Relative Yields of Antiparticles to Particles in Au+Au Collisions at 130 and 200 GeV per Nucleon Pair

by

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Abstract

Au+Au collisions at $\sqrt{s_{NN}} = 130$ and 200 GeV at the Relativistic Heavy Ion Collider have opened a new energy regime for studying nucleus-nucleus collisions. A search for the formation of a phase of matter in which quarks and gluons interact strongly over an extended deconfined volume is of primary importance.

The PHOBOS detector was constructed to measure many observables at RHIC. The primary subsystem used in this analysis is the PHOBOS spectrometer. The spectrometer is a multiplane silicon pad detector which detects particles traversing its pads as they propagate through a strong magnetic field. The hit positions are used along with energy loss information in the silicon pads to determine both the momentum and velocity of the particles, allowing for the identification of the particle species.

One of the most basic pieces of information to be determined is the baryochemical potential of the system. This information has been determined through the measurement of the relative yields of antiparticles to particles in RHIC collisions. The values of $\langle \bar{p} \rangle/\langle p \rangle$, $\langle K^- \rangle/\langle K^+ \rangle$, and $\langle \pi^- \rangle/\langle \pi^+ \rangle$ are measured as $0.60 \pm 0.04$(stat.) $\pm 0.06$(syst.), $0.91 \pm 0.07$(stat.) $\pm 0.06$(syst.), and $1.00 \pm 0.01$(stat.) $\pm 0.02$(syst.) at $\sqrt{s_{NN}} = 130$ GeV and $0.73 \pm 0.02$(stat.) $\pm 0.03$(syst.), $0.95 \pm 0.03$(stat.) $\pm 0.03$(syst.), and $1.025 \pm 0.006$(stat.) $\pm 0.018$(syst.) at $\sqrt{s_{NN}} = 200$ GeV, respectively. These relative yields indicate values of $\mu_B/T = 0.27 \pm 0.03$(stat.) and $0.17 \pm 0.01$(stat.) at $\sqrt{s_{NN}} = 130$ and 200 GeV, respectively. These values are compared to model predictions and are used to evaluate the contribution of baryon transport relative to particle production in determining the yields of baryon at midrapidity in central Au+Au collisions at RHIC energies.

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Contents

1 Introduction  .................................................. 9
   1.1 Quantum Chromodynamics ................................. 9
   1.2 Heavy Ion Collisions and the QCD Phase Diagram ........ 12
   1.3 Particle Ratios and $\mu_B$ ............................... 17
      1.3.1 Stopping and the Baryon Free Region ................. 18
   1.4 Microscopic Models of Heavy-Ion Collisions ............ 19
      1.4.1 RQMD ............................................. 20
      1.4.2 HIJING .......................................... 20
   1.5 Thermal Models ........................................... 20
   1.6 Thesis Goal ............................................... 21

2 Experimental Setup ............................................. 23
   2.1 The RHIC Accelerator Complex ............................ 23
      2.1.1 Tandem Van de Graaff ................................ 23
      2.1.2 AGS Booster ...................................... 25
      2.1.3 Alternating Gradient Synchrotron .................... 25
      2.1.4 Relativistic Heavy Ion Collider Ring ............... 26
   2.2 The PHOBOS Detector ...................................... 26
      2.2.1 Beampipe ......................................... 28
      2.2.2 Magnet ........................................... 28
      2.2.3 Triggering systems ................................ 28
      2.2.4 Multiplicity Array ................................ 33
      2.2.5 Vertex Detectors .................................. 34
      2.2.6 Spectrometer ...................................... 34
Chapter 1

Introduction

1.1 Quantum Chromodynamics

The basic constituents of matter are electrons, protons, and neutrons. The strong force binds the constituents of nucleons together and binds nucleons into atomic nuclei. The strong force counteracts the electrostatic charge repulsion of protons in nuclei. The theory of strong interactions is governed by quantum chromodynamics (QCD). The structure of this theory is analogous, in many ways, to the theory governing the electromagnetic force, quantum electrodynamics (QED). However, QCD has key distinguishing features from QED.

The discovery of a plethora of particles in the 1960s necessitated a theory to bring some order to the series of particles. Ne’eman and Gell-Mann proposed, simultaneously, a classification scheme for the spectrum of known particles[1, 2]. This theory is generally known by Gell-Mann’s title, “The EightFold Way”, where particles were classified according to their strangeness and isospin. In 1964, Gell-Mann went further to introduce the concept of spin 1/2 quarks with a flavor property (up, down, or strange), fractional charge, and baryon number[3]. This theory brought order to the spectrum of known particles and possessed predictive power for the discovery of new particles. The previously discovered $\Delta^{++}[4]$ presented a problem for this theory. The quark content of a $\Delta^{++}$ is $u \uparrow u \uparrow u \uparrow$. This particle appears to violate the Pauli principle since it is symmetric under interchange of its quarks. An additional quantum number termed color was added which could take on three values (red, green, and blue). The addition of this quantum number allowed for
the existence of the $\Delta^{++}$ without violating the Pauli principle[5, 6] ($\Delta^{++}$ content is now $u_r \uparrow u_g \uparrow u_b \uparrow$). However, in the original theory, quarks were not believed to actually exist, but were assumed to be a mathematical tool in the classification scheme of particles.

Experiments on deep-inelastic scattering of electrons on hydrogen targets, performed in collaboration between SLAC and MIT, showed that, at high transverse momentum of the electrons, the cross-sections were consistent with scattering from a point-like rather than diffuse object[7]. This provided evidence that quarks may actually exist inside hadrons. Building on this understanding, the full theory of QCD was proposed in the early 1970s[8, 9]. The result is summarized in the QCD Lagrangian.

$$L_{QCD} = -\frac{1}{4} F_{\mu\nu}^a F^{a\mu\nu} + \sum_j \bar{\psi}_j (i\gamma^\mu D_\mu - m_j) \psi_j$$

$$F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a - g_s f_{abc} A_\mu^b A_\nu^c$$

$$D_\mu = \partial_\mu + ig_s A_\mu^a \frac{\lambda_a}{2}$$

$g_s = \sqrt{4\pi \alpha_s}$ where $\alpha_s$ is the strong coupling constant characterizing the strength of the strong force. $\psi_j$ are the quark fields of flavor $j$ and mass $m_j$. $A_\mu^a$ are the gluon fields of color $a$. $f_{abc}$ are the SU(3) structure constants and $\lambda_a$ are the Gell-Mann matrices.

The SU(3) symmetry of QCD assigns color to the gluons (unlike QED where photons have no charge). This allows gluons to interact with each other via 3 and 4 gluon interactions (see Fig. 1-1), the analogues of which do not occur in QED. These additional interactions have a profound impact on the theory. The QCD coupling between colored objects increases with distance due to anti-screening effects from the additional gluon vertices. A cloud of quark-antiquark pairs and gluons with color around a quark amplifies the color force seen at large distances. In QED, a cloud of positively and negatively charged particles surrounding a charged particle reduces the charge seen at large distances.

The coupling constants for QED and QCD are shown in Eqs. 1.4 and 1.5, respectively, to second order in perturbation theory as a function of four momentum transfer, $\mu$.

$$\alpha(\mu^2) = \alpha(0) \left[ 1 + \frac{\alpha(0)}{3\pi} \ln(\mu^2/m_e^2) \right]$$

(1.4)
\[ \alpha_s(\mu^2) = \frac{12\pi}{(33 - 2n_f) \ln(\mu^2/\Lambda^2)} \] (1.5)

\( \alpha(0) \) is 1/137 and \( m_e \) is the mass of an electron (511 keV). Eq. 1.4 is only valid for the high \( \mu^2 \) limit and, therefore, does not give \( \alpha(0) \) for \( \mu^2 = 0 \). For the strong coupling constant, \( n_f \) is the number of quark flavors with mass less than four momentum transfer \( \mu \) and \( \Lambda \) is a constant that depends on which flavor thresholds have been surpassed. \( \Lambda \) is between 100 MeV and 500 MeV. The formulae demonstrate the concept of asymptotic freedom [10, 11] for the strong force. Whereas \( \alpha \) is large at large momentum transfers (small distances), \( \alpha_s \) is small. In the opposite regime (small momentum transfers, large distances), the strong coupling constant becomes large. The effective potential for a strong field has an extra \( k \cdot r \) term, phenomenologically, relative to the QED potential, where \( k \) is a constant and \( r \) is the distance from the quark, which causes the energy required to separate a quark and an antiquark to rise (boundlessly). At large distances enough energy has been added to the system to create a quark-antiquark pair allowing for the formation of new particles. This is a demonstration of the concept of confinement where quarks are always bound.
within colorless composite systems (hadrons normally). Fig. 1-2 shows the different shape of the field for the strong force versus the electromagnetic force. The field lines for the charged particles are spread out. The extra gluon interactions condense the field lines in a tube between the quark-antiquark pair when the pair is pulled apart (and the $k \cdot r$ term dominates).

![Field lines diagram](image)

Figure 1-2: Field lines between charged particles for the electromagnetic force (part a) and colored quarks for the strong force (part b) are shown. The strong force has much more condensed field lines as a result of anti-screening by gluons.

### 1.2 Heavy Ion Collisions and the QCD Phase Diagram

The discovery of a large number of particles with masses greater than the pion mass for mesons and greater than the proton and neutron masses for baryons in the 1960s indicated that an ever-increasing number of resonances might exist. In 1965, Rolf Hagedorn introduced a model based upon this ever-increasing multitude of resonances in higher mass ranges. At low masses a discrete number of states existed while at higher masses so many resonances existed that they could be represented by a continuum spectrum of states. Hagedorn's model necessitated the existence of a limiting value of temperature for a hadron gas[12]. The addition of more energy into the system at this temperature would produce more particles in newly accessible mass ranges (and increase the entropy of the system) but
not increase the temperature of the system. The inputs to this statistical bootstrap model were the volume of a cluster (hadron) and the mass of the cluster. All other particles were generated by the model. The concept was that a system of clusters compressed into a volume where the net density is the same as the cluster's density will act as an extended cluster. Hagedorn used this model to derive the density of states per unit mass as \( \rho(m) \propto m^{-3}e^{m/T_0} \) where \( T_0 \) is the temperature before hadrons would dissolve. Hagedorn performed a fit to the current particle spectrum and found that a limiting temperature of 158 MeV described the spectrum well.

It is now believed that a phase transition occurs beyond this limiting temperature (whereas Hagedorn's model would predict the production of heavier resonances in a hadron gas whose temperature remains below the limiting temperature). This new phase of matter, called the quark-gluon plasma, consists of strongly interacting quarks and gluons. The phase transition to a quark-gluon plasma depends on the thermodynamic variables of the system (baryochemical potential and temperature). A schematic diagram of the phase transition contour is shown in Fig. 1-3 [13]. To reach a phase of matter beyond the hadronic phase experimentally, we have to compress matter enough such that hadrons overlap significantly. Nuclear collisions provide a large volume for interactions and have multiple collisions per nucleon. For these reasons, colliding heavy ions is believed to be the optimal experimental technique to be applied towards verifying this theory. Heavy ion collisions probe the low baryon density, high temperature region of the diagram which is believed to have existed at about 1 \( \mu \)s after the big bang. The high baryochemical potential region of the phase diagram is difficult to probe experimentally. These conditions may exist in the core of neutron stars where gravitational forces are high enough to compress matter into the color superconductor phase. A possible astronomical observation exists searching for an increase in the angular frequency of neutron stars caused by the loss of their quark core[13, 14].

Insight into the behavior of a quark gluon plasma requires knowledge of the structure of hadrons. QCD possesses this information; however, it is difficult to extract. Most calculations performed with QED or QCD are based on perturbative expansions of the action in orders of the coupling constant. For QED, this coupling is small, \( \alpha \approx 1/137 \). For QCD, this coupling is larger, \( \alpha_s \approx 0.1 \). The running of this coupling for small momentum
Figure 1-3: A schematic drawing of the QCD phase diagram is shown[13]. The exact contours are not precisely known. The hadronic matter region contains a nuclear matter region with a phase transition between a nuclear matter liquid and a gas of nucleons. The continuous lines represent first order phase transitions. This first order phase transition stops at a critical endpoint at low baryochemical potential where a region of smooth crossover exists between the hadronic matter and the quark gluon plasma phases. Collider experiments probe the high temperature and low baryochemical potential region. A color superconductor phase exists at high baryochemical potential where quarks group to make Cooper pairs. The actual value of $\mu_B$ for RHIC collisions will be determined in this thesis given assumptions about the temperature of the phase transition.
transfers (bound states) makes the value of the strong coupling of order 1. Interactions of all orders will make non-negligible contributions to measurements and must be considered at these small momentum transfers. Wilson developed a regularization scheme that allowed for the computation of interactions at low momentum transfers numerically[15]. This lattice regularization scheme involves the discretization of the QCD action onto a 4 dimensional space-time lattice. Unlike perturbation theory, the full contribution of an interaction is evaluated. This method is computationally intensive. Errors arise from the finiteness of the lattice spacing and the volume of the lattice used. Lattice QCD calculations did show the onset of the phase transition at temperatures between 150 and 180 MeV depending on assumptions[16]. Fig. 1-4 shows the result of a lattice QCD calculation for the energy density normalized by temperature to the fourth power as a function of temperature. A steep rise at the critical temperature is seen as the transition to the deconfined state occurs. $\epsilon_{SB}$ is the energy density in the Stefan-Boltzmann limit of a gas of free quarks and gluons. Lattice QCD calculations show that residual interactions do not allow this limit to be reached.

The transition between a quark gluon plasma and a hadron gas at low baryochemical potential is believed to be smooth at zero chemical potential (as stated in the caption of Fig. 1-3). Lattice QCD has shown that the order of the phase transition at zero chemical potential depends on the masses of the quarks involved (up, down and strange quarks)[16]. If all three quarks are considered to be light and of the same mass, the phase transition at $\mu_B = 0$ is first order. If the effect of the strange quark is neglected and the up and down quarks are massless, the transition at $\mu_B = 0$ can be second order. The set of quark masses believed to best describe reality is when the up and down quarks are light and the strange quark is much heavier. Lattice QCD predicts a smooth cross-over (no discontinuity) for this realistic set of quark masses at $\mu_B = 0$[16]. This analysis will determine the value of $\mu_B$ in RHIC collisions. The order of the phase transition will not be determined.

If a quark gluon plasma is formed in heavy ion collisions, it must return to being hadronic matter after sufficient expansion of the system. The first stage of formation will be a pre-equilibrium phase where most hard collisions occur between the partons. On the time scale of 1 fm/c, the plasma may thermalize. Experimentally, evidence for thermal equilibrium can be obtained by measuring the azimuthal asymmetry in the number of particles emerging
Figure 1-4: The energy density of a system of quarks and gluons normalized by temperature to the fourth power as a function of temperature from lattice QCD calculations is shown[16]. The red line is when only 2 light quark flavors are considered (up and down). The green line shows the result if a heavy strange quark is added. The blue line is when all three quarks are light. A clear jump in energy density at the critical temperature is seen. The arrows in the upper right are the expected values for an ideal gas of quarks and gluons (Stefan-Boltzmann limit). The discrepancy is from residual interactions of quarks and gluons in a plasma limiting the particles’ freedom.

from a collision in the plane transverse to the beam direction. An initial spatial anisotropy of particles produced in the collision due to a non-zero impact parameter between the colliding nuclei would result in an anisotropic emission of particles in the transverse plane for a highly interacting system. This emission would be isotropic for a non-interacting system. Agreement between data and hydrodynamical models, which assume local thermal equilibrium, provide evidence for the amount of thermalization that has occurred. Expansion of the plasma is accompanied by a decrease in temperature. The plasma undergoes hadronization when the critical temperature is reached, at which time the partons form
1.3 PARTICLE RATIOS AND $\mu_B$

hadrons. This newly formed hadron gas then reaches chemical freezeout where inelastic collisions no longer occur and the yields of particles are determined. The hadrons continue to interact elastically until the onset of thermal freezeout, at which point the momentum distributions of the particles are determined. The energy density and entropy density of a quark gluon plasma is much higher than for a hadron gas.

1.3 Particle Ratios and $\mu_B$

The two parameters characterizing the phase of matter (as shown in Fig. 1-3) are temperature and baryochemical potential. These parameters tell us how a quark gluon plasma will enter the hadronic phase of matter and what order of transition is expected. The number of baryons expected to be produced in an equilibrated gas is given by Fermi-Dirac statistics. If strangeness is neglected, the relative yields of antibaryons to baryons in thermal equilibrium can be calculated as in Eq. 1.6.

$$\frac{\langle \bar{B} \rangle}{\langle B \rangle} = \frac{\int_0^\infty \frac{p^2 dp}{e^{(\sqrt{p^2+m^2+\mu_B})/T} + 1}}{\int_0^\infty \frac{p^2 dp}{e^{(\sqrt{p^2+m^2-\mu_B})/T} + 1}} \simeq e^{-\mu_B/T} \frac{\int_0^\infty p^2 e^{-\sqrt{p^2+m^2}/T} dp}{\int_0^\infty p^2 e^{-\sqrt{p^2+m^2}/T} dp} = e^{-2\mu_B/T} \quad (1.6)$$

A classical (Boltzmann) approximation is used to get to the last equations. This provides a simple way to estimate the ratio of baryochemical potential to temperature, $\mu_B/T = -0.5 \ln(\langle \bar{B} \rangle/\langle B \rangle)$. If we assume no isospin chemical potential exists (protons and neutrons have equal abundance), protons and antiprotons can be used to calculate $\mu_B/T$. More elaborate calculations which include strangeness can be used to calculate particle ratios in terms of $\mu_B$ and $T$ for collisions where conservation of quantities like strangeness is required. The assumption of equilibrium requires more information to test. Many particle ratios, especially ratios of different particles species, are needed to constrain models and test the applicability of the equilibrium hypothesis.
1.3.1 Stopping and the Baryon Free Region

Nucleons in colliding nuclei lose energy from multiple inelastic collisions. The amount a nucleon slows down as it traverses the other nucleus is called the nuclear stopping power [17]. When discussing the stopping power of nucleus-nucleus collisions, it is useful to define the rapidity variable as

\[ y = \frac{1}{2} \ln \left( \frac{E + p_L}{E - p_L} \right) \]  

where \( E \) is the total energy and \( p_L \) is the longitudinal momentum (momentum in the direction of the beam) of a particle emerging from the collision. Changes in rapidity due to interactions are the same in every frame of reference. Studies of stopping power in proton-nucleus collisions have shown that protons undergoing interactions while traversing a nuclear medium lose 2 units of rapidity on average [17]. The rapidity distribution of baryon number after a heavy ion collision is shown in Fig. 1-5. The top distribution is expected to result when the beam rapidity is much greater than 2. The original constituents of the nuclei (and, therefore, the baryon number) have a distribution of rapidity with peaks centered around \( \pm (|y_{beam}| - 2) \) which is not near midrapidity. Inelastic collisions of the original nucleons pair produce particles near midrapidity. The ratio of baryons to antibaryons at midrapidity is expected to be nearly 1 when \(|y_{beam}| \gg 2\). The bottom distribution shows the expected rapidity distribution of baryon number when the beam rapidity is approximately 2. The particles detected at midrapidity are a sum of baryon transport and pair production. The ratios of antibaryons to baryons is then small.

