 10^{-13}$ s potentially leaving open a certain lifetime window. It was pointed out [Kra86b] though that for short axion lifetimes, the coupling to the $b$ quark would be large, producing a branching ratio $B(Y(1S) \rightarrow \gamma + a) \approx 0.15$, making this the dominant decay mode of the $Y$. The subsequent decays, $e^+e^-$, would be two orders of magnitude greater than the QED background, excluding this possibility.

At CESR, Bowcock et al. [Bow86] searched for a short-lived particle decaying to $e^+ - e^-$ in the decay of the $Y(1S)$ using the CLEO detector. Limits on the branching ratio

$$B(Y(1S) \rightarrow \gamma a)B(a \rightarrow e^+e^-)$$

combined with previous results [Ala83] are plotted as mass vs. lifetime in Figure 2.17. For a 1.8 MeV/$c^2$ particle, the upper limit to the branching ratio is $\sim 3 \times 10^{-4}$. Mageras et al. [Mag86] using the CUSB detector, also at the Cornell Storage Ring, performed a similar experiment, and obtained a slightly better limit on the branching ratio of $\sim 2.2 \times 10^{-3}$.

2.2 Rare Kaon Decay

The decay $K^+ \rightarrow \pi^+e^+e^-$ also provides a constraint on a short-lived state coupling to $e^+ - e^-$. The measurement of the branching ratio $B(K^+ \rightarrow \pi^+e^+e^-)$ itself cannot produce a bound on the decay to 1.8 MeV/$c^2$ particles ([Kra86b], [Blo75]) because high energy cuts were made on the $e^+ - e^-$ invariant masses ($m_{e^+e^-} > m_a$). Axion searches by Asano et al. ([Asa81], [Asa82]) and Cence et al. [Cen74] in Kaon decay were similarly aimed at high energies and hence relatively insensitive near 1.8 MeV/$c^2$. Baker et al. [Bak87] executed a dedicated search for low energy objects decaying into $e^+ - e^-$ and obtained a 90% confidence level limit on the branching ratio of $4.5 \times 10^{-7}$ for a lifetime of $10^{-13}$, rising to $4 \times 10^{-6}$ at $10^{-11}$s (Figure 2.18).

2.3 Pion and Muon Decay

At PSI (formerly SIN), the SINDRUM collaboration [Eic86] searched for short-lived neutral particles in the decay of $\mu^+$ and $\pi^+$. In the Peccei and Quinn axion model, Suzuki [Suz86] calculated the transition $\pi^+ \rightarrow e^+\nu a$ to have a branching ratio of $2 \times 10^{-6}$. A search for $\pi^+ \rightarrow e^+ve^+e^-$ was started to look for particles a that would promptly decay
by a → e⁺e⁻. No such decays were found. From the contour plots (Figure 2.19a) of the upper limit as a function of mass and lifetime, we see that for a mass of ~ 1.8 MeV/c², the upper limit to the branching ratio is 10⁻¹⁰ to 3×10⁻⁹, for lifetimes from 10⁻¹⁰ s to 3×10⁻¹⁰ s, respectively. The search of μ⁺ → e⁺e⁻ provided even tighter limits, as shown in Figure 2.19b. The limits on the branching ratio Γ(μ⁺ → e⁺a)/Γ(μ⁺ → e⁺νν), at the 90% confidence level, are between 10⁻¹² and 10⁻¹⁰ for the region of interest.

2.4 Sigma Decay

Most of the experiments described above are sensitive to high energy e⁺ - e⁻, with a poor response at a few MeV/c². However, since we are interested in an e⁺ - e⁻ invariant mass near 1.8 MeV/c², a sensitivity to low energy lepton pairs is highly desirable. One such experiment, a search for flavor changing neutral weak current interactions, is the study of the decay Σ⁺ → p e⁺e⁻. The experiment dates from 1969 [Ang69], and used a hydrogen bubble chamber.

The event rate observed (3 events) is consistent with the background reaction Σ⁺ → p γ → p e⁺e⁻ and an upper limit for the production of objects decaying to e⁺ - e⁻ is therefore equal to the experimental upper limit:

\[ B(Σ⁺ → pα)B(α → e⁺e⁻) < 7 \times 10^{-6}. \]  \hspace{1cm} (2.26)

From this result Suzuki [Suz86] derived

\[ \left( \frac{g_s^2}{4\pi} \right) B(α → e⁺e⁻) < 2.1 \times 10^{-8}. \]  \hspace{1cm} (2.27)

If the branching ratio of α → e⁺e⁻ is assumed to be 1, we obtain \( \left( \frac{g_s^2}{4\pi} \right) < 2.1 \times 10^{-8} \), a strong limit on the coupling constant.

3. Beam Dump Experiments

Beam dumps provide an excellent environment to search for new particles. Using a beam of high energy electrons or protons, a particle, say the X, can be produced either by bremsstrahlung or Primakoff production from secondary photons (Figure 2.20). More complicated production mechanisms such as bremsstrahlung from secondary electrons or production using the Compton effect (Figure 2.20a, with a real photon) are also possible.
For proton beam dumps, π⁰'s from hadronic showers can decay to 2γ's, which can then in turn produce electrons and positrons or interact directly to emit an X. Axion production rates, energies and angular distributions were calculated by Tsai [Tsa86]. Schäfer [Sch88a] extended these calculations to scalar, vector and axial vector couplings. He showed the scalar and pseudoscalar production rates are almost equal, and that the axial vector and vector results, although close to each other, exceed the spin 0 predictions by up to three orders of magnitude. He also considered the case of an extended state, and concluded that the production rate for X⁰'s was only significantly affected when the radii were greater than a few hundred fermis.

### Table 2.2:
Summary of beam dump experiments

<table>
<thead>
<tr>
<th>Beam type</th>
<th>Beam Energy</th>
<th>Lifetime excluded (sec)</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>e⁻</td>
<td>275 GeV</td>
<td>4 × 10⁻¹⁶ → 4.5 × 10⁻¹²</td>
<td>[Bro91]</td>
</tr>
<tr>
<td>e⁻</td>
<td>9 GeV</td>
<td>8 × 10⁻¹⁵ → 10⁻⁹</td>
<td>[Rio87]</td>
</tr>
<tr>
<td>e⁻</td>
<td>2.5 GeV</td>
<td>2 × 10⁻¹³ → 10⁻⁷</td>
<td>[Kon86]</td>
</tr>
<tr>
<td>e⁻</td>
<td>1.5 GeV</td>
<td>6 × 10⁻¹⁴ → 9 × 10⁻¹¹</td>
<td>[Dav86]</td>
</tr>
<tr>
<td>e⁻</td>
<td>45 MeV</td>
<td>10⁻¹¹ → 10⁻⁵</td>
<td>[Bec79]</td>
</tr>
<tr>
<td>p</td>
<td>800 GeV</td>
<td>10⁻¹⁴ → 10⁻¹¹</td>
<td>[Bro86]</td>
</tr>
<tr>
<td>p</td>
<td>400 GeV</td>
<td>&lt; 3.7 × 10⁻²</td>
<td>[Ber85]</td>
</tr>
</tbody>
</table>

Several beam dump experiments were performed, and the more recent results are shown in Table 2.2. The latest experiment, by Bross et al. [Bro91], used 275 GeV electrons into a tungsten beam dump with a total length of 30 cm. As for most beam dump experiments, a new particle would be produced at high energy in the beam dump and then escape it, thanks to the large Lorentz boost to its lifetime. Behind the beam dump are detectors to eliminate charged particles, either secondary particles or straggling beam, and behind that are detectors to measure the potential decay products of neutral particles (Figure 2.21).

The domain of possible neutral particles excluded is conventionally plotted as a function of lifetime and mass, assuming a unit branching ratio into e⁺ - e⁻. Figure 2.22, from Bross et al. sows their results combined with those of Riordan et al. and Brown et al., as well as the domain excluded by the g-2 limits. Riordan et al. used a 9 GeV beam with a correspondingly shorter beam dump. The results by Buchis et al. and Bergsma et
al. originated from a standard axion search but are valid in this domain, whereas the results of Davier et al., Riordan et al., Konaka et al. and Brown et al. were specifically geared towards 1.8 MeV short-lived neutral particles searches.

4. Positive Result

It should be mentioned that one experiment [EIN88], using 4.5 GeV $^{12}$C and 4.2 GeV $^{22}$Ne beams on nuclear photoemulsion, did report a positive result for a new neutral particle. Depending on the analysis performed, the data, 13 events, provides evidence for either one or three new neutral bosons ([EIN88], [deB88]), with masses and lifetimes shown in Table 2.3.

de Boer et al. corroborate their result with 2 events from 40 year old $\pi^0$ measurements. These conclusions, obtained from a very limited data set and providing such different particle masses and lifetimes depending on the analysis, in conjunction with the wide variety of negative results described above, should be considered with due caution.

Table 2.3:
Mass and lifetime of new bosons for different analyses

<table>
<thead>
<tr>
<th>Analysis of</th>
<th>Mass (MeV/$c^2$)</th>
<th>lifetime ($x 10^{-16}$ s)</th>
</tr>
</thead>
<tbody>
<tr>
<td>El-Nadi et al.</td>
<td>1.60±.59</td>
<td>15±1</td>
</tr>
<tr>
<td>de Boer et al.</td>
<td>1.14±.1</td>
<td>13±3</td>
</tr>
<tr>
<td></td>
<td>2.1±0.4</td>
<td>3.0±0.5</td>
</tr>
<tr>
<td></td>
<td>9.2 ±1.4</td>
<td>6.0±1.0</td>
</tr>
</tbody>
</table>

E. LOW ENERGY $e^+ - e^-$ SCATTERING

An alternate and intrinsically clean method to look for a new neutral particle decaying to $e^+ - e^-$ is to perform the time reverse process, i.e. $e^+ - e^-$ scattering. This means carrying out experiments similar to the particle searches at CESR, SLC and LEP, but at much lower energies. Since low energy colliders are not currently available, the experiments must use beams on stationary targets. Two alternatives are possible: electron beams on positron targets, and positron beams on electron targets.

With the recent advances in particle trapping in Penning traps and similar devices, the idea of using positrons in a positron bottle as a target has become a reality. Such an
experiment is currently being built by Cowan and coworkers at Lawrence Livermore National Laboratory [Cow91].

All low energy $e^+ - e^-$ scattering experiments executed to date consisted of scattering positrons onto electrons. The source of positrons is typically a radioactive source and the positron energy is adjusted either by a spectrometer or by acceleration to a desired energy. The source of electrons for all experiments has been a foil of some material. Since any foil contains both nuclei and electrons, the positrons interact not only with the electrons (Bhabha scattering) but also with the nuclei (Mott scattering). In fact, Mott scattering cross sections are larger than Bhabha scattering cross sections, and hence the Mott scattering background, although not producing two outgoing particles, is a significant issue. The nature of the target material is also important because the atomic energy levels of the electrons in the target atoms have a sizable effect on the shape of an $e^+ - e^-$ resonance.

The Mott scattering cross section, the positron annihilation cross section, the Bhabha scattering cross section, the scattering kinematics and the effect the target electron momentum distribution on the shape and cross section of a resonance will be surveyed. The various low energy $e^+ - e^-$ scattering experiments already published will then be reviewed.

1. Mott Scattering

The most likely interaction for an $\text{MeV/c}^2$ positron in matter is Mott scattering. The cross section as written by Halzen and Martin [Hal84] is:

$$\frac{d\sigma}{d\Omega}_{\text{Mot}} = \frac{(Z\alpha)^2 E^2}{4 K^2 \sin^4 \frac{\theta}{2}} \left(1 - v^2 \sin^2 \frac{\theta}{2}\right)$$

or, rewritten in terms of known variables with constants included,

$$\frac{d\sigma}{d\Omega}_{\text{Mot}} = \frac{(Z\alpha \hbar c)^2}{4p^2 \beta^2 \sin^4 \frac{\theta}{2}} \left(1 - \beta^2 \sin^2 \frac{\theta}{2}\right)$$

where $Z$ is the atomic charge, $\alpha$ the fine structure constant, $\hbar$ is Planck's constant, $c$ the speed of light, $p$ the positron momentum, $\beta$ the ratio of the positron speed to $c$, and $\theta$ is the angle through which the positron was scattered.

The Mott scattering cross section for lithium between $15^\circ$ and $55^\circ$ is shown in Figure 2.23 and the differential cross section is shown in Figure 2.24.
2. Positron Annihilation

The probability of a positron annihilating with an electron, although large, is of little significance for this experiment, because the detectors used (the wire chambers) are rather insensitive to the produced gamma rays. The cross section can be written [Hei54] as

\[ \sigma = \frac{\pi r_0^2}{\gamma + 1} \left( \frac{\gamma^2 + 4\gamma + 1}{\gamma^2 - 1} \ln \left( \gamma + \sqrt{\gamma^2 - 1} \right) - \frac{\gamma + 3}{\sqrt{\gamma^2 - 1}} \right) \]  

(2.30)

where \( r_0 \) is the classical electron radius \((2.8 \times 10^{-15} \text{ m})\) and \( \gamma \) is the ratio of the positron energy to its rest mass. An expression for the differential cross section can be found in Kendall and Deutsch [Ken56], who derived

\[ \frac{d\sigma}{d\Omega} = \alpha^2 \left( \frac{\hbar c}{m c^2} \right) \frac{\gamma + 1}{\gamma \beta} \left[ \frac{1}{(1 + \gamma(1 - \beta \cos \theta))^2} + \frac{3 + \gamma}{2(\gamma + 1)\gamma(1 - \beta \cos \theta)} + \frac{(1 + \gamma(1 - \beta \cos \theta))^2}{2(\gamma + 1)^2 \gamma^2(1 - \beta \cos \theta)^2} \right] \]  

(2.31)

The annihilation cross section is shown in Figure 2.25 and the differential cross section between 15° and 55° is shown in Figure 2.26.

3. \( e^+ - e^- \) Scattering Kinematics

The process of \( e^+ - e^- \) scattering is possible by known non-resonant scattering, i.e. Bhabha scattering, and potentially by resonant scattering. The kinematics of the events would be identical, although the angular distributions and the cross sections would not. The kinematics will therefore be reviewed first, followed separately by a review of the Bhabha and resonant cross sections.

For a positron incident on an electron at rest, as shown in Figure 2.27, the invariant mass is given by:

\[ s = 2 m_e (E_+ + m_e) \]  

(2.32)

where \( m_e \) is the electron mass, \( E_+ \) the total positron energy, i.e. the accelerator energy plus the rest mass \((E_+ = \text{Beam Energy} + m_e)\). The center of mass energy is then simply given by \( \sqrt{s} \). This is the energy scale where on-mass condition must be met for resonant particle production to occur. Figure 2.28a shows the center of mass energy plotted against the beam energy.
Once a positron has scattered off the electron, the energy \( E' \) of the positron or electron is given by:

\[
E' = m_e \frac{1 + \beta_{cm}^2 \cos^2 \theta}{1 - \beta_{cm}^2 \cos^2 \theta}
\]

(2.33)

where \( \theta \) is the laboratory scattering angle of either particle and \( \beta_{cm} \) is the center of mass velocity

\[
(\beta_{cm} = \frac{\vec{p}_+}{E_+ + m_e}, \gamma_{cm} = \frac{E_+ + m_e}{\sqrt{s}}).
\]

(2.34)

Figure 2.28b shows the correlation between the scattering angles of both particles for \( e^+ - e^- \) scattering. If one particle is scattered at a small angle, the other one is scattered at a very large angle. Since the positron and electron have equal mass, the graph is symmetric, and the minimum opening angle is obtained for equal scattering angles in the laboratory frame of reference, corresponding to both particles being emitted at 90° in the center of mass reference frame.

The distribution of opening angles is shown in Figure 2.29. As the energy of the beam is increased, the boost from the center of mass to the laboratory frame of reference is increased, and the particles are more forward focused, leading to a decreased total opening angle, as shown in Figure 2.30.

4. Bhabha Scattering

The non-resonant scattering of \( e^+ - e^- \), illustrated by the Feynman diagrams of Figure 2.10, is called Bhabha scattering for Bhabha's original derivation [Bha35] of the cross section, which yielded:

\[
\left( \frac{d\sigma}{d\Omega} \right)_{\text{bhabha}} = \frac{r_e^2 m_e^2}{16 \epsilon^2} \left\{ \frac{(\epsilon^2 + p^2)}{p^4 \sin^4 \frac{\theta}{2}} - \frac{8 \epsilon^4 - m^4}{p^2 \epsilon^4 \sin^2 \frac{\theta}{2}} + \frac{12 \epsilon^4 + m^4}{\epsilon^4} - \frac{4 p^2 (\epsilon^2 + p^2) \sin^2 \frac{\theta}{2}}{\epsilon^4} + \frac{4 p^4 \sin^4 \frac{\theta}{2}}{\epsilon^4} \right\}
\]

(2.35)

where \( r_e \) is the classical electron radius, \( m \) is the electron mass, \( \epsilon \) is the positron energy, \( p \) is the positron momentum, and \( \theta \) is the angle of scattering of the positron. The differential cross section is given in the center of mass, and as for Rutherford scattering and Mott scattering diverges as \( \theta \) tends to 0. The differential cross section is shown in Figure 2.31,
and the cross section, integrated between 40° and 140°, is plotted as a function of beam energy in Figure 2.32.

5. Resonant Scattering

Resonant scattering can be calculated from the Feynman diagrams of Figure 2.11. The exchange diagram contribution to the cross section is proportional to \( \alpha_r \), which is known to be small, and will not exhibit any resonant behavior. It is therefore neglected. The annihilation diagram, on the other hand, will produce a resonance when the on-mass shell condition is met.

Before writing down the cross section, we will note the expression for the width of a resonance. If the particle is assumed to be an elementary pointlike boson described by a field \( \phi_X \) satisfying the Klein-Gordon equation, then the Lagrangian can be written [Rei87] as

\[
\mathcal{L}_{\text{int}} = g_i \int d^3x \\overline{\psi}_i \Gamma_i \psi_i \phi_X
\]  

(2.36)

where \( g_i \) is the coupling constant, \( \Gamma_i \) are the vertex factors shown in Table 2.4, and \( \overline{\psi}_i \Gamma_i \psi_i \) denotes one of the bilinear covariants of the Dirac field \( \psi_i \). The width of a resonance can then be written in terms of the ratio \( \rho \), where

\[
\rho = \frac{m_X^0}{m_e}
\]  

(2.37)

as

\[
\Gamma_{e^+e^-} = \alpha_i m_e G_i(\rho) \quad \text{with} \quad \alpha_i = \frac{g_i^2}{4\pi}
\]  

(2.38)

and the functions \( G_i(\rho) \) contain the coupling dependences, since

\[
G_S = \frac{1}{2}(p^2 - 4)^{3/2} / \rho^2
\]

\[
G_P = \frac{1}{2}(p^2 - 4)^{1/2}
\]

\[
G_V = \frac{1}{3}(p^2 - 4)^{1/2}(p^2 + 2) / \rho^2
\]

\[
G_A = \frac{1}{3}(p^2 - 4)^{3/2} / \rho^2
\]  

(2.39)

One should remember that if the particle was a composite object, then a form factor would have to be included, leading to a suppressed coupling constant. Another simplifying factor is that the branching ratio to \( e^+ - e^- \) is assumed to be 1. If other decays were to occur, the cross section would have to be multiplied by the branching ratio to \( e^+ - e^- \).
Table 2.4:
Elementary vertices for the coupling of spin 0 or spin 1 resonant states to electrons. For the spin 1 case three vertices are possible.

<table>
<thead>
<tr>
<th>i</th>
<th>$J^\pi$</th>
<th>$\Gamma_i^{(a)}$</th>
<th>$\Gamma_i^{(b)}$</th>
<th>$\Gamma_i^{(c)}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$S$ (scalar)</td>
<td>$0^+$</td>
<td>1</td>
<td>—</td>
<td>—</td>
</tr>
<tr>
<td>$P$ (pseudoscalar)</td>
<td>$0^-$</td>
<td>$i\gamma^5$</td>
<td>—</td>
<td>—</td>
</tr>
<tr>
<td>$V$ (vector)</td>
<td>$1^-$</td>
<td>$\gamma_\mu$ $\mu$</td>
<td>$(p_1 + p_2)_\mu$</td>
<td>$(p_1 - p_2)_\mu$</td>
</tr>
<tr>
<td>$A$ (axial vector)</td>
<td>$1^+$</td>
<td>$\gamma_\mu \gamma_5$</td>
<td>$i(p_1 + p_2)_\mu \gamma_5$</td>
<td>$i(p_1 - p_2)_\mu \gamma_5$</td>
</tr>
</tbody>
</table>

Returning to the cross section now, using the Feynman diagram rules explained in Sakurai [Sak67], the cross section of the annihilation diagram is:

$$
\left( \frac{d\sigma}{d\Omega} \right)_{CM} = \frac{1}{4} \frac{1}{(4\pi)^2} \frac{1}{E^2_{\text{total}}} \left| \frac{\vec{p}_{\text{final}}}{\vec{p}_{\text{initial}}} \right|^2 \frac{m_{\text{fermions}}}{\Pi_{\text{external}}} (2m_{\text{fermions}})
$$

(2.40)

In our case, we have $\vec{p}_{\text{initial}} = \vec{p}_{\text{final}}$, $E^2_{\text{total}} = s$ and there are four external fermions with $m = m_e$. The cross section can therefore be rewritten as

$$
\left( \frac{d\sigma}{d\Omega} \right)_{CM} = \frac{1}{64\pi^2 s} \left| \mathcal{M}_{fi}^{(s)} \right|^2 = \frac{m_e^4}{4\pi^2 s} \left| \mathcal{M}_{fi}^{(s)} \right|^2,
$$

(2.41)

where the problem has been reduced to solving for the matrix element.

Following the derivation of Reinhardt [Rei87], but executing the calculations in the center of mass, the matrix element $\mathcal{M}_{fi}$ for the annihilation graph is given by

$$
\mathcal{M}_{fi} = g_i^2 \left( \bar{u}(p'_2,s'_2) \Gamma_{fi}^{(s)} v(p'_1,s'_1) \Delta^{s'(s)}(P) \bar{v}(p_1,s_1) \Gamma_{fi}^{(s)} u(p_2,s_2) \right)
$$

(2.42)

where $g_i$ is the coupling constant, $\bar{u}$ and $v$ are the fermion and antifermion out of the interaction respectively, and $u$ and $\bar{v}$ are the fermion and antifermion into the interaction, again respectively. The vertex factors $\Gamma$ are given in Table 2.4 and $\Delta^{s'(s)}(P)$ is the boson propagator. The boson propagator is given by, using a complex mass $m_x \to m_x + \frac{i\Gamma}{2}$,
\[
\Delta(P) = \frac{1}{p^2 - \left(\frac{m_x - \frac{i\Gamma}{2}}{2}\right)^2}, \text{ for spin 0 coupling.} \tag{2.43}
\]

\[
\Delta_{\mu\nu}(P) = \left(-g_{\mu\nu} + \frac{P_\mu P_\nu}{m_x^2}\right) \Delta(P), \text{ for spin 1 coupling.} \tag{2.44}
\]

Therefore, we have
\[
|\Delta(P)|^2 = \frac{1}{\left|p^2 - \left(\frac{m_x - \frac{i\Gamma}{2}}{2}\right)^2\right|} = \frac{1}{(p^2 - m_x^2 + \frac{\Gamma^2}{4})^2 + m_x^2\Gamma^2}.
\tag{2.45}
\]

But \(p^2 = 2m_e(E + m_e) = s\), leading to
\[
|\Delta(P)|^2 = \frac{1}{\left(s - m_x^2 + \frac{\Gamma^2}{4}\right)^2 + m_x^2\Gamma^2}.
\tag{2.46}
\]

Since experimentally the electron and positron spins are not distinguished, the square invariant matrix element must be summed over the final and averaged over the initial states, i.e. we should substitute \(|\mathcal{M}^{(i)}_{fi}|^2\) for \(|\mathcal{M}^{(i)}_{fi}|^2\). Using Reinhardt's notation, this can rewritten as
\[
|\mathcal{M}^{(i)}_{fi}|^2 = \frac{1}{4} \sum \sum |\mathcal{M}^{(i)}_{fi}|^2 = \frac{1}{4} g_i^4 \frac{1}{\left|p^2 + \left(\frac{m_x - \frac{i\Gamma}{2}}{2}\right)^2\right|} \mathcal{S}_1(s,t) = 4\pi^2 \alpha_i^2 \frac{1}{\left|p^2 + \left(\frac{m_x - \frac{i\Gamma}{2}}{2}\right)^2\right|} \mathcal{S}_1(s,t)
\tag{2.47}
\]

where the functions \(\mathcal{S}_1(s,t)\) contain the spinor and Lorentz vector algebra in terms of the Mandelstam variables. For the various couplings they are:
\[
\mathcal{S}_s(s,t) = \frac{1}{4} (s - 4m^2)^2 / m^4
\]
\[
\mathcal{S}_p(s,t) = \frac{1}{4} s^2 / m^4
\]
\[
\mathcal{S}_v(s,t) = \frac{1}{2} s^2 + 4m^4 + (s - 4m^2)t + t^2) / m^4
\]
\[
\mathcal{S}_A(s,t) = \left(\frac{1}{2} s^2 - 4m^2 s + 12m^2 - 8m^4 s / m_x^2 + 4m^4 s^2 / m_x^4 + (s - 4m^2)t + t^2) / m^4.
\tag{2.48}
\]

Substituting for the matrix elements in the center of mass cross section, we obtain
Since \( s \) has no angular dependence \((s = 4 \, (k^2 + m^2))\) and \( t = -2k^2 \, (1 - \cos \theta) \)\) and \( t \) does, the pseudoscalar and scalar differential cross sections are constant, whereas the vector and axial vector cases have an angular dependence, as shown in Figure 2.33.

Integrating over the center of mass scattering angle for a total cross section and substituting back for \( \Gamma \) we see that, for scalars,

\[
\sigma_s(s) = \frac{4\pi \alpha_i^2}{m^2} \frac{4}{\rho^4} \frac{(s - 4m^2)^2}{(s - m_x^2 + \frac{\Gamma^2}{4})^2 + m_x^2\Gamma^2} \tag{2.50}
\]

The cross section can also be derived in the laboratory frame of reference, as Reinhardt did, and his result was

\[
\sigma_i = \frac{\pi \, \alpha_i^2}{4 \, m^2} \frac{m^2}{(E - E_R)^2 + \left(\frac{\rho \Gamma_i^2}{4}\right)^2} \mathcal{R}(E) \tag{2.51}
\]

where \( E_R \) is the laboratory resonance energy, i.e. \( E_R \) satisfies the equation

\[
s = m_x^2 = 2m(E_R + m) \]

leading to \( E_R = \frac{m_x^2}{2m} - m \). When the propagator is included, \( E_R \) becomes

\[
E_R = \frac{m_x^2 - 2m^2 - \frac{\Gamma^2}{4}}{2m} \tag{2.52}
\]

The function \( \mathcal{R}(E) \), like the \( \mathcal{R}(s,t) \) in the center of mass, contain the spinor algebra and are, when evaluated on resonance:

\[
\mathcal{R}_s(E_R) = \frac{(\rho^2 - 4)^2}{\rho^2} \]

\[
\mathcal{R}_p(E_R) = \rho^2
\]
\[ \mathcal{R}_V(E_R) = \frac{4 \left( \rho^2 + 2 \right)^2}{3 \rho^2} \]

\[ \mathcal{R}_A(E_R) = \frac{4 \left( \rho^2 - 4 \right)^2}{3 \rho^2} \]

By comparing these functions to equation 2.39 and integrating over energy, one obtains

\[ \int \sigma_i(E) \, dE = \frac{2\pi^2 \rho}{m^2(\rho^2 - 4)} \left( 2 J_i + 1 \right) \Gamma_i, \]  

using

\[ \frac{\mathcal{F}_i}{G_{i1}^2} = \frac{4 \left( 2 J_i + 1 \right) \rho^2}{\left( \rho^2 - 4 \right)}. \]

From equation 2.54, we see that the integrated cross section is proportional to the intrinsic resonance width and the spinor dependence is only the number of degenerate states factor, \( 2 J_i + 1 \). The various differential cross sections in the laboratory frame of reference are shown in Figure 2.34. Note the very large resonant cross sections. For the scalar and pseudoscalar couplings, the resonant cross section is \( \sim 2200 \) barns at \( \sqrt{s} = 1800 \text{keV} \).

As stated earlier, the cross sections are obtained by only considering the annihilation diagram of Figure 2.11. All four diagrams, the exchange and annihilation of a photon or a particle, should be considered together, leading to possible interference effects. However, Caldi et al. [Cal89] showed by explicit calculation that these effects are negligible when the electron motion broadening is included, as described in the next section.

6. Resonant Positron Scattering from Electrons in a Target

6.1 The Importance of the Electron Momentum

The experimental approach being discussed, i.e. the use of a positron beam on electrons, was needed because of the lack of low energy colliders or positron bottles. The source of electrons therefore cannot practically be a bottle or other trap either. This experiment, as all the others before it, used a thin foil as an electron target. Aside from the obvious background of positron-nucleus (Mott) scattering and the induced multiple scattering, the solid target has a profound effect on the resonant scattering experiment. The
energy and momentum distributions of the target electrons vary widely depending on the target material, and, as will be shown here, are the dominant effect in determining the amplitude and shape of the resonance cross section.

The importance of the electron motion can most easily be demonstrated by considering the cases of a positron beam with energy $E_+$ and momentum $p_+$ incident upon an electron (with an energy $E_-$ and momentum $p_-$) moving parallel and antiparallel to it, with a kinetic energy equal to the Fermi energy of the material in both cases. Dropping the terms in $p_+^2$, we calculate for the case of the parallel electron

$$s = (E_+ + E_-)^2 - (p_+ - p_-)^2 = 2m(E_+ + m) + 2p_+p_- = s_0 + 2p_+p_- \quad (2.56)$$

and the antiparallel electron

$$s = (E_+ + E_-)^2 - (p_+ + p_-)^2 = 2m(E_+ + m) - 2p_+p_- = s_0 - 2p_+p_- \quad (2.57)$$

where

$$s_0 = 2m(E_+ + m). \quad (2.58)$$

The difference in invariant mass for the same beam on target can therefore be as large as

$$\Delta\sqrt{s} = \sqrt{s_0 + 2p_+p_-} - \sqrt{s_0 - 2p_+p_-} = \frac{2p_+p_-}{\sqrt{s_0}}. \quad (2.59)$$

For $\sqrt{s}_0=1800$ keV, and a commonly used target of beryllium (with a Fermi energy of 12 eV), the spread in invariant mass is $\sim 10$ keV. This implies that the electronic properties of the target alone will broaden any potential resonance to $\sim 10$ keV, many orders of magnitude wider than the intrinsic width of at most $\sim 10$ meV derived from the $\alpha$-2 upper limits. The resonant cross section is therefore spread over a very wide range, leading to a reduction in the peak cross section of $\sim 6$ orders of magnitude. Additionally, the resonance is broadened by a factor $\rho = \left(\rho = \frac{m_X}{m_e}\right)$ when transforming from the center of mass reference frame to the laboratory frame of reference.

