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Global Variables of the $\Lambda$ Hyperon and $\pi^-$ Meson Production in Central $^{197}$Au $+$ $^{197}$Au Collisions at AGS

by

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A Thesis Submitted
in Partial Fulfillment of the Requirements for the Degree
Doctor of Philosophy

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Houston, Texas
April, 1997
Abstract

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S.V. Efremov

According to Quantum Chromodynamics Theory, highly compressed and heated nuclear matter should undergo a phase transition to a new state, called the Quark Gluon Plasma. These extreme conditions can probably be created in the collisions of relativistic heavy ions. We are looking for signatures of the phase transition to QGP such as enhancement of the strangeness production rates and reduction of the collective nuclear flow. The results of the invariant differential cross sections measurements for $\Lambda$ hyperon and $\pi^-$ meson production via track reconstruction in the Time Projection Chambers experiment BNL-AGS-E891 are presented. Comparison of our experimental data on the differential particle multiplicities for $\Lambda$ and $\pi^-$ with the cascade models ARC and RQMD is given. The shapes of the particles spectra indicate the strong collective expansion of nuclear matter. We employ the three dimensional expansion picture in order to parametrize the observed particle spectra and extract
the flow parameters together with the mean temperature of the particle gas after the nuclear fluid freeze-out.
Acknowledgments

I would like to take this opportunity to thank all the members of BNL-AGS-E891 collaboration, without which this thesis would be impossible. Special thanks are due to my advisor Gordon Mutchler, BNL-E891 spokesperson Edward Platner and Director of Bonner Nuclear Lab Billy Bonner. Thanks are also due to Ron Longacre (BNL), Al Saulys (BNL), Heinz Sorge (SUNY at Stony Brook), Bill Llope (Rice University), David Kahana (BNL) and Stratos Eftathiadis (University of Boston).
Contents

Abstract

Acknowledgments

List of Illustrations

List of Tables

1 Introduction 1

1.1 The Conventional Nuclear Physics 4

1.2 Heavy-Ion Physics, Search for the Fifth State of Matter 8

1.2.1 Traditional Interpretations of Heavy-Ion Experiments 10

1.2.2 Space-Time Evolution of Heavy-Ion Collisions 12

1.3 QCD Transition, Phase Diagram 18

1.4 Signals of the Quark-Gluon Plasma 20

1.4.1 Heavy Quarkonia 24

1.4.2 Dileptons 24

1.4.3 Direct Photons 25

1.4.4 Disoriented Chiral Condensate 26

1.4.5 Open Charm 27

1.4.6 Hanbury-Brown Twiss 27

1.4.7 Collective Nuclear Flow 29
1.4.8 Strangeness Enhancement ........................................... 33
1.5 Some Results of Current Heavy-Ion Experiments ................. 35
  1.5.1 Experiment BNL-E866 ........................................... 37
  1.5.2 Experiment BNL-E877 ........................................... 38
  1.5.3 Our Measurements and Analysis: Experiment BNL-E891 ... 43

2 Experimental Apparatus .................................................. 49
  2.1 The Experimental Layout ......................................... 50
  2.2 The MPS Magnet .................................................. 53
  2.3 The TPC Modules ................................................ 53
  2.4 The Electronics .................................................. 57
  2.5 The Gas System .................................................. 61
  2.6 The Au Beam ...................................................... 62
  2.7 The Trigger ....................................................... 62

3 Data Reduction and Analysis .......................................... 68
  3.1 Raw Data Analysis ................................................ 68
    3.1.1 Track Reconstruction Algorithm .................................. 69
    3.1.2 Detector Calibrations ......................................... 70
    3.1.3 \( V_0 \) Selection Criteria .................................... 81
    3.1.4 Production of the Data Summary Tapes ......................... 85
  3.2 Data Reduction .................................................... 86
3.2.1 Selection of Tracks and Momentum Resolution .......................... 88
3.2.2 Selection of Vertices ..................................................... 90
3.2.3 Selection of Central Events .............................................. 94
3.2.4 Monte Carlo Simulations of the Apparatus Effects ..................... 98
3.2.5 Acceptance Correction and Error Analysis ............................. 100
3.2.6 Search for Systematic Errors .......................................... 102

4 Results and Discussions .......................................................... 117

4.1 Introduction ................................................................. 117
4.2 $\Lambda$ Production in Central Au+Au Collisions ......................... 118
  4.2.1 Rapidity Dependence of Inverse Slopes for $\Lambda$ Production .... 118
  4.2.2 Global Fireball Parametrizations of $\Lambda$ Spectrum .............. 121
  4.2.3 Collective Effects and Flow Model Parametrization .............. 122
  4.2.4 Rapidity Distribution of $\Lambda$ Hyperons ....................... 134
4.3 $\pi^-$ Meson Production in Central Au+Au Collisions ................ 136
  4.3.1 Double Differential Spectra of $\pi^-$ Production ............... 138
  4.3.2 Rapidity Distributions of $\pi^-$ .................................. 141
  4.3.3 Mean Multiplicity of $\pi^-$ ...................................... 145
4.4 Comparison with Microscopic Cascade Models .......................... 145
  4.4.1 Cascade Model Calculations ...................................... 146
  4.4.2 Double Differential Spectra of Negative Pions .................. 147
4.4.3 Rapidity Distribution of Negative Pions .................................. 147
4.4.4 Mean Multiplicity of Negative Pions ..................................... 149
4.4.5 Double Differential Spectra of Lambdas ................................. 149
4.4.6 Rapidity Distribution of Lambdas ......................................... 150
4.4.7 Mean Multiplicity of Lambdas .............................................. 150
4.4.8 Conclusions ................................................................. 151

5 Summary and Conclusions .......................................................... 155

A BNL-AGS-E891 Collaboration ...................................................... 159

B The Vertex "N-tuple" .................................................................. 160

C The Track "N-tuple" ................................................................... 161

D Cascade Models of Nucleus–Nucleus Collisions .......................... 162
  D.1 RQMD ......................................................... 162
  D.2 ARC ......................................................... 163

E Occupation Probabilities of the Nucleon Resonances ................. 164

Bibliography ................................................................................. 166
Illustrations

1.1 Space-time evolution of hadronic matter in the one-dimensional hydrodynamic model. ................................. 19

1.2 Schematic QCD phase transition diagram as a function of $T$ and $\mu_B$. 21

1.3 Schematic QCD phase transition diagram as a function of $T$ and $n_B/n_B^0$. 22

1.4 Several key observables of the quark-gluon plasma. ................. 23

1.5 Macroscopic quantities of the collision system with and without the phase transition to QGP. ................................. 32

1.6 Proposed mechanism of strangeness enhancement for high-density quark matter. ........................................... 34

1.7 Transverse mass invariant spectra of protons in central $Au + Au$ interactions measured by BNL-E866 collaboration. ................. 39

1.8 Transverse mass invariant spectra of $\pi^+$ and $\pi^-$ measured by BNL-E866 collaboration. ................................. 40

1.9 Rapidity density distributions for $\pi^-$, $\pi^+$, $K^+$ and $K^-$, in central $Au + Au$ collisions, measured by BNL-E866 collaboration. ....... 41

1.10 Proton transverse mass spectra measured by BNL-E877 collaboration. 44

1.11 Pion transverse mass spectra measured by BNL-E877 collaboration. ... 45
1.12 Rapidity distributions for $\pi^+$, $\pi^-$ and protons measured by BNL-E877 collaboration. 46

1.13 Spectra of fragments measured by BNL-E877 collaboration. 47

2.1 The E891 plan view. 51
2.2 E891 Beam. 52
2.3 The MPS Magnet. 54
2.4 Front view of module showing principle components. 56
2.5 View of three modules in operating position as viewed from upstream end of the MPS magnet. 58
2.6 View of the main electrical connections to modules. 59
2.7 Scheme of individual electronics channel. 64
2.8 Schematic arrangement of electrodes in the detector end cap. 65
2.9 The Layout of a 128-channel hybrid printed circuit boards (one master and seven slaves). 66
2.10 AGS plan view. 67

3.1 Distortion corrections. 77
3.2 Drift Velocity Calibrations Results. 79
3.3 Offset Calibration Results. 80
3.4 The hit efficiency as a function of $x$ coordinate for the first row of the first TPC module.

3.5 The hit efficiency as a function of $y$ coordinate for the first row of the first TPC module.

3.6 Multiprocessing Scheme via Event Dispatcher.

3.7 Illustration to equation (??) for the saggita.

3.8 Illustration for the "fish tail" background combination.

3.9 The Azimuthal Angle Dependence of $M_{\pi^-p}$ Distribution.

3.10 The grid in the phase space used for acceptance calculations.

3.11 The proton’s direction dependence of $M_{\pi^-p}$ distribution for data.

3.12 The proton’s direction dependence of $M_{\pi^-p}$ distribution for MC.

3.13 $M_{\pi^-p}$ effective mass distribution integrated over azimuthal orientation angle of the decay plane $\Phi_{\pi p}$, 14,114 Data events.

3.14 $M_{\pi^-p}$ effective mass distribution integrated over azimuthal orientation angle of the decay plane $\Phi_{\pi p}$, 25,000 Monte-Carlo events.

3.15 $M_{\pi^-p}$ effective mass distribution, final sample. Normalization per one central event done in both data and Monte-Carlo.

3.16 Positive track reconstruction efficiency as a function of $x$ and $y$ coordinates of the track intercept with the front plane of the TPC.

3.17 Data to Monte-Carlo ratios of the $\pi^-$ counts as a function of the azimuthal angle $\Phi^-$ for all bins of rapidity used.
4.1 Results of local exponential and Boltzmann–like fits of the invariant
\Lambda spectrum. ......................................................... 120

4.2 Double differential spectra for the inclusive \Lambda production.
Comparison to the cascade models predictions and to the Flow Model
Parametrization. ......................................................... 123

4.3 Spectra of protons and pions normal to the beam in the c.m.s for
central \textit{Ne} + \textit{NaF} interactions at 800 \( A \) \( MeV \) in the lab, explained
by the \textit{blast wave} by P. Siemens and J. Rasmussen. ................. 127

4.4 Rapidity distributions for central 14.6 \( A GeV/c Si + Al \) collisions in
comparison to isotropic thermal distribution at \( T = 0.12 GeV \). .... 130

4.5 E891 rapidity distribution of \Lambda production. .................. 135

4.6 E891 rapidity distribution of the \( \pi^- \) production in \( Au + Au \).
Comparison with the Flow Model parametrization. ................... 142

4.7 E891 double differential spectra for \( \pi^- \) production in \( Au + Au \).
Comparison with the Flow Model parametrization. ................... 143

4.8 E891 double differential spectra for \( \pi^- \) production in \( Au + Au \).
Comparison with cascade models predictions. ...................... 148
Tables

2.1 Positions and sizes of the detectors in E891. Positions are measured relatively to the pivot point. .................................................. 50

4.1 Results of global fits of \( \Lambda \) spectrum. ........................................ 124

4.2 Results of the Flow Model parametrization of \( \Lambda \) and \( \pi^- \) spectra. ... 136

4.3 Rapidity distribution of \( \Lambda \) hyperons. ........................................... 137

4.4 Rapidity distribution of \( \pi^- \) mesons. ........................................... 144

E.1 Occupation probabilities of excited states of the nucleon and the relative densities in a thermal system at a temperature of 150 \( MeV \). 165
Chapter 1

Introduction

During the current century the fundamental interactions of matter were classified into the electromagnetic, strong, weak and gravitational, each acting on different space-time scale and involving different kinds of matter properties. Matter, in turn, was found to be subdivided into elementary constituents carrying various quantum numbers, which predetermine the way these constituents interact with each other. In all cases the interaction of different quanta literally means the change of their wave functions, annihilation or production of other quanta.

It was realized that electromagnetic interaction – interaction between the particles having electric charge – occurs through the massless and electrically neutral carrier – virtual gamma quanta. In space the electromagnetic force has an infinite range according to the Coulomb’s law $\sim 1/r^2$, which is due to the zero mass of its carrier. The method of describing the electromagnetic interaction developed into the quantum field theory, which thereafter has been generalized to include the weak interactions into a unified picture the electroweak interaction theory.

When physicists considered the structure of the atomic nucleus, they confronted a mystery. It was known that the nucleus is composed of nucleons. But how were
the large number of nucleons, protons and neutrons confined in a tiny volume? The electric repulsion between the protons is so great that the nucleus should explode.

There is only one solution to the problem: there must exist another basic force that keeps the nucleus together. This force is called the strong interaction force.

Unlike the electromagnetic interaction, observed in everyday life, the strong interaction is hidden. The strong interaction achieved its title on the basis of rate of changes produced by it. The scattering of particles (so called hadrons), governed by the strong interaction, sometimes accompanied by production of other particles, turns out to be the fastest among all known fundamental interactions – of the order of $10^{-23}$ seconds. The reason that the strong interaction is hidden is its short range, which is about $10^{-13}$ cm, or 1 fm (one Fermi), using conventional particle physics unit of length.

It is the small interaction radius and the physical analogy for the cause of the interaction that lead Hideki Yukawa to the idea that the strong interaction is caused by exchange of a virtual particle which he called a meson. Unlike the photon, the carrier of the strong interaction is required to have a mass of about $m_0 \sim 150 MeV/c^2$ in order to explain the limited radius of the strong interaction potential $V(r) = g^2/r \exp(-r/(\hbar/m_0c))$. Nowadays the description of the strong interaction in terms of exchange by various mesons is considered to be the effective way to parametrize the strong amplitudes, however it is believed that the fundamental theory of strong interaction is QCD, where the quark and gluon degrees of freedom are considered.
It is important for understanding of the strong interaction to create conditions when essentially fundamental constituents of matter, quarks and gluons determine the dynamics of the process. The quark-gluon plasma phase of matter [1, 22], a state consisting of weakly interacting deconfined quarks and gluons, was predicted to exist soon after the discovery [2] of the asymptotic freedom property of Quantum Chromodynamics (QCD). That novel phase of matter was predicted to occur perhaps deep in the cores of neutron stars [3] and during the first few moments (microseconds) of the Big Bang [4]. It was also suggested that perhaps the quark-gluon plasma could be formed in the laboratory in very high energy collisions of hadrons or nuclei [21]. Since the mid 1970's the experimental search for new phases of nuclear matter [5, 6] using heavy-ion reactions has been underway. Until 1984 only light ion reactions at incident momenta below 2 $A GeV/c$ [7] could be studied. In the past ten years, experiments with heavy-ion with energies up to 200 $A GeV$ have been investigated [8, 9, 10]. In 1999 a new era of experiments with heavy-ion beams in the collider mode at the Relativistic Heavy-Ion Collider (RHIC) at the Brookhaven National Laboratory (BNL) with center of mass energies up to 100 $A GeV$ will begin.

In order to experimentally study the phase transition of the quark-gluon plasma, we need to achieve very high energy density in the lab. It is also important to have a large interaction volume in order to approach the thermal equilibrium of the new phase. It is hoped that high energy heavy-ion collisions can provide us with these necessary conditions. Since the initial interaction volume and parton multiplicity turn
out to be large, the dense parton system may go through both thermalization and
chemical equilibration processes after its formation. At extremely high energies one
may study an almost fully-equilibrated plasma of quarks and gluons with zero baryon
density. At currently accessible energies, matter with high baryon density may be
studied as well. The QCD phase transition is also expected to occur at high baryon
density [22].

1.1 The Conventional Nuclear Physics

The goal of the nuclear physics is the study of nuclear structure and mechanisms
of nuclear reactions by comparing theoretical model predictions with the results of
experimental measurements. These two directions – structure and reaction studies
– contribute to each others progress. The first one is intended to determine the
laws governing the essentially quantum mechanical motion of the nucleons inside the
stable or excited nuclei, while the second explores the non-nucleonic and, further, non-
hadronic degrees of freedom. The validity of conventional nuclear physics nowadays
– the physics involving the hadronic degrees of freedom, is being challenged by the
precise measurements on the smaller space and time scale, when the importance of
the more fundamental matter constituents – quarks and gluons comes in effect.

An important tool to study both, nuclear structure and the fundamental hadronic
interactions is the scattering of various projectiles off nuclei at high values of mo-
momentum transfer, allowing one to look for the energy-momentum distribution of the
intranuclear nucleons and for collective correlations [11]. The results of limited types of experiments can be well described in terms of conventional nuclear physics without the need to incorporate the QCD degrees of freedom for its description. For example, extensive study of the nuclear reactions at intermediate energy region (from several hundred $MeV$ to a few $GeV$ beam energy) is consistent with the scenario of sequential interaction of the beam particle and its products with the intranuclear nucleons and the semiclassical propagation of the particles in the momentum-dependent mean-field potential between successive collisions.

Depending on the number of degrees of freedom and the typical momentum transfer involved, the scattering processes are classified into direct and the collective, soft and hard, respectively. While direct processes, involving small number of degrees of freedom play important role at low projectile energy and/or with one-particle projectiles, increasing energy over the thresholds for various particles production drowns the direct process under a background of abundantly occurring multi-step production and rescattering. Analytical models mainly intended for the description of the direct or a few step processes in hadron nucleus collisions or scattering of electromagnetic projectiles off the nuclei mostly known are the Impulse Approximation [16], First Collision Model [13], Folding Model [14] (for multi-step inclusive production) and the Glauber Model [15]. The models of other type, imitating the whole evolution of either one-particle distribution function $f(\vec{r},t)$ or the many-body Wigner distribution
function \( f(\vec{r}_1, \ldots, \vec{r}_N, t) \) are called the transport models and are numerically realized in the form of Monte-Carlo simulations.

The microscopic models based on empirical information for elementary processes occurring in the free space, such as elastic and inelastic scattering cross sections or exclusive cross sections for the particle production, in the majority of cases satisfactorily describe experimental data on the inclusive and exclusive differential cross sections for the particle production. Discrepancies can be attributed to the incomplete knowledge of a whole set of cross sections of the elementary processes involved and to in-medium modifications of the latter. While the lack of knowledge about some of the cross sections can be compensated via systematic theoretical description, based on meson exchange and the resonance models [17], many resonance widths, branching ratios and meson-baryon coupling constants have not been measured yet or have been measured with insufficient accuracy.

In the early stages of nuclear structure studies the empirical success was achieved by the shell-model, based on the mean-field concept of the independent motion of intranuclear nucleons in the self-consistent potential created by the rest of the nucleons, which can be described by the Hartree-Fock approach, originally applied in the atomic physics. The occupation numbers of each energy level – shell – within the accuracy of the experiment could be well explained in terms of the maximum occupation numbers determined by the angular momentum of the shell and the Pauli exclusion principle. However, some of the properties of the nuclei (the mass difference
of odd and even nuclei, existence of energy gap in the spectra of excited odd-odd nu-
clei and the absence of such a gap in the spectra of odd and odd-even nuclei, densities
of one-particle levels, etc) could not be described by the mean field approaches.

The possibility of probing shorter distances inside nuclei shed the light on higher
momentum and energy components of the nuclear wave function. The measurement
of the spectroscopic factors together with the missing mass spectrum in inclusive and
exclusive electron scattering experiments, revealed that only 80% of all nucleons in
the C\textsuperscript{12} nucleus occupy the shell-model levels, while the remaining 20% belong to the
high-momentum tail and have higher removal energy.

The mechanism of the high-energy-momentum depletion of the mean field levels
in He\textsuperscript{3} and in the ideal case of the infinite nuclear matter, observed in (e, e'p) ex-
periments was successfully described without the need to include partonic degrees
of freedom, but by the many-body calculations [19] based on the realistic NN pot-
tentials and, in case of intermediate nuclei, via convolution model for the nuclear
spectral function $S(E, \vec{p})$ *[18], taking into account the generation of the correlated
pair due to the short-range and tensor part of nuclear potential and the motion of
this pair relative to the remaining (A-2) nucleons of the nucleus. High momenta of
the nucleons of such pairs are thus due to the small relative distance within the pair
and the high removal energies are caused by the freeing the additional nucleon in the

* $S(E, \vec{p}) dE d\vec{p}$ is the probability that the struck target nucleon will have the removal energy, $E$, and
momentum, $\vec{p}$, in the infinitesimal interval, $dEd\vec{p}$, around $(E, \vec{p})$. 
process of nucleon knock-out. Even at the present stage however, only the part of the phase space accessible to \((e, e'p)\) experiments is satisfactorily described by the many-body calculations above, while the differences can be attributed to the multi-nucleon correlations.

1.2 Heavy-Ion Physics, Search for the Fifth State of Matter

New experimental results sometimes create challenges for the theoretical interpretations. However, often theory boosts the development of the new experimental techniques, needed to confirm or disprove one or another major theoretical hypothesis. The interfaces of different fields of science are always rich with the most critical findings.

While nuclear physics and physics of elementary particles are well established as independent vast branches of research, one of the major recent directions in nuclear physics was prompted by astrophysics, which in particular concentrates on the study of first moments of the universe. At some early time of the order of \(10^{-6}\) sec after the big bang matter probably occupied very small space and possessed extremely high pressure and temperature. The ambitious task of the physicists is to reproduce in the laboratory and study the conditions of extremely high energy density and temperature prevailing at the first stage of the universe evolution.

One of the motivations of the heavy-ion physics, as a separate direction in the nuclear physics, is the search for the non-hadronic degrees of freedom in the nuclear
reactions. Since the mean free path of hadrons is smaller than the radii of the nuclei, the nucleus-nucleus collision should exhibit the effects of collective motion of the nuclear matter, with essential destruction of the initial nuclear wave functions. The collective effects would of course manifest themselves stronger with increasing the initial energy of colliding nuclei and would be more important for the collision of the heavy nuclei.

In the collisions of the high-energy heavy-ions with each other the longitudinal energy of the participant nucleons is partially transformed into the transverse degrees of freedom and to the energies of the abundantly produced particles. While the physics of the hadron-nucleus reaction is essentially described by the concept of the hadron gas \(^1\), the heavy-ion collision with its collective effects would rather be characterized via the term hadron fluid \(^1\).

Essentially collective behavior of the participants in the heavy-ion collision, the rapid transformation of the collective motion energy of the incoming nuclei into the thermal energy and the transverse flow, is accompanied by creation of very high pressure and energy density, which could possibly lead to a sequence of phase transitions. Under certain conditions one would be able to encounter a new state of nuclear matter – the Quark-Gluon-Plasma. It is hoped that high energy heavy-ion collisions can provide us with these necessary conditions. Since the initial interaction volume and

\(^1\)Independent motion and binary collisions of hadrons in the mean-field potential generated by nucleonic matter.
\(^1\)Strongly correlated system of participating constituents having mean free path comparable to the relative distance between constituents and thus exhibiting the collective motion.
parton multiplicity turn out to be large, the dense parton system may go through both thermalization and chemical equilibration processes after its formation. At extremely high energies one may study an almost fully-equilibrated plasma of quarks and gluons with zero baryon density, and at currently accessible energies, matter with high baryon density may be studies as well. The QCD phase transition is also expected to occur at high baryon density [22].