The ratio of antiprotons to protons can be used as an estimate of the ratio of antibaryons to baryons. At AGS energies (\( \sqrt{S_{NN}} = 4.8 \) GeV) and SPS energies (\( \sqrt{S_{NN}} = 17.2 \) GeV), \( \langle p \rangle/\langle \bar{p} \rangle \) was measured to be 0.0001[18, 19] and 0.12[20], respectively. The majority of baryons at midrapidity were transported at these low energies. At RHIC, beam rapidity is approximately 5. This rapidity is much larger than the SPS beam rapidity of 3. It is, therefore, expected that rapidity distribution of baryon number is similar to the distribution in part a) of Fig. 1-5 at RHIC. Lattice QCD calculations can only be performed at \( \mu_B = 0 \). The predictions of lattice QCD should have increased applicability near midrapidity at
1.4  MICROSCOPIC MODELS OF HEAVY-ION COLLISIONS

Figure 1-5: A sketch of the rapidity distribution of baryon number resulting from two nuclei colliding is shown. When the energy is high (|y_{beam}| ≫ 2) the distribution in part a) results. The nuclei pass through each other leaving energy (and few baryons) in the midrapidity region. The baryon number remains at |y_{beam}|−2 leaving the midrapidity region baryon free. Part b) shows the resulting distribution if the beam rapidity is approximately 2. Many of the original colliding nucleons are transported to midrapidity making the midrapidity region baryon rich.

RHIC if this region is nearly baryon free as expected.

1.4  Microscopic Models of Heavy-Ion Collisions

Since the physics of heavy-ion collisions is difficult to calculate, models describing the interactions of heavy ions are employed. These models incorporate known physics determined from previously-studied data (at lower energies or from proton-proton interactions) and predict the final state observable quantities. The results of these models can be compared to
new measurements to gain understanding of which interactions are needed in the heavy-ion collisions and whether new physics is present.

1.4.1 RQMD

Relativistic Quantum Molecular Dynamics (RQMD[21]) is a hadronic transport model. Nucleons interact individually with other nucleons. Interactions occur by exciting resonances or strings. Strings are representations of soft interactions between the hadrons. The break up of strings causes particle production. RQMD attempts to treat interactions on a level smaller than a hadron, however the main mechanism of the model is to extrapolate the results of interactions from measured cross-sections. For collisions at 1 GeV/nucleon, this model solves transport equations for the nucleons and resonances. Above 10 GeV/nucleon, a parton treatment is used.

1.4.2 HIJING

Heavy Ion Jet Interaction Generator (HIJING[22]) is a partonic transport model. The constituent quarks of nucleons interact rather than hadrons. The main goal of HIJING is to properly model the effect of hard collisions between partons. HIJING uses perturbative QCD to calculate the result of hard interactions in the initial stage of heavy ion collisions for particles with transverse momentum above 2 GeV. Below 2 GeV string excitations are used to calculate the effect of soft interactions. HIJING’s physics is based upon simulating the effect of multiple high-energy nucleon-nucleon interactions and does not incorporate QGP formation or thermalization.

1.5 Thermal Models

Thermal models use statistical relationships and thermodynamics to predict the abundances of particles emerging from a collision. The dynamics of the collision are not considered. An equilibrium between the creation and annihilation rate of particles is assumed. An energy cost, called the chemical potential, exists for the addition of a particle to the system. Each particle species is associated with a different chemical potential. These chemical potentials
are related to the properties of the system: baryon number, charge, strangeness, etc. They characterize the interactions involved in particle creation. The values of these chemical potentials are difficult to calculate theoretically; therefore they are determined by finding the set of values which best describes the measured particle ratios or yields. The application of a thermal model to the description of many yields of particles can test whether the system appears equilibrated. An example of a thermal model (used in this analysis) is described in [23].

1.6 Thesis Goal

The commissioning of the RHIC accelerator opened a new energy regime for heavy ion physics. The purpose of this analysis is to characterize RHIC Au+Au collisions at $\sqrt{s_{NN}} = 130$ and 200 GeV by determining the relative yields of particles to antiparticles for the dominant charged species (pions, kaons, and protons). The baryochemical potential will be obtained from these results given reasonable assumptions about the chemical freezeout temperature. The level at which stopping contributes to the total number of baryons at midrapidity will be estimated to evaluate the extent of the approach to a baryon free region at midrapidity in RHIC collisions. The ratios of particle yields will also be compared to model predictions.
Chapter 2

Experimental Setup

2.1 The RHIC Accelerator Complex

The RHIC accelerator complex at Brookhaven National Laboratory accelerates Au ions to generate collisions to be analyzed by the four experiments on the RHIC collider ring: BRAHMS, PHENIX, PHOBOS, and STAR. Multiple machines are required to produce the fully accelerated Au$^{+79}$ ions for collisions. Fig. 2-1 shows the layout of these machines within the RHIC complex.

2.1.1 Tandem Van de Graaff

The ion source for RHIC is the Tandem Van de Graaff. This machine was completed in 1970 and has the ability to produce and accelerate more than 40 different ion species. Au$^{-1}$ ions are produced through ion sputtering. A 15 MV electrostatic field of the first Van de Graaff machine then accelerates these ions. The Au ions are sent through a thin foil where they lose electrons through collisions making them positively charged. A second van de Graaff machine (in tandem) accelerates these positive ions through a -15 MV electrostatic field. These more energetic ions are sent through another stripping foil. This process produces a Au$^{+32}$ beam with a kinetic energy of 1 MeV/nucleon. The gold ions are then transferred to the AGS Booster through the Tandem to Booster (TtB) transfer line for further acceleration and stripping.
Figure 2-1: A schematic diagram of the RHIC complex is shown. Au ions are produced through sputtering resulting in Au$^{-1}$ ions. A Tandem Van de Graaff accelerates these ions to 1 MeV per nucleon and 2 stripper foils remove 33 electrons resulting in Au$^{+32}$ ions. The Booster accelerates the ions to 95 MeV per nucleon after which the ions are stripped to a charge of +77. The AGS accelerates the ions to 10.8 GeV per nucleon after which the last 2 electrons are stripped from the Au ions. RHIC accelerates these gold ions to a top energy of 100 GeV per nucleon with counter rotating beams.
2.1.2 AGS Booster

The AGS Booster was constructed as an upgrade to the Alternating Gradient Synchotron (discussed in section 2.1.3). The Booster allows the AGS to produce higher intensity proton beams and allows for the use of heavier ion species (like Au) in the AGS. Before construction of the AGS Booster, silicon was the heaviest ion that could be used in AGS beams. At the energies the Tandem Van de Graaff could produce for ion beams, heavier ions still had a significant number of electrons. Beam gas interactions stripped some of these electrons changing the path they would follow in the AGS's magnetic fields. This resulted in beam significant beam loss for these heavy ion species. The AGS Booster has superior vacuum compared to the AGS and supplies more fully stripped higher energy ions to the AGS. After acceleration and a subsequent stripping during the transfer of the ions to the AGS, a Au$^{+77}$ beam exists with kinetic energy of 95 MeV/nucleon.

2.1.3 Alternating Gradient Synchotron

The Alternating Gradient Synchotron (AGS) was constructed in 1960. It was the first accelerator to implement the concept of strong focusing. Accelerators constructed before 1960 used bending magnets with flat pole tips whose field gradients all pointed in one direction. These magnets provided focusing of beam in the vertical direction, but lacked horizontal focusing. Beam loss would occur from particles dispersing horizontally and escaping from the magnet gap. The construction of higher energy accelerators of this type would require magnets with a larger gap (making the whole magnet much bigger). Strong Focusing, whereby both transverse dimensions of the beam are focused, was achieved at the AGS by making the magnet pole tips curved and alternating between magnets whose field gradient faced inward and outward. The curved pole tips provided a better field gradient. Passing the beam through magnets with alternating directions of the field gradients is equivalent to passing light through alternating focusing and defocusing lenses. Both transverse dimensions of the beam are focused. All modern accelerators employ this concept. Strong focusing allowed the magnet gaps in the AGS to be smaller than previous accelerators and thus use less steel and power. Au ions are accelerated to 10.8 GeV/nucleon in the AGS. Beam is then transferred through the AGS to RHIC transfer line where it is stripped once
more. Fully stripped $\text{Au}^{+79}$ ions then enter the RHIC ring.

2.1.4 Relativistic Heavy Ion Collider Ring

Construction of the Relativistic Heavy Ion Collider (RHIC) was finished in 2000. The accelerator uses superconducting magnets to maintain its orbiting beams. It is the first superconducting accelerator which must cross the transition energy while accelerating beam.

Particles circulating in an accelerator have a dispersion in their orbital periods from a synchronous path depending on their dispersion in revolution length and velocity or momentum dispersion. During acceleration asynchronous particles are kept in orbit by changing their momentum in radio frequency cavities. However, at transition, these changes in momentum do not affect the orbital periods of particles. An instability arises and beam loss can occur. Previous accelerators would inject beam above this transition energy so this instability would not occur. This scenario was not feasible at RHIC where transition occurs at $\gamma = 23$. To accelerate past transition at RHIC, the accelerator's optics are changed near the nominal transition energy to increase the value of $\gamma$ where transition occurs. This allows particles to be accelerated stably to the nominal transition energy. The optics are, then, abruptly changed so that the transition energy is below the current energy. The particles can then be accelerated up to the highest possible energies at RHIC.

The highest possible energy of $\text{Au}^{+79}$ beams at RHIC is 100 GeV/nucleon. Au beams with 65 GeV/nucleon were produced in 2000. RHIC's 2001 run produced $\text{Au}^{+79}$ beams at full energy. Six straight sections exist in the accelerator where collisions can occur. Heavy ion detectors exist in four of these interaction regions.

2.2 The PHOBOS Detector

The PHOBOS experiment is designed to measure the multiplicity of particles in nearly $4\pi$ of solid angle and the momentum of particles within a limited solid angle. PHOBOS consists primarily of silicon detectors with the exception of an array of trigger detectors and time of flight walls. Fig. 2-2 shows the components of the PHOBOS detector.
Figure 2-2: The PHOBOS detector is shown with its components labeled. Most of the components are silicon pad detectors. Zero degree calorimeters also exist beyond the Čerenkov detectors and are not shown. The PHOBOS coordinate system is also shown. The $x$-axis is horizontal, the $y$-axis is vertical, and the $z$-axis is oriented in the direction of the beampipe.
2.2.1 Beampipe

The beampipe in the interaction regions is made of Beryllium. Beryllium is a low-Z material used to reduce the probability of interactions occurring in the beampipe before reaching the detectors. The beampipe is made of three pieces, each 3.8 cm in radius, 1 mm thick, and 4 m long. All inactive material is minimized to reduce the probability of interactions. This includes the Beryllium bolts connecting different sections of the beampipe together and thin wire supports for the beampipe.

2.2.2 Magnet

A 2 Tesla quadrupole (double dipole) magnet surrounds the interaction region. The magnet bends the path of charged particles propagating horizontally from the interaction region. This allows the momentum of charged particles to be measured in the spectrometer (see section 2.2.6). The quadrupole magnet has low field intensity near the beampipe and the multiplicity detectors, which minimizes the effect of the magnetic field on the beam orbit and on particles detected in the multiplicity detectors. The magnet can be run in both positive and negative polarities.

2.2.3 Triggering systems

Paddle Counters

The paddle counters are two sets of 16 plastic scintillators located ±3.21 m away from the nominal interaction region along the beampipe. Fig. 2-3 shows a picture of a paddle counter forming an annulus around the beampipe. The pseudorapidity coverage of the paddle counters is $3 < |\eta| < 4.5$, ($\eta = -\ln \tan(\theta/2)$). Their annuli have an inner radius of 7 cm and an outer radius of 25.6 cm. The scintillators are connected to light guides which direct signals into photomultiplier tubes. These photomultiplier tubes are encased in a mu-metal magnetic shield [24].

The paddle counters have high efficiency (98%) for triggering and produce both pulse height (proportional to number of particles) and timing information. The timing resolution is 1 ns for an individual paddle. The counters can also withstand large amounts of radiation.
Figure 2.3: The paddle counters are shown mounted around the beampipe. The scintillator section is the flat front area approaching the beampipe. A light guide directs signals toward the photomultiplier tubes mounted on the outer edge of the support structure.
Zero Degree Calorimeters

The zero degree calorimeters (ZDCs) are instruments installed in all RHIC experiments to provide luminosity monitoring. The calorimeters detect forward moving neutral fragments originating from the collided Au ions. Fig. 2-4 part A shows the position of the ZDCs far away from the interaction region at ±18.6 m. This distance is beyond the bending (DX) magnets which sweep the Au beams back into separate pipes. Part B of Fig. 2-4 shows the position of forward moving neutrons and protons as they exit the interaction region. The DX magnets displace all charged particles away from the zero degree calorimeters leaving only neutrons to contribute to the signal.

The ZDCs are made of plates of tungsten 5 mm thick interspersed with plates of optical fiber[25]. The neutrons interact with the tungsten creating hadronic showers whose particles radiate Čerenkov light. The plates are angled at 45° to correspond with the Čerenkov angle. The optical fiber will only direct light propagating in the direction of its own orientation. The energy resolution of the ZDCs is 17.6% and the timing resolution of the ZDCs is better than 200 ps. The ZDCs are able to withstand large amounts of radiation. Their signals are used both for triggering of experiments and for beam tuning by the accelerator.

Čerenkov Counters

Čerenkov counters are located at ±5.5 m from the interaction region along the beampipe axis. These counters were not fully commissioned until the 2001 run and were not used for the analysis of $\sqrt{s_{NN}} = 130$ GeV data. The purpose of these counters is to provide real-time triggering information about the collision vertex position. The Čerenkov detectors consist of an array of 16 counters forming an annulus around the beampipe. Fig. 2-5 shows one of the two Čerenkov detectors. Each of the radiators are acrylic cylinders 4 cm long with a 1.25 cm radius. A photomultiplier tube is attached to the end of each radiator and is surrounded in mu-metal. The pseudorapidity acceptance of these counters is $4.5 < |\eta| < 4.7$. The counters are located radially 8.57 cm beyond the outer edge of the beampipe. The tubes are mounted in a frame that allows for fine timing adjustments within a 10 cm range along the beampipe direction. The signal to noise of these detectors is 6:1 and the timing resolution is 350 to 400 ps [26]. Position adjustments and delay cables were used to match the counters' peak
Figure 2.4: The location of the zero degree calorimeters is shown. Part A shows the position of the ZDCs relative to the interaction region. The ZDCs are located beyond the DX magnets which direct the two counter circulating beams into and out of the single beampipe for collisions. These magnets are beyond the furthest elements of the PHOBOS detector shown in Fig. 2.2. Part B shows the ZDCs' location between the beampipes for the counter rotating beams. Forward neutrons from an interaction are detected in the ZDCs while protons are bent away from these detectors.
time responses within 50 ps.

To find a vertex, the time difference between the first signal in the positive counter and the first signal in the negative counter is used. A vertex resolution of 4 cm (268 ps) is achieved. This processed vertex signal is available within 650 ns to be used for trigger logic.

Figure 2-5: The Čerenkov detectors are shown mounted around the beampipe. Each of the 16 tubes have its position along the beampipe fine-tuned within its frame to achieve a matching of the tubes' peak time response within 50 ps.
2.2.4 Multiplicity Array

The multiplicity array is a set of silicon detectors which measure the total number of charged particles produced in a Au+Au collision at RHIC. This set of detectors is designed to cover as much of pseudorapidity as possible.

Octagon

An octagonal cylinder of silicon sensors, shown in Fig. 2-6, surrounds the interaction region. The octagon detector is made of 92 sensors. This detector covers a large section of pseudorapidity ($|\eta| < 3.2$). The detector has $2\pi$ acceptance in $\phi$ except where openings exist for the vertex and spectrometer detectors. The octagon has a face to face diameter of 9 cm and a length of 1.1 m. Each sensor has 120 pads ($30 \times 4$). A sensor’s physical size is 3.4 cm $\times$ 8.1 cm. The octagon has relatively large pads (compared to other silicon detectors in PHOBOS) with dimensions 2.7 mm $\times$ 8.7 mm.

![Diagram of the Octagon Detector](image)

Figure 2-6: The octagon, ring and vertex detectors are shown. The octagon, ring, and vertex detectors have an acceptance that cover most of the solid angle with $|\eta| < 5.4$.

Rings

The ring detectors measure multiplicity beyond the octagon’s acceptance. The pseudorapidity acceptance of the ring counters is $3 < |\eta| < 5.4$. The ring counters are a set of three
rings of silicon sensors on each side of the octagon with each ring counter containing 8 ring sensors. The distances of the ring counters along the beam axis from the nominal interaction region is ±1 m, ±2 m, and ±5 m. A ring sensor is not rectangular (whereas all other silicon sensors in PHOBOS are rectangular). The pads of a ring sensor are characterized by Δφ and an inner and outer radius. Δφ is a little less than π/64. The radii of the pads increases from Δr ≈ 0.5 cm to 1.0 cm as the inner radii increase.

2.2.5 Vertex Detectors

The vertex detector consists of 4 layers of silicon sensors (2 layers above the beam pipe and 2 layers below the beampipe) which have high segmentation along the beam axis allowing for a precise determination of the z-component of the interaction vertex. The segmentation in the x-component of the vertex detector is large (1.2 cm for the inner vertex and 2.4 cm in the outer vertex) while the segmentation along the z-component of the vertex is only 0.47 mm. The inner vertex detector has 4 rows of 256 pads while the outer vertex detector has 2 rows of 256 pads. The vertex detector spans ±10 cm from the nominal interaction region. The inner vertex detectors are 5.6 cm above and below the beampipe center while the outer vertex detectors are 11.8 cm above and below the beampipe center. The pseudorapidity coverage of the vertex detectors is |η| < 1.54 for the inner vertex and |η| < 0.92 for the outer vertex. Both sets of vertex detectors cover 42.7° in φ.

2.2.6 Spectrometer

The spectrometer is the main detector used for this analysis. A picture of the spectrometer around the beampipe is shown in Fig. 2-7. The spectrometer has two arms, one on the positive x side of the beampipe and the other on the negative x side of the beampipe. Each arm has silicon sensors mounted on 8 frames. Each frame can have silicon sensors mounted on both sides allowing for 16 layers of silicon. Five different types of spectrometer sensors are used. The number of pads (N_{rows} × N_{columns}) and pad sizes for each sensor type are listed in Table 2.1. Type 1 sensors are mounted in the first two frames of each spectrometer arm. These sensors are located in a low magnetic field region. Their high granularity allows for an accurate determination of the initial trajectory of particles. Type 2 sensors are located
Table 2.1: The number of pads and pixel dimensions for each type of spectrometer silicon sensor are listed.

<table>
<thead>
<tr>
<th>Sensor Name</th>
<th>N&lt;sub&gt;columns&lt;/sub&gt;</th>
<th>N&lt;sub&gt;rows&lt;/sub&gt;</th>
<th>pad width [mm]</th>
<th>pad height [mm]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Type 1</td>
<td>70</td>
<td>22</td>
<td>1.000</td>
<td>1.0</td>
</tr>
<tr>
<td>Type 2</td>
<td>100</td>
<td>5</td>
<td>0.427</td>
<td>6.0</td>
</tr>
<tr>
<td>Type 3</td>
<td>64</td>
<td>8</td>
<td>0.667</td>
<td>7.5</td>
</tr>
<tr>
<td>Type 4</td>
<td>64</td>
<td>4</td>
<td>0.667</td>
<td>15.0</td>
</tr>
<tr>
<td>Type 5</td>
<td>64</td>
<td>4</td>
<td>0.667</td>
<td>19.0</td>
</tr>
</tbody>
</table>

in the next two frames. The pixel width of these sensors is smaller than the pixel width for type 1 sensors. These sensors can resolve deviations from a particle’s initial trajectory due to the increasing magnetic field. Type 4 sensors are primarily located in the fifth and sixth frames and type 5 sensors are located in the seventh and eighth frames. These sensors have pads with a small pixel width for good momentum resolution. The pad heights become increasingly larger near the back planes of the spectrometer to keep the total number of channels low while covering a larger surface area. Type 3 sensors are located in these last four frames of the spectrometer nearest to the beampipe. These sensors are more segmented than type 4 and 5 sensors in the y-direction since particle densities near the beampipe are expected to be high. Each arm has approximately 56000 pads in total.