The one positive consequence of the electron broadening of the resonance is that the relative contributions to energy broadening due to the accelerated beam energy spread and energy loss in the target are easily kept to a minimum.

Since the electron kinematics are the main consideration in determining the resonance shape, we will derive the actual resonance shape from the experimental electron momentum profiles. Reinhardt et al. [Rei87] calculated the shape of the resonance
assuming hydrogen-like wave functions and then expanded the calculations to account for atomic screening, producing a remarkably (Figure 2.35) accurate resonance line shape in beryllium, as will be seen later. Tsertos et al. [Tse89b] derived a rather crude Lorentzian line shape (Figure 2.36) from theoretical electron momentum densities. The derivation of the resonance shape using the actual electron momentum distribution from the method of Compton profile was performed by Henderson [Hen91a] and Scherding et. al. [Sch90], and will be summarized here.

6.2 The Compton Profile

To derive a realistic resonance cross section, we must integrate over all possible electron momenta, i.e.

\[ \langle \sigma(E+) \rangle = \int n(\vec{p}) \sigma(s) \, d^3 \vec{p} \]  

(2.60)

where \( n(\vec{p}) \) is the electron momentum distribution density. As a function of momentum and solid angle, this becomes

\[ \langle \sigma(E+) \rangle = \int n(\vec{p}) \, \sigma(s) \, p^2 \, dp \, d\Omega. \]  

(2.61)

To calculate the integral, the electron momentum distribution is needed. Experimentally, it is derived from the Compton Profile [Wil77], a measurement of the electron momentum density based on the Compton scattering of X-rays from electrons in motion. The wavelength shift \( \lambda' - \lambda \) of an X-ray scattered by a moving electron is

\[ \lambda' - \lambda = \frac{2h}{mc} \sin^2 \phi + \frac{2\lambda \sin \phi}{mc} p_z, \]  

(2.62)

where \( \phi \) is the angle of scattering and \( p_z \) is the projection of the electron momentum onto the scattering vector. The first term is the usual formula for Compton scattering from a free electron, and the second term contains the dependence on the electron momentum.

The line shape of Compton scattered monoenergetic X-rays from a sample at one particular angle of incidence and angle of scattering is then proportional to the electron momentum distribution. This spectrum, or energy profile, is called the Compton profile. The Compton profile for lithium, as measured by Eisenberger et. al. [Eis72], is shown in Figure 2.37.

Using the same notation as before, we can write the Compton profile \( J(p_z) \) as:
\[ J(p_z) = Z \int d\mathbf{p} \cdot n(\mathbf{p}) \] (2.63)

where the factor \( Z \) has been introduced to normalize the charge density, i.e.

\[ \int_{-\infty}^{\infty} J(p_z) \, dp_z = Z. \] (2.64)

For an isotropic charge distribution, \( J(p_z) \) can be rewritten as, in cylindrical coordinates,

\[ J(p_z) = 2\pi Z \int_0^\infty \frac{dp_r}{p_r} \, p_r \cdot n(\mathbf{p}) \] (2.65)

with \( p_r^2 = p_x^2 + p_y^2 \). In terms of the total momentum \( p \), the Compton profile is then

\[ J(p_z) = 2\pi Z \int_{|p_x|}^{\infty} \frac{dp}{p} \cdot p \cdot n(p). \] (2.66)

The Compton profile is in a close form to that desired for the calculation of the resonance cross section (equation 2.61), which can in fact be rewritten in a form containing it, as shown below.

6.3 Derivation of the Resonance Shape

The explicit calculation of the resonance shape is made possible by the realizing that the propagator for the intrinsically narrow resonances (a few meV wide) is essentially a delta function, and can hence be treated as such. In solving for the resonance shape,

\[ \left\langle \sigma(E_+) \right\rangle = \int n(\mathcal{p}) \sigma(s) \, p^2 \, d\mathcal{p} \, d\Omega \] (2.67)

we need \( \sigma(s) \), which was derived in equation 2.50.

When replacing the propagator by a delta function, proper normalization must be ensured. Since

\[ \int_0^\infty \frac{1}{(s - m_X^2)^2 + m_X^2 \Gamma^2} = \frac{\pi}{m_X \Gamma}, \] (2.68)

we replace

\[ |\Delta|^2 \to \frac{\pi}{m_X \Gamma} \delta(s - m_X^2) \] (2.69)

and hence the pseudoscalar cross section is
\[ \sigma_p(s) = \frac{4\pi \Gamma^2 s}{m^2(\rho^2 - 4)} \frac{\pi}{m_x \Gamma} \delta(s - m_x^2). \]  

(2.70)

Substituting \( m_x^2 \) for \( s \), this simplifies to

\[ \sigma_p(s) = \frac{4\pi^2 \rho \Gamma}{m(\rho^2 - 4)} \delta(s - m_x^2). \]  

(2.71)

Similarly, the scalar resonance cross section is found to be

\[ \sigma_s(s) = \frac{4\pi \Gamma^2 \rho^4}{m^2(\rho^2 - 4)} \frac{(s - 4m^2)^2}{s} \frac{\pi}{m_x \Gamma} \delta(s - m_x^2) \]  

(2.72)

\[ \sigma_s(s) = \frac{4\pi \Gamma^2 \rho^4}{m^2(\rho^2 - 4)^3} \frac{m^4(m_x^2 - 4)}{m^2} \frac{\pi}{m_x \Gamma} \delta(s - m_x^2) \]  

(2.73)

\[ \sigma_s(s) = \frac{4\pi^2 \rho \Gamma}{m(\rho^2 - 4)} \delta(s - m_x^2). \]  

(2.74)

The cross section for the scalar and pseudoscalar resonance are therefore identical. The process can be repeated for the vector and axial vector case to obtain

\[ \sigma_i(s) = \frac{4\pi^2 \rho \Gamma}{m(\rho^2 - 4)} (2J_i + 1) \delta(s - m_x^2). \]  

(2.75)

The cross section can then generally be rewritten as:

\[ \sigma_i(s) = \sigma_i^0 \delta(s - m_x^2) \]  

(2.76)

with

\[ \sigma_i^0 = \frac{4\pi^2 \rho \Gamma}{m(\rho^2 - 4)} (2J_i + 1). \]  

(2.77)

Solving for the resonance shape using the delta function notation, we get

\[ \langle \sigma(E_+) \rangle = \int n(p) \sigma_i^0 \delta(s - m_x^2) p^2 dp d\Omega, \]  

(2.78)

which when integrated over azimuthal angle, reduces to

\[ \langle \sigma(E_+) \rangle = 2\pi \sigma_i^0 \int n(p) \delta(s - m_x^2) p^2 dp d\cos \theta. \]  

(2.79)
We now need to express the integrand in terms of the variables \( p \) and \( \cos \theta \). If the kinematic symbols of Figure 2.27 are used, we find

\[
s = E^2 - p^2 = (E_+ + E_-)^2 - (p_\parallel^2 + p_\perp^2), \tag{2.80}
\]

where \( p_\parallel \) and \( p_\perp \) are the momenta parallel and perpendicular to the incoming positron. Rewriting in term of \( \cos \theta \), this becomes

\[
s = (E_+ + E_-)^2 - (p_\parallel + p_\perp \cos \theta)^2 - (p_\perp \sin \theta)^2. \tag{2.81}
\]

Using \( E^2 - p^2 = m^2 \) for both the electron and positron, \( s \) simplifies to

\[
s = 2m^2 + 2E_+ E_- - 2p_\parallel p_\perp \cos \theta, \tag{2.82}
\]

allowing for \( \cos \theta \) to be written in terms of \( s \) and \( p \) as

\[
\cos \theta = \frac{2E_+ E_- + 2m^2 - s}{2p_\parallel p_\perp}. \tag{2.83}
\]

Substituting \( x = \cos \theta \), and using the identity

\[
\delta(s - m_x^2) = \frac{1}{|ds/dx|_x} \delta(x - x'), \tag{2.84}
\]

where \( x' \) is evaluated at \( s - m_x^2 = 0 \), the resonance shape becomes

\[
\langle \sigma(E_+) \rangle = 2\pi \sigma^0 \int_0^\infty n(\bar{p}) p_\perp^2 \, dp_- \frac{1}{2p_\parallel p_-} \int_1^\infty \, dx \, \delta\left(x - \frac{2E_+ E_- + 2m^2 - m_x^2}{2p_\parallel p_-}\right). \tag{2.85}
\]

The integral over a delta function is an easy matter, producing a zero except at the value of \( x' \), where it is 1. But one must ensure that the value of \( x' \) is actually within the domain of integration, or the integral is zero. Therefore the limitation of \( -1 < x' < 1 \) must be investigated.

**First, for \( x'>1 \):**

we have \( 2E_+ E_- > -2 p_\parallel p_- + m_x^2 - 2m^2 \). \tag{2.86}

When substituting \( \mu = 2m^2 - m_x^2 \) and squaring, there are two possible inequalities:

\[
4 E_+^2 E_-^2 > 4 p_\parallel^2 p_-^2 + \mu^2 + 4 p_\parallel p_- \mu, \tag{2.87}
\]

if the inequality is between positive quantities and

\[
4 E_+^2 E_-^2 < 4 p_\parallel^2 p_-^2 + \mu^2 + 4 p_\parallel p_- \mu \text{ for negative quantities.} \tag{2.88}
\]
The right hand side of equation 2.87 can easily be checked to be > 0 for all region of interest, and therefore, using \( E_+^2 = p_+^2 + m^2 \), the inequality (2.87) can be rewritten as

\[ 4E_+^2 m^2 > 4(p_+^2 - E_+^2) p_-^2 + 4 p_+ \mu p_- + \mu^2 \]  

(2.89)

which can be reduced to

\[ p_-^2 - \frac{\mu}{m^2} p_+ p_- + E_+^2 - \frac{\mu^2}{4m^2} > 0. \]  

(2.90)

The roots of the equivalent equality are,

\[ p_- = \frac{1}{2} \left( p_+ \frac{\mu}{m^2} \pm E_+ \sqrt{\frac{\mu^2}{m^4} - 4} \right) \]  

(2.91)

and because \( \mu < 0 \), the minimum of the parabola is at negative \( p \), producing the left hand parabola of Figure 2.38a and 38b.

Similarly, for \( x' < 1 \), the inequality now is

\[ 2E_+ E_- < 2 p_+ p_- + m_{x'}^2 - 2m^2, \]  

(2.92)

and by analogy with the previous derivation, produces the right hand parabolas also shown in Figure 2.38 and the roots are

\[ p_- = \frac{1}{2} \left( -p_+ \frac{\mu}{m^2} \pm E_+ \sqrt{\frac{\mu^2}{m^4} - 4} \right). \]  

(2.93)

From Figure 2.38, with \( p > 0 \) and \( f_1(p) > 0 \) and \( f_2(p) < 0 \), we see that the upper limit is always

\[ p_- = \frac{1}{2} \left( -p_+ \frac{\mu}{m^2} + E_+ \sqrt{\frac{\mu^2}{m^4} - 4} \right). \]  

(2.94)

The lower limit in Figure 2.38b is clearly

\[ p_- = \frac{1}{2} \left( -p_+ \frac{\mu}{m^2} - E_+ \sqrt{\frac{\mu^2}{m^4} - 4} \right) \]  

(2.95)

when it is > 0.

When it is < 0,

\[ p_- = \frac{1}{2} \left( p_+ \frac{\mu}{m^2} + E_+ \sqrt{\frac{\mu^2}{m^4} - 4} \right) > 0 \]  

(2.96)

and \( f_1(p') > 0 \) and the situation is as in Figure 2.38a. The total acceptable domain for all cases can therefore be written as
\[
\left| \frac{1}{2} \left( p_+ \frac{\mu}{m^2} + E_+ \sqrt{\frac{\mu^2}{m^2} - 4} \right) \right| < p_- < \frac{1}{2} \left( -p_+ \frac{\mu}{m^2} + E_+ \sqrt{\frac{\mu^2}{m^2} - 4} \right)
\] (2.97)

or by definition as \(|q| < p_- < \chi\). The integral over the delta function is now

\[
\int_{-\infty}^{\infty} dx \delta \left( x - \frac{2E_+E_- + 2m^2 - m_*^2}{2p_+p_-} \right) = \Theta(p_- - |q|)\Theta(x - p_-)
\] (2.98)

where

\[
\Theta(x) = \begin{cases} 
0 & \text{for } x < 0 \\
1 & \text{for } x > 0
\end{cases}
\] (2.99)

The resonance is then

\[
\langle \sigma(E_+) \rangle = \frac{2\pi \sigma_{10}}{2p_+} \int_{-\infty}^{\infty} dp_- p_- n(p_-) \Theta(p_- - |q|)\Theta(x - p_-)
\] (2.100)

which reduces to

\[
\langle \sigma(E_+) \rangle = \frac{2\pi \sigma_{10}}{2p_+} \int_{|q|}^{\infty} dp_- p_- n(p_-).
\] (2.101)

But this functional form is the definition of the Compton profile. Using the \(J(q)\) function defined in the previous section, we can write

\[
\langle \sigma(E_+) \rangle = \frac{\sigma_{10}}{2p_+Z} (J(q) - J(\chi)).
\] (2.102)

When \(\chi\) is evaluated over the range of \(m_x\) considered, it is very large, and \(J(\chi) \to 0\), hence

\[
\langle \sigma(E_+) \rangle = \frac{\sigma_{10}}{2p_+Z} J(q).
\] (2.103)

Substituting back for \(\sigma_{10}\), we obtain the shape of the resonance in a material as

\[
\langle \sigma(E_+) \rangle = \frac{2\pi^2 p_+ \Gamma(hc)^2(2J_0 + 1)}{mZ(p^2 - 4)} \frac{J(q)}{p_+}
\] (2.104)

where

\[
q = \frac{1}{2} \left( (2 - p^2)p_+ + E_+ \left( (2 - p^2)^2 - 4 \right)^{1/2} \right)
\] (2.105)

### 6.4 Choice of Target

The final form for the resonance cross section (equation 2.104) is found to be proportional to the intrinsic width of the resonance, the spin of the coupling and the Compton profile.
We see that the resonance shape in the laboratory frame of reference is just the Compton profile of a function of the beam energy and particle rest mass, as given in equation 2.105.

On resonance, \( q=0 \) and the cross section is

\[
\langle \sigma(E_R) \rangle = \frac{4\pi^2\Gamma(\hbar c)^2(2J_1+1)}{mZ(p^2 - 4)^{3/2}} J(0). 
\]

(2.106)

The largest peak cross section will therefore be produced by the target with the largest ratio of \( \frac{J(0)}{Z} \). The \( Z^{-1} \) factor points toward low atomic number targets, and maximizing \( J(0) \) means maximizing low momentum density, i.e. small electron binding energies, also implying low \( Z \) targets. Since solid hydrogen and helium targets are not available, the preferred elements for targets are lithium and beryllium.

**Table 2.5:**
Comparison of different target materials. The normalized peaks of the Compton profiles, the corresponding peak cross section, and the full width at half maximum (FWHM) of the laboratory resonance are shown for each target material.

<table>
<thead>
<tr>
<th>Target Material</th>
<th>( \frac{J(0)}{Z} ) (a. u.)</th>
<th>( \sigma ) (mb)</th>
<th>( \Delta E ) (keV)</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>Li</td>
<td>0.625</td>
<td>2.46</td>
<td>20</td>
<td>[Dov83]</td>
</tr>
<tr>
<td>Be</td>
<td>0.460</td>
<td>1.82</td>
<td>33</td>
<td>[Man79]</td>
</tr>
<tr>
<td>Al</td>
<td>0.306</td>
<td>1.21</td>
<td>34</td>
<td>[Cau81]</td>
</tr>
<tr>
<td>Teflon</td>
<td>0.310</td>
<td>1.22</td>
<td>48</td>
<td>[Wei73]</td>
</tr>
<tr>
<td>Sc</td>
<td>0.277</td>
<td>1.09</td>
<td>42</td>
<td>[Paa76]</td>
</tr>
<tr>
<td>Ti</td>
<td>0.250</td>
<td>0.99</td>
<td>49</td>
<td>[Man76]</td>
</tr>
<tr>
<td>Polyethylyene</td>
<td>0.455</td>
<td>1.79</td>
<td>35</td>
<td>[Wei75]</td>
</tr>
</tbody>
</table>

To derive accurate resonance shapes and cross sections for different targets, experimental Compton profiles from the literature were utilized. Where the data was not sufficient at higher momenta, it was complemented by Hartree-Fock calculations. Henderson studied various possible targets for their predicted resonance characteristics, and his results are summarized in Figure 2.39 and Table 2.5. From Table 2.5, we see that the peak cross section using lithium targets is 35\% larger than with beryllium, causing a commensurate experimental increase in sensitivity for similar statistics. The detailed lithium Compton profile (Figure 2.40a) produced the resonance shape of Figure 2.40b
when evaluated for \( J=0 \) with a lifetime of \( 10^{-13}\) s at a rest mass of 1832 keV, corresponding to the 810 EPOS line.

Lithium is therefore clearly the target of choice, and was used in this and other experiments, with beryllium the best alternative. As \( Z \) increases, the resonance is broadened, and the peak cross section accordingly reduced. For this and other reasons such as multiple scattering, the choice of thorium targets in some of the initial experiments discussed below was rather unfortunate.

7. Review of Low Energy \( e^+ - e^- \) Experiments

The low energy \( e^+ - e^- \) particle searches can be categorized in two groups; an initial set of experiments, performed shortly after the publication of the EPOS results, using sources and high \( Z \) targets; and a second generation of experiments, based on collimated and selected energy positron beams on low \( Z \) targets.

7.1 Source Experiments

In order to study the production of narrow \( e^+ - e^- \) peaks in heavy ion systems, Erb et al. [Erb86] performed the first time reverse experiment with a \( ^{68}\text{Ge} - ^{68}\text{Ga} \) positron source (endpoint energy = 1.89 MeV) and a thorium target of 35-50 \( \mu \)m thickness. Thorium was chosen as a target to reproduce the conditions of the heavy ion experiments. With 2 mini-orange spectrometers at 90° to each other, they measured the sum energy distribution of \( e^+ - e^- \) coincidences, and found a peak at 670 keV with a width of 20-40 keV.

A similar experiment by Sakai et al. [Sak88] measured the electron and positron rates with a double focusing spectrometer, inverting the magnetic field to switch from electron counting to positron counting. Using a \( ^{118}\text{Te} - ^{118}\text{Se} - ^{118}\text{Sn} \) source, the group reports a peak in the electron to positron ratio with thorium and uranium targets, but no peaks were found with tantalum targets. The reported peaks had an energy of 330.8 ± 1 keV and a full width at half maximum of 3.7 ± 0.5 keV.

From the derivation of the previous section, the width of a resonance in targets of high atomic number can be calculated to be ~ 50 keV, in contradiction with the quoted results, particularly those of Sakai et al.

Wang et al. [Wan87] pointed out that for the particle decay to produce equal sharp energy peaks in the laboratory frame of reference, the decaying object would have to be nearly at rest. The decay would then be nearly back-to-back, and not at 90° as measured
before. They repeated the experiments with the mini-orange spectrometers at 180°, and after random coincidence subtraction, no peaks were found.

Table 2.6:
Experimental searches for low energy $e^+ - e^-$ resonances with radioactive source.

<table>
<thead>
<tr>
<th>Reference</th>
<th>$e^+$ Source</th>
<th>Target</th>
<th>Results</th>
</tr>
</thead>
<tbody>
<tr>
<td>[Erb86]</td>
<td>$^{68}$Ga (1.9 MeV)</td>
<td>Th (~ 30 keV)</td>
<td>Structure at 670 keV in $e^+e^-$</td>
</tr>
<tr>
<td>[Pec87]</td>
<td>$^{68}$Ga (1.9 MeV)</td>
<td>Th (58 mg/cm²)</td>
<td>no structure</td>
</tr>
<tr>
<td>[Wan87]</td>
<td>$^{68}$Ga (1.9 MeV)</td>
<td>Th (50 mg/cm²)</td>
<td>no structure</td>
</tr>
<tr>
<td>[Bar87], [Bar89]</td>
<td>$^{68}$Ga (1.9 MeV)</td>
<td>Th (30 mg/cm²)</td>
<td>structure at 640 keV in $e^+e^-$</td>
</tr>
<tr>
<td>[Sak88]</td>
<td>$^{118}$Te-$^{118}$Sb-$^{118}$Sn (2.66 MeV)</td>
<td>U, Th, Ta (~ 3 mg/cm²)</td>
<td>structure at 329 keV in $e^-$ for Th and U</td>
</tr>
<tr>
<td>[von87]</td>
<td>collimated $^{82}$Rb (3.35 MeV)</td>
<td>Polyethylene (0.75 mm)</td>
<td>structure at 710 keV in $e^+e^-$</td>
</tr>
<tr>
<td>[Mil87]</td>
<td>collimated $^{82}$Rb (3.35 MeV)</td>
<td>Be (1.5 mg/cm²)</td>
<td>$\int \sigma , dE = 14.5 \text{ barn keV}$</td>
</tr>
<tr>
<td>[Hal91]</td>
<td>collimated $^{82}$Rb (3.35 MeV)</td>
<td>Li</td>
<td>$\int \sigma , dE &lt; 1 \text{ barn keV}$</td>
</tr>
</tbody>
</table>

Yet another experiment, by Bargholtz et al. ([Bar87], [Bar89]), utilized a $^{68}$Ga source and a double focusing beta spectrometer. A 3 $\sigma$ enhancement near 320 keV was measured for both electrons and positrons, with a width of less than 40 keV. Although the reported upper limit to the width is not as far off as the other positive results, this experiments is based on just 5 energy settings for the electron and positron spectrometers, and the data analysis assumes a peaks structure between 300 and 400 keV, whereas the data shows the deviation from the simulated Bhabha spectrum increasing even below 300 keV.

All the above experiments used high Z targets. von Wimmersperg et al. [von87] reported the first experiment with low atomic number targets, namely polyethylene. Their detection system however was rather crude, consisting of a two segment annulus of plastic scintillator moved up and down the beam line to kinematically select certain energy Bhabha
scattering events. The achieved energy resolution with this method was ~ 80 keV for coincidences, much larger than the actual resonance width. The group reported a positive result of $14.5 \pm 16.8$ barn-keV cross section at a center of mass kinetic energy of 710 keV.

Shortly after von Wimmersperg et al., Mills and Levy [Mil87] presented results from an experiment also with a positron source and low-Z target, beryllium, but this time with high energy resolution of the measured coincidences. Bhabha events produced in the beryllium target were focused with a pair spectrometer (Figure 2.41) onto a shielded Si(Li) detector. Although the spectrometer provided only limited energy selection, it eliminated singles and allowed the coincidences to be measured with a sharp resolution. A 90% confidence level upper limit 1.8 MeV particles was placed at 1 barn-keV. Yet another experiment by Hallin et al. [Hal91], using positrons from a $^{27}$Si source (generated by the $^{27}$Al (p, n) $^{27}$Si reaction) and a double magnetic spectrometer obtained an upper limit of 2-4 beV for a resonance at beam energies between 2140 and 2280 keV.

### 7.2 Positron Beam Experiments

In a second generation of experiments, set-ups were refined by producing low energy spread positron beams instead of using the entire beta decay spectrum of positron sources. The well defined, tunable positron energies were obtained either with sources in conjunction with spectrometers, or by moderating positrons to a single energy and then reaccelerating them using electrostatic accelerators.

A great deal of excitement was generated when, after an initial set of negative results, Maier et al. ([Mai87], [Mai88], [Bos89]), equipped with a monoenergetic positron beam from the Stuttgart pelletron and low Z targets, reported a resonance in Be at a beam energy of 2263 keV, corresponding to the 810 line in heavy ion experiments. The energy integrated cross section was measured to be 30 barn-eV, and the full width at half maximum of the peak was 27±2 keV. These results were derived by measuring $e^+ - e^-$ coincidences in a truncated cone of plastic scintillator (Figure 2.42) with photomultiplier tubes located around the periphery of both edges of the cone. This detector allowed for the first time to measure not only the energy of a coincidence but also positional information on the event. Although the position reconstruction produced erroneous side bands (Figure 2.43), geometric cuts were placed on coincidence events to reduce the background. The quoted results seem to be just larger than the allowed cross section from the g-2 production limits of pointlike particles.
Table 2.7:
Experimental searches for low energy $e^+ - e^-$ resonances with positron beams.

<table>
<thead>
<tr>
<th>Reference</th>
<th>Target (mg/cm²)</th>
<th>Mass Range (MeV)</th>
<th>Lab Stepsize (keV)</th>
<th>Set Limit</th>
</tr>
</thead>
<tbody>
<tr>
<td>[Mai87]</td>
<td>Mylar (3)</td>
<td>1.80-1.86, 1.62-1.66</td>
<td>5-10</td>
<td>$\tau &gt; 1.8 \times 10^{-14}$ sec $\tau &gt; 3 \times 10^{-14}$ sec</td>
</tr>
<tr>
<td>[Kie87]</td>
<td>Mylar</td>
<td>1.52-1.7, 1.44-1.67</td>
<td>15</td>
<td>$\tau &gt; 3 \times 10^{-14}$ sec $\tau &gt; 3 \times 10^{-14}$ sec</td>
</tr>
<tr>
<td>[Tse88b]</td>
<td>Be (4.6)</td>
<td>1.79-1.86</td>
<td>5</td>
<td>$\tau &gt; 1.7 \times 10^{-13}$ sec $\tau &gt; 3 \times 10^{-14}$ sec</td>
</tr>
<tr>
<td>[Mai88]</td>
<td>Be (4.6)</td>
<td>1.81-1.86</td>
<td>7</td>
<td>$\tau = 3.8 \times 10^{-14}$ sec $\tau &gt; 3 \times 10^{-14}$ sec</td>
</tr>
<tr>
<td>[Lor88]</td>
<td>Be (9.2)</td>
<td>1.62-1.89</td>
<td>10</td>
<td>$\tau &gt; 0.3 \times 10^{-14}$ sec $\tau &gt; 3 \times 10^{-14}$ sec</td>
</tr>
<tr>
<td>[Tse89b]</td>
<td>Be (4.6), Mylar (6.4), Th (2.0)</td>
<td>1.80-1.86, 1.62-1.89, 1.51-1.65</td>
<td>5, 15</td>
<td>$\tau &gt; 2.2 \times 10^{-13}$ sec $\tau &gt; 2.1 \times 10^{-14}$ sec $\tau &gt; 5.3 \times 10^{-15}$ sec</td>
</tr>
</tbody>
</table>

In parallel with the Stuttgart group, experiments were performed at Giessen ([Kie87], [Koz89]), using positrons generated by pair production from a 26 MeV electron beam on a tungsten target and a magnetic monochromator; and at ILL in Grenoble ([Tse88a], [Tse89b]), where a titanium source, consistently activated by neutrons from a reactor, emits positrons that are selected with a double focusing $\beta$-spectrometer. Both the Giessen and ILL groups used Si(Li) detectors to measure $e^+ - e^-$ singles and coincidences, and both groups reported negative results from Mylar and beryllium targets (see Table 2.7). The detectors allowed for an accurate reconstruction of the energy of an event, and the kinematics (at ILL) were derived by assuming coplanarity and combining detector pairs at different angles. Tsertos et al. repeated the experiment with larger statistics ([Tse88b], [Tse89], [Tse90], [Tse91a]) and still reported a negative result.

The Stuttgart results were also tested at Munich by Lorenz et al. [Lor88], again with a double focusing spectrometer, but with positrons from a $^82$Sr source. The detection system consisted of an array of plastic scintillators, to produce both adequate energy and position resolution. Although the positron rate was rather weak (~ 1000 e+ /s), the group was able to place an upper limit to a resonance over a broad energy range.

In line with the improvements in experimental techniques, a collaboration between BNL, CCNY and Yale University set out to create an experiment, with a high intensity monoenergetic positron beam, capable of both accurate position and energy resolution to
perform high statistics and low background measurements to improve the sensitivity of
searches for $e^+ - e^-$ resonances. How this was executed is the subject of the next chapter.

**Table 2.8:**
Experimental searches for long-lived $e^+ - e^-$ resonances.

<table>
<thead>
<tr>
<th>Reference</th>
<th>Target (mg/cm$^2$)</th>
<th>Mass Range (MeV)</th>
<th>Set Limit</th>
</tr>
</thead>
<tbody>
<tr>
<td>[Kli88]</td>
<td>Lucite (120)</td>
<td>1.79-1.87</td>
<td>$\sigma dE &lt; 40$ barn eV for $10^{-12} \text{ sec} &lt; \tau &lt; 10^{-10} \text{ sec}$</td>
</tr>
<tr>
<td>[Jud90]</td>
<td>Be (4.6)</td>
<td>1.78-1.855</td>
<td>$\tau &lt; 4.5 \times 10^{-13}$ sec or $\tau &gt; 7.5 \times 10^{-12}$ sec</td>
</tr>
<tr>
<td>[Tse91b]</td>
<td>Be (4.6)</td>
<td>1.785-1.855</td>
<td>$\tau &lt; 4 \times 10^{-13}$ sec or $\tau &gt; 5 \times 10^{-11}$ sec</td>
</tr>
<tr>
<td>[Wid91]</td>
<td>380 $\mu$m Plastic</td>
<td>1.78-1.92</td>
<td>$\tau &lt; 2 \times 10^{-13}$ sec or $\tau &gt; 4 \times 10^{-12}$ sec</td>
</tr>
<tr>
<td>[Hen92]</td>
<td>Be (28)</td>
<td>1.50-1.86</td>
<td>$\tau &lt; 7 \times 10^{-13}$ sec or $\tau &gt; 5 \times 10^{-10}$ sec</td>
</tr>
</tbody>
</table>

It should be mentioned that after searches for resonances in Bhabha scattering reached the statistical limitations due to practical considerations (i.e. the length of the run) the longer lifetime domain was explored by alternative methods. The direct Bhabha scattering background was eliminated, while the experiments were still sensitive to longer lived decays. A first experiment was performed by van Klinken et al.[van88] using a thick Lucite target (see Table 2.8). Tsertos et. al. ([Tse90], [Jud90], [Tse91a], [Tse91b]) achieved this effect by designing an active collimator to shadow the direct interactions in the target, while our group ([Hen91a], [Hen91b]) actually used a platinum beam dump to block events originally from the target. Although our long-lived particle search will only be detailed in the final chapter, it is already included with the other experiments in Table 2.8.