1.2.1 Traditional Interpretations of Heavy-Ion Experiments

Because of the immaturity of current lattice QCD calculations, there exists no model nowadays capable of explicit treatment of the QCD degrees of freedom that are expected to play an important role during some part of heavy-ion collision. Therefore, current experimental data are interpreted in terms of hadronic models. Even though there are reasons to believe that space-time evolution of heavy-ion collisions involves a transition to QGP [12], the hadronic models provide a useful background for comparison with experimental data and search of new phenomena.

Following closely to [7] let us consider the length scales, encountered in nucleus-nucleus collisions in order to illustrate applicability of various hadronic models. Some of the relevant scales are the force range $\hbar/m_p c$, the nuclear radius $R \approx 1.2 A^{1/3}$, the de Broglie wavelength of nucleons $\hbar/k$ and the mean free path $\lambda$:

$$\lambda = \frac{1}{\sigma_{NN\rho}} \approx \frac{\hbar}{m_p c \rho},$$

(1.1)
with saturation nuclear density $\rho_0 = 0.15 \ h/k$. Usually the condition $R \gg h/k$ serves as motivation for the application of the classical methods in description of the $AA$ collision. This condition is obviously satisfied for nucleons on the early stage of relativistic heavy-ion collisions. The following conditions

$$\frac{\hbar}{m_p c} \ll \lambda \Rightarrow \rho/\rho_0 \ll 1, \quad (1.2)$$

$$\lambda \ll R \Rightarrow \rho/\rho_0 \gg A^{-1/3}, \quad (1.3)$$

$$\frac{\hbar}{m_p c} \ll \lambda \ll R \Rightarrow A^{-1/3} \ll \rho/\rho_0 \ll 1. \quad (1.4)$$

correspond to various scenarios of the collision process.

**Intranuclear Cascade Models:** Case (1.2) corresponds to the dilute gas system, when isolated collisions occur with negligible effect of the nuclear potential on the classical motion of particles. Such scenario corresponds the conventional INC model (Intranuclear Cascade Model) or to the Boltzmann equations. The condition of the (1.2) case is obviously not satisfied since in $AA$ collisions the nuclear density is much higher than normal.

**Viscous Hydrodynamics:** In case (1.3), $\lambda \ll R$, the mean free path is short compared to the size of the system, so that we can talk about temperature and pressure in each macroscopic element of the nuclear matter, in other words the local thermal equilibrium can be reached in such system. This is a case of viscous hydrodynamics. Potential effects are important here and can be introduced via the equation of state.
for nuclear matter $P(\rho, T)$. The important consideration for viscous hydrodynamics is that the necessary condition $\lambda \ll R$ is being violated for the particles in the periphery of the A+A system, therefore there always will be a part of the spectrum that can not be described by purely hydrodynamical model.

**Hydrodynamics of Ideal Gas, Thermal Models:** More specific is the case (1.4). Here the conditions of (1.3) are satisfied, but the importance of the potential effects is very small, so we encounter the hydrodynamics of an ideal gas. With the conditions of (1.4) the multiple collisions are most important, so that the final particle spectra completely lose their memory about the detailed structure of the individual baryon-baryon, meson-baryon and meson-meson collisions. In this case the thermal or statistical models are applicable.

None of the conditions above is satisfied in practice, but none of them is significantly violated, therefore the models of each kind eventually succeed in giving a reasonable semi-quantitative description to a limited number of observables in restricted regions of the phase space.

1.2.2 Space-Time Evolution of Heavy-Ion Collisions

A simple yet powerful picture of hadronic matter evolution, first suggested by J. Bjorken [83] is provided by one-dimensional hydrodynamic expansion in the scaling
variables:

\[ \tau = \sqrt{t^2 - z^2}, \]  \hspace{1cm} (1.5)

\[ \chi = \frac{1}{2} \ln \left( \frac{t + z}{t - z} \right). \]  \hspace{1cm} (1.6)

Here \( t \) and \( z \) are the time and spatial coordinates of one dimensional expansion in c.m.s. and \( \tau \) is the proper time. In these variables, using the initial velocity distribution determined by a rapidity plateau

\[ Y(\chi, \tau = \tau_i) = \begin{cases} 
  \text{const}, & |\chi| \leq Y_m, \\
  0, & |\chi| > Y_m,
\end{cases} \]  \hspace{1cm} (1.7)

the solution of the equations of one-dimensional hydrodynamic expansion is especially simple (scaling solution). The temperature turns out to be a function of only the proper time \( \tau \) of hadron fluid element, i.e.,

\[ T(\chi, t) = T(\tau), \]  \hspace{1cm} (1.8)

and not of the variable \( \chi \). The rapidities of the elements are in turn determined only by the variable \( \chi \):

\[ Y(\chi, t) = \chi. \]  \hspace{1cm} (1.9)

Correspondingly, the velocities of matter elements are given by

\[ v_l(\chi, t) = \frac{z}{t} = \tanh(\chi). \]  \hspace{1cm} (1.10)

A schematic picture of hadronic matter evolution in these variables is presented in Figure 1.1. In the region \( 0 \leq \tau \leq \tau_i \) the hadronic matter is in the pre-equilibrium
phase; then for $\tau_i \leq \tau \leq \tau_q$ there exists an equilibrium QGP phase, followed by a mixed phase region, and finally a hadron fluid (gas). At $\tau = \tau_f$ the system decays into observed secondary hadrons. From Figure 1.1 one can imagine how difficult it is to detect the QGP experimentally. The QGP formation signal is obscured by emission from pre-equilibrium, hadron and mixed phases.

Let us consider the heavy-ion collision in the center of mass system at very high energy. One will observe two highly Lorentz contracted heavy-ions moving toward each other at almost speed of light ($\gamma \gg 1$). A system of quanta is produced when the nuclei intersect in space. If the quanta produced rapidly reach local thermal equilibrium with each other, the motion of these quanta in the small but macroscopic fluid element seen from any longitudinally co-moving reference system will correspond to the equal distribution of energy between three orthogonal directions. If we now transfer to another system moving with a modest velocity in the longitudinal direction, relative to the former system, we must see just the same distribution of the quanta velocities, since the projectile and target pancakes again recede at the speed of light. This symmetry of the initial conditions relative to the Lorentz boost in the longitudinal direction in the initial condition of the fluid must lead to the symmetry in the whole evolution of the heavy-ion collision. As a result the spectra of the produced particles exhibit a plateau in the rapidity distribution.

The above picture of the high energy heavy-ion collision was first described by J. Bjorken and applied to the CERN SPS energy [83] and higher. It is important to
note that the boost invariant expansion of the hadronic fluid dictates the following linear dependence of the collective longitudinal velocity on longitudinal coordinate:

\[ v(z, t) \equiv \frac{dz}{dt} = \frac{z}{t}, \tag{1.11} \]

which is the only dependence invariant relative to the Lorentz transformations,

\[ z' = z \cosh(\eta) + t \sinh(\eta), \tag{1.12} \]
\[ t' = z \sinh(\eta) + t \cosh(\eta), \]

where \( \eta \) is a longitudinal boost angle related to the boost velocity as \( v = \tanh(\eta) \).

Namely, in the co-moving reference frame we will still have the same linear law of hadronic fluid expansion:

\[ v'(z', t') \equiv \frac{dz'}{dt'} = \frac{dz \cosh(\eta) + dt \sinh(\eta)}{dz \sinh(\eta) + dt \cosh(\eta)} = \frac{z'}{t'}, \tag{1.13} \]

according to (1.11) and (1.12). From the boost invariant expansion law (1.11) it follows that any fluid element moves at a constant longitudinal velocity:

\[ z(t) = z(t_0) \frac{t}{t_0} = v(z(t_0), t_0) t, \tag{1.14} \]

defined by the initial condition at the moment of time \( t_0 \). Hence the rapidity distribution of the fluid elements, - particle sources, is "frozen" during hadronic matter expansion, which preserves its boost invariance.

One dimensional model of nuclear matter evolution is, of course, only a first approximation to the real evolution in heavy-ion collisions. The expansion in the transverse direction progresses at later stages of the collision. Near the periphery of
the highly compressed disk of nuclear matter we may expect the particles to escape in the outward direction. At distances slightly less than peripheral there should be a rarefaction front propagating inward at the speed of sound $c_s$. At distances from the collision axis less than $c_s t$ (for $z = 0$) the fluid is not "aware" that the nucleus is of finite size and the expansion is yet one dimensional. The equation of the rarefaction front, to the extent that the sound velocity is constant, is

$$\rho(t) \approx R - c_s \sqrt{t^2 - z^2},$$

(1.15)

where $R$ is the radius of the nucleus.

The boost invariance scenario holds true only in the limit of very high energies, when the Lorentz factor of the heavy-ions in the center of mass system is much greater than unity. A relatively small gamma factor at AGS energy, $\gamma = \cosh(y_0 = 1.6) = 2.53$, together with the strong stopping power of protons at AGS, make the boost invariant picture generally not applicable.

The large number of particles involved in the collision dynamics of the gold ions at AGS, whether QGP transition did or did not take place, allows one to talk about temperature and pressure for the description of the process. The hadronic matter in turn can be described as the fluid (or gas) during the collision process, which evolves until the finally produced hadrons cease to interact with each other and freely flow to the detectors. The velocity profile of this fluid in the moment of freeze-out will thus be imprinted on the particle spectra as they are measured by the detectors.
Flow Excitation Function: If the local thermal equilibrium is reached by the moment of freeze-out, the temperature of the fluid element $T$ is a function of its velocity $\vec{v}$ and velocity is distributed according to a distribution function $w(\vec{v})$, which we will call the flow excitation function. A particle spectrum in this case would be described via superposition of the Boltzmann sources:

$$E \frac{dN}{d\vec{p}} = \int d\vec{v} w(\vec{v}) f_B(\vec{v}, \vec{p}, T(\vec{v})).$$  

Once the integration is performed in (1.16), the detailed information about the functions $w(\vec{v})$ and $T(\vec{v})$ is lost. Obviously the two unknown functions cannot be unambiguously extracted, from the measurements of the particle spectrum. At least one of the functions above must be known from the theory, which makes reconstruction of the other function model dependent. Apart from that, the local thermal equilibrium may be not reached or broken by the moment of freeze-out [26].

In relativistic regime it is convenient to replace the longitudinal velocity of the source by the longitudinal boost angle, defined as follows

$$\eta = \frac{1}{2} \ln \left( \frac{1 + v_z}{1 - v_z} \right).$$  

In terms of the boost angle the flow excitation function is related to the original $w(v_t, v_z)$ via the following equation:

$$w(v_t, \eta) = \frac{2\pi w(v_t, v_z)v_t}{\cosh^2(\eta)},$$  

(1.18)
where azimuthal symmetry of flow is assumed. Equation (1.16) can thus be rewritten as

$$\frac{dN}{dp^2} = \int \int d\nu_t d\eta \omega(\nu_t, \eta) f_B(\nu_t, \eta, \vec{p}, T(\nu_t, \eta)).$$  \hspace{1cm} (1.19)

Apparently the flow excitation function \( \omega(\nu_t, \nu_z) \) should fall with increasing \( \nu_t \), since the faster fluid elements gain their velocity due to the pressure from the slower ones. The concurring factor, \( \nu_t \), in (1.18) will thus create the maximum of \( \omega(\nu_t, \eta) \) at a value of transverse velocity, \( \nu_t = \nu_{t, \text{max}}(\eta) \). The particle spectrum (1.19) will therefore be determined mainly by the sources moving with \( \nu_{t, \text{max}}(\eta) \). Thus, as an approximation, it is possible to reduce the integration in (1.19) to the integration over the longitudinal boost angle \( \eta \) only:

$$\frac{dN}{dp^2} = \int d\eta \bar{\omega}(\eta) f_B(\nu_{t, \text{max}}(\eta), \eta, \vec{p}, T(\nu_{t, \text{max}}(\eta), \eta)), \hspace{1cm} (1.20)$$

where \( \bar{\omega}(\eta) = \omega(\nu_{t, \text{max}}(\eta), \eta) \).

In equation (1.20) the integration is performed over the flow surface instead of the flow volume. We shall use this, so called, blast wave approximation in this thesis for a parametrization of particle spectra in the Au + Au central interactions, where, the most bulk of the nuclear matter is expected to obey the laws of hydrodynamics.

### 1.3 QCD Transition, Phase Diagram

From lattice QCD numerical simulations, the QCD phase transition happens at \( T_c = 150 \pm 10 \ MeV \) [20] at zero net baryon density. The QCD phase transition is
Figure 1.1  Space-time evolution of hadronic matter in the one-dimensional hydrodynamic model.
also expected to occur at some high baryon density $n_c^B$ at zero temperature, when the hadrons are so dense that the partons from different nucleons start to overlap. At present, there is no lattice QCD calculation on phase transitions with finite baryon density $n_c^B$. In general one expects that the QCD phase transition occurs for a smooth curve on the plane of temperature $T$ and the baryon chemical potential $\mu_B$. In terms of baryon number densities, there is a jump from the hadronic phase to the QGP phase, just like the jump in the energy densities at zero baryon density. The phase diagram for the transition from hadronic phase to the QGP phase is roughly plotted in Figure 1.2 as a function of the temperature $T$ and the baryon chemical potential $\mu_B$ [25]. Figure 1.3 shows the phase diagram as a function of the temperature $T$ and the scaled baryon number density $n_B/n_B^0$ where $n_B^0 \approx 0.14/fm^3$ is the normal density in nuclear matter. The shaded area in Figure 1.2 represents the error from our current knowledge about the phase transition, and the shaded area in Figure 1.3 mostly represents the mixed phase.

1.4 Signals of the Quark-Gluon Plasma

In this section we discuss several key observables of the quark-gluon plasma, as sketched in Figure 1.4 borrowed from [25]. Among the probes of the early stages of quark-gluon plasma formation, electromagnetic probes such as direct photons [29, 30, 31] and dileptons [21, 32] are the most direct signals, since they do not experience the subsequent thermalization and final state interactions. Heavy quarkonia [33] and
Figure 1.2  Schematic QCD phase transition diagram as a function of $T$ and $\mu_B$. 
Figure 1.3 Schematic QCD phase transition diagram as a function of $T$ and $n_B/n_B^0$. 
Figure 1.4 Several key observables of the quark-gluon plasma.
heavy quark [28, 34] productions are also proposed to be the sensitive probes of the properties of the QGP. Recently new interesting experimental results such as low mass dilepton enhancement [35] and $J/\psi$ suppression in $A + B$ collisions [36] were presented.

1.4.1 Heavy Quarkonia

The production and suppression of heavy quarkonia bound states, such as $J/\psi$, was proposed by Matsui and Satz [33] to be an ideal test of the quark-gluon plasma production. The authors claimed that "there appears to be no mechanism for $J/\psi$ suppression in nuclear collisions except the formation of a deconfining plasma, and if such a plasma is produced, there seems to be no way to avoid $J/\psi$ suppression." The proposed suppression signature of QGP formation was first observed in 200 $A GeV O + U$ reactions at the CERN/SPS in 1987 [37]. However, far from settling the matter, the debate on the interpretation of those data and on the uniqueness of the QGP mechanism for suppression has only intensified since then.

1.4.2 Dileptons

Dileptons are one of the direct electromagnetic probes. Thermal dileptons, for example, can tell us about the thermal profile of the dense plasma. In [31] it was proposed that in the hot-glue scenario, where the plasma is less chemically equilibrated
but hotter than in the chemically equilibrated plasmas, dileptons of a few GeV energy will be enhanced over the standard scenario by a factor of 2.

However, in order to observe the interesting thermal signals, one has to deal with a large dilepton background. One such background is the dileptons from heavy flavor meson pairs decays. In a study by Vogt et al. [38], it was suggested that dileptons from open charm decay would be about an order of magnitude above the thermal and Drell-Yan signals in the few GeV region, and therefore it would be very difficult to subtract this large background and observe the desired plasma signal. However, Shuryak [75] suggested that one important effect was missing in the arguments of that study [38], which was the energy loss of the charm quark in the dense medium formed after high energy $A + A$ collisions. Once one assumes that the charm quark would lose a finite amount of energy along its path (typically 2 GeV/fm), then the dilepton spectrum from open charm will be pushed to the low mass region, and dileptons with a few GeV energy would be suppressed by over an order of magnitude.

1.4.3 Direct Photons

Direct photons are also one of the direct probes which could provide pristine information on the quark-gluon plasma. In addition, since the direct photon production in nucleon-nucleon collisions [40] is understood, its production in nucleus-nucleus collisions could provide constraints on the little-known gluon structure functions in the nucleus via the Compton process $gq \rightarrow q\gamma$. 
Data on single photons in $S+Au$ collisions [41] was claimed to be understood only if the quark-gluon plasma is formed, because an overly simplified hadronic fireball model led to yields a signal an order of magnitude [42, 43] above the data. The problem with this interpretation is that the fireball model is inconsistent with the observed rapidity and transverse momentum distributions.

The large background is another problem with the direct photon measurements. The photon background, which mainly comes from $\pi^0$ and $\eta$ decays, is usually much higher [39] compared to dileptons. Direct photon signatures are therefore rather difficult to disentangle [44].

1.4.4 Disoriented Chiral Condensate

The restoration of chiral symmetry at high temperatures is one of the most basic features of the QCD matter. According to lattice QCD calculations the critical temperature for chiral restoration coincides with that for the deconfinement transition. Thus the chiral phase transition in QCD is expected from both the theory study and lattice numerical simulations [23, 45]. A possible signal for the chiral transition is the formation of a coherent pion field called the disoriented chiral condensate (DCC) [46, 47, 48]. In equilibrium, a DCC can not grow larger than $\sim 1/T_c \sim 1\, fm$ [47]. However, in non-equilibrium it is possible to grow larger. A large domain of DCC will manifest itself by the striking fluctuation in the spectrum of neutral and charged
pions in terms of the ratio \( n_{\pi^0} / n_{\pi^0}^{\text{total}} \). There also exists some cosmic ray data [49] with that ratio far from the average value 1/3.

1.4.5 Open Charm

One of the main uncertainties in heavy-ion physics is the initial conditions of the dense parton system. The initial momentum distribution, energy density and chemical composition determine the fate of the QGP. Recently it is found that minijets lead to initial conditions characterized by large fluctuations of the local energy density and of the collective flow field. The "hot spots" [74] corresponding to higher local energy density may change the usual picture from cylindrically symmetric, homogeneous quark-gluon plasma to turbulent, volcanic plasma. Certain signals of QGP may be sensitive to this change of the initial condition. Initial conditions are themselves determined by the parton structure functions from the two colliding nuclei, mostly by the gluon structure function [25].

Open charm production is well suited to measure the initial gluon structure function and its nuclear modifications. At collider energies \( \sqrt{s} > 200 \, A \, GeV \) the initial minijet plasma is mostly gluonic [29, 24] with a quark content far below its chemical equilibrium value. Because open charm is produced mainly through gluon fusion, it can provide information on both the incoming gluon structure functions and the later-formed dense minijet plasma [25].
1.4.6 Hanbury-Brown Twiss

Correlation functions for identical particles probe the source size and freeze-out time of the dense medium [67, 68]. The two-particle correlation function in momentum space is defined as

$$ C_2(\vec{p}_1, \vec{p}_2) \equiv \frac{P(\vec{p}_1, \vec{p}_2)}{P(\vec{p}_1)P(\vec{p}_2)}. \tag{1.21} $$

For classical particles without dynamical correlations, the above function is always equal to 1. For identical quanta with bosonic statistics [69], symmetrization introduces a dependence on two momenta and $C_2(\vec{p}_1, \vec{p}_2) = 2$ without dynamical correlations. More generally, for a source described by seven-dimensional phase space density function $S(x, \vec{p})$, the semi-classical expression for the correlation function is given by [68, 70]

$$ C_2(\vec{K}, \vec{q}) = 1 + \frac{\int d^4x d^4x' S(x, \vec{K}) S(x', \vec{K}) e^{iq(x-x')}}{\int d^4x S(x, \vec{p}) \int d^4x' S(x', \vec{p})}, \tag{1.22} $$

where $\vec{K} = (\vec{p}_1 + \vec{p}_2)/2$, $q = (E_1 - E_1, \vec{p}_1 - \vec{p}_2)$, and $\int d^4x S(x, \vec{p}) \equiv P(\vec{p})$.

The following parametrization is used experimentally in order obtain information about the space-time dimension and expansion velocity of the emitting source from the fitted radii [72]:

$$ C_2(\vec{K}, \vec{q}) = 1 + \lambda \exp \left[ -q_{\text{side}}^2 R_{\text{side}}^2 - q_{\text{out}}^2 R_{\text{out}}^2 - q_{\text{long}}^2 R_{\text{long}}^2 \right], \tag{1.23} $$

where $q_{\text{long}}$ is defined to be the component of $\vec{q}$ along the beam direction, $q_{\text{out}}$ is the transverse component which is in the plane with the beam direction and the $K$ vector, and $q_{\text{side}}$ is the other orthogonal transverse direction.
The chaoticity parameter $\lambda < 1$, is introduced to account for correlations other than HBT and resonance decays. For charged particles such as pions and kaons, the Coulomb interactions are important and must be included. Data from $p + Pb$, $S + Pb$ and $Pb + Pb$ collisions [73] show that the fitted radii grow with the size of the system. Pions have larger radii and smaller chaoticity than kaons since the resonance decays are more important in pion production. Because of the transverse expansion, the fitted radii do not correspond directly to the source size. Nevertheless, pion interferometry can serve as a powerful signal of the QGP formation. Given a rapid crossover region around $T_c$, the speed of sound $c_s^2$ acquires a local minimum. As matter cools, the minimum of $c_s^2$ causes the transverse expansion to stall for some time. This leads to the time delay signature of the phase transition that could be seen in detailed systematics of the correlation function $C_2(q_{out}, q_{side})$ [65].

### 1.4.7 Collective Nuclear Flow

Recently, the study of collective flow in the nuclear collisions at relativistic energies has attracted an increased attention of both experimentalists and theorists. Although all forms of flow are interrelated and represent a different part of the global picture, usually people discuss various forms of collective flow, such as longitudinal expansion, radial transverse expansion, directed flow and elliptic flow [76, 77]. Flow introduces strong space-momentum correlations in particle production. Particles with a given rapidity and the transverse momentum are produced mainly by a portion of the source
moving with rapidity close to that of particles in the midrapidity region, or slightly less than that for particles in the fragmentation region.

Transverse flow from azimuthal anisotropy was observed recently in $Au + Au$ collisions at 10.8 $A GeV$ at the AGS [54]. It was first discovered in 1984 at much lower energies at Bevalac [59], and has been the subject of intense studies since then. In general, flow can provide information on the compressibility of the dense baryonic system and on the equation of state. It can also assist us in studying the role of the mean field in nuclear collisions.