### 2.2.7 Silicon Sensors

PHOBOS silicon sensors were produced by Miracle Technology Co. in Taiwan under the direction of National Central University (NCU). The sensors are made from a 300 μm thick n-type silicon wafer. Pads are made by doping sections to make p-n junctions[27]. Fig. 2-8 shows the cross-section of a silicon sensor. A sensor pad has a bias bus that applies reverse-bias to the pn-diode making a current flow when the passage charged particles creates electron-hole pairs. A readout line is AC coupled (through an oxide-nitride-oxide capacitive layer) to the p-n junction to extract the signal of a traversing particle. A grounded guard ring exists around the whole active area of the sensor which protects the active area from any stray currents around the edges of the sensor.

The silicon wafers were tested at UIC and MIT. Probing was performed on testkeys and
Figure 2-7: The spectrometer detector is shown with a portion of the beampipe. 8 frames exist in each arm of the spectrometer with the ability to mount sensors on both sides of each frame. The location of different type sensors within the spectrometer is also shown. The two arms possess mirror symmetry across the beampipe except for one set of type 4 sensors mounted in the eighth frame of the +x spectrometer arm.
bonding pads to determine the properties of each sensor. A critical piece of information collected from the testkeys is the full-depletion voltage for the wafer. This was used to determine the operating voltage to be applied during the sensor’s usage. Measurements of the leakage current through the guard and active areas, the capacitance of the ONO layers, the resistance of the polysilicon layer, and a scan for broken readout lines, shorted pads, and shorted capacitive layers were done for quality assurance and determining broken channels. Once the silicon sensors were mounted on a hybrid with a readout chip, the signal to noise of the sensors was determined. Fig. 2-9 shows the results of these tests. The signal to noise of all sensors is above 15.

2.2.8 Time of Flight Detectors

Two time of flight walls exist on the negative x side of the beampipe. Each wall consists of 120 scintillators which are 20 cm tall and are 8 mm × 8 mm in cross-section. Fig. 2-10 shows a time of flight wall. The scintillator has a 1.8 ns decay constant and a 1.6 m attenuation length. Two photomultiplier tubes are connected to each scintillator (on top and bottom). The time difference between the two signals is used to determine the vertical position of the traversing particle. The resulting time resolution is 75 ps. The spatial resolution of the TOF is 10 mm based on time difference and 37 mm based on the ratio of pulse heights in top and bottom photomultiplier tubes.

2.2.9 Raw Data Collection

When a collision occurs, trigger information must be evaluated to determine whether to readout information in all other detectors. The trigger information consists of a small number of channels which must be readout and processed rapidly. The silicon detectors consist of over 100000 channels of information. The large difference in data sizes combined with the speed at which the data must be processed makes data collection for these two sets of information different.

The collection of trigger and TOF data begins by digitizing all data in a FASTBUS crate located in the PHOBOS counting house electronics room and sending this information to custom VME modules which act as level 0 and level 1 triggers. The trigger information
Figure 2-8: The cross-section of a silicon sensor is shown. A bias pad applies reverse-bias to the pn junction. Signals from particles traversing the pn-junction are detected through AC coupled readout lines. The white regions are the bulk silicon and p-n junction. Green and blue regions are metal (aluminum) layers. The red region is a 5 MΩ polysilicon resistor. The yellow layers are oxide-nitride-oxide (ONO) capacitive layers.

Figure 2-9: The signal to noise for all silicon sensors. The values range from 16 to 24.
is ready for analysis in under 1 μs. The VME modules use the FASTBUS information to determine whether the trigger criteria were met. These modules also keep track of whether a previous event is still being processed, in which case a new event should be ignored. If the trigger criteria are met, a signal is issued to retrieve the silicon data and store the digitized trigger information.

Fig. 2-11 shows a schematic of the readout of silicon data. Silicon sensors are mounted on hybrid boards which have one or more readout chips. When trigger criteria have been met, the readout chips receive a hold signal. A set of front-end controllers readout the channels and digitize the signals. A Data Concentrator has Data Multiplexing Units read out the digitized information from the front-end controllers in parallel at 25 MB/s and stores the information in pipelines (FIFOs). A Multiplexor Distributor Controller then directs the
Data Multiplexing Units to transfer their information to a fiber interface which sends the data over two optical fibers at 100 MB/s to a VME Crate in the counting house. The VME crate has a RACEWAY crossbar switch on its backplane for the purpose of distributing sections of data to a set of worker CPUs for processing. Processing of data could include calibration of the raw signals and reduction of the data size by discarding empty channels and compressing the data. Neither of these processes had been implemented in the 2000 or 2001 data runs. The processed data is then sent to an Event Builder board which organizes the data into its storage format and writes this information to a disk array at 40 MB/s. The data is eventually transferred over gigabit ethernet at 30 MB/s to tape in RHIC’s HPSS storage system.

Figure 2-11: The flow of silicon information through its readout chain is shown. Trigger information enters this readout chain in the Data Concentrator.
Chapter 3

Signal Processing

Raw data is retrieved from readout chips connected to the silicon sensors. Properties of the readout electronics must be analyzed to be able to extract a signal representative of the energy deposited in a silicon sensor pad. The signals must also be calibrated to determine the conversion from the electronics output signal to the energy deposition in a pad.

3.1 Pedestal Determination

Leakage currents through the reverse biased silicon sensors cause offsets from zero (pedestals) in the raw silicon readings when no signal exists. These offsets are relatively constant over the length of a run or store (several hours). The pedestal was determined in two ways. The first way was to readout data when no beam was circulating or being injected in the RHIC ring. This data was analyzed by averaging the readings in each channel over many events and recording these values. These pedestals were subtracted from the readings when collisions were occurring.

This method became problematic when the time between stores shortened. A new method was developed to determine pedestals from real data. The first 200 events of a data run had their readings averaged as in the previous method with pedestal runs. To remove signals, a second pass was made over these events averaging only readings whose values deviate by less than ±100 ADC counts from the values determined in the first pass. The average of these readings determined the pedestal. This method automated pedestal
processing and insured that the pedestals are determined as recently as possible for each run. Fig. 3-1 part A shows the pedestal distribution for a small set of channels in the spectrometer. The pedestal values can be as large as 1000 ADC counts on a 4096 channel (12 bit) ADC.

3.2 Common Mode Noise and Random Noise Determination

Fluctuations around the pedestal value for a channel arise from two sources. The first source is random noise from the electronics. This noise determines the resolution with which a signal can be read out. Another source of noise arises from fluctuations in the input voltage to the chip. This type of noise is common to all channels within a chip. Fig. 3-1 part B shows the contribution of both common mode noise and random noise to fluctuations in values around the pedestal. The values is usually less than 5 ADC counts for channels in the spectrometer.

Common mode noise can be corrected for event by event. The correction is determined by averaging the pedestal subtracted readings over all channels within a chip every event and discarding large outliers from the mean (which could be signals). This average is subtracted from all readings removing the common mode noise for that event. Fig. 3-1 part C shows the level of random noise remaining after common mode noise is accounted for. Random noise in spectrometer channels is usually about 2 ADC counts.

3.3 Gain Calibration

The relationship between the energy deposited in a pad and the signal response of the pad's channel must be determined. Calibration runs are taken when no beam is circulating in RHIC. A signal is artificially injected into the readout chip's input and the ADC response is recorded for each channel. The response is very linear until high values (above 3000 ADC counts) is reached. The slope of this curve for each channel is recorded producing a relationship between output ADC response and input signal (called the 'gain'). The input signal comes from a digital to analog converter. The net number of electrons supplied to the readout chip can be calculated as a function of DAC setting. This information combined
Figure 3-1: The contributions of pedestal, common mode noise, and random noise to a signals readout from the silicon sensor are shown. Part A shows the pedestals determined for each channel of this sensor. Part B shows the rms of values around this pedestal when no signal exists. A fluctuation between 3 and 4 ADC counts is seen on average. A shift in this value is seen across the chips. This change is due to a higher common mode noise in the first chip. Part C shows the rms of values around the pedestal when the common mode noise is subtracted. The uncorrelated random noise is about 2 ADC counts.
with the amount of energy required to ionize silicon produces a conversion of 2.1 keV per calibrated ADC count. The gain of each channel was found to be very stable with time.

3.4 Hit Merging

Particles can traverse a pad at multiple angles depositing more energy on average than a normal incidence particle and possibly crossing multiple pads in a single sensor. A single pad may also have the particle traverse only its edge depositing less than 1 MIP of energy. Hit merging groups together pads in a silicon sensor which are likely to have energy deposited by a single particle. Hit merging in the spectrometer is performed only along the bending plane where segmentation is high. Except for the first two layers of spectrometer sensors, pad heights are large enough that different particles could traverse pads adjacent in the vertical plane. Each row is scanned for pads whose energy exceeds 0.15 MIP (12 keV), just above the edge of the noise peak. Up to 8 adjacent pads containing this energy have their signals summed and their coordinates energy weighted to determine the center of where the particle traversed the sensor. If the total energy of a cluster of pads exceeds 0.4 MIP, the hit coordinates and total energy is recorded for further analysis. Combining more than 3 pads into a single hit is rare due to the low occupancy of these sensors.
Chapter 4

Event Characterization

4.1 Event Selection

When a collision occurs, all detectors signals must be readout. The paddle counters are used in PHOBOS to trigger the occurrence of a collision because they have high efficiency and acceptance. A coincidence between signals in both paddle counters produces the initial indication of a possible collision. Fig. 4-1 shows the distribution of the time difference between the most prompt signal in each paddle counter. The time for a single particle moving at the speed of light to hit both paddle counters is 21.4 ns. Collisions between the Au ions and gas in the beampipe, known as beam-gas events, can produce particles. The peaks at 21.4 ns in Fig. 4-1 are caused by these background events. The peak in the middle (centered around 0 ns) is caused primarily by collisions. The time difference between signals in the paddle counters ($|t_{P_{dl}}^N - t_{P_{dl}}^P|$) was required to be less than 4 ns to select collisions close to the nominal interaction region.

Some events within the central peak in Fig. 4-1 can be caused by the random coincidence of two particles from separate beam-gas collisions. Timing information from the ZDCs allow these double beam gas events to be rejected. Triggers resulting from collisions have a characteristic time difference between the primary trigger (from the paddle coincidence) and the arrival of a signal from each ZDC. This time difference ($t_{zdc}^P$ for the positive ZDC and $t_{zdc}^N$ for the negative ZDC) is between 500 and 600 ns (including cable delay time). Events with ZDC signals are required to have these timing characteristics. Extremely central events,
Figure 4-1: The distribution of time differences between the most prompt signal in the positive and negative paddle counters is shown. 21.4 ns is the time required a single particle traveling at the speed of light to traverse the distance between the paddle detectors. The large peaks centered around ±21.4 ns are caused by these non-collision events. The peak around 0 ns is from collisions.
4.2. CENTRALITY DETERMINATION

Table 4.1: The trigger cuts for identifying collisions (minimum-bias trigger) are shown for both 130 GeV and 200 GeV data. A (-) represents a logical AND; a (+) represents a logical OR.

<table>
<thead>
<tr>
<th>Trigger Cut</th>
<th>Value for $\sqrt{s_{NN}} = 130$ GeV</th>
<th>Value for $\sqrt{s_{NN}} = 200$ GeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>PaddleTiming</td>
<td>$</td>
<td>t_{F_{dl}}^{N} - t_{F_{dl}}^{P}</td>
</tr>
<tr>
<td>ZDCN</td>
<td>$t_{zd_{dc}}^{N} &gt; 552$ ns</td>
<td>$t_{zd_{dc}}^{N} &gt; 520$ ns</td>
</tr>
<tr>
<td>ZDCP</td>
<td>$t_{zd_{dc}}^{P} &gt; 560$ ns</td>
<td>$t_{zd_{dc}}^{P} &gt; 520$ ns</td>
</tr>
<tr>
<td>ZDCSum</td>
<td>$(t_{zd_{dc}}^{P} + t_{zd_{dc}}^{N}) &lt; 1240$ ns</td>
<td>$(t_{zd_{dc}}^{P} + t_{zd_{dc}}^{N}) &lt; 1240$ ns</td>
</tr>
</tbody>
</table>

However, may have no ZDC signal (explained in section 4.2). These central events should produce a large paddle signal. The ZDC timing criteria are, therefore, ignored if very high paddle signal is observed. The criteria for identifying a collision from PHOBOS triggers is listed in Table 4.1.

4.2 Centrality Determination

The geometry (impact parameter) of a collision determines how many nucleons interact. Fig. 4-2 shows two nuclei as they approach each other with a nonzero impact parameter. Nucleons in the overlapping region, called participant nucleons, will collide with the nucleons in the other nucleus. Many measurable quantities in a heavy-ion collision, such as the total number of particles produced, scale with the centrality of the collision. The number of participant nucleons and the impact parameter are not directly measurable quantities. A detector signal must be used to estimate the centrality of a collision.

The average of the 12 lowest signals from the 16 signals in a paddle counter (a truncated mean) is used to determine the event centrality. Secondary particles traversing the paddles at high angles cause the energy loss distribution from a single paddle to be asymmetric towards higher losses. Fig. 4-3 shows this simulated paddle counter response in relation to the number of participating nucleons. The monotonic relationship between the paddle counters' signals and the number of participants, combined with the paddle counters' high
Figure 4-2: A geometrical picture of two nucleons colliding with nonzero impact parameter is shown. The overlapping region contains participant nucleons which interact as the two nuclei traverse each other. The non-overlapping region contains spectator nucleons which don't interact and continue propagating in the beam direction.
4.2. CENTRALITY DETERMINATION

efficiency, provide a good estimate for the centrality of a collision.

![Graph showing the relationship between N_{participants} and the Truncated Paddle Mean (a.u.)](image)

Figure 4-3: The number of participants, $N_{part}$, and the paddle signal are shown to correlate monotonically in simulated events.

This response of the paddles is the inverse of the response of the ZDCs. The ZDCs detect spectator neutrons. When many spectators exist, few participants interact, creating a smaller signal in the paddles. Inversely, when many participants exist, few spectators are detected. In very central collisions, no spectators may be detected in the ZDCs. Fig. 4-4 shows the correlation between the ZDC and paddle signals. An inverse relationship exists for sufficiently central collisions. For peripheral collisions, the relationship breaks down. The ZDCs are inefficient for peripheral collisions where spectators consist of large fragments of
the nucleus containing both protons and neutrons. The DX magnets sweep these fragments out of the ZDC acceptance producing a signal representative of fewer spectators.

The distribution of the truncated mean of signals from the paddle counters is shown in Fig. 4-5 for both data and simulation. The simulated distribution includes gaussian smearing of the signal to represent detector noise. The qualitative structure of the distribution is reproduced in the simulation. A large discrepancy exists between the two distributions arising from the trigger inefficiency at low paddle signals. The difference between data and simulation at high centrality represents a small portion of the total cross-section.

The determination of the missing cross-section in data was obtained by simulating the inefficiency of the trigger as a function of the number of paddles hit. Fig. 4-6 shows the distributions of number of paddles hit in simulations as compared to data (data distribution has a trigger requirement that greater than two paddles in each paddle counter must be hit). In simulations, a paddle with more than 0.5 MIP deposited by a particle was considered ‘hit’. The simulated and data distributions are normalized to have the same area between 15 and 22 paddles hit where high efficiency is expected. The ratio of the areas under the full distributions (data divided by simulation) is the net paddle efficiency. For a minimum bias trigger requiring at least one hit in each paddle counter, the efficiency is 97% for both 130 GeV and 200 GeV data. A trigger requiring three or more hits in each paddle counter produces an efficiency of 86% for 130 GeV data and 88% for 200 GeV data. The percentage of total cross-section above a paddle cut is determined by scaling the distribution by the inefficiency and accounting for completely missing events below the minimum requirements for the trigger cut.

Simulated events have the number of participating nucleons recorded. The correlation between paddle signal and number of participants is used to determine $N_{\text{participants}}$ from the simulated distribution. The smearing of paddle signals in simulations produces an error in the determination of number of participants of a few percent (3%). The main source of error in the determination of $N_{\text{part}}$ arises from uncertainty in the missing cross-section for data events and inefficiency of the paddle detectors. The missing cross-section was varied to investigate the effect on the number of participants at a given percentage of the total cross-section. For peripheral events, this has a large effect (up to 20%). For central events,
Figure 4-4: The correlation between ZDC signal and paddle counter signal is shown. A clear inverse relationship is seen except at low paddle signal. These signals are from peripheral collisions where large spectator fragments exist and are swept away from the ZDCs by the DX magnets. The black points are from particles with good ZDC timing representing collisions. The gray points have ZDC timing information representative of double beam-gas events.
Figure 4-5: The distribution of the truncated mean of paddle signals from data and simulation is shown. The two distributions are normalized to have the same area. Simulations reproduce the qualitative features of the data distribution. The discrepancy between the two distributions arises from the inefficiency in triggering on events with low paddle signal.
Figure 4-6: The distribution of number of paddles hit for both data and simulation is shown. The data distribution shown requires at least three hits in each paddle counter. A disagreement exists when few paddles are hit due to paddle inefficiency and the paddle counter trigger cut.
4.3 Vertex Determination

The event vertex is determined separately by multiple subdetectors. A brief description of each algorithm follows. A more detailed description can be found in [28].

4.3.1 SpecMainVertex

The SpecMainVertex relies on the reconstruction of straight tracks in the first 6 planes of the spectrometer. Straight track reconstruction is described in section 5.2. The minimum distance and midpoint between each pair of tracks in an event is computed. The average of the coordinates of these midpoints for track pairs whose minimum distance is less than 5 mm is used as an estimate of the vertex position. This algorithm gives vertex information for all three spatial coordinates of the vertex. It has a limited acceptance along Z axis (nominally between ±20 cm). This vertexing routine is efficient (greater than 95% efficient) for the upper 30% of the cross-section. At lower centrality, the vertex becomes increasingly more inefficient due to the small acceptance of the spectrometer and the requirement that at least two tracks be found.

4.3.2 SpecVertex

The SpecVertex also uses straight tracks from the spectrometer. The point in space which has the minimum total distance from all tracks is used as an estimate of the vertex position. The routine is slightly more inefficient as a function of centrality as compared to the SpecMainVertex, but produces more accurate results.

4.3.3 ZVertex

The ZVertex uses hits in the vertex detectors above and below the beampipe to determine the vertex position. Adjacent hits are joined into clusters and the distribution of pads hit as a function of their Z-position is determined. Clusters in the inner and outer vertex detectors are connected by a straight line and have the Z-component of the point on the
line where Y=0 recorded. The distribution of these Z-components is searched for a peak within a range determined by the histogram of pad density along the Z-axis. This peak is used as an estimate for the Z-position of the vertex. A similar procedure is run for the Y and X components of the vertex. This routine can only determine the Z and Y components of the vertex position accurately. The large segmentation along the X-axis in the vertex detector makes the X-component estimate worse than the spectrometer vertices.