The details of the set up for the direct short-lived resonance search is described in the next chapter and the additional beam dump hardware will be outlined in a later chapter.

**7.3 Search for Photon Final States**
The experiments described above all search for the $e^+ - e^-$ decay of a resonance produced in $e^+ - e^-$ scattering. However, there is no reason to assume that the branching ratio for the resonance decay into an $e^+ - e^-$ pair would be exactly 1. If coupling to $J = 0$ states, the resonance could also decay to 2 photons and a vector or axial vector ($J = 1$) resonance could decay into 3 photons. The possible reactions can therefore be summarized as:

$$
e^+e^- \rightarrow X^0 \rightarrow e^+e^- \quad (J = 0 \text{ or } J = 1)
$$

$$
e^+e^- \rightarrow X^0 \rightarrow 2\gamma \quad (J = 0)
$$

$$
e^+e^- \rightarrow X^0 \rightarrow 3\gamma \quad (J = 1)
$$

The $2\gamma$ and $3\gamma$ decay channels must be measured in separate experiments from the $e^+ - e^-$ resonance searches because the lepton detectors are not very sensitive to gamma rays. By replacing the leptons detectors by gamma ray detectors the two other decay channels can be explored.

Connell et al. [Con88] reported on the first search for resonances in the electron-positron annihilation in flight photon spectrum. The group used positrons from a $^{27}$Si source (from $^{27}$Al (p,n) $^{27}$Si reactions) and two 3" x 3" NaI detectors to derive an upper limit of $\sim 120$ barn-electron-Volts (beV) for decays to $2\gamma$ in the energy range of the peaks in the heavy ion experiments (1640-1830 keV). A separate series of experiments by Bosch ([Bos90], [Bos91]), Kramp [Kra90] and Widmann [Wid91] reduced the upper limit to $\sim 10$ beV. The improved sensitivity was achieved with the Stuttgart pelletron as a source of monoenergetic positrons and a pair of BGO's as gamma ray detectors.

A recent search for $3\gamma$ final states by Skalsey et al. [Ska92], performed with a $^{68}$Ga source and an array of Ge detectors, obtained a 95% confidence level upper limit of $\sim 11$ beV for this decay strength.

The recent progress in searches for the photon decays of resonances in $e^+ - e^-$ scattering have improved the sensitivity to the same level as the $e^+ - e^-$ decay channel.
Figure 2.1: Atomic binding energies as a function of nuclear charge. From [Mül87].
The positron singles line family for various systems (183 ≤ Zₑ ≤ 188) as observed with the EPOS spectrometer. The spectra are ordered columnwise for increasing line energies with mean values around ∼320 keV, ∼360 keV, and ∼430 keV. Each row stems from the collision system indicated but the individual spectra are differently selected for windows in the heavy-ion scattering angles positioned symmetric around θᵣ - θₛ = 0°.

Figure 2.2 : From [Bok91].
Schematic view of the positron spectrometer system EPOS in a version optimized in addition for conversion electron spectroscopy together with a perspective drawing of the main components (only one out of two heavy ion detectors are shown for reasons of clarity).

Figure 2.3: From [Bok87].
The sum-energy line family as observed with the EPOS spectrometer for $^{235}\text{U} + ^{237}\text{Th}$ (up) and $^{238}\text{U} + ^{187}\text{Ta}$ (down) collisions at beam energies given in Table 2. The sum (left) and associated difference (right) energy spectra are integrated over the total experimental angular range of $20^\circ \leq \theta \leq 70^\circ$ and gated with different lepton time-of-flight windows as exemplified in Fig. 4, sect. II.2.

Figure 2.4: From [Sal90].
Figure 2.5: From [Bok91].

Sum-energy spectra of coincident electron positron emission in $^{238}U + ^{232}Th$ collisions associated with a beam energy range of 5.86 to 5.90 MeV/u (left) and 5.86 to 5.90 MeV/u (right). For the two-dimensional event distributions on top, plotted as functions of the lepton energies and the lepton times-of-flight, respectively, see text. The sum-energy spectra are differently gated with the two-dimensional windows as indicated (see also text).
Hemisphere correlations for the $^{239}\text{U} + ^{197}\text{Ta}$ 625-keV (a), the 805-keV (b), and the 748-keV (c,d) sum-energy lines of fig. 2: The left column (only a,b,c) corresponds to $e^-$ (open arrow) and $e^+$ emitted into opposite, the right into equal hemispheres as indicated; (d) is associated with $e^-$ emitted forward (left) and with $e^+$ emitted backward (right), irrespective of the emission direction of the $e^+$.  

Figure 2.6: From [Bok90].
Schematic view of the experimental setup in the orange-type $\beta$-spectrometer. Positrons are focused by the toroidal magnetic field onto a position sensitive proportional gas counter which surrounds a plastic scintillator. Scattered particles are detected with an annular parallel plate avalanche counter with 12 anode rings.

Figure 2.7: From [Koz87].

Schematic view of the TORI spectrometer and its detectors

Figure 2.8: From [Krä87].
Figure 2.9: Projections of U+Th coincidence data (a-d) and simulations of the two body decay of a 1.8 MeV/c² neutral object (i-l) and IPC of a 1.8 MeV E1 transition in the scattered nuclei (m-p). Parts (e)-(h) are the average of similar gates adjacent to the gates used. From [Cow87].

Figure 2.10: Feynman diagrams of e⁺ - e⁻ interaction by photon exchange. Figure 2.10a represents the photon exchange process and Figure 2.10b the annihilation process.
Figure 2.11: Feynman diagrams of e⁺ - e⁻ interaction by particle exchange. Figure 2.11a represents the X⁰ exchange process and Figure 2.11b the annihilation process.

Feynman diagrams for (a) the QED radiative corrections of the magnetic moment of a lepton (to the order of α²) and (b) the first order correction to the photon-lepton vertex from the virtual X exchange. Straight line: lepton, broken line: X, curly line: photon.

Figure 2.12: From [Aso89].
The coefficients $K_i(m_x/m)$ describing the contribution of the postulated new particle to the anomalous magnetic moment of a lepton of mass $m$.

Figure 2.13: From [Rei86].

Lifetime limits imposed by the electron $g$-factor. The vertex couplings are marked S, P, V and A. The regions below the curves (short lifetimes) are excluded.

Figure 2.14: From [Hen91a].
Figure 2.15: Feynman diagrams for Delbrück scattering. From [Aso89].

Figure 2.16: Composite contour plot of the experimental limits on low mass particles as a function of lifetime and mass. From [Gre91]
Figure 2.17: Upper limit on $B(\Upsilon (1S) \rightarrow \gamma a)B(a \rightarrow e^+ e^-)$ at 90\% CL as a function of the lifetime (solid line). The branching ratio $B(\Upsilon (1S) \rightarrow \gamma a)$ from a previous measurement is shown as the broken line. From [Bow86].

Figure 2.18: Upper limit (90\% CL) to the branching ratio of Kaons decaying to a pion and neutral particle as a function of the particle lifetime. From [Bak87].
Figure 2.19: Contour plot of the 90% CL upper limit for branching ratios as a function of the assumed masses and lifetimes. Figure 2.19a is from muon decay, Figure 2.19b from pion decay. From [Eic86].

Figure 2.20: Feynman diagrams of the possible axion production mechanisms. Figure 2.20a is axion bremsstrahlung by an electron, Figure 2.20b is Primakoff production.
EXPERIMENTAL APPARATUS

Figure 2.21: From [Bro91].

Figure 2.22: Contour plot of the 90% CL upper limit for pseudoscalar production as a function of the assumed masses and lifetimes, as measured by Bross et al. [Bro91] The results of Brown et al., Riordan et al. and the g-2 limits are also displayed.
Figure 2.23: Mott Scattering Cross section between 15° and 55°.

Figure 2.24: Scattering angle distribution for Mott scattered positrons.
Figure 2.25: Annihilation in flight cross section (in barns) versus positron energy (in keV).

Figure 2.26: Differential cross section for the annihilation in flight (in barns) versus angle (in degrees).
Figure 2.27: Coordinates for positron-electron scattering.

\[ (E_+, p_+) \rightarrow (E'_+, p'_+) \]

\[ e^+ \rightarrow \theta_+ \]

\[ (E_-, p_-) \rightarrow (E'_-, p'_-) \]

Figure 2.28: Kinematic relations for $e^+ - e^-$ scattering. Figure 2.28a shows the relationship between the center of mass energy and the laboratory positron beam energy; Figure 2.28b shows the correlation between the scattering angles for $e^+ - e^-$ and $\gamma \gamma$ final states. From [Hen91a].
Figure 2.29: Opening angle sum distribution for $e^+ - e^-$ scattering. The peak at low angle corresponds to 90° scattering in the center of mass frame of reference producing equal scattering angles for both particles in the laboratory frame of reference.

Figure 2.30: Minimum sum opening angle versus beam energy. From [Hen91a].
Figure 2.31: Opening angle distribution of Bhabha scattered positrons. The electron distribution is 180° minus the positron distribution since the positron and electron scatter back to back in the center of mass.

Figure 2.32: Bhabha scattering cross section in the center of mass integrated between 40° and 140°.
Figure 2.33: Cross section for particle decays in the center of mass reference frame. Figure 2.33a shows the vector coupling dependence, Figure 2.33b the axial vector dependence.

Figure 2.34: Cross section limits for new pointlike particles derived from the g-2 results for scalar (S), pseudoscalar (P), axial vector (A), and vector (V) coupling. Cross sections larger than the limit are not allowed. From [Hen91a].
Figure 2.35: Resonance cross section for beryllium (dashed line) and carbon (solid line) from hydrogen-like wave functions with screening taken into account. From [Rei87].

Figure 2.36: Lorentzian-like resonance shape as derived by Tsertos et al. [Tse89b].
Figure 2.37: Experimental Compton Profile of lithium as measured by Eisenberger et al. [Eis72].

\[
s_{\text{solid}} = a p^2 + b p + c = f_1(p) \\
\text{dash} = a p^2 - b p + c = f_2(p)
\]

Figure 2.38: Acceptable domain of integration. Three conditions must be met: \( p > 0, f_1(p) > 0 \) and \( f_2(p) < 0 \), leading to the range shown, for the case where the parabolas do not overlap at zero and for the case where they do.
Calculated resonance curves for Li, Be, Al, and Ti, assuming $J = 0$, $M_x = 1832$ keV and $\tau = 1 \times 10^{-13}$ sec.

Figure 2.39: From [Hen91a].

Figure 2.40: Compton profile of lithium (Figure 2.40a) used in the calculations and derived resonance cross section (Figure 2.40b), for the case $J = 0$, $M_x = 1832$ keV and $\tau = 1 \times 10^{-13}$ s. From [Hen91a].
Figure 2.41: Pair spectrometer set up used by Mills et al. [Mil87].

Figure 2.42: Positive result reported by Maier et al. [Mai88]. The data shows the experimental $e^+ - e^-$ scattering cross section versus beam energy. The fitted gaussian yields a mean of $2263.6 \pm 4$ keV and a FWHM of $27 \pm 2$ keV.
Figure 2.43: Difference of the azimuthal angles for a coincidence. Note the side lobes. Figure 2.43a is before cuts, Figure 2.43b after cuts. From [Mai88].
3. EXPERIMENTAL APPARATUS

From the discussion in the last chapter, it is clear that the $e^+ - e^-$ scattering experiments need further refinements to improve the sensitivity to resonances. Several ameliorations to the experiments seemed feasible and were pursued. They included a better beam, an improved target and high position resolution detectors.

The positron beam should be intense, monochromatic and tunable. To achieve this, a radioactive source in combination with an accelerator were selected. The experiment was performed with the 3 MeV Dynamitron electron accelerator at Brookhaven National Laboratory. A program was started to manufacture more intense $\beta^+$-emitting sources than commercially available and the accelerator was modified to accelerate positrons. The monoenergetic beam was achieved by using the "positron moderation" technique in combination with a low energy positron filter to eliminate any high energy contamination in the beam. A new beam line was installed to produce a beam of small spatial and angular spread.

The significant detection enhancement of possible resonant peaks by the use of lithium targets over beryllium targets justified the complications associated with handling lithium. Flat and thin self supporting lithium was obtained commercially and provisions were made to make the targets and transfer them to the target arm and the beam position in inert atmospheres or vacuum.

An important improvement in the quality of the experimental data is possible with high position resolution detectors. Detectors specifically designed for this experiment were built with a dedicated data-acquisition system. The detectors consisted of wire chambers for the position measurement and plastic scintillator for the energy measurement. The wire chambers were required to achieve the desired position resolution over an area too large to cover with position sensitive semiconducting detectors affordably. The wire chambers also act as very effective veto against gamma rays since they are almost insensitive to them. The accurate determination of the position and energy of the scattered particles allows for the full kinematic reconstruction of the event. The experiment can therefore check the data on an event by event basis. This allows the elimination of any events that do not qualify as good coincidences or good singles and is very important when looking for small deviations from a large background. For any increase in the normalized coincidence rate, the kinematics of all the coincidences can be checked.
The combination of a more intense monoenergetic positron beam, a lithium target and good position resolution allowed for the execution of a much more sensitive experiment than the experiments that merely count coincidences and singles and allowed for a detailed investigation of the origin of any possible peaks. The description of the experimental apparatus will be divided along the following lines: the beam line, the detectors and the data-acquisition system.

A. THE BEAM LINE

The beam line is shown in Figure 3.1. The positrons, emitted from a $^{22}$Na source in the source deck of the accelerator, are moderated by a tungsten foil in front of the source. Slow positrons are selected by a low energy filter and are transmitted into the beam tubes of the Dynamitron accelerator. With some possible steering, the beam is focused by a solenoid onto an active object aperture. After a 90° bend in a double focusing bending magnet onto an image aperture, the beam is headed for the target chamber, with some final focusing before the target chamber. After interacting with the target, positrons scatter into the detectors or continue onto the beam dump, interacting with the beam dump scintillators or the germanium detector.

The vacuum in the beam line was typically in the $10^{-6}$ to $10^{-7}$ torr range and was obtained by using four cryopumps: one at the exit of the accelerator, one near the bending magnet, one under the target chamber, and one above the beam dump germanium detector. The beam line was roughed out either with sorption pumps or with a mechanical roughing pump. The last solenoid, the target chamber and the beam dump were in a separate target room from the accelerator, allowing work to proceed in the target room while the accelerator was at high voltage.

The individual components of the beam line are discussed in greater detail below.

1. The Source Deck

The positron beam is generated by a positron emitting $^{22}$Na radioactive source. A small fraction of the emitted positron are thermalized by a tungsten moderator. These low energy positrons are selected by a filter and are injected in the Dynamitron.
1.1. The Source and Moderator

As a source of positrons, various $\beta^+$-emitting isotopes were considered, namely $^{64}\text{Cu}$ ($\tau_{1/2} \approx 12.8$ hours), $^{58}\text{Co}$ ($\tau_{1/2} \approx 71$ days) and $^{22}\text{Na}$ ($\tau_{1/2} \approx 2.6$ years). Since the source had to be installed in an accelerator that needs at least 24 hours to pump out and gas up, $^{64}\text{Cu}$ was not feasible. To ensure the source would last through various shut downs of the accelerator and various experimental runs, $^{22}\text{Na}$ was considered to be a more appropriate.

The source consists of 250 mCi of $^{22}\text{Na}$ held in a metal capsule (see Figure 3.2). The source material, an aqueous solution of sodium acetate [Huo89], was deposited on ashless filter paper to obtain a homogeneous strength source. The filter paper was held in a titanium body with a tantalum plug to maximize the back-scattering of positrons. For safety reasons, the source was sealed with a 5-6 $\mu$m window welded to the capsule. The self absorption in the source and the solid angle of the capsule combined to produce a total positron emission yield of 16%. The active area of the source was 6 mm in diameter.

The positrons were emitted from the source with the characteristic beta decay spectrum with an endpoint energy of $\sim 550$ keV, as shown in Figure 3.3. For the experiment however, a monoenergetic beam is required. To obtain and select a fraction of positrons with small energy and kinematic spread, a "moderator" [Sch88b] was used. A moderator, typically an annealed crystalline material, is used to thermalize positrons and then emit them with a discrete energy from its surface. This effect is caused by the work function of certain materials and the charge of the positron.

The work function of a solid is the minimum energy required to take a charge from inside the material to a vacuum outside the material. It has two components, a bulk contribution (for positrons, the positron chemical potential $\mu^+$) and a surface contribution (the surface dipole barrier $D$). The work function $\varphi_+$ for positrons is therefore

$$\varphi_+ = \mu_+ + D.$$  \hfill (3.1)

The bulk contribution is the sum of the repulsive potentials from ion cores and the attraction of electrons and is positive. The surface dipole barrier is caused by the tailing of the electron distribution into the vacuum. It is positive (directed into the solid) for electrons, but negative for positrons (i.e. out of the solid). This sign difference causes the work function for positrons to be negative in some cases, producing a net potential out of the solid. Under this potential, thermalized positrons near the surface would be
preferentially emitted monoenergetically and normally out of the material. The energy spread is ~75 meV [Lyn85] and is caused by the room temperature energy of the positrons, as seen in Figure 3.4. The width in the angular distribution of emitted positrons, $\theta_{1/2}$, can be estimated from

$$\theta_{1/2} = \left[ \frac{k_B T}{\varphi_+} \right].$$

The measured angular distribution of moderated positrons is shown in Figure 3.5.

The moderator used for this experiment was a single crystal of tungsten of 1 μm thickness. The material and thickness had been determined (see Gramsch et al., [Gra87]) to be the most suitable for this application. Unfortunately, the efficiency of positron moderators is very low, typically $10^{-4}$ to $10^{-3}$. Strong radioactive sources are therefore needed to obtain beam of $10^5$ to $10^6$ positrons, hence the 250 mCi source used in this experiment.

Using a moderator, a monoenergetic source of positrons is obtained with low angular spread. However, there is still a large component of unmoderated positrons present, as shown in Figure 3.6. These positrons must be eliminated and this will be done by using a filter.

1.2. The Positron Filter

The aim of the positron filter is to select the monoenergetic low energy positrons and to eliminate the high energy unmoderated positrons that might contaminate the beam. The chosen method, for spatial, weight and other practical consideration related to installing this filter inside a 3 MeV source deck, was an electrostatic filter ([Aso89], [Huo89b]), as shown in Figure 3.7.

The filter works by focusing the low energy positrons through a small aperture (~6 mm) at the bottom of the the filter into the beam tubes of the Dynamitron. The electrostatic focusing is obtained with a modified Soa immersion lens and a pair of symmetric Einzel lenses, all at a few hundred volts. The high energy positrons are therefore almost unaffected by the filter voltages and the fraction of them entering the accelerator is given by the solid angle subtended by the 6 mm aperture. The filter was designed to be as long as possible and its length was 750 mm, leading to a filter factor of $10^{-4}$. 
The trajectories through the filter of slow positrons emitted from the moderator are shown in Figure 3.8. The exit aperture of the filter leads to the Dynamitron beam tubes by way of another set of electrostatic lenses originally installed on the source deck for electron beams. These lenses, called extractor and lens, focus the beam and inject it into the accelerator, as shown in Figures 3.9 and 3.10.

2. The Dynamitron

The Dynamitron is a 3 MeV electrostatic accelerator built by Radiation Dynamics designed for high power electron beams. It was converted to accelerate positrons from the source described in the previous section.

The principle of operation of the Dynamitron [Cle60] is similar to that of a Cockroft-Walton accelerator: a set of capacitors and rectifiers are fed an AC current and are applied in series to produce a large total voltage difference. The unique feature of the Dynamitron is that the rectifier system is driven in parallel from an RF oscillator. As shown schematically in Figure 3.11, the corona rings act as spark suppressors and provide the capacitance to couple the RF to the rectifiers. The other electrode to form the capacitors are large semi-cylindrical "antennas" fixed against the inside of the pressure vessel. These electrodes, in parallel with an inductance, form a resonant LC circuit tuned to the frequency of the RF generator.

Since the accelerator is powered by rectified RF, there is a small RF ripple voltage on top of the large DC voltage. The specifications of the machine put the ripple at < 1 keV, and initial measurements seemed to confirm this figure. However, later tests proved that the beam width, as measured using the width of the peak in the beam dump germanium detector, was inaccurate because the beam optics caused dispersion in the beam and the size of the germanium detector acted as an aperture selecting a certain energy part of the beam. To obtain the true energy width of the beam, the germanium detector was gated on different phases of the RF and the beam tuned onto the detector. The results for different energies are shown in Figure 3.12. The energy width of the beam at 2250 keV was therefore estimated to be ~ 9 keV.

A study of single-quantum annihilation of positrons with shell-bound electrons [Pal91] using the dynamitron provided a sensitive measure of the beam width. The group derived a beam width of less than 1 keV at 1 MeV. This is not necessarily in contradiction with the previous measurement since that method measured the beam energy extremes but
did not measure the relative intensities at the various energies. It is therefore likely we have a 9 keV energy spread with most of the beam within a central 1 keV.

3. The Beam Line

Once the beam leaves the accelerator, it travels down the evacuated beam line towards the target chamber, undergoing focusing, steering and being restricted by active collimators.

Since the beam consists of positrons, the focusing elements needed require only small magnetic fields and the quadrupole doublets or triplets typically used in nuclear accelerators were not usable. Instead, focusing solenoids were installed.

The focusing effect of a solenoid is produced by the fringe fields at either end of the solenoid: for a charged particle of mass $m$ with energy $eV$ traveling through an axial field $B_z$ at a radial distance $r$ from the beam axis, it can be shown [Elk70] that the equation of motion is given by:

$$\frac{d^2r}{dz^2} = -\left(\frac{e}{8mV}\right)B_z^2r. \quad (3.3)$$

The fringe field force on the particle is proportional to the distance from the beam axis and negative, hence producing a focusing effect.

Three focusing solenoids were installed, one at the exit of the accelerator, one after the bending magnet and one in front of the target chamber, as shown in Figure 3.13.

A 90° bending magnet was needed to bend the beam because the Dynamitron is a vertical accelerator (see Figure 3.1). Using its dispersion property, a bending magnet is also the standard method of selecting certain energy particles from a beam. The remaining high energy contamination of the beta decay spectrum passing through the filter was eliminated by the bending magnet using apertures.

The bending magnet used in the beam line was a double focusing [Car87] bending magnet. As sketched in Figure 3.14a, a standard bending magnet focuses the beam in the bending plane but not in a plane perpendicular to the bending magnet because the dipole magnetic field is constant between the poles of the magnet, neglecting edge effects. By rotating the pole face to a non-normal angle with respect to the reference trajectory, as shown in Figure 3.14b, the focusing effect of the bending magnet can be diminished. The
fringe fields of the rotated pole faces add a focusing effect in the plane perpendicular to the bending plane. The beam profile can be preserved by choosing angles $\beta_1$ and $\beta_2$ such that the focusing in both planes is equal. When $\beta_1 = \beta_2 = 26.57^\circ$, the resulting magnet is a double focusing bending magnet.

If the solenoid focuses the beam to a point at a distance $2p$ from the entrance of the bending magnet of radius $p$, an image point will be formed at a distance $2p$ from the exit of the magnet. As in optics, apertures can be positioned at these foci to eliminate stray beam and particles of different energy that, because of dispersion in the magnet, will not focus at the image point.

The object and image apertures are located at the entrance and exit focus of the bending magnet, respectively. They are active collimators, made of plastic scintillator with different sized holes for the beam. The plastic scintillator is connected to a photomultiplier tube by a lucite light pipe (see Figure 3.15). By minimizing the count rate on the object and image slits, the beam can be reproducibly directed through the apertures.

To ensure the beam follows the reference trajectory, an additional aperture was located above the first solenoid. The fraction of the beam hitting this collimator was monitored by a NaI detector located on the outside of the beam line.

<table>
<thead>
<tr>
<th>Table 3.1: Parameters of the Dynamitron positron beam</th>
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<tbody>
<tr>
<td><strong>$^{22}$Na SOURCE ACTIVITY</strong></td>
</tr>
<tr>
<td>e$^+$ YIELD FROM SOURCE</td>
</tr>
<tr>
<td>MODERATOR EFFICIENCY</td>
</tr>
<tr>
<td>TRANSPORT EFFICIENCY</td>
</tr>
<tr>
<td>BEAM INTENSITY ON TARGET</td>
</tr>
<tr>
<td>HIGH ENERGY CONTAMINATION</td>
</tr>
<tr>
<td>BEAM SIZE, FWHM</td>
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<tr>
<td>BEAM DIVERGENCE HALF ANGLE</td>
</tr>
<tr>
<td>BEAM ENERGY WIDTH, FWHM</td>
</tr>
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</table>

The beam was tuned through the apertures using a combination of electrostatic means (the filter, lens and extractor voltages) and magnetic means (the focusing solenoids, the bending magnet, and the magnetic steerers).
Not shown in Figure 3.7 are a pair of electrostatic steerers in the middle of the central filter lens. They are the first steering element in the beam line. As shown in Figure 3.13 there are later sets of steerers to tune the beam through the apertures and to hit the target on axis. To ensure that the beam always followed the same trajectory, the beam was tuned to minimize the amount of "doglegging" needed with the steerers. The beam was tuned until the same steerer settings (except for the centering on the target) could be conserved over the entire energy range. The steerers were then left unchanged during the run.

Using a beam optics program, the beam profile through the beam line can be calculated. This was done using the TRANSPORT program. The resulting beam envelope is shown in Figure 3.16.

A summary of the beam properties is shown in Table 3.1. From Table 3.1, one can see that rigorous beam requirements have been accomplished. The on-target beam intensity was routinely over 500 000 positrons per second, in a continuous mode. Noteworthy is the very low level of contamination achieved by a combination of the moderator, the low energy filter and the beam apertures. The high energy positron rate was measure to be less than 1 per second, down to the parts per million range. This was achieved while keeping the beam spot and divergence small.

4. The Target Chamber

The target chamber (Figure 3.17) was designed with a diameter large enough to support the detectors, and long enough to fit a large cryopump on its bottom. It was fitted with three large ports for easy access, one for target viewing, and five for various arms, holders or view ports. The detectors used to perform the experiment were mounted on the target chamber, as was ancillary equipment. The detectors will be described in detail in a later section. The other equipment in the target chamber is discussed here.

To allow the scattered electrons and positrons to pass from the vacuum of the beam pipe into the detectors kept at atmosphere, thin vacuum windows were needed. Thin windows made of light elements were required to minimize the energy loss of the positron and electrons, and more importantly, to minimize multiple scattering. Since the detectors demand matching area windows, it was realized that some mechanical support would be required for the thin windows. This was provided by 0.8 mm steel wires strung every
12.7 mm inside the vacuum across the front of the vacuum window. The window material was Kapton, chosen for its strength and low average atomic number.

Vacuum tight windows of 25 μm thickness were made successfully, but 50 μm thick windows were installed as a safety precaution. The energy loss in the window was typically 10 keV. Because Kapton is rather permeable to water vapor the vacuum in the target chamber was typically 5 x 10^{-6} torr.

As discussed earlier, the best target for this experiment is a lithium target. Given the beam intensity, a target thickness of ~ 1 mg/cm² was required to obtain an adequate event rate. This thickness could not easily be evaporated and would have had to be rolled. Luckily, thin lithium is now available industrially for use in batteries, and the Foot Minerals Company provided some sample of 1.5 mg/cm² rolled lithium. This was the thinnest rolled lithium available and was used as the target in the experiment.

The target was mounted on an adjustable vertical target arm. The repeatability of the vertical position was checked and found to be better than 0.1 mm. Since the lithium target could not be exposed to air, it was prepared in an argon atmosphere in a glove box. It was then mounted onto the target arm that had been lowered through a port into the glove box. The target arm was then raised into its vacuum tight container and sealed with a gate valve. This assembly was transferred onto the target chamber and attached to an airlock. The argon and the air from the airlock were pumped out and the gate valves opened.

A separate port pointing to the target position was used to insert and extract the scintillator calibration source by means of another airlock (see Figure 3.17).

One of the large side ports (not shown in Figure 3.17) housed a retractable beam profile imager. Conventional glowers, like quartz could not be used because the low positron beam intensities (< 10^{-4} nA) produced insufficient light. A micro-channel plate with a phosphor glower was used. The beam was sent into the channel plate, scattered electrons were multiplied in the channel plate, and the secondary electrons produced a glow on the phosphor. By adjusting the voltages on the channel plate, the brightness of the image could be adjusted and blooming of the beam spot prevented.

After each energy change, the channel plate was rotated into position a few centimeters upstream of the target and the focus of the last solenoid was fine tuned by viewing the spot and minimizing its size. At any other time the channel plate was rotated out of the beam.
5. The Beam Dump

When the beam has passed the target and detectors, it enters the beam dump. The obvious function of the beam dump is to absorb the beam without affecting the experiment, but it was also used to monitor the beam intensity and beam energy. The energy was measured by a germanium detector on axis at the end of the beam line. Even if a large germanium detector caught the entire beam, it could not handle the rate. Therefore, it was preceded by plastic scintillators acting as active collimators.

The originally projected beam intensities were $2.5 \times 10^6$ positrons/sec, requiring multiple detectors to count the positrons. There were two planes of plastic scintillators acting as successive collimators. Each plane consisted of a layer of scintillator with a central hole to allow transmission of the central part of the beam. The first plane was designed to catch the large angle scatterings and had a relatively large aperture ($\varnothing = 19$ mm) to limit the count rate. A diagram of the large aperture beam dump scintillator is shown in Figure 3.18. The second plane had a small aperture ($\varnothing = 3$ mm) that only let through beam that would project onto the germanium detector. Each plane was divided into two wedged scintillator-photomultiplier tube assemblies to reduce the event rate for each detector. One plane was wedged horizontally, one vertically, providing some positional information about the beam.

The final element of the beam line was a 7 mm thick by 80 mm$^2$ windowless ion-implanted high-purity intrinsic germanium (HPGe) detector manufactured by ORTEC. It was used to monitor the beam energy of the positron beam. The detector was cooled by a liquid nitrogen cold finger and was connected to a preamplifier located outside the detector vacuum flange. The signals from the preamplifier were amplified in an ORTEC 572 Amplifier, digitized in a LeCroy 3512 Buffered ADC, and histogrammed in a LeCroy 8810 Histogramming Memory [LeC92].

Before each data point, the target would be raised and the beam energy measured in the detector after the beam was tuned. The count rate was then typically $\sim 20,000$ counts / second. During a run the beam energy, minus the energy loss in the target, could be monitored. The absolute energy of the beam was derived from a detector calibration.

The germanium detector was calibrated daily using a variety of electron and gamma sources, as given in Table 3.2. A typical calibration spectrum is shown in Figure 3.19.
The $^{137}$Cs source was positioned in a cross in the beam line pointing towards the germanium detector and was not removed during the running of the experiment. The $^{56}$Co and $^{228}$Th sources were positioned on top of the germanium, outside the beam line, and were removed after every calibration. The difference in energy between the electron and the gamma ray from the cesium source was used to monitor the build up of materials on the face of the detector. As the ice on the cooled germanium crystal thickens, the electron lose more energy in the ice layer and the difference between the electron line and the gamma ray increases.