The azimuthal distribution of energy and multiplicity would be isotropic if there were no directed transverse flow. Directed transverse flow manifests itself via the anisotropy in the azimuthal distribution on an event-by-event basis. Directions of the particles produced in a collision are correlated with the orientation of the impact parameter. The direction of the impact parameter is unknown a priori: it is estimated event by event from particles produced in the collision. These estimates are generally subject to large statistical fluctuations due to the finite multiplicities. To study the anisotropy, one can Fourier-transform the azimuthal distribution [60]:

$$f(\phi) = \sum_i f_i e^{in\phi} = \sum_i v_i e^{-i\psi_n} e^{in\phi}. \quad (1.24)$$

Anisotropy gives non-zero values of $v_n$ ($n \geq 1$). Variables $v_1$ and $\psi_1$ for the first harmonic represent a shift of the center of the isotropy circle, with $\psi_1$ being the azimuthal angle of the reaction plane, which is the plane formed by the beam axis and
the impact parameter vector. The variable $v_2$ represents the eccentricity of the ellipse in the reaction plane. The above analysis is complicated by fluctuations resulting from finite multiplicities in a nuclear collision [61]. Therefore the reaction plane can not be exactly determined event by event, and the values of $\psi_n$ suffer from this uncertainty.

From the experiment E877 at AGS [62], clear transverse flow signal is observed in mid-central events. For protons, transverse flow grows linearly as a function of $p_T$, and reaches 20% of the total transverse momentum at $p_T = 1\, GeV$. By comparing with different transport models such as RQMD [63] and ARC [64], one can study the effect of the mean field in the nuclear collisions. The most striking collective flow signature of a QGP formation has been predicted to be a local minimum of the flow excitation function at the so-called "softest point" of the nuclear matter [66]. This proposed signature is being looked for at AGS. On Figure ?? brought from [66] the dependence of the speed of sound squared, $c_s^2 = dp/d\epsilon$, is shown in particular. The hydrodynamic expansion is driven by the pressure gradients. The speed of sound is quantity relevant to the ability of the system to transform the thermal energy into the mechanical work. In the mixed phase with the transition interval, $\Delta T = 0$, the speed of sound vanishes and the matter does not expand and cool for its own account. For a finite width of the transition region, $\Delta T = 0.1T_c$, the ability of the system to expand is still reduced compared to the hadronic gas without the phase transition to QGP.
Figure 1.5 Macroscopic quantities of the collision system with and without the phase transition to QGP.
1.4.8 Strangeness Enhancement

It was shown in [79] that one of the possible signatures of the transition to the QGP in the course of the heavy-ion collision is the increased rate for the strangeness production as compared to the superposition of the hadron-hadron collisions.

The world is composed mostly of the protons and neutrons – the stable hadrons – the constituents of nuclei. In terms of QCD the nucleons are the combinations of three quarks of two lightest flavors – $u$ and $d$ valence quarks and the sea quark antiquark pairs. The lightest quark among the heavy flavors is the strange quark $s$. When the energy of the collision of two hadrons exceeds the threshold for the strange quark-antiquark pair production, strange matter is excited via the reaction process by strong or electromagnetic interactions. The formation mechanism of strange quark production depends on the mechanism governing hadronic collisions. In our current understanding abundant strangeness generation in nuclear collisions may be a useful signature of QGP formation. It is likely that under extreme conditions of high hadron density and excitation in heavy-ion collisions, various fundamental properties of the strange quark production might be altered. However another conventional picture is also possible, involving collisions of hadrons and hadronic resonances.

The first situation is the one we are looking for, in order to confirm the prediction of QCD. The principal difference between the two mechanisms is that deconfinement in the quark matter is accompanied by abundance of highly excited gluons, which are especially efficient in generating strange $q\bar{q}$ pairs. In case of Chiral Symmetry
Figure 1.6 Proposed mechanism of strangeness enhancement for high-density quark matter [80].
restoration, the quarks lose their mass, thus reducing the threshold for the strange quark production. Detection of strangeness thus provides a unique way to indirectly observe the QGP formation.

Generally gluons produce light quarks more copiously than the strange quarks because of the heavier mass of the latter. To illustrate the mechanism of the strangeness enhancement at zero temperature let us refer to the Figure 1.6. Due to the large mass of the strange quarks the strangeness production in normal nuclear matter is suppressed as compared to the production of non-strange hadrons. At normal densities the Pauli exclusion principle in the dense QGP when the quark densities are high fermi levels of the light quarks increase (see illustration on Figure 1.6). Due to Pauli exclusion principle the energy states of $u$ and $d$ quarks below the fermi levels can not be created. Production of $s$ quark takes as much energy as creating the light quarks $u$ and $d$ [80]. It is possible however that the strangeness enhancement destroyed during hadronic stage of the collision.

1.5 Some Results of Current Heavy-Ion Experiments

Since the mid 1970's the experimental search for new phases of nuclear matter [5, 6] using heavy-ion reactions has been underway. Until 1984 only light ion reactions at energies below 2 A GeV (GeV per incident baryon) [7] were possible. In the past ten years, experiments with heavy-ion with energies up to 200 A GeV have been investigated [8, 9, 10]. In 1999 a new era of experiments with heavy-ion beams in
the collider mode at the Relativistic Heavy-Ion Collider (RHIC) at the Brookhaven National Laboratory (BNL) with center of mass energies up to 100 $A$ GeV will begin.

Current heavy-ion experiments at GSI-SIS, BNL-AGS and CERN-SPS \(^5\) are aimed at the creation of the Quark Gluon Plasma and measurement of its properties. Parallel theoretical effort is being made towards understanding of the reaction dynamics of heavy-ion collisions and its evolution as a function of the mass of the colliding system, in terms of hadronic degrees of freedom [97, 99]. Among the beams available nowadays in heavy-ion accelerators, especially favorable conditions for the creation of the dense and hot nuclear matter are provided by the gold beam at AGS and lead beam at SPS. Therefore a number of experiments studying $Au + Au$ collisions [91, 54, 55] and $Pb + Pb$ collisions [57, 58] were performed and analyzed during the last few years.

In this thesis we present the first measurements of the $\Lambda$ hyperon production cross sections and the measurements of the $\pi^-$ meson production cross section in the region of the phase space complementary to the existing data from experiments BNL-E866 and BNL-E877. We therefore review the results and conclusions of above experimental collaborations concerning the differential multiplicities of the particles produced in the central $Au + Au$ interactions.

\(^{\text{5}}\)

- SIS at GSI (Darmstadt, Germany) is a heavy-ion synchrotron used to accelerate heavy-ions from $^{12}C$ to $^{238}U$ up to the momentum of 2 $A$ GeV/c.
- AGS at BNL (Upton, NY, USA) is an Alternating Gradient Synchrotron capable of accelerating heavy-ions up to $^{197}Au$. The current maximum beam momentum is 11.6 $A$ GeV/c.
- SPS at CERN (Geneva, Switzerland) now provides a $^{207}Pb$ beam with 160 $A$ GeV/c.
1.5.1 Experiment BNL-E866

The results of the particles multiplicities measured by E866 collaboration have been presented at the Quark Matter '95 conference [92]. The double differential cross sections and rapidity distributions were measured for the following particles: $p, \pi^\pm, K^\pm$. Some of the experimental data are shown on Figures 1.7, 1.8 and 1.9. The central events with the cross section of 210 mb corresponding to $\approx 4\%$ of the geometric Au + Au cross sections were used for analysis.

The E866 experimental acceptance for $\pi^\pm$ is limited in the range of rapidities, $1.5 < y < 2.7$, and in the range of the transverse masses, $0.0 < m_t - m_0 < 1.2 \text{ GeV}/c^2$. The deviation from a single exponential scaling, a well known low $m_t$ enhancement below $200 \text{ MeV}/c^2$, is observed. The low $m_t$ component is found to constitute approximately 25% of the rapidity density value of pions. The departure from the single exponential scaling is more noticeable for $\pi^-$ than that for $\pi^+$. The maximum value of rapidity density, $dN/dy$, is found to be $\approx 55$ for $\pi^+$ and $\approx 70$ for $\pi^-$, under assumption that low $m_t$ component independent on rapidity is included.

The total yields of pions were found using assumption of a gaussian distribution in rapidity: $N_{\pi^+} \approx 115$ and $N_{\pi^-} \approx 160$. It was noted, that the yield and shape of $\pi$ meson spectra in the low $m_t$ region are in qualitative agreement with the assumption of excessive creation of baryonic resonances, in particular $\Delta$ resonances, in a baryon rich regime at AGS. The difference of low $m_t$ shapes for the $\pi^+$ and $\pi^-$ was attributed
to either feed-down from hyperons or to the mean field modifications of the spectra, in particular due to the Coulomb effects.

The protons spectra in the central (4%) Au + Au interactions were measured by E866 collaboration in the rapidity interval, \(0.7 < y < 2.1\), and the transverse mass interval, \(0.0 < m_t - m_p < 1.2 \text{ GeV}/c^2\). The invariant transverse mass distributions for protons are shown on Figure 1.7. An interesting phenomenon of the low \(m_t\) suppression was observed with the effect increasing at midrapidity. The transverse mass distributions exhibit a convex shape, with curvatures increasing towards midrapidity, which can be explained as a collective transverse expansion of the proton gas. The spherical expansion model was employed in order to parametrize the experimental data on invariant particle multiplicities at midrapidity. However the spherical expansion picture fails to account for anisotropy of the particle spectra.

### 1.5.2 Experiment BNL-E877

The recently published data from experiment E877 [55] contain the invariant differential multiplicities for \(p, \pi^+, \pi^-, K^+, K^-\) and the preliminary data for deuterons in the central (4%) Au + Au collisions. The acceptance coverage for protons is in the range, \(2.2 < y_p < 3.5\), and, \(0.0 < m_t - m_p < 0.75 \text{ GeV}/c^2\). The pion acceptance coverage is \(2.8 < y_\pi < 4.5\) and \(0.0 < m_t - m_\pi < 0.5 \text{ GeV}/c^2\) for \(\pi^+\) and \(0.0 < m_t - m_\pi < 0.85 \text{ GeV}/c^2\) for \(\pi^-\).
Figure 1.7 Transverse mass invariant spectra of protons in central (4%) \( Au + Au \) interactions measured by BNL-E866 collaboration. Overlaid histograms represent calculations of cascade models. Circles: data taken with the Forward Spectrometer. Frilled squares: data taken with the Henry Higgins spectrometer. Dashed histograms: RQMD. Dotted histogram: ARC. The successive spectra are for different rapidity (\( y \)) bins, from \( y = 0.7 \) to \( y = 2.1 \) (bottom) with intervals of 0.2. Each successive spectrum is divided by 10. The error bars indicate the statistical errors only.
Figure 1.8 Transverse mass invariant spectra of $\pi^+$ and $\pi^-$ measured by BNL-E866 collaboration. The data are shown for rapidity bins, $y = 1.5$ to $y = 2.7$, with a bin size of 0.2. For each subsequent bin in rapidity the cross section has been multiplied by 0.1 before being plotted.
Figure 1.9  Rapidity density distributions for $\pi^-$, $\pi^+$, $K^+$ and $K^-$, in central $Au + Au$ collisions, measured by BNL-E866 collaboration. The $\pi^-$ data have been multiplied by 10 before plotting. The data with crosses are from E866 analysis [92], while the remaining data are from [93]. The full drawn symbols are the measured data and the open circles the data reflected around $y_{NN}$. 
The BNL-E877 collaboration has observed an enhancement of pions above a pure exponential for \( m_t - m_\pi < 0.2 \text{ GeV}/c^2 \) as shown on Figure 1.11. This effect was already observed for the lighter systems [56]. It was noted in [55] that the enhancement of pions systematically increases toward midrapidity. The effect of low \( m_t \) enhancement was found to be stronger for negative pions in accord with E866 observation. However, it was noted that for high values of the transverse mass the \( \pi^+ / \pi^- \) ratio is consistent with unity. It was suggested in [55] that the charge asymmetry for the pion production and its rapidity dependence could be attributed to the Coulomb potentials seen by each charge type at freeze-out.

The rapidity distributions of pions were obtained via integration of the transverse mass spectra. The proton rapidity distributions in principle allows one to study the energy deposition and the nuclear stopping power. It was claimed in [55] that rapidity densities of pions and protons have approximately the same widths in \( Au + Au \) central collisions unlike the light systems [84], where the proton's rapidity distribution was wider. As can be seen from Figure 1.12, the experiment E877 basically measures the projectile-like rapidities for both protons and pions, therefore the conclusion about the shapes of the rapidity distributions can not be unambiguously made.

First measurement of the invariant differential spectra for the composite particles, deuterons was presented in [55]. The study of the composite particles production in the cascade models has just started recently. Due to their mass the composite particles are a good probe for the collective effects in the heavy-ion collisions.
It was observed by E877 that the transverse mass distributions of deuterons are very flat and have a rather weak dependence on the rapidity. The data from reference [55] for deuterons are presented on Figure 1.13 together with cascade model predictions. It is remarked in [55] that the mean field effects are very important to account for the shapes of the fragments spectra in the Au + Au central collisions.

1.5.3 Our Measurements and Analysis: Experiment BNL-E891

Experiment E891 (see Appendix A) is aimed at the measurements of the particle spectra in central Au + Au collisions at AGS by means of the track reconstruction in three Time Projection Chambers placed in 10 kGauss magnetic field. There is no particle ID measurement in the experiment, therefore it is feasible only to reconstruct the Λ and K^0_s spectra from the two particle decays into charged hadrons by testing kinematical hypothesis about the masses of the daughter particles and to measure the spectra of negative pions, since the yields of the other negative hadrons in Au + Au collisions at AGS are greatly suppressed. The experimental acceptance of our apparatus is basically in the forward hemisphere of the center of mass system of Au + Au collision. However the midrapidity region is also covered for both Λ and π^-.

A first measurement of the invariant differential spectrum of Λ hyperons and their rapidity distribution was reported by our collaboration in [90, 89]. In this thesis we perform an extended analysis of the Λ production cross sections. The new data on the invariant differential cross sections for the Λ production near midrapidity, namely
Figure 1.10 Proton transverse mass spectra measured by BNL-E877 collaboration. The data are presented in rapidity bins of 0.1 unit widths successively multiplied by increasing powers of 10 for decreasing rapidity. Errors are statistical. Also shown are exponential fits which for $y \geq 2.9$ exclude $m_t - m_0 < 0.2 \text{GeV/c}^2$. 
Figure 1.11 Pion transverse mass spectra measured by BNL-E877 collaboration. The data are presented in rapidity bins of 0.1 unit widths successively multiplied by increasing powers of 10.
Figure 1.12  Rapidity distributions for $\pi^+$, $\pi^-$ and protons measured by BNL-E877 collaboration. An RQMD calculation is overlaid (solid line) for comparison.
Figure 1.13 Spectra of fragments measured by BNL-E877 collaboration. Transverse mass spectra for deuterons are shown on the left. The data are presented in rapidity bins of 0.1 widths successively multiplied by increasing powers of 10. Different RQMD predictions for rapidity bin $1.4 < y < 1.8$ are compared to deuteron (circles) and proton (triangles) data on the right. The dashed histograms correspond to predictions of standard RQMD while the full histograms include the effect of mean field. The full line is a Boltzmann distribution fitted to the large $p_t$ region of RQMD result for deuterons.
for $1.4 < y < 2.0$ will be presented here. An interesting phenomenon of the low $m_t$ suppression of the $\Lambda$ transverse mass distribution is observed in this region. Since the $\Lambda$ hyperons with a small transverse mass may suffer from systematic errors caused by the high ionization of the TPC gas via the gold beam, it is important to study these effects in great detail in order to obtain a quantitative estimate for the resulting systematic error. We therefore intend to investigate the influence of dynamic field distortions on the shapes of the particle spectra.

Experiment E891 provides a wide acceptance coverage for the $\pi^-$ spectra. Our TPC's are capable to measure some of the regions of the phase space which were not previously covered by either experiment E866 or by E877. Thus we perform the analysis of the invariant double differential cross sections for the negative pion production in this thesis in the range of $m_t - m_\pi < 1.0 \text{ GeV}/c^2$ and $1.5 < y < 3.5$.

It was suggested in [65] that the so-called "softest point" of the nuclear matter or a local minimum of the flow excitation function caused by the first order phase transition to the QGP may occur somewhere between Bevalac and AGS energies in $Au + Au$ collisions. In order to extract the quantities of collective nuclear flow and compare them to analogous quantities from the cascade model simulations we will introduce a three dimensional expansion picture, called Flow Model parametrization.
Chapter 2

Experimental Apparatus

For the measurements of the particle spectra in gold on gold collisions the experiment E891 employed a set of three TPC's (time projection chambers) previously used by experiment E810 for the measurements with the Si beam. The TPC, vividly called "electronic bubble chamber" has significant advantages in the measurements with the crowded particle environment, because of its ability to measure the three dimensional points along the tracks of the charged particles. (Fig 2.4, 2.5, 2.6)

The E891 experiment used the 11.6 GeV/A Au beam, provided by the AGS (Alternating Gradient Synchrotron) at BNL (Brookhaven National Laboratory) to bombard a Au target. The produced particles passed through the three TPC modules placed in the strong magnetic field (10 kGauss) of the MPS (Multi Particle Spectrometer) magnet. The momentum dependent curvatures of the tracks in the magnetic field allows one to calculate the momenta for each well measured track. The tracking program is used in order to reconstruct the tracks basing on the set of 3D points – hits left by the charged particles in the process of ionization inside the TPC. The neutral particles $\Lambda$ and $K^0$, decayed into the charged couples of tracks mostly upstream of the first TPC module. The vertex finding algorithm was used to reconstruct and to measure such decays.
2.1 **The Experimental Layout**

The experimental layout and the beam measurement detectors are shown in Figures 2.1 and 2.2 respectively. The Table 2.1 shows the positions and the sizes of the detectors. The beam went along z axis and was measured via two proportional wire chambers A5 and A6. Cherenkov counter C1 – C2 was used to select the Au ions. Halo veto scintillation counter S2 was used to reject events with interactions upstream from the target; while the coincidence of S1 and S3 determined the fact that the beam particle went to the target. In order to reduce the number of fake triggers, the multiplicity counter S4 was used to select events in which the beam particle interacted with the target. The counter S6 served to veto the fragments with $Z > 6$ in order to trigger on the most central events.

<table>
<thead>
<tr>
<th>Detector</th>
<th>Position z + x cm</th>
<th>Size z x y x z</th>
</tr>
</thead>
<tbody>
<tr>
<td>S1</td>
<td>-1013+0</td>
<td>6.5&quot;x6.5&quot;x3/8&quot;</td>
</tr>
<tr>
<td>S2</td>
<td>-350+0</td>
<td>12&quot;x12&quot;x3/8&quot;</td>
</tr>
<tr>
<td>S3</td>
<td>-25+0</td>
<td>1.5&quot;x1.5&quot;x1/32&quot;</td>
</tr>
<tr>
<td>S4</td>
<td>-52.5+0</td>
<td>4&quot;x4&quot;x1/4&quot;</td>
</tr>
<tr>
<td>S6</td>
<td>696+213</td>
<td>9&quot;x9&quot;x1/2&quot;</td>
</tr>
<tr>
<td>C1-C2</td>
<td>-924+0</td>
<td>8&quot;x6&quot;x1/16&quot;</td>
</tr>
<tr>
<td>Target</td>
<td>+2.5+0</td>
<td>6&quot;x6&quot;x1/100&quot;</td>
</tr>
<tr>
<td>PWC1</td>
<td>-998+0</td>
<td>12&quot;x6&quot;x6&quot;</td>
</tr>
<tr>
<td>PWC2</td>
<td>-360+0</td>
<td>12&quot;x6&quot;x6&quot;</td>
</tr>
</tbody>
</table>

**Table 2.1** Positions and sizes of the detectors in E891. Positions are measured relatively to the pivot point.
E891 Plan View

MPS Magnet 10 kGauss $\otimes$

TPC Modules

Drift Chambers

Multiplicity Counter S4

Target S3

Beam

Halo Veto S2

$Z^2$

Counter S6

BEAM = S1 * S2 * $\overline{C1}$ * S3

TRIGGER = BEAM * S4 * $\overline{S6}$

Figure 2.1 The E891 plan view.
Figure 2.2  E891 Beam.
2.2 The MPS Magnet

The MPS magnet is a wide solid angle acceptance magnet spectrometer (Figure 2.3). The field volume between the poles is 4.60 meters long × 1.80 meters wide × 1.20 meters high, the magnetic field is 10 kilogauss and the weight of the magnet is 650 tons. The magnet rests on the hydraulic pads which allow a 30° rotation about a pivot point located near the front end of the magnet. The origin of the MPS coordinate system is the pivot point (see Figure 2.1); z axis points along the beam; x axis is horizontal and points towards bending of the positive particles trajectories; y axis points up and magnetic field is directed against y axis.

2.3 The TPC Modules

The TPC system (Figures 2.4, 2.5, 2.6) consists of three separate modules having the shape of rectangular boxes. In order to maximize the acceptance, the beam is directed to the middle of the first chamber, which limits the maximum intensity of the useful beam. Each TPC module is 60 cm high × 65 cm wide × 47 cm long box filled with a low–diffusion gas mixture (see Section 2.5).

A uniform electric field parallel to the magnetic field of the MPS is created inside each TPC to induce a downward drift of the electrons produced by the ionization of the gas mixture by the charged particles to be measured. Since the HV power supplies could not hold the voltage of 30 kV the cathode voltage was set to 20 kV. The TPC’s are placed in the vertical magnetic field of 10 kGauss generated by the
Figure 2.3  The MPS Magnet.
MPS magnet. The electrons drift down in the superposition of $E$ and $B$ fields, while the parallelity of these fields suppresses the diffusion. The readout is achieved by the 36 rows (12/module spaced 3.8 cm apart in $z$-direction) of 256 short anode wires (2.54 mm spacing in $x$-direction), with the wire number giving the $(x, z)$ coordinates and the drift time, giving the $y$ coordinate for unambiguous 3D points along the tracks. Electrons produced by the charged particles in the drift region must pass through a gating grid and a wire cathode in order to reach the anode. A second, solid copper cathode, below the anodes helps to achieve the high gain (Fig 2.8).

The beam rate of about 1000 ions/spill (1 spill lasts 1 second) was chosen to compromise the number of triggers achievable during the very short running time given (two weeks including the beam commission time and various adjustments of DAQ) with the limitations on the beam intensity. The beam intensity limitations are defined by the effect of the positive ions build-up in the drift region and in the amplification region of TPC. For the ions, created in the amplification region, the effective gating grid was used to prevent the transfer of drifting electrons from drift region to the amplification region and that of positive ions in the reverse direction except for selected events. Gating is achieved by applying different potentials to adjacent wires as shown on Fig 2.8, collecting the electrons on the more positive wires. When the trigger system indicates that event of interest has occurred the gate wires are shorted together for 40 $\mu$s, a little longer than the maximum drift time of electrons (the drift velocity of $e^-$ is 1.6 $cm/\mu s$), allowing ionization from the tracks
Figure 2.4  Front view of module showing principle components.
to reach the amplification region. Because of their slower drift velocity, positive ions from amplification at the anode do not reach the gating grid before the gate is closed, further improving the gating of the TPC.