4.3.4 OctMainVertex

The OctMainVertex finds the region in the octagon detector (along the Z-axis) with the maximum density of hits with energy above 0.5 MIP and energy below 2 MIPS. The lower energy threshold exists to remove noise and particles that only traversed a corner of a pad. The high energy threshold removes hits that travel through pads at high angles. The Z distribution of these hits is fit to a gaussian. The sigma of the fit is required to be less than 10 cm for this vertex to be valid. This vertexing routine only produces an estimate of the Z-component of the vertex. Due to the large acceptance of the octagon detector, the vertex can be determined between ±55 cm of the nominal interaction region. This vertex is only used as a cross-check due to its poor resolution (approximately 10 cm), however.

4.3.5 PaddleVertex

The PaddleVertex uses timing information from each paddle counter to determine the vertex position. The position of the vertex (in Z) is estimated as \((c/2) \cdot (t_{P,N}^F - t_{P,d}^F)\), where position is measured from the midpoint between the two detectors. This vertexing routine only gives information about the Z-component of the vertex position, but the vertex always exists (since the paddles trigger the event). The timing resolution of the paddles limits this vertex's resolution to 15 cm. It is used as a cross-check against the other vertices.

4.3.6 RMS Vertex Selector

The RMS vertex selector chooses the best spatial components of the vertex based on event conditions and the reliability of the previously mentioned vertex finders in the spatial range of the vertex. Reliability was determined from simulations of each algorithm. The ZVertex
gives the best Y and Z component estimates within a range of ±20 cm of the nominal interaction region. Outside of this region, SpecVertex becomes the best and, further out, OctMainVertex is the best. The X-component of the vertex can only be determined well by the spectrometer. The SpecVertex is used for this estimate. This vertex selector produces a vertex with resolutions in X, Y, and Z of 250 µm, 200 µm, and 100 µm, respectively.
Chapter 5

Track Reconstruction

Particles emerging from a Au+Au collision in the PHOBOS interaction region leave hits in the spectrometer detector. Fig. 5-1 shows primary tracks and hits in the PHOBOS spectrometer from a simulated Au+Au event at $\sqrt{s_{NN}} = 200$ GeV. A large number of hits from multiple primary and secondary particles is detected. A track reconstruction algorithm must first recognize which hits result from a single track and associate these hits together. Then the properties of the particle (mass, charge, momentum, etc) must be determined from the hit pattern.

5.1 Magnetic Field Configuration

A magnetic field is imposed on the spectrometer detector to bend the trajectory of charged particles. The deflection of a particle in a magnetic field depends on the particle's momentum, $p$, and charge, $q$. In a high energy collision, the majority of particles created have a charge of $-1e$, $0e$, or $1e$. The bending radius of a particle is used to measure the particle's momentum while the direction a particle bends in a magnetic field indicates the charge sign of the particle.

The PHOBOS magnetic field is oriented perpendicular to the spectrometer plane (in the $y$ direction) so that the maximum deflection of a particle in the magnetic field can be traced by the PHOBOS spectrometer. The distribution of the $y$ component of the magnetic field (at $y = 0$) in relation to the spectrometer sensors is shown in Fig. 5-2. The magnitude
Figure 5-1: A simulated Hijing Au+Au event generated from a collision with $\sqrt{s_{NN}} = 200$ GeV is shown. Only primary tracks possessing at least one spectrometer hit are drawn. Hits without tracks drawn through them derive from secondary particles.
of the field strength varies from 0T up to 2.18T[28]. The field strengths in the \( x \) and \( z \) directions are on average 4% of the field strength in the \( y \) direction in the \( y = 0 \) plane. The magnet's configuration minimizes the total field in the collision region. The effect of the magnetic field on the beam orbit is, therefore, minimized. This also allows physics analyses with the Octagon and Vertex detectors to neglect the effect of this field. Since the particle has constant velocity in the \( y \) direction, this component of the particle's trajectory is reconstructed separately from the \( x - z \) components.

Figure 5-2: The \( \hat{y} \) component of the magnetic field in the region of the spectrometer detector at \( y = 0 \) in the PHOBOS coordinate system is shown. The magnet is a double dipole creating a positive field in one spectrometer arm, a negative field in the other spectrometer arm, and a negligible field between the two arms where collisions occur.
5.2 Straight Track Reconstruction

The PHOBOS magnetic field has low intensity near the first six layers of the spectrometer. The field strength is less than 0.35T in these layers. Particles have small deviations from their initial path in these layers. Therefore, the initial trajectories of tracks is approximated by straight lines. A simple analytic algorithm is used to reconstruct these trajectories quickly.

Many tracks leave hits in the spectrometer. The algorithm, therefore, begins by connecting all of these hits in the first two spectrometer layers with the known vertex (previously determined). Segments connected to the first layer’s hits (layer 0) have their $\theta$ and $\phi$ compared to segments connected to the second layer’s hits (layer 1). If their difference in $\theta$ and $\phi$ is less than a limiting value, $\Delta\theta_{\text{max}}$ and $\Delta\phi_{\text{max}}$, respectively, a track seed is made with the vertex and these two hits. These track seeds are then extended through the next four spectrometer layers by creating segments connected to hits in the next layer and recording hits whose segments agree within $\Delta\theta_{\text{max}}$ and $\Delta\phi_{\text{max}}$ with the segment made with the hit in the $i^{th}$ layer ($0 < i < 5$). One missing hit is allowed between layers 2 and 5 by creating segments connected to hits in layer $i + 2$.

Track candidates with 5 or 6 hits in the first six layers of the spectrometer now exist. Straight lines are now fit to the hit coordinates of the track candidates. Least squares (unweighted) linear fits are evaluated separately in the $xz$ plane and the $yz$ plane with $z$ as the independent variable. The $xz$ plane fit only uses the first four layers to minimize the effect of the magnetic field. The $yz$ plane uses all six layers. The calculation of $\chi^2$ is based on the distance in the $xz$ and $yz$ planes to the track weighted by the corresponding pixel dimensions. Since effects like multiple scattering are not incorporated in the fitting procedure, a loose cut keeping tracks with a $\chi^2$ fit probability greater than 0.05% is applied. Tracks which share more than 1 hit are not considered independent. In this case, the track with the highest $\chi^2$ probability is kept.
5.3 Detector Effects

Three detector-induced effects can cause error in the determination of the initial trajectory of a track. The size of the pixels of the detector limits the accuracy of the initial trajectory determination for a track. This effect dominates at high momentum. At low momentum, the scattering of particles as they traverse detector material causes error in the determination of the initial trajectory of a track. The third effect (also larger for low momentum tracks) is the effect of the residual field in the first layers of the spectrometer.

5.3.1 Pixelization

The granularity of the PHOBOS detector limits the accuracy of straight track reconstruction. Pads in the first four layers of the spectrometer have a small pixel height of 1 mm. Hits gathered in these first four layers provide the most information about the $\phi$ angle of the track as it emerges from the vertex. The $\theta$ angle of the track is determined by the initial trajectory in the bending plane ($xz$ plane) of the spectrometer. The pixel width is also 1 mm in pads in the first four layers of the spectrometer and it is decreased to 0.427 mm in the next two layers to accurately determine $\theta$.

5.3.2 Multiple Scattering

Multiple scattering occurs from Coulomb interactions between particles and detector material. The effect is described by the Molière theory of scattering[29]. Particles can scatter into two transverse directions relative to the particle’s propagation direction. The central 98% of the distribution of scattering angles in each of these transverse directions is described by Eq. 5.1. The scattering angle in the transverse direction to propagation is $\theta_p$. The parameter $\theta_0$, determining the width of the Gaussian depends, on particle properties ($\beta$ is the particle velocity divided by the speed of light, $p$ is the particle’s momentum, and $z$ is the particle’s charge) and material properties ($x$ is the thickness of material being traversed and $X_0$ is the radiation length of the material). Multiple scattering has a larger effect on low momentum tracks. This is represented by the widening of the distribution as momentum decreases. As the thickness of the material increases, the distribution also becomes wider.
The probability of a large angle scattering occurring in a single collision makes the tails of the distribution larger than the Gaussian tails. In straight track reconstruction, a $\chi^2$ fit probability greater than 0.05% is required. This low probability cut compensates for neglecting multiple scattering while fitting hits to a straight line. Track fitting, used to evaluate track parameters using all layers of the spectrometer, will account for multiple scattering more directly.

$$\frac{dP}{d\theta_p} = \frac{1}{\sqrt{2\pi}\theta_0} e^{-\frac{q^2}{2\theta_0^2}} \text{ where } \theta_0 = \frac{13.6 \text{ MeV}}{\beta c p} z \sqrt{\frac{x}{X_0}} \left[ 1 + 0.038 \ln \left( \frac{x}{X_0} \right) \right]$$ \hspace{1cm} (5.1)

5.3.3 Effect of the Magnetic Field on Straight Tracks

The magnetic field strength in the region of the first six layers of the spectrometer is non-zero. Fig. 5-3 shows the path of a single track coming from the primary vertex at various momenta as it propagates through the PHOBOS spectrometer. There is a small deflection in the $xz$ plane (momentum dependent) in the region of the first six layers of the spectrometer. The deflection is 2 mm in the sixth layer for a track with a momentum of 200 MeV. As discussed in section 5.2, the impact of this effect on the determination of the initial trajectory is reduced by only using hits in the first four layers while performing straight line fitting in the $xz$ plane. The field strength remains below 0.2T in these first four layers.

5.4 Curved track Reconstruction

The determination of the momentum (and charge sign) of a particle leaving a track in the spectrometer requires the ability to associate the particle's resulting hits as its trajectory curves in the magnetic field. The use of an analytic function for a particle's trajectory is not possible due to the complicated geometric configuration of the PHOBOS magnetic field. Therefore, a general Hough transform method[30] is used to map variables determined from a set of hits to variables that characterize a track. Single simulated tracks are used to obtain this mapping.

Tracks with $0.25 < \theta_0 < 1.75$ and $250 \text{ MeV} < p < 1.2 \text{ GeV}$ are simulated. Each set
5.4. CURVED TRACK RECONSTRUCTION

Figure 5-3: The paths of a charged pion traversing a spectrometer arm at momenta of 0.2, 0.25, 0.333, 0.5, 1.0, and $\infty$ GeV/c are shown (the path bending most has momentum 0.2 GeV/c and the straight path has momentum $\infty$ GeV/c). The subpanels show the paths as they cross each of the first 6 layers of the spectrometer (the lower right subpanel is a legend describing which subpanel corresponds to which spectrometer layer). In the first 4 layers, a track with momentum as low as 0.2 GeV/c is only displaced by a single pixel. In the fifth and sixth layers, a track with a momentum of 0.2 GeV/c can be displaced by as many as 4 pixels. This results from the higher intensity of the magnetic field in the fifth and sixth layers combined with the smaller pixel width in these layers (pixel width = 0.427 mm) as compared to the first 4 layers (pixel width = 1 mm).
of hits in adjacent layers together with the vertex can be converted to an angle from the PHOBOS origin, \( \theta \), and the difference between this angle and the angle from the first hit to the next hit, \( \alpha \). A look up table that allows for the conversion from \((\theta, \alpha)\) to \((\theta_0, p)\) is made. The definitions of these angles where the hit pair comes from layers 14 and 15 is shown in Fig. 5-4.

![Diagram of track reconstruction](Image)

\((\pi^+, p=0.25 \text{ GeV}, \theta_0=0.35 \pi)\)

Figure 5-4: The conversion between track parameters and hit pair parameters used in the Hough tables is shown. The subscript on hit pair parameters is the layer of the first hit. This figure uses hits from layers 14 and 15 as an example.

Once the Hough table has been constructed, hit pairs in different layers (in the high...
field region) created by a specific track must be combined. All sets of hit pairs on the same track should have the same value for \((\theta_0, p)\). The hit pairs are grouped together according to their values of \((\theta_0, p)\) to form a curved trajectory. These curved trajectories are matched to the straight trajectories by comparing their values of \(\theta, \phi\). A single track has energy loss characteristics which depend on the particle's velocity. An additional comparison is made between the energy deposited by hits, \(E_{\text{dep}}\), in the detector in both straight and curved trajectories. Once good candidates for a full track in the spectrometer have been found, a track fit is done to determine the probability of hits belonging to a single track.

5.5 Curve Fitting

The matching of curved and straight elements of a track by their \(\theta, \phi\), and \(E_{\text{dep}}\) produces a set of track candidates likely to come from a real track. A full track fit is then performed to get a \(\chi^2\) estimate of the hit pattern probability conjointly with the best momentum and \(\theta_0\) for the track.

Since the magnetic field has no symmetries, the trajectory of the track is estimated numerically through an iterative procedure. Tracks (pions) are simulated at varying \(p\) and \(\theta_0\) and their deviations from the idealized track (a track without interactions in the detector material) in each spectrometer layer are recorded. The \(xz\)-plane deviations \((\delta x_z)\) and the \(yz\)-plane deviations \((\delta y_z)\) are used to generate a covariance matrix. This matrix is used to calculate the \(\chi^2\) for a candidate track with its \(xz\)-plane and \(yz\)-plane deviations. The \(\chi^2\) formula is shown below.

\[
\chi^2 = \delta x_z^t V_{xz}^{-1} \delta x_z + \delta y_z^t V_{yz}^{-1} \delta y_z
\]

\[
V_{qz} = \begin{pmatrix}
(\delta qz_0 \delta qz_0) & (\delta qz_0 \delta qz_1) & \cdots & (\delta qz_0 \delta qz_n) \\
(\delta qz_1 \delta qz_0) & (\delta qz_1 \delta qz_1) & \cdots & (\delta qz_1 \delta qz_n) \\
\vdots & & \ddots & \vdots \\
(\delta qz_n \delta qz_0) & (\delta qz_n \delta qz_1) & \cdots & (\delta qz_n \delta qz_n)
\end{pmatrix}, \quad q = x \text{ or } y
\]
Figure 5-5: The $\chi^2$ probability distribution from reconstructed Monte Carlo tracks is shown. The probability increases rapidly below 5%, which is the cut used for allowing tracks into the analysis.

The determination of the residuals of a track to its hits requires numerical estimation since the trajectory of a charged particle in the PHOBOS magnetic field does not have a simple functional form. A simulated particle is propagated through the magnetic field until it gets close to the hit. Then separate circle fits in the xz and yz planes are used to approximate the local trajectory. The distance of the track to a hit in this approximation is $d = \sqrt{(x_p - x_0)^2 + (z_p - z_0)^2} - R$. A sign is attached to this residual depending on which side of the trajectory the hit is located in that plane.

The diagonal elements of $V_{xx}$ and $V_{yz}$ estimate the magnitude of displacement of real tracks from the idealized track due to interactions in detector elements. The off diagonal elements estimate correlations between displacements in separate spectrometer layers. Multiple scattering in early detector layers results in a persistent deviation from the ideal
5.5. CURVE FITTING

trajectory.

The iterative method of calculating $\chi^2$ involves finding the best set of parameters, $(1/p, \theta_0, \phi_0, z_0, y_0)$, for the track (the vertex constraint is removed in the final fit and $z_0, y_0$ are defined as the values of $z$ and $y$ when $x = 0$). A simplex minimization routine is used to evaluate the $\chi^2$. The simplex method does not require the calculation of partial derivatives during minimization, thus saving computing time.

The $\chi^2$ probability distribution for tracks in the spectrometer is shown in Fig. 5-5. Ideally, a flat distribution between 0 and 1 is expected. The slight slope in the distribution (at probabilities higher than 35%) is due to the simplified assumptions used in the computation of the covariance matrices. The distribution shows a large number of tracks with a $\chi^2$ probability less than 5%. The probability indicates that the hit matched together are incorrect and are not from a single track. A 5% probability cut is applied in this analysis to remove these incorrect hit combinations (ghost tracks).

Particles lose energy as they traverse detector material. This loss of energy causes the particle to slow down. This results in particles having a smaller momentum in the last planes of the spectrometer as compared to their initial momentum. This effect depends on the mass and momentum of the particle. Deviations from the ideal track are, therefore, biased towards a smaller bending radius. A correction to the found momentum of a particle results. Fig. 5-6 shows the ratio of reconstructed momentum to true momentum for protons as a function of the reconstructed momentum arising from the proton's energy loss characteristics.

5.5.1 Efficiency and Resolution

The efficiency of the tracking algorithm is evaluated by reconstructing simulated tracks. Studies have shown that the difference between reconstructing a single track in an event and a track embedded in a full event (containing multiple tracks) lowers the efficiencies by as much as 14% for high centrality events while causing a negligible change in the resolution of evaluating track parameters[31]. The efficiencies and resolutions quoted here are obtained from single track reconstruction. The efficiency is shown in Fig. 5-7. The maximum efficiency (for single track reconstruction) is about 90% at 500 MeV transverse
Figure 5-6: The ratio of reconstructed momentum to true momentum for protons as a function of reconstructed momentum is shown. The disagreement arises from the different energy loss characteristics between pions and protons. The proton momentum acceptance in data (for the collected statistics in the 130 and 200 GeV runs) extends as low as 300 MeV/c.

The correction to the proton momentum is about 10% at 300 MeV/c and quickly becomes negligible.
Figure 5-7: The single track reconstruction efficiency as a function of $p_T$ for both bending directions is shown. For the particles considered in this analysis, the single track reconstruction efficiency is above 85% except at very low momentum.
momentum and is similar for particles bending toward and away from the beampipe in the magnetic field.

The momentum resolution of tracks with a $\chi^2$ probability greater than 5% is shown in Fig. 5-8. A 1% resolution is achieved at the lowest momentum (250 MeV) and rises roughly linearly to just under 3.5% at 4 GeV. The $\eta$ and $\phi$ resolutions are shown in Fig. 5-9 and Fig. 5-10, respectively. The resolution is less than 1% for values of $\eta$ within the acceptance of the spectrometer. The resolution in $\phi$ is less than 0.01 radians within the acceptance of the spectrometer.

Figure 5-8: The momentum resolution of the tracking algorithm is shown as a function of momentum. A 1% resolution is achieved at 250 MeV momentum and rises approximately linearly to 3.5% at a momentum of 4 GeV.
Figure 5-9: The $\eta$ resolution is shown as a function of $\eta$. It is below 1% within the acceptance of the spectrometer.
Figure 5-10: The $\phi$ resolution is shown as a function of $\phi$. The resolution is less than 0.01 radians within the acceptance of the spectrometer. The resolution shown is for the negative spectrometer arm centered at $\phi = \pi$. 
Chapter 6

Particle Identification

The hits that particles produce in a silicon sensor provide both momentum information (determined from the position of the hit) and energy loss information (determined from the ionization produced by the particle). The different energy loss characteristics of pions, kaons, and protons can be used conjointly with momentum to identify the particle type of a track.