The five gamma ray peaks were integrated, and the centroids were calculated. A line was fitted to the energy versus channel plot (see Figure 3.20), providing calibration coefficients.

Table 3.2: Sources used for germanium detector calibration

<table>
<thead>
<tr>
<th>Source</th>
<th>Particle</th>
<th>Energy (keV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{137}$Cs</td>
<td>e$^{-}$</td>
<td>624.2</td>
</tr>
<tr>
<td>$^{137}$Cs</td>
<td>$\gamma$</td>
<td>661.6</td>
</tr>
<tr>
<td>$^{56}$Co</td>
<td>$\gamma$</td>
<td>846.8</td>
</tr>
<tr>
<td>$^{56}$Co</td>
<td>$\gamma$</td>
<td>1238.6</td>
</tr>
<tr>
<td>$^{228}$Th</td>
<td>$\gamma$</td>
<td>1592.5</td>
</tr>
<tr>
<td>$^{228}$Th</td>
<td>$\gamma$</td>
<td>2614.5</td>
</tr>
</tbody>
</table>

B. THE DETECTORS

The main detectors used in the experiment were scintillators and multi-wire proportional chambers, as shown in Figure 3.21. Ideally, one would use one detector to measure the position and energy of the particles, and recently, this technology has become available with the development of position sensitive semiconducting detectors. Since these were not available at the time of the design of this experiment, a large array of small semiconducting detectors was suggested. The position sensitivity would be provided by the detector segmentation. To obtain the desired resolution ($<1^\circ$ in $\theta$ and $\varphi$), an excessive number of detectors would have been required. A combination of wire chambers and scintillators was therefore chosen.
The wire chambers were used to measure the position of a particle while not affecting its energy. Electrons and positrons would typically deposit a few keV of energy while crossing the wire chamber, enough for the position to be determined with great accuracy. The wire chambers being gas filled detectors were very insensitive to gamma rays and hence acted as a filter for gamma rays, an advantage over semiconducting detectors. The scintillators were located behind the chambers and measured the energy of the particles by stopping them in plastic scintillator.

There were four planes of detectors, assembled in a pyramid geometry to maximize the solid angle covered (see Figure 3.22). The wire chambers were therefore trapezoidal, as were the scintillator detectors. As shown in Figure 3.22, each scintillation plane was split into two individual detectors to optimize the plastic scintillator to light collection area ratio. The detector specification, operation and calibration is detailed below.

1. The Scintillators

There were eight plastic scintillation detectors, four identical ones backing one half of the chambers, and four mirror image detectors backing the other half. Each detector consisted of 10 mm of plastic scintillator, glued onto a lucite light pipe coupling to a Hamamatsu R1307 photomultiplier tube (see Figure 3.23). Each scintillator assembly was covered by a thin (12.5 μm) aluminized mylar foil to improve the light collection and the sides of the detectors made light tight with black masking tape. The light pipe was optically coupled to the photomultiplier tube with optical grease, and each pair of scintillators was held by a common aluminum frame. The detectors were built by Bicron and were operated near +1200 Volts.

The scintillators were used mainly as energy measuring detectors, but also provided the timing information for coincidences.

1.1 The Energy Measurement

A typical raw scintillator $^{207}$Bi calibration spectrum is shown in Figure 3.24. The sharp peak near channel 30 is the pedestal and must be corrected for. The two broad peaks, the $^{207}$Bi electron conversion lines, were used to calibrate the detector and corrections were made for nonlinearity in the energy measurements and for varying light collection across the face of the detector.
The ADC used to measure the scintillator pulse height is a charge integrating ADC. If an ADC channel has a small dc offset, a certain charge can build up while the input gate is open, without a real signal being applied to the input. Worse, if a signal is of the opposite polarity of the dc offset, the capacitor could end up with little or no net charge from the real pulse. To avoid such events, a small input offset current (0.03 pC per ns in this case) is applied to ensure that zero net charge corresponds to a small channel number, i.e. the pedestal. This pedestal is easily determined from running the ADC without any signals entering it. The channel with all the events will then be the pedestal (see Figure 3.25).

To eliminate any possible effects of drifts in the input offset current or dc offset in the amplifiers, the pedestal is subtracted before a channel number is used to determine an energy. The pedestal subtraction is also performed for the other ADC signals used for event reconstruction in the experiment (anode position and front window).

1.2 The Bismuth Calibration

The scintillators are calibrated using a 1 µCi ISOTOPE PRODUCTS 207Bi calibration source. 207Bi is used for its 481.7 and 975.6 keV electron-conversion lines. During calibration, a source is inserted at the target position as shown in Figure 3.26. It ensures that the electrons from the source travel through the same target chamber, wire chamber, and scintillator windows as positrons and electrons from the target. The only error as far as energy loss is concerned is the thickness of the source and source window. The source is coated with a few hundred µg/cm² of clear lacquer, producing a systematic error of ~ 0.5 keV, small compared to the resolution of the scintillators (> 150 keV above 2 MeV).

A typical calibration spectrum is shown in Figure 3.24. The 481.7 and 975.6 keV peaks are clearly visible around channels 230 and 390. As shown in Table 3.3, the 481.7 and 975.6 lines are the K-shell conversion of the 569.7 and 1063.6 keV gamma rays, respectively. The L-shell electrons from these gamma rays (554.4 and 1048.1 keV) also contribute to the spectrum. However, given the relatively poor resolution of plastic scintillators, the K-shell and L-shell electrons cannot be discerned. The 1682.2 keV line is too weak for daily calibrations, but it was measured in long calibration runs. The flat background between the peaks is due to self-absorption in the source and contamination in the bismuth 207 source itself. The slow rise around channel 200 is not the tail of the lower peak but is the discriminator cut off.
Table 3.3: Bismuth 207 electron energies and intensities.

<table>
<thead>
<tr>
<th>$E_e$ (keV)</th>
<th>$E_\gamma$ (keV)</th>
<th>Shell in Pb</th>
<th>$I_e$ relative</th>
</tr>
</thead>
<tbody>
<tr>
<td>481.7</td>
<td>569.7</td>
<td>K</td>
<td>21.7</td>
</tr>
<tr>
<td>554.4</td>
<td>L</td>
<td></td>
<td>6.08</td>
</tr>
<tr>
<td>975.6</td>
<td>1063.6</td>
<td>K</td>
<td>(100)</td>
</tr>
<tr>
<td>1048.1</td>
<td>L</td>
<td></td>
<td>25.3</td>
</tr>
<tr>
<td>1682.2</td>
<td>1770.2</td>
<td>K</td>
<td>0.344</td>
</tr>
<tr>
<td>1754.4</td>
<td>L</td>
<td></td>
<td>0.056</td>
</tr>
</tbody>
</table>

A spectrum is obtained for each of the eight scintillators. These spectra are then fitted to extract the peak channels, as shown in Figure 3.27. The fitting algorithm used is a Marquard least squares fit, as defined in Numerical Recipes. The best fit was found using four gaussians (one for each of the four most intense electron lines in Table 3.3), a Compton distribution for the 1063.5 γ line, and a flat background. The function fitted to has 9 parameters, defined as follows:

- $a_0$ is an overall normalization.
- $a_1$ is the 975.6 line gaussian mean.
- $a_2$ is the 1048.1 line gaussian mean.
- $a_3$ is the 481.7 line gaussian mean.
- $a_4$ is the 554.4 line gaussian mean.
- $a_5$ is the resolution at 1000 keV.
- $a_6$ is the amplitude of the 1063.6 keV Compton.
- $a_7$ is the linear background intercept.
- $a_8$ is the linear background slope.

The width of each gaussian is scaled from the resolution at 1000 Kev using the $1/\sqrt{E}$ scaling law for scintillator resolution, where $E$ is the particle energy. The relative intensities of the four gaussians are fixed as given in Table 3.3, and the functional form for the Compton distribution is, from Siegbahn [Sie65],
\[
\frac{d\sigma}{dT} = \frac{\pi r_0^2 mc^2}{(hv - T)^2} \left\{ \left[ \frac{mc^2 T}{(hv)^2} \right]^2 + 2 \left[ \frac{hv - T}{(hv)} \right]^2 + \frac{hv - T}{(hv)^3} \left[ (T - mc^2)^2 - (mc^2)^2 \right] \right\}
\]  

(3.4)

with \( T \) the electron energy, \( hv \) the gamma ray energy, \( mc^2 \) the electron rest mass, and \( r_0 \) the classical electron radius.

The fit is performed from below the edge of the 481.7 peak to above the 975.6 peak, to avoid distortions from discriminator cut off at the low end and high energy background at the high end. Once the parameters are calculated, the calibration for each scintillator is the straight line determined by 481.7 and 975.6 keV vs \( a1 \) and \( a3 \) respectively.

1.3 The Corrections

If a scintillator is calibrated as described above, the spectrum of Figure 3.24 is obtained. If however, a scintillator is calibrated with a cut on the position of the event in the wire chamber so that only events near the middle of the scintillator are accepted, a different spectrum is obtained. The new spectrum has narrower peaks, meaning better resolution. The improvement implies that the scintillation detector has a different response for different positions, causing the smearing of both peaks once electrons hitting all parts of the scintillator are histogrammed together. This uneven response is most likely due to the changing efficiency in light collection at the photo-multiplier tube due to the odd shape of the scintillator and light pipe (see Figure 3.23).

To remedy this nonuniformity in the detectors, their response was mapped using bismuth 207 spectra and cutting on one small section of the detector at a time, using the positional information provided by the wire chambers. This produced a two dimensional table of scintillator pulse height versus position as shown in Figure 3.28. One such map was made for each of the scintillators, and each event's energy was subsequently adjusted given its position. The final spectra are as shown in Figure 3.29.

The accuracy and linearity of the energy measurement can readily be checked using the beam itself. Since the beam energy width is small compared to the scintillator resolution, one can perform a wide energy scan and monitor the scintillator energy of Mott scattered beam positrons and compare it to the beam energy as measured by the germanium detector. Since the events, the detectors, the acquisition system and the analysis program
used in this measurement are the same as the actual experiment, the linearity of the entire system is tested.

As shown in Figure 3.30, the system is noticeably nonlinear, with a significant worsening at higher energies. It was realized that the high energy leveling off was caused by the punch-through of positrons in the plastic scintillator. The plastic was measured to be 1 cm thick as supplied by Bicron, compared to the 1.27 cm specified. To correct the nonlinearity, a cubic was fitted to the curve of Figure 3.30 and all energies were adjusted accordingly.

1.4 The Time Measurement

The timing information of an event is provided by the timing output of the Bicron scintillator photomultiplier tubes and is recorded in a CAMAC TDC. The start is provided by the trigger, discussed in section C.2, the stop by each scintillator signal. For a coincidence, the time difference between both events is the difference between the TDC outputs of the scintillators that fired.

A time difference histogram for coincidences is shown in Figure 3.31. The peak has a full width at half maximum (FWHM) of 11 ns, and there clearly are very few accidental or background coincidences. The accidental rate is expected to be low, considering the singles rate is ~ 300 Hz. With a 45 ns gate, the accidental rate should be ~0.004 Hz, leading to a dozen accidental coincidences per data point.

For singles events, the timing signal measures the delay between a trigger and its own stop. Since the trigger is generated by the scintillator signal itself, the singles timing is merely a measure of the delays and processing times introduced in the electronics and has no physical significance. Therefore, timing information is only relevant for coincidences.

2. The Wire Chambers

The position sensitive detectors were four trapezoidal multi-wire proportional chambers [Leo87] and used a combination of resistive division and delay line readouts to provide two dimensional positional information.

In order to maximize the solid angle subtented by the detectors, the sides of the chambers were kept to a minimum of ~ 12 mm. Since the nonparallel sides of adjacent wire chambers are very close to each other when installed in position on the pyramid-shaped target chamber, there is no room for feedthroughs or connectors from individual
wires to outside electronics. The anode wires, whose contacts are on the nonparallel sides, were therefore read out by resistive division, a method that requires only two outside contacts. That is the origin of the unusual combination of resistive division and delay line read-outs.

Each chamber consisted of a G-10 frame covered by thin aluminized mylar windows (8 μm). The chamber frame actually consisted of three frames glued together, two of which supported the sensing wires and contacts for the read-out electronics (see Figure 3.32). The frames were glued together without O-rings or provisions for replacing broken wires or electronics in order to keep the sides of the chambers narrow, as described above.

The front window served as a front cathode and was used to monitor the chamber gain and to provide a trigger of the chamber firing. Since the window was also a cathode, a constant anode-window spacing had to be maintained to produce an equal electric field across the face of the detector. To minimize any bowing of the window, it was reinforced with a second window and a thin nylon support wire along the axis of the symmetry of the detector, matching the dead region caused by the junction of the scintillators. The front window was grounded through the preamplifier.

The anode plane consisted of 39 gold plated tungsten wires of 20 μm diameter with a 1.95 mm (1/13 inch) spacing between them. The first and last wire were replaced by larger (62.5 μm copper-beryllium) wires to provide a smooth electric field transition to the chamber frame. The position sensing occurred by resistive division and the charge from either end of the resistive line was collected in a charge sensitive preamplifier. The ratio of the collected charges is a measure of the position of an event. The anode was operated at a voltage of +2000 Volts that was fed through the preamplifier. This was the highest voltage the chambers could be operated at without sparking. The leakage current was typically ≤ 1 nA and the power supplies trip points were set at ~ 20 nA.

The second cathode plane consisted of 118 sensing wires (of 75 μm diameter copper-beryllium) perpendicular to the anode wires. These wires were bonded in threes onto readout contacts feeding into a 1 μs delay line. Both ends of the delay line were read out and the relative time difference between the signals provided the position information. The second cathode was grounded through the preamplifiers at the end of the delay line.

Between the second cathode plane and the back window was a 6.25 mm (1/4 inch) frame. This thickness was needed for the gas input and output lines, entering laterally into
the frame. The chamber was operated just over ambient pressure with a gas mixture of 80%
% argon and 20% carbon dioxide. A constant flow of gas was kept to flush the chamber
of any contamination. The flow was regulated by bubblers built at Brookhaven National
Laboratory, a set of Matheson flowmeters, and an MKS 250B flow regulator.

The back window was designed to be operated at +200 Volts to sweep away the
charge collected in the region between the second cathode and the back window. This was
desirable to minimize the effective thickness of the detector to maximize the position
resolution for particles entering the detector diagonally. The effective thickness of the
chamber was the separation between the cathodes of 5mm. This corresponds to
0.9mg/cm² of the Ar/CO₂ gas mixture. The energy deposited by a minimum ionizing
electron is then ~ 1.4 keV, and since ~ 26 eV are needed to produce an ion pair in argon,
50~60 ionizations are generated in the gas. Optimal operation of the chamber with the
anode at 1900 V, 0.1 pC was observed at the preamplifier, demonstrating a chamber gain of
~ 12000. It was found desirable to increase the size of the pulses collected to minimize the
effect of lower level discriminators, both in hardware and in software. This was done by
increasing the effective thickness of the wire chamber. The back window voltage was
changed from +200 V to -1000 V, sending the charge created in the back of the chamber to
the anode, producing an increase in collected charge of near 50% (see Figure 3.33).

2.1 The Anode Plane

The anode plane senses the position along the height of the wire chamber (see
Figure 3.34). The position of an event is determined by resistive division [Char79]. This
means each wire is connected to a small resistor, turning the anode plane into a resistor
chain. Typically, the resistive chain is connected to one side of the contacts of the wire
plane. If a wire were to break, the detector would still work but with one less wire. Since
the detectors were designed conservatively but without provisions to replace wires
anyways, the resistors (chip resistors) were imbedded in the chamber frames on opposite
sides of the frames. The printed circuit board tracks are shown in Figure 3.35. Although
the loss of a wire or resistor would prove catastrophic for the chamber, the gain in space on
the side of the detector was considered to outweigh the risks. No chamber was lost in over
two years of operations.

When a particle crosses the detector, it ionizes Argon gas along its path. Freed
electrons cascade in the electric field between the anode and cathodes, resulting in a
measurable pulse on the anode. Secondary electrons collect almost exclusively on the
anode wire nearest the particle trajectory. The collected charge causes a pulse on a particular anode wire. The voltage pulse, $V$, has two paths to the relative ground (the anode reference voltage), currents $I_1$ and $I_2$, one along each direction of the resistor chain, with resistance $R_1$ and $R_2$ respectively, as shown in Figure 3.36. Using Ohm's law,

$$V = I_1 \cdot R_1 = I_2 \cdot R_2$$

and the fact that both pulses are integrated until all the charge is collected at time $T$ say, we obtain, with $I_i = C_i / T$, where $i = 1$ or 2,

$$C_1 \cdot R_1 / T = C_2 \cdot R_2 / T, \quad \text{or} \quad C_1 / C_2 = R_2 / R_1,$$

or

$$C_1 / (C_1 + C_2) = R_2 / (R_1 + R_2).$$

The resistances, $R_1$ and $R_2$, are determined by the position of an event, since $R_1 = n_1 \cdot R$ and $R_2 = n_2 \cdot R$ where $R$ is the value of an individual resistor, and $n_1$ and $n_2$ are the numbers of wires from the edge, as illustrated in Figure 3.36. Therefore, knowledge of the resistances is a measure of the position of an event. The integrated charges, $C_1$ and $C_2$, are the physical observables recorded to derive $R_1$ and $R_2$. Substituting for $R_1$ and $R_2$,

$$C_1 / (C_1 + C_2) = n_1 \cdot R / (n_1 \cdot R + n_2 \cdot R) = n_1 / (n_1 + n_2).$$

Hence, the above ratio of charges measures the fractional distance along the anode of an event.

For an event near an edge, the charge collected, from say the first wire, is $1/40$th as large on one side as on the other side. For a 2000 channel ADC, a pulse near the overflow channel on one side corresponds to a pulse near channel 50, barely above the pedestal, on the other end. To remedy this problem, resistors are added at each end of the resistive chain. The resistor chain consists of 40 800 Ω resistors in series leading to a total impedance of 32 kΩ. By adding $4$ kΩ at each end, the total impedance increases to 40 kΩ, and an event near the edge has a 4 kΩ impedance to its relative ground. The range of charge division is now 10% to 90% instead of the previous 2.5% to 97.5%. The dynamic range of the signals has been reduced by a factor of four.

Besides the range of pulses obtained by charge division, there is a wide distribution of the total chamber signal itself. It is equal to the distribution of energy deposited in the chamber, which is a Landau distribution, as measured by the charge collected on the front cathode. The Landau distribution will be discussed in greater detail in section 4.3, but it suffices to notice that it is a very broad distribution. The total charge collected on the anode is the sum of the charge from both outputs, and is the mirror charge of the cathode charge. The total charge collected, the denominator in equation (3.8), therefore follows the same distribution as the front cathode, as shown in Figure 3.37. To maximize chamber
efficiency and minimize effects from lower- and upper- end cutoffs, large dynamic range ADCs are needed.

Ideally, logarithmic ADCs or logarithmic amplifier into linear ADCs would be used. Another possibility is to use ADCs with a high number of channels. Neither of the above was practical, instead the following scheme was designed.

Each signal is amplified with a delay line amplifier into a pulse between 0 and 8 Volts. The unipolar pulse is read by an ADC that senses between 0 and 2 Volts (see Figure 3.38), whereas the bipolar pulse is read by an ADC that senses between 0 and 8 Volts (Figure 3.39). Small signals are accurately recorded with the 2000 channels between 0 and 2 Volts. If a signal is larger than 2 Volts, it will produce an overflow in the 0-2 V ADC, but will produce a large signal in the 0-8 V ADC. Using two ADC allows high sensitivity at the low end combined with a large dynamic range.

When using two ADCs to sense one pulse, a cross calibration between the two ADCs is needed, as is a means of choosing an ADC signal for the analysis. The 0-2V ADC channel is preferred unless it overflows, in which case the 0-8V ADC channel replaces it in the calculations. To cross calibrate the ADCs, the two channels measuring the same signal are plotted one versus the other using pulses from real data (Figure 3.40). The distributions do not converge at the origin because of the pedestal phenomenon in ADCs, as discussed in section 1.1. The pedestal channels are subtracted before calculations using the ADC channels. A line is fitted to the two dimensional distribution and the fitted slope and intercept are the cross calibration coefficients.

By measuring the charge from both outputs of the resistor chain, the anode position is calculated using equation (3.8). The ratio of one ADC channel over the sum of both ADC signals measures the fractional distance along the anode of an event. The ratio is multiplied by a large number to allow using integer channel numbers as the uncalibrated anode position (Figure 3.41). The channel numbers transform into anode positions with the calibrations discussed in the next section.

2.2 The Anode Calibration and the Linearity

It is desirable to calibrate the chamber with position markers from the chamber itself to avoid alignment or other errors from an outside calibration. As shown in Figure 3.41, when the wire chamber is uniformly illuminated with $^{55}\text{Fe}$ x-rays, the anode anode wires are clearly visible. Thirty five wires, with long exposures 37, can be counted. The outer
two wires are not visible because they are slightly thicker than the other wires, leading to smaller charge avalanches, leading to final pulses below the preamplifier sensitivity.

The inner 35 wires calibrate the anode readout. Each of the 35 peaks is integrated and the physical wire position plotted vs the centroid channel of the peak (see Figure 3.42). The resulting curve has the typical S shape of resistive division readouts [Fra81]. To calibrate the anode, a cubic was fitted to the S shaped curve for every chamber, and the resulting coefficients were the calibration coefficients.

2.3 The Cathode Plane

The cathode plane measures the position along the base, or long axis of the trapezoidal chamber, as shown in Figure 3.43. A delay line readout is used and the time difference between the pulses at each end of the delay line is a measure of the position of an event along the cathode [Char79]. The printed circuit tracks that were glued into the cathode frame are shown in Figure 3.44. The wire pads were of the appropriate width and spacing to glue three wire to a pad, and the delay line contacts specified to fit into pre-existing 1 μs delay lines.

During the initial testing of the detectors, some events with an erroneous position reconstruction were noticed. These erroneous positions were associated with early stops arriving at the TDC's, also leading to a wrong total delay time. The early stops were caused by small pulses preceding the real pulses and were called the "prepulses". The time difference between the prepulses and the real pulses was found to be related to the anode position of an event. This led to an understanding that these prepulses were due to the peculiar geometry of the detector as explained below.

Since the anode plane is at ~ + 2000 V and the cathode plane is grounded, there is a negative charge induced onto the cathode wires facing the anode wires. The negative charge from an event on a certain anode wire will cause a momentary dip of the voltage on that wire because of the significant impedance to ground caused by the resistance in the resistive division. The change in voltage of the anode wire causes a change in induced charge on the matching cathode wires, leading to a small pulse being generated by the zero-cross amplifiers, as shown in Figure 3.45. When the main pulse arrives, it also causes a signal, but of opposite polarity. In a normal TDC, the real pulse would be ignored because the TDC would already be busy with the prepulse and would have to be cleared to be ready for another pulse. The problem was solved by using multiple hit TDCs that can sense
many stops from one source. The first two stops from each side of the delay line were recorded.

To determine the position of an event, the correct stops must be identified. Since there is no a priori knowledge of an event having pre-pulsed or not, both TDC stops for each signals are always recorded. The first stop (Figure 3.46) contains most good stops, whereas the second stop (Figure 3.47) contains some stops from events with a pre-pulse and a large number of after-pulses. If an event has a good first stop and no second stop, the TDC overflow channel is recorded as the second stop.

In an event where both first stops are the good stops, the sum of both times will be equal to the delay line time, 1 μs, plus an offset for the cables and electronics. This is because for all events, the time for the good pulse to travel to one end of the delay line plus to the other end is the time constant of the delay line itself, i.e. 1 μs (see Figure 3.48). The timing sum can therefore be used as a criterion for finding the good stop. Figure 3.49 shows the sum of first stops for chamber zero. Events with both good stops as first stops are in the main peak, whereas events with pre-pulses are in the low early region between channels 300 and 450. To pick the correct stops, the four sum combinations of the stops are calculated (see Figure 3.50), i.e.:

- the sum of the first stop left and the first stop right.
- the sum of the first stop left and the second stop right.
- the sum of the second stop left and the first stop right.
- the sum of the second stop left and the second stop right.

The correct pair of stops is the one whose sum is closest to channel 500 (see Figure 3.51). The event position is calculated with these stops.

With the correct stops identified, the cathode position is the difference between these two stops. This is shown in Figure 3.52, where a constant is added to ensure that all channel numbers are positive. The physical cathode position can be determined with a timing calibration of the channels, but the position sensing of the chamber itself was used as a more thorough calibration.

2.4 The Cathode Calibration and the Linearity

The linearity of the delay line readouts was measured using a highly collimated stepping motor mounted X-ray source, with 1μm position accuracy and 25 μm beam size. It was scanned, in 1 mm steps, across the bottom of the chamber to measure the linearity
across the entire readout. At each step, a gaussian distribution with the resolution of the readout, 350 µm, was recorded. The channel mean of each gaussian was then plotted vs the stepping motor position. The result is shown in Figure 3.53, demonstrating the linearity of the system over the entire range. Therefore, only a linear calibration is needed.

The uniform illumination of a wire chamber with x-rays yields a distribution that reflects the steel wires in front of the vacuum windows, the pads holding groups of three wires, and the individual wires themselves (see Figure 3.54). Since these have known physical positions, they can be used in calibrating the readout. The pattern of the pads holding three wires is not sharp, hence not suitable for an accurate calibration. The steel wires and the chamber wires each produce many sharp features that could be used for calibrating the slope of the linear calibration, but an absolute reference is needed for the intercept. The position of forty central wires was used to determine the slope (see Figure 3.55).

Since the wire chambers are held in position on the steel vacuum chamber with nylon screws, their position has some uncertainty due to the flexibility of the nylon screws. On the other hand, the steel wire location is accurately known from the position of the holes in the steel vacuum chamber pyramid. Therefore, the steel wires were used in calibrating the intercept. The channel corresponding to the center of each wire was determined, and a line was fitted to the position versus channel line. The resulting coefficients led to an absolute reference position.

2.5 The Chamber Gain

The chamber gain is an important factor in setting the wire chamber efficiency, because it determines the absolute amplitude of all pulses. If the gain decreases significantly, less pulses will pass the discriminators, and if the gain is increased significantly, more pulses will overflow in the ADCs. Given the maximum operating voltage of the wire chambers, amplifier gains are set to optimize the efficiency. Any drift in chamber gain during the experiment would alter the efficiency for recording singles and coincidences and jeopardize the Mott to Bhabha normalization in the experiment.

The chamber gain is adjusted by changing the anode voltage, as shown in Figure 3.32. As can be seen in the figure, an important factor affecting the chamber gain is the bowing of the front window. Even with a reinforced front window, the pressure of the Ar/CO₂ mixture affects the gain. The gain dependence on gas pressure is shown in Figure 3.56. The chambers were therefore operated at a flow rate slightly above the lowest flow
rate possible to avoid gas contamination, or \( \sim 1.5 \) bubbles per second. With such a flow rate, the gain variation with anode wire is very small, as shown in Figure 3.57.

The coincidence efficiency is not necessarily the square of the singles efficiency because singles and coincidences have different angular distributions, leading to a different fraction of events in the dead regions around the edge of the chamber. Pulses near an edge of the anode have small signals on one end of the resistive division and large signals near the other end. This makes them prone to lowered reconstruction efficiencies for both increases and decreases in gain. Since singles and coincidences have different angular distributions, the fraction of events near the edges of the anode are different, leading to differently changing efficiencies.

To monitor the chamber gain, the front cathode signal is used. As discussed in section 2.1, both the anode signal sum and the front cathode provide a measurement of the energy deposited in the chamber. The anode sum can show artifacts due to one of the two signals added being a pedestal or an overflow (see Figure 3.58). The artifact can be corrected, but the front cathode is still the more direct signal to use. After corrections for pedestals and overflows, the anode sum distribution provides a double check of the front cathode measurement. The delay line readout cathode signals cannot be used, since they mostly measure timing, apart from the discriminator monitoring signals that only monitor the very bottom of the Landau distribution and would not provide a good measure of the overall gain.

C. THE DATA-ACQUISITION SYSTEM

The data acquisition system can be considered as consisting of three separate parts: the front end electronics that processes the signals from the detectors, the CAMAC units that digitize the signals, and the computer that analyzes the digitized data and writes it to tape. The system was designed and assembled by Sean McCorkle, Jim McDonough and Stuart Henderson.

1. Signal Processing

Before the signals from the detectors can be digitized, they might have to be amplified, shaped, delayed or inverted. Logic signals for trigger and counting purposes must also be produced. The scintillator and wire chamber signals were handled differently, as explained below.
A schematic diagram of the scintillator pulse electronics is shown in Figure 3.59. The scintillator output was large enough and did not need preamplification. It was fed into a LeCroy 428 Linear Fan-In/Fan-Out to produce two identical pulses. One was delayed and digitized for the pulse height, the other one was used for the timing information and the logic circuit.

The second signal was amplified using a LeCroy 612 AM Photomultiplier Amplifier to suit better in the range of the discriminator used, a LeCroy 621 Leading Edge Discriminator. One discriminator output was sent into the TDC to provide the timing information of the event. Another discriminator output was sent to the input register and to a scaler to count the event rate in the scintillator. The other discriminator outputs were used for the trigger logic discussed separately.

The wire chamber signals were processed quite differently from the scintillator signals (see Figure 3.60). All five signals are fed into preamplifiers before being amplified.

The cathode signals, used to reconstruct position from the time difference, are amplified in an amplifier zero-cross discriminator designed and built by the instrumentation division of Brookhaven National Laboratory. These units amplify the preamplifier output into a bipolar signal and produce a NIM pulse at the time the bipolar signal crosses zero. The NIM pulse is used in the TDC’s. The units also had an output providing the amplified signal and this was recorded to monitor the stability of the discriminator level in the unit. The discriminator was typically set at ~ 30 mV.

The anode signals were amplified in Ortec 460 delay line amplifiers and then fed into the ADC’s. Because the low gain and high gain ADC’s had opposite polarities, inverters, Ortec 433 A Inverters, had to be used on the bipolar output of the amplifiers.

The front cathode was used to monitor the chamber gain and to provide a logic pulse if the chamber fired. The preamplifier output was amplified in a timing filter amplifier, an Ortec 474 Timing Filter Amplifier, and one output used to record the pulse height in the ADC. The other signal, put through a LeCroy 621 Leading Edge Discriminator, produces a logic signal used for the timing information, the input register and the scalers.
2. The Trigger Logic

The event triggers were provided by the scintillator signals. The data acquisition system was set up to accept two different triggers: a singles trigger and a coincidence trigger (see Figure 3.61).