2.4 The Electronics

Readout electronics uses custom hybrids to produce a very compact, cost-effective system. The TPC's required electronics to give the time information only, since \( dE/dx \) is not measured in our system. In order to read out the many closely spaced short anode wires at minimum cost, a circuit was designed that allows each individual channel to be mounted directly on the chamber and to occupy only the anode wire spacing. The circuit for each channel is shown of Fig 2.7. Sixteen channels are assembled in a hybrid package (LeCroy HTD 161S, M) with inputs spaced 2.54 \( mm \) apart. Shaping–time constants were selected to match the anode waveform (Fig 2.7). An integration time constant of 50 \( ns \) and a differentiation time constant of 200 \( ns \) were chosen in order to minimize noise and maximize the two–track separation. Time information is recorded with memory driven by a counter to serve the function of a shift register. Each hybrid contains memory capable of recording 1024 time samples from each channel. To reduce a possible local noise pickup, all logic levels are ECL (Emitter Control Logic) to quickly release the charges accumulated in the base. The clock and address lines are terminated differential ECL.
Figure 2.5  View of three modules in operating position as viewed from upstream end of the MPS magnet.
Figure 2.6  View of the main electrical connections to modules.
To minimize the control circuitry two kinds of hybrids are used – MASTER and SLAVE. Master, which in addition to the amplifier–shaper–comparator and memory plus 16 channels has the counter used to address the memory plus the control logic necessary to read out itself plus seven slave hybrids. The slave hybrids contain only the amplifier–shaper–comparator and memory required for 16 channels. The printed circuit boards are assembled with one master and seven slaves (Fig. 2.9) which is sufficient for 128 anode readouts and is externally accessed by one clock line, one data line and 14 control lines. The data line is common to the two hybrids cards in a row and the other lines are common to the whole TPC. Thus only 27 lines are used to connect over 3072 channels per module to the external electronics. The time samples are obtained at 32 MHz and the data is sequentially read from each channel’s memory at 16 MHz *. This serial string of data bits is encoded into wire number, drift time and cluster size by FASTBUS modules in which all rows are encoded simultaneously. Consequently the entire readout is completed in 10 ms (dead time). This amount of dead time is acceptable because of the other limitations of the trigger rate. The on–chamber performance of the hybrid allows input sensitivity of less than 1 μA.

*The clock frequency was adjusted to the cathode voltage of 20 kV in order to maximize the fiducial volume of the TPC.
2.5 The Gas System

In order to have good event reconstruction in our high multiplicity events it is important to optimize the two-track resolution by having a large number of independent readout elements. We have used rows of short (≈ 1 cm) anode wires parallel to the beam direction. On the other hand, use of short anode wires precludes using $dE/dx$ for particle identification, which in the kinematic region of interest and at the particle densities in these events is a very expensive and extremely difficult task, and probably would have limited overall efficiency.

It is also essential that the gas used in the chambers be chosen to give the stable operation, safety and long chamber life. Various gas mixtures were studied for this purpose. We chose a mixture of a low-diffusion gas (79% Argon, 16% Isobutane and 5% dimethoxymethane) known to be stable at high gain in order to achieve this fine segmentation.

Tests made on these gases give good efficiency characteristics and linearity of the drift time versus the drift distance. In addition, this gas mixture has the further advantage of not being able to sustain combustion by itself — a very important consideration for such a large system. However, the high Argon content will result in a high rate of multiple scattering.
2.6 The Au Beam

Experiment E891 utilized the 11.6× A GeV/c Au beam provided by the AGS as shown in Figure 2.10. Au ions are prepared by Tandem Van De Graaf Heavy Ion Transfer Line (HITL) before being injected into AGS. The Au ion beam is transported to the experimental area after it was accelerated to the momentum of 11.6×A GeV/c. The beam intensity used ranged from 500 to 2000 Au ions per second.

2.7 The Trigger

The plan view of the E891 apparatus is shown in Figure 2.1. The Au beam at 11.6 × A GeV/c is incident on a target just upstream of a set of three TPC modules. Au beam is identified by a plastic Cherenkov counter C1 upstream of the target (Figure 2.2) †. The beam passes through the middle of the TPC's, which are used to reconstruct the angles and momenta of charged particles produced in the forward hemisphere of NN center of mass system (y ≥ 1.6). The main trigger element is a scintillation counter S6 in the beam downstream of the MPS magnet, which is used to reject events in which fragments with Z > 6 remain in the beam region (the pulse height in a scintillation counter is recorded for every event so that tighter cuts can be applied by the software). The counter S4 just upstream of the TPC is used to select interactions from the region of the target. Multiplicities of reconstructed tracks often

†The beam particles have the same momentum, Cherenkov counter determines Z^2 to sort out light ions.
exceed 200, with an estimated reconstruction efficiency >90% for long tracks. There
are no $dE/dx$ measurements in the TPC.
Figure 2.7  Scheme of individual electronics channel.
Figure 2.8  Schematic arrangement of electrodes in the detector end cap.
Figure 2.9 The Layout of a 128-channel hybrid printed circuit boards (one master and seven slaves).
Figure 2.10  The AGS plan view.
Chapter 3

Data Reduction and Analysis

3.1 Raw Data Analysis

The raw data analysis is the most CPU time consuming part. At this stage the data, recorded in the raw format containing number of the wire and the number of the time bin is being converted to the form suitable for the physical analysis. Before this conversion, the experimental apparatus should be calibrated. Step by step raw data analysis consists of the following procedures:

- Calibrations of TPC modules:
  - $x$ offsets and beam angle $\Theta_B$ calibrations,
  - the drift velocity $v_d$ and $y$ offsets calibrations,
  - $\vec{E}$ and $\vec{E} \times \vec{B}$ distortions calibrations,
  - the hit efficiency calibrations;

- Off-line processing:
  - production of Data Summary Tapes,
  - production of "N-tuples".
Since the TPC's are placed on the metal support their coordinates change slightly in time following the change of temperature. In addition the drift velocity inside the TPC modules is affected by the outside atmospheric conditions. The value of these parameters as a function of time is being determined by the calibrations. The $\vec{E}$ and $\vec{E} \times \vec{B}$ field distortions are caused by edge effects resulting in nonuniform $\vec{E}$ field and hence, a nonuniform drift velocity map. As a result, before the corrections are applied, the hits are reconstructed in the positions shifted relative to the true ones. We will return to this in more detail later in subsection 3.1.2. Since the detector calibrations are based on the information from the Pattern Recognition, we devote the next subsection to this topic.

3.1.1 Track Reconstruction Algorithm

The data is first translated from MPS DAQ format to the FASTBUS format and the latter is decoded as the drift time, cluster size and wire number via the Local Pattern Recognition algorithm, giving the 3D coordinates of the hits. i.e. ionization clusters. At the second stage the track reconstruction algorithm associates individual hits with the tracks. This phase of Global Pattern reconstruction starts with the hits found in the last row of the last TPC module and progresses upstream, first reconstructing the helical track segments in each TPC and then associating the segments from different modules into the global tracks. The parameters of these found tracks are further fitted for the minimum deviation squared from the corresponding
hits. The tracks with less than 8 hits are discarded, since it is impossible to reconstruct parameters of such tracks with sufficient accuracy.

Among the tracks there are those coming from the target and the tracks due to the decay products of the neutral particles. Among the latter ones suitable for reconstruction is the decay $\Lambda \rightarrow p\pi^- \ (BR = 69\%, \ c\tau = 7.9 \ cm)$. Another decay $K^0 \rightarrow \pi^+\pi^- \ (BR = 64\%, \ c\tau = 2.7 \ cm)$, is also visible in our apparatus, though with smaller acceptance because of the shorter lifetime and relative softness of the positive pion as compared to the proton. In addition to these possibilities, due to the multiple scattering of charged particles in the detector, the pattern recognition algorithm finds tracks which neither can be projected to the target nor can be associated with the charged decay products of neutral particles from the target.

3.1.2 Detector Calibrations

In order to perform the kinematical analysis, the pattern recognition algorithm applied to the event by event track and vertex reconstruction requires the experimentally determined parameters, such as positions of the TPC modules and the target, the field distortions in the drift volume of the TPC as a function of the uncorrected hit position and the time dependence of the drift velocities and vertical offsets of the TPC modules.

We have used the tracks created by the charged particles, in order to perform the calibrations. Obviously to make adjustments one needs to measure the parameters
of the tracks using tools other than the TPC's. We employed four MPS drift chambers located downstream of the TPC's in order to have independent measurement of the track parameters. The beam tracks were independently defined by the two proportional wire chambers placed upstream of the target (see Chapter 2).

The TPC calibrations required for the track reconstruction include distortion calibrations, alignment, drift velocity and offsets calibrations repeated iteratively until the determined parameters converge to the true values. Upon completion of the calibrations above, we perform the efficiency calibrations. The hit reconstruction efficiencies are needed as an input for the detector simulation program used for acceptance calculation (see subsections 3.2.4 and 3.2.5).

**Distortions.** Electrons in the drift region move in the superposition of the electric and magnetic fields along trajectories described by the following equation of motion:

$$\frac{d\vec{r}_e}{dt} = -\mu_e \vec{E}(\vec{r}_e) + \mu_e^2 [\vec{E}(\vec{r}_e) \times \vec{B}(\vec{r}_e)], \quad (SI) \quad (3.1)$$

where $\mu_e$ is mobility of electrons, defined as

$$\vec{v}_d(\vec{r}) = -\mu_e \vec{E}(\vec{r}) \quad (3.2)$$

for drift of electrons without magnetic field*. Due to the nonuniform electric field $\vec{E}(\vec{r})$, the drift velocity $\vec{v}_d(\vec{r})$ is coordinate dependent. This affects the pattern recognition algorithm in such a way that the hits are reconstructed in positions shifted rel-

---

*Drift velocity has a linear dependence on electric field applied for the field intensities used in our experiment [86].
atively to their true positions, before appropriate corrections are applied. Therefore the distortion calibrations are necessary in order to take into account the spatial variations of drift velocities inside each TPC module. Spatial variations of the drift velocity are small compared to the mean vector

\[ < \bar{v}_d > = \frac{1}{W_{TPC} \cdot H_{TPC} \cdot L_{TPC}} \int \int \int \bar{v}_d(\vec{r}) d\vec{r}, \]  

(3.3)

which, in turn, varies with time from run to run. We will describe the procedure of mean drift velocity, \( < \bar{v}_d > \), calibration in the subsequent paragraph (Drift Velocities and Vertical Offsets). The distortions calibration is a part of an overall iterative calibration procedure. The distortions in all three directions \( x, y \) and \( z \) are important and have to be determined independently. The values of the distortions are maximal at the edges and minimal for the middle rows of the TPC's. The real tracks in the full magnetic field of 10 \( kGauss \) were used for the fitting procedure.

Transverse \( x \) and \( y \) distortions were fitted using the upstream calibration data, i.e. events triggered on interactions of the beam particle upstream from the target with the full magnetic field applied. The tracks found in all 3 TPC's were used. Because of the small \( \theta \) angles of the tracks in the upstream calibration data, these tracks allow one to look for basically transverse distortions, not longitudinal (figure 3.1). The fitting procedure was started from the last TPC module closest to the DC's. Namely, the deviations of hits in the last TPC from the tracks defined by the first and second modules and the drift chambers were fitted, then the fitting procedure has
been repeated for the second module on the basis of data from DC's and the other TPC modules and finally for the first one. It took a few steps of iteration before the stable values of the $x, y$ distortion were found.

The distortions are three dimensional and to determine the longitudinal $z$ components one needs tracks forming large angles with respect to the beam axis, which can be only found in the closest in the TPC module closest to the target. Since distortions for the middle rows of the TPC are small, these rows were used to fix the coordinates of the tracks for further fitting the residual $y$ deviations, $\Delta y_{res}$, of hits after $xy$ correction were done. Then, the $z$ distortion, $\Delta z$, is defined as follows:

$$\Delta z = \frac{\Delta y_{res}}{\tan \theta},$$

(3.4)

which is illustrated on Figure 3.1. For the other modules, $z$ distortions were taken to be the same as for the first one, since the angles of tracks crossing these modules are small.

Significant deterioration of efficiency of the track reconstruction is caused by the dynamic distortions due to the intensive positive ion buildup in the beam region. For evaluation of the effect let us recall that for argon the mean energy loss of a minimum ionizing particle at normal temperature and pressure is of $(dE/dz) = 1.519 \text{ } MeV \cdot g \cdot cm^2$ and the mean energy of ion–electron pair creation is $V_0 = 26 \text{ } eV$ [78]. Thus with the density of argon, $\rho_{Ar} = 1.7 \times 10^{-3} \text{ } g \cdot cm^{-3}$, the average number of electron
ion pairs produced by one gold ion is

\[ \frac{dN_i}{dz} \approx \frac{79^2}{V_0} (dE/dz) \approx 6.1 \cdot 10^5 \ \text{cm}^{-1}. \]  

(3.5)

We had a typical beam intensity of 1000 ions/spill, which corresponds to a mean duration between two consecutive ions of \( t_i = 1 \) ms. Studies of the beam caused distortions by experiment E810 \cite{87} have shown that an ion sheet develops in the drift region for about 300 ms, the time comparable to the drift time of the positive ions from the beam up to the cathode.

Let us estimate the electric field, \( \Delta E_{dyn} \), and the corresponding \( x \) distortion, \( \Delta x_{dyn} \), assuming that the positive charge is uniformly distributed over the \( yz \) plane.

The surface charge density can be evaluated as follows:

\[ \frac{dq}{dz} = e(dN_i/dz) = 9.8 \cdot 10^{-12} \ C \cdot m^{-1}, \]  

(3.6)

\[ v_d' \approx \frac{H_{TPC}}{(2 \cdot 300 \ \text{ms})} \approx 1 \text{ m/s}^{-1}, \]  

(3.7)

\[ \frac{dq}{ds} = \frac{(dq/dz)}{(v_d' \cdot t_i)} = 9.8 \cdot 10^{-9} \ C \cdot m^{-2}. \]  

(3.8)

The electric field caused by the positive ion cloud is pointed in \( x \) direction. The estimate for its value is:

\[ \Delta E_{dyn} = \frac{1}{\varepsilon_0} \frac{dq}{ds} \approx 10 \ V \cdot \text{cm}^{-1}. \]  

(3.9)

Since the distortion field is small compared to the applied vertical field of \( E_0 = 330 \ V \cdot \text{cm}^{-1} \), the corresponding \( x \) distortion for the hits in the beam region on top
of TPC can be evaluated as follows:

\[
\Delta x_{dyn} = \frac{\Delta E_{dyn}}{E_0} \cdot (H_{TPC}/2) \approx 1 \text{ cm.}
\]

The above estimate of the \(x\) distortion agrees with experimental observation.

The electric field inside a TPC can be represented as sum of a constant vector and a small distortion field, which, in turn is a sum of a static and dynamic components:

\[
\vec{E}(\vec{r}) = \vec{E}_0 + \Delta \vec{E}_{stat}(\vec{r}) + \Delta \vec{E}_{dyn}(\vec{r}, \tau).
\]

The positive ion cloud charge density evolves on the time scale comparable to the duration of a spill (300 ms versus 1 s). Since a trigger time, \(\tau\), is randomly distributed relative to the beginning of the spill and the number of beam particles per spill fluctuates, the distortion map for any event is statistical. For distortion calibration we averaged the deviations of hits from tracks over a large sample of events. As a result, the obtained distortions are uncertain by a random value proportional to the dynamic part of the distortion:

\[
\Delta \vec{E}_{dyn}(\vec{r}, \tau) = \frac{1}{t_{spill}} \int_0^{t_{spill}} < \Delta \vec{E}_{dyn}(\vec{r}, \tau') > d\tau',
\]

the part caused by the positive ion cloud. The averaging over all beam intensities is performed in (3.12). We will approximately take this effect into account in the Monte–Carlo simulations of apparatus effects (see subsection 3.2.4).

One should note also that the \(\vec{E} \times \vec{B}\) part of the distortion (see equation (3.1)) relates to that of \(\vec{E}\) by the factor of \(B\mu\), which for the values of \(B = 1 \text{ T}\) and
\( \mu \approx \frac{v_d}{E_0} = 0.5 \text{ m}^2/(\text{V} \cdot \text{s}) \), is approximately 0.5. Thus, the cross product distortions are comparable with electric field distortions.

**Alignment.** The alignment runs, serving for the \( x \) offset and the beam angle determination, are the runs without a target and with the magnetic field off. The straight beam tracks are fitted using the hits left in four drift chambers and two PWC's. The beam angle and the \( x \) offsets of the TPC's are determined to give the best approximation to the beam tracks defined by the hits in the TPC modules.

**Drift Velocities and Vertical Offsets.** In order to calibrate the drift velocities, \( v_d \), and offsets in \( y \) direction, \( y_{off} \), in each TPC, we used the Au target runs with full magnetic field (10 kGauss). The problem is that the above parameters are in strong negative correlation with each other and due to the unusually large dynamic distortions in the beam region, the corresponding \( \chi^2 \) resulting from fitting a large number of tracks going through the various regions of the TPC's does not have the unique minimum in the parameter space \( (v_d^{(1)}, y_{off}^{(1)}, v_d^{(2)}, y_{off}^{(2)}, v_d^{(3)}, y_{off}^{(3)}) \).

The procedure of \((v_d, y_{off})\) calibrations was to minimize the deviations of the hits from the tracks measured via the drift chambers. Because of the large distortions it is important to choose the right first approximation for the drift velocities and offsets so that this approximation converges to the true values of above parameters. An independent criterion, which has been used for choosing the first approximation is a position of the peak of the effective mass \( M_{\pi-p} \) distribution – it should match the
DISTORTION CORRECTIONS

Figure 3.1 Distortion corrections. Using large angle tracks from the target for $z$ distortion correction, part A), and upstream calibration data for $y$ distortion correction, part B).
true value of the $\Lambda$ mass. Kinematically there is a dependence of the position of the maximum in the peak of above distribution on $v_d$ used for its calculation. Should $v_d$ be overestimated, it leads to the shift of the target image coordinate $x_t$ in the positive direction and, hence, to the overestimation of the slope angles (angles of tracks with the horizontal plane $xz$). Slope angle overestimation, in turn, results in the shift of the $M_{\pi\rho}$ distribution peak to the higher values of effective mass. Once the first approximation has been chosen, the iterative procedure was repeated several times until the corrections became very small.

It is important to note that due to the overload of the drift chambers by a large number of tracks we selected low multiplicity events with $N_{ch} < 100$, in order to perform the above iterations. The drift velocities and offsets are shown on Figures 3.2 and 3.3 as a function of the date (day number in October).

**Hit Efficiencies:** The hit efficiencies are calculated for each TPC row as functions of $x$ and $y$ coordinates. Upstream calibration data was used. The $xy$ planes for each row are divided into a grid $16 \times 16$ and the numbers of tracks $N_{tracks}$ together with the number of clusters (hits) associated with these tracks $N_{hits}$ are counted in each of 256 rectangles. The efficiency is defined as the ratio:

$$E_{hit}(i, j) = \frac{N_{hits}(i, j)}{N_{tracks}(i, j)},$$

(3.13)

where numbers, $i$ and $j$, correspond to the $x$ and $y$ coordinates, respectively. The tracks having at least 18 hits and length exceeding 100 cm were used for the calcula-
Figure 3.2 Drift Velocity Calibrations Results.
Figure 3.3  Offset Calibration Results.
tion. This selection was necessary in order to discard numerous short tracks, present due to the large dynamic distortions, which are absent in the Monte–Carlo simulation.

On Figures 3.4 and 3.5 the hit efficiency is shown as a function of \(x\) and \(y\) coordinates respectively. One can see the deterioration of the hit efficiency in the region of the beam. One should not confuse the hit efficiencies with the track reconstruction efficiencies. In fact the track reconstruction efficiency in the beam region becomes as low as 40% (see subsection 3.2.6).

3.1.3 \(V_0\) Selection Criteria

\(V_0\)'s are the vertices formed by two tracks left by particles of the opposite charges. In order to be reconstructed, the point of the tracks intersection, – vertex, need not necessarily be located inside one of the TPC modules but is required to be far enough from the target (15 cm downstream and more) to distinguish from the production vertex. In fact, because of the small lifetimes of \(\Lambda\) and \(R_S^0\) (\(ct < 8\) cm), most of them decay between the target and the first TPC.

The search for \(V_0\)'s starts when the track reconstruction is finished for a given event. At this stage the answer is known whether the particular track does or does not originate in the target vicinity. Namely, using the well measured 3D map of magnetic field, \(\vec{B}\), in the MPS magnet, each track is extrapolated back to the target region. First the production vertex is reconstructed: the coordinate \(z_T\) is found such that the dispersion of the \(x\) and \(y\) coordinates around the central point is the minimal,
Figure 3.4  The hit efficiency as a function of $x$ coordinate for the first row of the first TPC module. $j = 9$ in formula (3.13).
**Figure 3.5** The hit efficiency as a function of y coordinate for the first row of the first TPC module. \( i = 9 \) in formula (3.13).
i.e.

\[
\sum_{i}^{N_{tr}} \left[ \frac{(x_i - x_{cen})^2}{\sigma_{x}^2} + \frac{(y_i - y_{cen})^2}{\sigma_{y}^2} \right]_{z=const} \rightarrow \text{min},
\]  \hspace{1cm} (3.14)

where \( x_{cen} = (\Sigma_i x_i)/N_{tr} \), \( y_{cen} = (\Sigma_i y_i)/N_{tr} \). These values for \( z = z_T \) are taken as the \((x_T, y_T)\) coordinates of the production vertex. The value thus determined is consistent with measured target position \( z_T = 2.5 \text{ cm} \). The peripheral tracks with \(|(\Delta x, \Delta y)| > 7 \text{ mm} \) are considered as possible daughters of \( V_0 \)'s; this selection is required in order to reduce combinatorial background consisting of fake vertex combinations formed by the tracks from the target. For each "positive" track from this category the "negative" track is searched for among the tracks–candidates with the distance of closest approach not exceeding 3 mm. If the criterion of the closest approach is satisfied, the effective mass,

\[
M^2(\pm, -) = (p_+ + p_-)^2 = m_+^2 + m_-^2 + 2E_+E_- - 2p_+p_-,
\]  \hspace{1cm} (3.15)

of two particles forming the vertex combination is calculated. The following three assumptions about the particles masses:

\[
m_+ = m_p, \hspace{0.5cm} m_- = m_\pi,
\]  \hspace{1cm} (3.16)

\[
m_+ = m_\pi, \hspace{0.5cm} m_- = m_\pi,
\]  \hspace{1cm} (3.17)

\[
m_+ = m_\pi, \hspace{0.5cm} m_- = m_\rho,
\]  \hspace{1cm} (3.18)

are used to test the three decay hypotheses: \( \Lambda \rightarrow p\pi^- \), \( K^0_S \rightarrow \pi^+\pi^- \) and \( \bar{\Lambda} \rightarrow \bar{p}\pi^+ \).