6.1 Energy Loss in Silicon

Heavy particles traversing matter lose energy by ionizing the medium they are passing through. For particles traversing thin materials, where the energy lost through ionization is small compared to the particle’s total energy, the probability of losing an amount of energy between \( E \) and \( E + dE \) is described by the Landau distribution[32].

\[
\frac{dP}{dE} = \frac{\phi\left(\frac{E-E_0}{\sigma}\right)}{\sigma}
\]

\[
\phi(x) = \frac{1}{\pi} \int_0^\infty e^{-u \ln(u) - uz \sin(\pi u)} du
\]

\( E_0 \) and \( \sigma \) depend on the properties of the material being traversed and on the mass and momentum of the particle traversing that medium. The Landau distribution is shown in Fig. 6-1. A small but finite probability exists for the transfer of a large amount of
energy in a single collision. This produces the long tail at high energy losses in the Landau distribution. The assumption that the energy lost is small compared to the total energy of the particle makes the Landau distribution unbounded for high energy losses. The tail of the Landau distribution is so large that the calculated average energy loss is infinite. For very high energy transfers, the Vavilov distribution[33], which introduces a limit on the maximum energy transfer, is a better approximation. The Landau description is sufficient for describing the characteristics of energy loss in the PHOBOS silicon sensors. The value usually used to characterize the Landau (or Vavilov) distribution is the most probable energy loss. For a Landau distribution this value is \( E_{mp} = E_0 + \sigma x_{mp} \), where \( x_{mp} \approx -0.2228 \). For PHOBOS silicon detectors with 300 \( \mu \)m thick sensors, the Landau distribution describes the energy loss of approximately 100 keV per hit from particles with a minimum momentum (to be reconstructed) of 250 MeV/c.

Figure 6-1: A Landau distribution is shown \((E_{mp} = 1 \text{ and } \sigma = 0.065)\). The distribution is asymmetric with only about a 25% probability of losing energy below the most probable energy loss independent of \( E_{mp} \) and \( \sigma \).
6.2 Estimating Energy Loss Parameters

The energy loss of the hits on a track is used to estimate the parameters of the Landau distribution for the particle traversing the silicon sensors. Since a variation exists in sensor thickness and particles can traverse the sensors at varying angles, the first step in estimating energy loss parameters is to determine the actual thickness traversed for each hit and normalize the energy loss to a uniform thickness (300 μm - normal incidence in PHOBOS silicon sensors). In the presence of the magnetic field, the trajectory of the track through the detector changes. A quadratic fit (in the x-z plane) to the hits on a track in the high field region is performed to estimate the angle at which the particle traverses a silicon sensor.

![Graph showing Landau distribution, Mean Distribution, and Truncated Mean Distribution](image)

Figure 6-2: An example Landau distribution is shown. This distribution is sampled 12 times (as occurs for the hits on a track). The “Mean Distribution” is the distribution the average of these 12 values would have. This resulting distribution has the same $\sigma$ as the Landau distribution, but a most probable value that is shifted up by $\sigma \ln 12$. The “Truncated Mean Distribution” is the distribution of the average of the 8 lowest values sampled. This distribution is much narrower than the Landau distribution with a most probable value slightly larger than the Landau distribution’s most probable value.
Ideally, a maximum likelihood method would be used to determine the Landau distribution parameters from the normalized energy loss information. This method is computationally intensive and, therefore, an analytical method is desired. If the energy loss distribution was Gaussian (as would be true for a thick absorber), the most probable energy loss would be given by the average of the energy losses of the \( N \) hits with an error of \( \sigma/\sqrt{N} \). However, the distribution of average energy loss of \( N \) samples from a Landau distribution has the same \( \sigma \) as the original Landau distribution with the most probable value shift up by \( \sigma \ln N \). Therefore, to suppress the effect of the Landau tail, the average of the 7 or 8 hits with the lowest energy losses (a truncated mean) is computed for tracks with 11 or 12 hits, respectively. Fig. 6-2 shows an example Landau distribution and the distributions of the mean energy loss and truncated energy loss for a track with 12 hits. The truncated energy loss distribution is more peaked than the Landau distribution. The long tail of the Landau distribution is suppressed in the truncated energy loss distribution, but still exists in the mean energy loss distribution. The most probable value of the truncated energy loss distribution is higher than the most probable value of the Landau distribution. This deviation is not critical for particle identification, but the smaller width of the truncated energy loss distribution allows for an accurate characterization of the energy loss properties with a limited number of samples.

6.3 Energy Loss Momentum Dependence

The mean energy loss of a particle traversing a material as a function of the material's properties and the particle's momentum is described by the Bethe-Bloch equation[34].

\[
- \frac{dE}{dx} = 2\pi N \sigma^2 c^2 m_e c^2 \rho \frac{Z Z^*}{A \beta^2} \left[ \ln \left( \frac{2m_e c^2 \gamma^2 \beta^2 W_{\text{max}}}{I^2} \right) - 2\beta^2 - \delta - 2 \frac{C}{Z} \right]
\]
where
\[ r_e: \text{ classical electron radius} \]
\[ (2.817 \times 10^{-13} \text{ cm}) \]
\[ m_e: \text{ electron mass (511 keV)} \]
\[ N_a: \text{ Avogadro's number (6.022 \times 10^{23})} \]
\[ I: \text{ mean excitation potential} \]
\[ (173 \text{ eV for silicon}) \]
\[ Z: \text{ atomic number of absorbing material} \]
\[ (14 \text{ for silicon}) \]
\[ A: \text{ atomic weight of absorbing material} \]
\[ \rho: \text{ density of absorbing material} \]
\[ (2.33 \text{ g/cm}^2 \text{ for silicon}) \]
\[ z: \text{ charge of incident particle in units of } e \]
\[ \beta: \text{ velocity of incident particle divided by the speed of light} \]
\[ \gamma: \frac{1}{\sqrt{1 - \beta^2}} \]
\[ \delta: \text{ density correction} \]
\[ C: \text{ shell correction} \]
\[ W_{\text{max}}: \text{ maximum energy transfer in a single collision} \]

The maximum energy transfer allowed for a particle of mass \( M \) traversing a material is

\[
W_{\text{max}} = \frac{2m_e c^2 \beta^2 \gamma^2}{1 + 2 \gamma \frac{m_e}{M} + \frac{m^2_e}{M^2}}
\]

The density correction, \( \delta \), arises from polarization of the medium through which the particle is traveling. This polarization decreases the interactions the particle has with increasingly far electrons, reducing the amount of energy lost in the medium. The effect is momentum (and material) dependent becoming more significant at higher momenta. The shell correction is a modification to the Bethe-Bloch formula at low momentum. It arises from a breakdown in the assumption that the electrons being interacted with are stationary. The effect is very small. Detailed information on these corrections can be found in a textbook by W.R. Leo[34].

To detect energy loss, the electrons ionized in the silicon sensors must remain in the detector medium. Since the PHOBOS spectrometer is made of thin silicon sensors, some of the electrons can escape from the active area of these sensors. The Bethe-Bloch formula is, therefore, modified to neglect energy lost when the energy transfer is above a cut off \( W_{\text{cut}} \). The addition of another parameter into the Bethe-Bloch formula \( W_{\text{upper}} \) equal to the minimum of \( W_{\text{cut}} \) and \( W_{\text{max}} \) accounts for this. This modification flattens out the
relativistic rise of the Bethe-Bloch equation. This restricted energy loss formula[29] is

\[- \frac{dE}{dx} = 2\pi N_a r^2 \alpha_e c^2 \rho Z \beta^2 \left[ \ln \left( \frac{2m_e c^2 \gamma^2 \beta^2 W_{\text{upper}}}{I^2} \right) + \left( 1 + \frac{W_{\text{upper}}}{W_{\text{max}}} \right) \beta^2 - \delta - 2 \frac{C}{Z} \right] \]

A picture of the detected energy (truncated mean of hit energies) with the expected Bethe-Bloch curves overlaid is shown in Fig. 6-3.

Figure 6-3: The energy lost versus momentum for detected particles is shown. The theoretical Bethe-Bloch equations (with all discussed corrections) are overlaid on the data. The Bethe-Bloch equation lines shown have been converted to show the most probable energy loss instead of the mean energy loss.

6.4 Particle Identification cuts

The truncated mean estimate of the most probable energy loss of a particle agrees well with the values predicted by the Bethe-Bloch formula. Deviations arise when describing particles with high energy losses. Also, the width of the energy loss distribution changes
with particle momentum. The construction of particle identification criteria (cut bands) using these $dE/dx$ curves for particles which are clearly a pion, kaon, or proton must, in principle, take into account the changing of the average $dE/dx$ versus momentum and the change in the width of the Landau distributions around that value. These curves are not easy to compute functionally. Since good separation exists between the curves when the energy loss increases at low momentum, the exact functional form of the cut bands is not critical. These bands were, therefore, parameterized with the simple functional form:

$$\frac{-dE}{dx} = \Delta N_{u,l} + N_0 \frac{(1 + \beta_0)^2}{(\beta' + \beta_0)^2}$$

$$\beta' = \frac{p}{\sqrt{p^2 + M^2 + \Delta m_{u,l}^2}}$$

$N_0$ and $\beta_0$ are parameters identical for all particle types. $\Delta N_{u,l}$ are the upper and lower limits around the 1 MIP peak for the truncated energy distribution. $\Delta m_{u,l}^2$ is a particle dependent parameter controlling the change in the envelope around the central band as the energy loss increases towards low momentum. $M$ in this formula is $m_K$ or $m_p$ for kaons and protons, respectively. The energy loss of reconstructed pions does not increase significantly from 1 MIP. The pion cut is, therefore, constructed with two flat lines around the 1 MIP peak. Fig. 6-4 shows these cut bands. The number of particles around the 1 MIP peak is much larger than the number of particles in the accepted kaon or proton regions. An additional cut at 1.6 MIPs is, therefore, applied to limit the number of pions that could be wrongly identified as a kaon or proton. Lastly (for 200 GeV data only) a cut was put at 700 MeV in momentum on the pion band. For 130 GeV data, the pion band extends up to approximately 1 GeV but only accepts pions with anomalously low $dE/dx$ values (below the 1 MIP peak) to remove kaon contamination from the sample of accepted identified pions.

### 6.5 Acceptance

Once cuts for accepted pions, kaons, and protons have been established, the acceptance of identified particles in transverse momentum and rapidity can be determined. Fig. 6-5 shows an outline of the resulting acceptance for each particle type. The contour shown is defined
Figure 6-4: The particle identification bands based on $dE/dx$ and momentum information for accepted pions, kaons, and protons are shown.

to be where the acceptance of particles has fallen below 10% of the maximal number of particles accepted per unit transverse momentum and rapidity. The acceptance at high $p_T$ is controlled by the particle identification cuts. At low $p_T$, the contour is controlled by track reconstruction requirements (specifically, the requirement for the track to have hits in the last six layers of the central portion of the spectrometer detector). The difference between the acceptances for each bending direction is largest for the lightest particles, pions, while these two acceptances are similar for the heaviest particles, protons and antiprotons.
Figure 6-5: The acceptance of each particle type is shown versus transverse momentum and rapidity. The left plot is for particles bending toward the beampipe ($h^+, B^+$ or $h^-, B^-$) and the right plot is for particles bending away from the beampipe ($h^+, B^-$ or $h^-, B^+$).
Chapter 7

Raw Ratios Determination

The relative yields of identified particles can be determined by counting the number of particles within a specified acceptance. The relative detector acceptance for each particle must be taken into account and any other detector effects which determine the relative efficiency of detection must be factored in.

7.1 Event Selection

For this analysis, central events from the top 12% of the cross-section are used. To aid in identifying primary particles a 3.5 mm cut on the distance of closest approach to the primary vertex is imposed. For 130 GeV data the whole vertex range was used for the determination of particle ratios, \(-16 \text{ cm} < v_z < 10 \text{ cm}\). Fig. 7-1 shows the acceptance in \(v_z\) of reconstructable particles bending both directions in the magnetic field. Additional statistics in the 200 GeV data set allowed for a stricter vertex region to be selected. A vertex selection of \(-6 \text{ cm} < v_z < 6 \text{ cm}\) was selected where acceptance exists for both bending directions to reduce any vertex dependent systematic effects.

7.2 Counting

Fig. 7-1 shows four sets of particles. Positive particles collected at positive polarity have the same distribution of particles per event versus \(v_z\) as negative particles collected at negative
Figure 7-1: The distribution of vertices along the beampipe direction (based on tracks) is shown for each magnetic field polarity and track charge sign. Tracks bending the same direction in the magnetic field have the same acceptance.
polarity. Likewise, the distribution is the same for negative particles collected at positive polarity and positive particles collected at negative polarity. A charged particle’s trajectory in a magnetic field has the same trajectory as that particle’s antiparticle in the opposite magnetic field (when the momentum is identical). The acceptance and efficiency of a particle type and its associated antiparticle in each of these data sets are the same. Antiparticle to particle ratios can be taken in each of these data sets with the acceptance and efficiency canceling. Asymmetric interactions of particles versus antiparticles in the detector must be taken into account, but it is not necessary to calculate the absolute yields. This symmetry between particle and antiparticle is also evident in the particles’ momentum and $dE/dx$ distributions. Fig. 7-2 shows these distributions overlaid. The distributions agree within errors.

![Graphs showing momentum and truncated energy deposition distributions for the two acceptances in the PHOBOS spectrometer. The distributions are identical within errors.](image-url)
7.3 Ratio Computation

Ratios are determined for each direction a particle can bend in the PHOBOS spectrometer under the imposition of the magnetic field. This produces two independent measures of the same particle ratios allowing for a cross-check of each value. The cancellation of acceptance and efficiency require symmetry in the conditions with which data is collected. At $\sqrt{s_{NN}} = 200$ GeV, further divisions of the total data set are made to investigate this symmetry. The first additional division separates particles by which arm of the spectrometer they propagated through. During the collection of 130 GeV data only the negative spectrometer arm was fully installed. This cross-check was only possible with the 200 GeV data. This division can reveal if systematics exist due to misalignment of the spectrometer arms. Beam conditions should also be the same during the collection of data. In the 200 GeV data set, a change in beam conditions which caused a shift in the average $y$ coordinate of the interaction vertex was evident. Fig. 7-3 shows the distribution of the $y$ component of the vertex as a function of run number. Around run 7970, an abrupt shift in the $y$ component of the vertex occurred. The data set was divided into data before this run and data after this run. No significant shift was evident in the 130 GeV data set.

Eight independent measurements of a single antiparticle to particle ratio are constructed with the three divisions of the 200 GeV data set. The 130 GeV data set has 2 independent measures from the bending direction of particles. The total numbers of each particle type reconstructed and identified are listed in Table 7.1. Ratios are computed with values in the same division that have the same acceptance. Each value in Table 7.1 is normalized by the number of events in which the particles were produced and the values of $(h^-, B^+)/ (h^+, B^-)$ and $(h^-, B^-)/(h^+, B^+)$ are computed where the superscript on $h$ represents the charge of the hadron while the superscript on $B$ represents the magnetic field polarity. The results are reported in Table 7.2. Two numbers exist in each division for each ratio arising from the use of two different vertexing routines (for 200 GeV ratios). One value is computed from particles reconstructed when a primary vertex was computed from tracks in both spectrometer arms (called a global vertex). The other value is computed from tracks reconstructed when the primary vertex was computed separately for each arm of the spectrometer and
Figure 7-3: The average $y$ component of the vertex is shown as a function of run number. Around run 7970 there was a shift in this position which produced an effect on the ratios if ratios were not calculated separately for runs above and below this occurrence.
that vertex was used to find full tracks within that spectrometer arm (called a local vertex). These values are correlated since they derive from the same data. These values and the statistical errors are averaged to produce a single value for each particle type and division of the data set. The weighted average of these values is taken to determine the final raw ratio with the smallest statistical error. A graph of ratios determined in each data set along with the weighted average is shown in Fig. 7-4. The dark gray band is the final statistical error while the light gray band is the systematic error (discussed later). The antiparticle to particle ratios at 130 GeV were determined by an unweighted average of the values from each bending direction. The final (uncorrected) ratios are shown at the end of this chapter after the method of computing systematic errors is discussed.

Figure 7-4: A comparison of all independent ratios for each particle type at $\sqrt{s_{NN}} = 200$ GeV is shown. The central horizontal line represents the weighted average of the points (weighted only by statistical errors). The dark gray band is the final statistical error on the weighted average. The light gray band is the systematic error on the weighted average.
### 7.3. RATIO COMPUTATION

Table 7.1: The number of accepted events, accepted identified particles, and average transverse momentum \((p_T)\) in MeV/c at each magnetic field polarity is shown. Data is divided by spectrometer arm (only the full \(-x\) arm existed while collecting 130 GeV data) and average \(y\) component of the primary vertex (for 200 GeV data only). \((v_y)_{1}\) includes runs before run 7970 and \((v_y)_{2}\) includes runs after run 7970. The number in parentheses is the number of particles found using the local vertex obtained from tracks only using a single spectrometer arm while the number not in parentheses uses a global vertex obtained from tracks using both spectrometer arms. Only a local vertex was used while analyzing 130 GeV data.