A single event is defined as one scintillator firing with no scintillator signals from any other quadrant within 45 ns. A scintillator is considered to have fired if the discriminator connected to one of the amplified scintillator outputs send out a NIM pulse. The discriminator threshold levels are adjusted so that the signal from particles with more than 350 Kev causes it to produce a pulse. Within one quadrant, both scintillator discriminators are allowed to fire. In fact both discriminator outputs enter a logic unit which produces an output if either discriminator sent out a pulse. This logic is used to take into account the cases where a positron enters one scintillator, leaves enough energy to fire its discriminator, and then multiple-scatters into the adjacent scintillator, also setting off that one.

The coincidence trigger provides a signature for Bhabha or resonant events. Since in these events both the electron and positron are scattered back-to-back, the coincidence trigger requires two scintillators firing in different quadrants within 45 ns. Using the back-to-back property, coincidence events are later checked to be coming from opposite quadrants, hence triggers from adjacent quadrants are only used in background studies. Coincidence hits in the same quadrant cannot kinematically be Bhabha or resonant events even after including multiple scattering in the target.

At certain times during the experiment it is necessary to disable the electronics in order to allow signals already in the acquisition system to be processed or to avoid saturating the computer. In this experiment, the statistical limitation is the number of coincidences observed, whereas singles are used for normalization. It can therefore be advantageous to limit the number of singles recorded so as to limit the amount of dead time, while still recording enough so as to keep the statistical limit in singles small compared to that in coincidences. This is done by inhibiting the singles trigger for 70 μsec every 100 μsec (see Figure 3.62). In order to monitor the fraction of events inhibited by this so-called pre-scaling process, the singles rates, with and without the pre-scaling, are measured in a scaler (see previous section). As a double check, a clock rate is also continuously measured with and without the pre-scaling. These two ratios are usually within 0.01% of each other.
When two events occur in rapid succession, the data acquisition system might still be handling one event when the signals from the next event already enter the ADC's or TDC's. This would lead to erroneous event reconstruction. In order to avoid such events, the CAMAC crate's Starburst sends out a busy signals that inhibits the triggers until all CAMAC units have processed the event. This leads to a certain fraction of dead time, which is monitored in an identical fashion to the pre-scale fraction discussed above.

3. The CAMAC Electronics

Once an event trigger has been generated, it is sent to a Starburst unit in the CAMAC crate. The Starburst determines what CAMAC units must be read out. The ADC's, TDC's and scalers are then read out and buffered in the Starburst before being sent to the data acquisition and analysis computer. The Starburst, containing a DEC J-11 CPU microprocessor, acts as a buffer between the ADC's and TDC's and the data acquisition computer and is programmed like PDP11/70. It has two external interrupts which are used for the singles and coincidence triggers, respectively. For each of these triggers, the starburst is programmed to read the appropriate ADC's and TDC's and store the values in a buffer. When a buffer is full, the starburst starts filling a second buffer, and the first buffer is sent to the μVAX.

3.1 Signals Recorded

Depending on the trigger, different signals from the CAMAC crate are recorded. In addition to recording singles and coincidence events, the data acquisition also monitors the scalers and the germanium detector.

For a singles event, only the signals needed to reconstruct that event are recorded. That means:
- the input register where gets written which scintillators and which chambers fired.
- the scintillator ADC from which the energy will be calculated.
- the scintillator TDC providing timing information.
- the anode ADC signals, both left and right side, high gain and low gain from which the position along the anode will be obtained.
- the ADC signals monitoring the discriminators on the zero cross amplifier of the cathode read out.
- the front cathode ADC from which the chamber gain will be obtained.
- the cathode TDC signals, first and second stop, form the left and the right side of the delay line.
For a coincidence trigger, all signals from the event are recorded. The list can be summarized as follows:

- the input register.
- all scintillator ADC signals.
- all scintillator TDC signals, from which the relative timing between two events in a coincidence can be obtained.
- all anode ADC signals.
- all cross amplifier monitor ADC signals.
- all four front cathode signals.
- all cathode TDC signals.

The chamber signals recorded are the same as for a single event but all four chambers are recorded instead of just one.

While the data acquisition is recording singles and coincidence events, the scalers in the CAMAC crate monitor the rate of various detectors firing. After every 10 coincidences the scalers get read out. The rates monitored are:

- the beam dump scintillator rates.
- the rate of the sum of both scintillators in each quadrant.
- the object and image slit rates.
- the total beam dump rate and clock rate with and without the data acquisition dead time inhibit.
- the unscaled and pre-scaled trigger rates.
- the coincidence trigger rate.
- the rate of a NaI detector near the target.

After every ten scaler readouts, i.e. after every one hundred coincidences, the data acquisition system records the histogram compiled in the histogramming ADC of the spectrum from the germanium detector.

3.2 Digitizing Hardware

The Starburst collects signals for an event from the appropriate CAMAC units. These were ADC's and TDC's of various kinds, scalers and an input register. Four different analog to digital converters were used, each one needed for a specific application.
The analog signal from the scintillator was digitized by a charge sensing LeCroy 2249A ADC after an 80 nsec delay to allow for the production of a trigger by the logic electronics. The energy of an event was derived from this ADC output.

The voltage sensing LeCroy 2259B ADC was used to convert most wire chamber signals: the front window amplitude, the anode signals, and the bipolar cathode signal from the amplifier zero cross modules of the delay line readout. The front window amplitude was used to monitor the chamber gain, the anode signal were used for position reconstruction, and the delay line signals were used to monitor the discriminator of the amplifier zero-cross NIM modules.

To increase the dynamic range available, the anode signals were also digitized by an Ortec AD811, which accept pulses up to 8 Volts as opposed to 2 Volts for the LeCroy 2259b.

The germanium detector signals, with their high rate, were digitized in a separate ADC, a LeCroy 3512 Buffered ADC, and histogrammed in hardware in an adjoining LeCroy 8810 Histogramming Memory.

The time to digital conversion was performed by two different types of TDC's. Regular TDC's, LeCroy 2228a in this case were used to digitize the scintillator signals and the front cathode signals from the wire chambers. For the delay line readout signals of the wire chambers, multiple hit TDC's were needed, and LeCroy 4208 Multiple Hit TDC's were used. These TDC's recorded both the first and second stop from the same signal.

4. The Data Analysis Computer

The buffers of data from the Starburst, buffers of 4096 16 bit words, are sent to a MicroVAX II via a DRV11 parallel interface using the Q-bus output of the Starburst. Upon arrival in the μVAX, the buffer of data is placed in shared memory. This is an area of computer memory that is set up to be accessible, or shared, between different programs. This allows the data to be analyzed while it is also being written to tape.

When a run is started, the commands to set up the Starburst control program is sent from a dedicated Starburst terminal. The main μVAX control program sets up the communications to the Starburst, the area of shared memory can be cleared if desired, and the data acquisition can be turned on. The optional taping (on 6250 bpi open reel magnetic tapes) of the data is also controlled from the main control program. The entering data is
automatically analyzed by the online analysis program. The control program can set cuts on the histograms generated by the online analysis as well as perform various operations on the data.

The data acquisition software was written in the C programming language by Sean McCorkle from Brookhaven National Laboratory.

D. EXPERIMENTAL PROCEDURE

The previous sections outlined the hardware and software used to perform the experiment. This sections discusses how the experiment was actually conducted and which procedures were followed.

Before the experiment could be started, a number of adjustments were made. The voltage on the wire chambers and scintillators were set to provide large and equal gains while keeping the voltage on the wire chambers sufficiently low to avoid any sparking. All the amplifier gains were equalized for similar signals and tuned to maximize the used dynamic range of the ADC's.

The wire chamber position reconstruction in both directions was calibrated and used in the online data analysis program. The data acquisition system was tuned so that the rate of scaler read-outs and histogrammed germanium read-outs was low to minimize dead time but high enough to avoid overflows in any of the monitored channels.

As mentioned before, the beam was tuned until the same steerer settings (except for the centering on the target) could be conserved over the entire energy range. The steerers were then left unchanged during the entire run. The accelerator was conditioned to ensure that the highest accelerator energies could be reached. The experiment was performed in order of decreasing energy to take full advantage of the optimized initial condition of the machine.

Once the experimental run was started, it was run twenty four hours a day, with an interruption every twenty four hours for calibrations and a check on the accelerator. The scintillators and the germanium detector were calibrated using radioactive sources, as described elsewhere.

In the mean time, the accelerator voltage was increased until it became unstable. It is desirable to have a margin of safety of 100 –200 KV between the operating voltage of the
accelerator and the limit of stability. When this margin was no longer present, the accelerator was conditioned until it was reestablished. The accelerator voltage would then be lowered to the desired energy and allowed to stabilize while the detector calibration was continuing. The entire calibration and condition operation were typically performed in ~3 hours.

1. Procedure at each Energy

After the initial tuning of the beam, the magnets were scaled from energy to energy using the theoretical scaling laws. At each energy, the voltage was lowered to fit the beam through the apertures and the exact beam energy was measured using the germanium detector. The final solenoid was fine tuned to minimize the size of the beam spot on the channel plate. The beam was centered with respect to the detectors by equalizing the scintillator count rate in the four quadrants.

Once the beam was tuned, it was tested for halo before and after each run with a target, and both halo runs were required to be without stray beam for the data point to be accepted. This was almost always the case.

If the halo run was satisfactory, the target would be lowered, and the data acquisition system would be started. The data would be recorded on tape and analyzed on line, with one person controlling the data acquisition system and monitoring a variety of histograms. The accelerator and the beam optics would be controlled by another person who would also monitor on line data displayed on a terminal in the accelerator control room.

When the halo run after a data point had shown that no stray beam was present, the magnets were scaled to the next desired energy and the process was repeated.

2. The "Halo" Run

During a halo run, the target was raised and data was acquired as in a real run. This was performed to determine if there was any stray beam directly interacting with the detectors, i.e. if there was a halo to the beam. The histogrammed data from a halo run is shown in Figure 3.63.

The coincidence rate is very small and is primarily caused by cosmic rays. Because cosmic rays cross the detectors mostly at tangential angles, they deposit a lot of energy in the wire chambers and the scintillation detectors. This causes the ADC signals to overflow.
and the event to be rejected. The experiment therefore has a large intrinsic cosmic rejection efficiency. For this experiment, the coincidence rate from cosmic rates is negligible and will not be considered. In the longer run of this experiment however, cosmic ray veto counters installed for the long lifetime experiment were used and lowered the cosmic ray coincidence rate even further.

The singles rate is isotropic azimuthally and the scattering angle distribution matches the acceptance of the detector. The singles energy spectrum shows two features: an exponential decay at lower energies and a bump at high energies. The high energy bump is due to minimizing ionizing particles that did not overflow any of the signals. The exponential decay is caused by continuous gamma conversion. The good cosmic singles rate was ~1 Hz. A contamination of as little as 0.1 Hz of stray beam in the detectors could be noticed (see Figure 3.64) because it would produce a peak at beam energy where very little background was present.
Figure 3.1: Drawing of the beam line showing the source, the accelerator, the beam line, the target chamber and the beam dump.

Figure 3.2: Detail drawing of the $^{22}$Na source. From [Aso39].
Figure 3.3: Beta decay spectrum from the decay of $^{22}$Na.

Figure 3.4: Energy distribution of positrons emitted from a Ni (100) surface.

From [Sch88b].
Figure 3.5: Angular distribution of positrons emitted from a W (110) moderator. From [Fis86].

Figure 3.6: Comparison of yields for moderated and unmoderated positrons from a $\beta^+$-source, $^{58}$Co in this case. From [Sch88b].
Figure 3.7: Drawing of the positron filter.

Figure 3.8: Calculated trajectories of positrons through the filter. From [Aso89].
Figure 3.9: Calculated trajectories through the extractor and lens assembly.
From [Aso89].

Figure 3.10: Calculated trajectories into the top dynodes of the accelerator.
From [Aso89].
Figure 3.11: Schematic diagram of the operation of the Dynamitron.

Figure 3.12: Beam energy width at different accelerator energies.
Figure 3.13: Detail drawing of the beam line showing the location of the solenoids and the steerers. From [Hen91a].

Figure 3.14: Figure 3.14a shows the trajectories in a standard bending magnet, Figure 3.14b shows the angles $\beta_1$ and $\beta_2$ that can correct the focusing effect. From [Aso89].
Figure 3.15: Schematic diagram of the object or image aperture.

Figure 3.16: Envelope of the positron beam through the accelerator at 2.2 MeV. From [Aso89].
Figure 3.17: Schematic diagram of the target chamber.

Figure 3.18: Diagram of one of the large aperture beam dump scintillators. A second identical scintillator closes the circle and fills the wedge of the first scintillator.
Figure 3.19: Typical calibration spectrum obtained by the HPGe detector.

Figure 3.20: Plot of the energy of a calibration source versus the measured peak centroid of the calibration source line. The straight line fit to the points is the germanium calibration.
Figure 3.21: Schematic diagram of the detectors.

Figure 3.22: Blown up view of the detection system.
Figure 3.23: The scintillator to light pipe to photomultiplier tube assembly. From [Hen91a].

Figure 3.24: Raw calibration spectrum of a scintillator.
Figure 3.25: Typical pedestal, in this case from scintillator number zero.

Figure 3.26: Position of the bismuth source during calibrations. It is at the same position as the target is during a run. From [Hen91a].
Figure 3.27: Fit of the nine parameter function to the bismuth spectrum.

Figure 3.28: Two dimensional scintillator pulse height correction map. Note that the origin is suppressed on the z-axis.
Figure 3.29: Scintillator ADC spectrum after corrections are applied.

Figure 3.30: Energy measured by the scintillator versus beam energy. The nonlinearity at higher energy is due to punch-through in the scintillators.
Figure 3.31: Time difference between the scintillators signals of a coincidence event.

Figure 3.32: Frontal view and cross section of the wire chamber.
Figure 3.33: Charge on the anode of the wire chamber versus applied voltage. Note the increase when the second window was added and when the back window voltage was changed.

Figure 3.34: Resistive division readout for the anode position measurement.
Anode PC Mask

Figure 3.35: Drawing of the printed circuit board tracks for the anode wires and the resistors.

Figure 3.36: Diagram of the resistance to ground for an event on a given wire.
Figure 3.37: Pulse height distribution of signals from the front cathode and histogram of the sum of the anode signals for chamber zero. Note that both distributions are identical, as expected, except for the gain.

Figure 3.38: Pulse height distributions of anode signals measured with the small range ADC. Figure 3.38a is from one end of the resistive chain of chamber zero, Figure 3.38b is from the other end.
Figure 3.39: Pulse height distribution of anode signals measured with the large range ADC. Figure 3.39a is from one end of the resistive chain of chamber zero, Figure 3.39b is from the other end.

Figure 3.40: Two dimensional distribution of pulse height as measured by the two different ADCs. The straight line fit to this distribution is used as the cross calibration between ADCs. Figure 3.40a plots ADC1 and ADC4 from chamber zero, Figure 3.40b is from ADC2 and ADC5.
Figure 3.41: Uncalibrated anode position histogram for uniform X-ray illumination.

Figure 3.42: Plot of the channel of the centroid of each wire peak versus the wire number for chamber zero. The fitted cubic was used to correct the nonlinearity.
Figure 3.43: Delay line readout diagram for the cathode position measurement.

Figure 3.44: Drawing of the printed circuit board tracks for the cathode wires.
Figure 3.45: Diagram of the generation of the prepulses. Figure 3.45a shows the position of a hit, Figure 3.45b shows the voltages generated on the cathode, and Figure 3.45c shows the pulses after the preamplifiers.

Figure 3.46: Distribution of first stops from the left side of the delay line readout for chamber zero.
Figure 3.47: Distribution of second stops from the left side of the delay line readout for chamber zero.

For any position of a hit, if it takes \( x \) nsec to travel to one end of the delay line, it takes \((1000-x)\) nsec to travel to the other end. The total time is therefore always 1000 nsec.

Figure 3.48: Diagram showing the time taken by each signal to travel through the delay line. For any event position, the sum of both times is 1 \( \mu \)sec.
Figure 3.49: Distribution of sum of first stops for chamber zero. Events with pre-pulses are between channels 300 and 450, whereas events with both good stops are in the main peak.

Figure 3.50: The four distributions represent the four sum combinations of the two stops from each side, for chamber zero. The correct stops are the ones whose sum is closest to channel 500.
Figure 3.51: Distribution of sum of correct stops for chamber zero. Note that events with early stops in Figure 3.50 have now been corrected.

Figure 3.52: Uncalibrated cathode position histogram for uniform X-ray illumination. The distribution reflects the acceptance of the trapezoidally shaped detector.
Figure 3.53: Linearity test of the cathode position. The measured position is plotted vs the stepping motor position.

Figure 3.54: The cathode position shows the individual wires with dips for the steel wires. The pad pattern is a weak modulation of groups of three wires.
Figure 3.55: Position of the channel number of the center of the cathode wires versus the wire number. The fitted line was used to provide the slope calibration.

Figure 3.56: Dependence of the chamber gain on the gas pressure. The measure of the pressure was the flow rate through a gas bubbler.
Figure 3.57: Gain variation as a function of the anode wire.

Figure 3.58: Histogram of the sum of the anode signals for chamber zero. Note the bump near the middle of the distribution due to one of the channels having an overflow. A corrected distribution was shown in Figure 3.37.
Figure 3.59: Diagram of the electronics for the scintillator detector. From [Hen91a].

Figure 3.60: Diagram of the electronics for the wire chamber detector. From [Hen91a].
The trigger logic. ML1 stands for a majority logic level of 1 (any one signal "TRUE"), and ML2 for a level of 2 (any two signals "TRUE"). Also shown is the single event scale-down.

Figure 3.61: Diagram of the trigger electronics. From [Hen91a]

Figure 3.62: Prescale timing diagram. Singles triggers are processed for a preset time (usually 30 μsec), then inhibited until a total of 100 μsec have elapsed. The process then repeats itself.
Figure 3.63: Data from a halo run. Figure 3.63a shows the singles energy, Figure 3.63b the singles scattering angle distribution, Figure 3.63c the singles azimuthal angle distribution and 63d the coincidence sum energy.

Figure 3.64: Singles energy spectrum of a halo run with some beam contamination present near 1650 KeV. From [Hen91a].
4. DATA ANALYSIS

Once a buffer has been shipped from the STARBURST to the µVAX, the data is stored in shared memory where it is used by the on-line analysis program and from where it is written to magnetic tape for more detailed off-line analysis. The on-line and off-line analysis programs are very similar, and most of the following discussion is valid for both programs, although obviously most of the analysis work is performed off-line.

Since a buffer of data can contain singles events, coincidence events, scaler readout events and germanium readout events, the first step in the analysis is to recognize what type of data the next event is. Germanium spectra are added to the running total germanium spectrum. Similarly, the scalers are added to the running total of the scalers, although the last scalers readout acquired is saved separately on-line, as a means of monitoring recent changes.

Singles and coincidence events are analyzed as follows. After a check of the input register to determine which detectors fired and whether they fired correctly, the event is reconstructed in two separate parts. The energy and timing information are obtained from the scintillator, whereas the position information is obtained from the wire chambers. Once the energy and the position of a particle are determined, full kinematic reconstruction is possible.

All events undergo some cuts on the raw signals, and once the events are reconstructed, additional kinematics can be required. After some corrections, the normalized coincidence cross section where a resonance would appear is then simply the sum of all coincidences that passed the cuts over the sum of all singles that passed the cuts, at each energy point.

A. INPUT REGISTER CHECK

The first word in a singles or coincidence event is from the input register. The input register has 12 inputs. Eight of the inputs are used to register which scintillators fired. The signals originate in the scintillator discriminators also used as triggers, as discussed earlier. Depending on which bit is turned on, one knows which scintillator, or scintillators for a coincidence, fired. The other four channels are used to monitor which chamber fired. As a signature of a chamber firing, the front cathode (the window) signal
is used. It is amplified and sent through a discriminator set above the noise. The discriminator output for each of the four chambers is then sent to the input register.

For a singles event the program checks that the scintillator, or adjacent scintillators, and chamber firing are from the same quadrant, and that no other chambers fired. Similarly, for a coincidence, the program checks that both chambers of the same opposite quadrants as the scintillators triggering the coincidence fired. It also checks that no other chambers or scintillators fired.

Once the input register is checked, events with nonstandard input register have been eliminated and hence the input register sort is the first cut on the data. The rest of the analysis can therefore assume standard singles or coincidence events.

B. THE EVENT RECONSTRUCTION

The main part of the analysis is to reconstruct the position and the energy of events from the detector signals recorded. For singles, the energy and position of one particle have to be reconstructed. For coincidences the energy and position of both particles are reconstructed. The process of extracting the position and energy of a hit from the detectors is closely related to how the detectors function. The event reconstruction below follows directly from the operation of the detectors explained in detail in the previous chapter.

The raw energy signal is the ADC signal coming from the scintillator photomultiplier tube. The pedestal of the appropriate channel is subtracted from the signal amplitude. This new amplitude is then corrected for the position dependence of the scintillator and a separate histogram is collected (see Figure 4.1). From the daily scintillator calibration, linear calibration coefficients are available to transform the corrected ADC distribution into an energy distribution. The energy distribution is then corrected for the energy dependence of the scintillator response and the final energy measurement is obtained. The energy of a single event is the energy measured in one scintillator, whereas the energy for a coincidence event is the sum of the energies of the two scintillators that fired.

The position of an event is measured by the wire chamber. The two sensing planes are perpendicular and produce a signal at the point where a particle crossed the plane of the detector. Since one plane uses resistive division and the other plane a delay line readout, the respective signals are processed completely differently.
The anode sensing plane uses resistive division to measure the position of an event. The charge from both ends of the resistive chain is detected in both a small and a large range ADC, as explained in the previous chapter. The first step of the position reconstruction is to determine which ADC value to use and to scale the values appropriately using a cross calibration between the two ADC's. The ratio of the signal from one side to the sum from both sides is a measure of the position of a hit, as shown in Figure 4.2. The distribution of ratios is then transformed into a position distribution using the S-shaped anode position calibration.

The cathode plane senses position with a delay line and produces timing signals that are recorded in TDC's. The first and second stops from each side of the delay line are recorded. By choosing the set of timing stops that most closely add up to 1 μs, the correct stops are identified. The difference between these stops is then a measure of position along the cathode plane (Figure 4.3). A cathode position is obtained from the time difference distribution using the linear cathode calibration coefficients.

C. **THE KINEMATIC RECONSTRUCTION**

When the energy and position of an event in the detectors has been derived, the full kinematics of the event can be reconstructed. For singles events, the Mott characteristics are reproduced, and for Bhabha events both tracks are combined to yield azimuthal back-to-backness and center of mass scattering angle back-to-backness.

1. **Singles**

   For a singles event, the position of a hit in the chamber allows the scattering angle and azimuthal angle of a Mott scattered position to be calculated, as shown in Figure 4.4, while the energy of the event is determined by the scintillator signal.

   Using the procedure described earlier, the energy of Mott events is calculated, and histogrammed for each scintillator (Figure 4.5). The eight scintillators are also histogrammed together (Figure 4.6).

   The main peak contains the positrons that Mott scattered into the detectors. The high energy shoulder on the peak is due to Mott scattered positrons where one of the positron annihilation 511 keV gamma rays Compton scattered in the plastic scintillator. The measured energy is then the sum of the positron kinetic energy and the Compton energy.
The low energy shoulder on the peak is caused by misidentified Bhabha events. If one of the particles in a Bhabha event is measured in the scintillators and the other, due to kinematics or multiple scattering, escapes the detectors, then the event will be registered as a singles event. If the low energy, large angle particle escapes, then the event will look like a singles event with a few hundred keV of energy missing, producing the low energy shoulder of the Mott peak. If the large energy, small angle particle escapes, the event will miss most of the energy and contribute to the low energy background. The slow rise towards smaller energies (around 500 keV) in Figure 4.6 illustrates this effect.

The tail from the main peak down to the discriminator has multiple origins. A small fraction of it is due to positron energy loss in the steel wires supporting the KAPTON vacuum windows. The steel wires are 0.8 mm in diameter, with a spacing of 12.7 mm. They therefore degrade the energy of at most 5% of the incoming particles. The second cause of low energy tails is back-scattering in the plastic scintillator. Extrapolating from Knoll, the back-scattering fraction at 2.25 MeV in plastic is ~2%. A third cause of low energy tails is the dead region between adjacent scintillators in a same quadrant. Positrons loosing energy in this region are recorded as low energy positrons.

The main component of the low energy tail is a secondary process. Low angle Mott scattered positrons miss the detectors and interact with the steel vacuum chamber pyramid as shown in Figure 4.7. Since steel has a much larger atomic number than plastic, the back-scattering fraction is enhanced to ~15% for 2 MeV positrons. Some of these back-scattered positrons enter the detectors and deposit their remaining energy in the scintillators, causing the low energy tail. The large amplitude of this effect is caused by the sharp angular dependence of Mott scattering.

To verify the above explanation, the positron beam was sent directly into the detectors and the pyramid. As shown in Figure 4.8, when the beam enters the detectors directly, there is little low energy tail, as expected, whereas when the beam is directed onto the steel pyramid, a spectrum similar to the low energy tail is recorded.

The azimuthal distribution for a typical run is shown in Figure 4.9. The four regions correspond to the four wire chambers on the target chamber pyramid. The lack of events at certain azimuthal angles is caused by the detectors not covering the steel edges of the pyramid. The physical processes in the experiment are azimuthally symmetric, the detectors are not. The sharp dip in the middle of each chamber is also a detector effect. Each chamber is backed by two scintillators. Where the two scintillators are joined, there is
a small dead region in the scintillator. Since the scintillator is the trigger for the experiment, there is a dead region in the wire chamber distributions, even though the chamber is uniformly efficient.

The scattering angle distribution is histogrammed in Figure 4.10. The main feature is the Mott distribution, with decreased counts at larger angles. It is not an exact Mott distribution because the low energy tail events discussed in section 5.1.1 do not follow a Mott distribution. The low end and high end drop off reflect the acceptance of the detectors. The discontinuity at 45° corresponds to the largest scattering angle measured in the middle of a chamber, with the higher angle events originating in the corners of the wire chambers, as shown in Figure 4.11.

2. Coincidences

For a coincidence event, the total energy is obtained by summing the energy of the scintillators that fired. The positional information leads to the azimuthal back-to-backness and the scattering angle distributions (see Figure 4.12). The scattering angle distribution is boosted to the center of mass reference frame to measure its back-to-backness. Other kinematic distributions are also monitored.

In a coincidence the initial beam energy is distributed among the two scattered particles, but the total energy is still conserved. Hence, the energy spectrum of individual particles in a coincidence is a broad distribution as shown in Figure 4.13. Histogramming the total energy, i.e. the sum energy of the two scintillators that triggered the coincidence, produces a distribution peaked at beam energy as expected (Figure 4.14). The energy sum graph shows the low energy tail due to energy loss in the steel support wires, back-scattering from the scintillator, and energy loss in the junction of the scintillators. The large contribution due to positrons back-scattered from the pyramid has been eliminated by the demand for a coincidence. As shown in section 3, the accidental rate is very low, making it unlikely for back-scattered events from the pyramid to be interpreted as coincidences. The high energy shoulder on the peak is caused by Compton scattering of the 511 keV gamma ray from the annihilation of the positron, as discussed for the singles energy spectrum.

Both the positron and the electron in a coincidence can be histogrammed as two singles events at some azimuthal angles (Figure 4.15). The partial azimuthal coverage is again clearly visible in the histogram, as it was for the singles azimuthal distribution. Since Bhabha scattering is a two body interaction, it must take place in a plane. Both the scattered particle trajectories and the point of interaction must be in a plane. The difference
in azimuthal angle between both trajectories must then be 180° as shown in Figure 4.12a. To monitor this, the azimuthal angle difference is calculated (see Figure 4.16). The coplanarity is clearly established, as is the low background.

The width of the gaussian distribution is caused in order of importance by multiple scattering in the target, a finite beam size (~1 mm) blurring the point origin of the interaction, and the detector resolution. To study the detector response with a minimum of multiple scattering, coincidences were collected from Bhabha scattering off a 100 μm beryllium foil (see Figure 4.17).

The coincidence scattering angles can also be monitored by the scattering angles of the positron and the electron, as for singles. The resulting rate (Figure 4.18) decreases with angle, as expected from Bhabha scattering. The rate then increases again, counter to Bhabha scattering predictions. The electron and positron are indistinguishable in the scintillator, since the positron annihilation radiation is not measured. For every positron Bhabha scattered at small angle, there must be an electron scattered at large angle, and since positrons scatter mainly at small angles, the electrons mainly scatter at large angles. The final distribution is the sum of the positron and electron distributions and hence is double peaked. In the center of mass, the electron and positron are truly indistinguishable, and the observed distribution is a symmetrized Bhabha distribution (Figure 4.19).

The laboratory angular distribution of Figure 4.18 is not symmetric because the Lorentz boost from the center of mass to the laboratory frame of reference produces an asymmetric change to the angle of the particle scattered along the direction of the boost from the change to the angle of the particle moving in the opposite dimension to the boost (Figure 4.12c). This can be illustrated in two ways. The sum of both scattering angles is a function of the individual scattering angles, causing an opening angle distribution much broader than the system resolution (see Figure 4.20).

The opening angle sum dependence on the individual scattering angles is shown by plotting the scattering angle difference versus the opening angle sum (Figure 4.21). The largest sum opening angle is at zero angular difference. This corresponds to equal scattering angles for the electron and positron, meaning equal angles with respect to the Lorentz boost, i.e. 90° scattering in the center of mass reference frame. The opening angle sum can also be boosted to the center of mass. The resulting distribution measures the back-to-backness of scattering angle in the center of mass, shown in Figure 4.22.
D. **Cuts**

Besides the input register selection discussed in section 1, many cuts are applied to the data. There are two levels of cuts. The raw cuts are needed to eliminate events where not all the chamber signals were collected or where the signals collected would lead to an erroneous position reconstruction. The raw cuts were already applied in sections 4 and 5 to obtain the position and angle spectra displayed. The second level cuts are on angle, energy, or other physical observables.

1. **Raw Cuts**

When a scintillator triggers the data acquisition system, chamber ADCs and TDCs are read out irrespective of what signals went into them. The input register cut ensures that the front cathode fired, as a criterion of the chamber firing, but there are many events where the front cathode fired and other signals are unusable. The following cuts, or checks are applied (see Figure 4.23):

- the anode signals from the resistive division as measured by the small range ADC must be between channels 40 and 2500. The low end cutoff is needed to eliminate events where the ADC did not register above the pedestal. The overflows are accepted because the analysis program automatically switches to the the larger range ADCs.
- the anode signals from the resistive division as measured by the large range ADC are limited to channels 0 to 1900. The events in or near the pedestal are handled by the low range ADC, so are not eliminated here. The overflow signals are not accepted because the overflow channel is not the real value of the pulse, just a lower limit. Using the overflow value would lead to an erroneous position reconstruction.
- the first stops from each side of the delay line must be less than the length of the delay line, i.e. 1 μs. The corresponding channel in the TDC timing the signals is channel 1000. Events where no TDC stop arrived are registered as overflows in a high channel. The limits on first stops are therefore channels 0 to 1000.