At least one of the following conditions should be met:

\[
1.08 \text{ GeV}/c^2 < M_{(3.16)} < 1.16 \text{ GeV}/c^2,
\]  \hspace{1cm} (3.19)
\[ 0.43 \text{ GeV}/c^2 < M_{(3,17)} < 0.57 \text{ GeV}/c^2, \]  
(3.20)

\[ 1.08 \text{ GeV}/c^2 < M_{(3,18)} < 1.16 \text{ GeV}/c^2, \]  
(3.21)

in order for the (+, −) couple to be put to the vertex bank.

Since we do not measure the particle mass in our apparatus, the only way to identify the type of the decay is by looking at the effective mass distribution of the decay products, which should exhibit a peak at the true value of the neutral decaying particle mass.

The the cuts above listed are basically all used during of the raw data analysis. These cuts have intentionally been made rather loose in order to enable one to look for the possible systematic errors on the later stage of the data reduction. We will return to the selection criteria for the vertices in subsection 3.2.2.

3.1.4 Production of the Data Summary Tapes

At this stage we are all set with the parameter files resulting from the calibrations, containing the \( \vec{E} \) and \( \vec{E} \times \vec{B} \) distortion map as a function of \((x, y, z)\), drift velocity, \(v_d\), and \(x\) and \(y\) offsets for each individual run. The new task is to convert the raw data, recorded during data taking, into data files containing the parameters of the reconstructed tracks and vertices.

The large multiplicity of central \(Au + Au\) events dictates a large volume of information, therefore the time required for event processing is relatively long. In order to speed up the processing we employed four processors of SGI computers at BNL to
process our data in parallel mode. Namely, we ran four analyzers (analysis programs) simultaneously and had the Event Dispatcher reading the raw data from the files and distributing the spill records sequentially to the analyzers according to the scheme shown on Figure (3.6).

The Event Dispatcher (ED) reads a number of spill records from the raw data files and establishes connection with running analyzers through the communication port. As soon as connection is established, the spill records are sequentially sent to all available analyzers (A1, A2, A3 and A4). The ED periodically checks for one of the analyzers to complete the processing of the spill record – at this point another spill record is sent by ED to the analyzer and erased from the ED memory. This way, of course, there is no data word processed twice by different analyzers.

3.2 Data Reduction

When the data is converted to the Data Summary Files containing the parameters of all reconstructed tracks for each recorded event, it is easy to transform this information into the kinematical quantities of interest. We write out kinematical variables into a row-wise "N-tuple" on event by event basis *.

This stage of "N-tuple" production is relatively fast. The decision about the list of quantities to be stored in "N-tuple" is based on the task to be solved. Since we are

*"N-tuple" is a data file in compact Zebra format of CERN Library, accessible via Physics Analysis Workstation (PAW) and HBOOK packages.
Figure 3.6 Event Dispatcher and Multiprocessing Scheme.
particularly looking for the $\Lambda$ double differential spectrum, the variables of interest are the parameters of tracks forming the $V_0$ vertices such as

- 3–momenta of the positive and negative particles–daughters,
- track lengths,
- number of hits on the tracks,
- effective mass of $V_0$ calculated under $\Lambda \to \pi^- p$ decay hypotheses,
- 3D coordinates of the $V_0$,
- 3D coordinates of the production vertex ($Au + Au$ interaction point),
- number of tracks on the production vertex – charged multiplicity $N_{ch}$.

The number of particles emerging from the target is needed for the selection of the most central events, possessing the largest multiplicities. Monte-Carlo studies show that the impact parameter $b$ is in negative correlation with the multiplicity. Since the multiplicities often exceed 200, $N_{ch}$ is a good statistical number for the selection of the central events. Information contained in a vertex and track "N-tuples" is described in more detail in Appendices B and C respectively.

3.2.1 Selection of Tracks and Momentum Resolution

The selection of tracks to be considered in the analysis is based on the accuracy of their measurement. The TPC allows one to reconstruct the tracks parameters
based on the coordinates of the hits spaced by finite distances. The finite distances between the hits together with the limited track length limit the accuracy of the track measurement.

The calculation of the particle's momenta is based on the measurement of the tracks curvatures and the slope angles. For example,

$$p_{zz} = \frac{eB}{c} r_{zz} = \frac{r_{zz}}{\alpha},$$  \hspace{1cm} (3.22)

where the radius of curvature in the $xz$ projection is measured in $cm$, the particle momentum, projected to $xz$ plane is in $GeV/c$ and $\alpha = 333.80$ $cm$ $(GeV/c)^{-1}$. The vertical momentum component is defined as

$$p_y = \tan(\theta)p_{zz},$$  \hspace{1cm} (3.23)

where $\theta$ is the slope angle of the particle trajectory relative to the $xz$ plane.

Given the limited track lengths and curvatures of the tracks, one can see that the value associated with the accuracy of the track's curvature measurement is the saggita (the height of the circular segment inside the three TPC modules, Figure 3.7):

$$h \approx \frac{R \varphi^2}{8} = \frac{L_z^2}{8\alpha p_{zz}} \propto \frac{L_z^2}{p},$$  \hspace{1cm} (3.24)

where $L_z$ is the track length in $z$ projection and $\varphi$ is the angular size of the circular segment.
We select tracks with saggitas greater than \( h > 7.5 \text{ mm} \), which also implies the selection of the long tracks: \( L_z > \sqrt{2000 \cdot p_{xx}} \approx 40 \text{ cm} \). Since we have short anode wires (\( l_z = 1 \text{ cm} \)) the relative error of the track length measurement (\( l_z/L_z \approx 3\% \)) is rather small compared to saggita measurement error \( \Delta h/h \approx 10\% \) and can be neglected. Thus the momentum resolution for the tracks is:

\[
\Delta p \propto \frac{p^2}{L_z^2} \cdot \Delta h,
\]

which is dependent on the track length. Monte–Carlo studies show that momentum resolution for the full length tracks is \( \Delta p = 0.01 (\text{GeV/c})^{-1} p^2 \).

### 3.2.2 Selection of Vertices

A few preliminary selection steps for \( V_0 \) have been described in subsection 3.1.3. In order to sharpen the momentum resolution for the measured vertices and eliminate the contamination by the combinatorial background, a number of additional cuts needs to be applied. Let us particularly concentrate on the \( \Lambda \) hyperon measurements.

The way the \( \Lambda \) double differential spectrum \( d^2N/dydm_{\pi^-}^2 \) is measured is based on the observation of a peak in the invariant mass spectrum of \( \pi^-p \) pairs (to be precise \( ++ \) pairs, where positive particle is considered to be a proton and the negative particle is thought to be a \( \pi^- \)).

It is known that \( \Lambda \) hyperon has a strangeness \( S = -1 \) while the products of its decay are non–strange baryon and meson. Because of the strangeness conservation in
Figure 3.7 Illustration to equation (3.24) for the saggita. On the Figure $h$ denotes the saggita, $R$ is the radius of circular projection of the track to $xz$ plane and $L_z$ is the track length inside the detector. The box corresponds to the three TPC's volume.
strong interactions, \( \Lambda \) decays via the weak interaction and the corresponding lifetime is \( \tau_{\Lambda} = 2.632 \pm 0.020 \times 10^{-10} \text{ s} \), corresponding to the very small natural width, \( \Delta M(\pi^-p) \sim \hbar/\tau_{\Lambda} \approx 2.5 \mu\text{eV} \). The width of the measured invariant mass distribution is dictated by the experimental momentum resolution and amounts to \( 10 - 12 \text{ MeV} \).

The momentum resolution for the \( \Lambda \)'s is determined by the accuracy of measurements for tracks formed by the decay products of the \( \Lambda \)'s. We therefore apply the same sagitta cut as described in section 3.2.1 to select the final \( \Lambda \) sample for better momentum resolution.

The coordinates of the production vertex are limited in the cube

\[
-3.5 < x_{p.v.} < 3.5,
-3.5 < y_{p.v.} < 3.5,
0.0 < z_{p.v.} < 7.0,
\]

in order to reduce the contribution from background interactions.

Due to the large multiplicities in \(^{197}\text{Au} + ^{197}\text{Au} \) central events there is a large probability for the particles escaping the target, mostly protons and pions, to approach close enough to each other on their way to the TPC's, so that our pattern reconstruction program confuses these crossing tracks with the real \( \Lambda \) decays. We therefore should expect an increased background as compared to the collision of light nuclei in experiment E810. To partially reduce this background we require that the \( \Lambda \) candidate momentum points to the production vertex, namely the sum of momenta
of the two intersecting tracks is allowed to miss the production vertex coordinate \((x_T, y_T)\) by no more than 2.5 \(mm\).

Unlike the true \(\Lambda\)'s, combinatorial background can only occur for certain orientation of tracks in space. Namely, let us define the decay plane by the point of the two tracks intersection, \(\vec{V}_0\), and the two vectors – momenta of the charged positive and negative particles at the above point, \(- \vec{p}(+)\) and \(\vec{p}(-)\). Instead of having the two momenta in the definition of the decay plane it is sufficient to know only the unit vector in the direction of their cross product, \(\vec{n}_0 = (\vec{p}(+) \times \vec{p}(-))/|\vec{p}(+) \times \vec{p}(-)|\) – the normal vector to the decay plane. In a spherical coordinate system the normal vector can be expressed through the polar angle, \(\theta\), and the azimuthal angle, \(\Phi_{\rho p}\), so that

\[
\vec{n}_0 = (n_{0x}, n_{0y}, n_{0z}) = (\sin \theta_{\rho p} \sin \Phi_{\rho p}, \sin \theta_{\rho p} \cos \Phi_{\rho p}, \cos \theta_{\rho p}). \tag{3.27}
\]

Given the aperture of our apparatus, the polar angle \(\vec{n}_0\) encountered among the reconstructed vertex combinations is close to 90\(^0\), since we are measuring the particles going forward in the lab system, while the azimuth, \(\Phi_{\rho p}\), can range in the full region from \(-180^0\) to \(180^0\). Let us express \(x\), \(y\) and \(z\) projections of the vector \(\vec{n}_0\) in terms of respective components of momenta, \(\vec{p}(+)\) and \(\vec{p}(-)\):

\[
\begin{align*}
\vec{n}_{0x} &= [p_y(+)p_z(-) - p_y(-)p_z(+)]/|\vec{p}(+) \times \vec{p}(-)|, \\
\vec{n}_{0y} &= [p_z(+)p_x(-) - p_z(-)p_x(+)]/|\vec{p}(+) \times \vec{p}(-)|, \\
\vec{n}_{0z} &= [p_x(+)p_y(-) - p_y(-)p_x(+)]/|\vec{p}(+) \times \vec{p}(-)|. \tag{3.28}
\end{align*}
\]
Since the radii of the measured tracks \((p > 1 \text{ GeV/c})\) are large compared to the length of the detector (see (3.22)), the trajectories of the tracks projected to \(yz\) plane are almost straight lines, collinear with the respective projections of the particles momenta. Positive and negative particles should approach each other as close as 3 \(mm\) or less, in order to pass the vertex cut (see subsection 3.1.3), therefore in \(yz\) projection their momenta should be almost collinear:

\[
\frac{p(+)_y}{p(+)_z} \approx \frac{p(-)_y}{p(-)_z},
\]  

(3.29)

which, according to (3.28), makes the projection \(n_{0x}\) be small compared to unity. From (3.27) this can happen for \(\Phi_{\pi p}\) around 180° or 0°. In case of \(\Phi_{\pi p} \approx 0^\circ\) however, the tracks bend away from the target and therefore can not be found in abundance.

The topology of the "fish tail" background combination is illustrated on figure 3.8. The experimental distribution of observed \(\Lambda\) candidates over the angle, \(\Phi_{\pi p}\), and effective mass \(M_{\pi p}\) is shown on Figure 3.9. This distribution exhibits large wings of combinatorial background around \(\Phi_{\pi p} = \pm \pi \ \text{rad}\) which overshadow any possible real \(\Lambda\) signal in this region, while in the region of \(\Phi_{\pi p}\) around 0° the two–dimensional histogram has a maximum at the true value of the \(\Lambda\) mass, \(m_\Lambda = 1.1156 \ \text{Gev/c}^2\), and relatively good signal to background ratio.

### 3.2.3 Selection of Central Events

In order to select the most central events we have employed the charged multiplicity cut, \(N_{ch} > N_{ch}^{cut}\), which corresponds to the central event sample with integrated
Figure 3.8 Illustration for the "fish tail" background combination. Note that for the tracks from the target the normal vector to the decay plane $\mathbf{n}_0$ should point down or have the $\Phi_{\pi}$ angle close to $180^\circ$. 
Figure 3.9  The Azimuthal Angle Dependence of $M_{\pi^-p}$ Distribution.
cross section $\sigma_{cen} = 270\ mb$. The number of the central events, $N_{cen}$, and the cut-off value, $N_{ch} = 220$, has been found using the following relation:

$$
\sigma_{cen} = I_{trig} \cdot \left( \frac{N_{cen}}{N_{trig}} \right) \cdot \frac{A_{Au}}{\rho_{Au} w_T N_{Au}}.
$$

(3.30)

We have obtained $N_{cen} = 14114$ corresponding to the integrated cross section above while the total number of triggers recorded was $N_{trig} = 348403$. The trigger rate, $I_{trig} = N_{trig}/N_{beam}$, was 1.01% with the negligible statistical fluctuations from run to run. The target thickness was $w_T = 0.254\ mm$. Atomic number of gold is $A_{Au} = 197$ and the density is $\rho_{Au} = 19.3\ g\cdot cm^{-3}$.

The geometric cross section for $A_1 + A_2$ collision is defined as

$$
\sigma^{geom}_{A_1+A_2} = \pi \left[ r_0 \cdot \left( A_1^{1/3} + A_2^{1/3} \right) \right]^2,
$$

(3.31)

where $r_0 = 1.2\ fm$. For $^{197}Au + ^{197}Au$ interactions, $\sigma^{geom}_{Au+Au} = 6.127\ b$. We therefore measure approximately 4.4% of the geometric cross section.

For further comparison of our experimental results with cascade model calculations let us evaluate the geometric cross section corresponding to an impact parameter cut, $b_{imp} < b$:

$$
\sigma^{geom}(b) = \pi b^2.
$$

(3.32)

We obtain the cross section of 283\ mb for the cut, $b = 3\ fm$, which is very close to our measured cross section after the multiplicity cut specified earlier.
3.2.4 Monte Carlo Simulations of the Apparatus Effects

The four-momenta of all particles produced in the Au + Au collisions were simulated using the HIJET event generator with additional strange particles. We used 25,000 events with the strangeness yield normalized to the charged kaon yield from experiment E866. The shapes of rapidity and transverse mass distributions for $K^0_S$ and $\Lambda$ particles were defined to be the same as in E810 Si+Si data [88], i.e., they were generated according to the global fits:

\[
\frac{d^2 N_{K^0_S}}{dy dm_t^2} \propto \exp \left[ -(3.015 + 2.248 \cdot \cosh(y - y_0)) \cdot m_t \right],
\]

\[
\frac{d^2 N_{\Lambda}}{dy dm_t^2} \propto \exp \left[ -(3.784 + 1.035 \cdot \cosh(y - y_0)) \cdot m_t \right],
\]

where $y_0 = 1.6$ is midrapidity corresponding to the beam momentum, $p_0 = 11.6 \ GeV/c$.

The simulation of the known apparatus effects was performed by means of program GEANT, which performs the tracking in medium, simulates the decays of particles with experimentally determined branching ratios, allows one to define the detector positions, sizes and to describe the detailed structure of detecting elements.

The parameters used in simulations include:

- hit efficiencies from our "efficiency calibration" runs,

- cluster size distributions derived from the data,

- hit resolution based on deviations of hits from fitted tracks,

- tracks merging based on above,
• drift velocity variations with $\sigma = 1\%$,

• distortion correction miscalibration based on adding a distortion scaled by the measured distortion with the scale factor determined randomly event-by-event with $\sigma = 30\%$,

• distortion correction miscalibration based on adding a distortion scaled by the measured distortion with the scale factor determined randomly row-by-row with $\sigma = 20\%$.

After propagating tracks in the TPC's, the code digitizes hits according to the TPC readout elements positions, clock frequency and experimental cluster size distribution. † The TPC electronics output is simulated via the software package called TPSIGNAL. The input of TPSIGNAL is coordinate and momentum vector of track intercept of TPC detector plane. The package generates Landau-distributed charge in the detector gap, diffuses charge, and produces digitized electronics readout. The background hits are added to all TPC rows. A constant noise signal over all buckets due to sinusoidal pickup is added to the anode signal. The resulting Monte-Carlo data is written out in the FASTBUS format in order to further be analyzed by the same tracking program as the real data.

†Cluster is a superposition of hits read out on adjacent wires which overlap in time. Physically, the cluster is a collection of electrons resulting from the primary and secondary ionization of the gas in the drift volume via relativistic charged particles to be tracked.
3.2.5 Acceptance Correction and Error Analysis

We ultimately need to measure the double differential spectrum \((1/N_{\text{ev}})d^2N/dydm_t^2\) for the \(\Lambda\) and \(\pi^-\) production as a function of rapidity and transverse mass:

\[
y = \frac{1}{2} \ln \left( \frac{E + p_t}{E - p_t} \right), \quad m_t = \sqrt{m_i^2 + p_t^2},
\]

in the region of the phase space, \((y, m_t)\), visible by our detectors. Due to the limited efficiency of the track and vertex reconstruction and the limited aperture of the detectors, the number of particles counted in each \((\Delta y, \Delta m_t)\) bin is much smaller than the actual number of particles occurring to be in this bin \(^1\). We define the acceptance in a usual way:

\[
A(y, m_t) = \frac{d^2N_{\text{DATA}}(y, m_t)}{dydm_t^2} \Big/ \frac{d^2N_{\text{Au+Au}}(y, m_t)}{dydm_t^2}.
\]

Since the denominator in formula (3.36) is unknown, one can only approximate the actual acceptance via the value determined from Monte-Carlo simulations:

\[
A^{MC}(y, m_t) = \frac{d^2N_{\text{MC}}(y, m_t)}{dydm_t^2} \Big/ \frac{d^2N_{\text{Hijing}}(y, m_t)}{dydm_t^2}.
\]

For the acceptance calculation we have used 25,000 HIJET central events with added strange particles as the input for the simulation code GEANT. The output of the program GEANT in FASTBUS format was further analyzed by the same tracking program as the real data. In case of \(\Lambda\) the phase space, \((y, m_t)\), was subdivided into inhomogeneous grid, shown on Figure 3.10, and the numbers of lambdas have been

\(^1\)the typical ratio is 2-5%.
counted in each bin for the raw HIJET files and for the analyzed GEANT output. For the $\pi^-$ we have used homogeneous grid with the full rapidity range of $y \in (1.4; 3.6)$, $m_t - m_0 \in (0; 1.0) \text{ GeV}/c^2$, where rapidity range was subdivided into 11 bins, 0.1 unit of rapidity each, and $m_t - m_0$ interval was divided into 60 bins of equal widths.

Since the effective mass selection criterion does not allow one to fully sort out the background fake vertex combinations in the region within the peak of the $M_{\pi^-p}$ distribution, we have performed the background subtraction using the tails of invariant mass distribution. The $\Lambda$ signal was defined requiring the effective mass of $V_0$ to be in the range, $M_{\pi^-p} \in (1.104; 1.128)$. The tails of effective mass distribution, $M_{\pi^-p} \in (1.080; 1.092)$ and $M_{\pi^-p} \in (1.148; 1.160)$, were used for background subtraction. This choice of the signal and the background intervals is based on equal area under the curve of background in the region of the peak and tails. The acceptance is numerically calculated as follows:

$$A_{MC}(i,j) = \frac{N_{MC}^{sig}(i,j) - N_{MC}^{bkg}(i,j)}{N_{HIJET}(i,j)} \left[ 1 \pm \frac{\Delta A_{MC}(i,j)}{A_{MC}(i,j)} \right] , \quad (3.38)$$

where $i$ and $j$ are the numbers of $y$ and $m_t$ bins, respectively. The relative statistical error in acceptance is given by

$$\frac{\Delta A_{MC}(i,j)}{A_{MC}(i,j)} = \sqrt{\frac{N_{MC}^{sig}(i,j) + N_{MC}^{bkg}(i,j)}{(N_{MC}^{sig} - N_{MC}^{bkg})^2} + \frac{1}{N_{HIJET}(i,j)}} . \quad (3.39)$$

Acceptance corrected data, thus equals

$$N_{DATA}^{corr}(i,j) = \frac{N_{DATA}(i,j)}{A_{MC}(i,j)} \left[ 1 \pm \frac{1}{N_{DATA}(i,j)} + \left( \frac{\Delta A_{MC}(i,j)}{A_{MC}(i,j)} \right)^2 \right] . \quad (3.40)$$
The $\Lambda$ production double differential cross spectrum normalized per one central interaction is expressed as follows:

$$
\left( \frac{1}{N_{ev}} \right) \frac{d^2 N}{dy dm_t^2} = \left( \frac{1}{N_{ev}} \right) \frac{N_{\text{corr}}^{\text{DATA}}(i,j)}{y_{\text{bin}} m_{t_{\text{bin}}}(j)},
$$

(3.41)

where $N_{ev} = N_{\text{cen}}$ is the number of central events (see subsection 3.2.3).

The calculation is even more trivial in case of $\pi^-$, where no background subtraction needs to be done.

### 3.2.6 Search for Systematic Errors

A systematic error may be introduced at the stage of acceptance corrections, provided that simulation does not take into account a significant apparatus effect. Generally the acceptance is a function of many variables, thus in the value $A(y, m_t)$ used for the raw data correction, an integration is performed over all the variables that can vary for a given point in the phase space $(y, m_t)$:

$$
A(y, m_t) = \int \ldots \int A(y, m_t, x_1, \ldots x_n) dx_1 \ldots dx_n.
$$

(3.42)

If the apparatus effects are imitated properly in the simulation program, the dependence on the variables $(x_1, \ldots x_n)$ is identical for $A(y, m_t)$ and $A_{MC}(y, m_t)$, otherwise, a bias is introduced via integration over $(x_1, \ldots x_n)$. One may scan the ratio $(dN_{\text{DATA}}/dN_{\text{MC}})$ as a function of $(y, m_t, x_1, \ldots x_n)$ in order to discover discrepancies. The above ratio is expected to be independent of any variable $x_i$ with the fixed values of $(y, m_t)$. To be more specific about variables $(x_1, \ldots x_n)$, let us note that the
Figure 3.10 The grid in the phase space used for acceptance calculations.
efficiency of the reconstruction for the Λ vertex is defined by the trajectories in the TPC’s, along which the decay products travel after the Λ decay. Mathematically these trajectories are defined by the point of the Λ decay, and the orientation of decay products in space.