<table>
<thead>
<tr>
<th>Arm</th>
<th>Particle</th>
<th>(N_{\text{particles}} @ B^+)</th>
<th>(\langle p_T \rangle @ B^+)</th>
<th>(N_{\text{particles}} @ B^-)</th>
<th>(\langle p_T \rangle @ B^-)</th>
</tr>
</thead>
<tbody>
<tr>
<td>(\sqrt{s_{NN}} = 130\text{ GeV})</td>
<td></td>
<td>26509 events</td>
<td></td>
<td>41850 events</td>
<td></td>
</tr>
<tr>
<td>(-x)</td>
<td>(\pi^+)</td>
<td>6208</td>
<td>412 ± 1</td>
<td>23223</td>
<td>254 ± 1</td>
</tr>
<tr>
<td></td>
<td>(\pi^-)</td>
<td>14783</td>
<td>253 ± 1</td>
<td>9679</td>
<td>410 ± 1</td>
</tr>
<tr>
<td></td>
<td>(K^+)</td>
<td>136</td>
<td>358 ± 4</td>
<td>256</td>
<td>248 ± 4</td>
</tr>
<tr>
<td></td>
<td>(K^-)</td>
<td>146</td>
<td>255 ± 5</td>
<td>198</td>
<td>366 ± 4</td>
</tr>
<tr>
<td></td>
<td>(p)</td>
<td>223</td>
<td>581 ± 10</td>
<td>331</td>
<td>475 ± 8</td>
</tr>
<tr>
<td></td>
<td>(\bar{p})</td>
<td>109</td>
<td>481 ± 12</td>
<td>218</td>
<td>560 ± 8</td>
</tr>
<tr>
<td>(\sqrt{s_{NN}} = 200\text{ GeV}, (v_y)_{1})</td>
<td></td>
<td>23725 events</td>
<td></td>
<td>20965 events</td>
<td></td>
</tr>
<tr>
<td>(-x)</td>
<td>(\pi^+)</td>
<td>8532 (8286)</td>
<td>418 ± 1</td>
<td>12501 (12935)</td>
<td>289 ± 1</td>
</tr>
<tr>
<td></td>
<td>(\pi^-)</td>
<td>14298 (14692)</td>
<td>288 ± 1</td>
<td>7643 (7472)</td>
<td>416 ± 1</td>
</tr>
<tr>
<td></td>
<td>(K^+)</td>
<td>200 (195)</td>
<td>379 ± 4</td>
<td>241 (255)</td>
<td>259 ± 4</td>
</tr>
<tr>
<td></td>
<td>(K^-)</td>
<td>243 (250)</td>
<td>253 ± 4</td>
<td>208 (198)</td>
<td>373 ± 3</td>
</tr>
<tr>
<td></td>
<td>(p)</td>
<td>458 (472)</td>
<td>580 ± 6</td>
<td>396 (405)</td>
<td>481 ± 6</td>
</tr>
<tr>
<td></td>
<td>(\bar{p})</td>
<td>319 (324)</td>
<td>480 ± 7</td>
<td>293 (291)</td>
<td>596 ± 8</td>
</tr>
<tr>
<td>(+x)</td>
<td>(\pi^+)</td>
<td>7139 (7273)</td>
<td>414 ± 1</td>
<td>13172 (13423)</td>
<td>279 ± 1</td>
</tr>
<tr>
<td></td>
<td>(\pi^-)</td>
<td>14961 (15344)</td>
<td>278 ± 1</td>
<td>6432 (6569)</td>
<td>413 ± 1</td>
</tr>
<tr>
<td></td>
<td>(K^+)</td>
<td>215 (223)</td>
<td>373 ± 3</td>
<td>280 (290)</td>
<td>256 ± 3</td>
</tr>
<tr>
<td></td>
<td>(K^-)</td>
<td>308 (310)</td>
<td>256 ± 3</td>
<td>179 (182)</td>
<td>373 ± 3</td>
</tr>
<tr>
<td></td>
<td>(p)</td>
<td>409 (427)</td>
<td>584 ± 6</td>
<td>424 (433)</td>
<td>485 ± 6</td>
</tr>
<tr>
<td></td>
<td>(\bar{p})</td>
<td>327 (346)</td>
<td>474 ± 7</td>
<td>267 (264)</td>
<td>573 ± 7</td>
</tr>
<tr>
<td>(\sqrt{s_{NN}} = 200\text{ GeV}, (v_y)_{2})</td>
<td></td>
<td>28785 events</td>
<td></td>
<td>13824 events</td>
<td></td>
</tr>
<tr>
<td>(-x)</td>
<td>(\pi^+)</td>
<td>10471 (10248)</td>
<td>417 ± 1</td>
<td>7407 (7759)</td>
<td>290 ± 1</td>
</tr>
<tr>
<td></td>
<td>(\pi^-)</td>
<td>16607 (17267)</td>
<td>291 ± 1</td>
<td>5083 (4990)</td>
<td>417 ± 1</td>
</tr>
<tr>
<td></td>
<td>(K^+)</td>
<td>304 (302)</td>
<td>376 ± 3</td>
<td>137 (154)</td>
<td>257 ± 5</td>
</tr>
<tr>
<td></td>
<td>(K^-)</td>
<td>278 (305)</td>
<td>257 ± 3</td>
<td>135 (124)</td>
<td>373 ± 4</td>
</tr>
<tr>
<td></td>
<td>(p)</td>
<td>612 (606)</td>
<td>577 ± 5</td>
<td>248 (262)</td>
<td>492 ± 8</td>
</tr>
<tr>
<td></td>
<td>(\bar{p})</td>
<td>398 (407)</td>
<td>484 ± 6</td>
<td>195 (193)</td>
<td>573 ± 9</td>
</tr>
<tr>
<td>(+x)</td>
<td>(\pi^+)</td>
<td>8322 (8574)</td>
<td>414 ± 1</td>
<td>8746 (8946)</td>
<td>280 ± 1</td>
</tr>
<tr>
<td></td>
<td>(\pi^-)</td>
<td>18680 (19181)</td>
<td>276 ± 1</td>
<td>4229 (4350)</td>
<td>413 ± 1</td>
</tr>
<tr>
<td></td>
<td>(K^+)</td>
<td>238 (249)</td>
<td>369 ± 3</td>
<td>202 (193)</td>
<td>260 ± 4</td>
</tr>
<tr>
<td></td>
<td>(K^-)</td>
<td>396 (398)</td>
<td>257 ± 3</td>
<td>106 (104)</td>
<td>378 ± 4</td>
</tr>
<tr>
<td></td>
<td>(p)</td>
<td>542 (565)</td>
<td>584 ± 5</td>
<td>289 (298)</td>
<td>479 ± 7</td>
</tr>
<tr>
<td></td>
<td>(\bar{p})</td>
<td>418 (434)</td>
<td>491 ± 6</td>
<td>179 (197)</td>
<td>584 ± 9</td>
</tr>
</tbody>
</table>
Table 7.2: The antiparticle to particle ratios for each division of the data is shown.

<table>
<thead>
<tr>
<th>Arm</th>
<th>Ratio</th>
<th>((h^-/h^+)/B^+)/(h^+/B^-)</th>
<th>((h^-/h^+)/B^-)/(h^+/B^+)</th>
</tr>
</thead>
<tbody>
<tr>
<td>(\sqrt{s_{NN}} = 130) GeV</td>
<td>(\pi^-/\pi^+)</td>
<td>1.005 ± 0.011</td>
<td>0.988 ± 0.016</td>
</tr>
<tr>
<td>(K^-/K^+)</td>
<td>0.90 ± 0.09</td>
<td>0.92 ± 0.10</td>
<td></td>
</tr>
<tr>
<td>(\bar{p}/p)</td>
<td>0.52 ± 0.06</td>
<td>0.62 ± 0.06</td>
<td></td>
</tr>
<tr>
<td>(\sqrt{s_{NN}} = 200) GeV, (v_2)</td>
<td>(\pi^-/\pi^+)</td>
<td>1.011 ± 0.012</td>
<td>1.014 ± 0.016</td>
</tr>
<tr>
<td>(K^-/K^+)</td>
<td>0.89 ± 0.08</td>
<td>1.18 ± 0.12</td>
<td>(1.15 ± 0.12)</td>
</tr>
<tr>
<td>(\bar{p}/p)</td>
<td>0.71 ± 0.05</td>
<td>0.72 ± 0.05</td>
<td>(0.70 ± 0.05)</td>
</tr>
<tr>
<td>(\sqrt{s_{NN}} = 200) GeV, (v_2)</td>
<td>(\pi^-/\pi^+)</td>
<td>1.004 ± 0.012</td>
<td>1.020 ± 0.018</td>
</tr>
<tr>
<td>(K^-/K^+)</td>
<td>0.97 ± 0.08</td>
<td>0.94 ± 0.10</td>
<td>(0.92 ± 0.09)</td>
</tr>
<tr>
<td>(\bar{p}/p)</td>
<td>0.68 ± 0.05</td>
<td>0.74 ± 0.06</td>
<td>(0.70 ± 0.05)</td>
</tr>
</tbody>
</table>

7.4 Systematic Errors

The computation of systematic errors on particle ratios from the 130 GeV and 200 GeV data are quite different. The statistics of the 130 GeV data set was small with few cross-checks available for investigating and resolving systematic deviations. The 200 GeV data has enough statistics to allow for many independent cross-checks with the ability to determine if systematic effects exist.

The systematic error on the uncorrected values of antiparticle to particle ratios at 130 GeV was calculated as half the deviation between the two values for each bending direction. This value is 0.009, 0.01, and 0.05 for \(\pi^-/\pi^+\), \(K^-/K^+\), and \(\bar{p}/p\) respectively. \(K^-/K^+\) has a very large statistical error. A large systematic effect within this ratio would be difficult to resolve. Therefore, the systematic error obtained from the difference in the two values for \(\bar{p}/p\) is also assigned to \(K^-/K^+\).

For the 200 GeV data, the two numbers for each of the eight independent ratios that
7.4. **SYSTEMATIC ERRORS**

correspond to the use of a different vertex finder have half of their deviation assigned as a systematic error. A final systematic error contribution for vertex finding is assigned by computing the standard deviation of all of these individual systematics for each of the three final ratios being sought. This computes to 0.003 for $\pi^-/\pi^+$, 0.018 for $K^-/K^+$, and 0.013 for $\bar{p}/p$.

Eight statistically independent values for each ratio now remain. For each of the three divisions of data, the four values corresponding to the other two division are averaged. The resulting values are examined to check whether their difference is less than their errors added in quadrature. This cross-check is verified for all particle types and data divisions except for the division arising from the shift in the $y$ component of the vertex for $\pi^-/\pi^+$. Here, these two values disagree by 0.030, but are only allowed to disagree by 0.013. A systematic error is assigned as half the difference between these two values (0.015). Since the other cross-checks are verified, a default systematic error is assigned to each ratio for that cross-check as the value computed from $\pi^-/\pi^+$ since this value has the ability to resolve systematic errors most precisely. The systematics assigned to all ratios for deviations based on bending direction, spectrometer arm, and shift in the $y$ component of the vertex are 0.015, 0.001, and 0.001 respectively. These systematic errors are added in quadrature with the systematic errors derived from using different vertex finding routines to arrive at the uncorrected ratios shown below for the 200 GeV data set along with the values determined from the 130 GeV data set.

<table>
<thead>
<tr>
<th>Ratio</th>
<th>$\sqrt{s_{NN}} = 130$ GeV</th>
<th>$\sqrt{s_{NN}} = 200$ GeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\langle \pi^-/\pi^+ \rangle$</td>
<td>0.996 ± 0.010(stat.) ± 0.009(syst.)</td>
<td>1.025 ± 0.006(stat.) ± 0.015(syst.)</td>
</tr>
<tr>
<td>$\langle K^-/K^+ \rangle$</td>
<td>0.91 ± 0.07(stat.) ± 0.05(syst.)</td>
<td>0.95 ± 0.03(stat.) ± 0.02(syst.)</td>
</tr>
<tr>
<td>$\langle \bar{p}/p \rangle$</td>
<td>0.57 ± 0.04(stat.) ± 0.05(syst.)</td>
<td>0.71 ± 0.02(stat.) ± 0.02(syst.)</td>
</tr>
</tbody>
</table>
Chapter 8

Corrections

Interactions of particles with detector material and the particle identification method introduce distortions to the determination of primary particle ratios. Corrections must be applied to the raw ratios obtained in chapter 7 to represent the actual relative yields of particles in Au+Au collisions. This chapter describes these corrections.

8.1 Electron Contamination

The energy lost by electrons and positrons traversing matter is similar to the energy lost by pions except at very low momentum (below 250 MeV). The two particles are indistinguishable within the cut bands developed in chapter 6. Electron and positron contamination to the pion ratio was studied using GEANT[35] simulations on HIJING events. A correction was derived as follows.

\[
\frac{\pi^\pm_{\text{raw}}}{\pi^\pm_{\text{raw}}} = \frac{\pi^\pm + e^\mp}{\pi^\pm + e^\mp} = \frac{\pi^-}{\pi^+} \left( \frac{1 + \frac{e^-}{\pi^-}}{1 + \frac{e^+}{\pi^+}} \right) \Rightarrow \frac{\pi^-_{\text{raw}}}{\pi^+_{\text{raw}}} = \frac{\pi^-_{\text{raw}}}{\pi^+_{\text{raw}}} \left( \frac{1 + \frac{e^+}{\pi^+}}{1 + \frac{e^-}{\pi^-}} \right)
\]

correction = \frac{1 + \frac{e^+}{\pi^+}}{1 + \frac{e^-}{\pi^-}}

where the electrons and positrons must satisfy the same particle identification cuts and track quality cuts as the pions. Contributions from secondary electrons and positrons misidentified as primary pions are reduced by the requirement that they appear to originate
within 3.5 mm of the primary vertex. The relative number of \( e^+ \)'s to \( \pi^+ \)'s or \( e^- \)'s to \( \pi^- \)'s for electrons and positrons that satisfy the pion identification cuts is below 2%. The correction depends on an asymmetry between \( e^+ / \pi^+ \) and \( e^- / \pi^- \) which is small since the number of positrons was found to be approximately equal to the number of electrons in simulated events (similar to the relative number of charged pions). The net correction to \( \langle \pi^- \rangle / \langle \pi^+ \rangle \) was found to be less than 0.2%.

### 8.2 Pion Contamination

Particles in the MIP band of the \( dE/dx \) versus momentum curves (primarily pions) constitute the majority of reconstructed particles. Landau fluctuations toward higher energy losses combined with the much larger number of reconstructed pions relative to kaons or protons can allow a significant number of pions to be misidentified as kaons or protons. The identification bands were constructed to minimize the probability of this misidentification without discarding too many real kaons or protons. The net effect of this misidentification was estimated.

The \( dE/dx \) versus momentum curve was divided into momentum bins 100 MeV wide. Each \( dE/dx \) distribution was fit in the region of its 1 MIP peak to a truncated mean distribution of 12 hits with the top 4 hits truncated as shown in Fig. 8-1. The inclusion of tracks with 11 hits into the fit had a negligible effect on the estimate. The truncation process conserves the property that any Landau distribution, \( L \), can be evaluated from a universal function, \( \phi(x) \), such that \( L(E, E_0, \sigma) = \phi((E - E_0)/\sigma)/\sigma \) (see Appendix A). Accordingly, any truncated mean distribution from \( N + M \) sampled hits from a Landau distribution truncating the top \( M \) hits, \( L^{\text{trunc}}_{N,M} \), can be evaluated from a universal function, \( f_{N,M}(x) \), such that \( L^{\text{trunc}}_{N,M}(E, E_0, \sigma) = f_{N,M}((E - E_0)/\sigma)/\sigma \). The function \( f_{S,A}(x) \) was computed to fit \( L^{\text{trunc}}_{S,A}(E, E_0, \sigma) \) to the MIP peaks in each momentum bin and estimate the number of possible misidentified \( \bar{p} \)'s, \( p \)'s, \( K^- \)'s, and \( K^+ \)'s. A pion contamination correction was derived as follows.

\[
\frac{P}{P^{\text{raw}}} = \frac{P^{\text{raw}} - P^{\text{cont}}}{P^{\text{raw}} - P^{\text{cont}}} = \frac{P^{\text{raw}}}{P^{\text{raw}}} \left( \frac{1 - P^{\text{cont}}}{P^{\text{raw}}} \right) \left( \frac{1 - P^{\text{cont}}}{P^{\text{raw}}} \right)
\]
A conservative estimate of this correction to the proton ratio is less than 0.8%. This correction for the kaon ratio is less than 0.1%.

Figure 8-1: The $dE/dx$ distribution is fit in a single momentum bin ($800 \text{ MeV} < p < 900 \text{ MeV}$). The fit is used to determine the number of pions that could be misidentified as kaons or protons. In this momentum range kaons are indistinguishable from pions.

8.3 Secondary Correction

Secondary protons are generated from interactions in the beampipe and detector material. These can be reconstructed and identified even though they are not created from the Au+Au interaction directly and can cause the number of reconstructed protons to be too large. A
correction to $\bar{p}/p$ for primary particles is derived to account for this effect.

$$\frac{\hat{P}_{\text{primary}}}{P_{\text{primary}}} = \frac{\hat{P}_{\text{primary}}}{P_{\text{found}}} \cdot \frac{P_{\text{found}}}{P_{\text{primary}}} = \frac{\hat{P}_{\text{primary}}}{P_{\text{found}}} \cdot \frac{P_{\text{primary}} + P_{\text{secondary}}}{P_{\text{primary}}} = \frac{\hat{P}_{\text{primary}}}{P_{\text{found}}} \cdot \left(1 + \frac{P_{\text{secondary}}}{P_{\text{primary}}}\right)$$

$$\downarrow$$

correction $= 1 + \frac{P_{\text{secondary}}}{P_{\text{primary}}}$

HIJING events were processed by GEANT and the ratio of findable (reconstructable) secondary protons relative to primary protons was recorded. Most secondaries are produced at low momentum. These secondaries experience high multiple scattering which decreases the likelihood of being reconstructed. The magnetic field deflects these secondary protons more than primary protons, causing them to miss the last layers of the spectrometer and therefore sweeping them out of the acceptance (hits in the last layers of the spectrometer are required for reconstruction). An exponential was fit to the simulated $p_T$ distribution of secondary protons per primary proton. The net correction was computed as the average of the correction function weighted by the $p_T$ distribution of found protons. This correction was 0.8% at 130 GeV and 0.7% at 200 GeV. The estimated systematic error on the correction arising from the Monte Carlo statistics available was 1.0%.

8.4 Feeddown Correction

Hyperon decay (primarily from $\Lambda$ and $\bar{\Lambda}$) to protons and antiprotons causes an error in the determination of $\bar{p}/p$ for primary particles. A correction is made to account for these
8.4. FEEDDOWN CORRECTION

feeddown particles.

\[
\frac{\bar{p}_{\text{primary}}}{p_{\text{primary}}} = \text{correction} \frac{\bar{p}_{\text{primary}} + \bar{p}_{\text{accepted from } \Lambda}}{p_{\text{primary}} + p_{\text{accepted from } \Lambda}}
\]

\[
\downarrow
\]

\[
\text{correction} = \frac{1 + p_{\text{accepted from } \Lambda}}{p_{\text{primary}}}
\]

\[
\downarrow
\]

\[
\text{correction} = \frac{1 + br \cdot \epsilon_{\text{acc}} \frac{\Lambda_{\text{primary}}}{p_{\text{primary}}}}{1 + br \cdot \epsilon_{\text{acc}} \frac{\Lambda_{\text{primary}}}{p_{\text{primary}}}}
\]

The correction depends on the values $\Lambda/p$, $\bar{\Lambda}/\bar{p}$, the branching ratios for $\Lambda \to p\pi^-$ and $\bar{\Lambda} \to \bar{p}\pi^+$ ($br = 0.64$), and an acceptance factor accounting for the probability of a decay proton being identified as primary ($\epsilon_{\text{acc}}$).

$\epsilon_{\text{acc}}$ should only account for the probability of detecting a decay proton relative to a primary proton and should not include detector acceptance. To evaluate this factor, simulated $\Lambda$'s (or $\bar{\Lambda}$'s) are produced and propagated by GEANT where they decay in flight. If the decay produces a proton (or antiproton), information about the number of hits left in the detector and the distance of closest approach to the primary vertex is recorded. To factor out the geometrical acceptance, decay protons (or antiprotons) are propagated through the detector from the primary vertex recording whether the proton would have satisfied the tracking requirements if it had been primary. Fig. 8-2 shows the result of this study as a function of $p_T$. The decay length, $c\tau$, for a $\Lambda$ is 7.9 cm. The first plane of the spectrometer (where a hit is required) is 10 cm from the nominal interaction region. High momentum $\Lambda$'s decay later (on average) than low momentum $\Lambda$'s due to time dilation and have a reduced probability of being reconstructed. Primary particles are required to propagate within 3.5 mm of the primary vertex. Low momentum $\Lambda$'s produce protons which are less likely to satisfy this condition. These competing effects make $\epsilon_{\text{acc}}$ nearly constant as a function of $p_T$. The average of $\epsilon_{\text{acc}}$ over the $p_T$ range is 0.465. A systematic error of 0.050 is assigned to this value to account for the scatter in the $p_T$ range.