The above requirements must all be satisfied for a chamber to reconstruct the position of an event. If any of the cuts fails, the event is not processed further.

2. **Physical Cuts**

Once raw cuts have been applied, all the remaining events are reconstructed events. To ensure constant acceptance and efficiency, further cuts are applied to the data. They are
mainly intended to ensure against drifts in hardware discriminators or edge effects in the detectors.

The singles energy spectrum has a lower level cut set at 500 keV (see Figure 4.24). It keeps the lower level discriminator constant even if the gain in the scintillator changes. Such gain drifts result in different energy lower level discriminators (up to 50 keV), because the hardware discriminators cut off at a given voltage. Since the low energy tail stretches down to the cut off, it is important to avoid these drifts, even though they are 1.5 MeV away from the main peak. The coincidence spectrum has no low energy tail (see Figure 4.14), hence the integral is not as sensitive to lower level cuts. The low energy sensitivity in coincidences is due to the event kinematics. Low energy particles come from scattering at large angles. By cutting at a certain angle, events with energies lower than the one corresponding to the given angle are eliminated. Therefore no energy cut is applied to coincidence events. The angular cuts are discussed in the next section.

As discussed earlier, the experiment is statistically limited by the number of coincidences and not singles. The number of singles can therefore be reduced to an optimized subset still large enough to normalize to without a significant contribution to the statistical error. The singles events are limited to the central region of each scintillator. This optimizes the scintillator resolution, as well as minimizing the amount of energy lost to adjacent scintillators. By avoiding small scattering angles, the amount of low energy positrons in the low energy tail of the singles energy spectrum is also reduced. Any edge effects or change in acceptance in the wire chambers are also avoided.

The singles cuts applied are:
- for any quadrant the azimuthal angle is limited to 22.5° to 42.5°, and 47.5° to 67.5° (see Figure 4.25).
- the scattering angle must be between 22° and 43° as shown in Figure 4.26.

The effective area of a wire chamber used for singles events after the combined angular cuts is shown in Figure 4.27.

Coincidence events are restricted by similar cuts (see Figure 4.28). The cuts are maximized for acceptance while still avoiding the wire chamber and scintillator edges. The limits are:
- each coincidence azimuthal angle must be, for a quadrant, between 20.0° and 42.5° or between 47.5° and 70° (see Figure 4.29), to avoid edges, as for singles.
- The scattering angle must be between 62° and 118° in the center of mass (see Figure 4.30). The scattering angle cuts are fixed in the center of mass to keep a constant acceptance despite the change in Lorentz boost causing decreasing scattering angles for increasing energies. The cut is symmetric around 90° since a positron emitted at 62° will be accompanied by an electron at 118° in the center of mass, and both particles are indistinguishable in the detectors.

The scattering angle cut is also an energy cut because of the Bhabha scattering kinematics. The 118° cut correspond to a 600 keV cut in energy (for a beam energy of 2260 keV) on low energy scattered electrons or positrons.

E. ECorrections

In principle, the coincidence to singles ratio is found by integrating singles and coincidences passing all cuts and dividing both numbers. In practice, a number of corrections have to be made for misalignment, changing efficiencies, and changes in acceptance because of target motion. The misalignment corrections described below were already performed for the distributions shown in the previous sections.

1. Target position

The data analysis program assumes that the target position is at the very position specified by the design. Since for each particle there is only one point to reconstruct the track, the origin is assumed to be the point of intersection between the beam line and the target. If the distance between the target and the detectors is not correct, erroneous event reconstruction follows.

Figure 4.31 shows the center of mass scattering angle and scattering angle difference distributions with the target position uncorrected. The scattering angle difference is not centered at 180° which is unphysical. The center of mass scattering angle distribution is discontinuous at 90° because the analysis mis-identifies which particle scattered below 90° and which above. Using the back-to-backness of Bhabha scattering in the center of mass, the true target position was found. The target position was varied along the beam axis until the center of mass scattering angle difference was nearest 180° for all energies.

Once the position of the target along the beam axis has been fixed, the point of interaction is varied. It is possible for the beam not to enter the apparatus along the
theoretical beam axis because of the way the beam is centered. The centering is achieved
by equalizing the total scintillator count rate from each of the four quadrants. If in one
quadrant scintillators are slightly misaligned or less efficient in a corner at small angles, the
centering procedure would mis-align the beam. Figures 4.32 and 4.33 show the effect of
moving, in software, the point of interaction 1 cm horizontally and 1 cm vertically,
respectively. The difference between horizontal and vertical movement stems from the fact
that the azimuthal angle origin was chosen horizontally, breaking the azimuthal symmetry.
As seen in Figure 4.34, a horizontal movement also causes discontinuities in the center of
mass scattering angle and a shift in the azimuthal back-to-backness, whereas a vertical
movement causes a splitting in the azimuthal back-to-backness (see Figure 4.33). By
monitoring the azimuthal and scattering angle distributions, the angle difference centroids
and widths, the interaction point between the target and the beam was determined in three
dimensions.

In the extended run of the experiment described in the last chapter, the target
position was monitored with an alignment telescope and no target corrections were
required.

2. Chamber misalignment

When the target position has been optimized, there remains a discontinuity and
offset in the angular distributions (see Figure 4.34). The data analysis program assumes
that the detectors are located exactly where they were designed to be. The cathode read
outs are calibrated in absolute position by the steel wires used in their calibrations, but the
anode position is not absolutely calibrated, neither is the distance between the wire chamber
and the steel pyramid. It is therefore necessary to correct for any chamber misalignments.
Each chamber is allowed a small correction in both the wire chamber sensing directions and
along the beam axis, the correction to be set using the data. This is done by monitoring the
same distributions used to determine the target position.

Since there are more parameters than degrees of freedom, there is some obvious
indetermination, e.g. the four chambers and the target can be located anywhere along the
beam axis as long as the relative spacing remains constant. Therefore, the target position is
first optimized and fixed, then chamber positions are optimized. The procedure is iterated
to obtain the best relative positions between the four wire chambers and the beam
interaction point. The final position correction values are given in Table 4.1.
Table 4.1: Final corrections on the position of the target and wire chambers

<table>
<thead>
<tr>
<th>Correction applied to</th>
<th>Size of correction (in mm)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>in X</td>
</tr>
<tr>
<td>target position</td>
<td>-1.3</td>
</tr>
<tr>
<td>wire chamber number zero</td>
<td>0.15</td>
</tr>
<tr>
<td>wire chamber number one</td>
<td>0.0</td>
</tr>
<tr>
<td>wire chamber number two</td>
<td>-0.15</td>
</tr>
<tr>
<td>wire chamber number three</td>
<td>0.0</td>
</tr>
</tbody>
</table>

3 Chamber gain

The chamber gain influences the singles and coincidence efficiencies. To avoid any systematic errors due to chamber efficiencies, the gain was measured during each run by collecting pulses from the front cathode. The histogram obtained from each run was fitted with a Landau distribution [Lan44] defined by:

\[
\begin{cases}
  a_0 \ e^{-\frac{1}{2}(x+\lambda)^2} & \text{for } \lambda \leq 0 \\
  a_0 \ e^{-\frac{1}{2}(x+\alpha)^2} & \text{for } \lambda > 0 \\
\end{cases}
\]

with \( \lambda = \frac{x - a_1}{a_2} \) (4.1)

where \( x \) is the channel number, \( a_1 \) and \( a_2 \) are the fitting parameters. The second parameter, \( a_1 \), is the channel number of the maximum of the distribution (see Figure 4.35) and is used as a measure of the gain of a chamber.

A reference gain was chosen and for each run a relative gain was calculated for the four wire chambers. The ADC signals were then multiplied by the relative gain to obtain a constant gain throughout the experimental run. Care must be taken with the lower and upper end cutoffs. A signal whose gain was low and just failed the discriminator cannot be recovered by this procedure. Similarly, a signal whose gain was high and was recorded as an overflow cannot be reduced to a channel number.

Since the aim of gain equalization was to eliminate effects from overflow and discriminator cut off changes, more restrictive cuts have to be applied to the data. The lower level cut on signals is raised so that a pedestal multiplied by the maximum relative
gain still fails the cut. Some events with a small relative gain will be rejected even though they can be reconstructed, but a constant fraction efficiency will result.

The lower level cut was applied to the front cathode signal. It was placed high enough to be more restrictive than any hardware discriminator (i.e. the zero-cross amplifier discriminators on the delay line readout and the front cathode discriminator itself) and to reject all pedestals. The front cathode signals were limited to channels 450 and up, as shown in Figure 4.36.

Upper level cuts are most relevant for the anode readout where the pulse height is needed to reconstruct the event, whereas the front cathode and the delay line readout only need a pulse large enough to pass a discriminator. The upper level cut is therefore applied to the high range ADC signals, i.e. ADC4 and ADC5 for the four chambers. As opposed to the low end, the upper end cuts must be set low enough to prevent any overflow channel from scaling down into the allowed region when multiplied by the relative gain factor. This is the reason the cuts on ADCs 4 and 5 were set low in the previous section. The limits were originally set just under the overflow channel, but were then lowered to channel 1900 when the gain scaling was introduced.

4. Chamber Efficiency

The corrections applied to the wire chamber gain forced conservative cuts to be applied to the Landau distribution of charge deposited in the wire chamber. This led to a reduction in overall efficiency. The chamber efficiency can best be estimated by using coincidence events. If, for a coincidence trigger, a tight cut is placed on a centered event in one chamber (see Figure 4.37), the fraction of good events in the other chamber equals the chamber efficiency. The measured efficiency was 90.0%.

5. Effect on the Bhabha to Mott Ratio

To show the effect of the position and gain corrections on the final results, the data was analyzed without corrections, with only the positional corrections, with only the gain corrections, and with all corrections. Uncorrected results are shown in Figure 4.38. Large fluctuations are visible in the coincidences/singles ratio, and there is an offset in the center of mass scattering angle difference.

When the gain corrections are applied, the excitation function is smoother but the kinematic reconstruction is still not correct (see Figure 4.39). On the other hand, position
corrections without gain corrections produce more correct kinematics without seemingly improving the overall excitation function (Figure 4.40).

**F. EXPERIMENT MONITORING**

To ensure that fluctuations in the data were not due to beam tuning, detector performance, etc., various parameters were monitored during the experiment. These can be plotted versus energy to look for any trends. The influence of the known systematic effects that were corrected for can also be displayed.

**Angular Parameters**

The mean and standard deviation of the azimuthal angle difference were monitored and are plotted in Figures 4.41 and 4.42. The small range of the mean shows the constancy of the tune of the beam and the reliability of the positional reconstruction of the wire chambers. The slight offset from 180° is caused either by a small but consistent error in the position of the point of interaction, or by the not perfectly round shape of the beam spot. The range of the azimuthal difference is also small, again pointing to the accurate beam tuning and position reconstruction.

The mean of the laboratory opening angle sum is plotted as a function of the beam energy in Figure 4.43. The forward focusing of the Lorentz boost is clearly visible in the decrease of the opening angle with increasing energy, as expected. The standard deviation of the opening angle sum is constant and featureless (Figure 4.44).

The mean and standard deviation of the scattering angle center of mass difference are shown in Figures 4.45 and 4.46. As the energy increases, the mean slowly shifts to lower values. Since this is a center of mass plot, the Lorentz boost has already been taken into account and the center should not vary with energy. However, since a boost has to be performed to the center of mass frame, an incorrect position reconstruction could lead to such an effect. This was investigated using the Monte Carlo simulation.

By varying various offsets, it was found that these plots were most sensitive to a drift or change along the anode direction. An effective change in position of 0.56mm was needed to produce the change in mean between large and small energies. Looking at the data, one notices that there must have been a drift, since the repeated points, between 2200 and 2270 keV, are distinctly lower than the original scan. There seems to be an additional trend towards lower means with increasing energy.
A drift in reconstructed anode position could be created by a drift in the gain of one of the delay line amplifiers at each end of the resistive division. A drift in chamber gain would not cause such a change in position, since only relative amplitudes are used. To check for such a drift, the anode distributions at the beginning and end of the scan were overlapped (see Figure 4.47). A small but significant shift between the individual wire peaks is clearly noticeable. Using the anode calibration, the position drift of 4 channels was calculated to be 0.19 mm, matching the shift between the first scan and the repeated points. This does not explain the overall decrease of the mean though.

There is still a need to find a cause for an anode offset slowly changing with energy. At different energies, the coincidences with minimal total opening angle are detected by different parts of the detector. Since the total opening angle is smaller at higher energies, the anode position will be lower. The problem can therefore be rephrased as one of changing position offset with position, i.e. a non-linearity in the anode read-out.

The nonlinearity of the anode is well known [Fra81] and is shown in Figure 4.48. The difference between reconstructed anode position and real position is a characteristic S-shaped curve (see Figure 4.49). The difference in the nonlinearity of the anode position corresponding to the mean opening angle at 2150 and 2350 keV can be calculated. It came to 0.38 mm, accounting for the additional drift in the center of mass scattering angle difference. These drifts, although measurable, were considered steady and small, therefore not requiring correction.

2. Detector Performance

After gain corrections are made, the fraction of events failing the raw signal cuts should be constant. These fractions were monitored for each ADC and the performance of the four chambers, as measured by one of the ADC’s, is shown in Figure 4.50. Although a few points seem to have nonstatistical deviations, the overall fractions seemed constant.
Figure 4.1: Corrected ADC channel number for all events in scintillator number zero.

Figure 4.2: Reconstructed anode distribution for all events in chamber zero.
Figure 4.3: Reconstructed cathode distribution for all events in chamber zero.

Figure 4.4: Kinematics of a Mott scattered positron. Figure 4.4a is a view looking downstream and Figure 4.4b is a side view.
Figure 4.5: Energy spectrum of single events in scintillator number zero. The regularly spaced dips are due to the channel round off of integers multiplied by a real number.

Figure 4.6: Energy spectrum of single events from all scintillators.
Figure 4.7: Diagram illustrating a possible path for positrons back-scattered from the steel pyramid.

Figure 4.8: Energy spectrum for direct beam in a scintillator. Note that the low energy tail is highly suppressed.
Figure 4.9: Azimuthal distribution of single events. The four groups of angles correspond to the four chambers, with the dip in the middle of each one due to the joint between scintillators.

Figure 4.10: Scattering angle distribution of single events. The fall-off at 17° and 45° is due to the acceptance of the chamber.
Figure 4.11: Angular acceptance of wire chambers.

Figure 4.12: Kinematics for Bhabha scattered positrons. Figure 4.12a is a view looking downstream, Figure 4.12b a side view, and Figure 4.12c a center of mass side view.
Figure 4.13: Energy spectrum from individual particles in a coincidence.

Figure 4.14: Total energy in a coincidence. The energy is the sum of the energy of the two particles in a coincidence.
Figure 4.15: Azimuthal scattering angle of individual particles in a coincidence.

Figure 4.16: Azimuthal angular difference between the two particles in a coincidence.
Figure 4.17: Azimuthal angular difference for positrons Bhabha scattered from 100 µm Beryllium foil.

Figure 4.18: Scattering angle for individual particles in a coincidence.
Figure 4.19: Center of mass scattering angle for individual particles in a coincidence.

Figure 4.20: Total opening angle between the two particles in a coincidence. The total opening angle is the sum of the individual scattering angles.
Figure 4.21: Two dimensional histogram of the scattering angle difference versus opening angle sum for a coincidence.

Figure 4.22: Difference of the center of mass scattering angles.
Figure 4.23: Raw cuts applied to the data. The dashed lines show the used cuts. Figures 4.23a, 4.23b, 4.23c and 4.23d show the cuts on ADC1, ADC2, ADC4 and ADC5 respectively. Figure 4.23e and 4.23f show the cuts on TDC1a and 2a respectively. Chamber zero signals were chosen as an example.
Figure 4.24: Singles energy spectrum with dashed line showing applied cut.

Figure 4.25: Azimuthal angle distribution for singles events with dashed line showing applied cut.
Figure 4.26: Scattering angle distribution for singles with dashed line showing applied cut.

Figure 4.27: Two dimensional distribution of selected singles hits in wire chamber zero.
Figure 4.28: Two dimensional distribution of selected coincidence hits in wire chamber zero.

Figure 4.29: Azimuthal angle distribution for coincidence events with dashed line showing applied cut.
Figure 4.30: Scattering angle distribution of coincidence events in the center of mass with dashed line showing applied cut.

Figure 4.31: Center of mass scattering angle (Figure 4.31a) and scattering angle difference (Figure 4.31b) with the target position uncorrected. Note the discontinuity at 90° in the scattering angle and the offset from 180° in the difference.
Figure 4.32: Azimuthal angular distributions with the origin offset by 1 cm horizontally in software. Figure 4.32a shows the azimuthal distribution for individual events in a coincidence, Figure 4.32b shows the azimuthal difference.

Figure 4.33: Azimuthal angular distributions with the origin offset by 1 cm vertically in software. Figure 4.33a shows the azimuthal distribution for individual events in a coincidence, Figure 4.33b shows the azimuthal difference.
Figure 4.34: Center of mass scattering angle distributions with the origin offset by 1 cm horizontally in software. Figure 4.34a shows the scattering angle distribution for individual events in a coincidence, Figure 4.34b shows the scattering angle difference.

Figure 4.35: Pulse height distribution of the front cathode signals with fitted Landau. Parameter a1 is used to measure the gain of the wire chamber.
Figure 4.36: Front cathode pulse height for all events with dashed line showing applied cut.

Figure 4.37: Chamber efficiency cuts on coincidences. Figure 4.37a shows the cut applied to coincidences in one chamber. Figure 4.37b shows the matching events in the opposite chamber. The ratio of the total number of counts if the chamber efficiency.
Figure 4.38: uncorrected data. Figure 4.38a shows the coincidences to singles ratio, Figure 4.38b the mean of the scattering angle center of mass difference for coincidences versus beam energy.
Figure 4.39: Data corrected for chamber gain. Figure 4.39a shows the coincidences to singles ratio, Figure 4.39b the mean of the scattering angle center of mass difference for coincidences versus beam energy.
Figure 4.40: Data corrected for position. Figure 4.40a shows the coincidences to singles ratio, Figure 4.40b the mean of the scattering angle center of mass difference for coincidences versus beam energy.
Figure 4.41: Mean of the azimuthal angle difference for coincidences versus beam energy.

Figure 4.42: Standard deviation of the mean of the azimuthal angle difference for coincidences versus beam energy.
Figure 4.43: Mean of the opening angle sum for coincidences versus beam energy.

Figure 4.44: Standard Deviation of the mean of the opening angle sum for coincidences versus beam energy.
Figure 4.45: Mean of the scattering angle center of mass difference for coincidences versus beam energy.

Figure 4.46: Standard deviation of the mean of the scattering angle center of mass difference for coincidences versus beam energy.
Figure 4.47: Anode distribution from lowest and highest energy points. A small shift of the anode wires is visible between the two runs.

Figure 4.48: Nonlinearity of the anode. The reconstructed position is plotted versus the real position in units normalized between -1 and 1.
Figure 4.49: Difference between the real and reconstructed position of the anode, displaying the S-shape of the nonlinearity.
Figure 4.50: Fraction of events failing the cuts on ADC5 as a function of energy, for each chamber.
5. MONTE CARLO SIMULATION

A Monte Carlo simulation of the experiment was performed. This chapter outlines the experimental parameters entering in the simulation and the physical phenomena included. Although some of the parameters were extracted from the data, the simulation shows how the experimental data can be reproduced using known cross sections and distributions.

A. OVERVIEW

The aim of the Monte Carlo simulation is to simulate the experiment in as complete a fashion as possible while still being able to run the program for equivalent statistics to the data and over the energy range scanned by the experiment. This is not possible with the complete simulation approach of GEANT or EGS4 because of the prohibitive computer time needed for each event.

Instead a program was written that included the relevant physical phenomena but only considered events scattered in or near the detectors. This eliminates the main part of the beam that is only slightly scattered in the target but does not interact with the detectors. The reduction in events handled allows the simulation to be run with realistic statistics and over the desired range of energies. However, the effect of back-scattered beam from the steel pyramid is not simulated and must be artificially introduced.

The simulations tried to reproduce the experiment as closely as possible. This meant including beam properties, target properties, wire chamber properties and scintillator properties. The simulation is run by sending a number of positrons through the simulated experiment and then reconstructing the event as the data acquisition and analysis system would. The "measured" distributions are histogrammed to be compared with the experimental data. The same cuts as in the data analysis are used and the coincidence to single ratio is then calculated. The excitation function can then be compared to the real data.

The Monte Carlo simulation has the unique advantage of being able to artificially turn on or off any chosen physical phenomenon and hence allowing for a study of the effects of the chosen phenomenon. Such studies were executed and are described below.
B. **THE BEAM**

The simulation aims to reproduce the experimental rates by sending the real number of positrons through the apparatus. It therefore needs as an input the beam intensity, target material and target thickness. It uses the given target thickness to calculate the probability of a positron interacting in the experiment. Since a large fraction of the beam crosses the target chamber without interacting with the detectors, one only needs consider the fraction of positrons that do interact with the detectors. This probability is then scaled up and the integrated beam scaled down by a same factor to improve the efficiency of the simulation.

The beam energy is set to whatever accelerator voltage is to be simulated. At any given voltage, there is a certain amount of radio frequency ripple on top of the accelerating voltage. This causes a broadening of the mono-energetic beam into a ~ 5 keV wide beam for this experiment. The beam broadening was simulated by a uniform distribution around the chosen energy. The broadening was typically set at the measured 5 keV.

The positron source size and the accelerator beam optics determine the size and position of the beam-target interaction region. Since the beam is neither point like nor perfectly centered, these effects must be included. The beam can be given any size and the simulation will create a uniform circular distribution of the given radius. The beam radius was typically set to 0.75 mm, a measured beam spot size. The beam spot position is designed to be on the axis of symmetry of the detection system. However, from reconstruction of the experimental data, it is known that there are some slight offsets. A two dimensional offset in the plane of the target was therefore included. The offsets were typically 1.5 mm in X and 1.5 mm in Y where X and Y are an orthogonal axis system centered on the detector axis of symmetry.

C. **THE TARGET**

The simulation accepts lithium or beryllium targets of any thickness and calculates the cross sections and probabilities accordingly. When a target material is chosen, the program looks up the atomic charge and the atomic weight. The atomic charge is used to calculate the Mott scattering cross section and the number of target electrons for Bhabha or resonant scattering. The choice of target material also affects the momentum distribution of the target electrons and hence the energy broadening of resonant positron-electron scattering.
The electron momentum distribution in a material is experimentally determined by measuring the Compton profile of a sample. The electron momentum distribution along one axis is measured by Compton scattering X-rays in a sample of material. The simulation uses tables with the experimental momentum distributions for lithium and beryllium to assign a momentum amplitude to target electrons. The electron momentum is assumed to be isotropic. The lithium electron momentum distribution is shown in Figure 5.1.

The dominating factor in the shape of a resonance in positron-electron scattering is the broadening due to the motion of target electrons. Since the electron momentum distributions are different in lithium and beryllium, the resonance shapes are different. The simulation uses the calculated resonance shapes from tables. The lithium and beryllium resonance shapes are shown in Figure 2.39.

The target thickness is needed to determine the absolute event probabilities, or event rates in the experiment. It is entered in mg/cm² and is combined with the atomic weight and the number of electrons per nucleus to obtain the number of scatterers. The target position can be offset along the beam axis relative to the detection system. The target was typically located 3mm upstream from its design location as observed in the experiment.

D. THE WIRE CHAMBERS AND POSITION

Once an event has been generated at the target, the program checks whether the particle or particles registered in the detectors. Detector size, offsets resolution and efficiency are important in realistically reproducing the experiment and are therefore included.

1. Geometric Effects

The geometry of the wire chambers is identical to the experimental set up. Since the geometry is crucial in determining acceptances, the wire chamber frame position was taken over from the blue prints. Besides the position and size of the chambers, the offsets and dead region of the chambers were also included in the simulation.

By correcting the experimental data to produce back-to-back Bhabha events in azimuthal and center of mass scattering angles, the relative offsets of each wire chamber were obtained. These offsets, along with the beam and target offsets were included in the simulation using the values from Table 4.1.
The wire chambers measure the position of charged particles by detecting the charge caused by an avalanche of ionization near the anode wires. The avalanche is produced when electrons, freed in collisions of the detected particle with chamber gas, accelerate in the increasing radial field near the anode wires. Near the edge of the detectors, the electric field are altered by the presence of the detector frame. This creates a small dead region all around the edge of the detector. The region was measured in detector X-ray tests to be ~5mm. Such a dead region was included in the simulation.

Since the beam line and the accelerator are at a vacuum of ~2 x 10^{-6} torr and the detectors are at atmosphere, a vacuum window transparent to the electrons and positrons was needed. This was obtained using 50 μm thick kapton windows. To prevent excessive bowing of the kapton however, steel reinforcement wires were installed on the vacuum side of the windows. The circular cross section wires were spaced every 12.7 mm and were 0.8 mm in diameter. The steel wires were thick enough to stop a large fraction of incoming particles and hence produced partially dead regions across the detectors (see Figure 3.54).

The simulation incorporated the wires by assuming square wires of equal cross sectional area, i.e. 0.7 mm a side. When a particle hits the wire, it assumed to stop in it and not trigger any detector. The resolution of the detectors is taken into account after this process, producing a smearing of these masked regions, as seen Figure 5.2.

2. Resolution

The wire chambers combine a resistive division with a delay line read-out technique to provide the two dimensional position of a particle. Each method has a distinct resolution caused by the technique.

When a particle crosses the middle of the wire chamber, all the ionization is produced along a line parallel to the drift field lines in the chamber, causing the charge to be projected on and collected at a single point on the anode. On the other hand, a charged particle traversing the detector diagonally will deposit charge along a broad anode region.

The perceived resolution near the edge of the chamber is therefore noticeably worse than the resolution measured with X-ray sources or a perpendicular electron source. This is illustrated in Figure 4.2. The anode resolution, as measured with an X-ray source, is limited by the anode wire spacing because the resistive division can only reconstruct which wire the charge originates from. The position is therefore quantized in ~1.9mm steps.
However, a hit at a wire position will be correctly reconstructed, whereas a hit between two wires will be offset by 0.95 mm. The mean offset from the correct position is therefore close to 1 mm, and this was considered the resolution in the center of the chamber.

The resolution worsens near the edges. As can be seen from Figure 4.2, the wires become indistinguishable near the top and bottom of the wire chamber. The resolution near the edge was hard to estimate, but for the simulation anode resolution an overall average value of 2.0 mm was typically chosen.

As for the anode, the cathode resolution worsens from the center of the chambers to the edges and an overall average value was determined (see Figures 3.54 and Figure 4.3). The resolution of the cathode was ~ 450 μm for X-rays. The β particle resolution in the center of the chamber is approximately equal to the X-ray resolution since since individual wires, with a 1.27 mm spacing, are clearly visible. A typical cathode resolution in the simulation was 1.0 mm.

3. Efficiency

The simulation assumed that the chamber efficiency was identical for singles and coincidences, and that it was uniform across the chamber except for the dead region around the edge of the wire chambers. The efficiency was typically set at 84.5%. This value was chosen to reproduce the experimental values of the Bhabha to Mott ratio. The difference between the measured efficiency and the value used in the simulation can be caused by the approximations made in the steel wire and the dead region approximation. The dead region is not 100% dead in the real detector, and the exact value of the dead region varied around different edges of the detectors and from detector to detector.

To show the effect of a chamber efficiency, the simulation was also run with an efficiency of 82.5%, and the results are compared in Figure 5.3.

E. THE SCINTILLATORS AND ENERGY

After crossing the wire chambers, the scattered positrons and electrons enter and deposit their energy in the plastic scintillators. For elastically scattered positrons, as in Mott events, the scintillator pulse height distribution is a typical Gaussian with the characteristic resolution of the scintillator. However, since the scattered particles are positrons, they also emit annihilation radiation. These 511 keV gamma rays have a small probability of Compton scattering in the scintillator, depositing some additional energy.
There is also a small probability for the positrons to back-scatter out of the scintillator or to lose some energy by hitting the steel wires supporting the vacuum window.

1. Resolution

To study the performance of the scintillators, some direct beam was steered into a scintillator. The resulting energy distribution is shown in Figure 5.4. By fitting a gaussian to the main peak, the resolution of the scintillator was calculated. This resolution can then be scaled using known scintillator behavior. The resolution of a plastic scintillator scales as $1/\sqrt{E}$. This behavior was checked in our detectors and indeed provided a good description of the energy behavior of the scintillator. The simulation used the fitted gaussian ($\sigma = 64$ keV at 2200 keV) as a reference resolution and then scaled it using $1/\sqrt{E}$.

The described method produces a good imitation of the scintillator performance in the center of the scintillator. The energy spectra obtained closely resemble experimental spectra when the geometric cuts are applied. However, when no geometric cuts are applied, the experimental peaks broaden due to the poor resolution of the edges of the scintillator. This was not simulated since most of the analysis is performed with the geometric cuts in effect and since no energy dependant cuts are applied besides the lower level cut off that is already appropriately simulated.

2. Compton Scattered Annihilation Radiation

As can be seen from the shoulder above beam energy in Figure 5.4, the fraction of positrons whose annihilation radiation Compton scatters in the scintillator is significant. The probability of Compton scattering was obtained by fitting both the main peak and the Compton peak and calculating the ratio of the Compton integral to the integral of the main peak (ratio = 0.35). This probability was used in the simulation whenever a positron was detected in the scintillator. Events where a Compton scattering event was generated were given an additional energy of 100 keV plus a gaussian with a standard deviation of 170 keV, the distribution that best fit the experimental shoulder above the main energy peak.

The probability of detecting annihilation radiation was zero for electrons.
3. Low Energy Tail

The large low energy tail in the experimental singles spectrum is due to back-scattering from the steel pyramid. However, back-scattering from the scintillator itself or energy loss in the steel widow support wires are also possible. These can not be distinguished from the pyramid back-scattered positrons, but can be detected if the beam is sent directly into the detector as for Figure 5.4.

The low energy tail of Figure 5.4 was fitted and was installed in the simulation. The same low energy tail was assumed for electrons and positrons.

4. Lower Level Discriminators

There were two lower level discriminators applied to the scintillator energy signals in the experiment: one hardware discriminator and one software discriminator. For coincidences, only the hardware discriminator was in effect at 350 keV. This was simulated in the experiment by a 350 keV lower cut-off on both electron and positron energy signals.