Miscalibration of the distortions makes it impossible to reproduce the detector efficiencies and momentum resolution when the tracks intersect the affected TPC region. One may expect discrepancies of data and Monte–Carlo Λ counts, when the proton momentum points in the direction corresponding to the positive ion cloud.

**Dynamic Distortions and the $M_{\pi-p}$ Distribution:** As it was mentioned in subsection 3.1.2, our calibration procedures do not allow one to do the distortion corrections on event by event basis, instead the averaged distortion map is used. In reality the dynamic part of the field distortions caused by copiously produced ions in the beam region changes from event to event and depends on the position of the triggered event in the spill and on the beam intensity for a given spill. The dynamic part of the distortions results in coherent shift of hits in the neighborhood of the ion sheet in the outward direction. This leads to the mismeasurement of the tracks crossing the beam region and to the reduced efficiency of the hit and track reconstruction there.

We observe the deterioration of the momentum resolution and a loss of efficiency in case of Λ’s with the protons crossing the region of the maximum dynamic distortions.
On Figure 3.11 we show the two-dimensional histogram illustrating the distribution of the $\Lambda$ candidates over effective mass, $M_{\pi^- p}$, and the azimuthal angle $\Phi_p$ of the proton's momentum for the real data. On Figure 3.12 we show the analogous distribution resulting from the Monte-Carlo simulation. One can see that there is significant difference of the Monte-Carlo and data distributions in the angular region, $-0.9 < \Phi_p < 2.6$, namely for the data the signal in this region is dispersed over the wide range of effective masses, indicating the considerable mismeasurement of the positive tracks with such $\Phi_p$. As a result the $\Lambda$ signal in this region is practically lost.

However we do not observe the analogous effect with the negative tracks, and it can be explained as follows: just from kinematics alone in the decay $\Lambda \rightarrow \pi^- p$ the proton takes most of the momentum of the $\Lambda$ thus producing a stiff track curved in the same direction as the beam track, while the $\pi^-$ is relatively soft and has an opposite curvature. Therefore $\pi^-$ originating from $\Lambda$ decay spends just a little time in the distortion region and crosses it at relatively large angle, which suggests only insignificant influence of the beam caused distortions on the measurements of the soft negative tracks. It will be shown later that the distortions introduce only a minor effect on the measurement of spectrum for $\pi^-$ mesons from the target.

**Combinatorial Background Abundance in the Data:** In order to illustrate the abundance of combinatorial background in the data, on Figures 3.13 and 3.14 we have plotted the invariant mass, $M_{\pi^- p}$, distributions for the data and Monte-
Carlo respectively. The data sample contains 14,114 central events and the Monte–
Carlo sample consists of 25,000 central events. On both Figures we have shown
three histograms corresponding to different cuts. It was explained in subsection 3.2.2
that combinatorial background of the "fish tail" type can be encountered only for
azimuthal angle of the decay plane $\Phi_{\pi p}$ around $180^0$, we therefore discard the region
of $|\Phi_{\pi p}| > 2.7$ to leave only the region of a good signal to background ratio (dashed
histogram on both figures above). Before this cut is made one can see that amount
of background in the data greatly exceeds that of the Monte–Carlo.

The explanation for the excessive combinatorial background in the data is the
following: the protons from the target that get into the distortion region are recon-
structed with significantly mismeasured parameters, so that they can not be projected
to the target within 7 mm cut (see subsection 3.1.3) and thus may fall into $V_0$ daugh-
ters category – be assigned to the vertex bank; abundance of $\pi^-$’s production favors
to creation of the fake vertex combinations.

To assure that the data and the Monte-Carlo have similar momentum resolution
and efficiencies, we therefore have to ban proton going through the distortion region,
which is done as explained in the previous paragraph, ($\Phi_p$ cut).

After applying the combinatorial background cut and the distortion region angular
cut, we come to fairly good agreement of the backgrounds in data and Monte–Carlo,
which means that detector is simulated properly everywhere but in the distortion
region \(^5\). The resulting effective mass distributions for data and Monte–Carlo normalized per one central event are shown on Figure 3.15.

**Calibration of Track Efficiencies Using ARC model output:** The angular cuts mentioned above, however do not guarantee that all proton tracks in the Monte–Carlo will be reconstructed with the same efficiencies as in the data. In other words, some of the tracks left over after the angular cuts above have been applied, may still get into the distortion region. Since there is a correlation between momentum of the \(\Lambda\) and the region of the TPC where its daughter tracks go, efficiency miscalibration can lead to the change of the measured \(\Lambda\) spectrum shape, which is crucial for determining the flow parameters (see section 4). We therefore need to perform additional calibration of the track reconstruction efficiency in the region of the overlap above.

Fortunately the \(p\) and \(\pi^+\) invariant differential spectra have been measured in a wide range of rapidities and transverse momenta by experiment E866 [91] and are found to be in a fairly good agreement with the model predictions by "A RelativisticCascade Model" (ARC) [99]. We will use the ARC model output for the positive particle spectra in order to determine the positive track reconstruction efficiencies as a function of \(x, y\) coordinates of the track intercept with the front plane of the TPC:

\[
E_{\text{tracks}}(i, j) = \frac{N_{\text{tracks}}^{\text{DATA}}(i, j)/N_{\text{events}}^{\text{DATA}}}{N_{\text{tracks}}^{\text{MC}}(i, j)/N_{\text{events}}^{\text{MC}}}, \tag{3.43}
\]

\(^5\)The insignificant discrepancy of backgrounds can be explained by the relatively small spatial overlap of the TPC regions defined by the left over angles, \(\Phi_p\) and \(\Phi_{\pi p}\), and miscalibrated TPC regions.
Figure 3.11 The proton's direction dependence of $M_{\pi-p}$ distribution for data. Dashed lines indicate the $\Phi_p$ cuts.
Figure 3.12 The proton's direction dependence of $M_{\pi-p}$ distribution for MC.
Figure 3.13 $M_{\pi-p}$ Effective Mass Distribution Integrated over azimuthal orientation angle of the decay plane $\Phi_{\pi p}$, 14,114 data events. The solid histogram corresponds to the integration over full interval ($-\pi < \Phi_{\pi p} < \pi$); the dashed histogram – integration over interval ($-2.7 < \Phi_{\pi p} < 2.7$); the dotted histogram – the same as dashed but the positive particles from $V_0$'s are not allowed to have the azimuthal angle $\Phi_p$ outside the interval ($-0.9 < \Phi_p < 2.6$).
Figure 3.14 $M_{\pi^-p}$ effective mass distribution integrated over azimuthal orientation angle of the decay plane $\Phi_{\pi p}$, 25,000 Monte-Carlo events. The solid histogram corresponds to the integration over full interval $(-\pi < \Phi_{\pi p} < \pi)$; the dashed histogram – integration over interval $(-2.7 < \Phi_{\pi p} < 2.7)$; the dotted histogram – the same as dashed but the positive particles from $V_0$'s are not allowed to have the azimuthal angle $\Phi_p$ outside the interval $(-0.9 < \Phi_p < 2.6)$. 
Figure 3.15 $M_{\pi^+p}$ effective mass distribution, final sample, $(1.4 < y < 3.2)$. Normalization per one central event done in both data and Monte-Carlo. Solid histogram - data; dashed histogram - Monte-Carlo.
where \( N_{\text{DATA}}^{\text{tracks}} \) is number of tracks in the real data, \( N_{\text{MC}}^{\text{tracks}} \) is that of the Monte-Carlo with ARC used as input and the numbers \((i, j)\) indicate position of the 2 cm \( \times \) 2 cm square to be calibrated. We perform the cut on the impact parameter, \( b_{\text{imp}} < 3 \text{ fm} \) on the ARC event sample, to achieve the best description of our centrality selection criterion and to match the measured interaction cross section.

Figure 3.16 shows the results of the efficiency calculation according to (3.43). One can see the significant deterioration of the positive track reconstruction efficiency in the neighborhood of the positive ion sheet. To illustrate a considerable overlap of the region of reduced efficiency of the positive track reconstruction with the area corresponding to the protons from \( \Lambda \) decay, on Figure 3.16 we also show the \( \Lambda \) counts (background subtracted) as a function of \( x, y \) coordinates in the front plane of the TPC. The significant fraction of overall number of \( \Lambda \)'s should be affected via reduced efficiency as can be seen from the Figure above. In the region of \((x > -6 \text{ cm})\) and \((y > -4 \text{ cm})\) the track efficiency drops below 80\%, while the number of \( \Lambda \)'s going through this region and surviving all cuts is about 1160.

The calculation shows that \( \Lambda \)'s affected by the track efficiency miscalibration have large rapidities and/or low transverse momenta. Performing additional calibration of the track reconstruction efficiency using ARC model output we obtain slightly increased yield of \( \Lambda \) hyperons mainly in the region of high rapidities. The systematical errors for the \( \Lambda \) spectra do not exceed the statistical ones.
Systematic Errors for $\pi^-$ Multiplicity Measurements: To determine the $\pi^-$ spectrum we select the region of the positive angles $\Phi^-$, where the ratio of data counts to those of Monte-Carlo becomes flat. These ratios for all rapidity bins are shown on Figure 3.17. This selection criterion leads to the small systematic error which is estimated via tightening the $\Phi^-$ cut to the region $\Phi^- \in (0.5; 2.5) \text{ rad}$. The differences in $\pi^-$ counts in each $(m_t, y)$ bin are regarded as systematic errors. Varying the saggita cut for the tracks left by negative pions we come to a noticeable sensitivity of the acceptance corrected results for the $\pi^-$ invariant spectra and the rapidity density in the region of rapidities, $y < 2.0$. By varying the saggita cut from 2000 to 4000 we find that the systematic error for the $\pi^-$ mesons exceeds the statistical error near midrapidity and may be as large as 10%.
Figure 3.16  Positive track reconstruction efficiency as a function of $x$ and $y$ coordinates of the track intercept with the front plane of the TPC. Beam impact coordinates are $x = 3.1 \, \text{cm}$, $y = 0.0 \, \text{cm}$. 
Figure 3.17  Data to Monte-Carlo ratios of the $\pi^-$ counts as a function of the azimuthal angle $\Phi^-$ for all bins of rapidity used.
Chapter 4

Results and Discussions

4.1 Introduction

In this chapter, we present the experimental results from the analysis procedure described in Chapter 3, for the production of $\Lambda$ hyperons and the negative pions in the central $Au + Au$ collisions, corresponding to 4.4% of the geometric cross section. The centrality is defined by the charged multiplicity cut as described in the previous Chapter.

The double differential cross sections are presented as a function of transverse kinetic energy, $m_t - m_0$, motivated by the fact that the thermal, invariant Boltzmann distribution is given by equation

$$
\mathcal{E} f_{\text{bollz}}(\vec{p}) \equiv \mathcal{E} \frac{d^3N_{\text{bollz}}}{d\vec{p}} = m_t \cosh(y - \eta) e^{-m_t \cosh(y-\eta)/T},
$$

where $\mathcal{E} = m_t \cosh(y - \eta)$ is the particle energy measured in the reference frame of the emitting source moving in the longitudinal direction with the velocity $v = \tanh(\eta)$.

To emphasize the difference between purely thermal distribution (4.1) and observed particle spectra we present double differential cross sections divided by the transverse mass, $m_t$. In order to interpret the observed particle spectra we introduce
a Flow Model parametrization allowing one to obtain a quantitative picture of the three dimensional expansion of hot nuclear matter. The rapidity distributions are also presented in this chapter with extrapolations to the region of high \( m_t \) based on the Flow Model parametrization used to evaluate the double differential particle spectra. Experimental results for the double differential spectra of \( \Lambda \) and \( \pi^- \), the rapidity distributions and total yields of these particles are compared to the predictions of microscopic cascade models, ARC and RQMD.

### 4.2 \( \Lambda \) Production in Central Au+Au Collisions

On Figure 4.2 we present acceptance corrected double differential spectra for the inclusive \( \Lambda \) production in central \( Au + Au \) interactions with the integrated cross section of \( \sigma_{\text{cen}} = 270 \text{ mb} \). In the interval of our acceptance the double differential spectra for \( \Lambda \) hyperon production change from the exponential behavior in the interval \( y \in (2.0; 3.2) \) to the convex spectral shapes for \( y \in (1.4; 2.0) \). From Figure 4.2 one can see that the spectrum becomes less steep as \( y \) approaches midrapidity \( y_0 = 1.6 \). As a first step to quantitatively characterize the spectrum we perform local (separately in each rapidity bin) fits.

#### 4.2.1 Rapidity Dependence of Inverse Slopes for \( \Lambda \) Production

Since experimental spectra for \( \Lambda \) production in the central \( Au + Au \) interactions have exponential shapes in the rapidity region \( y \in (2.0; 2.9) \) as functions of the
transverse mass, it is legitimate to extract such values as inverse slopes $B(y)$ from parametrizations of the Λ spectra via exponentials:

$$\frac{1}{N_{ev} \ dy dm_t^2} \frac{d^2N}{dy dm_t^2} = e^{[A_{exp}(y) - B_{exp}(y) m_t]}, \quad (4.2)$$

and via Boltzmann–like shapes:

$$\frac{1}{N_{ev} \ dy dm_t^2} \frac{d^2N}{dy dm_t^2} = m_t e^{[A_{boll}(y) - B_{boll}(y) m_t]}, \quad (4.3)$$

for further comparison with analogous measurements of exponential inverse slopes by E810 experiment in case of Si + Si collisions in case of (4.2) and to compare with the spectrum of the thermal source in case of (4.3).

On Figure 4.1 we present the results of the local exponential and Boltzmann–like fits of the invariant Λ spectrum $(1/N_{ev})d^2N/dy dm_t^2$. On Figure 4.1 a) the exponential inverse slopes for various rapidities $y$ measured by E891 and by E810 experiments are shown, where one can observe a general tendency of the inverse slopes to increase for heavier system created in Au + Au collision. Figure 4.1 b) presents the Boltzmann–like inverse slopes, $B(y)$, from E891, while the dashed line is a function corresponding to the rapidity dependence of inverse slopes of a thermal source at $y = y_0$ with the temperature $T = 0.230 \ GeV$, as one can check with equation (4.1). On Figures 4.1 c) and 4.1 d) we present parameters, $A(y)$, resulting from the fits according to the formulas (4.2) and (4.3) respectively. On Figure 4.1 d) the dashed line is a function $9.5 + \ln [\cosh (y - y_0)]$, corresponding to a thermal source at $y = y_0$ of an arbitrary temperature. One can see a slight increase of $A_{boll}(y)$ as compared to $A_{exp}(y)$ and
Figure 4.1 Results of local exponential and Boltzmann-like fits of the invariant Λ spectrum. Figure a): exponential inverse slopes measured by E891 and by E810. Figure b): Boltzmann-like inverse slopes from E891; the dashed line corresponds to a thermal source at $y = y_0$ with the temperature $T = 0.230$ GeV. Figures c) and d): parameters $A$ in the fits according to the formulas (4.2) and (4.3) respectively. On Figure d) the dashed line corresponds to a thermal source at $y = y_0$ of an arbitrary temperature. Errors are statistical only.
a reduction of $B_{\text{boltz}}(y)$ compared to $B_{\text{exp}}(y)$, which is a result of $m_t$ factor in (4.3). The quality of the fits of both kinds is however approximately the same.

### 4.2.2 Global Fireball Parametrizations of $\Lambda$ Spectrum

To make a comparison of the $\Lambda$ spectra in $Au+Au$ collisions with the corresponding spectra in $Si+Si$ collisions, we perform the fit via the shapes used previously: The experimental data for $\Lambda$ production cross section in $Si+Si$ collisions measured by E810 collaboration were found to be described by the following dependence [88]:

$$\left(\frac{1}{N_{ev}}\right) \frac{d^2 N}{dy dm_t^2} \propto e^{-(3.7842+1.035 \cosh(y-y_0))m_t}, \quad (4.4)$$

where the transverse-mass dependence is exponential for any rapidity.

**Exponential Fireball Parametrization:**

Figure 4.2 shows that the invariant differential distributions for the $\Lambda$ production in $Au+Au$ interactions for rapidities $y > 2$ exhibit exponential dependence from the transverse mass, $m_t$. We therefore parametrize these data via the following function:

$$\left(\frac{1}{N_{ev}}\right) \frac{d^2 N}{dy dm_t^2} = e^{[A-(B+C \cosh(y-y_0)+D(y-y_0)^2/m_t)m_t]. \quad (4.5)$$

Comparing with the equation (4.2) and referring to Figure 4.1 d) one can see that the reason for introducing the additional parameter $D$ to the parametrization (4.4) is a growth of parameter $A_{\text{exp}}$ as a function of rapidity $y - y_0$ in the center of mass system.
Boltzmann–like Fireball Parametrization:

Due to the fact that the interval of the transverse masses within our experimental acceptance is small compared to the $$\Lambda$$ mass, $$m_0 = 1.1156 \text{ GeV}/c^2$$, introducing additional factor, $$m_t$$, does not reduce the quality of the fit in terms of $$\chi^2$$ as can be seen in Table 4.1.

\[
\left( \frac{1}{N_{ev}} \right) \frac{d^2N}{dy dm_t^2} = m_t \cosh (y - y_0) e^{[A - (B + C \cosh(y - y_0) + D(y - y_0)^2/m_t - m_t)]}. \tag{4.6}
\]

Again, parameter $$D$$ was needed in order to obtain a better fit to the data.

The possibility to describe the invariant differential $$\Lambda$$ spectrum by parametrizations above shows that the spectra $$\Lambda$$ hyperons produced in central $$Si + Si$$ and $$Au + Au$$ collisions at AGS differ from those expected of a single thermal source. The difference may be attributed for example to the collective expansion of the nuclear matter involving $$\Lambda$$ hyperons. The disadvantage of these parametrizations is that the relation of the parameters used with the parameters of the expansion is unknown. Moreover, none of parametrizations above is suitable for the description of the $$\Lambda$$ particle spectrum in the region near midrapidity, $$y < 2.0$$, where deviations from a single exponential scaling is observed.

4.2.3 Collective Effects and Flow Model Parametrization

An interesting phenomenon observed in relativistic $$Au + Au$$ collision by E891 is low $$m_t$$ suppression in the double differential spectra of $$\Lambda$$ hyperons around midrapidity,
Figure 4.2 Double differential spectra for the inclusive $\Lambda$ production in central $Au + Au$ interactions from experiment E891. The spectra correspond to the integrated cross section of $\sigma_{\text{cen}} = 270 \text{ mb}$. Dotted lines – RQMD model predictions. Dashed lines – the ARC prediction. Solid lines – results of the Flow Model parametrization as described in the text. Errors are statistical only.
Table 4.1  Results of global fits of $\Lambda$ spectrum.

<table>
<thead>
<tr>
<th></th>
<th>A±ΔA</th>
<th>B±ΔB</th>
<th>C±ΔC</th>
<th>D±ΔD</th>
<th>χ²/NDF</th>
</tr>
</thead>
<tbody>
<tr>
<td>Eq. (4.6)</td>
<td>5.81±0.40</td>
<td>-3.19±0.82</td>
<td>6.55±0.61</td>
<td>-3.97±0.47</td>
<td>24.74/33</td>
</tr>
<tr>
<td>Eq. (4.5)</td>
<td>5.11±0.40</td>
<td>-4.07±0.82</td>
<td>6.67±0.60</td>
<td>-4.45±0.47</td>
<td>25.38/33</td>
</tr>
</tbody>
</table>

which is a clear deviation from the superposition of $N + N$ collisions. The spectra shown on Figure 4.2, indicate non-exponential behavior, namely convex shape for $y < 2.0$. No such behavior was observed for lighter systems by E810 [88].

The comparison of the invariant $\Lambda$ spectrum measured by E891 collaboration with a thermal source Boltzmann spectrum reveals significant differences, even in the region of rapidities $y > 2$, where the transverse-mass distributions are exponential. While the rapidity dependence of inverse slopes on Figure 4.1 b) can lead one to the conclusion that the $\Lambda$ spectrum corresponds to that of the thermal source with the temperature of about 230 $MeV$, the dependence of normalization factors $A_{\text{bolts}}(y)$ shown on figure 4.1 d) disagrees with those describing a single thermal source.

Let us review the existing parametrizations for the invariant differential spectra of particles based on different scenarios of expansion and construct our own. Each parametrization is of course schematic and based on the simplified idea of representing a particle spectrum as superposition of the classical moving thermal sources distributed over a hypothetical flow surface. This approximation has often been called a blast wave approximation (see for example Reference [82]).
Defining the momentum distribution function of particles as \( f(\vec{p}) \), we note that it changes under the Lorentz transformations, \((E, \vec{p}) \rightarrow (E', \vec{p}')\), as follows:

\[
\mathcal{E}' f'(\vec{p}') = \mathcal{E} f(\vec{p}),
\]  

while the product \( \mathcal{E} f(\vec{p}) \) is invariant under the Lorentz transformations. This relation will be used hereafter to obtain the momentum distribution functions from moving sources.

**Boltzmann Source:**

The invariant differential spectrum from the thermal source at rapidity \( \eta \) is described by expression (4.1). In the center of mass frame of the source, the integration of the expression (4.1) over transverse momentum gives the rapidity distribution:

\[
\frac{dN_{\text{boltz}}}{dy} = \pi \int_{m_0}^{\infty} \mathcal{E} f(\vec{p}) dm_t^2 = 2\pi T m_0^2 (1 + 2\chi + 2\chi^2) e^{-1/\chi},
\]  

\[
\chi = \frac{T}{m_0 \cosh(y)},
\]

while the total number of particles emitted by the source (4.1) is

\[
N_{\text{boltz}}(m_0, T) = 2\pi T m_0^2 \int_{-\infty}^{\infty} (1 + 2\chi + 2\chi^2) e^{-1/\chi} dy.
\]

**Spherically Symmetric Expansion:**

For the explanation of the different inverse slopes in the momentum distributions for protons and pions at Bevalac in the central \( N_e + NaF \) collisions at the lab energy of 800 MeV per nucleon Siemens and Rasmussen [82] suggested that after thermal
equilibration the nuclear matter experiences the spherically symmetric expansion (see Figure 4.3). Due to the fact that protons are heavier than pions the momentum generated by a collective flow velocity and thus an increase of the inverse slope is more significant for protons than it is for pions.