The values of $\Lambda/p$ and $\bar{\Lambda}/\bar{p}$ were unknown during the determination of $\bar{p}/p$ at $\sqrt{s} = $
Figure 8-2: The fraction of protons produced by $\Lambda$'s that meet the tracking criterion are shown as a function of $p_T$. The distribution is relatively flat until low transverse momentum is reached.
8.5. ABSORPTION CORRECTION

130 GeV. To reduce the number of unknown quantities the concept of quark coalescence[36] was invoked. Quark Coalescence asserts that the number of particles produced is proportional to the product of the number of constituent quarks. The proportionality is assumed to be equal for a particle and its antiparticle (and divides out when determining the ratios of antiparticles to particles). In the following relation using quark coalescence, a particle’s symbol represents the number of those particles that exist during hadronization.

\[ \frac{\bar{\Lambda}}{\Lambda} = \frac{\bar{u}d\bar{s}}{uds} = \frac{\bar{u}\bar{d}}{u\bar{d}} \cdot \frac{\bar{u}s}{\bar{u}s} = \frac{\bar{p}}{p} \cdot \frac{K^+}{K^-} \]

\[ \downarrow \]

\[ \frac{\Lambda}{p} = \frac{K^-}{K^+} \cdot \frac{\bar{\Lambda}}{\bar{p}} \]

This relation allows only an estimate of \( \Lambda/p \) or \( \bar{\Lambda}/\bar{p} \) to be required for determining the feeddown correction since \( K^-/K^+ \) is a measured value. During the analysis of 130 GeV data, a value for \( \Lambda/p \) of 0.3 \( \pm \) 0.3 was used. Taking into account the errors on all values in the correction formula, the correction was estimated to be 0.992 \( \pm \) 0.020(syst.). While the 200 GeV data was being analyzed, a value of 0.95 \( \pm \) 0.09(stat.) \( \pm \) 0.22(syst.) for the \( \bar{\Lambda}/\bar{p} \) at \( \sqrt{s_{NN}} = 130 \) GeV was published by the PHENIX experiment at RHIC[37]. This value was used to estimate the 200 GeV correction. The correction evaluated to 0.988 \( \pm \) 0.022(syst.).

8.5 Absorption Correction

Antiprotons are preferentially absorbed in the beampipe and detector material as compared to protons. A correction must be applied to the measured ratio to account for this detector effect. This effect was simulated by propagating protons and antiprotons through the detector using GEANT. Hadronic interactions were turned off to determine if a particle’s acceptance would allow it to be reconstructed. The interactions were then turned on to see if the particle would no longer satisfy reconstruction requirements. The non-absorbed fraction was determined as a function of \( p_T \).

This correction was investigated during the 130 GeV analysis using GEANT’s default hadronic interactions package, GHEISHA[38]. The correction to \( \bar{p}/p \) was determined to be 4.9% \( \pm \) 0.5%(syst.) averaged over the \( p_T \) of found protons and antiprotons (the systematic
error is determined from simulated statistics). The same investigation with the 200 GeV data yielded a correction of 5.1% ± 0.5%(syst.). The 200 GeV analysis possessed much higher statistics (and a smaller statistical error) than the 130 GeV analysis. This smaller error necessitated an in-depth study of the systematics of the absorption correction. An alternative package for hadronic interactions (FLUKA[39]) was used to again to simulate this effect. The $p_T$ dependence of the absorbed fraction of protons and antiprotons for both hadronic interactions packages is shown in Fig. 8-3. The two packages are consistent in describing the absorption of protons in PHOBOS detector material. However, a disagreement exists in the absorption of antiprotons in detector material. GHEISHA reports about double the number of absorbed antiprotons as FLUKA. The deviation is due to different methods of determining cross-sections for interactions[40]. The average correction for FLUKA was found to be 2.2% ± 0.5%(syst.).

![Graph showing the $p_T$ dependence of the absorption of protons and antiprotons in detector material for both hadronic packages (FLUKA and GHEISHA) is shown.](image)
8.6. FINAL RATIOS

Analytic calculations were done to estimate this correction from a simple model using the cross-sections for interactions and the actual amount of material traversed. Proton-deuteron and antiproton-deuteron cross-sections were scaled up to estimate the cross-sections for protons and antiprotons on Beryllium, Silicon, and air. The scaling used was

$$\sigma_{(p,\bar{p})A} = \sigma_{(p,\bar{p})d} \left( \frac{A}{2} \right)^{\frac{2}{3}}$$

When the $pd$ and $\bar{p}d$ cross-sections were unknown (at low momentum), the sum of $(pn+pp)$ and $(\bar{p}n+\bar{p}p)$ cross-sections were used. The resultant correction to $\bar{p}/p$ was determined to be 3.7%. This is halfway between the FLUKA and GHEISHA corrections. The final correction used for the 200 GeV analysis was the average of the FLUKA and GHEISHA corrections with the assignment of a large systematic error due to model uncertainty plus the uncertainty of the individual corrections. This value was $3.7\% \pm 1.6\%(\text{syst.})$.

8.6 Final Ratios

The 130 GeV corrections for $\pi^-/\pi^+$ and $K^-/K^+$ had no values which were greater than 1%. $\bar{p}/p$ had three correction factors, $1.008 \pm 0.010(\text{syst.})$ for the secondary correction, $1.049 \pm 0.005(\text{syst.})$ for the absorption correction, and $0.992 \pm 0.020(\text{syst.})$ for the feeddown correction. These factors were multiplied by the raw value of $0.57 \pm 0.04(\text{stat.}) \pm 0.05(\text{syst.})$ (with systematic errors computed appropriately) to arrive at a value of $0.60 \pm 0.04(\text{stat.}) \pm 0.05(\text{syst.})$. An additional systematic error contribution of 0.01 was added linearly to all ratios at 130 GeV as an estimate of a systematic error for all other effects on these ratios. This additional systematic error contribution is larger than the underestimation of the absorption correction found during the 200 GeV analysis. The final values for the 130 GeV particle ratio are shown at the end of this chapter. These values were published in [41].

The 200 GeV corrections for $\pi^-/\pi^+$ and $K^-/K^+$ were also negligible. A correction factor of $1.000 \pm 0.010(\text{syst.})$ was multiplied by these values to account for all other effects on these ratios (this includes the pion contamination correction for $K^-/K^+$). For $\bar{p}/p$ four corrections exist, $1.007 \pm 0.010(\text{syst.})$ for the secondary correction, $1.037 \pm 0.016(\text{syst.})$ for the absorption correction, $0.988 \pm 0.022(\text{syst.})$ for the feeddown correction, and $1.000 \pm$
0.008(syst.) for the pion contamination correction. These factors were multiplied by the raw value of $0.71 \pm 0.02\text{(stat.)} \pm 0.02\text{(syst.)}$ (with systematic errors computed appropriately) to arrive at a value of $0.73 \pm 0.02\text{(stat.)} \pm 0.03\text{(syst.)}$. This produces the final values for the 200 GeV particle ratios shown below. These values were published in [42].

<table>
<thead>
<tr>
<th>Ratio</th>
<th>$\sqrt{s_{NN}} = 130$ GeV</th>
<th>$\sqrt{s_{NN}} = 200$ GeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\langle \pi^- \rangle / \langle \pi^+ \rangle$</td>
<td>$1.00 \pm 0.01\text{(stat.)} \pm 0.02\text{(syst.)}$</td>
<td>$1.025 \pm 0.006\text{(stat.)} \pm 0.018\text{(syst.)}$</td>
</tr>
<tr>
<td>$\langle K^- \rangle / \langle K^+ \rangle$</td>
<td>$0.91 \pm 0.07\text{(stat.)} \pm 0.06\text{(syst.)}$</td>
<td>$0.95 \pm 0.03\text{(stat.)} \pm 0.03\text{(syst.)}$</td>
</tr>
<tr>
<td>$\langle \bar{p} \rangle / \langle p \rangle$</td>
<td>$0.60 \pm 0.04\text{(stat.)} \pm 0.06\text{(syst.)}$</td>
<td>$0.73 \pm 0.02\text{(stat.)} \pm 0.03\text{(syst.)}$</td>
</tr>
</tbody>
</table>
Chapter 9

Results and Model Comparisons

9.1 Final Ratios

The measured relative yields of antiparticles to particles are listed below.

\[
\begin{array}{ccc}
\sqrt{s_{NN}} = 130 \text{ GeV} & \sqrt{s_{NN}} = 200 \text{ GeV} \\
\langle \pi^- \rangle / \langle \pi^+ \rangle & 1.00 \pm 0.01\text{(stat.)} \pm 0.02\text{(syst.)} & 1.025 \pm 0.006\text{(stat.)} \pm 0.018\text{(syst.)} \\
\langle K^- \rangle / \langle K^+ \rangle & 0.91 \pm 0.07\text{(stat.)} \pm 0.06\text{(syst.)} & 0.95 \pm 0.03\text{(stat.)} \pm 0.03\text{(syst.)} \\
\langle \bar{p} \rangle / \langle p \rangle & 0.60 \pm 0.04\text{(stat.)} \pm 0.06\text{(syst.)} & 0.73 \pm 0.02\text{(stat.)} \pm 0.03\text{(syst.)} \\
\end{array}
\]

\(\langle \bar{p} \rangle / \langle p \rangle\) and \(\langle K^- \rangle / \langle K^+ \rangle\) have increased dramatically from lower energies. Fig. 9-1 compares these ratios to values at AGS[43, 44, 45] and SPS[46, 47] energies.

Particles are either transported from the original nucleons or they are pair produced from inelastic collisions. Since the original colliding nuclei contain no antiparticles (and the baryon number is conserved), the number of antibaryons must be equal to the number of baryons resulting from pair production. The majority of baryons created (and the only baryons transported) are protons and neutrons. If we assume protons and neutrons behave identically when recombining from pair produced quarks and stopping, the antiproton to proton ratio can be used as an estimate of the fraction of baryons resulting from pair production in a collision.

\[
\frac{B_{\text{produced}}}{B_{\text{produced}} + B_{\text{transported}}} = \frac{\bar{B}}{B_{\text{produced}} + B_{\text{transported}}} \approx \frac{\bar{p}}{p_{\text{produced}} + p_{\text{transported}}} = \frac{\bar{p}}{p}
\]
The fraction of baryons resulting from pair production near midrapidity at $\sqrt{s_{NN}} = 130$ GeV is 60% and at $\sqrt{s_{NN}} = 200$ GeV is 73%. SPS collisions at $\sqrt{s_{NN}} = 17.2$ GeV measured this fraction to be 12%. For the first time, at RHIC, the midrapidity baryon yield in heavy-ion collisions is not dominated by baryon transport.

Figure 9-1: $\bar{p}/p$ and $K^-/K^+$ are shown as a function of center of mass energy. A large increase in $\bar{p}/p$ is seen between the SPS data and RHIC data suggesting a significant step towards producing collisions nearly free of baryons near midrapidity.

The yields of particles in collisions depend on the abundance of quarks during the hadronization process. The coalescence of quarks has been discussed in section 8.4. A coalescence model postulates that the number of particles of a certain type created is proportional to the product of the number of constituent quarks available. For a baryon, $B_i = a_B q_1 q_2 q_3$, where $q_1$, $q_2$, and $q_3$ are the number of quarks available during hadronization to make the baryon, $B_i$ is the number of baryons of type $i$, and $a_B$ is the probability...
of the baryon forming if the quarks are close enough to coalesce. An analogous equation
holds for mesons, \( M_i = a_{M_i} q_1 \bar{q}_2 \). Since the strong interaction is symmetric for particles and
antiparticles, \( a_{B_i} = a_{\bar{B}_i} \) and \( a_{M_i} = a_{\bar{M}_i} \). This model allows us to relate \( \bar{p}/p \) and \( K^-/K^+ \).

\[
\frac{\bar{p}}{p} = \left( \frac{\bar{q}}{q} \right)^3, \quad \frac{K^-}{K^+} = \left( \frac{\bar{q}}{q} \right) \left( \frac{s}{\bar{s}} \right) \mu_s^{-40} \Rightarrow \frac{\bar{p}}{p} = \left( \frac{K^-}{K^+} \right)^3
\]

\( q \) is the abundance of light quarks (no distinction is made between up and down quarks).
\( s \) is the abundance of strange quarks. If \( \mu_s \) is vanishingly small (which is expected since
the colliding nuclei have no net strangeness), the number of strange antiquarks is equal to
the number of strange quarks. The value of \( K^-/K^+ \) is, then, dominated by the relative
abundance of light antiquarks to quarks. In this limit, \( K^-/K^+ \) should be equal to \( (\bar{p}/p)^{1/3} \).

These values are shown below at both RHIC energies. They agree within errors.

\[
\sqrt{s_{NN}} = 130 \text{ GeV} \quad \sqrt{s_{NN}} = 200 \text{ GeV}
\]

\[
\frac{\langle K^- \rangle}{\langle K^+ \rangle} : \quad \begin{array}{c}
0.91 \pm 0.07 \text{(stat.)} \pm 0.06 \text{(syst.)} \\
0.95 \pm 0.03 \text{(stat.)} \pm 0.03 \text{(syst.)}
\end{array}
\]

\[
\frac{\langle \bar{p}/p \rangle}{(1/3)} : \quad \begin{array}{c}
0.84 \pm 0.02 \text{(stat.)} \pm 0.03 \text{(syst.)} \\
0.90 \pm 0.01 \text{(stat.)} \pm 0.01 \text{(syst.)}
\end{array}
\]

### 9.2 Chemical Potential

A simple estimate of the baryochemical potential can be calculated from particle ratios
as described in section 1.3: \( \mu_B/T = -0.5 \ln(\bar{p}/p) \). This formula uses classical statistics
and provides a baseline expectation for the baryochemical potential. The result of this
calculation for both RHIC energies is shown below. For comparison with later calculations
using a thermal model to estimate \( \mu_B \), a reasonable freezeout temperature of 165 MeV has
been used. \( \mu_B \) decreases by about one third as the center of mass collision energy increases
from 130 to 200 GeV per nucleon.

\[
\begin{array}{c}
\mu_B/T : \quad \begin{array}{c}
\sqrt{s_{NN}} = 130 \text{ GeV} \\
\sqrt{s_{NN}} = 200 \text{ GeV}
\end{array} \\
\mu_B [\text{MeV}] @ T=165 \text{ MeV} : \quad \begin{array}{c}
0.26 \pm 0.03 \text{(stat.)} \\
0.16 \pm 0.01 \text{(stat.)}
\end{array} \\
42 \pm 6 \text{(stat.)} \\
26 \pm 2 \text{(stat.)}
\end{array}
\]
9.3 Thermal Model

A more elaborate thermal model used in [23] includes the electric chemical potential, $\mu_Q$, and the strange chemical potential, $\mu_s$, and uses quantum statistics. The model calculates each particle yield via Eq. 9.1

$$\langle n_i \rangle = \frac{(2J_i + 1)}{(2\pi)^3} \frac{V}{\gamma_s} \int d^3p \frac{1}{\gamma_s e(E_i - \mu_Q Q_i - \mu_s S_i - \mu_B B_i)/T \pm 1}$$

where

- $\langle n_i \rangle$: the yield of hadron $i$
- $J_i$: the spin of hadron $i$
- $V$: the volume in which the hadrons are created
- $\gamma_s$: strangeness suppression factor
- $s_i$: the sum of strange and antistrange quarks in hadron $i$
- $E_i = \sqrt{p^2 + m_i^2}$
- $m_i$: mass of hadron $i$
- $\mu_Q$: electric chemical potential
- $Q_i$: charge of hadron $i$
- $\mu_s$: strange chemical potential
- $S_i$: strangeness of hadron $i$
- $\mu_B$: baryochemical potential
- $B_i$: baryon number of hadron $i$
- $T$: temperature
- $\pm$: + for fermions, - for bosons

The free parameters, determined by fits to measured particle ratios or yields, are $V$, $\gamma_s$, $\mu_Q$, $\mu_s$, $\mu_B$, and $T$. The strangeness suppression factor, $\gamma_s$, accounts for incomplete strangeness equilibration when few strange quarks are produced. This has occurred in low energy collisions where strange particles are too heavy to be created abundantly. Full strangeness equilibration would result in $\gamma_s = 1$.

The original constituents of a heavy-ion collision have no net strangeness, and, therefore, the final state should also have no net strangeness. Often, $\mu_s$ is set equal to zero in thermal
models to account for this. This method, however, allows for systems in the ensemble average to have a distribution of net strangeness around zero. A stricter requirement (applied in [23]) restricts the ensemble to states with zero net strangeness. This exact strangeness conservation method removes $\mu_s$ as a free parameter. This constraint strictly applies to $4\pi$ integrated yields, but can be applied in limited acceptances at energies where strange particles are abundant.

The electric chemical potential accounts for the original colliding constituents having a net charge (+79 for central Au+Au events in each nucleus). The value of this parameter is constrained by setting the net charge to baryon number to that of the original colliding nuclei (0.401 = 79 protons/197 nucleons for central Au+Au collisions).

This model gives a relationship between $\mu_B$, $T$, and both $\bar{p}/p$ and $K^+/K^-$ for $\gamma_s = 1$. Based on a temperature of 159 MeV determined from fits to SPS data, a temperature range of 160 to 170 MeV was considered at RHIC energies. $\mu_B/T$ was found to depend primarily on $\bar{p}/p$ or $K^+/K^-$ with a weak residual dependence on $T$. A value of $T = 165$ MeV was chosen. The result is shown in Fig. 9-2 along with the measured values of $\bar{p}/p$ and $K^+/K^-$ at $\sqrt{s_{NN}} = 130$ and 200 GeV. Using this thermal model, $\mu_B/T$ is $0.27 \pm 0.03$(stat.) and $0.17 \pm 0.01$(stat.) at 130 and 200 GeV, respectively. These values correspond to $\mu_B = 45 \pm 5$ MeV at 130 GeV and $\mu_B = 27 \pm 2$ MeV at 200 GeV when a temperature of 165 MeV is used. These values are quite close to the estimates using classical statistics and neglecting chemical potentials besides $\mu_B$.

Thermal fits in [48] were performed with the inclusion of more measured particle ratios at RHIC at $\sqrt{s_{NN}} = 130$ GeV. A baryochemical potential of 46 MeV and temperature of 174 MeV were found to describe the data well. Measurements of baryochemical potential and temperature in heavy ion collisions from SIS to SPS energies have been shown to be consistent with chemical freezeout occurring when the average energy per particle is between 1.0 and 1.1 GeV[23]. Fig. 9-3 shows this result. Using $\langle E \rangle / \langle N \rangle = 1$ GeV and a phenomenological fit to the energy dependence of the baryochemical potential ($\mu_B = 1.3$ GeV/$\left(1 + \sqrt{s}/4.5$ GeV$)\right)$, the chemical potential and temperature at $\sqrt{s_{NN}} = 200$ GeV were predicted to be $29 \pm 8$ MeV and $177 \pm 7$ MeV, respectively. This produces a prediction of 0.16 for $\mu_B/T$. The $\bar{p}/p$ ratio was predicted to be 0.752 and the $K^-/K^+$ ratio was
Figure 9-2: The dependence of $\langle \bar{p}/(p) \rangle$ and $\langle K^+/K^- \rangle$ as a function of $\mu_B/T$ is shown using the model in [23]. The top plot shows the extracted value of $\mu_B/T$ using the 130 GeV measurement whereas the bottom plot shows this for 200 GeV data. In both cases, $\langle \bar{p}/(p) \rangle$ provides the stronger constraint on $\mu_B/T$. Only statistical errors are used in error bands for the measured values of $\langle \bar{p}/(p) \rangle$ and $\langle K^+/K^- \rangle$. 
predicted to be 0.93. All of these predictions agree with the measured results presented here. Since the 200 GeV Au+Au run at RHIC, many more ratios have been determined at 130 and 200 GeV. Fig. 9-4 shows these ratios along with best values from a thermal fit[49]. A multitude of ratios at RHIC can be described with only two free parameters (T and μB). This provides evidence for the production of a system in RHIC Au+Au collisions consistent with being in equilibrium at chemical freezeout.