For singles triggers, an additional software cut-off was placed at 500 keV, to avoid any systematic errors caused by drifts in the hardware lower level discriminators. This was not necessary for coincidences because unlike singles, very few events were near the lower level. The software discriminator was taken over in the simulation and singles events were cut off below 500 keV.

F. INTERACTIONS

The main part of the programme involves simulating the interactions between the positrons in the beam and the target. Positrons can just multiple-scatter in the target and continue on to the beam dump; they can interact with the nucleus to Mott scatter into the steel pyramid or the detectors; they can interact with the target electrons to Bhabha scatter or possibly exhibit resonant scattering behavior.

As discussed previously, not all positrons in the beam are taken through the experimental simulation. This corresponds to dropping the first interaction from the list above. The other interactions were included and are discussed in greater detail below.
1. Mott Scattering

The predominant type of events detected are Mott scattered positrons. These are positrons that scatter off the lithium nuclei in the target. Most of them are scattered directly into the detectors from the target, while some are back-scattered from the steel pyramid holding the detectors. The simulation must reproduce the correct cross-sections and angular distributions for these events.

The cross section for Mott scattering used in the simulation was as written in equations (2.28) and (2.29) and displayed in Figures 2.23 and 2.24 in Chapter 2. The simulation integrated the Mott cross section between 15° and 55°. This range covers the detectors and ignores the largely unscattered beam that hits the beam dump detectors. The Mott scattering cross section, like the Rutherford scattering cross section, diverges for small angles and hence does not adequately describe such scattering anyway.

The probability of a positron Mott scattering is given by

\[ \text{prob}_{\text{Mott}} = \int d\sigma \cdot N_a \cdot \frac{\text{thick}}{\text{mass}} \]

(5.1)

where \( \int d\sigma \) is the integrated cross section, \( N_a \) is Avogadro's number, \( \text{thick} \) is the thickness in mg/cm\(^2\) and \( \text{mass} \) is the molar mass in g/mole. The simulation calculates this probability and generates Mott events accordingly. To improve the efficiency of the program however, a factor multiplies all probabilities to minimize the amount of generated positrons that do not interact. This factor is then corrected for by correspondingly dividing the integrated beam by it.

The angular distribution of the Mott scattered positrons is also given by the differential cross section above and was reproduced in the simulation. The azimuthal distribution is uniform and the detected azimuthal distribution is determined by the detector acceptance as shown in Figure 5.5. The distributions are altered when multiple scattering is included in the simulation.

The Mott scattering angular distributions will also contain events caused by coincidences detected as singles. This is most noticeable at large scattering angles, as will be seen later.
2. Back-scattered Positrons

Since Mott scattering varies as $1 / \sin^4 \theta/2$ a large number of positrons are scattered onto the steel pyramid supporting the detectors, as can be seen in Figure 5.6. Most of these positrons annihilate in the metal and are not detected by the experiment. However, electrons and positrons have a significant probability of back-scattering out of a material, and the probability increases with the atomic charge of the material, as shown in Figure 5.7 ([Tab71], [Kno79]).

The back-scattering probability $\eta$ is given [Tab71] by:

$$\eta = \frac{1}{4} a_1 \left(1 - \tan \left[ \frac{1}{2} (\ln a_2 + a_3 \ln \tau) \right] \right),$$

(5.2)

where for iron $a_1 = 0.254$, $a_2 = 0.107$, and $a_3 = 1.39$; $\tau$ is given by $T / mc^2$ with $T$ the beam energy and $mc^2$ the rest mass of the electron in energy units. The back-scattering probability as a function of energy for iron, a good approximation to the steel used in the experimental setup, is shown in Figure 5.8. The back-scattering fraction for 2.2 MeV positrons, assuming it is identical to electrons, is therefore $\sim 14\%$. To check that back-scattering could produce an effect of the amplitude seen in the experiment, a rough calculation was performed.

Mott scattering was integrated between 6.1 and 15.9 degrees, corresponding to the scattering angles of the inner diameter of the beam pipe at the top of the pyramid and the scattering angle at the steel edge of the top of a wire chamber. This was considered to be the cross section (24.5 barns) for hitting the pyramid. The reflection coefficient was taken to be 14%, and the acceptance of the detectors for back-scattered positrons was estimated from a point at the middle of the top of a wire chamber. The back-scattered positrons were assumed to be isotropically distributed. This produced an acceptance estimate of 22%, leading to a back-scattering detection probability of 3.1%. The detected back-scattering cross section reduces to 0.76 barns, as compared to a Mott scattering cross section of 2.35 barns at the same energy.

This rough calculation shows that back-scattering can account for the significant tail in the singles energy spectrum. More detailed calculations would require knowledge of the angular distributions and energy spectra of back-scattered positrons, neither of which was available. The simulation was therefore based on our experimental data.
The experimental effect of the back-scattered positrons is to produce a low energy tail on the singles (Mott scattered positrons) energy distribution. The tail is caused by detecting, using the plastic scintillators, back-scattered positrons that lost energy in the scattering process. The tail is therefore mainly an addition to the Mott scattering rate in the detectors and not an energy degradation of the positrons Mott scattered directly into the scintillators. The addition must be included in the simulation. This was done by modeling the spectrum with a functional form as shown in Figure 5.9.

The fit was performed with a ten parameter function consisting of 2 gaussians, one for the main peak and one for the Compton shoulder, and a linear low energy tail that falls off at high energy as a half gaussian at a variable point. The parameters of the tail were obtained over several energies and a straight line was fitted to them, as shown in Figure 5.10. The straight line fit parameters provided the energy dependance of the tail used in the simulation.

3. Bhabha Scattering

Positrons that interact with electrons in the target can Bhabha scatter or possibly exhibit resonant scattering behavior. This section deals with Bhabha scattering, the next one with resonant scattering.

The differential cross section for Bhabha scattering in the center of mass is given in equation 2.35. The cross section was integrated and simulated for angles between 40° and 140° in the center of mass and over all azimuthal angles. This angular range completely covers the detectors. As for Mott scattering, a probability was calculated from the cross section and then scaled for the simulation. The scattering angle distribution for Bhabha scattered positrons is used in conjunction with the electron momentum distribution to determine the final distribution in the center of mass. Multiple scattering is also added and will be discussed later.

The positron and electron scattering angle angular distribution in the center of mass is shown in Figure 5.11. The distribution is symmetric around 90°, since the positron electron pair are emitted back-to-back on the center of mass. The back-to-backness can be checked by subtracting the positron scattering angle by the electron scattering angle (see Figure 5.12). The angles are boosted to the laboratory frame of reference to test the detector geometry and to apply cuts. The laboratory scattering angle distributions for the positrons and electrons are shown in Figure 5.13 and 5.14. Since in the simulation the positron can be distinguished from the electrons, separate distributions can be monitored.
The experimental situation corresponds to seeing the sum of both distributions, as shown in Figure 5.15.

The sum opening angle in the laboratory frame of reference is shown in Figure 5.16. The sharp peak and the low cut-off are created by positron-electron pairs scattering at equal angles, generating the smallest possible opening angle. The wide distribution is caused by the range of laboratory sum angles kinematically allowed. It is a kinematic effect unrelated to any broadening due to resolution, multiple scattering, etc...

The azimuthal angular distributions for the positron and electron are uniform and hence any features are introduced by the detector acceptance (see Figure 5.17). The azimuthal back-to-backness is monitored and shown in Figure 5.18.

4. Resonant Scattering

To simulate resonant scattering, the shape of a resonance in lithium or beryllium must be established, the amplitude must be calculated, and the correct kinematics must be worked out. These depend on the lifetime and type of particle is assumed to cause the resonant behavior. For most work a pseudo-scalar was assumed.

As discussed previously, the electron motion in the target broadens the intrinsically narrow width of a point-like particle resonance. The resonance shapes for lithium and beryllium, derived by using the experimental electron momentum distributions, are stored in tables and used in the simulation.

The resonance amplitude depends on the type of particle and lifetime assumed. If the resonant shape is assumed to be Lorentzian, then the energy integrated cross section is

\[ \int \sigma_x \, dE = \frac{\pi \, \sigma^0_x \, \hbar}{2 \, \tau}, \]  

(5.3)

where \( \sigma^0_x \) is the intrinsic resonance height and \( \tau \) is the lifetime. For a 1.8 MeV center of mass resonance, this can be reduced to

\[ \int \sigma_x \, dE = \frac{2.28 \times 10^{-12} \, \text{barn eV sec}}{\tau}. \]  

(5.4)

The amplitude is therefore inversely proportional to the lifetime of a particle. A reference lifetime of \( 10^{-13} \) seconds, a rough approximation to the \( g-2 \) limits, was typically chosen.
The resonant scattering amplitude also depends on the type of particle assumed. The simulation accommodates scalar or pseudo-scalar particles, who have an identical resonant cross section for a given lifetime. The vector and axial vectors couplings were added, and for these cases the cross sections have to be recalculated. The kinematics are also affected, as discussed below.

Once a positron has resonantly scattered off an electron, the kinematics are identical to Bhabha scattering. There is no way of distinguishing between Bhabha scattering and resonant scattering. All distributions and histograms are therefore shared with Bhabha scattering events. If one wants to study either Bhabha scattering or the resonance individually, either can be turned on or off.

The difference in kinematics is introduced in the scattering process itself. Whereas positrons are Bhabha scattered according to the differential cross section, resonantly scattered positrons have an angular distribution that depends on the spin of the resonant state. For \( J=0 \), as in the case of the scalar or pseudo-scalar particles, the positron distribution is isotropic in the center of mass reference frame. For \( J=1 \), an angular distribution has to be introduced in the center of mass. The vector and axial vector angular distributions are shown in Figure 2.33.

5. Multiple Scattering

Positrons or electrons undergoing any of the above interactions are simultaneously interacting with the bulk of electrons in the target to multiple scatter. Multiple scattering causes a broadening of all angular distributions. It increases as the square root of the thickness and the square of the atomic charge. Unfortunately, the thicknesses typical to our experiment are outside the domain of validity of multiple scattering or Moliere Scattering. The mean multiple scattering angle can therefore not be established from theoretical principles.

To simulate multiple scattering, all scattered particles undergo a gaussian distributed additional scattering. The standard deviation of the gaussian is provided to the programme. The used width of the gaussian was derived from the data by measuring the broadening of the azimuthal back-to-backness and the opening angle sum in the laboratory and center of mass reference frames. For a 1.5 mg target of lithium, a 1.4° multiple scattering angle was measured.
The main effect of multiple scattering is to broaden all angular distributions: the center of mass and azimuthal back-to-backness widths are mostly due to multiple scattering (see Figures 5.19 and 5.20). The laboratory opening angle sum is broadened enough to lose the sharp low angle peak, although the asymmetric feature of the distribution is still noticeable (see Figure 5.21).

Another, less intuitive effect of multiple scattering is on the center of mass scattering angle distribution. Since multiple scattering is a laboratory phenomenon, the effects on small and large scatterings will be different in the center of mass reference frame (see Figure 5.22). Multiple scattering on large angles has a smaller effect in the center of mass, causing less broadening of the large angle peak, as seen in Figure 5.22.

6. Coincidence as single

When a coincidence event, either a Bhabha scattering event or a resonant scattering event, produces a particle scattered near the edge of a detector, there is the possibility that the accompanying particle misses the opposite detector. In the experiment, the event would be recorded as a singles event. This feature was included in the simulation.

Coincidence events detected as singles are most easily recognized in the singles energy spectrum when no geometric cuts are applied. Since these events are caused near the edge of a chamber, they usually consist of a very small angle or a very large angle scattering, with an energy close to the beam energy or a small energy, respectively. Figure 5.23 displays both features: a rise at low energies from large angle scatterings, and a shoulder on the left side of the Mott peak from small angle scatterings.

G. Cuts

There are two types of restrictions on the data in the experiment. There are hardware discriminators used to decide whether a signal is larger than noise implying that a detector fired, and there are software cuts applied to the data.

Since this is a simulation, no hardware cuts can nor need be applied. But they must be simulated to reproduce the experiment. The scintillator hardware discriminators are simulated by requiring that a particle must have more than 400 keV of energy to register as a hit, as described earlier. In addition, singles events are required to have more than 500 keV to be included in the analysis, as was required with the experimental data. The
chamber discriminators are simulated by checking the detector for its efficiency, which is caused by the discriminators on all the chamber signals.

The geometric cuts applied are identical to the ones applied to the experimental data. They are, for singles:
- the scattering angle must be between 22° and 43°.
- the azimuthal angle is limited to 22.5° to 42.5°, and 47.5° to 67.5° for any quadrant.

For coincidences:
- the center of mass scattering angle must be between 62° and 118°.
- the azimuthal angle must be between 20° and 42.5° or 47.5° and 70° for any quadrant.

H. COMPARISON OF SIMULATED AND EXPERIMENTAL HISTOGRAMS

The Monte Carlo simulation results can now be compared to the experimental results. Since all relevant physical phenomena and variables were included, the simulation results should closely reproduce the experimental data. In this section, the distributions at one particular energy will be examined. The Bhabha to Mott ratios will be compared in the next section. The histograms were obtained by simulating a 50 minute run of a 500000 positrons per second beam at 2350 keV on a 1.5 mg/cm² lithium target.

1. Singles Histograms

The singles distributions are a combination of the Mott scattered positrons and the Bhabha scattered particles from coincidences detected as singles.

The simulated singles energy distribution is shown in Figure 5.24a and the experimental energy distribution is shown in Figure 5.24b. The main Mott peak is reproduced with the correct resolution, as was the right hand shoulder produced by the capture of the Compton scattered annihilation radiation. The back-scattered positrons and the mis-identified Bhabha scattering produced the long tail on the left of the main peak. The singles events from Bhabha scatterings are best seen when no geometric cuts are applied, as shown in Figures 5.25a and 5.25b.

The angular distributions are most affected by multiple scattering in the target. The chamber and target position, the resolution and beam size also affect the angular distributions. The simulated scattering angle distribution is shown in Figure 5.26a and the experimental scattering angle distribution is shown in Figure 5.26b. The sharp $1/\sin^4\theta$ feature is clear in both distributions. The individual anode wires are not present in the
simulation, causing the unevenness they produced in the scattering angle distribution to be absent. The distributions without the 22° to 43° cut are shown in Figures 5.27a and 5.27b. The sharp lower cut-off in the simulation is caused by the dead region in the detector to be considered 0 % efficient, whereas the real detectors have a gradually dropping efficiency. The discontinuity near 44° is caused by the acceptance of the detectors in both cases.

The simulated azimuthal angle distribution is shown in Figure 5.28a and the experimental azimuthal angle distribution is shown in Figure 5.28b. The repeated cuts are applied to avoid all detector boundaries. The uncut distributions are shown in Figures 5.29a and 5.29b. The four broad distributions represent the four chambers. The simulation does not have the small drop in the middle of each chamber because the thin dead layer causing the dip at the boundary of the scintillators was not included.

2. Coincidences histograms

The coincidence distributions contain Bhabha events, resonant events, or both depending on what was chosen. The distributions below compare the data to a simulation with only Bhabha events selected.

The simulated sum energy distribution is shown in Figure 5.30a and the experimental energy sum is shown in Figure 5.30b. A remarkable agreement is obtained between the simulated and experimental distributions. This is more noteworthy than with singles events because the information used to set up the simulation was extracted from direct beam into the detectors.

The simulated and experimental energy distributions for individual particles in coincidences are shown in Figures 5.31a and 5.31b respectively. The lower peak is produced mainly by large angle low energy scattered electrons and the higher peak is produced mostly by small angle high energy scattered positrons.

The simulated opening angle sum distribution is shown in Figure 5.32a and the experimental opening angle distribution is shown in Figure 5.32b. As described earlier, the sharp features of the opening angle sum distribution are blurred by multiple scattering. The multiple scattering value used in the simulation reproduces the width of both the opening angle sum and difference, in the laboratory and in the center of mas reference frames respectively, and the azimuthal angular difference, as will appear later. Both were fit with gaussians and the results are shown in Table 5.1. The simulated and experimental
center of mass scattering angle difference distributions are shown in Figures 5.33a and 5.33b respectively. Again, the width is closely reproduced, as can be seen in Table 5.1.

The scattering angle distributions of individual particles in a coincidence can also be compared. Although the simulation can distinguish electrons from positrons, this is not possible in the experiment. The experimental data is therefore compared to the sum of the electron and positron distributions. The simulated and experimental center of mass scattering angle distributions are shown in Figures 5.34a and 5.34b respectively. The positron and electron peaks are clearly visible. The simulated and experimental laboratory scattering angle distributions are shown in Figures 5.35a and 5.35b. Again both peaks are reproduced in the simulation. Since the distributions of individual particles in a coincidence are rather featureless, not much can be learned about the quality of the simulation by studying them, but they were included for the sake of completeness.

<table>
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</table>

The simulated azimuthal difference distribution is shown in Figure 5.36a and the corresponding experimental distribution in Figure 5.36b. Both are typically centered at 0° and have similar widths. The individual angular distributions are azimuthally symmetric and hence only reproduce the cuts, if applied (see Figures 5.37a and 5.37b), or the detector acceptance (see Figures 5.38a and 5.38b).

3. Detector Histograms

As was done in the experiment, the position of an event in the chamber can be histogrammed before it is is reconstructed into an angle. In the simulation, the anode and cathode hits were monitored.
The simulated and experimental cathode distributions are displayed in Figures 5.39a and 5.39b. The main shape of both distributions is a combination of the detector acceptance and the angular dependence of the scattering processes. The distributions have small fluctuations caused by the steel wires supporting the vacuum window.

The simulated and experimental anode distributions (see Figures 5.40a and 5.40b) have the same overall forward angle peak caused by the Mott and Bhabha angular dependence. The experimental distribution has narrow peaks caused by the discrete nature of the resistive anode readout. This was not included in the simulation. The smooth drop-off at the edges of the anode is a combination of slowly decreasing efficiency, worsening resolution, and pedestals taking on an increasing importance as one nears an edge. Since these effects are difficult to separate and since the detector edges were not used in the experiment, these effects were not simulated.

I. STUDIES OF INDIVIDUAL EFFECTS

As stated earlier, various physical phenomena can artificially be turned on or off. The various contributions to the broadening of the angular distributions can be illustrated and studied. In the following section, one experimental consideration at a time will be added, and the change to the reconstructed distributions will be monitored. The later sections will focus on studying other individual phenomena.

1. Peak Broadening

Initially, the simulation was run without position offsets, beam spot size, detector resolution and dead region or multiple scattering. The resulting distributions are shown in Figure 5.41.

When the wire chambers are positioned to simulate the experimental conditions, as opposed to an ideal alignment, the distributions from Figure 5.42 are obtained. As expected, the azimuthal difference is broadened from a single channel at 180° to a gaussian distribution, as shown in Figure 5.42c. The next most influenced distribution is the center of mass back-to-backness, where any small change in position causes a larger change in angle due to the relativistic boost to the center of mass (see Figure 5.42e).

If the beam spot is allowed to grow from a point to a 1.5 mm diameter spot, the only distribution to be significantly affected is the azimuthal difference. This is due to the small radial distance between the beam axis and the point of detection of the particles. Any
deviation from the theoretical beam axis causes a significant change in reconstructed angle, as seen in Figure 5.43c.

The detector resolution is the second most significant contribution to the broadening of the angular distributions (see Figure 5.44). The azimuthal difference (Figure 5.44c), mostly determined from the cathode position in a chamber, is less broadened because of the better resolution in the cathode direction. The steel support wires now appear as broad dips in the cathode distribution (Figure 5.44h), as opposed to narrow dead regions as in Figure 5.43h.

The opening angle sum (Figure 5.44d) loses its distinctive shape and appears like a gaussian with an overly wide top. The center of mass difference in scattering angle (Figure 5.44e) is broadened more than the laboratory opening angle sum due to the relativistic effects discussed above.

Adding multiple scattering broadens the angular distributions still further. Again, the effect on the azimuthal difference (Figure 5.45c) is larger than on the opening angle sum (Figure 5.45d) or the center of mass difference (Figure 5.45e) because, for a given cone of multiple scattering, a larger azimuthal angle than opening angle will be subtended.

The relativistic boost to the center of mass is not linear. A small difference in reconstructed angle at low scattering angle will be boosted differently from one at large scattering angle. This causes a different width, and hence height of the center of mass scattering angle peaks (Figure 5.45f). This effect was already present due to earlier angular broadening, and it becomes more noticeable as the angular reconstruction worsens.

When the beam and target misalignment are included, the distributions of Figure 5.46 are produced. The uneven acceptance of the individual detectors are displayed in the singles azimuthal angle distribution (Figure 5.46b). There is also a shift in the centroid of the azimuthal angle difference, the opening angle sum and the center of mass scattering angle difference. This is the case if the point of interaction offsets are included in the simulation but not included in the analysis.

The only real situation is of course when all the effects are acting simultaneously, and this is what should be compared to the experimental data. The previous section performed exactly this comparison. The set of simulations displayed in this section, although unphysical, was executed to compare the relative effects of the different contributions to the uncertainty in the position of an event.
2. Low Energy Tail

To illustrate the contribution to the tail in the singles energy spectrum from Bhabha scattering events, the back-scattered positrons cross section was turned off. The resulting singles spectrum is shown in Figure 5.47. As discussed earlier, the low energy bump is due to scattering events detected at large angles and the shoulder on the left of the Mott peak is caused by detections at small angle where the second particle of the coincidence was not detected.

3. Resonance without Bhabha Scattering

As a check to the simulation, one can eliminate Bhabha scattering and only allow for resonant scattering. The characteristic resonance shape should then be reproduced, with appropriate errors. Such a run is shown in Figure 5.48.

J. SUMMARY

By writing a Monte Carlo simulation of the experiment, a full understanding of the experiment was gained, the experimental results were reproduced from basic theoretical principles, and the method of extracting resonance amplitudes was successfully tested.

The Monte Carlo provided the insight in understanding the assymmetry between the two peaks in the center of mass scattering angle. It also was a means of investigating the effects of changing efficiencies and geometries. The loss in acceptance at different scattering angles from the combinations of the steel wires and multiple scattering could only be investigated with a full simulation.

The main achievement of the simulation was that it reproduced the measured energy and angular distributions for singles and coincidences, and hence the Bhabha to Mott ratios. These results were obtained by using the established theoretical cross sections for Mott scattering and Bhabha scattering, the scattering kinematics and experimental geometry, and by including back-scattering and multiple scattering and the wire chamber efficiencies.

The possible existence of resonances was included in the simulation. This allowed us to investigated the expected coincidences to singles ratios for different types of coupling and for different lifetimes. The inclusion of resonances also made it possible to test the
algorithms used to derive the amplitude of a potential resonance from the data, as will be shown in the next chapter.
Figure 5.1: Electron momentum distribution in lithium.

Figure 5.2: Simulated cathode distribution. The regularly spaced dips are produced by the steel support wires.
Figure 5.3: Simulated excitation functions with different chamber efficiencies. The top curve was generated with an efficiency of 84.5 %, the bottom curve with an efficiency of 82.5 %.

Figure 5.4: Energy distribution of direct beam sent into the scintillator. Note the low energy tail and the high energy shoulder produced by Compton Scattered 511 keV gamma rays.
Figure 5.5: Azimuthal angular distribution of Mott scattered positrons. The distribution reflects the acceptance of the four wire chambers.

Figure 5.6: Schematic path of back-scattered positrons.
Figure 5.7: Energy dependence of the back-scattering fraction of electrons incident on different materials. From [Tab71].

Figure 5.8: Back-scattering energy dependence of electrons incident on iron.
Figure 5.9: Fit to the experimental singles spectrum to parametrize the tail.

Figure 5.10: Energy dependence of the singles energy tail parameter.
Figure 5.11: Center of mass scattering angle distribution of both positrons and electrons.

Figure 5.12: Difference in center of mass scattering angle for a coincidence. $0^\circ$ corresponds to a back-to-back scattering.
203

**Figure 5.13**: Simulated laboratory scattering angle distribution for positrons in a coincidence.

**Figure 5.14**: Simulated laboratory scattering angle distribution for electrons in a coincidence.
Figure 5.15: Combined simulated laboratory scattering angle distribution for positrons and electrons in a coincidence.

Figure 5.16: Sum opening angle for a coincidence in the laboratory frame of reference.
Figure 5.17: Azimuthal angle distribution for positrons in a coincidence. The electron distribution is identical.

Figure 5.18: Azimuthal angle difference for a coincidence. 180° corresponds to a back-to-back scattering (in the experiment, this was set at 0°).
Figure 5.19: Center of mass scattering angle difference after spot size, resolution and multiple scattering are included.

Figure 5.20: Azimuthal angle difference after spot size, resolution and multiple scattering are included.
Figure 5.21: Laboratory opening angle sum after spot size, resolution and multiple scattering are included.

Figure 5.22: Center of mass scattering angle for positrons and electrons. The asymmetry between the small and large angles is caused by the Lorentz boost.
Figure 5.23: Singles energy spectrum with low energy tail and mis-identified Bhabha events included. The coincidence events detected as singles appear near the lower level discriminator and on the low energy shoulder of the Mott peak.

Figure 5.24: Singles energy distribution with cuts. Figure 5.24a was simulated, Figure 5.24b was experimentally obtained.
Figure 5.25: Singles energy distribution without cuts. Figure 5.25a was simulated, Figure 5.25b was experimentally obtained.

Figure 5.26: Singles scattering angle distribution with cuts. Figure 5.26a was simulated, Figure 5.26b was experimentally obtained.
Figure 5.27: Singles scattering angle distribution without cuts. Figure 5.27a was simulated, Figure 5.27b was experimentally obtained.

Figure 5.28: Singles azimuthal angle distribution with cuts. Figure 5.28a was simulated, Figure 5.28b was experimentally obtained.
Figure 5.29: Singles azimuthal angle distribution without cuts. Figure 5.29a was simulated, Figure 5.29b was experimentally obtained.

Figure 5.30: Coincidence sum energy distribution. Figure 5.30a was simulated, Figure 5.30b was experimentally obtained.
Figure 5.31: Energy distribution of individual particles in a coincidence. Figure 5.31a was simulated, Figure 5.31b was experimentally obtained.

Figure 5.32: Opening angle sum distribution for coincidences. Figure 5.32a was simulated, Figure 5.32b was experimentally obtained.
Figure 5.33: Center of mass scattering angle difference distribution for coincidences. Figure 5.33a was simulated, Figure 5.33b was experimentally obtained.

Figure 5.34: Center of mass scattering angle distribution for both positrons and electrons in coincidences. Figure 5.34a was simulated, Figure 5.34b was experimentally obtained.
Figure 5.35: Laboratory scattering angle distribution for both positrons and electrons in coincidences. Figure 5.35a was simulated, Figure 5.35b was experimentally obtained.

Figure 5.36: Azimuthal angle difference distribution for coincidences. Figure 5.36a was simulated, Figure 5.36b was experimentally obtained.
Figure 5.37: Azimuthal angle distribution for both positrons and electrons in coincidences, with cuts. Figure 5.37a was simulated, Figure 5.37b was experimentally obtained.

Figure 5.38: Azimuthal angle distribution for both positrons and electrons in coincidences, without cuts. Figure 5.38a was simulated, Figure 5.38b was experimentally obtained.
Figure 5.39: Distribution of all hits along the cathode readout of the wire chambers. Figure 5.39a was simulated, Figure 5.39b was experimentally obtained.

Figure 5.40: Distribution of all hits along the anode readout of the wire chambers. Figure 5.40a was simulated, Figure 5.40b was experimentally obtained.
Figure 5.41: Histograms of run without corrections. Figure 5.41a is the singles scattering angle distribution, 5.41b the singles azimuthal distribution, 5.41c the coincidence azimuthal difference, 5.41d the coincidence sum opening angle, 5.41e the coincidence center of mass difference, 5.41f the coincidence center of mass scattering angle distribution, 5.41g the anode distribution, 5.41h the cathode distribution and 5.41i the coincidence scattering angle distribution.
Figure 5.42: Histograms of run with the wire chamber position corrections. Figure 5.42a is the singles scattering angle distribution, 5.42b the singles azimuthal distribution, 5.42c the coincidence azimuthal difference, 5.42d the coincidence sum opening angle, 5.42e the coincidence center of mass difference, 5.42f the coincidence center of mass scattering angle distribution, 5.42g the anode distribution, 5.42h the cathode distribution and 5.42i the coincidence scattering angle distribution.
Figure 5.43: Histograms of run with the beam spot size added. Figure 5.43a is the singles scattering angle distribution, 5.43b the singles azimuthal distribution, 5.43c the coincidence azimuthal difference, 5.43d the coincidence sum opening angle, 5.43e the coincidence center of mass difference, 5.43f the coincidence center of mass scattering angle distribution, 5.43g the anode distribution, 5.43h the cathode distribution and 5.43i the coincidence scattering angle distribution.
Figure 5.44: Histograms of run with the detector resolution added. Figure 5.44a is the singles scattering angle distribution, 5.44b the singles azimuthal distribution, 5.44c the coincidence azimuthal difference, 5.44d the coincidence sum opening angle, 5.44e the coincidence center of mass difference, 5.44f the coincidence center of mass scattering angle distribution, 5.44g the anode distribution, 5.44h the cathode distribution and 5.44i the coincidence scattering angle distribution.
Figure 5.45: Histograms of run with multiple scattering added. Figure 5.45a is the singles scattering angle distribution, 5.45b the singles azimuthal distribution, 5.45c the coincidence azimuthal difference, 5.45d the coincidence sum opening angle, 5.45e the coincidence center of mass difference, 5.45f the coincidence center of mass scattering angle distribution, 5.45g the anode distribution, 5.45h the cathode distribution and 5.45i the coincidence scattering angle distribution.
Figure 5.46: Histograms of run with beam and target misalignment added. Figure 5.46a is the singles scattering angle distribution, 5.46b the singles azimuthal distribution, 5.46c the coincidence azimuthal difference, 5.46d the coincidence sum opening angle, 5.46e the coincidence center of mass difference, 5.46f the coincidence center of mass scattering angle distribution, 5.46g the anode distribution, 5.46h the cathode distribution and 5.46i the coincidence scattering angle distribution.
Figure 5.47: Singles energy spectrum without back-scattered positrons. Notice the contribution from Bhabha events detected as singles near 500 keV and below beam energy.

Figure 5.48: Excitation function with Bhabha scattering turned off.
6. RESULTS

This chapter describes how a physical upper limit to a resonance is obtained from the data. It then shows how the experiment was monitored and the systematic studies were performed. The results are then compared to those obtained with the Monte Carlo simulation.