Boosting the invariant Boltzmann distribution (4.1) in the direction of the unit vector, \( \vec{n} \), by velocity, \( v \), according to formula (4.7) we get *:

\[
\mathcal{E} f(\vec{p}, \vec{n}) = \gamma (\mathcal{E} - \vec{n} \vec{p} v) e^{-\gamma (\mathcal{E} - \vec{n} \vec{p} v) / T},
\]

(4.11)

\[
\mathcal{E} f_{sph}(\vec{p}) = \int d\vec{n} \mathcal{E} f(\vec{p}, \vec{n}).
\]

(4.12)

The integral above can be easily taken analytically and leads to the following expression for the invariant spectrum from the spherically expanding source:

\[
\mathcal{E} f_{sph}(p) = 4\pi T e^{-\gamma \mathcal{E} / T} \left\{ \left( 1 + \frac{\gamma \mathcal{E}}{T} \right) \frac{\sinh \alpha}{\alpha} - \cosh \alpha \right\},
\]

(4.13)

\[\alpha = \gamma v p / T, \quad \gamma = 1 / \sqrt{1 - v^2},\]

(4.14)

where the energy, \( \mathcal{E} \), and momentum, \( p \), are measured in the center of mass frame. For evaluation of the total number of particles emitted by this source, let us note that the Lorentz transformation does not change the integrated yield of particles, thus,

\[
N_{sph}(m_0, T) = \int d\vec{n} N_{bottz}(m_0, T) = 4\pi N_{bottz}(m_0, T).
\]

(4.15)

At high values of momentum one obtains the asymptotic values of inverse slopes, or apparent temperatures, to be \((-d\ln \sigma / d\mathcal{E})^{-1} = T \gamma^{-1} (1 - v \mathcal{E} / p)^{-1}\). Thus pions with \( p \approx \mathcal{E} \) appear to be cooler than protons (\( p < \mathcal{E} \)).

*Hereafter we express expansion velocities in the units of the speed of light.
Figure 4.3 Spectra of protons and pions normal to the beam in the c.m.s for central Ne + NaF interactions at 800 A MeV in the lab, explained by the blast wave by P. Siemens and J. Rasmussen [82]. The dashed line is a Boltzmann distribution extrapolated from the high-energy proton cross section. The solid curve is a blast wave parametrization by equation (4.13).
Application of the spherical expansion picture to the experimental data on the particle spectra measured by E891 gives unsatisfactory results, since the expansion at AGS energy is anisotropic.

**Transverse Radial Expansion:**

It is of practical interest to obtain an expression for a spectrum resulting from superposition of sources moving in the transverse directions distributed uniformly over the azimuthal angle, \( \varphi \). One of such infinitesimal sources will produce a spectrum

\[
\mathcal{E}f(p, \varphi) = \gamma (E - p_t v \cos \varphi) e^{-\gamma (E - p_t v \cos \varphi)/T},
\]

while altogether the sources define the invariant distribution given by the integral,

\[
\mathcal{E}f_i(p) = \int_0^{2\pi} d\varphi \mathcal{E}f(p, \varphi),
\]

\[
\mathcal{E}f_i(p) = 2\pi \gamma e^{-\gamma E/T} \{ \mathcal{E}I_0(\alpha) - p_t v I_1(\alpha) \},
\]

\[
\alpha = \gamma p_t v / T, \quad \gamma = 1/\sqrt{1 - v^2}.
\]

where \( I_0(x) = (2\pi)^{-1} \int_0^{2\pi} e^{x \cos \varphi} d\varphi \) and \( I_1(x) = (2\pi)^{-1} \int_0^{2\pi} \cos \varphi e^{x \cos \varphi} d\varphi \) are modified Bessel functions of zeroth and first order respectively. The integrated yield of the source, \( f_i(p) \), is given by the equation,

\[
N_i(m_0, T) = 2\pi N_{\text{boltz}}(m_0, T).
\]

The curvature of the logarithmic plot of \( \mathcal{E}f_i(p)/m_t \) as well as its sign is uniquely determined via the temperature, \( T \), expansion velocity, \( v \), and the particle mass \( m_0 \).
For example, for $T = 90\, MeV$ and $v = 0.6$, the $\Lambda$ hyperons and protons must exhibit the negative curvature, i.e. convex shape of the spectrum, while for $\pi$ mesons the invariant transverse mass distribution divided by $m_t$ must be concave on a logarithmic scale.

**Boost Invariant Longitudinal Expansion:**

According to the Bjorken scenario of one dimensional longitudinal expansion [83], the spectra of particles produced in ultra-relativistic heavy-ion collisions are boost invariant. Such a spectrum can be represented as superposition of the thermal sources distributed uniformly over a limited interval of the longitudinal boost angle, $\eta$ [84] \(^1\):

$$E_{f_i}(\vec{p}) = \int_{-\eta_{max}}^{\eta_{max}} d\eta m_t \cosh(y - \eta) e^{-m_t \cosh(y - \eta)/T}, \quad (4.21)$$

where rapidities $y$ and $\eta$ are measured in the center of mass frame. The limits in the boost invariance interval, $[-\eta_{max}, \eta_{max}]$, are confined within the projectile and target rapidities. The integrated yield of the source (4.21) is evaluated as follows:

$$N_i(m_0, T) = 2\eta_{max} N_{bolts}(m_0, T). \quad (4.22)$$

Assuming that the longitudinal and the transverse expansion can be decoupled, expression (4.21) integrated over the transverse mass squared was used in [84] to evaluate the rapidity distributions of different hadrons in $Si + Al$ interactions at AGS energy (see Figure 4.4). Rather good description of the rapidity distributions was

\(^1\)The longitudinal boost angle, $\eta = 0.5 \ln[(1 + v_z)/(1 - v_z)]$, is essentially the rapidity of the moving source.
Figure 4.4  Rapidity distributions for central 14.6 \( A \, GeV/c \) Si + Al collisions [84] in comparison to isotropic thermal distribution at \( T = 0.12 \, GeV \) — solid lines and distributions for source at the same temperature expanding with \(< v_t > = 0.52\) — dashed lines.
achieved with the temperature, $T = 120 \ MeV$, and the mean longitudinal expansion velocity, $< v_l > = 0.52$. However the authors of [84] did not explore the whole parameter space of $T$ and $< v_l >$. In addition, no attempt to parametrize the double differential distributions in a wide range of rapidities was made in this article. For the description of invariant differential multiplicity at midrapidity the authors applied expression (A8) from [85], which is intended for the description of the transverse mass distribution integrated over rapidity, not the invariant differential distribution!

Unlike the CERN-SPS energies, at AGS energy domain there seems to be no particular reason for the expansion to be boost invariant. In addition, nonzero curvatures of the particle spectra, $E f(\vec{p})/m_i$ on a logarithmic plot, indicate the presence of a transverse radial expansion. The longitudinal and the transverse expansion may be decoupled accurately only in the limit of nonrelativistic transverse velocities. Therefore we believe that the correct procedure of determining the flow velocity profile is a parametrization of the double differential spectra of particles. In this procedure the rapidity distributions and integrated transverse mass distributions will be described automatically.

**Three Dimensional Expansion – the Flow Model Parametrization:**

We represent the inclusive spectrum for any particle as a continuous superposition of the flowing Boltzmann sources with a temperature, $T$. The flow velocity profile is chosen to be an azimuthally symmetric surface in the space of the longitudinal boost
angle and the transverse two dimensional velocity. The flow surface in our approach belongs to the interval, \([-\eta_{\text{max}}; \eta_{\text{max}}]\), where both limits are measured in \(Au + Au\) center of mass system.

\[
\mathcal{E}f_{3D}(\vec{p}) = \int_{-\eta_{\text{max}}}^{\eta_{\text{max}}} d\eta \mathcal{E}f(y, \eta, m_t),
\]

(4.23)

\[
\mathcal{E}f(y, \eta, m_t) = \frac{1}{\pi} A(\eta) e^{\frac{m_0^2}{T^2} - a_1 [a_1 I_0(a_2) - a_2 I_1(a_2)]},
\]

(4.24)

\[
a_1 = \frac{\gamma(\eta) m_t \cosh (y - \eta)}{T}, \quad a_2 = \frac{\gamma(\eta) p_t v_t(\eta)}{T}.
\]

(4.25)

Choosing the flow velocity profile and the distribution of the power of the source along the flow surface we employ the following simple functions:

\[
v_t(\eta) = v_0 (1 - \eta^2/\eta_{\text{max}}^2)^{\alpha_1}, \quad A(\eta) = A_0 (1 - \eta^2/\eta_{\text{max}}^2)^{A_1},
\]

(4.26)

\[
\frac{\gamma_t(\eta)}{\sqrt{1 - v_t^2(\eta)}},
\]

(4.27)

where \(A_1, v_1 \geq 0\). Rapidities, \(y\) and \(\eta\), are measured in the center of mass frame and the transverse velocity, \(v_t\), is measured in a system co-moving with \(\eta\). The choice of these functions is somewhat arbitrary, but the functions chosen allow a reasonable flexibility, are even functions of the source rapidity as required by the forward–backward symmetry of the collision and approach zero as rapidity of the source, \(\eta\), approaches limiting values, \(\pm \eta_{\text{max}}\). The functions above reflect the fact of increase of the transverse expansion velocity towards midrapidity, which is seen from
the increase of the curvature of the double differential cross section on a logarithmic
plot (see Figure 4.2). Mean rapidity and the transverse velocity are given by the
following expressions:

\[ < \eta > = \frac{\int_0^1 x(1-x^2)^{A_1} dx}{\int_0^1 (1-x^2)^{A_1} dx} \eta_{max} = \frac{\Gamma(A_1+3/2)}{\sqrt{\pi} \Gamma(A_1+2)} \eta_{max}, \]

\[ < v_t > = \frac{\int_0^1 (1-x^2)^{A_1+n_1} dx}{\int_0^1 (1-x^2)^{A_1} dx} v_0 = \frac{\Gamma(A_1+n_1+1)}{\Gamma(A_1+n_1+3/2)} \frac{\Gamma(A_1+3/2)}{\Gamma(A_1+1)} v_0, \] (4.28)

where averaging is performed over the forward hemisphere in the center of mass
system. The rapidity distribution and the total number of particles emitted by this
source are

\[ \frac{dN_{3D}(y)}{dy} \approx \frac{e^{m_0^2}}{\pi T} \int_{\eta_{max}}^{-\eta_{max}} d\eta A(\eta) \frac{dN_{bolts}(y-\eta)}{dy}, \] (4.29)

\[ N_{3D}(m_0, T) = \frac{A_0 \eta_{max} e^{m_0^2}}{\sqrt{\pi} T} \frac{\Gamma(A_1+1)}{\Gamma(A_1+3/2)} N_{bolts}(m_0, T), \] (4.30)

respectively, where \( \Gamma(x) = \int_0^\infty t^{x-1} e^{-t} dt, (x > 1) \) is well known gamma function.

The expression (4.23) represents the convolution of the Boltzmann distribution
with the flow excitation function,

\[ w(v_t, \eta) = A(\eta) \delta(v_t - v_t(\eta)), \] (4.31)

which was constructed by us for a schematic description. The delta-functional form of
the flow excitation function is motivated by the fact that the product of the actual flow
excitation function by the differential element \( 2\pi v_t dv_t \) will lead to the selection of the
non-zero transverse velocities defined by the maximum of the product \( 2\pi v_t w(v_t, \eta) \),
which arises due to the competition of its factors.
The physical meaning of the function $A(\eta)$ is the distribution of the particle emission power over the boost angle, $\eta$, along the flow surface. In case of boost invariant scenario, a constant power distribution must take place, which is achieved with the parameter $A_1$ equal to zero in formula (4.23). The parameter $v_1$ in our parametrization defines the flow velocity profile or the shape of the flow surface. In case of the boost invariant scenario the flow surface is a cylinder with the axis along $\eta$, which is achieved with the parameter $v_1$ equal to zero in (4.23). Finally, for low energies, when the expansion is spherically symmetric [82], our formula (4.23) is applicable with the parameters, $v_1 = 1/2$ and $A_1 = 1$.

We have applied the formulas listed in this paragraph to obtain the expansion parameters as well as the mean temperature of $\Lambda$ hyperons after the freeze-out. The obtained parameters $A_0$, $A_1$, $v_0$, $v_1$, $T$, $\eta_{mas}$ are listed in Table 4.2 together with the mean source rapidity, $< \eta >$, the mean transverse velocity, $< v_t >$, and the integrated yield, $N_{3D}$ derived from above six parameters.

4.2.4 Rapidity Distribution of $\Lambda$ Hyperons

Since experimental acceptance of our apparatus for $\Lambda$ hyperons is limited to the interval, $m_t - m_0 < 0.7 \text{ GeV}/c^2$, we extrapolate the measured double differential cross sections to the region of higher values of the transverse mass in order to present the results for rapidity density $dN/dy$. The results of such calculations are listed in
Figure 4.5 E891 rapidity distribution of $\Lambda$ production. Dotted curve is RQMD model prediction. Dashed line – ARC model prediction. Solid curve is the result of integration of the Flow Model parametrization (4.23) over the transverse mass squared. Errors are statistical only.
Table 4.2 Results of Flow Model parametrization of the $\Lambda$ and $\pi^-$ spectra.

<table>
<thead>
<tr>
<th>main parameters</th>
<th>$\Lambda$ E891</th>
<th>$\Lambda$ RQMD</th>
<th>$\pi^-$ E891</th>
<th>$\pi^-$ RQMD</th>
</tr>
</thead>
<tbody>
<tr>
<td>$A_0 ,(GeV/c^2)^{-2}$</td>
<td>4.62 ± 0.67</td>
<td>4.30 ± 0.55</td>
<td>405.7 ± 6.7</td>
<td>345.7 ± 17</td>
</tr>
<tr>
<td>$v_0$</td>
<td>0.50 ± 0.02</td>
<td>0.34 ± 0.02</td>
<td>0.68 ± 0.01</td>
<td>0.67 ± 0.007</td>
</tr>
<tr>
<td>$A_1$</td>
<td>0.62 ± 0.38</td>
<td>0.98 ± 0.32</td>
<td>0.00</td>
<td>0.20 ± 0.09</td>
</tr>
<tr>
<td>$v_1$</td>
<td>0.67 ± 0.30</td>
<td>0.98 ± 0.36</td>
<td>0.33 ± 0.02</td>
<td>0.64 ± 0.04</td>
</tr>
<tr>
<td>$T ,(MeV/c^2)$</td>
<td>90 ± 10</td>
<td>136 ± 10</td>
<td>78 ± 11</td>
<td>98 ± 1</td>
</tr>
<tr>
<td>$\eta_{\text{max}}$</td>
<td>1.41 ± 0.13</td>
<td>1.1 ± 0.09</td>
<td>1.19 ± 0.08</td>
<td>1.36 ± 0.04</td>
</tr>
<tr>
<td>$\chi^2/NDF$</td>
<td>39/51</td>
<td>58/58</td>
<td>1495/481</td>
<td>1007/605</td>
</tr>
</tbody>
</table>

| derived parameters | |
|--------------------|----------------|----------------|--------------|--------------|
| $< v_t >$ | 0.41 ± 0.02 | 0.27 ± 0.02 | 0.57 ± 0.01 | 0.52 ± 0.01 |
| $< \eta >$ | 0.58 ± 0.05 | 0.43 ± 0.03 | 0.59 ± 0.01 | 0.63 ± 0.02 |
| $N_{3D}$ | 20 ± 3 | 18 ± 2 | 160 ± 3 | 183 ± 9 |

Table 4.3. The extrapolations were done according to the flow model parametrization (4.23-4.27). These results are also shown on Figure 4.5 together with RQMD model predictions. The larger width of experimentally measured rapidity distribution of $\Lambda$ hyperons indicates that RQMD underpredicts the longitudinal flow. The flow in RQMD 1.08 is generated by pressure due to the resonance gas only, additional pressure caused by the repulsive mean field potential, is present in reality [100].

4.3 $\pi^-$ Meson Production in Central Au+Au Collisions

The pions are produced copiously in the heavy-ion collisions at AGS and are usually considered as a good probe of the dynamics of heavy-ion collisions. Because of the large cross section of interaction pions are expected to thermalize relatively soon
Table 4.3  Rapidity distribution of $\Lambda$ hyperons.

<table>
<thead>
<tr>
<th>Rapidity</th>
<th>$dN_\Lambda/dy$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$1.7 &lt; y &lt; 2.0$</td>
<td>$9.33 \pm 0.59$</td>
</tr>
<tr>
<td>$2.0 &lt; y &lt; 2.3$</td>
<td>$7.73 \pm 0.41$</td>
</tr>
<tr>
<td>$2.3 &lt; y &lt; 2.6$</td>
<td>$7.22 \pm 0.51$</td>
</tr>
<tr>
<td>$2.6 &lt; y &lt; 2.9$</td>
<td>$5.34 \pm 0.91$</td>
</tr>
</tbody>
</table>

after their production. If the temperature of the system is high enough, the small pion mass makes an influence of the transverse expansion to be minimal. The energy of the transverse expansion however comes from the thermal kinetic energy of the constituents of the collision system, therefore expansion must lead to the reduction of the temperature of the hadrons escaping the system. The influence of the transverse expansion on the shape of the pion transverse mass distribution becomes significant when the temperatures are low enough. So, for example, we have mentioned earlier that pion transverse mass distributions acquire a concave shape for the temperature of about $T \sim 90 \text{ MeV}$ and the transverse expansion velocity $v_t \sim 0.6$. It is likely however that significant fraction of pions entering detectors originates from the decay of hadronic resonances, which survived until the freeze-out. Such pions may have not had a chance to thermalize with the rest of the system.
4.3.1 Double Differential Spectra of $\pi^-$ Production

Let us consider the spectra of negative pions measured in our apparatus. The transverse mass spectra of $\pi^-$ mesons in the central $Au + Au$ interactions, shown on Figure 4.8, have a concave shape, which can be attributed, for example, to the transverse radial expansion of the pion gas, increased role of the secondary interactions and to the attractive mean field potential of the pions in the dense nuclear medium. The departure from the exponential behavior has been established at the Bevalac [101] and in SPS experiments at CERN [103]. The influence of resonance production on transverse momentum spectra in heavy ion collisions has been noted by Sollfrank, Koch and Heinz [104]. The deviation of the pion spectra from the single exponential scaling was observed also by E810 experiment in case of lighter systems, which is reported in [102]. It was suggested there that the deviation from exponential scaling in the $\pi^+$ and $\pi^-$ spectrum can be explained by the excessive production of nucleon resonances in the AGS energy domain.

The effect of the nucleon resonance production on the transverse mass distributions of the pions in relativistic heavy-ion collisions at AGS was estimated by G. Brown et al. in [105]. It was pointed out, that because of effects of flow it is impossible to unambiguously determine the freeze-out temperatures reached in relativistic heavy-ion collisions, however the fraction of the $\Delta(1230)$ resonances among the nucleons in the excited nuclear matter depends through the Boltzmann factors from the temperature. The authors claimed that the best description of $\pi^-$ spectra in
the $Si + Au$ interactions, measured by E810 collaboration, can be achieved with the freeze-out temperature, $T \geq 150 \text{ MeV}$.

At the temperatures about 150 MeV the nucleons occupy the ground states with the probability of roughly 40%, while the rest 60% of the nucleons are in excited, $\Delta$ and $N^*$ resonance states [105] (see also Appendix E). Being the lightest, $\Delta(1230)$ are most common among the nucleon resonances and constitute 34% of the total number of nucleons at this temperature. The largest contribution to the negative pion spectrum from two particle decays comes from these lightest resonances:

$$\Delta^- \rightarrow \pi^- + n, \quad BR = 1 \cdot 1.00,$$
$$\Delta^0 \rightarrow \pi^- + p, \quad BR = 1/3 \cdot 1.00.$$  \hspace{1cm} (4.32)

A comparable contribution is made by the higher lying $\Delta$ and $N^*$ resonances which have significant probability of decay to the three particle final states,

$$\Delta, N^* \rightarrow N\pi\pi,$$  \hspace{1cm} (4.33)

producing the effect of even cooler pions.

Let us make a connection of these statements to the expansion picture used in this thesis. Secondary interactions of the type,

$$N\pi \rightarrow \Delta\pi, N^*\pi,$$  \hspace{1cm} (4.34)

lead to the cooling of the pions due to the negative $Q$-value of the reaction, while the pion absorption and subsequent disintegration of the produced the lightest nucleon
resonances,

\[ N\pi \rightarrow \Delta(1230), \quad \Delta(1230) \rightarrow N\pi(\pi), \]  \hspace{1cm} (4.35)

generate transverse flow. The origin of the transverse flow effect in the decays of the \( \Delta \) resonances is the bias from the integration over the transverse momentum, which makes the lightest resonances with nonzero transverse velocity contribute the most to the final pion spectrum. To clarify, let us note that the thermal Boltzmann distribution is a probability function falling with the transverse momentum of the resonance, while the \( 2\pi p_t \) factor, coming from the integration over \( p_t \), rises. As a result nonzero transverse velocities are selected as the most probable for the resonances. The transverse expansion of the resonances will further increase the boost of pions in the transverse direction. Both factors, the cooling of pions and the excitation and subsequent decay of the \( \Delta(1230) \) resonances and heavier nucleon resonances via three particle modes must increase the low transverse mass yield of the pions. Expansion of nuclear matter is accompanied by its cooling. According to the cascade model simulations the expansion time for a central \( Au+Au \) collision constitutes 10–15 \( fm/c \) [12], significantly longer than the lifetime of the nucleon resonances. The expansion must lead to a reduction of the resonance occupation probabilities. Thus produced pions, having a large interaction cross section and the small mass, have a good chance to thermalize with the rest of the nuclear matter.

Due to the consistency of the concave shape of invariant spectra of pions with the presence of the transverse expansion of the pion gas we employed the three dimen-
sional expansion picture described via equations (4.23–4.27) in order to parametrize the spectra of pions in a wide range of rapidities. The results of such parametrization are presented in Table 4.2 and are shown by the solid curves on Figure 4.7. One can see that three dimensional expansion picture provides a reasonably good description of the invariant $\pi^-$ multiplicities in a wide range of rapidities. However, a more detailed inspection shows that the high $m_t$ region for the bins near midrapidity is underestimated, which for example can be attributed to the different global variables of the thermalized pions and the pions originating from the nucleon resonances after the freeze-out. Monte-Carlo simulations show that, the pions produced in the two particle decays of resonance $N^*(1440)$ and heavier, have a flat transverse mass distribution and are capable at producing the effect of stiff pions. Another possible reason for this small disagreement is the blast wave approximation itself, namely the replacement of a continuous flow field by the flow surface. We also estimate an increasing systematic error for the pion spectrum in the region of midrapidity. Measurement of higher transverse masses in this region together with $\pi p$ correlation studies may clarify the relative role of the primary pions and of those coming from the heavy resonance decays.