9.4 Conclusions

The ratios of charged antiparticles to particles at RHIC energies clearly show that we have entered a new regime of heavy-ion physics where pair production of baryons is the dominant component of the total baryon yield near midrapidity, 60% at \( \sqrt{s_{NN}} = 130 \) GeV and 73% at \( \sqrt{s_{NN}} = 200 \) GeV. The large increase in \( \bar{p}/p \) from lower energies suggests that a significant approach towards the baryon free region at midrapidity has occurred. This impression is based on the assumption that the chemical freezeout temperature has remained relatively unchanged while \( \mu_B \) has decreased significantly from SPS energies due mainly to less transport of baryons to midrapidity. This assumption is confirmed by measurements by the BRAHMS Collaboration[50] which show directly a drop in the midrapidity net proton yield from SPS energies and the movement of most of the net protons to higher rapidities due to the larger beam rapidity at RHIC.

The ratios of particles in Au+Au collisions at RHIC are consistent with chemical equilibrium including complete strangeness equilibration. However, this is not universally accepted as proof of a system which has equilibrated on such short time scales since \( e^+e^- \) collisions can also produce yields consistent with thermal models[51] where, however, complete strangeness equilibration is not achieved. The values extracted for \( \mu_B \) and \( T \) from \( \sqrt{s_{NN}} = 130 \) GeV to \( \sqrt{s_{NN}} = 200 \) GeV show a small increase in temperature along with a large decrease in baryochemical potential, consistent with the phase transition contour in Fig. 1-3 and an average energy per particle at chemical freezeout of between 1.0 and 1.1 GeV.

New possible RHIC runs at 63 GeV per nucleon can provide intermediate data on the
Figure 9-3: Contours in $\mu_B$ and $T$ space corresponding to a fixed average energy per particle of 1.0 and 1.1 GeV as computed from a thermal model[23] are shown. Measurements of $\mu_B$ and $T$ from SIS, AGS, and SPS energies are also plotted. The measurements are consistent with this chemical freezeout condition. The values determined in this analysis are also shown. The error bars on temperature represent the range of temperatures considered in [23] and not an actual measurement of temperature.
9.4. CONCLUSIONS

Figure 9-4: A thermal fit to a multitude of particle ratios measured at both RHIC energies is shown[49]. The thermal model is successful in describing particle ratio at both energies with only 2 free parameters.

The evolution of the freezeout parameters between SPS and full RHIC energies. The LHC should be very close to the baryon free region at about $\mu_B = 1$ MeV if the energy scaling in [48] of the baryochemical potential remains accurate. This will provide an environment where production accounts for the vast majority of particles detected around midrapidity and where lattice QCD predictions will become increasingly valid as a description of heavy-ion collisions.
Appendix A

The Truncated Mean Distribution

A truncated mean distribution is generated by sampling a single distribution \(N + M\) times and averaging the lowest \(N\) values sampled. This distribution arises, in this analysis, from estimating the most probable energy loss of a track in our spectrometer. The original distribution sampled is a Landau distribution. Since a Landau distribution is not symmetric around its most probable value and has a large tail towards higher energy losses, the scatter in the mean energy loss of the tracks remains large. A maximum likelihood estimator of the most probable energy loss is ideal, but requires significant computing time since it must be estimated iteratively. A truncated mean is requires little computing time and produces a good characterization of the tracks energy loss since the truncated mean distribution suppresses the Landau tail. The derivation of the truncated mean distribution where \(N + M\) values are sampled from a single distribution and the lowest \(N\) values are averaged is presented here. The derivation applies to any sample distribution until specific properties of the Landau distribution are invoked.

The \(N+M\) dimensional distribution of values from a sample distribution \(S(E)\) is the product of distributions from each individual sample.

\[
\frac{d^{N+M}P}{dE_1 \cdots dE_N \cdots dE_{N+M}} = \prod_{i=1}^{N+M} \frac{dP}{dE_i} = \prod_{i=1}^{N+M} S(E_i)
\]

To evaluate the effect of truncation, the samples must be ordered. A transformation of variables is, therefore, made which depends on the order of the samples. Since \((N +
$M)!$ permutations exist, $(N + M)!$ transformations mapping \{\(E_1, \ldots, E_N, \ldots, E_{N+M}\)\} to \{\(E_{L_1}, \ldots, E_{L_N}, \ldots, E_{L_{N+M}}\)\} are made where the range of each variable in the new space is constrained such that \(E_{L_1} < E_{L_2} < \cdots < E_{L_N} < \cdots < E_{L_{N+M}}\). However, the truncated values must only be greater than \(E_{L_N}\) and need not be ordered. There are \(M\)! permutations of the truncated values.

\[
\frac{d^{N+M}P}{dE_{L_1} \cdots dE_{L_N} \cdots dE_{L_{N+M}}} = \frac{(N + M)!}{M!} \prod_{i=1}^{N+M} S(E_{L_i})
\]

where \(E_{L_1} < \cdots < E_{L_N}\) and \(E_{L_N} < E_{L_{N+k}}\) for \(1 \leq k \leq M\).

The \(M\) truncated values can be integrated over leaving the \(N\) dimensional distribution of values to be averaged.

\[
\frac{d^{N}P}{dE_{L_1} \cdots dE_{L_N}} = \frac{(N + M)!}{M!} \prod_{i=1}^{N} S(E_{L_i}) \left[ \int_{E_{L_N}}^{\infty} S(\vec{E})d\vec{E} \right]^{M}
\]

where \(E_{L_1} < \cdots < E_{L_N}\).

A change of variables is made such that the average is one of the values (the lowest value) and the order of the variables is preserved.

\[
E = \frac{1}{N} \sum_{j=1}^{N} E_{L_j} \quad \text{and} \quad E_i' = \frac{1}{N} \left\{ E_{L_i} + \sum_{j=2}^{N} E_{L_j} \right\} \quad \text{for} \quad 2 \leq i \leq N
\]

\[
E_{L_1} < E_{L_2} < \cdots < E_{L_N}
\]

\[
\downarrow
\]

\[
E_{L_1} + \sum_{j=2}^{N} E_{L_j} < E_{L_2} + \sum_{j=2}^{N} E_{L_j} < \cdots < E_{L_N} + \sum_{j=2}^{N} E_{L_j}
\]

\[
\downarrow
\]

\[
\frac{1}{N} \sum_{j=1}^{N} E_{L_j} < \frac{1}{N} \left\{ E_{L_2} + \sum_{j=2}^{N} E_{L_j} \right\} < \cdots < \frac{1}{N} \left\{ E_{L_N} + \sum_{j=2}^{N} E_{L_j} \right\}
\]

\[
\downarrow
\]

\[
E < E_2' < \cdots < E_N'
\]
The transformation can be written in matrix form so that it may be inverted.

\[
\begin{bmatrix}
E \\
E_2' \\
\vdots \\
E_N'
\end{bmatrix} = \frac{1}{N}
\begin{bmatrix}
1 & \cdots & \cdots & 1 \\
0 & 2 & \cdots & 1 \\
\vdots & \vdots & \ddots & \vdots \\
0 & 1 & \cdots & 2
\end{bmatrix}
\begin{bmatrix}
E_{L_1} \\
E_{L_2} \\
\vdots \\
E_{L_N}
\end{bmatrix}
\]

In any matrix shown, lines within the matrix mark off a region of a symmetric submatrix. Large bold numbers on off diagonals denote the value of all off diagonal elements. The inverted transformation is shown below.

\[
\begin{bmatrix}
E_{L_1} \\
E_{L_2} \\
\vdots \\
E_{L_N}
\end{bmatrix} = \begin{bmatrix}
N & -1 & \cdots & -1 \\
0 & N-1 & \cdots & -1 \\
\vdots & \vdots & \ddots & \vdots \\
0 & -1 & \cdots & N-1
\end{bmatrix}
\begin{bmatrix}
E \\
E_2' \\
\vdots \\
E_N'
\end{bmatrix}
\]

i.e. $E_{L_1} = N E - \sum_{j=2}^{N} E_j'$ and $E_{L_i} = N E_i' - \sum_{j=2}^{N} E_j'$ for $2 \leq i \leq N$

To transform to the new coordinates, the determinant of Jacobian transformation matrix, $J$, must be evaluated. The Jacobian is the matrix of the inverse transformation above.

\[
|J| = \begin{vmatrix}
N & -1 & \cdots & -1 \\
0 & N-1 & \cdots & -1 \\
\vdots & \vdots & \ddots & \vdots \\
0 & -1 & \cdots & N-1
\end{vmatrix} = N \begin{vmatrix}
N-1 & -1 \\
-1 & N-1
\end{vmatrix}
\]

Cofactor expansion along the first column reduces the determinant to $N$ times a symmetric submatrix. For notational ease below in evaluating the determinant of the submatrix, the N-1 is replaced by $d$ for the submatrix's dimension and the submatrix is then referred to as
D. A matrix, $P$, which diagonalizes $D$, is introduced.

$$
P = \begin{bmatrix}
1 & 1 & \cdots & 1 \\
\vdots & \vdots & \ddots & \vdots \\
0 & \cdots & \cdots & 0 \\
1 & -1 & \cdots & -1
\end{bmatrix} \Rightarrow P^{-1} = \frac{1}{d} \begin{bmatrix}
1 & \cdots & \cdots & 1 \\
2 & -1 & & -1 \\
\vdots & & \ddots & \vdots \\
-1 & & 2 & -1
\end{bmatrix}
$$

and

$$
P^{-1} \begin{bmatrix}
d & -1 \\
\vdots & \vdots \\
-1 & d
\end{bmatrix} P = \begin{bmatrix}
1 & 0 \\
0 & d + 1
\end{bmatrix}
$$

$$
detD = \frac{1}{detP} \cdot detD \cdot detP = detP^{-1} \cdot detD \cdot detP = detP^{-1} DP = (d + 1)^{d-1}
$$

The Jacobian is evaluated by setting $d = N - 1$ and multiplying by the $N$ from the asymmetric portion of the Jacobian matrix, $|J| = N \cdot N^{N-2} = N^{N-1}$. The distribution after the transformation is below.

$$
\frac{d^N P_M}{dE_2 \cdots dE_N} = \frac{(N + M)!}{M!} N^{N-1} \frac{S(NE - \sum_{j=2}^{N} E_j')}{\prod_{i=2}^{N} S(NE_i' - \sum_{j=2}^{N} E_j')} \times \\
\left[ \int_{\sum_{j=2}^{N} E_j'}^{\infty} S(\tilde{E}) d\tilde{E} \right]^M
$$

Integration over all variables besides $E$ produces the truncated mean distribution.

$$
\frac{dP_{N,M}}{dE} = \frac{(N + M)!}{M!} N^{N-1} \int_{E_2'}^{\infty} \cdots \int_{E_N'}^{\infty} \frac{S(NE - \sum_{j=2}^{N} E_j')}{\prod_{i=2}^{N} S(NE_i' - \sum_{j=2}^{N} E_j')} \times \\
\left[ \int_{\sum_{j=2}^{N} E_j'}^{\infty} S(\tilde{E}) d\tilde{E} \right]^M dE_N' \cdots dE_2'
$$

This function is difficult to evaluate, but an important property arises when the Landau distribution is used as the sample function. The Landau distribution is defined by a universal
function $\phi(x)$.

$$L(E, E_0, \sigma) = \frac{1}{\sigma} \phi \left( \frac{E - E_0}{\sigma} \right)$$

This substitution produces the truncated mean distribution specific for the Landau distribution, $L_{N,M}^{\text{trunc}}$.

$$L_{N,M}^{\text{trunc}}(E, E_0, \sigma) = \frac{(N + M)!}{M!} \frac{N^{N-1}}{\sigma^{N+M}} \int_{E}^{\infty} \cdots \int_{E_{N-1}}^{\infty} \phi \left( \frac{NE - \sum_{j=2}^{N} E_j - E_0}{\sigma} \right) \times$$

$$\prod_{i=2}^{N} \phi \left( \frac{NE_i - \sum_{j=2}^{N} E_j - E_0}{\sigma} \right) \left[ \int_{NE_i - \sum_{j=2}^{N} E_j}^{\infty} \phi \left( \frac{\tilde{E} - E_0}{\sigma} \right) d\tilde{E} \right] dx_N \cdots dx_2$$

A transformation of variables, $x_i' = (E_i - E_0)/\sigma$ and $\tilde{x} = (\tilde{E} - E_0)/\sigma$, produces a form of $L_{N,M}^{\text{trunc}}$ where $E$, $E_0$, and $\sigma$ appear together as $(E - E_0)/\sigma$.

$$L_{N,M}^{\text{trunc}}(E, E_0, \sigma) = \frac{(N + M)!}{M!} \frac{N^{N-1}}{\sigma} \int_{x_2}^{\infty} \cdots \int_{x_{N-1}}^{\infty} \phi \left( \frac{NE - E_0}{\sigma} - \sum_{j=2}^{N} x_j' \right) \times$$

$$\prod_{i=2}^{N} \phi(N x_i' - \sum_{j=2}^{N} x_j') \left[ \int_{N x_i' - \sum_{j=2}^{N} x_j'}^{\infty} \phi(\tilde{x}) d\tilde{x} \right] dx_N \cdots dx_2$$

Like the Landau distribution, the truncated mean distribution from a sampled Landau distribution is defined by a universal function.

$$L_{N,M}^{\text{trunc}} = \frac{1}{\sigma} f_{N,M} \left( \frac{E - E_0}{\sigma} \right)$$

where

$$f_{N,M}(x) = \frac{(N + M)!}{M!} N^{N-1} \int_{x}^{\infty} \cdots \int_{x_{N-1}}^{\infty} \phi(N x - \sum_{j=2}^{N} x_j) \prod_{i=2}^{N} \phi(N x_i' - \sum_{j=2}^{N} x_j') \times$$

$$\left[ \int_{N x_i' - \sum_{j=2}^{N} x_j'}^{\infty} \phi(\tilde{x}) d\tilde{x} \right] dx_N \cdots dx_2$$

The functional form of $f_{N,M}(x)$ relative to $\phi(x)$ is exactly the same as the functional form of the truncated mean distribution relative to the Landau distribution. $f_{N,M}(x)$ can be generated by sampling values from $\phi(x)$ and evaluating the truncated mean of those values. $f_{N,M}(x)$ is used to generate a function to fit the pion energy deposition distributions to and
determine the probability of misidentifying a pion as a kaon or proton.
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List of Figures

1-1 QED and QCD vertices .................................................. 11
1-2 Field lines between charged and colored particles .................. 12
1-3 QCD phase diagram ..................................................... 14
1-4 $\epsilon/T^4$ near $T_c$ from lattice QCD ............................. 16
1-5 Rapidity distribution of baryon number after a collision ......... 19

2-1 RHIC complex ............................................................ 24
2-2 PHOBOS detector ....................................................... 27
2-3 Paddle counters ......................................................... 29
2-4 Zero degree calorimeter locations .................................... 31
2-5 Čerenkov detectors ..................................................... 32
2-6 Octagon and ring detectors .......................................... 33
2-7 Spectrometer ............................................................ 36
2-8 Silicon sensor cross-section ......................................... 38
2-9 Signal to noise for silicon sensors .................................. 38
2-10 Time of flight wall .................................................... 39
2-11 Silicon data readout chain .......................................... 40

3-1 Pedestal, common mode noise, and random noise for a small set of channels 43

4-1 Paddle time difference distribution .................................. 46
4-2 Geometrical picture of colliding nuclei .............................. 48
4-3 Correlation between $N_{part}$ and paddle signal .................. 49
<table>
<thead>
<tr>
<th>Figure</th>
<th>Description</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>4-4</td>
<td>Correlation between ZDC signal and paddle signal</td>
<td>51</td>
</tr>
<tr>
<td>4-5</td>
<td>Truncated paddle mean distributions for data and simulation</td>
<td>52</td>
</tr>
<tr>
<td>4-6</td>
<td>Paddle inefficiency</td>
<td>53</td>
</tr>
<tr>
<td>5-1</td>
<td>Hits in the spectrometer from a simulated HIJING Au+Au event</td>
<td>58</td>
</tr>
<tr>
<td>5-2</td>
<td>Magnetic field configuration around spectrometer</td>
<td>59</td>
</tr>
<tr>
<td>5-3</td>
<td>Paths of particles in the spectrometer at varying momenta</td>
<td>63</td>
</tr>
<tr>
<td>5-4</td>
<td>Hough tracking coordinate conversion</td>
<td>64</td>
</tr>
<tr>
<td>5-5</td>
<td>$\chi^2$ probability distribution from reconstructed simulated tracks</td>
<td>66</td>
</tr>
<tr>
<td>5-6</td>
<td>Ratio of reconstructed momentum to true momentum for protons</td>
<td>68</td>
</tr>
<tr>
<td>5-7</td>
<td>Single track reconstruction efficiency</td>
<td>69</td>
</tr>
<tr>
<td>5-8</td>
<td>Momentum resolution</td>
<td>70</td>
</tr>
<tr>
<td>5-9</td>
<td>$\eta$ resolution</td>
<td>71</td>
</tr>
<tr>
<td>5-10</td>
<td>$\phi$ resolution</td>
<td>72</td>
</tr>
<tr>
<td>6-1</td>
<td>Landau distribution (general properties)</td>
<td>74</td>
</tr>
<tr>
<td>6-2</td>
<td>A Landau distribution, mean distribution, and truncated mean distribution</td>
<td>75</td>
</tr>
<tr>
<td>6-3</td>
<td>$dE/dx$ vs momentum for reconstructed particles with Bethe-Bloch curves</td>
<td>78</td>
</tr>
<tr>
<td>6-4</td>
<td>Particle identification cut bands</td>
<td>80</td>
</tr>
<tr>
<td>6-5</td>
<td>Acceptance for each particle type</td>
<td>81</td>
</tr>
<tr>
<td>7-1</td>
<td>Vertex distribution along beampipe direction based on tracks</td>
<td>84</td>
</tr>
<tr>
<td>7-2</td>
<td>Comparison of momentum and energy deposition distributions</td>
<td>85</td>
</tr>
<tr>
<td>7-3</td>
<td>Average $y$ component of the vertex vs run number</td>
<td>87</td>
</tr>
<tr>
<td>7-4</td>
<td>Comparison of independent ratio measurements for each particle type</td>
<td>88</td>
</tr>
<tr>
<td>8-1</td>
<td>Pion contamination fit</td>
<td>95</td>
</tr>
<tr>
<td>8-2</td>
<td>Fraction of feeddown protons which are reconstructed vs $p_T$</td>
<td>98</td>
</tr>
<tr>
<td>8-3</td>
<td>$p_T$ dependence of absorption for protons and antiprotons</td>
<td>100</td>
</tr>
<tr>
<td>9-1</td>
<td>Energy dependence of $\langle \bar{p} \rangle/\langle p \rangle$ and $\langle K^- \rangle/\langle K^+ \rangle$</td>
<td>104</td>
</tr>
<tr>
<td>9-2</td>
<td>Dependence of $\langle \bar{p} \rangle/\langle p \rangle$ and $\langle K^+ \rangle/\langle K^- \rangle$ on $\mu_B/T$</td>
<td>108</td>
</tr>
<tr>
<td>9-3</td>
<td>Chemical freezeout contours</td>
<td>110</td>
</tr>
</tbody>
</table>
9.4 Results of a thermal fit to many particle ratios ................. 111
List of Tables

2.1 Spectrometer sensor pad dimensions ........................................ 35

4.1 Minimum-bias trigger cuts ...................................................... 47

7.1 Information about statistics used in this analysis ......................... 89

7.2 Antiparticle to particle ratios for each division of the data set .......... 90