A. THE BHABHA/MOTT RATIO

After the events are reconstructed, cuts are applied and corrections are made. The normalized coincidence spectrum is then simply the ratio of coincidences to singles at each energy point. The ratio extracted from the data is shown in Figure 6.1. There are no obvious peaks in the data. This means one has to look for a small perturbation to the large Bhabha scattering background. It is done by fitting the data.

One can fit the overall shape of the distribution, determined by Bhabha scattering, and look at the residuals. If a peak is observed, it is an indication of a resonance. If no statistically significant peaks are observed, an upper limit on the size of any resonance can be derived.

Alternatively one can fit the background and any potential resonances at the same time. If the fitting consistently finds a significant peak, it would be an indication of a resonance. And, as for the residuals, if no peaks are found, upper limits can be set. Fitting the peak and the background simultaneously is a more correct treatment of the data, although the residuals are the most obvious visual display of the data.

1. General Fit

The simplest fit to the data is a linear fit. It is shown in Figure 6.2. The reduced $\chi^2$ for the fit is 1.08, suggesting a line to be a good description of the excitation function. A quadratic polynomial was also fitted to the data and is shown in Figure 6.3. Since the quadratic fit has a larger $\chi^2$ than the linear fit, the linear fit was preferred. The linear nature of the excitation function is also independently confirmed by the Monte Carlo simulation.

By subtracting the linear fit from the data, one eliminates the trends generated by the Bhabha and Mott scattering distributions and one is left with the statistical fluctuations and
any potential resonances. The residual spectrum is shown in Figure 6.4. Again, no peaks of the expected width are immediately noticeable. A quantitative search for peaks will therefore be carried out.

2. Fit to the Resonance

To confirm the presence or absence of peaks, the data were fit with a linear background and peaks of varying center and amplitude. The width was fixed at the expected width produced by the electron motion in the target, to avoid unphysical fits. No consistent peaks were found.

To obtain an upper limit for a peak, the data were fitted with a linear background and a gaussian of known width centered at each point, as shown in Figure 6.5 for the 2176 keV data point. The amplitude of the gaussian at each point is an estimate of the size of a resonance at that point. The error in the amplitude is a measure of the significance of the amplitude of a peak.

The amplitudes and errors of the gaussian fits are shown in Figure 6.6. The first and last two points were not fit because there would have been no data more than one standard deviation away from the center to fit to. All points are within two standard deviations from zero, again indicating no positive sign of a peak. The error bars are smaller near 2230 keV because there were more data points in that region and some with high statistics.

The shape of a resonance in lithium was derived in Chapter 2 and was shown in Figure 2.40. The sharp peak is caused by the valence electrons, whereas the broad peak is the contribution of the core electrons. The peak amplitude fitting described above should be performed with the derived resonance shape and not with a gaussian. However, the fitting algorithms used fits to analytical functions and not to tables. The resonance was fitted with several functions to determine the best analytical description of the shape (see Figure 6.7a for a gaussian fit, Figure 6.7b for a fit with a Lorentzian, Figure 6.7c for a fit with a gaussian and a Lorentzian, and Figure 6.7d for a fit to 2 gaussians). The combination of two gaussians with the same center produced the best analytical description of the resonance shape.

With the resonance shape fit, the relative amplitudes and the widths of the two gaussians can be fixed, leaving just one degree of freedom, the overall amplitude. This amplitude can then be fit, with a linear background, to the data, just as was done in section
2. The resulting amplitudes are shown in Figure 6.8. The main features of the spectrum are similar to the single gaussian fit, but the amplitudes are somewhat different. Since the fit with the two gaussians is more accurate than the one gaussian fit, the data from Figure 6.8 will be used in the rest of the analysis.

3. Upper Limit to the Amplitude

With the amplitude and error of the resonance calculated, an upper limit can be derived. A problem arises with negative amplitude resonances. Different, more accurate upper limits are obtained if a positive amplitude physical constraint is introduced.

A two or three standard deviation upper limit, corresponding to a 95.4 % or 99.7 % confidence level respectively, can be calculated by adding two or three standard deviations to the fit amplitude. The results are shown in Figures 6.9 and 6.10. The two standard deviation upper limit still contains some unphysical negative amplitudes. The three standard deviation upper limit contains no negative amplitudes but has a very small upper limit near 2210 keV.

The above upper limits were derived using both positive and negative resonance amplitudes obtained from the fits. However, the resonance can only be zero or a positive number. This information must be included in the analysis. By using the negative amplitudes in the same way positive amplitudes are, a false impression of a powerful upper limit is created. To remedy this problem, the procedure outlined in the Review of Particle Properties [Rev90] is followed.

The gaussian probability distribution is renormalized so that the integral between zero and infinity is 1.0 (see Figure 6.11). The distribution is then again integrated until the point where the integral equals the desired confidence level. This point is then the upper limit at the desired confidence level. This method converges with the method used in the previous section when the physical limit vanishes or is much less than the measured amplitudes. As stated in the Review of Particle Properties, this is a conservative upper limit.

The resulting upper limits to the amplitude of a resonance are shown in Figure 6.12 and 13, for a 90 and 95 percent confidence level, respectively. By comparing the 2σ unnormalized upper limits with the 95% confidence level normalized upper limits, the general agreement for positive amplitudes is clear. The upper limits for negative amplitudes are increased, as expected. One informal argument for these limits is to consider they can
not be smaller than the errors in the amplitudes on top of the physical lower limit to the amplitude, 0 in our case. Actually, this method is sometimes also used to quote an upper limit with a physical constraint.

4. Upper Limit to the Cross Section

The upper limits calculated so far are in units of the experimental Bhabha/Mott ratio. This has to be converted into physical units.

4.1 Limit in Barns

The first step in obtaining a cross section in barns is to calculate the Bhabha scattering cross sections for the angular acceptance of the experimental cuts. The Bhabha to Mott ratio can then be equated with a certain physical cross section leading to the fitted amplitude of the resonance being given an absolute physical value, i.e. the amplitude of the fitted line at each point is equated to the Bhabha scattering cross section at each energy. The upper limit to the cross section can then be derived by using

\[
\frac{A_{\text{upper limit}}}{A_{\text{fit line}}} = \frac{\sigma_{\text{upper limit}}}{\sigma_{\text{Bhabha}}},
\]

where \(A_{\text{upper limit}}\) is the upper limit to the amplitude derived in the previous section, \(A_{\text{fit line}}\) is the amplitude of the fit background line, \(\sigma_{\text{Bhabha}}\) is the integrated Bhabha cross section and \(\sigma_{\text{upper limit}}\) is the desired upper limit to the cross section of a resonance. The integrated Bhabha cross section is given by

\[
\int_{118^\circ}^{180^\circ} \left(\frac{d\sigma}{d\Omega}\right)_{\text{Bhabha}} d\Omega,
\]

where \(\left(\frac{d\sigma}{d\Omega}\right)_{\text{Bhabha}}\) is the differential cross section from equation (2.35) and the limits of integration are the same as the cuts applied to the data. The upper limit derived from the data is shown in Figure 6.14, in barns, assuming an isotropic distribution of positron-electron pairs in the center of mass.

The cross section of the amplitude calculated is the cross section into the angular acceptance of the applied cuts assuming an isotropic distribution. To calculate the total cross section, the angular distribution of the resonance must be known, since different angular distribution would produce different numbers of coincidences in the accepted
range. The angular distribution depends on the nature of the resonance. The angular dependence for the different types of coupling were given in equation (2.48).

The angular dependence is introduced by the $t$ variable and is present for the vector and axial vector cases. The angular distributions $S_V$, the vector resonance scattering angle distribution in the center of mass, and $S_A$, the axial vector scattering angle distribution in the center of mass, are shown in Figure 2.34. The angular distributions for the scalar and pseudo-scalar cases are isotropic in the center of mass. Since the angular distributions are different, the angular acceptance of the detectors will depend on the type of resonance being simulated. Whereas the scattering angle integrals have to be performed for each distribution, the azimuthal distributions are uniform and do not depend on the type of resonance involved. The scattering angle integral ratio

$$
\frac{\int_{62^\circ}^{118^\circ} \mathcal{F} \sin \theta \, d\theta}{\int_{0^\circ}^{180^\circ} \mathcal{F} \sin \theta \, d\theta}
$$

(6.3)

gives the fraction of the theta integral (with a maximum integral of 2) that is within the applied cuts. Integration over the theta angle yields values of

- 0.939 for the scalar and pseudo-scalar cases,
- 0.877 for the vector case, and
- 0.759 for the axial vector case,

out of a total cosine theta integral of 2. Since the $S$ functions are a function of $s$, the squared center of mass energy, the integrals must be performed at one energy. The integrals were performed for $\sqrt{s} = 1832$ KeV, although the difference in the integrals over the range of energies of interest was less than 0.3 % for the vector case, with no difference for the other cases. The azimuthal integral was $\pi$ out of $2\pi$.

The total cross section is determined by integrating the angular distributions over the experimental range and scaling the total cross section by $4\pi$ over the integral. The total cross section upper limits for each of the four cases are shown in Figures 6.14 through 6.16. The variation between the cases are the scale factors shown above.

### 4.2 Limit in Barn-ElectronVolts/Steradian

Upper limits for a resonance search are sometimes quoted in units of barn-electronVolts per steradian. The advantage of this method is that one does not worry about
the type of particle and different angular distributions. The measured upper limit to the cross section, $\sigma_{\text{max}}$ from the last section, is divided by the angular acceptance to yield barns per steradian and is then multiplied by the width of the resonance to yield the quoted upper limit in barns-electronVolts/steradian. The upper limit $U$ is therefore given by

$$
U = \frac{\Delta E \times \sigma_{\text{max}}}{\Delta \Omega},
$$

(6.4)

where $\Delta E$ is the width of the resonance as determined from the full width at half maximum from the derived resonance shape in conjunction with the energy loss in the target and the beam energy width. The solid angle $\Delta \Omega$ was calculated in the last section. The limits in barns-electronVolts/steradian are shown in Figure 6.17.

4.3 Limit on resonance lifetime

Another way of expressing the upper limit to a resonance is to quote the lifetime of the resonance cross section corresponding to the upper limit. Three cases have to be considered. The scalar and pseudo-scalar couplings, having identical angular distributions and spin, are bound by the same lifetime limits. The vector and axial vector couplings, while both $J = 1$ couplings, produce distinct lifetime limits because of the different total cross sections generated by the variations in the angular distributions. The lifetime limits are derived from the cross section limits using equation (2.106). From the equation we see that the resonance cross section is proportional to the width of the resonance and to the spin by a factor of $(2J+1)$. However, there is also an energy dependent factor that must be corrected for, the $(p^2 - 4)^{3/2}$ factor appearing in the denominator. This energy dependence, Figure 6.18, shows a decrease in the peak cross section of a resonance as the energy of the resonance is increased.

With the energy and spin dependent corrections included, the lifetime upper limits are as shown in Figure 6.19.

B. COMPARISON WITH THE MONTE CARLO

Since the energy and angle distributions are well reproduced, one can consider the Monte Carlo simulation to be an adequate model of the experiment. One can now consider what the excitation function would look like in the simulation. First, one must consider the overall rates, since they have not been discussed yet.
1. Rates

The simulation was run for 50 minutes with 500000 positrons per second on a 1.5 mg/cm$^2$ lithium target, as was the experiment. The produced singles and coincidence rates actually matched the experimental values closely. If one allows the chamber efficiency to be adjusted to match one experimental point, then the excitation function calculated closely matches the experimentally measured one, as discussed below.

2. Excitation Function

The Monte Carlo simulation was run in intervals of 5 keV from 2150 to 2350 keV. In addition, to simulate the high statistics points in the experiment, it was run for double the statistics between 2207.5 and 2277.5 also in 5 keV steps. The resulting excitation function is shown in Figure 6.20. The experimental result is shown in Figure 6.21. Both sets of points are clearly described by a line with the same slope and intercept. For comparison, the ratio of analytic Bhabha cross section to Mott cross sections is shown in Figure 6.22.

The simulation has therefore successfully reproduced the experimental Bhabha to Mott ratio and one can be confident of understanding the physical processes, including any possible resonances, in the experiment. One can perform the search for a resonance with the simulated excitation function as was done for the experimental data. This will be the subject of the next section.

3. Upper Limit

Since no peaks can be found in fits to the simulation, it will be analyzed for an upper limit, as was the data. This means it was fitted at every point with a functional form of the line shape and the fitted amplitudes and errors recorded (see Figures 6.23 and 6.24).

The normalized upper limits are obtained using the method used for the experimental data. The upper limits are shown in Figure 6.25. A tighter upper limits is set in the high statistics region as for the data. This is caused by the smaller errors in the fit amplitudes. The statistical fluctuations are of the same size as the data, implying that the systematic errors have been reduced to less than the statistical errors.

4. Resonances

The simulation successfully reproduced the experimental data. But it can also be used to predict the shape of the excitation function for runs with different targets and
statistics and for resonances of different lifetimes and coupling types. Such simulated runs are discussed below.

The excitation function for a run with four times the statistics is shown in Figure 6.26. This would correspond to statistical accuracy of the long experimental run performed at the dynamitron. The Monte Carlo was run without any resonances turned on. Figures 6.27 through 6.29 show similar runs with a $10^{-13}$ second lifetime resonance of scalar, pseudo-scalar, vector and axial vector coupling type. The lifetime dependence is illustrated in Figures 6.30 and 6.31 by similar runs with $5 \times 10^{-14}$ second and $10^{-14}$ second lifetimes respectively.

The simulated peaks were fitted to determine the "measured cross section", as would be done experimentally. This was done to check that the correct cross section is extracted from the excitation function by the analysis method. The simulated and measured cross sections are compared in Table 6.1.

<table>
<thead>
<tr>
<th>Type of resonance</th>
<th>lifetime $\times 10^{-14}$ seconds</th>
<th>simulated cross section $\times 10^{-2}$ barns</th>
<th>calculated cross section $\times 10^{-2}$ barns</th>
</tr>
</thead>
<tbody>
<tr>
<td>axial vector</td>
<td>1</td>
<td>7.17</td>
<td>7.14 $\pm$ 0.1</td>
</tr>
<tr>
<td>vector</td>
<td>1</td>
<td>7.17</td>
<td>7.05 $\pm$ 0.09</td>
</tr>
<tr>
<td>pseudo scalar</td>
<td>1</td>
<td>2.39</td>
<td>2.28 $\pm$ 0.07</td>
</tr>
<tr>
<td>pseudo scalar</td>
<td>5</td>
<td>0.478</td>
<td>0.455 $\pm$ 0.06</td>
</tr>
</tbody>
</table>
Figure 6.1: Ratio of coincidences to singles. A resonance would appear as a ~20 KeV wide peak on top of the linearly increasing background.

Figure 6.2: Linear fit to the data.
Figure 6.3: Quadratic fit to the data.

Figure 6.4: Residuals of data with linear background subtracted.
Figure 6.5: Fit of a line and a gaussian with a fixed mean and standard deviation. The mean was fixed at each data point and the best fit of the amplitude determined.

Figure 6.6: Amplitude and error in the amplitude of the fit gaussian at each energy.
Figure 6.7: Fits to the resonance. Figure 6.7a is a gaussian fit, Figure 6.7b a lorentzian fit, Figure 6.7c a gaussian and a lorentzian fit and Figure 6.7d shows a fit of two gaussians to the resonance profile.
Figure 6.8: Amplitude and error in the amplitude of the fit gaussians at each energy.

Figure 6.9: Two standard deviation upper limit to the data in Figure 6.8.
Figure 6.10: Three standard deviation upper limit to the data in Figure 6.8.

Figure 6.11: With a lower bound at zero \( (\mu_{\text{min}}) \), the distribution is renormalized so that the integral between zero and infinity \( (= \text{the shaded area}) \) is one.
Figure 6.12: Normalized upper limit to the amplitude of a fitted resonance to the data with 90% confidence level.

Figure 6.13: Normalized upper limit to the amplitude of a fitted resonance to the data with 95% confidence level.
Figure 6.14: Upper limit to the cross section of a fitted resonance to the data assuming isotropic distribution of positron-electron pairs in the center of mass, i.e. for scalar and pseudo-scalar coupling. The triangles represent the 90% confidence level, the squares the 95% confidence level.

Figure 6.15: Upper limit to the cross section of a fitted resonance to the data for a vector resonance. The triangles represent the 90% confidence level, the squares the 95% confidence level.
Figure 6.16: Upper limit to the cross section of a fitted resonance to the data for an axial vector resonance. The triangles represent the 90% confidence level, the squares the 95% confidence level.

Figure 6.17: Upper limit to the cross section of a fitted resonance to the data in barns-electronVolts/steradian. The triangles represent the 90% confidence level, the squares the 95% confidence level.
Figure 6.18: Energy dependence of the peak resonance cross section. Note the decrease in peak cross section with increasing energy for a given resonance lifetime.
Figure 6.19: Lifetime upper limit to the data for the different resonance couplings. Figure 6.18a is the limit for a scalar resonance or a pseudo-scalar resonance, Figure 6.18b the limit for a vector resonance, and Figure 6.18c the limit for an axial vector resonance. The triangles represent the 90% confidence level, the squares the 95% confidence level.
Figure 6.20: Simulated excitation function.

Figure 6.21: Experimental excitation function.
Figure 6.22: Ratio of Bhabha cross section (integrated between 40° and 140° in the center of mass) to Mott cross section (integrated between 15° and 55°). Both were integrated over 360° azimuthal angle.

Figure 6.23: Fit of the resonance shape and a linear background to the simulated excitation function. Such a fit was performed at each energy to determine the best value of the amplitude of the resonance at that energy.
Figure 6.24: Plot of the fitted amplitudes of a resonance versus energy for the simulated data set.

Figure 6.25: Upper Limit to a resonance for the simulated data. The upper curve shows a 95% confidence level, the bottom curve a 90% confidence level.
Figure 6.26: Simulation of an excitation function with .25% errors in the good coincidence rate.

Figure 6.27: Simulation of an excitation function with .25% errors in the good coincidence rate with a $10^{-14}$ second scalar or pseudo scalar resonance at 2260 keV.
Figure 6.28: Simulation of an excitation function with .25% errors in the good coincidence rate with a $10^{-14}$ second vector resonance at 2260 keV.

Figure 6.29: Simulation of an excitation function with .25% errors in the good coincidence rate with a $10^{-14}$ second axial vector resonance at 2260 keV.
Figure 6.30: Simulation of an excitation function with .25% errors in the good coincidence rate with a $5 \times 10^{-14}$ second scalar or pseudo scalar resonance at 2260 keV.

Figure 6.31: Simulation of an excitation function with .25% errors in the good coincidence rate with a $10^{-13}$ second scalar or pseudo scalar resonance at 2260 keV.
7. CONCLUSIONS AND OUTLOOK

Having derived an upper limit to a resonance from our data, the results can be compared to those of other groups. The need for an extended and improved data set then becomes clear. The experiment was repeated and the results, based on an identical data analysis, will be discussed. To investigate the possibility of longer lived particles, the Bhabha background would have to be eliminated. By constructing a thin active beam dump and using it in conjunction with the same detector system as described in this thesis, an experiment was performed to search for longer lived particles. Combining the results of the extended Bhabha run and the beam dump experiment, a large portion of the mass-lifetime domain was excluded. Further experiments are still warranted though, and some currently being executed will soon provide interesting results. Finally, the peaks in heavy ion interactions are being investigated by a new generation of high sensitivity experiments.

A. GENERAL CONCLUSIONS

The collaboration of Brookhaven National Laboratory-CCNY-Yale successfully refurbished and converted the 3 MeV Dynamitron electron accelerator into a positron accelerator. The modified accelerator produced high intensity (> $5 \times 10^5$ e+/s) positron beams tunable from 1.0 MeV to 2.5 MeV. The beams were of small energy width (keV) and could be focused into 1mm beam spots. The beam, incident on thin lithium foils allowed the collaboration to perform a sensitive search for resonances in e+-e- scattering.

The scattering of positrons on electrons was detected by a set of wire chambers and scintillators designed and built for this application. The wire chambers successfully provided high spatial resolution reconstruction of all scattering events as well as a highly efficient filter against gamma rays. Using the accurate kinematic reconstructions possible with these detectors, a clean, low background unambiguous set of measurements was made possible. A dedicated data acquisition system was built and analysis software was written.

A complete analysis of the data from a first run yielded the upper limits to a resonance derived in the last chapter. These upper limits are compared with results from other experiments in the next section.

Although the upper limits from this run are not significantly different from those derived by other groups, the experiment has many features that make it stand out:
- It is the only experiment performed with lithium, the most sensitive target material for resonance searches.
- The experiment has the highest intensity monoenergetic positron source in this energy range, allowing for runs with higher statistics.
- The beam has a very low contamination of high energy positrons and can be steered and focused down to small spot sizes.
- The wire chambers provided significantly better position resolution than the other experiments, producing cleaner spectra and allowing a full kinematic reconstruction of the events.
- The beam energy was monitored in real time by a germanium detector serving as the beam dump as data was being acquired.

B. LIMITS FROM THIS RUN

As was shown in the last chapter, no peaks were found in our search for resonances in $e^+ - e^-$ scattering. An upper limit for the lifetime was derived and is shown in Figure 7.1 (repeating figure 6.19, for $J = 0$).

For a better comparison with experimental results, the lifetime limit can be plotted in terms of the lifetime versus the invariant mass, as in Figure 7.2. The other experimental limits ([Lor88], [Tse89b], [Tse91b], [Wid91]) and the g-2 theoretical limit can be displayed on a similar graph, and a composite is shown in Figure 7.3. This experiment therefore excludes particles ($J = 0$) for lifetimes shorter than $\sim 1 \times 10^{-13}$ seconds in the mass range $1805 \text{ keV} < m_X < 1855 \text{ keV}$.

Although competitive with other results, we see from Figure 7.3 that an obvious improvement to the results would be to extend the data over a larger energy range. The experimental upper limit derived from the data being bound by statistical errors, the results could also be significantly improved by acquiring much more data at each energy. The experiment described in this thesis was repeated, using the hardware, software, chamber calibrations, corrections and experience gained from the previous run. Learning from the first run, the target position was checked for each energy point with a telescope and is known not to have moved to within 0.1 mm. The pressure in the wire chambers, the principal factor affecting the chamber gain, was also more carefully controlled and monitored.
C. LIMITS FROM THE EXTENDED RUN

The experiment was executed with a 2.5 mg/cm² lithium target, acquiring about four times more data per point than the previous run or ~ 300000 coincidences per point. The beam was scanned from 1350 keV to 2350 keV in 3.5 keV steps, corresponding to center of mass energy of ~ 1560 keV to ~ 1860 keV. These data are superior not only in the extended energy range covered and the increased statistics, but also in that the position of the target was monitored with a telescope for each energy point. Any possible systematic errors introduced by such motion are therefore eliminated.

The analysis of this data was performed by X. Wu [Wu92], and his results are described here. The normalized coincidences to singles ratio, derived completely analogously to the shorter experiment, is shown, with error bars (of ~ 0.25%), in Figure 7.4. The ratio was fit with a third degree polynomial in separate 130 keV intervals, and the normalized residuals, i.e., \((R_{\text{measured}} - R_{\text{fit}}) / R_{\text{fit}}\), do not display any resonance structure (Figure 7.5).

Again, by fitting the amplitude at each point with a resonance shape, an upper limit to the resonance cross section can be derived, and from that an upper limit to the particle lifetime. The 90% confidence level upper limit to the lifetime from the extended run is shown in Figure 7.6, again compared to the other limits. From these results, the existence of a particle coupling \((J = 0)\) to \(e^+ - e^-\) can be excluded in the range 1560 keV < \(m_x\) < 1860 keV for lifetimes \(\tau \leq 2.5 \times 10^{-13}\) seconds.

From Figures 7.3 and 7.6 we see that even a relatively long run cannot probe lifetimes much longer than a few times \(10^{-13}\) seconds, since the improvement in statistical errors of the data is proportional to the square root of the data acquired, i.e. the square root of the length of data acquisition. If, however, the Bhabha scattering background was eliminated, much more sensitive experiments could be performed. Again using the same experimental set up, with the addition of an active target-beam dump system, Henderson et al. ([Hen91a], [Hen91b], [Hen92]) achieved a high sensitivity to long-lived particles.

D. THE BEAM DUMP EXPERIMENT

In order to eliminate the direct Bhabha scattering background, an active target-beam dump was designed. The beam dump must stop the direct scattering events, yet be thin enough to allow the escape of relatively long-lived \((10^{-13} \text{ s}-10^{-10} \text{ s})\) particles. The design
of Figure 7.7 was therefore adopted, with a total thickness of 2 or 3 mm, depending on which target-beam dump assembly was used.

The target in this experiment was beryllium, because the construction of the beam dump required multiple steps that could not be performed in an inert environment needed for a lithium target. Behind the target was a plastic scintillator to veto all events that did not become neutrals in the target. Behind the scintillator was the main beam dump, consisting of ~ 700 μm of platinum to stop the beam and any charged particles from the target. Behind the beam dump came another scintillator, to detect any secondary charged particles produced in the beam dump.

If a long-lived neutral particle was made in the target, it would not produce a signal in the first scintillator and if it decayed downstream from the beam dump could be detected as a coincidence event in the wire chambers and their associated scintillators. Such an event is called an unvetoed event. If a neutral particle was not produced in the target, then the front beam dump scintillator would fire, and possibly the back one too, generating a vetoed event. A resonance due to neutral particle production would therefore be measured as a significant increase in unvetoed events at a certain energy.

If a short-lived particle was produced, it would decay within the beam dump and hence not be detected as an unvetoed event. The short lifetime limit can be estimated to be the travel time of the neutral particle through the beam dump, leading to a cut off in sensitivity of ~ a few times 10^{-12} seconds. The long lifetime limit on the experiment is a combination of the falling production cross section with increasing lifetime and the reduced detection yield from particle decays outside the fiducial volume of the experiment. The sensitivity of the experiment is therefore as in Figure 7.8.

Since the direct beam has been eliminated, the experiment becomes a very low count rate experiment and is therefore dominated by the "natural" background. To reduce the background, large scintillator paddles were installed around the experiment and cosmic ray events were vetoed. The background unvetoed coincidence rate was thereby reduced to ~ 13.3 events per day. With this background, from Figure 7.8, the experiment would be sensitive to particle lifetimes from a few times 10^{-12} seconds to a few times 10^{-10} seconds.

The actual rate of coincidences was ~ 2.7 counts per hour and the excitation function is shown in Figure 7.9, where the unvetoed coincidence rate has been normalized to the incoming beam current (coincidences per 10^{10} incoming positrons). After the kinematic cuts
\[ 170^\circ < \Delta \phi < 190^\circ \]
\[ \sqrt{s} < \sqrt{2m(E_+ + m)} + 60, \]

where \( \Delta \phi \) is the azimuthal difference between the tracks in a coincidence and \( \sqrt{2m(E_+ + m)} \) the Bhabha scattering invariant mass, the unvetoed rates are reduced to the levels of Figure 7.10. From these low rates (~ 3 events per day), strict limits on the existence of a long-lived particle can be derived just as for the short-lived searches. The 90% confidence level limit on the lifetime excluded as a function of energy is shown in Figure 7.11.

Combining the active target experiment with the thin target experiment of the previous section and the other limits, a final composite graph with all the excluded domains is obtained (Figure 7.12). The existence of neutral particles coupling to e\(^+\) - e\(^-\) is therefore excluded in the mass range from \( \sim 1550 \text{ keV} \) to \( \sim 1850 \text{ keV} \) for a wide range of lifetimes, although the region just below \( 10^{-12} \) seconds and \( 10^{-9} \) seconds remain unexplored. The unexplored lifetime region between \( 2 \times 10^{-13} \) s and \( 8 \times 10^{-13} \) s could soon be measured by the experiment of Cowan et al. [Cow91].

E. FUTURE EXPERIMENTS

In the near future, the most promising results should originate from the experiment at Lawrence Livermore National Laboratory by Cowan and coworkers. By building a positron trap, they are planning for target leptons with a mean energy of 8 meV, as compared to a few eV for electrons in solid targets, leading to a correspondingly larger peak cross section for any lifetime particle. Since electron beams can be made much more intense than positron beams, the group will generate a high current (10 \( \mu A \)) electron beam incident on a positron plasma of \( \sim 10^{10} \text{ e}^+/\text{cm}^3 \). With this set up, a factor of \( \sim 100 \) improvement in sensitivity is expected in the lifetime range near \( 10^{-12} \) seconds. Their results are eagerly awaited.

While the low energy e\(^+\) - e\(^-\) searches have excluded a large domain of possible particles, the peaks in heavy ion scattering remain unexplained. An upgraded generation of heavy ion experiments is planned at GSI ([Ber92], [Bok92]), and a separate experiment (APEX) is being constructed at Argonne National Laboratory [Wol91]. The APEX experiment is expected to measure an e\(^+\) - e\(^-\) coincidence rate 20 times larger than EPOS, and the e\(^+\) and e\(^-\) kinematics will be reconstructed.
With the results from these experiments expected in the near future, additional progress toward solving the puzzle of the $e^+ - e^-$ peaks in heavy ion collisions is imminent.
Figure 7.3: Composite of this result (white region at the bottom) with those of other experiments ([Lor88], [Tse89], [Jud90]), the heavy ion fiducial volume limitation and the theoretical g-2 limits.

Figure 7.4: Coincidence to singles ratio versus invariant mass for the extended run. Note the size of the error bars. From [Wu92].
Figure 7.1: 90% confidence limit on the lifetime (in units of $10^{-13}$ seconds) for a $J = 0$ particle vs beam energy, as derived in the previous chapter.

Figure 7.2: Same as Figure 7.1 but plotted on a logarithmic scale as a function of the invariant mass.
Figure 7.5: Normalized residuals versus invariant mass. From [Wu92].

Figure 7.6: Composite of the extended thin target run and the other limitations. From [Wu92].
Figure 7.7: Diagram of the target-beam dump assembly. From [Hen91a].

Figure 7.8: Expected sensitivity in counts per hour to 1832 keV, J=0 particle versus lifetime of the particle. The simulation assumes $4 \times 10^5$ positrons per second and no kinematic cuts. From [Hen91a].
Figure 7.9: Unvetoed coincidence rate versus beam energy for the 3mm assembly (fig. 9a) and the 2mm assembly (fig. 9b), without cuts. From [Hen91a].

Figure 7.10: Unvetoed coincidence rate versus beam energy for the 3mm assembly (fig. 10a) and the 2mm assembly (fig. 10b), with cuts. From [Hen91a].
Figure 7.11: 90% confidence limit on the excluded lifetimes for a $J = 0$ particle vs beam energy (fig. 11a), for vector coupling (Figure. 11b) and for axial vector coupling (fig. 11c). From [Hen91a].

Figure 7.12: Composite of the limitations on a particle ($J = 0$) from low energy positron electron scattering, with the $g$-2 theoretical limits and the EPOS fiducial volume limits. From [Hen92].
REFERENCES