4.3.2 Rapidity Distributions of $\pi^-$

The wide coverage in the transverse momentum for pions allows us to calculate the rapidity distribution with relatively small errors caused by extrapolation to the
Au+Au – $\pi^-$ Rapidity Distribution, E891

Figure 4.6 E891 rapidity distribution of the $\pi^-$ production in Au + Au. Solid curve is the result of the Flow Model parametrization (4.23) integrated over the transverse mass squared. Dashed curve is the ARC prediction. Dotted curve is the RQMD prediction. Small error bars correspond to the statistical errors, while the large ones include the systematic errors.
$\pi^-$ Spectra Au+Au, E891

![Graph showing double differential spectra for $\pi^-$ meson production.](image)

**Figure 4.7** E891 double differential spectra for $\pi^-$ meson production. Dashed lines correspond to the Flow Model parametrization. Solid lines - the Flow Model parametrization. Starting with rapidity bin (1.4 < y < 1.6) each successive spectrum is divided by increasing powers of ten. Errors are statistical only.
higher transverse momenta. We present the results of the measurements for the rapidity distributions of $\pi^-$ in central $Au + Au$ collisions on Figure 4.6. Extrapolations to higher values of $m_t - m_0 > 1 \text{ GeV}/c^2$ are done according to the results of the parametrization (4.23-4.27) with the parameters given in Table 4.2. The tail contributions to the resulting rapidity density of $\pi^-$ for all rapidities are smaller than 4%. We, therefore, do not introduce a significant systematic error via addition of the extrapolated tail. We have performed the numerical integration of the Flow Model

<table>
<thead>
<tr>
<th>Rapidity</th>
<th>$dN_{\pi^-}/dy$</th>
<th>$\sqrt{\sigma_{stat}^2 + \sigma_{syst}^2}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$1.6 &lt; y &lt; 1.8$</td>
<td>$75.22 \pm 1.31$</td>
<td>$\pm 6.62$</td>
</tr>
<tr>
<td>$1.8 &lt; y &lt; 2.0$</td>
<td>$73.94 \pm 0.94$</td>
<td>$\pm 5.69$</td>
</tr>
<tr>
<td>$2.0 &lt; y &lt; 2.2$</td>
<td>$67.17 \pm 0.51$</td>
<td>$\pm 3.41$</td>
</tr>
<tr>
<td>$2.2 &lt; y &lt; 2.4$</td>
<td>$57.44 \pm 0.32$</td>
<td>$\pm 2.44$</td>
</tr>
<tr>
<td>$2.4 &lt; y &lt; 2.6$</td>
<td>$46.05 \pm 0.26$</td>
<td>$\pm 0.63$</td>
</tr>
<tr>
<td>$2.6 &lt; y &lt; 2.8$</td>
<td>$36.35 \pm 0.22$</td>
<td>$\pm 0.47$</td>
</tr>
<tr>
<td>$2.8 &lt; y &lt; 3.0$</td>
<td>$26.67 \pm 0.19$</td>
<td>$\pm 0.36$</td>
</tr>
<tr>
<td>$3.0 &lt; y &lt; 3.2$</td>
<td>$21.06 \pm 0.17$</td>
<td>$\pm 0.29$</td>
</tr>
<tr>
<td>$3.2 &lt; y &lt; 3.4$</td>
<td>$12.83 \pm 0.13$</td>
<td>$\pm 0.21$</td>
</tr>
<tr>
<td>$3.4 &lt; y &lt; 3.6$</td>
<td>$7.05 \pm 0.09$</td>
<td>$\pm 0.17$</td>
</tr>
</tbody>
</table>

Table 4.4 Rapidity distribution of $\pi^-$ mesons.

The results of numerical integration of (4.23-4.27) approximately equal the results of calculation according to the formula (4.29) and match the ones for the transverse expansion velocity equal to zero.
results on Figure 4.6 via the dashed line. It turns out that the Flow Model (4.29) underestimates the experimentally measured rapidity density in the region close to midrapidity, \( y_0 = 1.6 \). However, the difference may be attributed to the increased systematic error in this region. The total errors on Figure 4.6 including the systematic ones are indicated by the large error bars, while the small error bars correspond to the statistical errors only.

4.3.3 Mean Multiplicity of \( \pi^- \)

Since experimental acceptance for \( \pi^- \) covers most of the phase space in the forward hemisphere of the center of mass system and the collision is symmetric relative to the forward-backward transformation, we can obtain the integrated yield of \( \pi^- \) via summation of all \( dN/dy \) data for \( y > 1.6 \) and multiplying the result by two. Thus we obtain for the mean multiplicity of negative pions

\[
< N_{\pi^-} >= 170 \pm 0.71,
\]  

(4.36)

corresponding to the region of our acceptance, \(|y - y_0| < 2.0\). The error is statistical only. The contribution from the unmeasured tails, \(|y - y_0| > 2.0\), estimated via the results of the Flow Model parametrization, is rather small, \(2.84 \pm 0.05\).

4.4 Comparison with Microscopic Cascade Models

Cascade simulation codes have been developed and used for comparison with experimental data to study the mechanisms of relativistic heavy-ion collisions. Some
models are based on hadronic scattering and a string formation picture, and some on purely hadronic picture. The first type includes FRITIOF [94], HIJET [95], VENUS [96] and RQMD [97], while ARC [99] is one of purely hadronic models, treating low-mass resonances as its main ingredients. Among these models ARC and RQMD are particularly suitable for the AGS energy region. A brief description of these models can be found in Appendix D

4.4.1 Cascade Model Calculations

The RQMD 1.08 output was generated by Efstratios Efstathiadis with the impact parameter distributed evenly in two dimensions within the circle of radius of 4 fm. The ARC model output was generated by one of the ARC authors, David Kahana. The impact parameters below 4 fm were distributed evenly in two dimensions. Since our experimentally measured cross section equals 4.4% of the geometric cross section, we have used the events with the impact parameter not exceeding 3 fm in order to match experimentally measured interaction cross section (see Equation (3.2.3) in Chapter 3). 1089 events from RQMD and 3671 events from ARC have passed these cuts and were used for the comparison with the data. It is important to note here that the multiplicity cut, used in the data for selection of the most central events, is not exactly equivalent to the impact parameter cut used in the Monte-Carlo.

We have performed Monte-Carlo simulation of the resulting multiplicity distribution using ARC and RQMD outputs as inputs for the simulation program GEANT.
The outputs of GEANT have been analyzed by the same analysis program as the real data. The obtained multiplicity distributions for the charged tracks from the target are found to be in a good agreement with the data.

4.4.2 Double Differential Spectra of Negative Pions

Both cascade models, ARC and RQMD reproduce the double differential distributions of the $\pi^-$ production reasonably well as can be verified on Figure 4.8. The dashed histograms correspond to the predictions of ARC model while the dotted ones denote the RQMD predictions. The invariant transverse mass distributions are steeper in the region of high rapidities than the prediction of both models. RQMD seems to be better describing the high $m_t$ tails near midrapidity than it is done by ARC.

4.4.3 Rapidity Distribution of Negative Pions

Rather good description for the rapidity distributions of the negative pions is made by both cascade models as can be seen on Figure 4.6. The two models also agree with each other well in their description of the rapidity density for the $\pi^-$ mesons. A slight overestimation of the rapidity density is given by both models in the region of high rapidities.
Figure 4.8 E891 double differential spectra for $\pi^-$ meson production. Dashed line is the ARC prediction. Dotted line is RQMD prediction. Errors are statistical only.
4.4.4 Mean Multiplicity of Negative Pions

To compare the total yields of the negative pions let us consider only the pions within our acceptance interval of rapidities and use the forward-backward symmetry of the collision to reflect our measurements to the backward hemisphere of the center of mass system of the collision system. We have obtained the mean multiplicity of $\pi^-$ equal to $170 \pm 0.71$ as in (4.36). For the model ARC we obtain $N_{\pi^-} = 181.7 \pm 0.2$. For the model RQMD, $N_{\pi^-} = 175.8 \pm 0.4$. Errors quoted are statistical only.

4.4.5 Double Differential Spectra of Lambdas

The comparison of the cascade model predictions with the E891 measurements on the double differential cross sections for the $\Lambda$ production in the central $Au + Au$ interactions can be seen on Figure 4.2. While the model ARC seems to give better predictions in the region of high rapidities, the RQMD is closer to the experimental data in the region around midrapidity.

Both models considered in this thesis are purely cascade versions, where the mean field effects are not included. The $\Lambda N$ interaction, analogously to the $NN$ interactions, becomes repulsive at the nuclear densities above normal, which leads to the increasing transverse flow manifesting itself in the data via the deviation from a single exponential scaling in the midrapidity region, an effect which both cascade models clearly underestimate.
We have parametrized the invariant spectrum of $\Lambda$ hyperons from the RQMD using our Flow Model parametrization in order to compare amounts of the transverse and the longitudinal expansion observed in the data with those predicted via RQMD. The results of such parametrization are presented in Table 4.2. In terms of our parametrization the RQMD significantly underpredicts the strengths of both longitudinal and the transverse expansion of the $\Lambda$ gas and overpredicts the $\Lambda$ temperature.

4.4.6 Rapidity Distribution of Lambdas

The rapidity distributions of the $\Lambda$ hyperons are compared with the cascade model prediction on Figure 4.5. Again a better prediction for high rapidities is provided by the ARC and for the midrapidity region, by RQMD.

Both models predict much narrower distribution of $\Lambda$ hyperons over rapidity than the one observed in the experiment E891. This disagreement is associated with the different role of collective motion of the $\Lambda$ hyperons, which, as mentioned earlier, is probably caused by the mean field effects which are not included in both models.

4.4.7 Mean Multiplicity of Lambdas

Since we do not have a sufficient rapidity coverage to measure the mean multiplicity of $\Lambda$ hyperons without extrapolations, we will compare the mean multiplicity, given by the Flow Model parametrization, as the latter describes the double differen-
tial cross sections for the Λ production and the corresponding rapidity distributions fairly well. As listed in Table 4.2, the total number of Λ hyperons emitted via the flow surface of the thermal sources, $N_{3D}$, is found to be $20 \pm 3$. The ARC model predicts $<N_\Lambda> = 25.1 \pm 0.1$ per one central event. The model RQMD gives a prediction of $<N_\Lambda> = 17.9 \pm 0.1$, which is consistent with the result from the Flow Model parametrization of RQMD Λ spectrum presented in Table 4.2.

4.4.8 Conclusions

As discussed in the preceding sections, there exist some aspects of $Au + Au$ collisions at the AGS energy which can not be reproduced by the present cascade models, even though the models are generally successful in describing the $Au + Au$ ion collisions as well as lighter systems in the AGS energy region.

The aim of the current heavy ion physics experiments at AGS is a search for and identification of signals of the QGP formation. Among various suggested signals of the QGP in our experiment it is feasible to look at such features as strangeness enhancement and the collective nuclear flow. As in QGP the mechanism of the strange quark production may be altered, it was expected that the strange particle yields must increase in comparison to the superposition of the hadron-hadron interactions. It is expected also that the first order phase transition from QGP to the hadronic fluid must lead to the reduction of the collective nuclear flow as compared to the picture with hadronic degrees of freedom only. The hadronic models such as ARC and RQMD
thus provide us with a useful background for comparison of these observables with the experimental data.

In this thesis we have reported an extended analysis of the E891 experimental data on the $\Lambda$ production as compared to our original publication \cite{89}. In accord with the measurements for lighter systems \cite{88}, we do not observe an enhancement of the strange hyperon $\Lambda$ yield as compared to the superposition of hadronic collisions simulated via the cascade models ARC and RQMD. However, it is possible that the strangeness enhancement in the QGP phase is destroyed after the phase transition to the hadronic matter.

The evaluation of the collective flow thus seems to be a more promising signature of the QGP formation. To get a quantitative insight to the picture of expansion of the particle gases after the freeze-out of nuclear fluids, we have constructed the parametrization of the particle spectra, called here a Flow Model parametrization, according to the 3D expansion scenario. This parametrization allows one to take into account the change in the particle global variables due to the collective motion of nuclear matter. The flow velocity field is described in a blast wave approximation via the azimuthally symmetric surface in the space of the longitudinal boost angle and the transverse two dimensional velocity. The application of the Flow Model to the measured invariant differential and the rapidity distribution of particles lead us to the good description of the measured $\Lambda$ spectra including the interesting midrapidity effect of the low $m_t$ suppression and the spectra of $\pi^-$ mesons. The small disagreement
of the Flow Model parametrization with the data for \( \pi^- \) in the region near midrapidity may be attributed to the blast wave approximation used — approximation of a continuous flow field by the flow surface. It is also likely that the disagreement is caused by the difference of the global variables of the thermalized pions and the pions originating from the nucleon resonance decays after freeze-out. We, however, estimate an increasing systematic error in the midrapidity region.

For the comparison with the data we have also parametrized the RQMD predictions for the \( \Lambda \) and the \( \pi^- \) double differential spectra. The spectral shapes for \( \Lambda \) and \( \pi^- \) predicted by ARC are such that all attempts to parametrize the ARC predictions with our approach resulted in a poor convergence of the minimization procedures. Rather close parameters were obtained for \( \pi^- \) in the data and RQMD, however there are some differences, which in terms of our parametrizations may be expressed as follows: 1) the temperature of \( \pi^- \) mesons in RQMD is by 26\% larger than in the data, 2) the strength of longitudinal flow, \( \langle \eta \rangle \), is overpredicted by RQMD by 7\%, 3) the mean transverse flow velocity, \( \langle v_t \rangle \), is found to be by 10\% larger in the data than in RQMD.

The basic differences between the experimental data and the cascade model predictions are however in the \( \Lambda \) spectra, where more longitudinal and the transverse flow is observed than it is predicted by RQMD. Even though with our Flow Model we could not well parametrize the ARC results, the width of the \( \Lambda \) rapidity distribution from ARC is small compared to that from the data, which indicates underestima-
tion of the longitudinal flow in ARC. Instead of expected reduction of the collective nuclear flow for the $\Lambda$ hyperons we thus observe an increase of flow as compared to the cascade models. Both models, ARC and RQMD used in this thesis however did not include the effects of the mean field potentials which are believed to have a significant influence on the strength of expansion of nuclear matter [12]. Repulsive mean field potential for $\Lambda$ hyperons must lead to the increase of the flow strength for these particles compared to the pure cascade, while the good agreement of the flow parameters for $\pi^{-}$ mesons with the prediction of RQMD may indicate that the pion flow is mainly caused by the thermal pressure rather than by the mean field potential. The models thus need to be improved to make it possible to look for the flow signature of the QGP formation.
Chapter 5

Summary and Conclusions

Using Au beam of 11.6 A GeV/c at AGS we have performed AGS-E891 experiment aimed at the study of the spectra of particles produced in the central Au + Au collisions. To deal with the experimental difficulties concerning to the gold beam we have studied the important hardware effects associated with intensive ionization inside the Time Projection Chambers. Examining the consistency of the data with Monte-Carlo for the distributions of various kinematical and geometry variables we have found the regions of the phase space where the influence of dynamic distortions on the measurement of $\Lambda$ momentum is crucial. This allowed us to evaluate the systematic errors caused by dynamic field distortions and resulting hit position mismeasurement.

We have presented the results of the first measurement of the $\Lambda$ double differential cross section and rapidity distributions in central Au + Au collisions at AGS. In addition to our original publication [89], the data on the invariant differential distributions near midrapidity ($y < 2.0$) is analyzed. The complementary results for $\pi^-$ spectrum in the kinematical region overlapping with earlier measurements by AGS-E866 and AGS-E877 collaborations have been obtained. The experimental data were taken in October 1994 for a short period of two weeks and analyzed fully by December 1997.
Shapes of the spectra of particles produced in central \( Au + Au \) collisions, as observed in our experiment, differ significantly as compared to the mere superposition of nucleon-nucleon collisions and as compared to the lighter nuclei. We have observed

- Deviation from the single exponential scaling, namely, a convex shape of the \( \Lambda \) transverse mass distributions with the curvatures and inverse slopes increasing towards midrapidity. The effect is consistent with transverse radial expansion.

- A concave spectrum for \( \pi^- \) meson transverse mass distributions. The effect is consistent with the increased role of secondary interactions, production and decay of nucleon resonances and with transverse radial expansion.

- The flattening of the invariant transverse mass distributions for both \( \Lambda \) and \( \pi^- \) close to midrapidity.

The collective effects thus are confirmed to play more important role in the \( Au + Au \) collisions compared to light systems.

Wide rapidity coverage allowed us to study the trends of the inverse slopes and the curvatures of the transverse mass distributions for the particles of both types as functions of rapidity and to associate these trends with the picture of the three dimensional expansion of the particle gases after the nuclear fluid freeze-out.

Application of our 3D Flow Model to the parametrization of the invariant differential multiplicities and the rapidity densities of \( \Lambda \) hyperons has lead us to a good description of both distributions and to the possibility of measuring the values of the
transverse and longitudinal expansion for the gas of $\Lambda$ hyperons as it is observed in the detectors. A fairly good description is also achieved for the negative pions in a wide range of rapidities. There is however some disagreement in the description of the high $m_t$ tail and the rapidity density in the region near midrapidity. Due to the fact that at high temperatures reached in $Au + Au$ collisions the large fraction of the nucleons must be excited to the states of higher mass and momentum - the nucleon resonances, $\Delta$ and $N^*$. One of the possible explanations is thus a difference of the global variables for the thermalized pions and those originating from the resonance decays after the freeze-out. However our evaluation of the systematic errors in the region of midrapidity shows the difference may be attributed to the apparatus effects.

The results of our measurements for $\Lambda$ and $\pi^-$ double differential cross sections and rapidity distributions were compared to the predictions of cascade model ARC and RQMD. The reasonable agreement was found. However, there are some features of the observed particle spectra which can not be described via these cascade models. The most significant difference is underestimation of the transverse and the longitudinal flow for the $\Lambda$ production. One of the possible explanations for this disagreement is the significance of the mean field effects, which were not included in the cascade models. Namely, apart from the thermal pressure, causing the expansion of the nuclear matter, the repulsive mean field potential, increasing at the nuclear densities above normal, strengthens the transverse as well as the longitudinal expansion of baryons, the effect we observe in our experiment. Contrary to the $\Lambda$ hyperons the cascade model RQMD
provides a good prediction for the flow parameters of $\pi^-$ mesons found in our study.

It is thus possible to conclude that the flow of pions is mainly caused by the thermal pressure rather than by the mean field.
Appendix A

BNL-AGS-E891 Collaboration

*BNL/LBNL/CCNY/Rice*

*Brookhaven National Laboratory*

S.E. Eiseman, A. Etkin, K.J. Foley, R.W. Hackenburg, R.S. Longacre, W.A. Love

and A.C. Saulys

*Brookhaven National Laboratory/CCNY*

S.J. Lindenbaum *Lawrence Berkeley National Laboratory*

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C.S. Chan, E. Efstathiadis, M.A. Kramer,

K. Zhao and Y. Zhu

*Rice University*

S. Ahmad, B.E. Bonner, J.M. Clement, S.V. Efremov, G.S. Mutchler and E.D. Platner
Appendix B

The Vertex "N-tuple"

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<th>Description</th>
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<td>7</td>
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<td>Number of KOS candidates</td>
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<tr>
<td>17</td>
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<td>REAL</td>
<td>Y coordinate of the production vertex</td>
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<tr>
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Appendix C

The Track "N-tuple"

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<td>Event number</td>
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<td>Y component of particle momentum</td>
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<tr>
<td>16</td>
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<td>REAL</td>
<td>Z component of particle momentum</td>
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<td>Total momentum of particle</td>
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<td>VX</td>
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<td>VY</td>
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Appendix D

Cascade Models of Nucleus–Nucleus Collisions

The basic ideas and ingredients of the cascade models ARC and RQMD are summarized here. We compare the results of our experiment with the output of these cascade models in the discussions of this thesis.

D.1 RQMD

Relativistic Quantum Molecular Dynamics model has been developed at Frankfurt University, to be applied in the energy range from the AGS to the CERN-SPS. The model is a combination of classical propagation of hadrons and quantum effects such as stochastic scattering, particle decay, and Pauli blocking in collisions. Color string formation and fragmentation are also used to describe particle production from decay of high-mass resonances (above 2 GeV/c\(^2\) for non-strange baryons), mainly to handle the energy region of the SPS. A difference of RQMD from non-relativistic QMD model is Lorentz invariance; instead of the propagation with non-relativistic Hamiltonian in QMD the Lorentz invariant equations of motion are used in RQMD.

As a principle, the model includes all available experimental data on the reaction cross sections and decay probabilities. For reactions whose parameters are experimentally undetermined, the model assumes their cross section following from fundamental
invariance principles for the strong interactions, which however introduces a few free parameters to the model yet to be fixed.

D.2 ARC

A Relativistic Cascade is a simulation code developed at BNL to describe nucleus-nucleus collisions at the AGS energy based only on a hadronic picture. Particle production in the model is achieved through low-mass resonances such as $\Delta$ and $\rho$. This is motivated by examination of $p + p$ collisions at 12 and 24 $GeV/c$ in the laboratory frame, showing that a majority of inelastic collisions contains one or more such resonances in the final state. The finite lifetimes of the resonances effectively play the role of so-called formation time of secondary particles in other models. On the other hand, it is a possible cause of the less stopping of incident nucleons in the model compared to RQMD that ARC includes only low mass resonances and hence can not transform much kinetic into excitation energy. Input parameters of two-body reactions in the code are adjusted so that it reproduces experimental data of relevant cross sections and transverse-momentum distributions in $p + p$ and $p + A$ collisions at $12 - 24$ $GeV/c$. 
Appendix E

Occupation Probabilities of the Nucleon Resonances

In the following Table the occupation probabilities of the excited states of the nucleon at temperature of 150 MeV are listed together with relative densities. This Table is borrowed from reference [105]. The excitation energy $\epsilon_i$ is taken to the mass of $i$-th species. The normalization factors $\rho_i$ and $Z$ are the same over all states shown. The values of $\rho_i/\rho_t$ are independent of the total baryon density at this temperature (Boltzmann regime). The relative density of $i$-th state is determined according to the equation

$$\frac{\rho_i}{\rho_t} = \frac{1}{\rho_t} \frac{g_i}{2\pi^2} \int_0^\infty dp \frac{p^2}{\exp[\beta (\epsilon_i - \mu)] + 1},$$  \hspace{1cm} (E.1)$$

where $\beta = 1/T$, $\epsilon_i = (p^2 + m_i^2)^{1/2}$, $\rho_t = \sum_i \rho_i$. The chemical potential $\mu$ is determined self-consistently from a total density of $4\rho_0$, assuming the chemical equilibrium between all species.
Table E.1 Occupation probabilities of excited states of the nucleon and the relative densities in a thermal system at a temperature of 150 MeV.

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<tr>
<th>State</th>
<th>$J^P$</th>
<th>$g_i$</th>
<th>($g_i e^{-\xi_i/T}$)</th>
<th>$\rho_i/\rho_t$</th>
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